

# Squark Production in the R-Symmetric Supersymmetric Standardmodel

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## Summary

Abstract

English: flaws of the SM

motivation: aesthetic: Coleman-Mandula  $\rightarrow$  Haag-Lopuszanski-Sohnius-Theorem

Abstract

Deutsch

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# 1 Introduction

motivation: aesthetic: Coleman-Mandula  $\rightarrow$  Haag-Lopuszanski-Sohnius-Theorem

plots for exclusion of squarks in specific SUSY scenarios (from Michael)  $\rightarrow$  R-Symmetry could be possible explanation for that because:

MSSM-Lagrangian  $\rightarrow$  trafo rules for superfields under R-symm  $\rightarrow$  forbidden terms in MRSSM (write down Lagrangian for R-symmetric SUSYQCD)

suppression of squark production in MRSSM by less diagrams ( $m_{gluino}^{-4}$  suppression at low energies in MRSSM and only  $m_{gluino}^{-2}$  suppression in MSSM)

R-charges of all fields (show in diagram!)  $\rightarrow$  only if R-charges of final / initial particles are zero, a diagram is allowed in R-symm. model

references to build in

- "Matching Squark Pair Production at NLO with Parton Showers" from Gavin, Hangst, Krämer, Mühlleitner,.. for complete treatment of NLO calculation
- "DIRAC gAUGINOS IN susy - sUPPRESSED jETS + <met sIGNALS: a sNOWMASS WHITEPAPER" FROM kRIBS, mARTIN for Squark production at LO, allude to same result
- "Dirac Gaugino Masses and supersoft SUSY breaking" from Fox, Weiner, Nelson for the introduction of the MRSSM

## 2 The Standard Model

The Standard Model of particle physics is the commonly accepted theory describing the world's fundamental particles and their interactions. It is a gauge quantum field theory which is characterized by its invariance under symmetry groups. The Standard Model contains different fields, whose quantized excitations are interpreted as particles.

This chapter summarizes the most important aspects of the Standard Model.

### 2.1 Symmetries and Transformations

#### Space-time Symmetries

The Standard Model is defined on Minkowski space, whose coordinates are label with  $x^\mu$   $\mu \in \{0, 1, 2, 4\}$ . As a relativistic theory it is invariant under Poincaré transformations, i.e. it is invariant under Lorentz-transformations (with generators  $J^{\mu\nu}$ ) and translations (with generators  $P^\mu$ ) in spacetime. The set of all Poincaré transformations form the Poincaré group, which is a Lie group. Its generators obey the Poincaré-algebra

$$\begin{aligned} [P^\mu, P^\nu] &= 0 \\ [P^\mu, J^{\nu\rho}] &= i(g^{\mu\nu} P^\rho - g^{\mu\rho} P^\nu) \\ [J^{\mu\nu}, J^{\rho\sigma}] &= i(g^{\nu\rho} J^{\mu\sigma} + g^{\mu\sigma} J^{\nu\rho} - g^{\mu\rho} J^{\nu\sigma} - g^{\nu\sigma} J^{\mu\rho}). \end{aligned} \quad (2.1)$$

The fields of the Standard Model transform in different representations of the Poincaré-group [1].

#### Gauge Symmetries

In order to describe interactions of matter particles gauge theories are used. In the Standard Model matter fields are described by Dirac spinors. The Lagrangian of a free Dirac field reads

$$\mathcal{L}_{Dirac} = \bar{\Psi}(i\cancel{\partial} - m)\Psi. \quad (2.2)$$

To include interactions one imposes a local group symmetry (gauge symmetry) upon this Lagrangian. A spinor transforms under a generic gauge transformation like

$$\Psi(x) \rightarrow U(x)\Psi(x), \quad (2.3)$$

where  $U(x)$  is an element of the gauge group in question. Because the gauge group is a unitary matrix Lie group it can be written in the form  $U(x) = \exp(-igT^a\theta^a(x))$ . Here  $T^a$  are the

self-adjoint generators of the associated Lie algebra which obey

$$[T^a, T^b] = if_{abc}T^c \quad (2.4)$$

where  $f_{abc}$  are the structure constants of a Lie algebra,  $g$  is the coupling constant of the gauge group and  $\theta^a(x)$  are local parameters.

Because the parameters of the gauge group are local the derivative in 2.2 spoils the gauge invariance. In order to rectify gauge invariance of the Lagrangian one introduces a further field for each index  $a$  of the generators - the gauge vector  $G^{a\mu}$ . Defining the transformation of the matrix valued gauge vector  $G^\mu := G^{a\mu}T^a$  as

$$G^\mu(x) \rightarrow U^{-1}(x) \left( G^\mu(x) + \frac{i}{g} \partial^\mu \right) U(x) \quad (2.5)$$

and introducing the gauge covariant derivative

$$D^\mu = \partial^\mu + igT^a G^{a\mu} \quad (2.6)$$

one finds that the expression  $D_\mu \Psi(x)$  transforms as

$$D_\mu \Psi(x) \rightarrow U(x) D_\mu \Psi(x) \quad (2.7)$$

Therefore gauge invariance is restored in eq. 2.2 by replacing  $\partial_\mu$  with  $D_\mu$ . But if the gauge vector is interpreted as a physical field there must apart from the so far introduced interaction term also be a kinetic term associated with it. Using eq. 2.7 one defines the field strength tensor<sup>1</sup>  $F^{a\mu\nu}$  whose matrix valued form

$$F^{\mu\nu} = F^{a\mu\nu}T^a := \frac{1}{ig}[D^\mu, D^\nu] = \partial^\mu G^\nu - \partial^\nu G^\mu - gf_{abc}T^c G^{a\mu}G^{b\nu} \quad (2.8)$$

transforms as  $F^{\mu\nu} \rightarrow U(x)F^{\mu\nu}U^{-1}(x)$ . Using the cyclic property of the trace and the Dynkin index  $T(F)$  defined in eq. 10.6 in the appendix one can write down a gauge invariant kinetic term for the gauge vector:

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{2}\text{Tr}(F^{\mu\nu}F_{\mu\nu}) = -\frac{T(F)}{2}F^{a\mu\nu}F_{\mu\nu}^a \quad (2.9)$$

This completes the construction of a Lagrangian which is invariant under non-abelian gauge group transformations. The result is the famous Yang and Mills Lagrangian [3]

$$\mathcal{L}_{\text{Yang-Mills}} = \bar{\Psi}i\not{D}\Psi - \frac{1}{4}F^{a\mu\nu}F_{\mu\nu}^a \quad (2.10)$$

---

<sup>1</sup>An alternative construction of the field strength tensor makes use of the gauge invariant Wilson loop. This gives some insights into the geometry of gauge transformations [2].



This Lagrangian gives rise to spin- $\frac{1}{2}$  (matter) particles which interact with spin-1 (force mediator) particles. Furthermore if the gauge group is non-abelian, i.e.  $f_{abc} \neq 0$ , there are self interactions among the spin-1 particles.

The gauge group of the Standard Model is a direct product of the three gauge groups<sup>2</sup>:  $U_Y(1)$ ,  $SU_L(2)$  and  $SU_C(3)$ . The elements  $U(x)$  of those are given in table 2.1.

These gauge groups give rise to 3 forces: the strong force, the weak force and the electromagnetic force.

Gauge Group	Group Element
$U_Y(1)$	$U(x) = \exp\left(-ig_Y \frac{\hat{Y}}{2} \theta_Y(x)\right)$
$SU_L(2)$	$U(x) = \exp\left(-ig_w \vec{\tau} \cdot \vec{\theta}_w(x)\right)$
$SU_C(3)$	$U(x) = \exp(-ig_s T^a \cdot \theta_s^a(x))$

**Table 2.1:** The table lists the explicit element  $U(x)$  of the gauge groups  $U_Y(1)$ ,  $SU_L(2)$  and  $SU_C(3)$ . The hypercharge operator  $\hat{Y}$  gives the eigenvalue of the hypercharge of the field it is applied to (see table 2.3).  $\vec{\tau}$  and  $T^a$  are the generators of  $SU_L(2)$  and  $SU_C(3)$  respectively. In the fundamental representation  $\vec{\tau} = \frac{\vec{\sigma}}{2}$  where  $\vec{\sigma}$  has the 3 Pauli matrices as components and  $T^a = \frac{\lambda^a}{2}$  where  $\lambda^a$  are the 8 Gell-Mann matrices.  $\varepsilon_{abc}$  and  $f_{abc}$  are the structure constants of  $SU_L(2)$  and  $SU_C(3)$  respectively.

## 2.2 The Particles of the Standard Model

In the Standard Model different matter particles take part in different interactions, i.e. their corresponding spinor couples to different gauge vectors.

If a spinor couples to a certain gauge vector it transforms non trivially (like indicated in table 2.1) under the gauge group which is associated with this gauge vector<sup>3</sup>. This means if a particle couples to a certain force its charge which is associated with this force is nonzero.

The charges of particles for a force are defined as the eigenvalues of the generators which correspond to the force.

### The Quarks:

Quarks are strongly interacting fermions, which means their spinors transform non trivially under  $SU_C(3)$ . Because they transform in the fundamental representation of  $SU_C(3)$  this means a quark spinor is built up by 3 spinors each carrying another color. This splitting of the quark spinor in colors is often suppressed for the sake of simplicity. This convention is adopted throughout this thesis.

<sup>2</sup>The subscript stands for the associated charge of the groups respectively:  $Y$  for hypercharge,  $L$  for left handedness (weak Isospin  $I_3$ ) and  $C$  for color

<sup>3</sup>In the Standard Model all matter particles transform in the fundamental (or trivial) representation of gauge groups.

Furthermore the left handed component of quarks interact weakly, which means that their spinors<sup>4</sup> transform (in the fundamental representation) under  $SU_L(2)$  transformations meaning that 2 left handed quark spinors are assembled within a doublet.

Finally all quarks carry a hypercharge. In section 2.3 the mechanism of electroweak symmetry breaking is described. This mechanism explains how electromagnetism arises from the groups  $SU_L(2)$  and  $U_Y(1)$ . All quarks interact electromagnetically.

After all there are 6 quarks which are listed in table 2.2. They are categorized in 3 generations because their quantum numbers except for their masses reoccur in each generation. The two types of quarks within a generation which have distinct quantum numbers are referred to as up-type and down-type quarks. An up-type-quark and the down-type quark of the same generation built up a doublet.

### The Leptons:

Leptons do not interact strongly. They take part in the weak and the electromagnetic interaction, i.e. their spinors transform under the fundamental representation of  $SU_L(2)$  and  $U_Y(1)$ . As for the quarks only the left handed components interact weakly.

As for the quarks there are 6 leptons which are classified into 3 generations (see table 2.2). In each generation is a lepton with a negative electrical charged and an electrically neutral lepton. The latter ones are referred to as neutrinos. Right handed neutrinos have not been observed (yet) and are therefore absent in the SM. The former are called electron, muon and tau. Each left handed leptons with an electric charge is assembled with its neutrino in a doublet.

Particle	1 <sup>st</sup> generation	2 <sup>nd</sup> generation	3 <sup>rd</sup> generation
$u_i$ up-type-Quark	$u$ up-Quark	$c$ charm-Quark	$t$ top-Quark
$d_i$ down-type-Quark	$d$ down-Quark	$s$ strange-Quark	$b$ bottom-Quark
$e_i$ Charged Lepton	$e$ Electron	$\mu$ Muon	$\tau$ Tau
$\nu_i$ Neutrino	$\nu_e$ Electron Neutrino	$\nu_\mu$ Muon Neutrino	$\nu_\tau$ Tau Neutrino

**Table 2.2:** The matter particles of the SM. Listed are the symbol and the name of the particles.

Quarks and Leptons are the matter particles of the Standard Model. They are listed together with their charges for the different forces in table 2.3. There is the color for strong interactions, the third component of the weak isospin  $I_3$  for weak interactions (the eigenvalue of the third generator of the  $SU_L(2)$ ) and the half of the hypercharge  $\frac{Y}{2}$  to obtain the electric charge  $Q$  via the Gell-Mann–Nishijima formula:  $Q = I_3 + \frac{Y}{2}$ .

Because the left and right-handed parts of spinors transform differently under the  $SU_L(2)$  they

<sup>4</sup>The left handed part of a 4-spinor  $\Psi$  is projected out by the appropriate projector  $P_L$ . This is explained in Appendix 10.3.

are listed separately. All quarks occur with three different colors.

In the last row the Higgs-boson is listed. Its associated field is responsible for the mass of elementary particles. That is explained in section 2.3.

Particle	Symbol	color	$I_3$	$\frac{Y}{2}$	Q
Left handed Quarks	$Q_{iL} = \begin{pmatrix} u_{iL} \\ d_{iL} \end{pmatrix}$	red, green, blue	$\begin{pmatrix} +\frac{1}{2} \\ -\frac{1}{2} \end{pmatrix}$	$+\frac{1}{6}$	$\begin{pmatrix} +\frac{2}{3} \\ -\frac{1}{3} \end{pmatrix}$
Right-handed Quarks	$u_{iR}$	red, green, blue	0	$+\frac{2}{3}$	$+\frac{2}{3}$
	$d_{iR}$	red, green, blue	0	$-\frac{1}{3}$	$-\frac{1}{3}$
Left-handed Leptons	$\ell_{iL} = \begin{pmatrix} \nu_{iL} \\ e_{iL} \end{pmatrix}$	-	$\begin{pmatrix} +\frac{1}{2} \\ -\frac{1}{2} \end{pmatrix}$	$-\frac{1}{2}$	$\begin{pmatrix} 0 \\ -1 \end{pmatrix}$
Right-handed Leptons	$e_{iR}$	-	0	+1	+1
Higgs	$H$	-	$-\frac{1}{2}$	$+\frac{1}{2}$	0

**Table 2.3:** This table lists all matter particles in the Standard Model and the Higgs particle with their charges for all forces. This is the color, the weak isospin  $I_3$ , the half of their hypercharge and their electrical charge. The index  $i = 1, 2, 3$  labels the generation of the matter particles and is written out in table 2.2. If there are no colors specified or charges are zero this means that the fields in question transform trivially under the pertaining gauge transformation.

## The Force Particles

before EWSB			after EWSB		
group	coupling constant	gauge field	coupling constant	gauge field	Particle
$SU_C(3)$	$g_s$	$G_\mu^a$	$g_s$	$G_\mu^a$	Gluon
$SU_L(2)$	$g_w$	$W_\mu^b$	$g_W = \sqrt{2}g_w,$ $g_Z = \sqrt{g_w^2 + g_Y^2}$	$W_\mu^\pm,$ $Z_\mu^0$	$W^\pm,$ $Z^0$ Boson
$U_Y(1)$	$g_Y$	$B_\mu$	$e = g_Y \cdot c_w$	$A_\mu$	Photon

**Table 2.4:** The gauge fields and their coupling constants before and after electro weak symmetry breaking (EWSB). The Gluon field is not affected by EWSB.  $a = 1, \dots, 8$  and  $b = 1, 2, 3$  label the number of gauge fields.  $c_w$  is the cosine of the electroweak mixing angle defined in 2.3

The force particles are described by gauge fields. The gauge field of  $SU_C(3)$  is the gluon field. Because the  $SU_C(3)$  has 8 generators there are 8 gluons. Their coupling constant is denoted with  $g_s$ .

For the other force particles in the Standard Model - the  $W^\pm$  bosons, the  $Z_0$  boson and the

photon the situation is slightly more involved. They are obtained as a mixture of the  $W_\mu^b$  ( $b = 1, 2, 3$ ) and the  $B_\mu$  field which are the gauge fields of  $SU_L(2)$  and  $U_Y(1)$  respectively. This mixing procedure is explained in section 2.3.

For the moment being the coupling constants and gauge fields before and after this mixing are quoted in table 2.4.

The Lagrangian of the SM is built up by qualitatively different terms. Firstly there are the kinetic and minimal coupling terms of the matter fields

$$\mathcal{L}_{\text{matter}} = \sum_{i=1}^3 (\bar{\ell}_{iL} i \not{D} l_{iL} + \bar{e}_{iR} i \not{D} e_{iR} + \bar{q}_{iL} i \not{D} q_{iL} + \bar{u}_{iR} i \not{D} u_{iR} + \bar{d}_{iR} i \not{D} d_{iR}) \quad (2.11)$$

where  $i \in \{1, 2, 3\}$  labels the generations of matter. The gauge covariant derivative is given by

$$D_\mu = \partial_\mu + ig_Y \frac{\hat{Y}}{2} + ig_w \vec{\tau} \cdot \vec{W}^\mu + ig_s T^a G_a^\mu \quad (2.12)$$

where for each field the corresponding representation (fundamental or trivial) of the gauge group is to be inserted (see table 2.3). The hyper charge operator  $\hat{Y}$  gives the eigenvalue of the hypercharge of the field it is applied to. These can also be found in table 2.3. The kinetic terms of the gauge fields are given by

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4} F^{\mu\nu} F_{\mu\nu} - \frac{1}{4} W^{a\mu\nu} W_{\mu\nu}^a - \frac{1}{4} G^{a\mu\nu} G_{\mu\nu}^a. \quad (2.13)$$

## 2.3 Electroweak Symmetry Breaking

So far no mass terms like in the Dirac Lagrangian 2.2 have been introduced. The reason for this is that they are not gauge invariant for left- and right-handed spinors transform differently. The same argument forbids terms like  $-\frac{m^2}{2} A^\mu A_\mu$  for a generic gauge boson. Electroweak symmetry breaking (EWSB) ascribes masses to those particles [?, 4–9]. To this end one considers a complex scalar doublet

$$\Phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} \quad (2.14)$$

which receives a vacuum expectation value (VEV)  $\langle \Phi \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v \end{pmatrix}$  by the Higgs potential

$$V(\Phi^\dagger \Phi) = -\mu^2 \Phi^\dagger \Phi + \lambda (\Phi^\dagger \Phi)^2 \quad (2.15)$$

where  $\mu^2, \lambda > 0$ . The Higgs sector of the Standard Model reads

$$\mathcal{L}_{\text{Higgs}} = (D_\mu \Phi)^\dagger (D^\mu \Phi) - V(\Phi^\dagger \Phi). \quad (2.16)$$

The Higgs doublet couples to the gauge fields of  $SU_L(2)$  and  $U_Y(1)$  in the fundamental representation. Inserting an expansion<sup>5</sup> around the VEV  $\Phi = \begin{pmatrix} \phi^+(x) \\ \frac{1}{\sqrt{2}}(v + H(x) + i\sigma(x)) \end{pmatrix}$  one obtains quadratic terms, i.e. mass terms, for the gauge fields in question. In order to obtain mass eigenstates out of  $B_\mu$  and  $W_\mu^3$  and charge eigenstates for  $Q$  and  $I_3$  out of  $W_\mu^1$  and  $W_\mu^2$  one performs the transformation

$$\begin{pmatrix} A_\mu \\ Z_\mu^0 \end{pmatrix} = \begin{pmatrix} \cos \theta_w & \sin \theta_w \\ -\sin \theta_w & \cos \theta_w \end{pmatrix} \begin{pmatrix} B_\mu \\ W_\mu^3 \end{pmatrix} \quad W_\mu^\pm = \frac{1}{\sqrt{2}}(W_\mu^1 \mp iW_\mu^2), \quad (2.17)$$

where the electroweak mixing angle is given by  $\cos \theta_w = \frac{g_w}{\sqrt{g_w^2 + g_Y^2}}$ . These gauge fields acquire masses:

$$m_W = \frac{g_w}{2}v \quad m_Z = \frac{\sqrt{g_w^2 + g_Y^2}}{2}v \quad m_A = 0. \quad (2.18)$$

Apart from the massive bosons  $W_\mu^\pm$  and  $Z_\mu^0$  one obtains the massless photon  $A_\mu$ . As the photon is massless it is still associated with a gauge symmetry called  $U_{em}(1)$ . One therefore often writes EWSB as the breaking of the gauge group  $SU_L(2) \times U_Y(1)$  to  $U_{em}(1)$ .

Matter particles acquire mass via Yukawa couplings to the Higgs doublet. For up-type-quarks one uses that the charge conjugate of  $\Phi$ :  $\Phi^C = i\sigma^2 \Phi^*$  also transform as  $\Phi$ .

$$\mathcal{L}_{\text{Yukawa}} = - \sum_{i,j=1}^3 (y_{ij}^e \bar{\ell}_L \Phi e_R + y_{ij}^d \bar{q}_L \Phi d_R + y_{ij}^u \bar{q}_L \Phi^C u_R) + h.c. \quad (2.19)$$

where  $y^e, y^d, y^u$  are  $3 \times 3$  matrices in generation space. The fermion mass matrices are therefore:

$$m_{ij}^e = \frac{y_{ij}^e}{\sqrt{2}}v \quad m_{ij}^d = \frac{y_{ij}^d}{\sqrt{2}}v \quad m_{ij}^u = \frac{y_{ij}^u}{\sqrt{2}}v. \quad (2.20)$$

The quark mass matrices are not diagonal, which is the precondition of the violation of CP-invariance.<sup>6</sup> One therefore has to distinguish between interaction and mass eigenstates of the quarks. The corresponding transformation matrix is the well known CKM-matrix [11, 12]. Throughout this thesis, the CKM-matrix is approximated with the unit matrix as it will have

<sup>5</sup>The complex  $\phi^+(x)$  and the real  $\sigma(x)$  are the fields of the so called massless Goldstone bosons. These degrees of freedom can be absorbed in the longitudinally polarized degrees of freedom of the arising gauge bosons  $W^\pm$  and  $Z^0$ . This is referred to as unitary gauge [10]. The real  $H(x)$  is the Higgs field, whose excitation is the Higgs boson.

<sup>6</sup>The actual CP-violating term is the coupling term from quarks to the  $W$ -bosons.

a minor influence upon the results

The upshot of EWSB are masses for all matter particles except for the neutrinos and masses for the gauge bosons  $W^\pm$  and  $Z^0$ .

## 2.4 Quantization

The Quantization of Spin 0 and Spin  $\frac{1}{2}$  fields yield no complication in the Lagrangian formalism. To quantize Spin 1 fields it turns out that the usual gauge invariance needs to be replaced by the so called BRST invariance [13–15]. This results in two extra contributions in the Lagrangian. Firstly there are the gauge fixing terms:

$$\mathcal{L}_{R_\xi} = -\frac{1}{2\xi_A}(\partial^\mu A_\mu)^2 - \frac{1}{\xi_W}|\partial^\mu W_\mu^+ - im_W\xi_W\phi^+|^2 - \frac{1}{2\xi_Z}(\partial^\mu Z_\mu - m_Z\xi_Z\sigma)^2 - \frac{1}{2\xi_G}(\partial^\mu G_\mu^a)^2. \quad (2.21)$$

Here  $R_\xi$ -gauge is chosen, where the parameters  $\xi_i$  specify the gauge further. The two terms in the middle are modified with the Goldstone bosons from section 2.3. This is to cancel terms of the form  $V_\mu\partial^\mu\phi$  up to a total derivative arising from EWSB which would lead to non-diagonal propagators, where  $V$  stands for a gauge boson  $W^\pm$  or  $Z^0$  and  $\phi$  for a Goldstone boson.

Secondly there is a ghost Lagrangian

$$\mathcal{L}_{\text{ghost}} = -\bar{c}_a\partial^\mu(\partial_\mu c_a + g_s f_{abc}c_b G_{c\mu}) + \mathcal{L}_{\text{weak ghosts}}. \quad (2.22)$$

The ghost Lagrangian corresponding to the electroweak sector is not needed within this thesis and due to its lengthy form not quoted here. It can be found in [10]. In eq. 2.22  $c_a$  and  $\bar{c}_a$  are the Faddeev-Popov ghost and antighosts. These fields do not correspond to physical particles because they violate the spin–statistics theorem, i.e. they anticommute while being spin 0 fields. Ghost fields are an elegant way of accounting for an additional term in the Lagrangian of non-abelian gauge fields which is best seen in the path integral quantization [2].

## 2.5 Lagrangian of the Standard Model

The complete Lagrangian of the Standard Model reads

$$\mathcal{L}_{\text{SM}} = \mathcal{L}_{\text{matter}} + \mathcal{L}_{\text{gauge}} + \mathcal{L}_{\text{Higgs}} + \mathcal{L}_{\text{Yukawa}} + \mathcal{L}_{R_\xi} + \mathcal{L}_{\text{ghost}} \quad (2.23)$$

with the corresponding parts of the previous chapters. For further reading and the electroweak Lagrangian given completely in terms of physical fields, see [10].

### 3 The Minimal Supersymmetric Standard Model

Even though the Standard Model is a very well working machinery which has partially been tested to an impressive accuracy it is not the ultimate theory of nature. This is not only because of its incompleteness regarding gravity but also because of its missing incorporation of dark matter. This is that astrophysical observations [?] indicate that only 4.9% of the universe's mass energy consist of ordinary matter while the majority is represented by 26.8% of dark matter and 68.3% of dark energy. Another puzzle is the hierarchy problem which questions the large discrepancy between the Higgs boson mass and its quantum correction, see [?]. Finally the gauge groups of the Standard Model cannot be unified to a theoretically appealing single gauge group at some high energy scale for the three gauge couplings do not meet at any renormalization scale [?]. There are many other problems which are not mentioned here.

Supersymmetry provides a possible extension of the Standard Model. It constitutes a mathematically very aesthetic solution of the above mentioned problems<sup>7</sup>. This section explains briefly the mechanism of supersymmetry and the minimal supersymmetric extension of the Standard Model (MSSM).

#### 3.1 Supersymmetry as an Extension of the Poincaré Symmetry

As already seen in the previous chapter symmetries are of vital importance when studying particle physics. Coleman and Mandula [?] found that for any reasonable quantum field theory space-time and internal symmetries, i.e. gauge symmetries "cannot be combined in any but the trivial way" [?]. Ergo the symmetry groups of the Standard Model are the most general ones.

This statement was corrected by the Haag-Lopuszanski-Sohnius theorem [?] which weakened the condition of the quantum field theory by allowing also for anticommuting symmetry generators. This is by introducing the generators  $Q_\alpha$  and  $\bar{Q}_{\dot{\alpha}}$  with  $\alpha, \dot{\alpha} \in \{1, 2\}$  the Poincare-algebra eq. 2.1 can be nontrivially extended to the so called superalgebra (see Appendix 10.3 for the definition of the  $\sigma$ -matrices):

$$\begin{aligned} \{Q_\alpha, Q_\beta\} &= \{\bar{Q}_{\dot{\alpha}}, \bar{Q}_{\dot{\beta}}\} = 0, \\ \{Q_\alpha, \bar{Q}_{\dot{\alpha}}\} &= 2\sigma_{\alpha\dot{\alpha}}^\mu P_\mu, \\ [P^\mu, Q_\alpha] &= [P^\mu, \bar{Q}_{\dot{\alpha}}] = 0, \\ [Q_\alpha, J^{\mu\nu}] &= \frac{1}{2}(\sigma^{\mu\nu})_\alpha{}^\beta Q_\beta. \end{aligned} \tag{3.1}$$

As  $J^{\mu\nu}$  is the generalized angular momentum operator the last line in eq. 3.1 implies that  $Q_\alpha$  is a spin  $\frac{1}{2}$  operator. Applying  $Q_\alpha$  to a field will change the spin of it by  $\frac{1}{2}$ . This means

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<sup>7</sup>In its minimal extension to the Standard Model supersymmetry does not incorporate gravity but many quantum gravity models are supersymmetric

supersymmetry generators convert fermions to bosons and vice versa.

The superalgebra in eq. 3.1 is actually the smallest extension of the Poincare-algebra as it is possible to introduce further anticommuting symmetry generators. Their maximal number is fixed to  $N = 8$ . The MSSM deals with only one set of anticommuting symmetry generators and is therefore referred to be a  $N = 1$  supersymmetric theory.

A representation in the form of differential operators can be given by introducing superspace. Superspace is a manifold obtained by enlarging Minkowski space, whose coordinates are label with  $x^\mu$ , with four anticommuting numbers:  $\theta^\alpha$  and  $\bar{\theta}^{\dot{\alpha}}$ , where  $\alpha, \dot{\alpha} \in \{1, 2\}$  (see Appendix 10.4 for the definition of anticommuting numbers). A representation of the superalgebra reads

$$\begin{aligned} P^\mu &= i\partial^\mu \\ J^{\mu\nu} &= i(x^\mu\partial^\nu - x^\nu\partial^\mu) \\ Q_\alpha &= i(\partial_\alpha + i\sigma_{\alpha\dot{\alpha}}^\mu\bar{\theta}^{\dot{\alpha}}\partial_\mu) \\ \bar{Q}_{\dot{\alpha}} &= i(-\bar{\partial}_{\dot{\alpha}} - i\theta^\alpha\sigma_{\alpha\dot{\alpha}}^\mu\partial_\mu). \end{aligned} \quad (3.2)$$

### 3.2 A Generic Supersymmetric Model in Superspace Formulation

This chapter outlines the generic ingredients and terms of a supersymmetric model. To this end it is practical to work in the language of superspace and superfields. Superspace has just been introduced and superfields are functions on superspace.

The superfields<sup>8</sup> relevant for the MSSM are the chiral superfield  $\hat{\Phi}$ , the antichiral superfield  $\hat{\bar{\Phi}}$  and the vector superfield  $V$ . Chiral superfields are defined by the restriction  $\bar{\mathcal{D}}_{\dot{\alpha}}\hat{\Phi} = 0$ , antichiral superfields by  $\mathcal{D}_\alpha\hat{\bar{\Phi}} = 0$  (see Appendix 10.4 for the definition of chiral covariant derivatives) and vector superfields by the condition of being real  $V^\dagger = V$ .

All superfields can be decomposed into components. The component decomposition of the above fields read<sup>9</sup>

$$\begin{aligned} \hat{\Phi}(x, \theta, \bar{\theta}) &= A(x) + \sqrt{2}\theta\psi(x) + \theta\theta F(x) - i\theta\sigma^\mu\bar{\theta}\partial_\mu A(x) - \frac{1}{4}\theta\theta\bar{\theta}\bar{\theta}\partial_\mu\partial^\mu A(x) - \frac{i}{\sqrt{2}}\theta\theta\bar{\theta}\sigma^\mu\partial_\mu\psi(x) \\ \hat{\bar{\Phi}}(x, \theta, \bar{\theta}) &= A^\dagger(x) + \sqrt{2}\bar{\theta}\bar{\psi}(x) + \bar{\theta}\bar{\theta}F^\dagger(x) + i\theta\sigma^\mu\bar{\theta}\partial_\mu A^\dagger(x) - \frac{1}{4}\theta\theta\bar{\theta}\bar{\theta}\partial_\mu\partial^\mu A^\dagger(x) - \frac{i}{\sqrt{2}}\bar{\theta}\bar{\theta}\theta\sigma^\mu\partial_\mu\bar{\psi}(x) \\ \hat{V}(x, \theta, \bar{\theta}) &= \theta\sigma^\mu\bar{\theta}v_\mu + i\theta\theta\bar{\theta}\bar{\lambda}(x) - i\bar{\theta}\bar{\theta}\theta\lambda(x) + \frac{1}{2}\theta\theta\bar{\theta}\bar{\theta}D(x), \end{aligned} \quad (3.3)$$

where  $A(x)$  and  $F(x)$  are complex scalar fields,  $\psi(x)$  and  $\lambda(x)$  are left handed Weyl spinors and  $D(x)$  being a real scalar field.

<sup>8</sup>Superfield are throughout this thesis labeled with a hat.

<sup>9</sup>For the vector superfield Wess-Zumino-gauge is applied.



The superfields transform under a generic gauge transformation as

$$\begin{aligned}\hat{\Phi} &\rightarrow e^{-2ig\hat{\Lambda}}\hat{\Phi} \\ \hat{\bar{\Phi}} &\rightarrow \hat{\bar{\Phi}}e^{2ig\hat{\Lambda}} \\ e^{2g\hat{V}} &\rightarrow e^{-2ig\hat{\Lambda}}e^{2g\hat{V}}e^{2ig\hat{\Lambda}},\end{aligned}\tag{3.4}$$

where  $\hat{\Lambda} = \hat{\Lambda}^a T^a$  and  $\hat{V} = \hat{V}^a T^a$ .  $\hat{\Lambda}^a$  is an arbitrary chiral superfield and the  $T^a$  are the generators of the gauge group in question.  $g$  is the gauge coupling constant of the gauge group.

One can therefore construct the important gauge invariant term  $\int d^4\theta \hat{\bar{\Phi}}e^{2g\hat{V}}\hat{\Phi}$ . If one introduces the gauge covariant derivative  $D_\mu = \partial + igT^a v_\mu^a$  its component decomposition reads

$$\begin{aligned}\int d^4\theta \hat{\bar{\Phi}}e^{2g\hat{V}}\hat{\Phi} &= F^\dagger F + (D_\mu A)^\dagger (D^\mu A) + \bar{\psi}\bar{\sigma}^\mu i D_\mu \psi \\ &\quad - \sqrt{2}g \left( -i(A^\dagger T^a A)\lambda^a + i\bar{\lambda}^a (AT^a A^\dagger) \right) + g(A^\dagger T^a A)D^a.\end{aligned}\tag{3.5}$$

Therefore this term gives rise to the kinetic terms of the components of the chiral and antichiral superfields  $A$ ,  $A^\dagger$ ,  $\psi$  and  $\bar{\psi}$  as well as their minimal coupling to the gauge fields  $v_\mu^a$  and their superpartners  $\lambda^a$  and  $\bar{\lambda}^a$  and terms involving the auxiliary fields  $F$ ,  $F^\dagger$  and  $D$ .

With the field-strength chiral superfields  $\hat{W}_\alpha := -\frac{1}{4}\overline{\mathcal{D}}\mathcal{D}(e^{-2gV}\mathcal{D}_\alpha e^{2gV})$  one can write down a gauge invariant term yielding the kinetic terms of the gauge fields and their superpartners:

$$\int d^2\theta \frac{1}{16g^2} \hat{W}^{\alpha a} \hat{W}_\alpha^a + h.c. = \frac{1}{2} D^a D^a - \frac{1}{4} F_{\mu\nu}^a F^{a\mu\nu} + \frac{i}{2} \bar{\lambda}^a \bar{\sigma}^\mu (D_\mu \lambda^a) + \frac{i}{2} \lambda^a \sigma^\mu (D_\mu \bar{\lambda}^a).\tag{3.6}$$

A third generic term in a supersymmetric theory arises from the superpotential  $W(\hat{\Phi})$  which is a holomorphic function in the chiral superfields:

$$\int d^2\theta W(\hat{\Phi}).\tag{3.7}$$

A renormalizable superpotential is given by  $W(\hat{\Phi}) = c_i \hat{\Phi} + \frac{m_{ij}}{2} \hat{\Phi}_i \hat{\Phi}_j + \frac{g_{ijk}}{3!} \hat{\Phi}_i \hat{\Phi}_j \hat{\Phi}_k$ . The component decomposition of the corresponding terms is

$$\begin{aligned}\int d^2\theta \hat{\Phi}_1 &= F_1 \\ \int d^2\theta \hat{\Phi}_1 \hat{\Phi}_2 &= A_1 F_2 + F_1 A_2 - \psi_1 \psi_2 \\ \int d^2\theta \hat{\Phi}_1 \hat{\Phi}_2 \hat{\Phi}_3 &= F_1 A_2 A_3 + A_1 F_2 A_3 + A_1 A_2 F_3 - A_1 \psi_2 \psi_3 - \psi_1 A_2 \psi_3 - \psi_1 \psi_2 A_3.\end{aligned}\tag{3.8}$$

The Lagrangian for a supersymmetric theory is therefore given by

$$\begin{aligned}\mathcal{L}_{\text{SUSY}} &= \mathcal{L}_{\text{matter}} + \mathcal{L}_{\text{gauge}} + \mathcal{L}_{\text{superpot}} \\ &= \int d^4\theta \hat{\Phi} e^{2g\hat{V}} \hat{\Phi} + \left( \int d^2\theta \frac{1}{16g^2} \hat{W}^{\alpha a} \hat{W}_\alpha^a + h.c. \right) + \int d^2\theta W(\hat{\Phi})\end{aligned}\quad (3.9)$$

Observing the component decomposition eq.3.5, 3.6 and 3.8 of the 3 parts of this Lagrangian one observes that the  $F$  and  $D$  fields have no kinetic term and are therefore auxiliary fields which can be eliminated by their equation of motion  $\frac{\partial \mathcal{L}}{\partial \phi} = \partial_\mu \frac{\partial \mathcal{L}}{\partial (\partial_\mu \phi)}$  with  $\phi \in \{F, D\}$ . Doing this one obtains

$$\begin{aligned}\mathcal{L}_D &= \frac{1}{2} D^a D^a + g A^\dagger T^a D^a A \quad \Rightarrow \quad D^a = -A^\dagger T^a A \\ \mathcal{L}_D &= -\frac{1}{2} (A^\dagger T^a A)^2.\end{aligned}\quad (3.10)$$

and

$$\begin{aligned}\mathcal{L}_F &= F_i^\dagger F_i + \left( c_i F_i + m_{ij} F_i A_j + \frac{g_{ijk}}{2} F_i A_j A_k + h.c. \right) \quad \Rightarrow \quad F_i^\dagger = -\frac{\partial W(A)}{\partial A_i} \\ \mathcal{L}_F &= -\left| \frac{\partial W(A)}{\partial A_i} \right|^2\end{aligned}\quad (3.11)$$

### 3.3 The Minimal Supersymmetric Standard Model

After having introduced a generic supersymmetric model a possible realization of supersymmetry in nature - the Minimal Supersymmetric Standardmodel (MSSM) - is discussed in this subsection. The field content of the MSSM is summarized in table ???. In comparison to the Standard Model each particle has a superpartner which differs in its spin by  $\frac{1}{2}$ . The superpartners of fermions have Spin 0 and are referred to as sfermions. The name of each sfermion is given by the name of its Standard Model partner with an additional "s" in front of it, e.g. selectron or up-squark. The superpartners of bosons have spin  $\frac{1}{2}$  and get an additional "ino" at the end of their name, e.g. Higgsino as the superpartner of the Higgs or Zino as superpartner of the  $Z$ -boson.

Apart from this doubling up of the field content of the Standard Model there is one modification within the Higgs sector, i.e. there are two instead of one Higgs superfield. These are necessary to avoid an anomaly in the electroweak gauge symmetry and to give masses to both up- and down-type quarks [?, page 8].

Note that the "chirality" of sfermions is understood to be the chirality of their superpartners. Of course scalar particles have no handedness. The Lagrangian<sup>10</sup> for the MSSM reads

Superfield	Components	$SU_C(3) \times SU_L(2) \times U_Y(1)$
$\hat{\Phi}$	$A, \psi$	
$\hat{V}$	$\lambda, v_\mu$	
$\hat{Q}$	$\tilde{q}_L = \begin{pmatrix} \tilde{u}_L \\ \tilde{d}_L \end{pmatrix}, q_L = \begin{pmatrix} u_L \\ d_L \end{pmatrix}$	$(3, 2, \frac{1}{6})$
$\hat{U}$	$\tilde{u}_R^\dagger, u_R$	$(3^*, 1, -\frac{2}{3})$
$\hat{D}$	$\tilde{d}_R^\dagger, d_R$	$(3^*, 1, +\frac{1}{3})$
$\hat{L}$	$\tilde{l}_L = \begin{pmatrix} \tilde{\nu}_L \\ \tilde{e}_L \end{pmatrix}, l_L = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix}$	$(1, 2, -\frac{1}{2})$
$\hat{E}$	$\tilde{e}_R^\dagger, e_R$	$(1, 1, 1)$
$\hat{H}_d$	$H_d, \tilde{H}_d$	$(1, 2, -\frac{1}{2})$
$\hat{H}_u$	$H_u, \tilde{H}_u$	$(1, 2, +\frac{1}{2})$
$\hat{V}_Y$	$\lambda_Y, B_\mu$	$(1, 1, 0)$
$\hat{V}_w^a$	$\lambda_w^a, W_\mu^a$	$(1, 3, 0)$
$\hat{V}_s^a$	$\lambda_s^a, G_\mu^a$	$(8, 1, 0)$

**Table 3.1:** The table shows the field content of the MSSM in terms of the superfields and their component decomposition. The first two lines show the decomposition of the generic superfields (cf. subsection 3.3).

The third column shows the representation (for  $SU_C(3)$  and  $SU_L(2)$ ) in which the fields transform and the charges of the fields for  $U_Y(1)$ .

$$\begin{aligned}
\mathcal{L}_{\text{MSSM}} = & \int d^4\theta \left[ \hat{\bar{Q}} e^{2g_Y \hat{V}_Y + 2g_w \hat{V}_w + 2g_s \hat{V}_s} \hat{Q} + \hat{\bar{U}} e^{2g_Y \hat{V}_Y + 2g_w \hat{V}_w - 2g_s \hat{V}_s^T} \hat{U} + \hat{\bar{D}} e^{2g_Y \hat{V}_Y + 2g_w \hat{V}_w - 2g_s \hat{V}_s^T} \hat{D} \right. \\
& + \hat{\bar{L}} e^{2g_Y \hat{V}_Y + 2g_w \hat{V}_w} \hat{L} + \hat{\bar{E}} e^{2g_Y \hat{V}_Y + 2g_w \hat{V}_w} \hat{E} \\
& \left. + \hat{\bar{H}}_d e^{2g_Y \hat{V}_Y + 2g_w \hat{V}_w} \hat{H}_d + \hat{\bar{H}}_u e^{2g_Y \hat{V}_Y + 2g_w \hat{V}_w} \hat{H}_u \right] \\
& + \int d^2\theta \left[ \frac{1}{16g_Y^2} \hat{W}_Y^\alpha \hat{W}_{Y\alpha} + \frac{1}{16g_w^2} \hat{W}_w^{a\alpha} \hat{W}_{w\alpha}^a + \frac{1}{16g_s^2} \hat{W}_s^{a\alpha} \hat{W}_{s\alpha}^a \right] + h.c. \\
& + \int d^2\theta W_{\text{MSSM}} + h.c. \\
& + \mathcal{L}_{\text{soft}}.
\end{aligned} \tag{3.12}$$

Apart from the kinetic and minimal coupling terms in the first four lines of 3.12 there is a superpotential  $W_{\text{MSSM}}$ :

$$W_{\text{MSSM}} = y_d \hat{H}_d \hat{Q} \hat{D} + y_u \hat{H}_u \hat{Q} \hat{U} + y_e \hat{H}_d \hat{L} \hat{E} - \mu \hat{H}_d \hat{H}_u \tag{3.13}$$

After eliminating auxiliary fields a Higgs potential is generated. In contrast to the Standard Model the quadrilinear term in the Higgs potential is given in terms of the gauge couplings and therefore no independent parameter. From the Higgs potential both  $H_u$  and  $H_d$  acquire

<sup>10</sup>This is the Lagrangian on the classical level, i.e. there are neither gauge fixing nor ghost terms.

a vacuum expectation value:

$$\langle H_u \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_u \end{pmatrix} \quad \langle H_d \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} v_d \\ 0 \end{pmatrix} \quad (3.14)$$

whose quotient is an important parameter of the electroweak MSSM.

$$\tan \beta := \frac{v_u}{v_d} \quad (3.15)$$

As in the Standard Model this spontaneous symmetry breaking leads to Yukawa-terms (first three terms of eq. 3.13) which give masses to the fermions<sup>11</sup>. Apart from that there are also other terms like quark-squark-higgsino interactions. In addition there is the  $\mu$ -term which is responsible for Higgsino and Higgs masses.

In order to explain why supersymmetric particles have not been found so far supersymmetry has to be broken. Like in electroweak symmetry breaking supersymmetry is preserved at some high energy scale but its vacuum state is not supersymmetric. However there is no consensus on how supersymmetry is broken. But by ignoring the exact mechanism and parameterizing it by introducing appropriate terms in the Lagrangian one accounts for it in an absolutely viable way.

Terms which break supersymmetry should be soft, i.e. they should have coupling constants with positive mass dimension.

$$\begin{aligned} \mathcal{L}_{\text{soft}} = & -M_{\tilde{Q}}^2 |\tilde{q}_L|^2 - M_{\tilde{U}}^2 |\tilde{u}_R|^2 - M_{\tilde{D}}^2 |\tilde{d}_R|^2 \\ & - M_{\tilde{L}}^2 |\tilde{l}_L|^2 - M_{\tilde{E}}^2 |\tilde{e}_R|^2 - M_{H_d}^2 |H_d|^2 - M_{H_u}^2 |H_u|^2 \\ & + \frac{1}{2} (M_1 \lambda_Y \lambda_Y + M_2 \lambda_w^a \lambda_w^a + M_3 \lambda_s^a \lambda_s^a) + h.c. \\ & - \left( A_d y_d H_d \tilde{q}_L \tilde{d}_R^\dagger + A_u y_u H_u \tilde{q}_L \tilde{u}_R^\dagger + A_e y_e H_d \tilde{l}_L \tilde{e}_R^\dagger - B \mu H_d H_u \right) + h.c. \end{aligned} \quad (3.16)$$

The softly breaking terms in the MSSM comprise masses of the squarks, sleptons, Higgs-bosons and gauginos. The last line in eq. 3.16 resembles the superpotential in eq. 3.13. In fact the only two differences are that the superfields are replaced by their scalar components and that additional parameters with mass dimension one have been added.  $A_d$ ,  $A_u$ ,  $A_e$  are  $3 \times 3$  matrices in family space.

Many supersymmetric particles have no definite mass but mix to particles with definite mass. This is true for neutral higgsinos and gauginos which mix to neutralinos and charged higgsinos and gauginos which mix to charginos. However, for this thesis only the mass matrix for electrically charged sfermions is given because this will change significantly when introducing

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<sup>11</sup>Note the sum convention of  $SU_L(2)$  doublets:  $\hat{H}_u \hat{Q} = (\hat{H}_u)_\alpha \hat{Q}_\beta \epsilon^{\alpha\beta}$  to ensure gauge invariance.

R-symmetry in the next section. The Lagrangian describing these masses reads

$$\mathcal{L}_{m_{\tilde{q}}, m_{\tilde{e}}} = - \begin{pmatrix} \tilde{u}_L^\dagger & \tilde{u}_R^\dagger \end{pmatrix} M_{\tilde{u}}^2 \begin{pmatrix} \tilde{u}_L \\ \tilde{u}_R \end{pmatrix} - \begin{pmatrix} \tilde{d}_L^\dagger & \tilde{d}_R^\dagger \end{pmatrix} M_{\tilde{d}}^2 \begin{pmatrix} \tilde{d}_L \\ \tilde{d}_R \end{pmatrix} - \begin{pmatrix} \tilde{e}_L^\dagger & \tilde{e}_R^\dagger \end{pmatrix} M_{\tilde{e}}^2 \begin{pmatrix} \tilde{e}_L \\ \tilde{e}_R \end{pmatrix}. \quad (3.17)$$

where the mass matrix is given by

$$M_f^2 = \begin{pmatrix} M_{fLL}^2 & m_f X_f^* \\ m_f X_f & M_{fRR}^2 \end{pmatrix}. \quad (3.18)$$

The diagonal terms are given by the sum of the fermion mass squared  $m_f^2$ , the appropriate soft breaking term, e.g.  $M_Q^2$  for left handed squarks and a contribution from the elimination of  $D$ -fields. In fact this last term splits the masses of fermions and sfermions already without explicit supersymmetry breaking. This is because the generation of fermion masses requires the breaking of gauge symmetry which is not possible without breaking supersymmetry. However the off-diagonal terms are given by

$$X_f = A_f + \mu^* \{\cot \beta, \tan \beta\} \quad (3.19)$$

with  $\cot \beta$  for up-type squarks and  $\tan \beta$  for down-type sfermions

### 3.4 R-Parity

The MSSM does not include any terms which violate baryon (B) or lepton number (L)<sup>12</sup> like

$$W_{\Delta B=1} = \lambda \hat{U} \hat{D} \hat{D} \quad \text{or} \quad W_{\Delta L=1} = \lambda' \hat{L} \hat{Q} \hat{D} \quad (3.20)$$

in the superpotential although this would be possible from the perspective of gauge invariance and renormalizability. However, this has not been observed experimentally. The most famous constraint on them is the non-observation of proton decay which restricts at least one of the parameters  $\lambda$  or  $\lambda'$  to be extremely small [?]. Forbidding these terms by postulating a conservation of B and L might not be the best solution since both B and L are known to be violated in nature by non-perturbative effects [?].

A possible way out is the postulation of the conservation of R-parity

$$P_R = (-1)^{3(B-L)+2s} \quad (3.21)$$

where  $s$  is the spin of the particle on which the R-parity operator is applied to. R-parity is a multiplicative quantum number which is constructed to be +1 for each Standard Model

<sup>12</sup>The chiral superfield  $\hat{Q}$  carries baryon number  $\frac{1}{3}$ , whereas  $B = -\frac{1}{3}$  for  $\hat{D}$ ,  $\hat{U}$  and  $B = 0$  for all other superfields. The only multiplets carrying a non-zero lepton number are  $\hat{L}$  with  $L = 1$  and  $\hat{E}$  with  $L = -1$ .

particle and -1 for each supersymmetric partner. Apart from the the prohibition of the terms in eq. 3.20 this has important consequences:

- There has to be a lightest supersymmetric particle (LSP) which is absolutely stable. This could be a candidate for dark matter.
- At colliders like the LHC there must always be an even number of supersymmetric particles be produced, if at all.

The superpotential of the MSSM (see, eq. 3.13) is chosen in such a way that it is the only renormalizable and gauge invariant potential which respects R-parity.

## 4 R-Symmetry

R-symmetry is an additional continuous symmetry which extends supersymmetry in a non-trivial way. This is actually covered in the Haag-Lopuszanski-Sohnius theorem [?] but often concealed.

This section introduces R-symmetry and discusses briefly a minimal viable extension of the MSSM including R-symmetry, the MRSSM. Thereby some general features of the model are mentioned before the focus is laid on the strongly interacting sector which is the central topic of the thesis.

### 4.1 R-Symmetry Transformation

R-symmetry is a global  $U(1)$  symmetry. It must not be confused with R-parity which is a discrete  $\mathbb{Z}_2$  symmetry. A continuous global symmetry implies according to Noether's theorem a conserved charge. In the case of R-symmetry it is called R-charge and one therefore refers to R-symmetry as  $U_R(1)$ .

The defining property of  $U_R(1)$  is that the anticommuting coordinates  $\theta^\alpha$  and  $\bar{\theta}^{\dot{\alpha}}$  transform like

$$\theta \rightarrow e^{i\alpha}\theta \qquad \bar{\theta} \rightarrow e^{-i\alpha}\bar{\theta}, \quad (4.1)$$

where  $\alpha$  parametrizes the transformation. This in turn implies that the supersymmetry generators (see eq. ??) transform like

$$Q \rightarrow e^{-i\alpha}Q \qquad \bar{Q} \rightarrow e^{i\alpha}\bar{Q} \quad (4.2)$$

and R-symmetry does not commute with supersymmetry

$$[R, Q] = -Q \qquad [R, \bar{Q}] = \bar{Q}, \quad (4.3)$$

meaning that superpartners do not have the same R-charge. The transformation of chiral and vector superfields reads

$$\begin{aligned} \hat{\Phi}(x, \theta, \bar{\theta}) &\rightarrow e^{ir_{\hat{\Phi}}\alpha} \hat{\Phi}(x, e^{i\alpha}\theta, e^{-i\alpha}\bar{\theta}) \\ \hat{V}(x, \theta, \bar{\theta}) &\rightarrow \hat{V}(x, e^{i\alpha}\theta, e^{-i\alpha}\bar{\theta}), \end{aligned} \quad (4.4)$$

i.e. the chiral superfields transform with arbitrary R-charge where the R-charge of vector superfield is restricted to be zero by the condition that they are real. If one inserts the component decomposition eq. 3.3 of the superfields one can read off the R-charges of the component fields.

superfield		boson		fermion	
$\hat{\Phi}$	$r_{\hat{\Phi}}$	$A$	$r_{\hat{\Phi}}$	$\psi$	$r_{\hat{\Phi}} - 1$
$\hat{V}$	0	$v^\mu$	0	$\lambda$	+1

**Table 4.1:** This table shows the R-charges of a generic chiral and vector superfield.

## 4.2 The Minimal R-Symmetric Supersymmetric Standard Model

If one imposes R-symmetry upon the MSSM one is faced with a certain arbitrariness, i.e. the choice of the R-charges of the chiral superfields. A minimal R-symmetric extension of the MSSM is the minimal R-symmetric supersymmetric Standard Model (MRSSM) [?]. In this model the R-charges are chosen in such a way, that every Standard Model particle has R-charge zero. Following this one obtains the R-charges of all particles which are summed up in table 4.2.

Field	Superfield		Boson		Fermion	
Gauge Vector	$\hat{V}_s^a, \hat{V}_w^a, \hat{V}_Y$	0	$G_\mu^a, W_\mu^a, B_\mu$	0	$\lambda_s^a, \lambda_w^a, \lambda_Y$	+1
Matter	$\hat{L}, \hat{E}$	0	$\tilde{l}_l, \tilde{e}_R^\dagger$	+1	$l_L, e_R$	0
	$\hat{Q}, \hat{D}, \hat{U}$	+1	$\tilde{q}_L, \tilde{d}_R^\dagger, \tilde{u}_R^\dagger$	+1	$q_L, d_R, u_R$	0
H-Higgs	$\hat{H}_{d,u}$	0	$H_{d,u}$	0	$\tilde{H}_{d,u}$	-1
R-Higgs	$\hat{R}_{d,u}$	+2	$R_{d,u}$	+2	$\tilde{R}_{d,u}$	+1
Adjoint Chiral	$\hat{O}, \hat{T}, \hat{S}$	0	$\sigma^a, \omega^a, \rho$	0	$\tilde{\chi}_s^a, \tilde{\chi}_w^a, \tilde{\chi}_Y$	-1

**Table 4.2:** This table lists the R-charges of all superfields and their components in the MRSSM. The fields of the R-Higgs and the adjoint chiral superfields are not present in the MSSM

The gauge, matter and H-Higgs fields are the fields of the MSSM. Below the horizontal line one finds the fields which are not present in the MSSM, i.e. the R-Higgs and adjoint chiral fields. These occur for the following reason.

In the MSSM the gauginos are Majorana particles. Their mass terms reads in 2-spinor notation

$$\mathcal{L}_{\text{Majorana mass}} = -m\lambda\lambda + h.c. \quad (4.5)$$

which is not R-invariant because the Weyl fermion  $\lambda$  has R-charge +1. Giving no mass to the gauginos is phenomenologically not possible. The only way to account for a fermion mass is to write down a Dirac mass term.

$$\mathcal{L}_{\text{Dirac mass}} = -m\chi\lambda + h.c. \quad (4.6)$$

This means that the 4-spinor of every gaugino is a Dirac- instead of a Majorana-spinor, e.g.



for the gluino:

$$\tilde{g}_{\text{MSSM}} = \begin{pmatrix} -i\lambda_s \\ +i\overline{\lambda}_s \end{pmatrix} \quad \tilde{g}_{\text{MRSSM}} = \begin{pmatrix} -i\lambda_s \\ +i\overline{\chi}_s \end{pmatrix}. \quad (4.7)$$

In order to get a *R*-symmetric mass term one has to choose the *R*-charge of the new Weyl-spinor  $\chi$  to be the opposite of  $\lambda$ .

This explains the necessity of enlarging the field content if one imposes *R*-symmetry.

Of course the new Weyl-spinor  $\chi$  must have also a superpartner. One chooses this superpartner to be a scalar, i.e. the additional Weyl fermion comes from a chiral superfield. In order to maintain gauge invariance this chiral superfield has to transform in the adjoint representation, hence the name adjoint chiral in table 4.2. To fix notation the component decomposition of the 8 chiral supermultiplets associated to the gluons is given by

$$\hat{O}^a(x, \theta, \bar{\theta}) = \sigma^a + \sqrt{2}\theta i\chi_s^a + \dots \quad a = 1, \dots, 8. \quad (4.8)$$

The scalar components  $\sigma^a$  are referred to as scalar gluons and the Weyl spinors  $\chi^a$  are called octinos.

The same argument as for the adjoint chiral explains the existence of additional Higgs-superfields which are referred to as *R*-Higgs fields.

Apart from these extensions *R*-symmetry also forbids terms which are allowed by supersymmetry. For the above choice of *R*-charges the  $\mu$ -term in 3.13 and the *A*-terms in the last line of 3.16 are absent in the MRSSM. Inspecting eq. 3.19 implies that there is no squark mixing in the MRSSM. As a consequence terms which allow flavor violating processes like  $\mu \rightarrow e\gamma$  are forbidden in the MRSSM [?]. Quite interestingly Dirac gauginos can be heavier than Majorana gauginos without being less natural in view of the hierarchy problem [?, ?, ?]. Finally the production cross section of strongly interacting particles is suppressed [?, ?]. It is this last feature which is going to be observed further in this thesis. To this end the colored sector of the MRSSM is now considered in more detail.

### 4.3 The *R*-Symmetric Supersymmetric Quantum Chromodynamics

The subject of this thesis is the phenomenology of the strongly coupling sector of the MRSSM. The *R*-symmetric supersymmetric quantum chromodynamics (RSQCD) is therefore considered closer. Its Lagrangian for an arbitrary massless quark reads

$$\begin{aligned} \mathcal{L}_{\text{RSQCD}} = & \int d^4\theta \left( \hat{\bar{Q}}_L e^{2g_s \hat{V}_s} \hat{Q}_L + \hat{\bar{Q}}_R e^{-2g_s \hat{V}_s^T} \hat{Q}_R + \hat{\bar{O}} e^{2g_s \hat{V}_s^{\text{fund}}} \hat{O} \right) \\ & + \left( \int d^2\theta \frac{1}{16g_s^2} \hat{W}_s^{a\alpha} \hat{W}_{s\alpha}^a + h.c. \right) + \mathcal{L}_{\text{soft}} \end{aligned} \quad (4.9)$$

where  $\hat{Q}_L$  and  $\hat{Q}_R$  refer to the left and right handed superfield of the quark. Note that each field in the first line of eq. 4.9 transforms in a different representation of  $SU_C(3)$ , i.e. the fundamental, the antifundamental and the adjoint one. In terms of components fields these expressions are given by

$$\int d^4\theta \hat{\bar{Q}}_L e^{2g_s \hat{V}_s} \hat{Q}_L = F_L^\dagger F_L + (D_\mu \tilde{q}_L)^\dagger (D^\mu \tilde{q}_L) + \bar{q}_L \bar{\sigma}^\mu i D_\mu q_L \\ - \sqrt{2} g_s \left( -i(\tilde{q}_L^\dagger T^a q_L) \lambda^a + i \bar{\lambda}^a (\bar{q}_L T^a \tilde{q}_L) \right) + g_s \tilde{q}_L^\dagger T^a D^a \tilde{q}_L \quad (4.10)$$

$$\int d^4\theta \hat{\bar{Q}}_R e^{-2g_s \hat{V}_s^T} \hat{Q}_R = F_R^\dagger F_R + (D_\mu \tilde{q}_R)^\dagger (D^\mu \tilde{q}_R) + \bar{q}_R \bar{\sigma}^\mu i D_\mu q_R \\ + \sqrt{2} g_s \left( -i(\tilde{q}_R T^{*a} q_R) \lambda^a + i \bar{\lambda}^a (\bar{q}_R T^{*a} \tilde{q}_R^\dagger) \right) - g_s \tilde{q}_R T^{*a} D^a \tilde{q}_R^\dagger \quad (4.11)$$

$$\int d^4\theta \hat{\bar{O}} e^{2g_s \hat{V}_s^{\text{fund}}} \hat{O} = F_O^\dagger F_O + (D_\mu \sigma^a)^\dagger (D^\mu \sigma^a) + \bar{\chi} \bar{\sigma}^\mu i D_\mu \chi \\ - \sqrt{2} g_s \left( -i(\sigma_b^\dagger (-i f_{abc}) (-i \chi^c)) \lambda^a + i \bar{\lambda}^a (i \bar{\chi}_b (+i f_{abc}) \sigma^{c\dagger}) \right) \\ - i g_s \sigma^{b\dagger} f^{abc} D^a \sigma^c \quad (4.12)$$

where in the gauge covariant derivative  $D_\mu = \partial_\mu + i g_s T^a G_\mu^a$  the generator  $T^a$  needs to be replaced by  $-T^{*a}$  or  $-i f^{abc}$  if applied to a field transforming in the antifundamental or adjoint representation respectively.

The soft breaking Lagrangian accounts for the squark, gaugino and scalar gluon masses. These mass terms arise from a hidden sector spurion. For gauginos the D-type spurion is given by  $\hat{W}'_\alpha = \theta_\alpha D$  and mediates super symmetry breaking at the mediation scale  $M$ :  $\int d\theta^2 \frac{\hat{W}'_\alpha}{M} W_s^\alpha \hat{O}$ . After integrating out the spurion one obtains [?, ?]

$$\mathcal{L}_{\text{soft}} = -\frac{m_{\tilde{q}}^2}{2} (|\tilde{q}_L|^2 + |\tilde{q}_R|^2) \\ - m_\sigma^2 |\sigma^a|^2 - m_g (\lambda \chi - \sqrt{2} D^a \sigma^a + h.c.) \quad (4.13)$$

where the complex scalar gluons which are also referred to as sgluons  $\sigma = \frac{\phi_0 + i\sigma_0}{\sqrt{2}}$  constitute of two real ones with different masses. The equations of motion for the auxiliary fields are

$$D^a = -g_s \tilde{q}_L^\dagger T^a \tilde{q}_L + g_s \tilde{q}_R T^{*a} \tilde{q}_R^\dagger + i g_s \sigma^{\dagger b} f^{abc} \sigma^c - \sqrt{2} m_g (\sigma^a + \sigma^{\dagger a}) \quad (4.14)$$

$$F_i = 0 \quad \text{for} \quad i = L, R, O \quad (4.15)$$

where  $D^a$  is still real as the purely imaginary parts do not contribute by virtue of the antisymmetry of the structure constants. After eliminating the auxiliary fields the complete

Lagrangian reads in 4 spinor notation<sup>13</sup>

$$\begin{aligned}
\mathcal{L}_{RSQCD} = & |D_\mu \sigma|^2 + |D_\mu \tilde{q}_R|^2 + |D_\mu \tilde{q}_L|^2 + \bar{q} i \not{D} q + \bar{\tilde{g}}^a i \not{D} P_L \tilde{g}^a + \bar{\tilde{g}}^a i \not{D} P_R \tilde{g}^a - \frac{1}{4} (F_a^{\mu\nu})^2 \\
& - \sqrt{2} g_s \left( \bar{\tilde{g}}^a P_R (q^C T^a \tilde{q}_L) + (\tilde{q}_L^\dagger T^a \bar{q}^C) P_L \tilde{g}^a \right) \\
& + \sqrt{2} g_s \left( \bar{\tilde{g}}^a P_R (q T^{*a} \tilde{q}_R^\dagger) + (\tilde{q}_R T^{*a} \bar{q}) P_L \tilde{g}^a \right) \\
& - \sqrt{2} g_s \left( \bar{\tilde{g}}^a P_R (\tilde{g}^b (i f^{abc}) \sigma^c) + (\sigma^{\dagger b} (-i f^{abc}) \bar{\tilde{g}}^c) P_L \tilde{g}^a \right) \\
& - \frac{m_{\tilde{q}}^2}{2} (|\tilde{q}_L|^2 + |\tilde{q}_R|^2) - m_\sigma^2 |\sigma^a|^2 - m_g \bar{\tilde{g}}^a \tilde{g}^a \\
& - \frac{1}{2} \left( g_s \tilde{q}_L^\dagger T^a \tilde{q}_L - g_s \tilde{q}_R T^{*a} \tilde{q}_R^\dagger - i g_s \sigma^{\dagger b} f^{abc} \sigma^c + \sqrt{2} m_g (\sigma^a + \sigma^{\dagger a}) \right)^2 \quad (4.16)
\end{aligned}$$

Observe that there is no 3 sgluon vertex, because of the antisymmetry of the structure constants  $f^{abc}$ . As already alluded to there are two distinct sgluons with different masses. The mass of the pseudoscalar  $\sigma_0^a$  is given by the soft breaking parameter  $m_{\sigma_0} = m_\sigma$  whereas the mass of the scalar  $\phi^0$  is given by

$$m_{\phi_0} = \sqrt{m_{\sigma_0}^2 + 4m_{\tilde{g}}}. \quad (4.17)$$

Add table of particle content of RSQCD

<sup>13</sup>How a 4 spinor is composed of Weyl-spinors is given in the Appendix 10.3

## 5 Squark and Gluino Production at Tree Level

Comment on sgluon production: which channels, calculate cross section for at least one configuration of masses.

In this chapter the production of strongly interacting supersymmetric particles<sup>14</sup> in the MRSSM is considered. The various processes and their associated cross section is compared to their analogues in the MSSM.

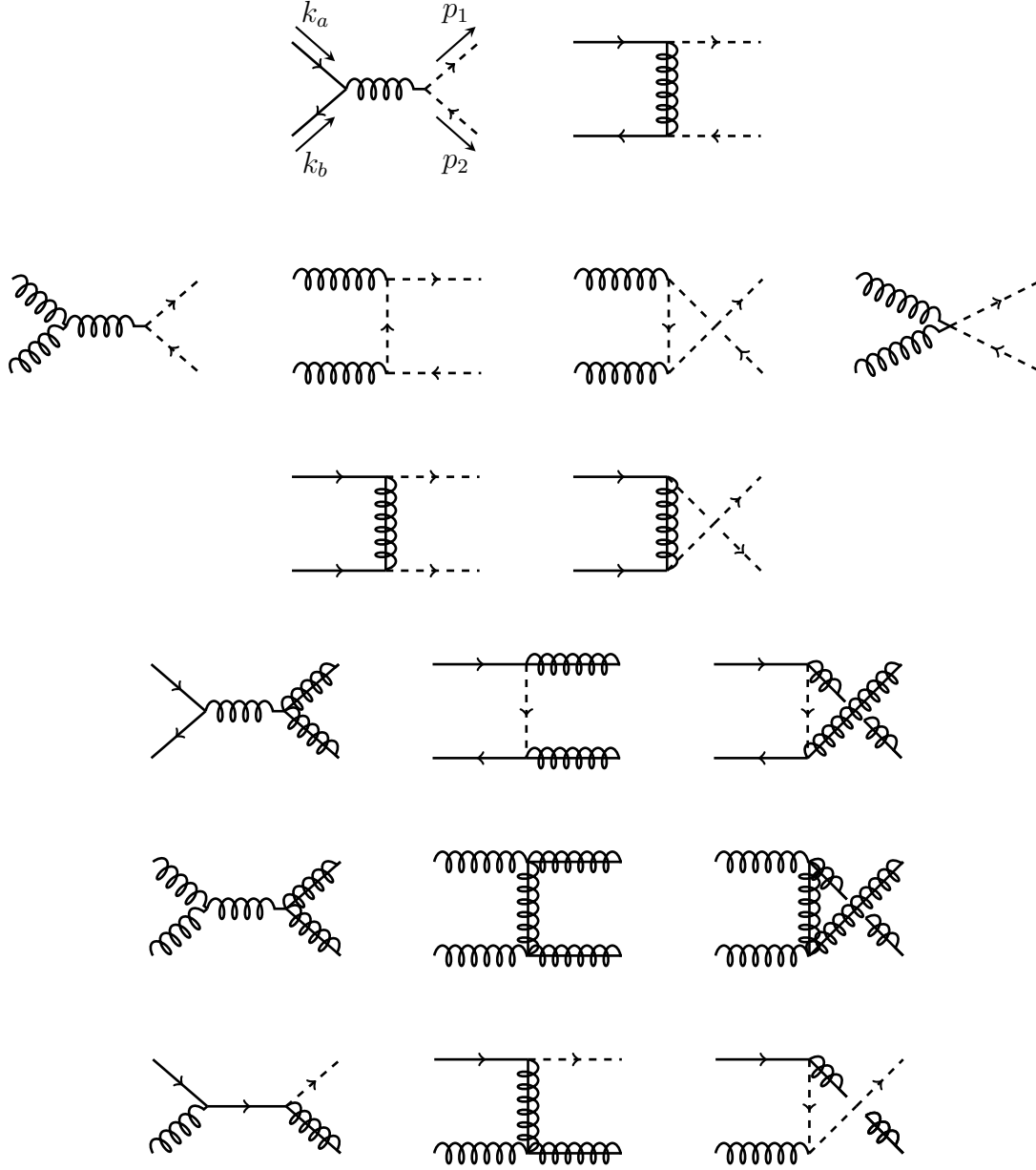
### 5.1 Partonic Processes and their Cross Section

For the following calculation the top quark is excluded from the initial state as it is too heavy to be significantly present in hadrons. For consistency reasons also the stop is excluded from the final states. One therefore deals with  $n_f - 1 = 5$  quark flavors. In view of the renormalization to be performed at the 1-loop level the results are given in  $D = 4 - 2\epsilon$  dimensions and it has been distinguished between the gauge coupling  $g_s$  from the gluon-quark-quark vertex and its supersymmetric analogue  $\hat{g}_s$  from the gluino-squark-quark vertex. The usual Mandelstam variables  $s, t, u$  and the following modifications of them are used<sup>15</sup>

$$\begin{aligned}
 s &= (k_a + k_b)^2 = (p_1 + p_2)^2 \\
 t &= (k_a - p_1)^2 = (k_b - p_2)^2 \\
 u &= (k_a - p_2)^2 = (k_b - p_1)^2 \\
 t_{\tilde{g}} &= t - m_{\tilde{g}}^2 & t_{\tilde{q}} &= t - m_{\tilde{q}}^2 \\
 u_{\tilde{g}} &= u - m_{\tilde{g}}^2 & u_{\tilde{q}} &= u - m_{\tilde{q}}^2
 \end{aligned} \tag{5.1}$$

<sup>14</sup>The production of sgluons is excluded from the analysis for their mass is chosen to be too large to significantly produce them.

<sup>15</sup>The kinematics of the process is like denoted in fig. 5.1.



**Figure 5.1:** Tree level diagrams for squark and gluino production at tree level in the MRSSM. The processes  $GG \rightarrow \tilde{q}\tilde{q}^*$  and  $qG \rightarrow \tilde{q}\tilde{g}$  are identical to those in the MSSM. The processes  $q\bar{q} \rightarrow \tilde{q}\tilde{q}^*$  and  $qq \rightarrow \tilde{q}\tilde{q}^*$  involve the production of less chiralities than in the MSSM. Also in the  $q\tilde{q} \rightarrow \tilde{g}\tilde{g}^*$  channel only half of the t-channel squark chiralities occur in the MRSSM. (Anti-)Gluino production via initial gluons  $GG \rightarrow \tilde{g}\tilde{g}^*$  proceeds via the same diagrams like in the MSSM but its cross section is twice as much in the MRSSM for gluino and antigluino are distinguishable particles.

Using the Feynman rules in the Appendix 10.5 one obtains the following sums over absolute squared Feynman amplitudes.

$$\begin{aligned} \sum |\mathcal{M}^B|^2(q_i \bar{q}_j \rightarrow \tilde{q} \tilde{q}^\dagger) &= \delta_{ij} \left[ 8N_c C(F) g_s^4 \frac{(n_f - 1)}{s^2} + 4N_c C(F) \hat{g}_s^4 \frac{1}{t_{\tilde{g}}^2} - 8C(F) g_s^2 \hat{g}_s^2 \frac{1}{t_{\tilde{g}} s} \right] (tu - m_{\tilde{q}}^4) \\ &\quad + (1 - \delta_{ij}) 4N_c C(F) \hat{g}_s^4 \frac{tu - m_{\tilde{q}}^4}{t_{\tilde{g}}^2} \end{aligned} \quad (5.2)$$

$$\begin{aligned} \sum |\mathcal{M}^B|^2(GG \rightarrow \tilde{q} \tilde{q}^\dagger) &= 4(n_f - 1) g_s^4 \left[ 2N_c^2 C(F) \left( 1 - 2 \frac{t_{\tilde{q}} u_{\tilde{q}}}{s^2} \right) - 2C(F) \right] \\ &\quad \left[ 1 - \epsilon - 2 \frac{s m_{\tilde{q}}^2}{t_{\tilde{q}} u_{\tilde{q}}} \left( 1 - \frac{s m_{\tilde{q}}^2}{t_{\tilde{q}} u_{\tilde{q}}} \right) \right] \end{aligned} \quad (5.3)$$

$$\begin{aligned} \sum |\mathcal{M}^B|^2(q_i q_j \rightarrow \tilde{q} \tilde{q}) &= \delta_{ij} 2 \hat{g}_s^4 N_c C(F) \left[ \frac{1}{t_{\tilde{g}}^2} + \frac{1}{u_{\tilde{g}}^2} \right] (tu - m_{\tilde{q}}^4) \\ &\quad + (1 - \delta_{ij}) 4 \hat{g}_s^4 N_c C(F) \frac{tu - m_{\tilde{q}}^4}{t_{\tilde{g}}^2} \end{aligned} \quad (5.4)$$

$$\begin{aligned} \sum |\mathcal{M}^B|^2(q \bar{q} \rightarrow \tilde{g} \tilde{g}) &= 8N_c^2 C(F) g_s^4 \left[ \frac{2m_{\tilde{g}}^2 s + t_{\tilde{g}}^2 + u_{\tilde{g}}^2}{s^2} - \epsilon \right] \\ &\quad + 4N_c^2 C(F) g_s^2 \hat{g}_s^2 \left[ \frac{m_{\tilde{g}}^2 s + t_{\tilde{g}}^2}{s t_{\tilde{q}}} + \frac{m_{\tilde{g}}^2 s + u_{\tilde{g}}^2}{s u_{\tilde{q}}} + \epsilon \left( \frac{t_{\tilde{g}}}{t_{\tilde{q}}} + \frac{u_{\tilde{g}}}{u_{\tilde{q}}} \right) \right] \\ &\quad + 2C(F) (N_c^2 - 1) \hat{g}_s^4 \left( \frac{t_{\tilde{g}}^2}{t_{\tilde{q}}^2} + \frac{u_{\tilde{g}}^2}{u_{\tilde{q}}^2} \right) \end{aligned} \quad (5.5)$$

$$\begin{aligned} \sum |\mathcal{M}^B|^2(GG \rightarrow \tilde{g} \tilde{g}) &= 16N_c^3 C(F) g_s^4 \left( 1 - \frac{t_{\tilde{g}} u_{\tilde{g}}}{s^2} \right) \\ &\quad \left[ \frac{s^2}{t_{\tilde{g}} u_{\tilde{g}}} (1 - \epsilon)^2 - 2(1 - \epsilon) + 4 \frac{m_{\tilde{g}}^2 s}{t_{\tilde{g}} u_{\tilde{g}}} \left( 1 - \frac{m_{\tilde{g}}^2 s}{t_{\tilde{g}} u_{\tilde{g}}} \right) \right] \end{aligned} \quad (5.6)$$

$$\begin{aligned} \sum |\mathcal{M}^B|^2(qg \rightarrow \tilde{q} \tilde{g}) &= 2g_s^2 \hat{g}_s^2 \left[ 2N_c^2 C(F) \left( 1 - 2 \frac{s u_{\tilde{q}}}{t_{\tilde{g}}^2} \right) - 2C(F) \right] \\ &\quad \left[ (-1 + \epsilon) \frac{t_{\tilde{g}}}{s} + \frac{2(m_{\tilde{g}}^2 - m_{\tilde{q}}^2) t_{\tilde{g}}}{s u_{\tilde{q}}} \left( 1 + \frac{m_{\tilde{q}}^2}{u_{\tilde{q}}} + \frac{m_{\tilde{g}}^2}{t_{\tilde{g}}} \right) \right] \end{aligned} \quad (5.7)$$

Having calculated the absolute squared Feynman amplitudes one obtains the partonic cross sections (see Appendix 10.7) via

$$\begin{aligned} \frac{d^2 \sigma^B}{dt du} &= \frac{K_{ab}}{s^2} \frac{\pi S_\epsilon}{\Gamma(1 - \epsilon)} \left[ \frac{tu - m_1^2 m_2^2}{\mu^2 s} \right]^{-\epsilon} \Theta(tu - m_1^2 m_2^2) \\ &\quad \Theta(s - 4m^2) \delta(s + t + u - m_1^2 - m_2^2) \sum |\mathcal{M}^B|^2 \end{aligned} \quad (5.8)$$

where  $m_1$  ( $m_2$ ) is the mass of the first (second) final state particle and  $m$  is their arithmetic

mean. The leading order cross sections are

$$\begin{aligned}\sigma^B(q_i \bar{q}_j \rightarrow \tilde{q} \tilde{q}^\dagger) &= \delta_{ij} \frac{g_s^4}{16\pi s} (n_f - 1) \left[ \frac{4}{27} - \frac{16m_{\tilde{q}}^2}{27s} \right] \\ &\quad + \delta_{ij} \frac{g_s^2 \hat{g}_s^2}{16\pi s} \left[ \left( \frac{4}{27} + \frac{8m_-^2}{27s} \right) \beta_{\tilde{q}} + \left( \frac{8m_{\tilde{g}}^2}{27s} + \frac{8m_-^4}{27s^2} \right) L_1 \right] \\ &\quad + \frac{\hat{g}_s^4}{16\pi s} \left[ -\frac{8}{9} \beta_{\tilde{q}} + \left( -\frac{4}{9} - \frac{8m_-^2}{9s} \right) L_1 \right]\end{aligned}\quad (5.9)$$

$$\sigma^B(GG \rightarrow \tilde{q} \tilde{q}^\dagger) = \frac{(n_f - 1)g_s^4}{16\pi s} \left[ \left( \frac{5}{24} + \frac{31m_{\tilde{q}}^2}{12s} \right) \beta_{\tilde{q}} + \left( \frac{4m_{\tilde{q}}^2}{3s} + \frac{m_{\tilde{q}}^4}{3s^2} \right) \ln \frac{1 - \beta_{\tilde{q}}}{1 + \beta_{\tilde{q}}} \right] \quad (5.10)$$

$$\sigma^B(q_i q_j \rightarrow \tilde{q} \tilde{q}) = \frac{\hat{g}_s^4}{16\pi s} \left[ -\frac{8}{9} \beta_{\tilde{q}} + \left( -\frac{4}{9} - \frac{8m_-^2}{9s} \right) L_1 \right] \quad (5.11)$$

$$\begin{aligned}\sigma^B(q\bar{q} \rightarrow \tilde{g} \tilde{g}) &= \frac{g_s^4}{16\pi s} \left[ \frac{16}{9} + \frac{32m_{\tilde{g}}^2}{9s} \right] \beta_{\tilde{g}} \\ &\quad + \frac{\hat{g}_s^2 g_s^2}{16\pi s} \left[ \left( -\frac{4}{3} - \frac{8m_-^2}{3s} \right) \beta_{\tilde{g}} + \left( \frac{8m_{\tilde{g}}^2}{3s} + \frac{8m_-^4}{3s^2} \right) L_2 \right] \\ &\quad + \frac{\hat{g}_s^4}{16\pi s} \left[ \left( \frac{32}{27} + \frac{32m_-^4}{27(m_-^4 + m_{\tilde{q}}^2 s)} \right) \beta_{\tilde{g}} - \frac{64m_-^2}{27s} L_2 \right]\end{aligned}\quad (5.12)$$

$$\sigma^B(GG \rightarrow \tilde{g} \tilde{g}) = \frac{g_s^4}{16\pi s} \left[ \left( -6 - \frac{51m_{\tilde{g}}^2}{2s} \right) \beta_{\tilde{g}} + \left( -\frac{9}{2} - \frac{18m_{\tilde{g}}^2}{s} + \frac{18m_{\tilde{g}}^4}{s^2} \right) \ln \frac{1 - \beta_{\tilde{g}}}{1 + \beta_{\tilde{g}}} \right] \quad (5.13)$$

$$\begin{aligned}\sigma^B(qG \rightarrow \tilde{q} \tilde{g}) &= \frac{g_s^2 \hat{g}_s^2}{16\pi s} \left[ \frac{\kappa}{s} \left( -\frac{7}{9} - \frac{32m_-^2}{9s} \right) + \left( -\frac{8m_-^2}{9s} + \frac{2m_{\tilde{q}}^2 m_-^2}{s^2} + \frac{8m_-^4}{9s^2} \right) L_3 \right. \\ &\quad \left. + \left( -1 - \frac{2m_-^2}{s} + \frac{2m_{\tilde{q}} m_-^2}{s^2} \right) L_4 \right]\end{aligned}\quad (5.14)$$

where the abbreviations [16]

$$\begin{aligned}\beta_{\tilde{q}} &= \sqrt{1 - \frac{4m_{\tilde{q}}^2}{s}} & \beta_{\tilde{g}} &= \sqrt{1 - \frac{4m_{\tilde{g}}^2}{s}} \\ m_-^2 &= m_{\tilde{g}}^2 - m_{\tilde{q}}^2 & \kappa &= \sqrt{(s - m_{\tilde{g}}^2 - m_{\tilde{q}}^2)^2 - 4m_{\tilde{g}}^2 m_{\tilde{q}}^2} \\ L_1 &= \ln \frac{s + 2m_-^2 - s\beta_{\tilde{q}}}{s + 2m_-^2 + s\beta_{\tilde{q}}} & L_2 &= \ln \frac{s - 2m_-^2 - s\beta_{\tilde{g}}}{s - 2m_-^2 + s\beta_{\tilde{g}}} \\ L_3 &= \ln \frac{s - m_-^2 - \kappa}{s - m_-^2 + \kappa} & L_4 &= \ln \frac{s + m_-^2 - \kappa}{s + m_-^2 + \kappa}\end{aligned}\quad (5.15)$$

are used. In the following the considered processes of the MRSSM are compared to those of the MSSM.

**The process  $q_i \bar{q}_j \rightarrow \tilde{q} \tilde{q}^\dagger$** 

The production of a squark and an antisquark through a quark and an antiquark in the initial state originates from two types of Feynman diagrams. The first one has an  $s$ -channel gluon and is the same in the MSSM and in the MRSSM. The second one exhibits a difference due to the  $t$ -channel gluino which is no Majorana particle in the MRSSM. To see this one can either



**Figure 5.2:** Tree level diagrams for  $q\bar{q} \rightarrow \tilde{q}\tilde{q}^\dagger$

just apply the Feynman rules in Appendix 10.5 or think of the conservation of R-charge: Left handed squarks have R-charge +1 and right handed squarks have R-charge -1. Antiparticles have the opposite R-charge of their corresponding particles. The final state particles have to meet the total R-charge zero from the initial state. Because of this one only has  $\tilde{q}_L \bar{\tilde{q}}_L$  and  $\tilde{q}_R \bar{\tilde{q}}_R$  as the final states in the MRSSM whereas in the MSSM one actually has four instead of two t-channel diagrams: The corresponding final states are  $\tilde{q}_L \bar{\tilde{q}}_L$ ,  $\tilde{q}_R \bar{\tilde{q}}_R$ ,  $\tilde{q}_L \bar{\tilde{q}}_R$  and  $\tilde{q}_R \bar{\tilde{q}}_L$ . Consequently contributions from the  $t$ -channel diagrams are suppressed in the MRSSM in comparison to the MSSM. As visible in fig. 5.5 this suppression grows with the masses. The reason for this lies in the gluino mass dependence of the  $t$ -channel diagrams which is explained in the discussion of squark production.

If the initial state quarks are of different flavor the  $s$ -channel diagram is absent.

**The process  $GG \rightarrow \tilde{q}\tilde{q}^\dagger$** 

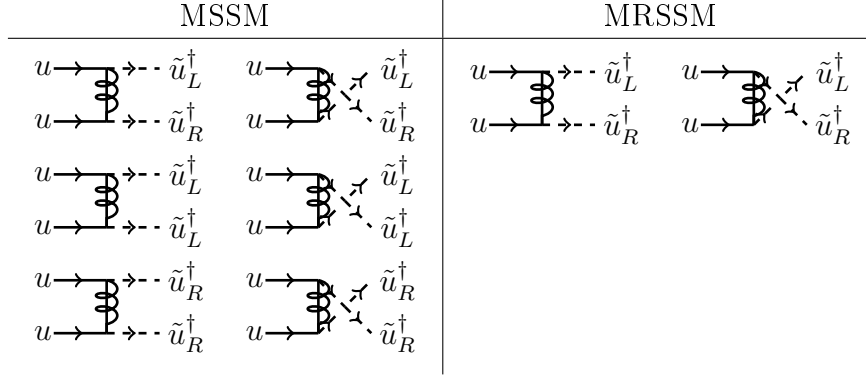
This process has the same cross section as in the MSSM for also in the MSSM only like chirality squark and antisquark  $\tilde{q}_A \bar{\tilde{q}}_A^\dagger$  with  $A \in L, R$  can be produced.

**The process  $q_i q_j \rightarrow \tilde{q}\tilde{q}$** 

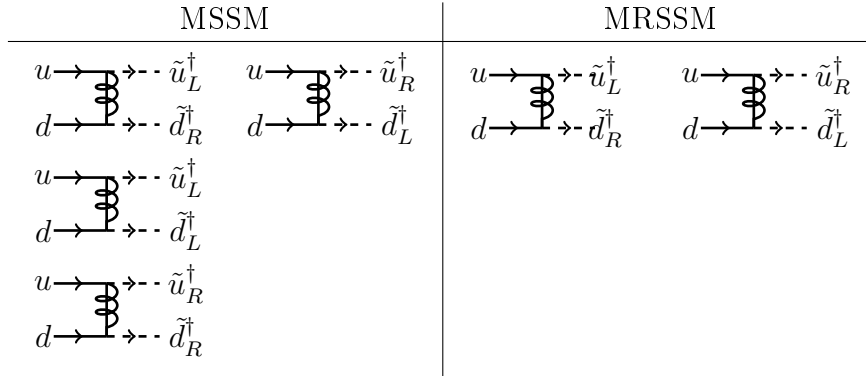
In the MRSSM only the production of unlike chirality squarks  $\tilde{q}_L \tilde{q}_R$  is allowed whilst in the MSSM also like chirality squarks can be produced. This is again a consequence of the conservation of R-charge. The upshot of this is a suppression of squark production in the MRSSM in comparison to the MSSM. To be more explicit the suppression of squark production in the MRSSM grows with the gluino mass. This can be understood as follows: As in the MRSSM a left handed squark needs to be produced with a right handed squark one can read off from the Feynman rules given in Appendix 10.5 that the gluino propagator  $i \frac{\not{p} + m_{\tilde{g}}}{p^2 - m_{\tilde{g}}^2}$  is sandwiched between the projectors  $P_L$  and  $P_R$  which leads to the cancellation of the gluino mass in the numerator. Therefore for small momenta of the gluino compared to the gluino mass one gets  $\mathcal{M} \sim \frac{1}{m_{\tilde{g}}^2}$  in the MRSSM while in the MSSM one finds a suppression proportional to  $\frac{1}{m_{\tilde{g}}}$ .



When considering squark production one has to sum over all  $|\mathcal{M}|^2$  with different helicity. The diagrams for those different final state helicities are shown for flavor like squark production in tab. 5.1 and for flavor unlike squark production in tab. 5.2 for the MSSM and the MRSSM. To avoid double counting one needs to weight  $|\mathcal{M}|^2(qq \rightarrow \tilde{q}_L \tilde{q}_L)$  and  $|\mathcal{M}|^2(qq \rightarrow \tilde{q}_R \tilde{q}_R)$  with a statistical factor of  $\frac{1}{2}$  as one integrates over all momenta of both particles to arrive at the cross section.



**Table 5.1:** All Feynman diagrams contributing to flavor like squark production in the MSSM and the MRSSM for the example of  $u$ -quarks. For the MSSM the absolute squared Feynman amplitudes from the diagrams with chirality like squarks need to be weighted with a factor of  $\frac{1}{2}$ .



**Table 5.2:** All Feynman diagrams contributing to flavor unlike squark production in the MSSM and the MRSSM for the example of  $u$ - and  $d$ -quarks.

Because of the absence of chirality like squarks in the final state in the MRSSM (and a missing interference of the two diagrams in tab. 5.1 in the column of the MRSSM) the partonic cross section of flavor like and unlike squarks is the same in the MRSSM, i.e. on the partonic level:

$$\sigma_{\text{Part, MRSSM}}^{\text{B}}(uu \rightarrow \tilde{u}_L \tilde{u}_R) = \sigma_{\text{Part, MRSSM}}^{\text{B}}(ud \rightarrow \tilde{u}_L \tilde{d}_R) + \sigma_{\text{Part, MRSSM}}^{\text{B}}(ud \rightarrow \tilde{u}_R \tilde{d}_L). \quad (5.16)$$

That is however not true in the MSSM because for flavor like squarks there is a non vanishing interference of the  $t$ - and  $u$ -channel diagram of the processes  $qq \rightarrow \tilde{q}_L \tilde{q}_L$  and  $qq \rightarrow \tilde{q}_R \tilde{q}_R$ . This

term is given in [16, p.7]

### The process $q\bar{q} \rightarrow \tilde{g}\tilde{g}$

In contrast to the MSSM no statistical factor of  $\frac{1}{2}$  is taken into account when turning from  $|\mathcal{M}|^2$  to  $\sigma$ . This is because gluino and antigluino are distinguishable particles. Still in comparison to the MSSM cross section [16, p.9] only the first line in 5.12 is doubled up as the other two lines originate from an  $t$  or  $u$  channel squark which occurs in only one instead of two chiralities. Furthermore an interference term from the  $t$  and  $u$  channel diagram which occurs in the MSSM is absent in the MRSSM.

### The process $GG \rightarrow \tilde{g}\tilde{g}$

As in the previous process the MRSSM cross section for  $GG \rightarrow \tilde{g}\tilde{g}$  receives no statistical factor of  $\frac{1}{2}$  like in the MSSM. As there are no further differences between MSSM and MRSSM in this channel, the MRSSM cross section is simply twice as large as in the MSSM.

### The process $qG \rightarrow \tilde{q}\tilde{g}$

This process is exactly the same in the MSSM and MRSSM for there are two differences between the models which exactly cancel each other: Firstly in the MRSSM there is not only the production of a squark and a gluino but also the production of a squark and an antigluino. But on the other hand R-charge needs to be conserved which dictates that a "left-handed" squark can only be produced with an antigluino and a "right-handed" squark can only be produced with a gluino (see table 4.2).

In order to account for the confinement of quarks and gluons within hadrons the hadronic cross sections for the considered processes are calculated in the next subsection.

## 5.2 Hadronic Cross Section

Quarks, antiquarks and gluons are no free particles but are confined within hadrons. As hadrons consist of a variety of the just mentioned partons which share the hadron's momentum one does not have a definite initial state in hadron collisions which is assumed in the previous subsection. Fortunately the hadronic cross section for the production of a final state  $X$ , e.g.  $X = \tilde{q}\tilde{q}$ , can be obtained by convolving the partonic cross section with parton density functions of the initial hadrons.

$$\sigma_{\text{Had}}^{\text{B}}(P_1 P_2 \rightarrow X) = \int dx_1 dx_2 f_{P_1/H_1}(x_1) f_{P_2/H_2}(x_2) \sigma_{\text{Part}}^{\text{B}}(P_1 P_2 \rightarrow X, s = x_1 x_2 S). \quad (5.17)$$

The parton density functions  $f_{P_i/H_j}(x)$  and the momentum fractions  $x_i$  are explained below. The hadrons center-of-mass energy is denoted with  $\sqrt{S}$  where  $\sqrt{s}$  denotes the partons center-of-mass energy.

As the production of the final state  $X$  may proceed via various initial partons one has to sum over all possible possibilities arising from the initial hadrons  $H_1$  and  $H_2$ :

$$\sigma_{\text{Had}}^B(H_1 H_2 \rightarrow X) = \sum_{i,j} \sigma_{\text{Had}}^B(P_i P_j \rightarrow X) \quad (5.18)$$

where the sum runs over all partons  $P_i$  ( $P_j$ ) which are in the hadron  $H_1$  ( $H_2$ ).

To get an intuitive idea of this factorization regard the hadrons as extended objects consisting of partons<sup>16</sup> which are permanently interacting with each other. Now consider two colliding hadrons in their center-of-mass frame. Due to Lorentz contraction the hadrons appear as thin discs and the parton's mutual interactions are time-delayed within this frame. This effectively means that a hadron at the time of collision is virtually frozen.

maybe picture of colliding disks

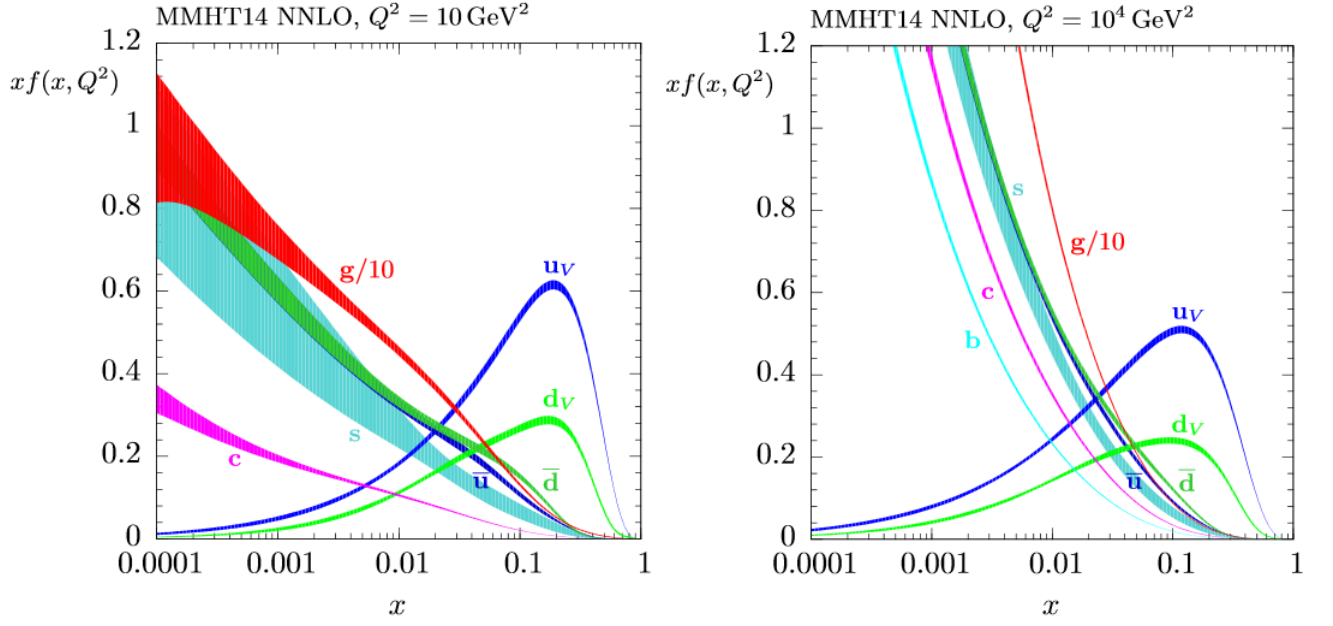
This in turn implies that the hadron consists at the time of collision of a definite number of partons which can be thought of as carrying a definite fraction of the hadron's momentum  $p_{\text{Part}} = x p_{\text{Had}}$  with  $x \in (0, 1)$ . The parton density function  $f_{P/H}(x)$  can therefore be understood as the probability of finding a parton  $P$  within the hadron  $H$  carrying  $x$  of its momentum.

They actually depend on an arbitrary energy scale referred to as the factorization scale  $\mu_F$ . The  $\mu_F$  evolution is given within the scope of perturbation theory by coupled integro-differential equations named after Gribov, Lipatov, Dokshitzer, Altarelli and Parisi: the DGLAP-equations. For further reading on parton density functions and factorization see [18] and [19]. See [20] for how to determine parton density functions.

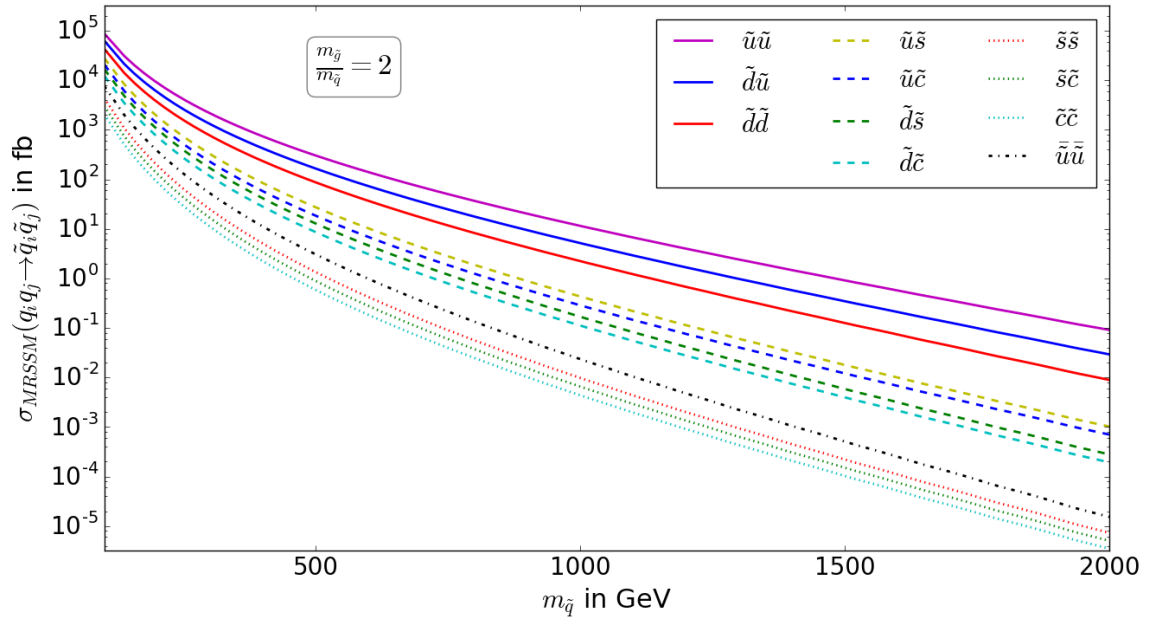
Fig. 5.3 shows a parton density function set for the proton. One can see that the partons most probable to carry a high fraction of momentum are the valence quarks. For lower values of  $x$  gluons constitute the major component of the proton.

Fig.5.5 shows the hadronic cross section for the production of various squarks in the MRSSM. These can only be produced by their non supersymmetric partner, e.g. two up-squarks can only be produced by two up-quarks. Because the partonic cross section for these processes are the same, cf. eq. 5.16, one can see the influence of the parton density functions quite nicely: The up-squark production is the dominant contribution to squark production. Apart from the solid lines in fig.5.5 which correspond to the production of first generation squarks, also the cross section of mixed first and second generation (dashed lines) and second generation (dotted lines) squarks is shown. The dashed dotted line shows the cross section of the charged conjugated particles of the dominant channel.

<sup>16</sup>The proton is considered as composed of three valence quarks: Two up-quarks and one down-quark. In addition there are gluons and seaquarks, i.e. virtual quark-antiquark pairs.



**Figure 5.3:** Parton density function of the proton, taken from MMHT2014 [17] with 68% confidence level. The parton density function is multiplied by  $x$  in order to counteract its peak at  $x \rightarrow 0$ . The parton density function is given at two different factorisation scales  $\mu_F = Q$



**Figure 5.4:** Hadronic cross section for squark production in the MRSSM at the LHC at  $\sqrt{S} = 13 \text{ GeV}$ . The ratio of the gluino and squark mass is fixed to 2. The parton densities used are MMHT2014 LO with  $\alpha_s(M_Z) = 0.135$  in the 5-flavor scheme [17]. As renormalization and factorization scale  $\mu_R = \mu_F = \frac{m_1 + m_2}{2}$  has been chosen, where  $m_i$  are the final state particle masses.

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The following paragraph discusses fig. 5.5, 5.6 and 5.7, i.e. the hadronic cross section of all considered processes for three successive mass ratios of the gluino and squark:  $\frac{m_{\tilde{g}}}{m_{\tilde{q}}} = 0.9, 2$  and 5. Also the corresponding fraction of final states is shown.

For the small mass ratio  $\frac{m_{\tilde{g}}}{m_{\tilde{q}}} = 0.9$  and small squark masses the dominant contribution comes from gluino (antigluino) production. This is even more dominant in the MRSSM which can be explained by the above mentioned factor of 2 difference in the  $GG \rightarrow \tilde{g}\tilde{g}(\bar{\tilde{g}})$  channel which comes from the distinguishable gluino and antigluino in the MRSSM. Furthermore there is the above mentioned destructive interference term in the  $q\bar{q} \rightarrow \tilde{g}\tilde{g}(\bar{\tilde{g}})$  channel which makes the cross section in the MRSSM dominant over the one in the MSSM.

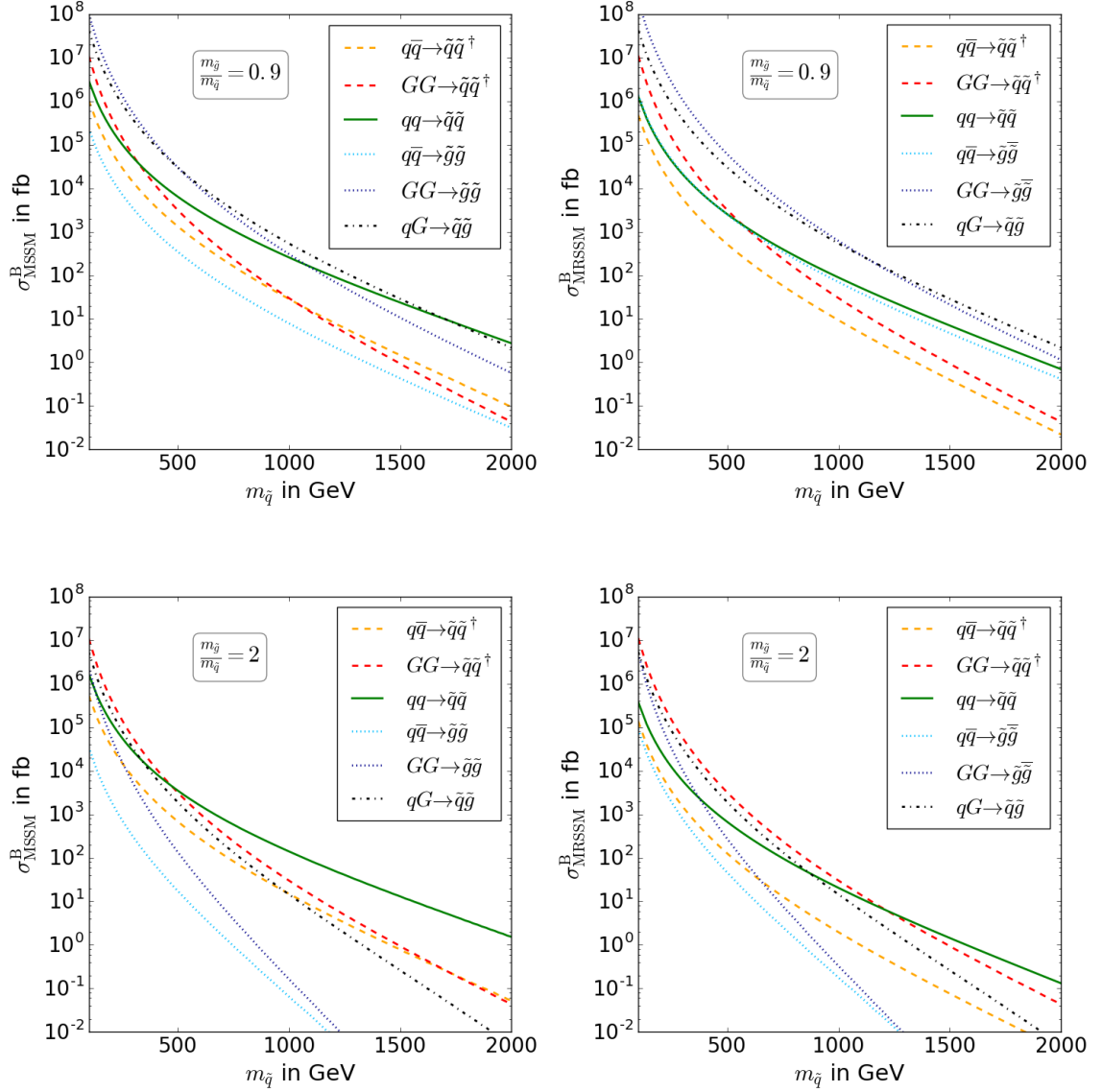
For a growing squark mass the previously subdominant gluino-squark production takes over and becomes dominant. A salient difference between the MSSM and the MRSSM is that squark production is strongly suppressed in the MRSSM. While in the MSSM it becomes dominant above about 1700 GeV it is subsubdominant for the whole displayed squark mass in the MRSSM. This suppression has also already been explained above.

Turning from  $\frac{m_{\tilde{g}}}{m_{\tilde{q}}} = 0.9$  to  $\frac{m_{\tilde{g}}}{m_{\tilde{q}}} = 2$  suppresses quite strongly the production of gluinos in all pertaining channels. For the MSSM one finds above 500 GeV a growing dominance of squark production. The same happens in the MRSSM but for a significantly higher squark mass of about 1300 GeV. This suppression of squark production in the MRSSM increases with a growing gluino masses which can be seen when going from  $\frac{m_{\tilde{g}}}{m_{\tilde{q}}} = 0.9$  to  $\frac{m_{\tilde{g}}}{m_{\tilde{q}}} = 2$  or even  $\frac{m_{\tilde{g}}}{m_{\tilde{q}}} = 5$  in fig. 5.7. This reflects the gluino mass dependence explained in detail in section 5.1. See also 5.9 for a visualization of this behavior.

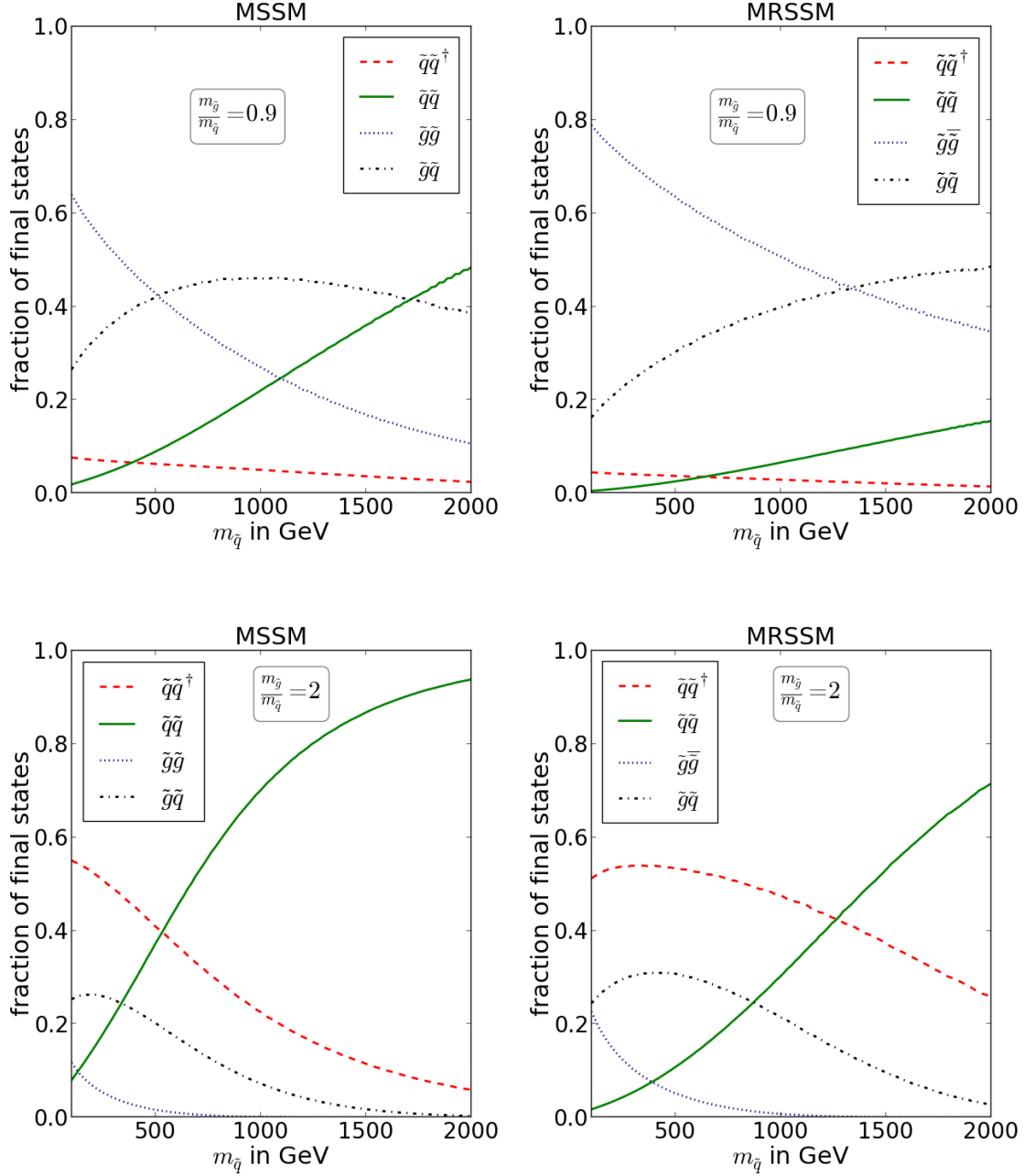
Note also that for sufficiently large squark masses squark production becomes the dominant process, see fig. 5.5 and 5.6. This is because the production of heavy particles requires large longitudinal momenta of the colliding partons, i.e. large  $x$ . Looking at the parton density functions in fig. 5.3 one sees that within the proton the partons most likely to carry large momentum are quarks. So looking at large  $x$  the proton basically consists of quarks and because the only process involving only quarks in the initial state is squark production, this channel dominates for high masses.

The total cross section in the MRSSM is increasingly smaller than in the MSSM in particular from about 500 GeV onwards. The is because for masses below that the  $GG \rightarrow \tilde{q}\tilde{q}^\dagger$  channel which is the same for both models is the dominant one.

When raising the gluino squark mass ratio further to  $\frac{m_{\tilde{g}}}{m_{\tilde{q}}} = 5$  the above described tendency continues. This is all processes including a gluino get even more negligible and squark production gets further suppressed. While in the MSSM it becomes dominant over  $GG \rightarrow \tilde{q}\tilde{q}^\dagger$  at

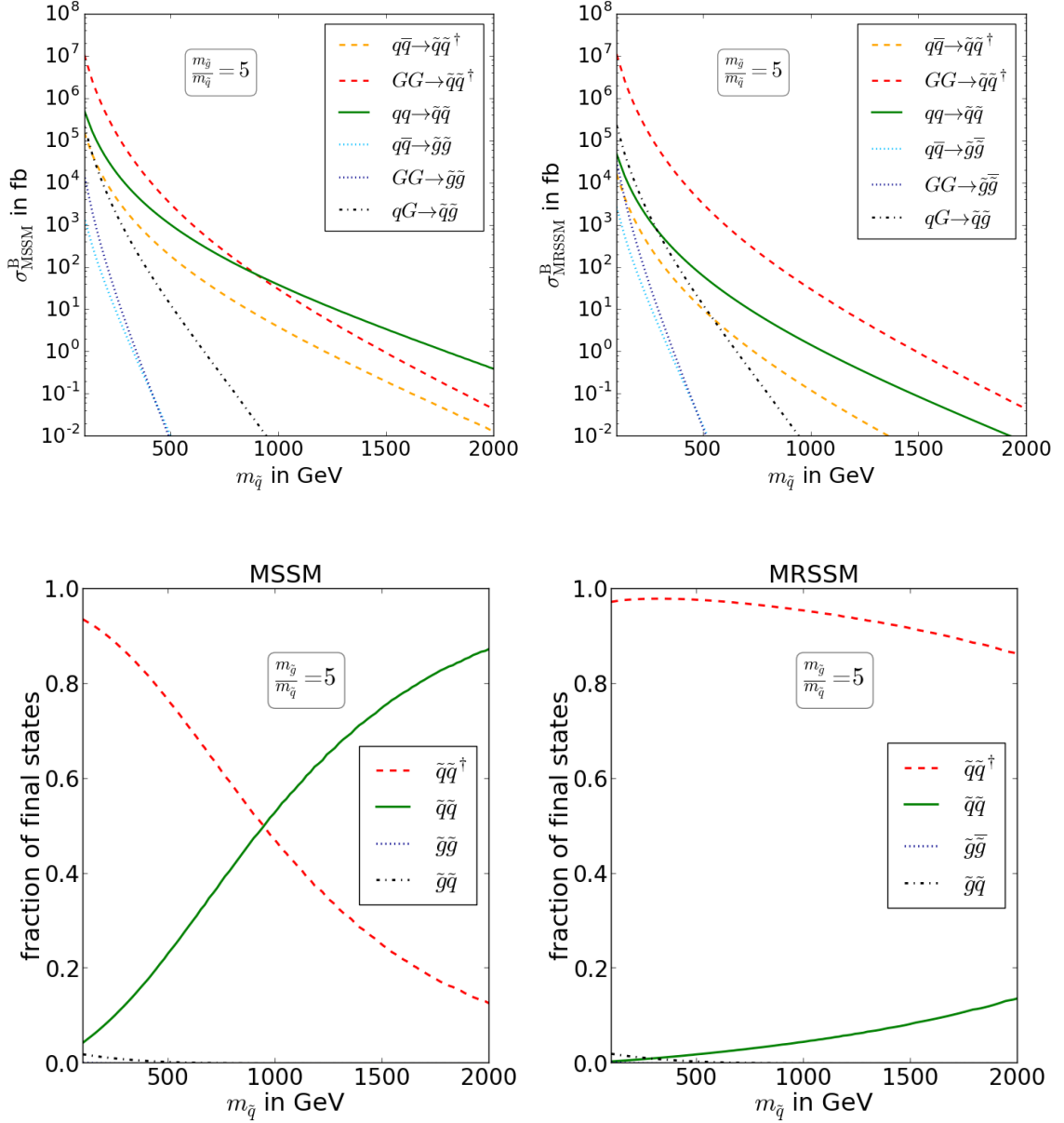


**Figure 5.5:** Hadronic cross section for squark and gluino production in the MSSM (left-hand side) and MRSSM (right-hand side) at the LHC with  $\sqrt{S} = 13 \text{ GeV}$ . The ratio of gluino and squark mass is fixed to 0.9 (first row) and 2 (second row). In the final state it has been summed over all squark flavors except for  $t$ -squarks. For the channels  $q q \rightarrow \tilde{q}\tilde{q}$  and  $q G \rightarrow \tilde{q}\tilde{g}$  also the charge conjugated process is included. The parton densities used are MMHT2014 LO with  $\alpha_s(M_Z) = 0.135$  in the 5-flavor scheme [17]. As renormalization and factorization scale  $\mu_R = \mu_F = \frac{m_1 + m_2}{2}$  has been chosen, where  $m_i$  are the final state particle masses.



**Figure 5.6:** Relative contributions of the indicated final states on the total hadronic cross section in the MSSM (left-hand side) and MRSSM (right-hand side) at the LHC with  $\sqrt{S} = 13$  GeV. The ratio of gluino and squark mass is fixed to 0.9 (first row) and 2 (second row). In the final state it has been summed over all squark flavors except for  $t$ -quarks. For the channels  $qq \rightarrow \tilde{q}\tilde{q}$  and  $qG \rightarrow \tilde{q}\tilde{G}$  also the charge conjugated process is included. The parton densities and the renormalization and factorization scale are chosen as in fig. 5.5.

$m_{\tilde{q}} \approx 1000$  GeV it is subdominant in the MRSSM in the whole displayed mass range. even though it does not fall that steeply as the dominant channel.



**Figure 5.7:** Hadronic cross section for squark and gluino production (in the first row) and relative contributions of the indicated final states on the total hadronic cross section (in the second row) in the MSSM (left-hand side) and MRSSM (right-hand side) at the LHC which  $\sqrt{S} = 13$  GeV. The ratio of gluino and squark mass is fixed to 5. In the final state it has been summed over all squark flavors except for staus. For the channels  $qq \rightarrow \tilde{q}\tilde{q}$  and  $qG \rightarrow \tilde{q}\tilde{g}$  also the charge conjugated process is included. The parton densities and the renormalization and factorization scale are chosen as in fig. 5.5.

The contour plots 5.8 underline the already announced difference between squark production



in the MRSSM and the MSSM. The main difference is a suppression of the cross section for higher gluino masses in the MRSSM. Fig.5.9 shows the Born cross section of squark production including all possible flavors for a fixed mass of the gluino  $m_{\tilde{g}} = 2000 \text{ GeV}$  (top) and the squark  $m_{\tilde{q}} = 1000 \text{ GeV}$  (bottom). The figure includes errors of the calculation which come from the parton density functions and the renormalization and factorization scale.

The 68% uncertainty from the central value pdf-set is obtained by performing the calculation with all eigenvector sets of the central value and taking the minimum and maximum at each point as deviations from the calculation with the central set.

The error from the renormalization scale  $\mu_R$  and the factorization scale  $\mu_F$  is obtained by varying both parameters with a factor of  $\frac{1}{2}$  and two. For each parameter combination the cross section has been calculated. Out of these nine numbers for each point the minimum and maximum has been chosen to estimate the error originating from the scale.

The total error is then obtained by Gauß's propagation of uncertainties:

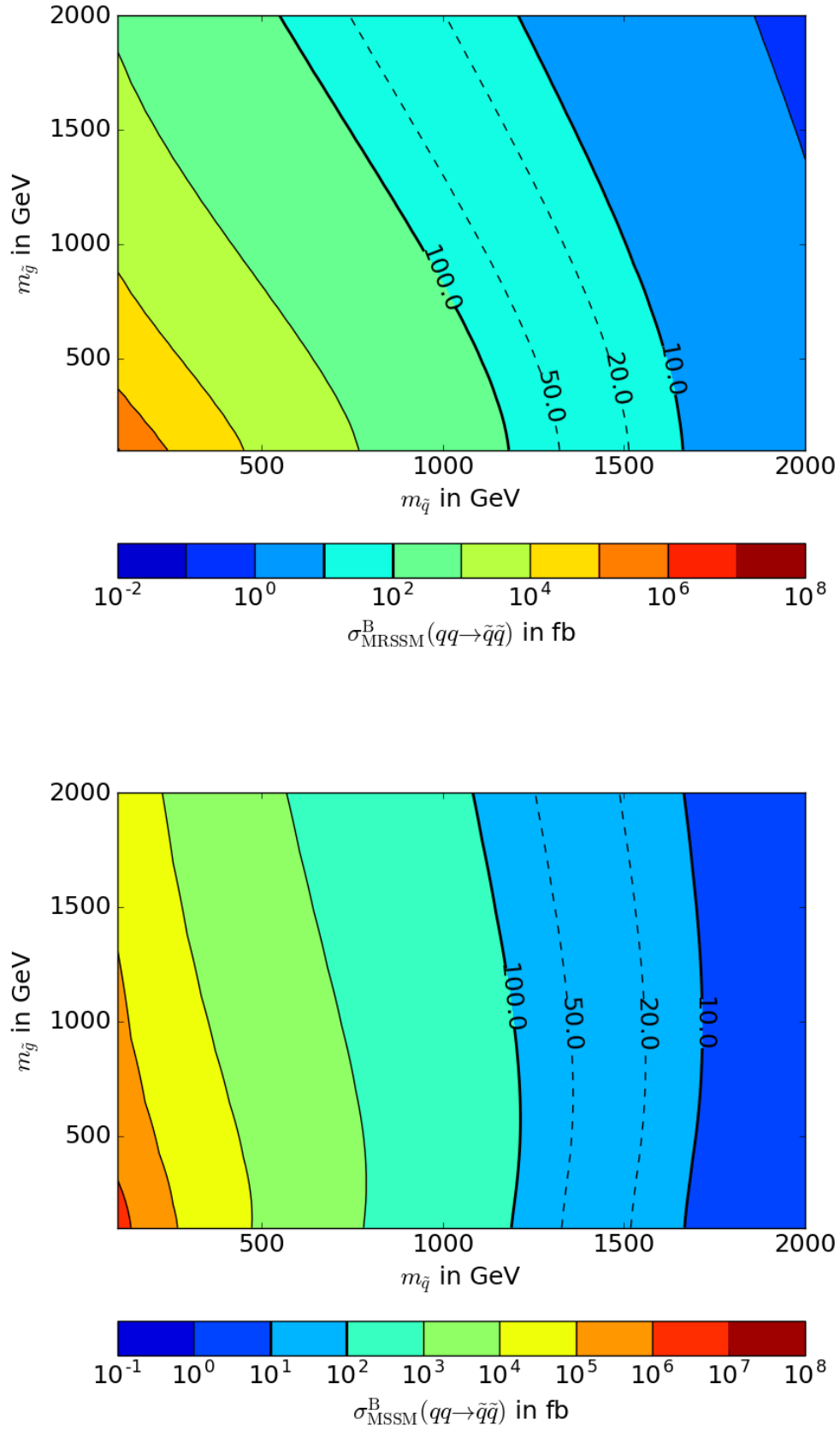
$$\Delta\sigma = \sqrt{(\Delta\sigma_{\text{pdf}})^2 + (\Delta\sigma_{\text{scale}})^2}. \quad (5.19)$$

The uncertainty coming from the scale variation has been found much greater than the one from the parton density functions. The error from the phase space integration is negligible, i.e.  $\Delta\sigma_{\text{int}} < 1\%$ .

As a summary of the comparison between MSSM and MRSSM gluino production one can say, that the gluino production is enhanced in the MRSSM due to distinguishable gluino and antigluino. In order to meet missing experimental evidence for supersymmetry this suggests that gluinos might be heavier than squarks in the MRSSM. Focusing on  $\frac{m_{\tilde{g}}}{m_{\tilde{q}}} \geq 2$  the most significant change between the considered models is the growing suppression of squark production for a increasing gluino mass. This is why this channel is going to be investigated at next-to-leading order within this thesis. The ensuing sections provide the prerequisites for this calculation.

$m_{\tilde{g}} \setminus m_{\tilde{q}}$	500	1000	1500
1000			
1500			
2000			

**Table 5.3:** Table



**Figure 5.8:** Contour plot for squark production cross section including all possible flavors (also charged ) in the MRSSM (top) and MSSM (bottom)

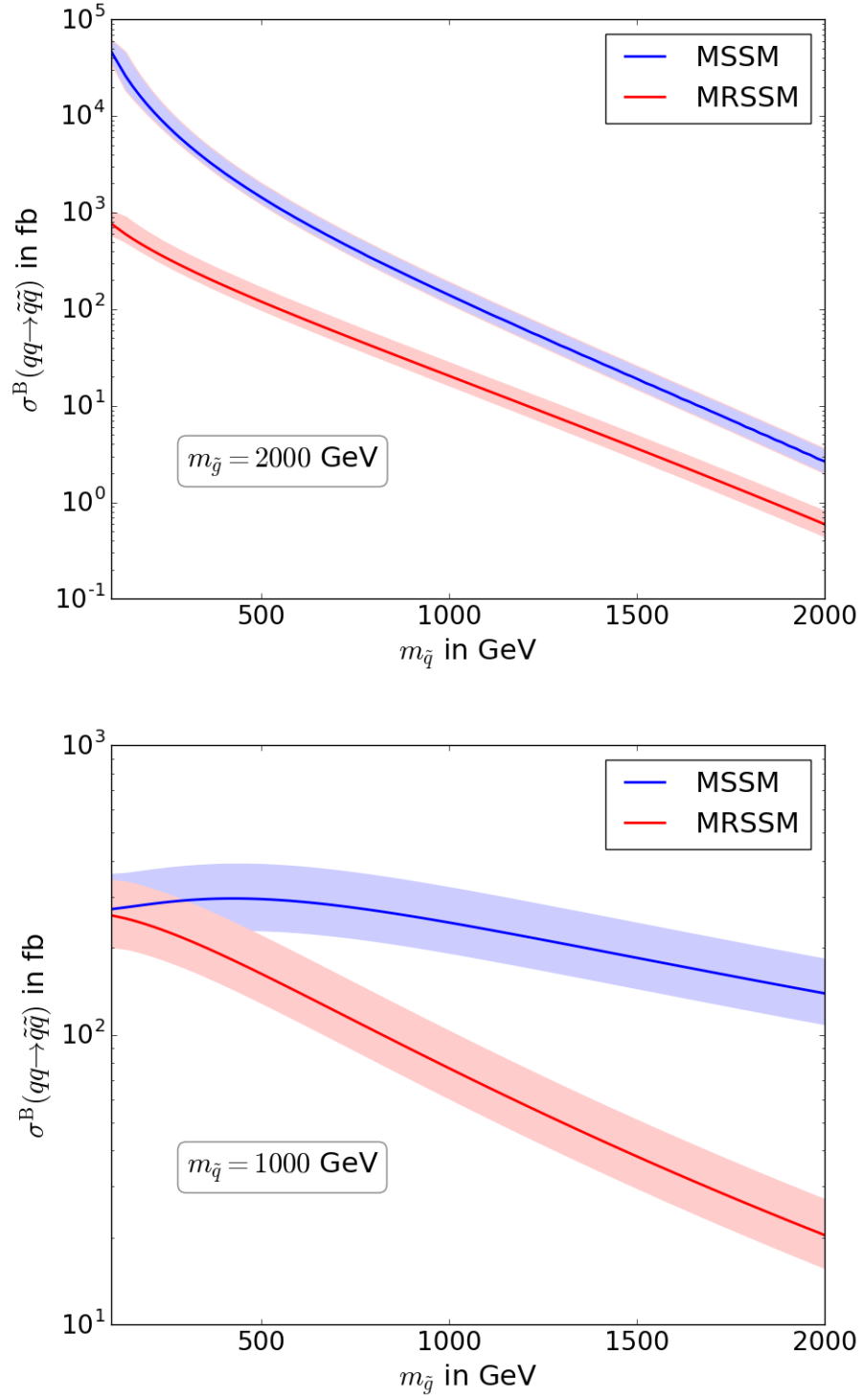


Figure 5.9: add error bars

## 6 Virtual and Real Corrections

The results calculated in the previous chapter are not the full theoretical predictions for the different processes but only the first term in a perturbation theory of  $\alpha_s$ .

This section describes the necessary steps in the calculation of the squark production cross section at next-to-leading-order, i.e. its  $\mathcal{O}(\alpha_s)$ -correction.

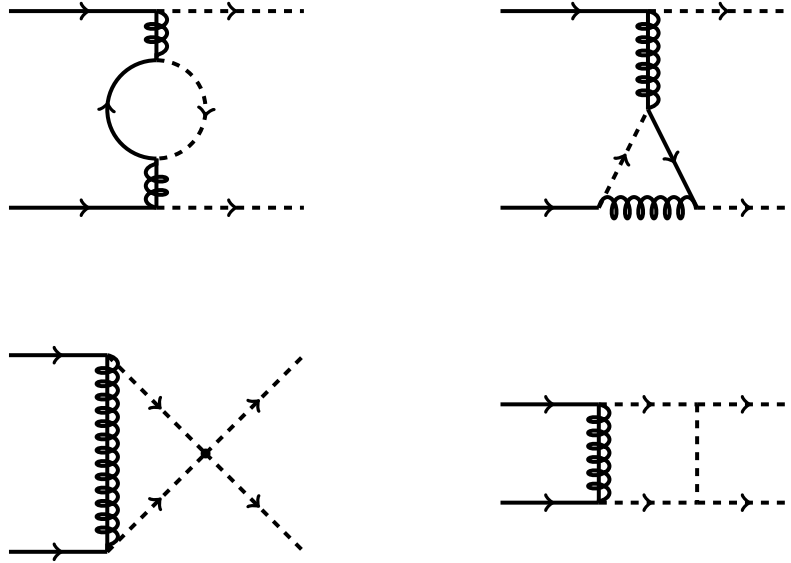
Reason validity of perturbation theory in this context.

### 6.1 Virtual Correction

According to quantum field theory the  $\mathcal{O}(\alpha_s)$  correction to a tree level process includes the computation of one-loop diagrams such as shown in fig. 6.1. To yield  $\mathcal{O}(\alpha_s)$  corrections in the cross section the interference term between the one-loop amplitude  $\mathcal{M}^{1L}$  and the Born amplitude  $\mathcal{M}^B$  needs to be considered. However the virtual amplitude

$$\mathcal{M}^V = \mathcal{M}^B \mathcal{M}^{1L*} + \mathcal{M}^{1L} \mathcal{M}^{B*} = 2\Re(\mathcal{M}^B \mathcal{M}^{1L*}) \quad (6.1)$$

contains divergences. These come from integrals over undetermined momentum and energy of particles running in loops. Due to their origin these divergences are referred to as ultra-



**Figure 6.1:** A selection of diagrams contributing to the 1-loop matrix element for squark production. Each diagram is a representative of a certain diagram type. In the first line an example for a self-energy and a vertex-correction diagram is shown. The second line lists typical box-diagrams: a three-point box and a four-point box.

violet and infrared divergences. Ultraviolet divergences occur when the loop momenta and energy tend to infinity which corresponds to arbitrary short distance interactions. However, these divergences are cured by first regularizing them and then introducing appropriate coun-

terterms<sup>17</sup> in the Lagrangian to cancel the extracted singularities. The second step is called renormalization and can be understood as a redefinition or rescaling of parameters and fields in the Lagrangian in the first place:

$$\begin{aligned} \phi^0/\sigma^0 &\rightarrow \sqrt{Z_{\phi^0/\sigma^0}} \phi^0/\sigma^0 & \tilde{q}_{L/R} &\rightarrow \sqrt{Z_{\tilde{q}_{L/R}}} \tilde{q}_{L/R} \\ P_L \tilde{g} &\rightarrow \sqrt{Z_{\tilde{g}}^L} P_L \tilde{g} & P_R \tilde{g} &\rightarrow \sqrt{Z_{\tilde{g}}^R} P_R \tilde{g} \\ q &\rightarrow \sqrt{Z_q} q & G_\mu &\rightarrow \sqrt{Z_G} G_\mu \end{aligned} \quad (6.2)$$

$$\begin{aligned} g_s &\rightarrow g_s \text{ bare} & m_{\tilde{q}}^2 &\rightarrow m_{\tilde{q}}^2 \text{ bare} \\ m_\sigma^2 &\rightarrow m_\sigma^2 \text{ bare} & m_{\tilde{g}} &\rightarrow m_{\tilde{g}} \text{ bare} \end{aligned} \quad (6.3)$$

In this thesis the computation is performed in  $D = 4 - 2\epsilon$  dimensions in order to regularize it. Ultraviolet divergences show then up as single poles in  $\epsilon$ . This procedure will be discussed in detail in section 7 and 8.1.

Having removed the ultraviolet divergences the matrix element is not free of divergences as it comprises infrared divergences. These split into soft and collinear (or mass) singularities<sup>18</sup> which cannot be removed by means of renormalization. These additional singularities show up as single and double poles of  $\epsilon$  in the matrix element  $\mathcal{M}^{\text{1L}}$  and the virtual cross section which is obtained after performing the 2-body phase space integration as in eq. 5.8 but over  $|\mathcal{M}^{\text{V}}|^2$ :

$$\begin{aligned} \frac{d^2\sigma^{\text{V}}}{dt du} &= \frac{K_{ab}}{s^2} \frac{\pi S_\epsilon}{\Gamma(1-\epsilon)} \left[ \frac{tu - m_1^2 m_2^2}{\mu^2 s} \right]^{-\epsilon} \Theta(tu - m_1^2 m_2^2) \\ &\quad \Theta(s - 4m^2) \delta(s + t + u - m_1^2 - m_2^2) \sum |\mathcal{M}^{\text{V}}|^2. \end{aligned} \quad (6.4)$$

The single poles correspond to soft or collinear divergences whereas the double poles correspond to the coincidence of soft and collinear divergences.

## 6.2 Real Corrections

A crucial thing to recognize in the calculation of the  $\mathcal{O}(\alpha_s)$  correction to the squark production cross section is that not only virtual corrections but also real corrections give rise to this physical observable. This is because one does not measure the final state partons but the jets in which they hadronize in due to confinement. This means that a radiation of a massless particle whose energy tends to zero or the radiation of a particle which is collinear to one of the final state particles is in experiment measured as a two jet event even though there are three particles in the final state of the calculation. However real corrections also contain divergences. Including these contributions to the cross section causes the cancellation of singles

<sup>17</sup>These counterterms have to be of the same  $\mathcal{O}(\hbar)$  as the loop diagrams they are constructed to cancel.  
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<sup>18</sup>The names differ in the literature. Often infrared divergences are used as a synonym for soft divergences.

poles coming from soft divergences and the cancellation of double poles. At this point the only singularities which are left in the cross section are the collinear ones.

To see how these are removed consider diagram the second digram in fig. 6.3. The propagator of the quark which radiates off a gluon is giving a contribution of the form

$$\frac{1}{(p_q + p_g)^2 - m_q^2} = \frac{1}{p_q \cdot p_g} = \frac{1}{2E_q E_g (1 - \beta_q \cos \theta_{qg})} \quad \text{with } \beta_q = \frac{\vec{p}_q}{E_q} \quad (6.5)$$

to the the diagram in question. Where  $p_q, p_g$  ( $E_q, E_g$ ) are the quark's and gluon's four-momentum (energy) and  $\theta_{qg}$  is the angle between the quark's and gluon's three-momentum. One observes two potentially singular limits

$$\text{soft gluon :} \quad E_g \rightarrow 0 \quad (6.6)$$

$$\text{collinear particles :} \quad \theta_{qg} \rightarrow 0 \quad \text{if } m_q \rightarrow 0. \quad (6.7)$$

Within these limits the inner quark can almost be considered as on-shell which means that it may travel a long distance between the emission of a gluon and taking part in the actual interaction. This suggest that some divergences may come from the incoming hadron and are actually not associated with the interaction. In fact this turns out to be the case: The remaining collinear divergences can be absorbed in to a redefinition of the parton density functions. This procedure is referred to as mass factorization. Figure 6.2 summarizes the required steps to render the next-to-leading-order cross section finite.

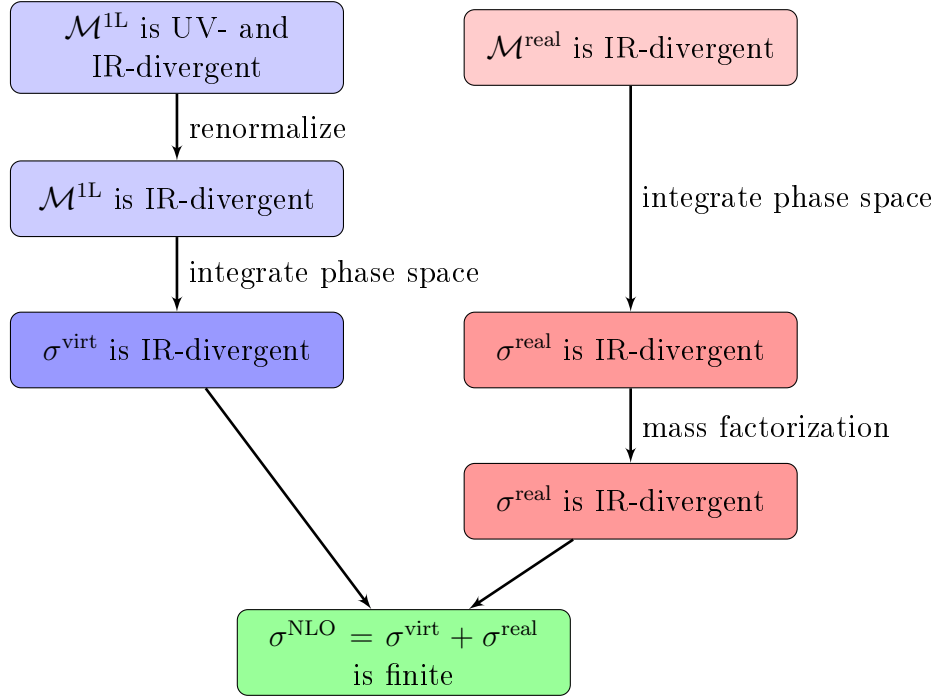
In fact the Kinoshita-Lee-Naunberg-Theorem [21], [22] guarantees the finiteness of sufficiently inclusive observables.

### 6.2.1 Real Gluon Radiation

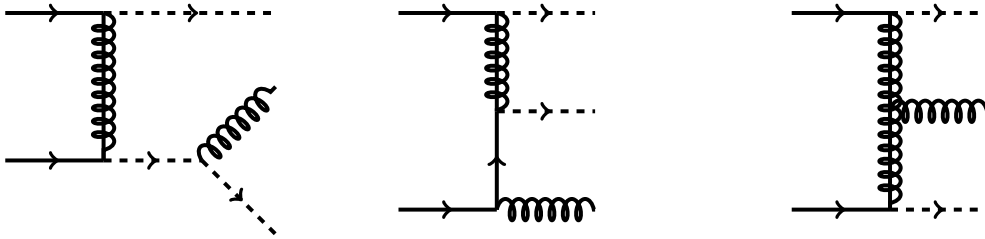
Representative Feynman-diagrams for real gluon emission are shown in fig. 6.3. In order to extract the singularities the so called two-cut method has been adopted. This is the phase space has been divided into different regimes: a soft part, a hard collinear part and a hard non-collinear part.

In the soft part the energy of the outgoing gluon is only integrated up to  $\delta_s \frac{\sqrt{s}}{2}$  where  $\delta_s$  is a cut on which the final result must not depend on.

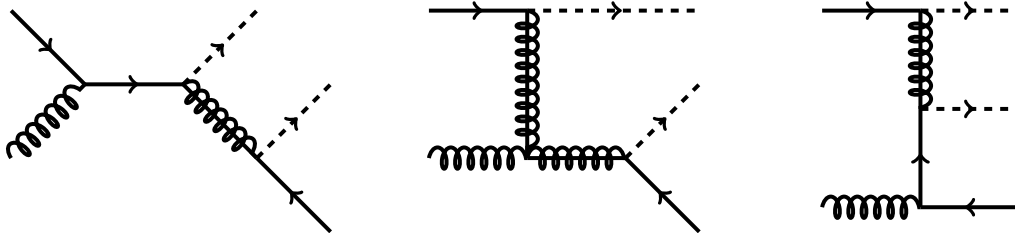
### 6.2.2 Real Quark Radiation



**Figure 6.2:** This scheme illustrates the necessary steps in the computation of the next-to-leading order cross section which are required to render it finite.



**Figure 6.3:** A selection of real gluon emission diagrams contributing to squark production. Again each diagram is a representative of a certain diagram type. The first one is final state gluon emission, the second one initial state gluon radiation and the third one is the radiation of a gluon from a virtual particle within the diagram.



**Figure 6.4:** A selection of real quark emission diagrams contributing to squark production. The first two diagrams have a gluino propagator which may go on-shell. In this regime of phase space these diagrams should not be included in the correction to squark-production but to a tree-level contribution to squark-gluino-production.

## 7 Renormalization of the MRSSM

In order to improve the prediction of the cross section of the previous chapter one has to take quantum corrections into account. These are associated with loops in the corresponding Feynman diagrams. Computing these loop diagrams one might encounter infinities which arise from certain momentum configurations of the unspecified loop momentum. These infinities can be classified due to their origin. Infinities which are associated with loop momenta which tend to infinity are referred to as ultraviolet(UV) divergences. Infinities arising from loop momenta approaching zero can occur in loops with massless particles and are called infrared(IR) singularities. Furthermore there are collinear singularities which occur when a massless particle splits into two massless collinear particles.

These infinities are not physical and must therefore be removed to get sensible predictions. To this end one regularizes them to extract them from the quantity in question. UV-divergences can be removed by means of renormalization, i.e. counterterms are inserted into the Lagrangian to cancel UV-divergences. Infrared and collinear divergences are removed by adding up all possible contributions which give rise to the considered observable.

### 7.1 Regularization Schemes

#### Dimensional Regularization(DREG)

Dimensional regularization(DREG) is a very common procedure for regularizing infinities which was devised by t'Hooft and Veltman [23]. In this scheme loop momenta, gamma- and epsilon-tensors, phase space and fields are defined in  $D$  dimensions. As in every regularization scheme a parameter with mass dimension needs to be introduced. In DREG that is the  $\mu$  parameter which ensures that the loop integrals still have mass dimension 4:

$$\int \frac{d^4 p}{(2\pi)^4} \rightarrow \mu^{4-D} \int \frac{d^D p}{(2\pi)^D}. \quad (7.1)$$

One often writes  $D = 4 - 2\epsilon$ . Then the divergences of the loop integral manifest in  $\frac{1}{\epsilon}$  poles. However DREG suffers a flaw in supersymmetry. As the degrees of freedom for a massless gauge boson are  $D-2$  but the degrees of freedom for its superpartner are 2 there is a mismatch if  $D \neq 4$ . As a consequence there are  $2\epsilon$  degrees of freedom associated with the gluon<sup>19</sup> which do not have a supersymmetric partner. Therefore DREG violates supersymmetry.

---

<sup>19</sup>These degrees of freedom are identified with scalars and are therefore referred to as  $\epsilon$  scalars



## Dimensional Reduction (DRED)

Dimensional reduction (DRED) was introduced to rectify the imperfections of DREG, i.e. it preserves supersymmetry<sup>20</sup>. DRED promotes only loop momenta to  $D$  dimensions. All other quantities which are  $D$  dimensional in DREG stay in 4 dimensions.

maybe refer to Collins: Renormalization

## 7.2 Regularization Scheme Dependences

It is useful to introduce the effective action  $\Gamma$  to discuss the subject of this and ensuing subsections. A formal introduction of  $\Gamma$  can be found in [2]. In short  $\Gamma$  can be viewed as a modification of the classical action  $\Gamma_{\text{cl}} = \int \mathcal{L}_{\text{cl}}$  by quantum effects:

$$\Gamma = \Gamma_{\text{cl}} + \mathcal{O}(\hbar) \quad (7.2)$$

This means that in addition to the vertices in the classical Lagrangian new vertices arise due to loop effects. As already alluded to loop corrections might a priori not be finite and then need to be rendered finite by the addition of counterterms. For  $\mathcal{O}(\hbar)$  corrections one writes

$$\Gamma^{(\leq 1)} \rightarrow \Gamma^{(\leq 1)} + \Gamma^{(1), \text{ ct}} \quad (7.3)$$

These counterterms depend on the regularization (and renormalization) scheme. If one chooses to work with DREG supersymmetry will not be preserved at 1-loop order, i.e.  $\Gamma_{\text{DREG}}^{(\leq 1)}$  is not supersymmetric. To maintain supersymmetry invariance of the renormalized effective action the counterterms will not only consist of supersymmetric counterterms  $\Gamma_{\text{DREG}}^{(1), \text{ ct, sym}}$  but also of counterterms restoring supersymmetry  $\Gamma_{\text{DREG}}^{(1), \text{ ct, restore}}$ .

$$\Gamma_{\text{DREG}}^{(1), \text{ ct}} = \Gamma_{\text{DREG}}^{(1), \text{ ct, sym}} + \Gamma_{\text{DREG}}^{(1), \text{ ct, restore}} \quad (7.4)$$

Fortunately a supersymmetry conserving regularization scheme (at 1-loop level) is given by DRED [24]. One way to acquire supersymmetry restoring counterterms is therefore given by

$$\Gamma_{\text{DRED}}^{(\leq 1)} + \Gamma_{\text{DRED}}^{(1), \text{ ct}} \stackrel{!}{=} \Gamma_{\text{DRED}}^{(\leq 1)} + \Gamma_{\text{DRED}}^{(1), \text{ ct}}. \quad (7.5)$$

Equating also the finite terms in  $\Gamma^{(1), \text{ ct, sym}}$  in DRED and DREG the choice of the supersymmetry restoring counterterms is fixed by:

$$\Gamma_{\text{DRED}}^{(1), \text{ ct, restore}} = \Gamma_{\text{DRED}}^{(\leq 1)} - \Gamma_{\text{DREG}}^{(\leq 1)}. \quad (7.6)$$

<sup>20</sup>It is not clear if DRED preserves supersymmetry at all orders in perturbation theory but it does preserve supersymmetry at the 1-loop level.

This way supersymmetry is preserved by the renormalization constants.

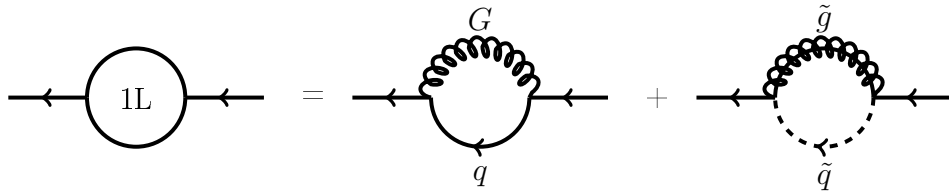
In the case of the MRSSM it will turn out that the only supersymmetry violation comes from corrections associated with the gluon as already alluded to in section 7.1. However supersymmetry restoring will always already be included in  $\delta Z^{\text{DREG}}$ . This is  $\delta Z^{\text{DREG}} = \delta Z^{\text{DREG, sym}} + \delta Z^{\text{trans}}$  where  $\delta Z^{\text{trans}} = \delta Z^{\text{DREG}} - \delta Z^{\text{DRED}}$  is the supersymmetry restoring renormalization constant. The only point where particularly care is required is the coupling: The gauge coupling  $g_s$  and the Yukawa coupling  $\hat{g}_s$  receive different supersymmetry restoring counterterms due to different loop diagrams. Therefore one has to make a difference between these couplings at 1-loop level. In order to match  $g_s$  to the experimentally measured coupling it is renormalized in the  $\overline{\text{MS}}$ -scheme (with an additional manipulation for heavy particles, which will be explained in section 7.4). The Yukawa coupling  $\hat{g}_s$  therefore needs to be added with the difference of the supersymmetry restoring counterterms from the  $q\tilde{q}^\dagger\tilde{g}$  vertex and the  $q\bar{q}G$  vertex at 1-loop in order to be renormalized the same way.

### 7.3 On-Shell Renormalization

A part of the computation of NLO processes is the calculation of renormalization constants. The field and mass renormalization constants have been calculated in DREG in the on-shell scheme. This has the advantage that when turning to the cross section no manipulation of the Green function to the S-matrix element has to be done.

#### The Quark Self-Energy

The quark self-energy splits into contribution from the SM as well as a supersymmetric analogue which is already present in the MSSM.



**Figure 7.1:** Feynman diagrams contributing to the self-energy of the quark at 1-loop level

The 1-PI diagrams evaluate to

$$i\Gamma_{q_i\bar{q}_j}^{1L} = i\frac{g_s^2}{16\pi^2}\delta_{ij}C(F)\left[2\left(B_0(p^2, 0, 0) + B_1(p^2, 0, 0) - \frac{1}{2}\right)\not{p} - 2B_1(p^2, m_{\tilde{g}}^2, m_{\tilde{q}}^2)\not{p}\right]. \quad (7.7)$$

With the counterterm Feynman rule

$$i\Gamma_{q_i\bar{q}_j}^{1L,ct} \hat{=} i \text{---}\leftarrow \text{---}\text{X}\text{---}\leftarrow j \hat{=} i\delta_{ij}\delta Z_q \not{p}$$

and the on-shell renormalization condition

$$\frac{\partial}{\partial \not{p}} \left[ \Re(\Gamma_{q_i\bar{q}_j}^{1L}) + \Gamma_{q_i\bar{q}_j}^{1L,ct} \right]_{p^2=0} = 0 \quad (7.8)$$

where  $\Re(\dots)$  denotes the real part of  $\dots$  one finds

$$\delta Z_q = 2C(F) \frac{g_s^2}{16\pi^2} \Re [B_1(0, m_{\tilde{g}}^2, m_{\tilde{q}}^2)] . \quad (7.9)$$

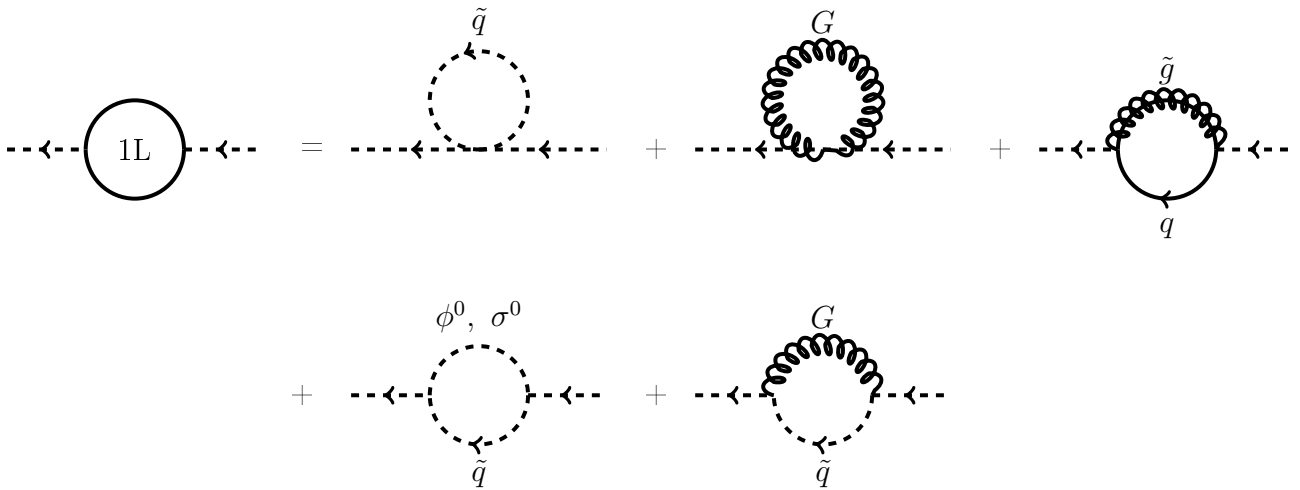
Note that the term  $\frac{1}{2}$  from eq. 7.7 has vanished because the self energy is taken on-shell which means that this factor is absorbed into the on-shell  $B_0$  function which is in DREG defined to be zero if it has no mass scale.

Doing the same calculation in DRED one finds that this very term  $\frac{1}{2}$  is absent. Therefore the transition counterterm between DREG and DRED is given by

$$\delta Z_q^{\text{trans}} = \delta Z_q^{\text{DREG}} - \delta Z_q^{\text{DRED}} = C(F) \frac{g_s^2}{16\pi^2} . \quad (7.10)$$

### The Squark Self-Energy

The contributions to the self-energy of the left- and right-handed squark are the same. Therefore to avoid unnecessary labeling  $\Gamma_{\tilde{q}\tilde{q}^\dagger}$  stands in the following for  $\Gamma_{\tilde{q}_L\tilde{q}_L^\dagger} = \Gamma_{\tilde{q}_R\tilde{q}_R^\dagger}$ .

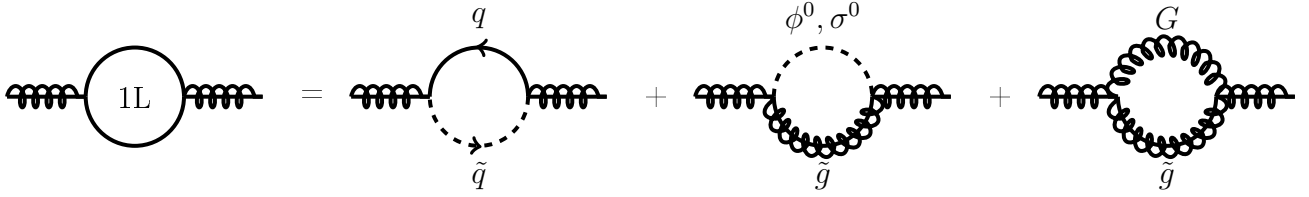


**Figure 7.2:** Diagrammatic contributions to the self-energy of the squark at 1-loop level



the gluino and may be referred to as the octino.

From the Lagrangian 4.16 one can see that the couplings of the two particles are quite distinct. This is reflected by different field renormalization constants of the left- and right-handed part of the gluino. As for the quarks the fermion (and momentum) flow is from the right to the



**Figure 7.3:** diagrammatic contributions to the self-energy of the squark at 1-loop level

left.

$$\begin{aligned}
 i\Gamma_{\tilde{g}^a \tilde{g}^b}^{1L} &= i \frac{g_s^2}{16\pi^2} \delta_{ab} \left[ -4T(F) \left( (n_f - 1)B_1(p^2, 0, m_{\tilde{q}}^2) + B_1(p^2, m_t^2, m_{\tilde{q}}^2) \right) P_L \not{p} \right. \\
 &\quad + C(A) \left( (B_0(p^2, m_{\tilde{g}}^2, m_{\phi^0}^2) - B_0(p^2, m_{\tilde{g}}^2, m_{\sigma^0}^2))m_{\tilde{g}} - (B_1(p^2, m_{\tilde{g}}^2, m_{\phi^0}^2) + B_1(p^2, m_{\tilde{g}}^2, m_{\sigma^0}^2))\not{p} \right) \\
 &\quad \left. + C(A) \left( (2 - 4B_0(p^2, 0, m_{\tilde{g}}^2))m_{\tilde{g}} - (1 - 2(B_0(p^2, 0, m_{\tilde{g}}^2) + B_1(p^2, 0, m_{\tilde{g}}^2)))\not{p} \right) \right] \quad (7.17)
 \end{aligned}$$

Where  $n_f = 6$  is the number of quark flavors.

The counterterm Feynman rule reads

$$i\Gamma_{\tilde{g}^a \tilde{g}^b}^{1L,ct} \hat{=} a \text{ (crossed wavy lines) } b \hat{=} i\delta_{ab} \left[ (\delta Z_{\tilde{g}}^L P_L + \delta Z_{\tilde{g}}^R P_R) \not{p} - \left( \frac{\delta Z_{\tilde{g}}^L + \delta Z_{\tilde{g}}^R}{2} m_{\tilde{g}} + \delta m_{\tilde{g}} \right) \right].$$

The on-shell renormalization conditions for the fields are

$$\frac{\partial}{\partial(P_L \not{p})} \left[ \Re(\Gamma_{\tilde{g}^a \tilde{g}^b}^{1L}) + \Gamma_{\tilde{g}^a \tilde{g}^b}^{1L,ct} \right]_{\not{p}=m_{\tilde{g}}} = 0, \quad \frac{\partial}{\partial(P_R \not{p})} \left[ \Re(\Gamma_{\tilde{g}^a \tilde{g}^b}^{1L}) + \Gamma_{\tilde{g}^a \tilde{g}^b}^{1L,ct} \right]_{\not{p}=m_{\tilde{g}}} = 0 \quad (7.18)$$

where the derivative of  $\Sigma = \Sigma^{VL} P_L \not{p} + \Sigma^{VR} P_R \not{p} + \Sigma^{SL} P_L + \Sigma^{SR} P_R$  with respect to  $P_A \not{p}$  ( $A \in \{L, R\}$ ) is defined by

$$\left. \frac{\partial}{\partial(P_A \not{p})} \Sigma \right|_{\not{p}=m} = \Sigma^{VA} + \frac{\partial}{\partial p^2} (m^2 \Sigma^{VL} + m^2 \Sigma^{VR} + m \Sigma^{SL} + m \Sigma^{SR}). \quad (7.19)$$

This leads to the following renormalization constants

$$\begin{aligned}
\delta Z_{\tilde{g}}^L = & \frac{g_s^2}{16\pi^2} \Re \left[ 4T(F) \left( (n_f - 1) B_1(m_{\tilde{g}}^2, 0, m_{\tilde{q}}^2) + B_1(m_{\tilde{g}}^2, m_t^2, m_{\tilde{q}}^2) \right) \right. \\
& + C(A) (B_1(m_{\tilde{g}}^2, m_{\tilde{g}}^2, m_{\phi^0}^2) + B_1(m_{\tilde{g}}^2, m_{\tilde{g}}^2, m_{\sigma^0}^2)) \\
& + C(A) (1 - 2(B_0(m_{\tilde{g}}^2, 0, m_{\tilde{g}}^2) + B_1(m_{\tilde{g}}^2, 0, m_{\tilde{g}}^2))) \\
& + 4T(F) m_{\tilde{g}}^2 \frac{\partial}{\partial p^2} \left( ((n_f - 1)) B_1(p^2, 0, m_{\tilde{q}}^2) + B_1(p^2, m_t^2, m_{\tilde{q}}^2) \right) \\
& - 2C(A) m_{\tilde{g}} \frac{\partial}{\partial p^2} \left( B_0(p^2, m_{\tilde{g}}^2, m_{\phi^0}^2) - B_0(p^2, m_{\tilde{g}}^2, m_{\sigma^0}^2) - B_1(p^2, m_{\tilde{g}}^2, m_{\phi^0}^2) - B_1(p^2, m_{\tilde{g}}^2, m_{\sigma^0}^2) \right) \\
& \left. - 4C(A) m_{\tilde{g}}^2 \frac{\partial}{\partial p^2} \left( -B_0(p^2, 0, m_{\tilde{g}}^2) + B_0(p^2, 0, m_{\tilde{g}}^2) \right) \right]_{p^2=m_{\tilde{g}}^2} \quad (7.20)
\end{aligned}$$

and

$$\begin{aligned}
\delta Z_{\tilde{g}}^R = & \frac{g_s^2}{16\pi^2} \Re \left[ C(A) (B_1(m_{\tilde{g}}^2, m_{\tilde{g}}^2, m_{\phi^0}^2) + B_1(m_{\tilde{g}}^2, m_{\tilde{g}}^2, m_{\sigma^0}^2)) \right. \\
& + C(A) (1 - 2(B_0(m_{\tilde{g}}^2, 0, m_{\tilde{g}}^2) + B_1(m_{\tilde{g}}^2, 0, m_{\tilde{g}}^2))) \\
& + 4T(F) m_{\tilde{g}}^2 \frac{\partial}{\partial p^2} \left( ((n_f - 1)) B_1(p^2, 0, m_{\tilde{q}}^2) + B_1(p^2, m_t^2, m_{\tilde{q}}^2) \right) \\
& - 2C(A) m_{\tilde{g}} \frac{\partial}{\partial p^2} \left( B_0(p^2, m_{\tilde{g}}^2, m_{\phi^0}^2) - B_0(p^2, m_{\tilde{g}}^2, m_{\sigma^0}^2) - B_1(p^2, m_{\tilde{g}}^2, m_{\phi^0}^2) - B_1(p^2, m_{\tilde{g}}^2, m_{\sigma^0}^2) \right) \\
& \left. - 4C(A) m_{\tilde{g}}^2 \frac{\partial}{\partial p^2} \left( -B_0(p^2, 0, m_{\tilde{g}}^2) + B_0(p^2, 0, m_{\tilde{g}}^2) \right) \right]_{p^2=m_{\tilde{g}}^2} . \quad (7.21)
\end{aligned}$$

As for the quark there are constant terms amid the Passarino-Veltman integrals. These arise only in DREG and not in DRED. The transition counterterms are

$$\delta Z_{\tilde{g}}^{A \text{ trans}} = \delta Z_{\tilde{g}}^{A \text{ DREG}} - \delta Z_{\tilde{g}}^{A \text{ DRED}} = C(A) \frac{g_s^2}{16\pi^2} \quad (7.22)$$

for  $A \in \{L, R\}$ . The gluino mass counterterm is ascertained by the condition

$$\left[ \Re(\Gamma_{\tilde{g}^a \tilde{g}^b}^{1L}) + \Gamma_{\tilde{g}^a \tilde{g}^b}^{1L, \text{ct}} \right]_{p=m_{\tilde{g}}} = 0 \quad (7.23)$$

which is equivalent to

$$\delta m_{\tilde{g}} = \Re \left( m_{\tilde{g}} \frac{\Sigma^{VL} + \Sigma^{VR}}{2} + \frac{\Sigma^{SL} + \Sigma^{SR}}{2} \right) \quad (7.24)$$

and yields

$$\begin{aligned} \delta m_{\tilde{g}} = & \frac{g_s^2}{16\pi^2} m_{\tilde{g}}^2 \Re \left[ -2T(F) \left( (n_f - 1) B_1(m_{\tilde{g}}^2, 0, m_{\tilde{q}}^2) + B_1(m_{\tilde{g}}^2, m_t^2, m_{\tilde{q}}^2) \right) \right. \\ & + C(A) \left( B_0(m_{\tilde{g}}^2, m_{\tilde{g}}^2, m_{\phi^0}^2) - B_0(m_{\tilde{g}}^2, m_{\tilde{g}}^2, m_{\sigma^0}^2) - B_1(m_{\tilde{g}}^2, m_{\tilde{g}}^2, m_{\phi^0}^2) - B_1(m_{\tilde{g}}^2, m_{\tilde{g}}^2, m_{\sigma^0}^2) \right) \\ & \left. + C(A) \left( 1 - 2B_0(m_{\tilde{g}}^2, 0, m_{\tilde{g}}^2) + 2B_1(m_{\tilde{g}}^2, 0, m_{\tilde{g}}^2) \right) \right]. \end{aligned} \quad (7.25)$$

Again there is a transition counterterm

$$\delta m_{\tilde{g}}^{\text{trans}} = \delta m_{\tilde{g}}^{\text{DREG}} - \delta m_{\tilde{g}}^{\text{DRED}} = C(A) \frac{g_s^2}{16\pi^2} m_{\tilde{g}}. \quad (7.26)$$

## 7.4 Renormalization of the Gauge Coupling

The gauge coupling  $g_s$  is renormalized in the  $\overline{\text{MS}}$ -scheme with the modification that additional logarithms are subtracted, i.e. light particles are treated in the  $\overline{\text{MS}}$ -scheme and heavy particles in the zero-momentum subtraction scheme. This is to decouple heavy particles from the running of  $\alpha_s = \frac{g_s^2}{4\pi}$ . This renormalization procedure allows to adopt the experimental values of  $\alpha_s$  from the parton density functions. The running due to effects of heavy particles is then encoded in the logarithms of  $\delta g_s$ .

Extracting  $\delta g_s$  from the quark-quark-gluon vertex requires not only the computation of  $i\Gamma_{q_i\bar{q}_j G_a^\mu}^{\text{1L}}$  but also the (re)evaluation of auxiliary field renormalization constants  $\delta Z_q^{\text{aux}}$  and  $\delta Z_G^{\text{aux}}$  in the above mentioned scheme. These will not be the same as the on-shell scheme.

### The Quark Self-Energy Revisited

The quark self-energy has two contribution which are shown in figure 7.1. The first one corresponds to light particles the second one to heavy particles:  $i\Gamma_{q_i\bar{q}_j}^{\text{1L}} = i\Gamma_{q_i\bar{q}_j}^{\text{1L,light}} + i\Gamma_{q_i\bar{q}_j}^{\text{1L,heavy}}$ . For light particles only the UV-divergent<sup>21</sup> part is kept. The self energy corresponding to the heavy particles is taken at zero momentum  $p^2 = 0$ .

$$i\Gamma_{q_i\bar{q}_j}^{\text{1L,light}} \Big|_{\text{UV-div}} = iC(F) \frac{g_s^2}{16\pi^2} \Delta_\epsilon \not{p} \delta_{ij} \quad (7.27)$$

$$i\Gamma_{q_i\bar{q}_j}^{\text{1L,heavy}}(p^2 = 0) = -iC(F) \frac{g_s^2}{16\pi^2} 2B_1(0, m_{\tilde{g}}^2, m_{\tilde{q}}^2) \not{p} \delta_{ij} \quad (7.28)$$

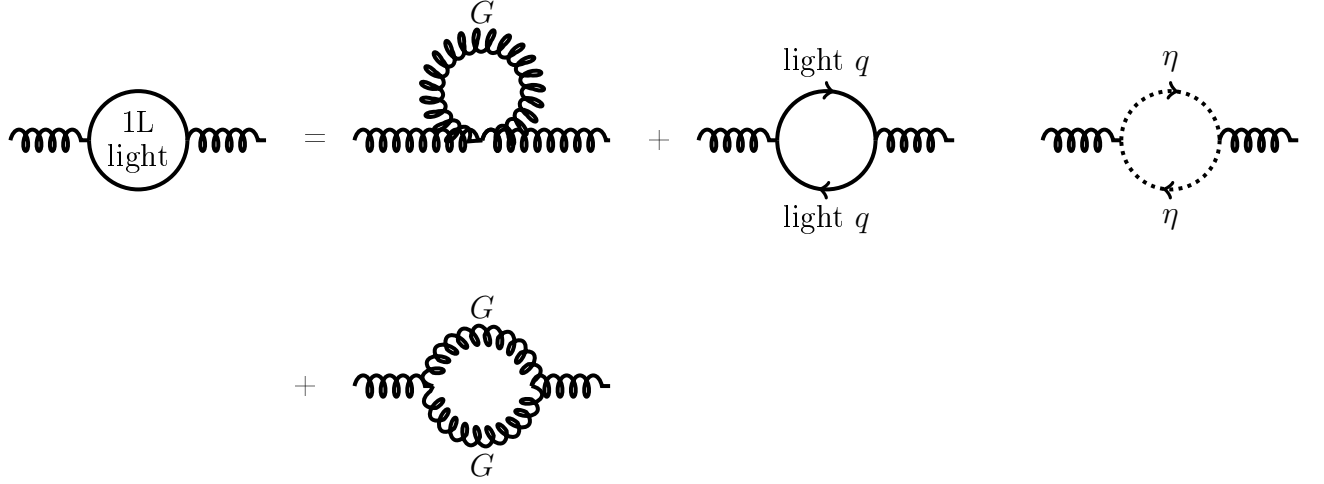
The renormalization constant for the evaluation of  $\delta g_s$  is determined by the condition

$$\frac{\partial}{\partial \not{p}} \left[ \Gamma_{q_i\bar{q}_j}^{\text{1L,light}} \Big|_{\text{UV-div}} + \Gamma_{q_i\bar{q}_j}^{\text{1L,heavy}}(p^2 = 0) + \Gamma_{q_i\bar{q}_j}^{\text{1L,ct}} \right] = 0 \quad (7.29)$$

<sup>21</sup>In contrast to the  $\overline{\text{MS}}$ -scheme in the  $\overline{\text{MS}}$ -scheme not only the pure ultraviolet divergence but also two additional transcendent numbers are subtracted. It is therefore common to define  $\Delta_\epsilon = \frac{1}{\epsilon} - \gamma_E + \ln 4\pi$ .

and computes to

$$\delta Z_q^{\text{aux}} = \frac{g_s^2}{16\pi^2} C(F) \left[ -\Delta_\epsilon + 2B_1(0, m_{\tilde{g}}^2, m_{\tilde{q}}^2) \right]. \quad (7.30)$$



**Figure 7.4:** Contribution to the self-energy of the gluon originating from light particles

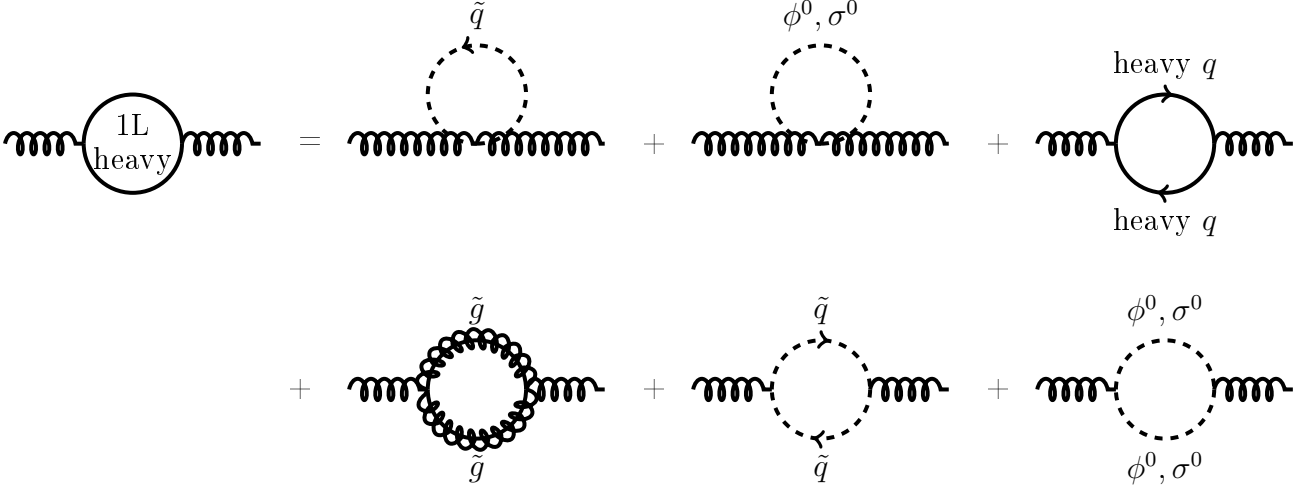
### The Gluon Self-Energy

As for the quark self-energy there are again contributions to the self-energy originating from light and heavy particles. Again these are differently dealt with.

$$\begin{aligned} i\Gamma_{G_\mu^a G_\nu^b}^{\text{1L,light}} \Big|_{\text{UV-div}} &= i \frac{g_s^2}{16\pi^2} \Delta_\epsilon \left[ 0 - \frac{4(n_f - 1)}{3} T(F) (p^2 g^{\mu\nu} - p^\mu p^\nu) + \frac{C(A)}{12} (p^2 g^{\mu\nu} + 2p^\mu p^\nu) \right. \\ &\quad \left. + \frac{C(A)}{12} (19p^2 g^{\mu\nu} - 22p^\mu p^\nu) \right] \delta_{ab} \\ &= i \frac{g_s^2}{16\pi^2} \Delta_\epsilon \left[ -\frac{4(n_f - 1)}{3} T(F) + \frac{5}{3} C(A) \right] (p^2 g^{\mu\nu} - p^\mu p^\nu) \delta_{ab} \end{aligned} \quad (7.31)$$

Both the contribution from the gluon loop and the one from the ghost loop are not proportional to  $(p^2 g^{\mu\nu} - p^\mu p^\nu)$  but their sum is. The heavy particle contributions are calculated in the zero-





**Figure 7.5:** Contribution to the self-energy of the gluon originating from heavy particles, in the last diagram either  $\phi^0$  or  $\sigma^0$  are running in the loop

momentum subtraction

$$\begin{aligned}
i\Gamma_{G_\mu^a G_\nu^b}^{1L, \text{heavy}}(p^2 = 0) &= i\frac{g_s^2}{16\pi^2}\delta_{ab} \left[ -4T(F)n_f \left( \Delta_\epsilon - \ln \frac{m_{\tilde{q}}^2}{\mu^2} \right) m_{\tilde{q}}^2 g^{\mu\nu} - C(A) \left( \Delta_\epsilon - \ln \frac{m_{\phi^0}^2}{\mu^2} \right) m_{\phi^0}^2 g^{\mu\nu} \right. \\
&\quad - C(A) \left( \Delta_\epsilon - \ln \frac{m_{\sigma^0}^2}{\mu^2} \right) m_{\sigma^0}^2 g^{\mu\nu} - \frac{4}{3}T(F) \left( \Delta_\epsilon - \ln \frac{m_t^2}{\mu^2} \right) (p^2 g^{\mu\nu} - p^\mu p^\nu) \\
&\quad - \frac{4}{3}C(A) \left( \Delta_\epsilon - \ln \frac{m_{\tilde{g}}^2}{\mu^2} \right) (p^2 g^{\mu\nu} - p^\mu p^\nu) \\
&\quad - \frac{2}{3}T(F)n_f \left( \Delta_\epsilon - \ln \frac{m_{\tilde{q}}^2}{\mu^2} \right) (p^2 g^{\mu\nu} - p^\mu p^\nu) + 4T(F)n_f \left( \Delta_\epsilon - \ln \frac{m_{\tilde{q}}^2}{\mu^2} \right) m_{\tilde{q}}^2 g^{\mu\nu} \\
&\quad - \frac{1}{6}C(A) \left( \Delta_\epsilon - \ln \frac{m_{\phi^0}^2}{\mu^2} \right) (p^2 g^{\mu\nu} - p^\mu p^\nu) + C(A) \left( \Delta_\epsilon - \ln \frac{m_{\phi^0}^2}{\mu^2} \right) m_{\phi^0}^2 g^{\mu\nu} \\
&\quad \left. - \frac{1}{6}C(A) \left( \Delta_\epsilon - \ln \frac{m_{\sigma^0}^2}{\mu^2} \right) (p^2 g^{\mu\nu} - p^\mu p^\nu) + C(A) \left( \Delta_\epsilon - \ln \frac{m_{\sigma^0}^2}{\mu^2} \right) m_{\sigma^0}^2 g^{\mu\nu} \right] \\
&= i\frac{g_s^2}{16\pi^2}\delta_{ab} \left[ -\frac{4}{3}T(F) \left( \Delta_\epsilon - \ln \frac{m_t^2}{\mu^2} \right) (p^2 g^{\mu\nu} - p^\mu p^\nu) \right. \\
&\quad - \frac{4}{3}C(A) \left( \Delta_\epsilon - \ln \frac{m_{\tilde{g}}^2}{\mu^2} \right) (p^2 g^{\mu\nu} - p^\mu p^\nu) \\
&\quad - \frac{2}{3}T(F)n_f \left( \Delta_\epsilon - \ln \frac{m_{\tilde{q}}^2}{\mu^2} \right) (p^2 g^{\mu\nu} - p^\mu p^\nu) \\
&\quad - \frac{1}{6}C(A) \left( \Delta_\epsilon - \ln \frac{m_{\phi^0}^2}{\mu^2} \right) (p^2 g^{\mu\nu} - p^\mu p^\nu) \\
&\quad \left. - \frac{1}{6}C(A) \left( \Delta_\epsilon - \ln \frac{m_{\sigma^0}^2}{\mu^2} \right) (p^2 g^{\mu\nu} - p^\mu p^\nu) \right] \tag{7.32}
\end{aligned}$$

The counterterm Feynman rule for the gluon propagator is

$$i\Gamma_{G_\mu^a G_\nu^b}^{1L,ct} \hat{=} a, \mu \text{---} \text{---} \text{---} b, \nu \hat{=} -i\delta Z_G (p^2 g^{\mu\nu} - p^\mu p^\nu) \delta_{ab}.$$

The renormalization condition for  $\delta Z_G^{\text{aux}}$  reads

$$\Gamma_{G_\mu^a G_\nu^b}^{1L,\text{light}} \Big|_{\text{UV-div}} + \Gamma_{G_\mu^a G_\nu^b}^{1L,\text{heavy}}(p^2 = 0) - \delta Z_G^{\text{aux}} (p^2 g^{\mu\nu} - p^\mu p^\nu) \delta_{ab} = 0 \quad (7.33)$$

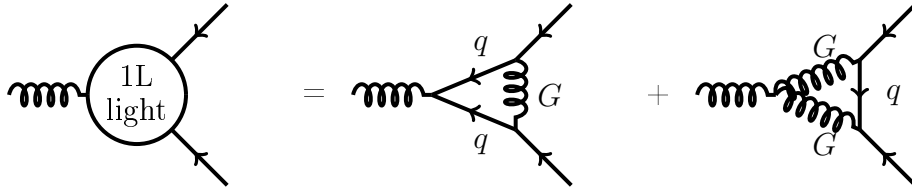
and yields

$$\begin{aligned} \delta Z_G^{\text{aux}} = \frac{g_s^2}{16\pi^2} \Big\{ & \left[ -\frac{4}{3}T(F)(n_f - 1) + \frac{5}{3}C(A) \right] \Delta_\epsilon + \left[ -\frac{4}{3}T(F) \left( \Delta_\epsilon - \ln \frac{m_t^2}{\mu^2} \right) \right. \\ & - \frac{4}{3}C(A) \left( \Delta_\epsilon - \ln \frac{m_g^2}{\mu^2} \right) - \frac{2}{3}T(F)n_f \left( \Delta_\epsilon - \ln \frac{m_q^2}{\mu^2} \right) \\ & \left. - \frac{1}{6}C(A) \left( \Delta_\epsilon - \ln \frac{m_{\phi^0}^2}{\mu^2} \right) - \frac{1}{6}C(A) \left( \Delta_\epsilon - \ln \frac{m_{\sigma^0}^2}{\mu^2} \right) \right] \Big\} \end{aligned} \quad (7.34)$$

where the first squared brackets come from light and the second one from heavy particles running in the loop.

### The $q\bar{q}G$ Vertex Correction

As for the self energies the vertex correction constitutes of loops from light particles which correspond to the Standard Model and additional heavy particle loops. The contributions

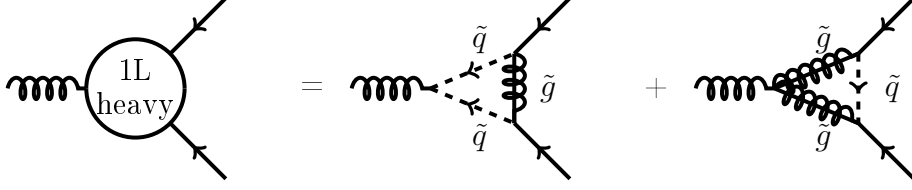


**Figure 7.6:** Contribution from light (Standard Model) particles to the  $q\bar{q}G$  vertex correction

from Standard Model particles to the  $q\bar{q}G$  yield

$$\begin{aligned} i\Gamma_{q_i \bar{q}_j G_\mu^a}^{1L,\text{light}} \Big|_{\text{UV-div}} &= -ig_s T_{ij}^a \gamma^\mu \frac{g_s^2}{16\pi^2} \left[ \left( C(F) - \frac{C(A)}{2} \right) 2(B_0(0,0,0) - 2C_{00}(0,0,0,0,0)) \right. \\ &\quad \left. + C(A)(B_0(0,0,0) + 2C_{00}(0,0,0,0,0)) \right] \\ &= -ig_s T_{ij}^a \gamma^\mu \frac{g_s^2}{16\pi^2} \Delta_\epsilon \left[ \left( C(F) - \frac{C(A)}{2} \right) + \frac{3}{2}C(A) \right]. \end{aligned} \quad (7.35)$$

The term in the curved brackets of eq. 7.35 corresponds to the first diagram in fig. 7.6 whereas the term with the prefactor  $C(A)$  corresponds to the second diagram. The 3-point-function



**Figure 7.7:** Contribution from heavy particles to the  $q\bar{q}G$  vertex correction.

with only heavy particles running in the loop evaluates to

$$\begin{aligned}
 i\Gamma_{q_i\bar{q}_j G_\mu^a}^{1L, \text{heavy}}(p_i^2 = 0) &= -ig_s T_{ij}^a \gamma^\mu \frac{g_s^2}{16\pi^2} \left[ \left( C(F) - \frac{C(A)}{2} \right) 4C_{00}(0, 0, 0, m_{\tilde{g}}^2, m_{\tilde{q}}^2, m_{\tilde{q}}^2) \right. \\
 &\quad \left. + C(A) (B_0(0, m_{\tilde{g}}^2, m_{\tilde{q}}^2) - 2C_{00}(0, 0, 0, m_{\tilde{g}}^2, m_{\tilde{g}}^2, m_{\tilde{q}}^2)) \right] \\
 &= -ig_s T_{ij}^a \gamma^\mu \frac{g_s^2}{16\pi^2} 2C(F) [B_0(0, m_{\tilde{g}}^2, m_{\tilde{q}}^2) + B_1(0, m_{\tilde{g}}^2, m_{\tilde{q}}^2)] \quad (7.36)
 \end{aligned}$$

The argument  $p_i^2 = 0$  of the vertex function denotes, that all external particles are taken on shell, i.e. at zero momentum.

In the second line identities from 10.6 have been used. The renormalization condition for the counterterm of the gauge coupling  $\delta g_s$  reads

$$i\Gamma_{q_i\bar{q}_j G_\mu^a}^{1L, \text{light}} \Big|_{\text{UV-div}} + i\Gamma_{q_i\bar{q}_j G_\mu^a}^{1L, \text{heavy}}(p_i^2 = 0) + \left[ -ig_s T_{ij}^a \gamma^\mu \left( \frac{\delta g_s}{g_s} + \delta Z_q^{\text{aux}} + \frac{\delta Z_G^{\text{aux}}}{2} \right) \right] = 0. \quad (7.37)$$

Finally one can read off  $\frac{\delta g_s}{g_s}$

$$\begin{aligned}
 \frac{\delta g_s}{g_s} &= \frac{g_s^2}{16\pi^2} \left[ \left( \frac{2}{3}T(F)(n_f - 1) - \frac{11}{6}C(A) \right) \Delta_\epsilon + \left( \frac{5}{6}C(A) + \frac{2}{3}T(F) + \frac{1}{3}T(F)n_f \right) \Delta_\epsilon \right. \\
 &\quad \left. - \frac{2}{3}C(A) \ln \frac{m_{\tilde{g}}^2}{\mu^2} - \frac{1}{3}T(F)n_f \ln \frac{m_{\tilde{q}}^2}{\mu^2} - \frac{2}{3}T(F) \ln \frac{m_t^2}{\mu^2} - \frac{1}{12}C(A) \left( \ln \frac{m_{\phi^0}^2}{\mu^2} + \ln \frac{m_{\sigma^0}^2}{\mu^2} \right) \right]. \quad (7.38)
 \end{aligned}$$

Note the only difference to the MSSM [16] is the appearance of the sgluon masses as well as a factor of two in front of the logarithm of the gluino mass which is due to twice as much degrees of freedom of the gluino in the MRSSM.

## 7.5 The Beta Function

The beta function describes the dependence of the gauge coupling  $g_s$  upon the energy scale  $\mu$ . Writing down the action of a theory in  $D$  dimensions one needs to introduce an energy scale  $\mu$  in order to keep the action dimensionless. But  $\mu$  is no physical parameter and can be absorbed

into the fields and parameters. To this end one defines

$$g_{sB}(g_s, \mu) = \mu^\epsilon g_s \left( 1 + \frac{\delta g_s}{g_s} \right) \quad (7.39)$$

which must not depend upon the unphysical scale  $\mu$ , ergo

$$0 = \frac{dg_{sB}}{d \ln \mu} = \frac{\partial g_{sB}}{\partial \ln \mu} + \beta \frac{\partial g_{sB}}{\partial g_s} \quad (7.40)$$

where the definition of the beta function  $\frac{\partial g_s}{\partial \ln \mu}$  has been inserted. Equation 7.40 serves to calculate  $\beta(g_s, \epsilon)$ . By equating coefficients and using the shortcuts

$$\begin{aligned} \frac{\beta_0^L}{2} &= \frac{2}{3}T(F)(n_f - 1) - \frac{11}{6}C(A) \\ \frac{\beta_0^H}{2} &= \frac{5}{6}C(A) + \frac{2}{3}T(F) + \frac{1}{3}T(F)n_f \\ L &= -\frac{2}{3}C(A) \ln \frac{m_{\tilde{g}}^2}{\mu^2} - \frac{1}{3}T(F)n_f \ln \frac{m_{\tilde{q}}^2}{\mu^2} - \frac{2}{3}T(F) \ln \frac{m_t^2}{\mu^2} - \frac{1}{12}C(A) \left( \ln \frac{m_{\phi^0}^2}{\mu^2} + \ln \frac{m_{\sigma^0}^2}{\mu^2} \right) \end{aligned}$$

so that

$$\frac{\delta g_s}{g_s} = \frac{g_s^2}{16\pi^2} \left( \frac{\beta_0^L}{2\epsilon_{UV}} + \frac{\beta_0^H}{2\epsilon_{UV}} + L \right) \quad (7.41)$$

one finds

$$\beta(g_s, \epsilon) = -\epsilon g_s \left( 1 + \frac{g_s^2}{16\pi^2} L \right) + \beta(g_s) + \mathcal{O}(2 - \text{loop}) \quad (7.42)$$

$$\beta(g_s) = \frac{g_s^3}{16\pi^2} \beta_0^L + \mathcal{O}(2 - \text{loop}). \quad (7.43)$$

This is the beta function of QCD first found by [Gross, Politzer, Wil]

## 7.6 Supersymmetry Restoring Counterterm

As already discussed in section 7.2 care is required in terms of supersymmetry restoring when renormalizing the gauge coupling  $g_s$  and the Yukawa coupling  $\hat{g}_s$ . In doing so one needs the already calculated transition counterterms of the quark, squark and gluino from 7.10, 7.15 and 7.22 as well as the transition counterterm of the gluon.

### The Gluon Self-Energy Revisited

The only regularization dependence of the gluon self-energy arises from the gluon loop, i.e. the last diagram in figure 7.4. With the definition of  $\Gamma_{\text{DREG}}^{(1),\text{ct},\text{restore}}$  in 7.6 one obtains

$$i\Gamma_{\text{DREG},G_\mu^a G_\nu^b}^{(1),\text{ct},\text{restore}} = -i\frac{1}{3}C(A)\frac{g_s^2}{16\pi^2}(p^2 g^{\mu\nu} - p^\mu p^\nu)\delta_{ab} \quad (7.44)$$

which translates to the transition counterterm

$$\delta Z_G^{\text{trans}} = \frac{C(A)}{3} \frac{g_s^2}{16\pi^2}. \quad (7.45)$$

### The $q\bar{q}G$ Vertex Correction Revisited

The diagrams from which the supersymmetry restoring contributions to the gauge coupling correction come from are shown in figure 7.6 and evaluate to

$$i\Gamma_{\text{DREG},q_i\bar{q}_j G_\mu^a}^{(1),\text{ct},\text{restore}} = -ig_s T_{ij}^a \gamma^\mu \frac{g_s^2}{16\pi^2} \left[ \left( C(F) - \frac{C(A)}{2} \right) + \frac{C(A)}{2} \right] \quad (7.46)$$

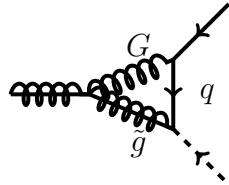
$$= -ig_s T_{ij}^a \gamma^\mu \left[ \frac{\delta g_s^{\text{trans}}}{g_s} + \delta Z_q^{\text{trans}} + \frac{\delta Z_G^{\text{trans}}}{2} \right] \quad (7.47)$$

where the second line shows the supersymmetry restoring counterterms. This yields

$$\frac{\delta g_s^{\text{trans}}}{g_s} = -\frac{C(A)}{6} \frac{g_s^2}{16\pi^2}. \quad (7.48)$$

### The $q\bar{q}^\dagger\tilde{g}$ Vertex Correction

The supersymmetry restoring corrections to the Yukawa coupling origin from the diagram below. The supersymmetry restoring part is given by



**Figure 7.8:** diagram of the supersymmetry restoring correction of the  $q\bar{q}\tilde{g}$  vertex

$$i\Gamma_{\text{DREG},q_i\bar{q}_j\tilde{g}^a}^{(1),\text{ct},\text{restore}} = -ig_s\sqrt{2}P_L T_{ij}^a \frac{g_s^2}{16\pi^2} C(A) \quad (7.49)$$

$$= -ig_s\sqrt{2}P_L T_{ij}^a \left[ \frac{\delta \hat{g}_s^{\text{trans}}}{g_s} + \frac{\delta Z_q^{\text{trans}} + \delta Z_{\tilde{q}}^{\text{trans}} + \delta Z_{\tilde{g}}^{\text{trans}}}{2} \right]. \quad (7.50)$$

The supersymmetry restoring part of the Yukawa renormalization constants is therefore

$$\frac{\delta \hat{g}_s^{\text{trans}}}{g_s} = -\frac{C(F) - C(A)}{2} \frac{g_s^2}{16\pi^2}. \quad (7.51)$$

As a consequence of the two different supersymmetry restoring parts of the coupling renormalization constants an additional renormalization constant  $\delta g_s^{\text{restore}}$  needs to be introduced. As described in section 7.2 it is given by

$$\frac{\delta g_s^{\text{restore}}}{g_s} = \frac{\delta \hat{g}_s^{\text{trans}}}{g_s} - \frac{\delta g_s^{\text{trans}}}{g_s} = \frac{g_s^2}{16\pi^2} \left( \frac{2C(A)}{3} - \frac{C(F)}{2} \right). \quad (7.52)$$

In short this means that the gauge coupling  $g_s$  is renormalized with  $\delta g_s$  given in 7.38 and the Yukawa coupling  $\hat{g}_s$  is renormalized with  $\delta \hat{g}_s = \delta g_s + \delta g_s^{\text{restore}}$ .

The finite correction  $\delta g_s^{\text{restore}}$  is the same as in supersymmetric QCD which should not surprise too much as all its contributions originate from loops with gluons.

## 7.7 $\overline{\text{MS}}$ - Renormalization

To quickly check for UV-finiteness of a Greenfunction or a physical observable it is useful to have the extracted UV-divergences of the renormalization constants at hand. How these are obtained is described in 10.6.

(In order to obtain these the Passarino-Veltman integrals need to be substituted by their  $\frac{1}{\epsilon}$  coefficient. These had been taken from [25] and checked with `FeynArts` and `FormCalc` [26], [?], [26].)

$$\begin{aligned} \delta Z_{\tilde{g}}^L &= -\frac{g_s^2}{16\pi^2\epsilon_{\text{UV}}} 2 [T(F)n_f + C(A)] & \delta Z_{\tilde{g}}^R &= -\frac{g_s^2}{16\pi^2\epsilon_{\text{UV}}} 2C(A) \\ \delta m_{\tilde{g}} &= \frac{g_s^2}{16\pi^2\epsilon_{\text{UV}}} [T(F)n_f - 2C(A)] m_{\tilde{g}} & & \\ \delta Z_{\tilde{q}} &= 0 & \delta m_{\tilde{q}}^2 &= 0 \\ \frac{\delta g_s}{g_s} &= \frac{g_s^2}{16\pi^2\epsilon_{\text{UV}}} [T(F)n_f - C(A)] & & \\ \delta Z_{\phi^0} &= 0 & \delta m_{\phi^0}^2 &= -\frac{g_s^2}{16\pi^2\epsilon_{\text{UV}}} [8T(F)n_f - 16C(A)] m_{\tilde{g}}^2 \\ \delta Z_{\sigma^0} &= 0 & \delta m_{\sigma^0}^2 &= 0 \\ \delta Z_G &= -\frac{g_s^2}{16\pi^2\epsilon_{\text{UV}}} 2T(F)n_f & & \end{aligned} \quad (7.53)$$

## 8 Squark Production at One-Loop

### 8.1 The LSZ Theorem

The LSZ theorem [27] or LSZ reduction formula prescribes how to obtain the S-matrix element, i.e. a physical observable, from the time ordered correlation function of the respective field operators. The time ordered correlation function of fields in an interacting theory can be calculated perturbatively with the aid of the Gell-Mann and Low theorem and the Wick theorem.

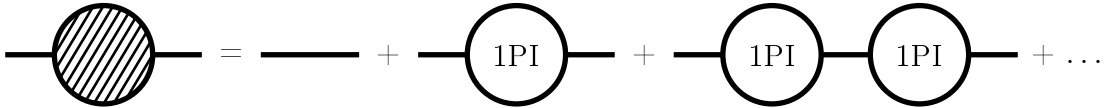
Considering a physical process with kinematics  $\vec{k}_1 \dots \vec{k}_n \rightarrow \vec{p}_1 \dots \vec{p}_m$  and taking for the sake of simplicity only one scalar field  $\phi$  the Fourier transform of the time ordered product of a correlation function is related to the corresponding S-matrix element like<sup>22</sup>

$$\begin{aligned} \prod_{i=1}^n \int dx_i e^{ip_i x_i} \prod_{j=1}^m \int dy_j e^{ik_j y_j} \langle \Omega | \mathcal{T} [\phi(x_1) \dots \phi(x_n) \phi(y_1) \dots \phi(y_m)] | \Omega \rangle \\ \sim \prod_{i=1}^n \frac{i\sqrt{Z}}{p_i^2 - m^2 + i\epsilon} \prod_{j=1}^m \frac{i\sqrt{Z}}{k_j^2 - m^2 + i\epsilon} \left\langle \vec{p}_1 \dots \vec{p}_n | S | \vec{k}_1 \dots \vec{k}_m \right\rangle. \end{aligned} \quad (8.1)$$

Here  $\mathcal{T}$  denotes the time ordering operator,  $|\Omega\rangle$  is the ground state of the interacting theory and  $\sqrt{Z}$  is the residue of the single particle pole in the two-point function

$$\begin{aligned} \int dx e^{ipx} \langle \Omega | \mathcal{T} [\phi(x) \phi(0)] | \Omega \rangle &= \frac{i}{p^2 - m_0^2} + \frac{i}{p^2 - m_0^2} \left( \frac{i\Sigma(p^2)}{p^2 - m_0^2} \right) + \frac{i}{p^2 - m_0^2} \left( \frac{\Sigma(p^2)}{p^2 - m_0^2} \right)^2 + \dots \\ &= \frac{i}{p^2 - m_0^2 - \Sigma(p^2)}. \end{aligned} \quad (8.2)$$

The parameter  $m_0$  is the tree-level mass and the quantity  $-i\Sigma(p^2)$  denotes the sum of all



**Figure 8.1:** Diagrammatic figure of the two-point function of a scalar field: The propagator in an interaction theory can be calculated in a perturbation series in the coupling constant.

one-particle-irreducible contributions to the particle's self-energy. From eq. 8.2 one can read off the physical mass  $m^2$  of the particle which is associated with the field  $\phi$ . It is determined as the value of  $p^2$  where the propagator has a pole, i.e.

$$[p^2 - m_0^2 - \Sigma(p^2)]_{p^2=m^2} = 0 \quad (8.3)$$

<sup>22</sup>In this subsection the  $\sim$  indicates that the poles on either side are the same provided that all momenta are close to their mass shell, i.e.  $p_i^0 \rightarrow E_{\vec{p}_i}$  and  $k_j^0 \rightarrow E_{\vec{k}_j}$ .

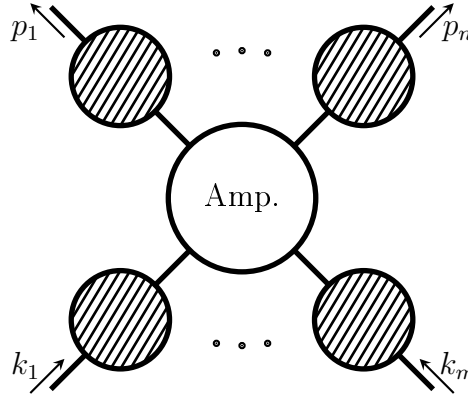
Close to the pole the denominator of eq. 8.2 can be expanded like

$$p^2 - m_0^2 - \Sigma(p^2) = (p^2 - m^2) \left( 1 - \frac{\partial \Sigma(p^2)}{\partial p^2} \right)_{p^2=m^2} + \mathcal{O}((p^2 - m^2)^2). \quad (8.4)$$

The residue of the propagator can therefore be written as

$$Z = \left( 1 - \frac{\partial \Sigma(p^2)}{\partial p^2} \right)_{p^2=m^2}^{-1}. \quad (8.5)$$

Now considers the full  $(n + m)$ -point function in scalar theory. One can decompose it into an amputated  $(n + m)$ -point function and "full" propagators like written in eq. 8.2 and depicted in fig. 8.1.



**Figure 8.2:** Diagrammatic figure of a "full"  $(n + m)$ -point function. Apart from the "full" propagators there is the "full" amputated  $(n + m)$ -point function

Inserting the expression

$$\frac{i}{p^2 - m_0^2 - \Sigma(p^2)} \sim \frac{iZ}{p^2 - m^2} \quad (8.6)$$

for the full propagators one notices the same singularity as in eq. 8.1. If one now compares the coefficients of these poles in eq. 8.1 one finds

$$\left\langle \vec{p}_1 \dots \vec{p}_n | S | \vec{k}_1 \dots \vec{k}_m \right\rangle = \sqrt{Z}^{n+m}.$$



Now it becomes visible why the renormalization in the on-shell scheme

$$\left. \frac{\partial \Sigma(p^2)}{\partial p^2} \right|_{p^2=m^2} \stackrel{!}{=} 0 \quad (8.7)$$

introduces no further manipulation when turning from the correlation function to the S-matrix-element because  $Z = 1$  in on-shell renormalization. Furthermore the on-shell condition for the mass renormalization

$$\Sigma(p^2)|_{p^2=m^2} = 0 \quad (8.8)$$

means that the physical mass equals the tree level mass.

## 8.2 The Squark Production Cross Section at Next-to-Leading Order

Calculate uncertainties for one point(pdf, scale, integration)

## 8.3 The Cross Section in the Limit of Large Sgluon Masses

The cross section for squark production does not exist in the limit of an infinitely large sgluon mass, instead it was found that it diverges logarithmically.

$$\lim_{m_{\sigma^0} \rightarrow \infty} \sigma(qq \rightarrow \tilde{q}\tilde{q}) \sim \ln \frac{m_{\sigma^0}^2}{\mu^2} \quad (8.9)$$

This is actually expected as an effective field theory of the MRSSM where the sgluon is integrated out is no longer supersymmetric. This is because the sgluon is together with the octino part of a supermultiplet. Integrating out only the sgluon means that the octino misses its superpartner in the effective field theory. In this case the decoupling theorem [28] does no longer hold.

Refer to super oblique correction and quantify difference of  $g$  and  $\hat{g}$  from eq 4 in [29]

append plot

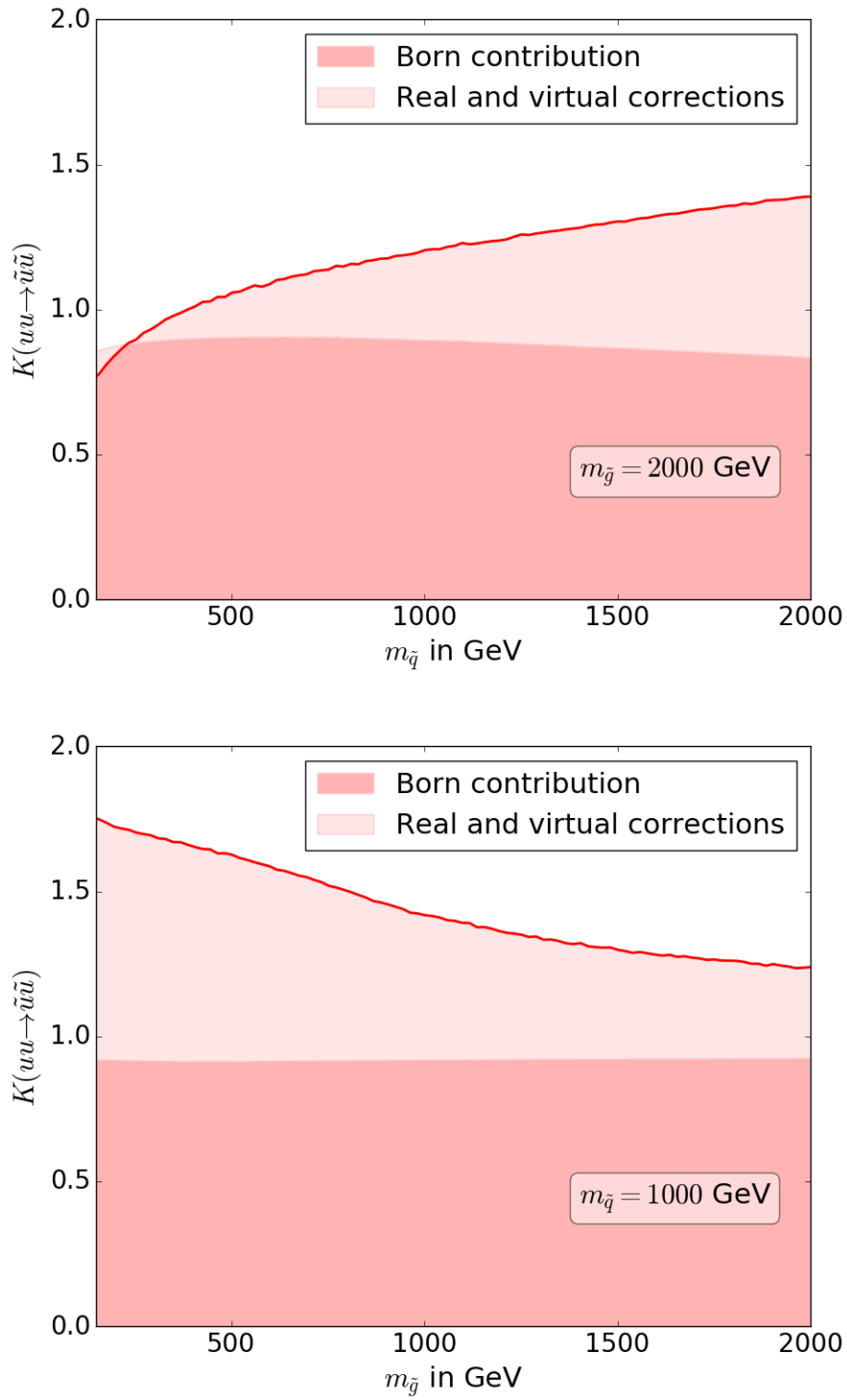


Figure 8.3

## 9 Summary and Outlook

## 10 Appendix

### 10.1 System of Units and Metric

In this thesis the natural units are used, i.e.  $c = \hbar (= k_B) = 1$ . Furthermore the Minkowski metric is chosen to be

$$g^{\mu\nu} = \text{diag}(1, -1, -1, -1). \quad (10.1)$$

### 10.2 Constants of the Colour Algebra $SU(N)$

[Marina von Steinkirch]

The Casimir operator  $C(R)\mathbb{1}$  of a semi-simple Lie algebra in the irreducible representation  $R$  is given by

$$g^{ab}T^a(R)T^b(R) = C(R)\mathbb{1} \quad (10.2)$$

where  $T^a(R)$  is the  $a$ -th generator of the matrix valued representation  $R$ ,  $g^{ab}$  is the metric of group,  $C(R)$  is the Quadratic Casimir invariant of the representation  $R$  and  $\mathbb{1}$  is the identity in the representation space.

Apart from  $C(R)$  it is common to define the Dynkin-Index  $T(R)$ :

$$\text{Tr} [T^a(R)T^b(R)] = T(R)\delta^{ab}. \quad (10.3)$$

The two constants are connected by

$$C(R) \cdot \dim(R) = T(R) \cdot \dim(G) \quad (10.4)$$

where  $\dim(G)$  is the dimension of the group and  $\dim(R)$  is the dimension of the irreducible representation  $R$ .

In the case of  $SU(N)$  one has a diagonal metric  $g^{ab} = \delta^{ab}$  and therefore 10.2 turns to

$$\sum_a (T^a(R))^2 = C(R)\mathbb{1}_{\dim(R) \times \dim(R)} \quad (10.5)$$

and one can write down the following useful formulae for the fundamental representation  $R = F$ :  $T_{ij}^a = \frac{\lambda_{ij}^a}{2}$  and the adjoint representation  $R = A$ :  $(T_{ij}^a)^{adj} = -if_{aij}$

$$\begin{aligned} T_{ik}^a T_{kj}^a &= C(F) \mathbb{1}_{ij} & \text{with } C(F) &= \frac{N^2 - 1}{2N} = \frac{4}{3} \\ f^{abc} f^{dbc} &= C(A) \delta^{ad} & \text{with } C(A) &= N = 3 \\ \text{Tr} [T^a T^b] &= T(F) \delta^{ab} & \text{with } T(F) &= \frac{1}{2} \end{aligned} \quad (10.6)$$

where  $\lambda_{ij}^a$  are for  $N_c = 3$  the Gell-Mann matrices and  $f_{abc}$  are the structure constants of  $SU(N_c)$ .

### 10.3 Weyl basis and 2-spinor notation

As representation of the  $\gamma$ -matrices the Weyl or chiral representation is chosen:

$$\gamma^\mu = \begin{pmatrix} 0 & \sigma^\mu \\ \bar{\sigma}^\mu & 0 \end{pmatrix} \quad \gamma_5 = \begin{pmatrix} -\mathbb{1}_2 & 0 \\ 0 & \mathbb{1}_2 \end{pmatrix} \quad (10.7)$$

with

$$\sigma^\mu = \begin{pmatrix} \mathbb{1}_2 & \sigma^i \\ 0 & 0 \end{pmatrix}, \quad \bar{\sigma}^\mu = \begin{pmatrix} \mathbb{1}_2 & -\sigma^i \\ 0 & 0 \end{pmatrix}, \quad (10.8)$$

where  $\sigma^i$  are the Pauli matrices and  $\mathbb{1}_n$  is the  $n \times n$  unit matrix. The left and right handed projectors are then given by

$$P_L = \frac{1}{2}(\mathbb{1}_4 - \gamma_5) \begin{pmatrix} \mathbb{1}_2 & 0 \\ 0 & 0 \end{pmatrix} \quad P_R = \frac{1}{2}(\mathbb{1}_4 + \gamma_5) \begin{pmatrix} 0 & 0 \\ 0 & \mathbb{1}_2 \end{pmatrix} \quad (10.9)$$

The representation of generators of the Lorentz group on 4-spinor space are composed of the above matrices. Because of the block form of those it is not surprising that the representation on 4-spinor space is reducible to two representations on 2-spinor (Weyl spinor) spaces. It is therefore sensible to decompose a 4 spinor into a left and a right handed Weyl spinor<sup>23</sup>

$$\Psi = \begin{pmatrix} \psi_\alpha \\ \bar{\chi}^{\dot{\alpha}} \end{pmatrix} \quad (10.10)$$

where  $\alpha, \dot{\alpha} \in \{1, 2\}$ . Left handed Weyl spinors are labeled with undotted and right handed Weyl spinors with dotted indices. One distinguishes 4 different Weyl spinors:

$$\psi^\alpha, \quad \bar{\psi}^{\dot{\alpha}} = (\psi^\alpha)^*, \quad \psi_\alpha = \epsilon_{\alpha\beta} \psi^\beta, \quad \text{and} \quad \psi_{\dot{\alpha}} = \epsilon_{\dot{\alpha}\dot{\beta}} \psi^{\dot{\beta}} = (\psi_\alpha)^*, \quad (10.11)$$

<sup>23</sup>The projectors in the chiral basis (eq. 10.9) explain the names left and right handed Weyl spinors.

where  $*$  denotes complex conjugation and indices are lowered with the antisymmetric  $\epsilon_{\alpha\beta}$  ( $\epsilon_{\dot{\alpha}\dot{\beta}}$ ), which obeys

$$\epsilon^{\alpha\beta} = \epsilon_{\beta\alpha}, \quad \epsilon^{\dot{\alpha}\dot{\beta}} = \epsilon_{\dot{\beta}\dot{\alpha}} \quad \text{and} \quad \epsilon_{12} = \epsilon_{\dot{1}\dot{2}} = 1. \quad (10.12)$$

By virtue of the antisymmetry of  $\epsilon$  one finds the Lorentz invariant products:

$$\begin{aligned} \psi\chi &:= \psi^\alpha\chi_\alpha = -\chi_\alpha\psi^\alpha = \chi^\alpha\psi_\alpha = \chi\psi, \\ \bar{\psi}\bar{\chi} &:= \bar{\psi}_{\dot{\alpha}}\bar{\chi}^{\dot{\alpha}} = -\bar{\chi}^{\dot{\alpha}}\bar{\psi}_{\dot{\alpha}} = \bar{\chi}_{\dot{\alpha}}\bar{\psi}^{\dot{\alpha}} = \bar{\chi}\bar{\psi} \end{aligned} \quad (10.13)$$

To make the index structure of the Pauli matrices explicit one writes  $\sigma_{\alpha\dot{\alpha}}^\mu$  and  $\bar{\sigma}^{\mu\dot{\alpha}\alpha}$  for the formulae in 10.8. For the definition of the supersymmetry algebra in section 3.1 the representation of the generators of the Lorentz group on the left and right handed Weyl spinor space are introduced:

$$\begin{aligned} \frac{1}{2}(\sigma^{\mu\nu})_\alpha{}^\beta &:= \frac{i}{4}(\sigma^\mu\bar{\sigma}^\nu - \sigma^\nu\bar{\sigma}^\mu)_\alpha{}^\beta \\ \frac{1}{2}(\bar{\sigma}^{\mu\nu})^{\dot{\alpha}}{}_{\dot{\beta}} &:= \frac{i}{4}(\bar{\sigma}^\mu\sigma^\nu - \bar{\sigma}^\nu\sigma^\mu)^{\dot{\alpha}}{}_{\dot{\beta}} \end{aligned} \quad (10.14)$$

With the definition of bared and charge conjugated 4-spinors<sup>24</sup>

$$\bar{\Psi} := \Psi^\dagger\gamma^0, \quad \Psi^C := i\gamma^2\gamma^0\bar{\Psi}^T \quad (10.15)$$

one obtains:

$$\begin{aligned} \Psi &= \begin{pmatrix} \psi_\alpha \\ \bar{\chi}^{\dot{\alpha}} \end{pmatrix}, & \bar{\Psi} &= \begin{pmatrix} \chi^\alpha & \bar{\psi}_{\dot{\alpha}} \end{pmatrix}, \\ \Psi^C &= \begin{pmatrix} \chi_\alpha \\ \bar{\psi}^{\dot{\alpha}} \end{pmatrix}, & \bar{\Psi}^C &= \begin{pmatrix} \psi^\alpha & \bar{\chi}_{\dot{\alpha}} \end{pmatrix}. \end{aligned} \quad (10.16)$$

The 4-spinor of an arbitrary quark  $q$  is given in terms of Weyl spinors  $q_L$  and  $\bar{q}_R$  by

$$q = \begin{pmatrix} q_L \\ \bar{q}_R \end{pmatrix} \quad (10.17)$$

whereas the 4-spinor of the Dirac gauginos is given in terms of the Weyl spinors  $\lambda$  and  $\bar{\chi}$ <sup>25</sup>

$$\tilde{g}^a = \begin{pmatrix} -i\lambda^a \\ i\bar{\chi}^a \end{pmatrix} \quad (10.18)$$

<sup>24</sup> $\Psi^T$  denotes the transpose of the spinor  $\Psi$  and  $\Psi^\dagger$  is the Hermitian adjoint of  $\Psi$ .

<sup>25</sup> $\lambda$  is the superpartner of the gluon, called the gluino and  $\bar{\chi}$  is the Weyl spinor of the chiral superfield which is associated with the gluon, called the octino.

## 10.4 Anticommuting numbers

Anticommuting numbers  $\theta^\alpha$  are also referred to as Grassmann numbers and are defined by  $\theta^\alpha\theta^\beta = -\theta^\beta\theta^\alpha$  and commute with ordinary numbers.

They occur in superspace formalism in the form of 2 tuples, i.e.  $\theta^\alpha$  with  $\alpha = 1, 2$ . The complex conjugate of this tuple is denoted with  $\bar{\theta}^{\dot{\alpha}}$ . Derivatives are defined by

$$\begin{aligned}\partial^\alpha\theta_\beta &:= \frac{\partial}{\partial\theta_\alpha}\theta_\beta := \delta_\beta^\alpha & \partial_\alpha\theta^\beta &:= \frac{\partial}{\partial\theta^\alpha}\theta^\beta := \delta_\alpha^\beta \\ \bar{\partial}^{\dot{\alpha}}\bar{\theta}_{\dot{\beta}} &:= \frac{\partial}{\partial\bar{\theta}_{\dot{\alpha}}}\bar{\theta}_{\dot{\beta}} := \delta_{\dot{\beta}}^{\dot{\alpha}} & \bar{\partial}_{\dot{\alpha}}\bar{\theta}^{\dot{\beta}} &:= \frac{\partial}{\partial\bar{\theta}^{\dot{\alpha}}}\bar{\theta}^{\dot{\beta}} := \delta_{\dot{\alpha}}^{\dot{\beta}}\end{aligned}\quad (10.19)$$

whereby one needs to be cautious as these definitions imply

$$\partial_\alpha = -\epsilon_{\alpha\beta}\partial^\beta \qquad \bar{\partial}_{\dot{\alpha}} = -\epsilon_{\dot{\alpha}\dot{\beta}}\bar{\partial}^{\dot{\beta}}. \quad (10.20)$$

Integrals are defined by:

$$\begin{aligned}\int d\theta_\alpha (a + b\theta^\beta + c\theta^\beta\theta^\gamma) &:= b\delta_\alpha^\beta + c(\delta_\alpha^\beta\theta^\gamma - \delta_\alpha^\gamma\theta^\beta) \quad \text{and} \\ \int d\theta_\alpha (a\bar{\theta}^{\dot{\beta}}) &:= (a\bar{\theta}^{\dot{\beta}}) \int d\theta_\alpha\end{aligned}\quad (10.21)$$

where the first line mirrors the claim of translation invariance. One furthermore introduces the shortcuts

$$\begin{aligned}\int d\theta^2 &:= \int \frac{1}{4}\epsilon_{\alpha\beta}d\theta^\alpha d\theta^\beta, & \int d\bar{\theta}^2 &:= \int \frac{1}{4}\epsilon_{\dot{\alpha}\dot{\beta}}d\theta^{\dot{\alpha}}d\theta^{\dot{\beta}}, & \text{and} \\ \int d^4\theta &:= \int d\theta^2 d\bar{\theta}^2.\end{aligned}\quad (10.22)$$

For the definition of (anti-)chiral superfields and the field strength chiral superfields it proves useful to introduce supersymmetry (or chiral) covariant derivatives

$$\begin{aligned}\mathcal{D}_\alpha &:= \frac{\partial}{\partial\theta^\alpha} - i(\sigma^\mu\bar{\theta})_\alpha\partial_\mu & \mathcal{D}^{\dot{\alpha}} &:= \frac{\partial}{\partial\bar{\theta}_{\dot{\alpha}}} - i(\bar{\sigma}^\mu\theta)^{\dot{\alpha}}\partial_\mu \\ D^\alpha &:= \epsilon^{\alpha\beta}D_\beta = -\frac{\partial}{\partial\theta_\alpha} + i(\bar{\theta}\bar{\sigma}^\mu)^\alpha\partial_\mu & \mathcal{D}_{\dot{\alpha}} &:= \epsilon_{\dot{\alpha}\dot{\beta}}\mathcal{D}^{\dot{\beta}} = -\frac{\partial}{\partial\bar{\theta}^{\dot{\alpha}}} + i(\sigma^\mu\theta)_{\dot{\alpha}}\partial_\mu.\end{aligned}\quad (10.23)$$

By construction they are supersymmetry invariant, i.e. they fulfill the following anticommutation relations with the supersymmetry generators from eq. 3.2:

$$\{Q_\alpha, \mathcal{D}_\beta\} = \{\bar{Q}_{\dot{\alpha}}, \mathcal{D}_\beta\} = \{Q_\alpha, \mathcal{D}_{\dot{\beta}}\} = \{\bar{Q}_{\dot{\alpha}}, \mathcal{D}_{\dot{\beta}}\} = 0 \quad (10.24)$$

## 10.5 Feynman rules for the RSQCD

The following Feynman rules are derived from 4.16. When compared with the Feynman rules of the supersymmetric QCD the diagrams involving scalar gluons are new. In addition the gluon-quark-squark vertex is different in RSQCD for the gauginos are Dirac fermions.

refer to paper cited in Beenakker(Appendix: Feynmanrules)

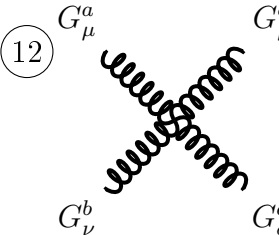


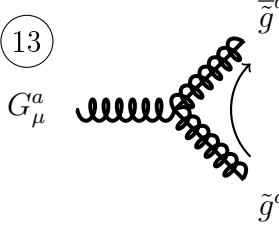
$$\begin{array}{ll}
\textcircled{1} & \textcircled{2} \\
(\sigma^0)^a \cdots (\sigma^0)^b \hat{=} \frac{i}{p^2 - m_{\sigma^0}^2 + i\varepsilon} \delta_{ab} & (\phi^0)^a \cdots (\phi^0)^b \hat{=} \frac{i}{p^2 - m_{\phi^0}^2 + i\varepsilon} \delta_{ab} \\
\\
\textcircled{3} & \textcircled{4} \\
\tilde{q}_{Ai} \cdots \leftarrow \cdots \tilde{q}_{Bj}^\dagger \hat{=} \frac{i\delta_{AB}}{p^2 - m_q^2 + i\varepsilon} \delta_{ij} & q_i \longrightarrow \cdots \bar{q}_j \hat{=} i \frac{\not{p} + m_q}{p^2 - m_q^2 + i\varepsilon} \delta_{ij} \\
\\
\textcircled{5} & \textcircled{6} \\
G_\mu^a \cdots G_\nu^b \hat{=} -i \frac{g^{\mu\nu}}{p^2 + i\varepsilon} \delta_{ab} & \tilde{g}^a \cdots \bar{\tilde{g}}^b \hat{=} i \frac{\not{p} + m_{\tilde{g}}}{p^2 - m_{\tilde{g}}^2 + i\varepsilon} \delta_{ab} \\
\\
\textcircled{7} & \\
c^a \cdots \leftarrow \cdots \bar{c}^b \hat{=} \frac{i}{p^2 + i\varepsilon} \delta_{ab} & \\
\\
\textcircled{8} & \textcircled{9} \\
G_\mu^a \cdots \begin{array}{l} \nearrow \bar{q}_i \\ \searrow q_j \end{array} \hat{=} -ig_s T_{ij}^a \gamma^\mu & G_\mu^a \cdots \begin{array}{l} \nearrow \tilde{q}_{Ai}^\dagger(-p_1) \\ \searrow \tilde{q}_{Bj}(p_2) \end{array} \hat{=} -ig_s (p_2 + p_1)^\mu T_{ij}^a \delta_{AB} \\
\\
\textcircled{10} & \\
G_\mu^a \cdots \begin{array}{l} \nearrow \tilde{q}_{Ai}^\dagger \\ \searrow \tilde{q}_{Bj} \end{array} \hat{=} ig_s^2 g^{\mu\nu} \{T^a, T^b\}_{ij} \delta_{AB} & \\
\\
\textcircled{11} & \\
G_\rho^c(p_c) \cdots \begin{array}{l} \nearrow G_\mu^a(p_a) \\ \searrow G_\nu^b(p_b) \end{array} \hat{=} -g_s f_{abc} [g_{\mu\nu}(p_a - p_b)^\rho + g_{\nu\rho}(p_b - p_c)^\mu + g_{\rho\mu}(p_c - p_a)^\nu] &
\end{array}$$

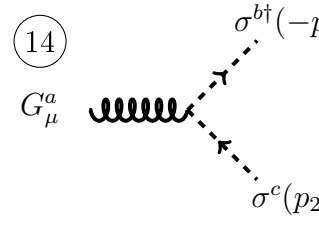
**Figure 10.1:** In the Feynman diagrams of the propagators the momentum is flowing from the right to the left hand side.

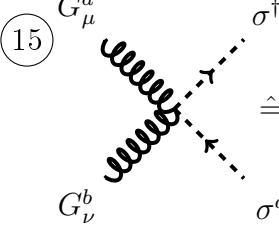
In the Feynman diagrams of the vertices all momenta flow towards the vertex.

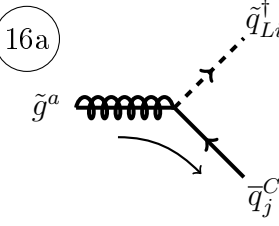
The indices  $A, B = L, R$  label the right/left-"handedness" of the squarks. The indices  $i, j = 1, 2, 3$  are the color indices in the (anti)fundamental representation where  $a, b, c, \dots = 1, \dots, 8$  are the color indices of the adjoint representation.

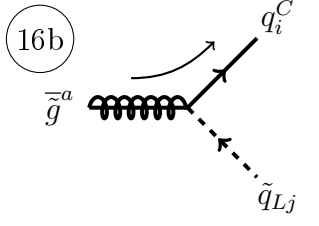
(12)   $\hat{=} -ig_s^2[f^{abe}f^{cde}(g^{\mu\rho}g^{\nu\sigma} - g^{\mu\sigma}g^{\nu\rho}) + f^{ace}f^{bde}(g^{\mu\nu}g^{\rho\sigma} - g^{\mu\sigma}g^{\nu\rho}) + f^{ade}f^{bce}(g^{\mu\nu}g^{\rho\sigma} - g^{\mu\rho}g^{\nu\sigma})]$

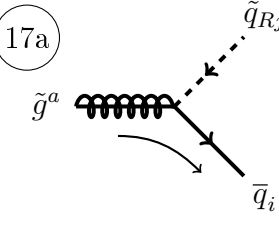
(13)   $\hat{=} -g_s f_{abc} \gamma^\mu$

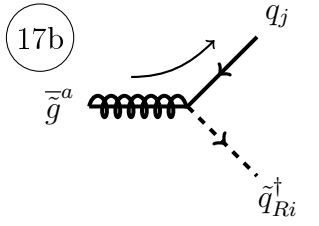
(14)   $\hat{=} -g_s(p_1 + p_2)^\mu f_{abc}$

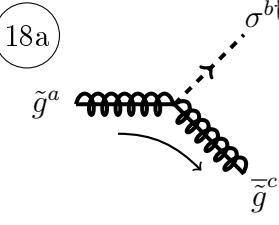
(15)   $\hat{=} +ig_s^2 g^{\mu\nu}[f^{aec}f^{bed} + f^{bec}f^{aed}]$

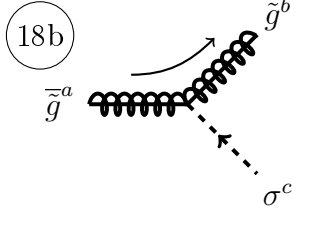
(16a)   $\hat{=} -i\sqrt{2}g_s T_{ij}^a P_L$

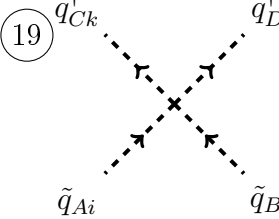
(16b)   $\hat{=} -i\sqrt{2}g_s T_{ij}^a P_R$

(17a)   $\hat{=} +i\sqrt{2}g_s T_{ij}^a P_L$

(17b)   $\hat{=} +i\sqrt{2}g_s T_{ij}^a P_R$

(18a)   $\hat{=} -\sqrt{2}g_s f^{abc} P_L$

(18b)   $\hat{=} +\sqrt{2}g_s f^{abc} P_R$

(19)   $\hat{=} -ig_s^2[T_{ki}^a T_{lj}^a(\delta_{AL}\delta_{CL} - \delta_{AR}\delta_{CR})(\delta_{BL}\delta_{DL} - \delta_{BR}\delta_{DR}) + T_{kj}^a T_{li}^a(\delta_{BL}\delta_{CL} - \delta_{BR}\delta_{CR})(\delta_{AL}\delta_{DL} - \delta_{AR}\delta_{DR})]$

**Figure 10.2:** The curved arrows indicate the fermion flow. The Feynman rules 16b, 17b and 18b are the complex conjugates of 16a, 17a and 18a respectively. Applying a flipping rule to a vertex one has to reverse the curved arrow, i.e. the fermion flow and replace  $\Psi$  with  $\bar{\Psi}^C$ . In addition one has to add a minus sign for Feynman rule 13.

$$\begin{aligned}
(20) \quad & \begin{array}{c} \tilde{q}_{Aj}^\dagger \quad \sigma^{b\dagger} \\ \diagdown \quad \diagup \\ \times \\ \diagup \quad \diagdown \\ \tilde{q}_{Ai} \quad \sigma^c \end{array} \hat{=} -g_s^2 T_{ij}^a f^{abc} (\delta_{AL} \delta_{CL} - \delta_{AR} \delta_{BR}) \\
(21) \quad & \begin{array}{c} \tilde{q}_{Ai}^\dagger \\ \diagdown \quad \diagup \\ \sigma^a + \sigma^{a\dagger} \quad \tilde{q}_{Bj} \end{array} \hat{=} -i\sqrt{2} g_s m_{\tilde{g}} T_{ij}^a (\delta_{AL} \delta_{BL} - \delta_{AR} \delta_{BR}) \\
(22) \quad & \begin{array}{c} \sigma^{\dagger b} \quad \sigma^{\dagger d} \\ \diagdown \quad \diagup \\ \times \\ \diagup \quad \diagdown \\ \sigma^c \quad \sigma^e \end{array} \hat{=} -g_s^2 (f^{abc} f^{ade} + f^{abc} f^{adc})
\end{aligned}$$

## 10.6 Passarino-Veltman Integrals

The definition of the Passarino-Veltman integrals in this thesis agrees with the one from `LoopTools` [30] convention. The original paper [31] uses slightly different conventions. A pedagogical introduction to the evaluation of one-loop integrals can be found in [32]

$$\begin{aligned}
\frac{i}{16\pi^2} A_0(m^2) &= \int_l \frac{1}{l^2 - m^2} \\
\frac{i}{16\pi^2} B_{0,\mu,\mu\nu}(p^2, m_1^2, m_2^2) &= \int_l \frac{\{1, l_\mu, l_\mu l_\nu\}}{[l^2 - m_1^2][(l+p)^2 - m_2^2]} \\
\frac{i}{16\pi^2} C_{0,\mu,\mu\nu}(p_1^2, p_2^2, (p_1+p_2)^2, m_1^2, m_2^2, m_3^2) &= \int_l \frac{\{1, l_\mu, l_\mu l_\nu\}}{[l^2 - m_1^2][(l+p_1)^2 - m_2^2][(l+p_1+p_2)^2 - m_3^2]}
\end{aligned} \tag{10.25}$$

with the shortcut  $\int_l = \mu^{2\epsilon} \int \frac{d^D l}{(2\pi)^D}$ . Furthermore there are suppressed  $\varepsilon$ 's which prescribe how the poles in the complex plane are avoided. They are hidden in the infinitesimal shift of the masses:  $m_i^2 \rightarrow m_i^2 - i\varepsilon$ .

The tensor integrals can be decomposed as

$$\begin{aligned}
B_\mu &:= p_\mu B_1 \\
B_{\mu\nu} &:= g_{\mu\nu} B_{00} + p_\mu p_\nu B_{11} \\
C_\mu &= p_{1\mu} C_1 + p_{2\mu} C_2 \\
C_{\mu\nu} &:= g_{\mu\nu} C_{00} + p_{1\mu} p_{1\nu} C_{11} + p_{2\mu} p_{2\nu} C_{22} + (p_{1\mu} p_{2\nu} + p_{2\mu} p_{1\nu}) C_{12}
\end{aligned} \tag{10.26}$$

In the the special case of vanishing momenta the integrals take a succinct form.

$$A_0(m^2) = m^2 \left( \Delta_\epsilon - \ln \frac{m^2}{\mu^2} + 1 \right) + \mathcal{O}(\epsilon) \quad (10.27)$$

where the typical UV-divergent constant  $\Delta_\epsilon = \frac{1}{\epsilon} - \gamma_E + \ln 4\pi$  is defined. It comprises the Euler-Mascheroni constant  $\gamma_E$ .

$$B_0(0, m_1^2, m_2^2) = \frac{A_0(m_1^2) - A_0(m_2^2)}{m_1^2 - m_2^2} = \Delta_\epsilon + 1 - \frac{m_1^2 \ln \frac{m_1^2}{\mu^2} - m_2^2 \ln \frac{m_2^2}{\mu^2}}{m_1^2 - m_2^2} + \mathcal{O}(\epsilon) \quad (10.28)$$

$$B_0(0, m^2, m^2) = \frac{\partial}{\partial m^2} A_0(m^2) = \Delta_\epsilon - \ln \frac{m^2}{\mu^2} + \mathcal{O}(\epsilon) \quad (10.29)$$

$$B_0(0, 0, 0) = 0 \quad (10.30)$$

$$B_1(0, m_1^2, m_2^2) = -\frac{\Delta_\epsilon}{2} + \frac{1}{2} \ln \frac{m_1^2}{\mu^2} + \frac{-3 + 4t - t^2 - 4t \ln t + 2t^2 \ln t}{4(t-1)^2} + \mathcal{O}(\epsilon) \quad (10.31)$$

$$B_1(0, 0, 0) = 0 \quad (10.32)$$

The scaleless integrals are defined to be zero. The parameter  $t$  is given by  $\frac{m_2^2}{m_1^2}$ . As can be seen from 10.31  $B_1$  is in contrast to  $B_0$  not symmetric in its masses but it can be shown [Romao] that

$$B_1(p^2, m_1^2, m_2^2) = -(B_0(p^2, m_2^2, m_1^2) + B_1(p^2, m_2^2, m_1^2)) \quad (10.33)$$

It can further be shown that

$$C_{00}(0, 0, 0, m_1^2, m_1^2, m_2^2) = -\frac{1}{2} B_1(0, m_1^2, m_2^2) \quad (10.34)$$

and that  $C_{00}(0, 0, 0, m_1^2, m_1^2, m_2^2)$  is a symmetric function of its masses.

From the generic  $\epsilon$ -expansion of the  $B_0$  integral [25]

$$B_0(p^2, m_1^2, m_2^2) = \Delta_\epsilon - \int_0^1 dx \ln \frac{-x(1-x)p^2 + xm_2^2 + (1-x)m_1^2}{\mu^2} \quad (10.35)$$

and Passarino-Veltman decomposition [31] one can determine the UV-divergent part of all  $B$  and  $C$  integrals. In chapter 7.7 this was necessary in order to obtain the renormalization constants in the  $\overline{\text{MS}}$ -scheme. Infrared and collinear singularities arise from the special case where one or multiple masses tend to zero. These poles are either regularized in terms of a small mass cutoff  $\Lambda$  or also dimensionally as  $\epsilon$ -poles. In the later case the integral first needs to undergo the limit to zero masses and than being evaluated.

The following list shows all necessary integrals needed to determine the renormalization con-

stants in 7.7.

$$A_0(m^2)|_{\text{UV-div}} = m^2 \Delta_\epsilon \quad (10.36)$$

$$B_0(p^2, m_1^2, m_2^2)|_{\text{UV-div}} = \Delta_\epsilon \quad (10.37)$$

$$B_1(p^2, m_1^2, m_2^2)|_{\text{UV-div}} = -\frac{1}{2} \Delta_\epsilon \quad (10.38)$$

$$C_{00}(p_1^2, p_2^2, (p_1 + p_2)^2, m_1^2, m_2^2, m_3^2)|_{\text{UV-div}} = \frac{1}{4} \Delta_\epsilon \quad (10.39)$$

The UV-divergent part of  $C_{11}$ ,  $C_{22}$ ,  $C_{12}$  equals zero. As can be seen already from the superficial degree of divergence also  $C_i|_{\text{UV-div}} = 0$  for  $i \in \{0, 1, 2\}$ .

## 10.7 Cross section and Phase Space Integration

Once the Feynman amplitude  $\mathcal{M}$  for a  $2 \rightarrow N$  body scattering<sup>26</sup> is computed one can calculate physical observables with it. The differential cross section for  $2 \rightarrow N$  scattering is given by

$$d\sigma = \frac{1}{F} d\Phi_N |\mathcal{M}|^2. \quad (10.40)$$

The flux factor is defined by  $F = 4\sqrt{(k_a \cdot k_b)^2 - (m_a m_b)^2}$  which equals  $F = 2s$  for massless initial state particles. The differential for the  $N$  body phase space in  $D$  dimensions is given by

$$d\Phi_N = (\mu^{2\epsilon})^{N-1} \left( \prod_{f=1}^N \frac{d^{D-1} p_f}{(2\pi)^{D-1}} \frac{1}{2E_f} \right) (2\pi)^D \delta^{(D)}(k_a + k_b - \sum_{f=1}^N p_f). \quad (10.41)$$

The factor  $\mu^{2\epsilon}$  is included to maintain the mass dimension of the cross section. As in this thesis the sum of  $|\mathcal{M}|^2$  over helicities and colors  $\sum |\mathcal{M}|^2$  has been calculated one can write

$$d\sigma = \frac{1}{2s} d\Phi_2 K_{ab} \sum |\mathcal{M}|^2 \quad (10.42)$$

where  $K_{ab}$  encodes the averaging over initial state helicities and colors. Specifying to the center-of-mass frame and assuming that  $\sum |\mathcal{M}|^2$  is only a function of the modulus of one of the final state particle's 3-momentum  $|\vec{p}_i|$  and the angle  $\theta$  between  $\vec{k}_a$  and  $\vec{p}_1$  one can write

$$\begin{aligned} \int d\Phi_2 &= \mu^{2\epsilon} \int \frac{d|\vec{p}_1| d\Omega_1^{D-1}}{(2\pi)^{D-2} 4E_1 E_2} |\vec{p}_1|^{D-2} \delta \left( k_a^0 + k_b^0 - \sqrt{m_1^2 + |\vec{p}_1|^2} - \sqrt{m_2^2 + |\vec{p}_1|^2} \right) \\ &= \frac{1}{(2\pi)^{D-2}} \frac{2\pi^{\frac{D}{2}-1}}{\Gamma(\frac{D}{2}-1)} \mu^{2\epsilon} \int_0^\infty d|\vec{p}_1| \int_0^\pi d\cos\theta \frac{1}{4E_1 E_2} p_1^{D-2} \sin^{D-4}\theta \\ &\quad \delta \left( k_a^0 + k_b^0 - \sqrt{m_1^2 + |\vec{p}_1|^2} - \sqrt{m_2^2 + |\vec{p}_1|^2} \right). \end{aligned} \quad (10.43)$$

<sup>26</sup>with kinematics  $k_a + k_b \rightarrow p_1 + \dots p_N$

In the second line the integral over the D-dimensional hypersphere

$$\int d\Omega^D = \int_0^{2\pi} d\phi \prod_{i=1}^{D-2} \int_0^\pi \sin^i \theta_i d\theta_i = \frac{2\pi^{\frac{D}{2}}}{\Gamma(\frac{D}{2})} \quad (10.44)$$

has been used. Because  $\sum |\mathcal{M}|^2$  is calculated in terms of Mandelstam variables

$$\begin{aligned} t &= (k_a - p_1)^2 \\ t &= -2 \left( |\vec{k}_a| \sqrt{m_1^2 + |\vec{p}_1|^2} - |\vec{k}_a| |\vec{p}_1| \cos \theta \right) + m_1^2 \end{aligned} \quad (10.45)$$

$$\begin{aligned} u &= (k_a - p_2)^2 \\ u &= -2 \left( |\vec{k}_a| \sqrt{m_2^2 + |\vec{p}_1|^2} + |\vec{k}_a| |\vec{p}_1| \cos \theta \right) + m_2^2 \end{aligned} \quad (10.46)$$

it is useful to perform a change of coordinates yielding

$$d|\vec{p}_1| d\cos \theta = -\frac{E_1 E_2}{4|\vec{k}_a|^2 |\vec{p}_1|^2 (E_1 + E_2)} du dt. \quad (10.47)$$

Inserting 10.47 into 10.43 and using  $2|\vec{k}_a| = \sqrt{s} = E_1 + E_2$  gives

$$\begin{aligned} \int d\Phi_2 &= \frac{1}{s} \frac{\pi^{-\frac{D}{2}+1}}{2^{D-3} \Gamma(\frac{D}{2}-1)} \int du dt \left( \frac{tu - m_1^2 m_2^2}{\mu^{2\epsilon} s} \right)^{\frac{D-4}{2}} \\ &\quad \frac{1}{4} \Theta(tu - 4m_1^2 m_2^2) \delta(s + t + u - m_1^2 - m_2^2) \end{aligned} \quad (10.48)$$

where the  $\Theta$ -function comes from the bounds of  $|\vec{p}_1|$  and  $\theta$  visible in 10.43 and the combination of 10.45 and 10.46. Working in  $D = 4 - 2\epsilon$  dimensions and inserting  $\Theta(s - 4m^2)$  with  $m = \frac{m_1 + m_2}{2}$  to account for the production threshold one finds

$$\begin{aligned} \frac{d^2\sigma}{dt du} &= \frac{K_{ab}}{s^2} \frac{\pi S_\epsilon}{\Gamma(1-\epsilon)} \left[ \frac{tu - m_1^2 m_2^2}{\mu^2 s} \right]^{-\epsilon} \Theta(tu - m_1^2 m_2^2) \\ &\quad \Theta(s - 4m^2) \delta(s + t + u - m_1^2 - m_2^2) \sum |\mathcal{M}|^2 \end{aligned} \quad (10.49)$$

where  $S_\epsilon = (4\pi)^{-2+\epsilon}$  as defined in [16]. The averaging factors  $K_{ab}$  are given by

$$K_{qq} = \frac{1}{4N_c^2} \quad K_{GG} = \frac{1}{4(1-\epsilon)^2(N_c^2-1)^2} \quad K_{qG} = \frac{1}{4(1-\epsilon)N_c(N_c^2-1)}. \quad (10.50)$$

## 11 References

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## **Erklärung**

Hiermit erkläre ich, dass ich diese Arbeit im Rahmen der Betreuung am Institut für ??  
Physik ohne unzulässige Hilfe Dritter verfasst und alle Quellen als solche gekennzeichnet habe.

Vorname Nachname

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