

Supersymmetry

Lecture notes for FYS5190/FYS9190 Supersymmetry

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December 1, 2021

Preface

The goal of these lecture notes is to introduce the basics of low-energy models of supersymmetry (SUSY) using the Minimal Supersymmetric Standard Model (MSSM) as our main example. The notes are based on lectures given at the University of Oslo in 2011, 2013, 2015, 2017, and 2019, and lectures at the NORDITA Winter School on Theoretical Particle Physics in 2012. This document owes a particular gratitude to Paul Batzing, who took notes during the 2011 lectures, forming the start of this document.

Oslo, May 2021
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Introduction

Rather than the traditional approach of starting with the current problems of the Standard Model of particle physics, and how supersymmetry can solve these, in these notes we will focus on the algebraic origin of supersymmetry in the sense of an extension of the symmetries of Einstein's special relativity (SR). This was the original motivation for work on what we today call supersymmetry.

We first need to introduce some basic mathematical concepts used in physics for exploring symmetries, mainly groups and Lie algebras, which we will take care of in Chapter 1. In Chapter 2 we will then study the symmetries of special relativity, through the Poincaré group, and look at how these symmetries can be extended into the super-Poincaré group by adding so-called supercharges to the Poincaré Lie algebra. We will then introduce superspace, a coordinate system where supersymmetry is manifest, and use this to derive differential representations for the supercharges in Chapter 3. Here, we also define superfields, which function as representations of the super-Poincaré group, and are the building blocks of the supersymmetric Lagrangians we construct in Chapter 4. In Chapter 5 we discuss how to break supersymmetry using spontaneous symmetry breaking.

From Chapter 6 and onwards we turn more towards phenomenology, constructing the minimal realisation of a supersymmetric Standard Model, before we discuss its phenomenology in Chapter 7, before ending with a discussion of supersymmetric Dark Matter candidates in Chapter 8.

Note that sections marked with an asterisk are somewhat tangential to the main argument of the text and can be read for light entertainment only. Solutions to some of the exercises can be found in Appendix A.

Chapter 1

Groups and algebras

The study of symmetries plays a central part in theoretical physics, and the mathematical language we use is that of groups. The action of the group elements on our (quantum mechanical) states describes the transformation of the symmetry operation, while the invariance of the physical properties of the system under that transformation is the symmetry itself. For example, rotations in three-dimensions can be carried out by the application of a 3×3 rotation matrix on the coordinates. These matrices form the group called $SO(3)$. For a sphere, which is invariant under these rotations, $SO(3)$ is then the symmetry group.

Of special interest to us are the Lie groups, which are the groups that represent continuous transformations, such as the rotations. The properties of Lie groups can be further studied by finding their generators which form a (Lie) algebra. The generators *almost* – in a very specific sense of the word almost – describes the whole group, and allows us to reconstruct the group elements by what is called the exponential map.

Here, we will begin by defining groups and looking at some of their most important properties. What is crucial in physics are the *representations* of groups, meaning what the operators of the transformations on the states actually look like. Returning to the rotation example these are 3×3 matrices, but with some restrictions on their elements. After discussing representations we will move on to defining Lie groups, before we end on a discussion of their generators and corresponding algebras.

1.1 Group definition

A group is an abstract mathematical structure that consists of a set of objects (elements), and a multiplication rule acting between pairs of these objects. We define a group as follows.

Definition: The set of elements $G = \{g_i\}$ and operation $\circ$ (sometimes called multiplication) form a **group** if and only if for $\forall g_i \in G$:

- i) $g_i \circ g_j \in G$, (closure)
- ii) $(g_i \circ g_j) \circ g_k = g_i \circ (g_j \circ g_k)$, (associativity)
- iii) $\exists e \in G$ such that $g_i \circ e = e \circ g_i = g_i$, (identity element)
- iv) $\exists g_i^{-1} \in G$ such that $g_i \circ g_i^{-1} = g_i^{-1} \circ g_i = e$. (inverse)

Below, where no confusion can occur, we will often drop the multiplication symbol for the group multiplication (and other abstract multiplications), writing $g_i \circ g_j = g_i g_j$.

A straight forward example of a group is $G = \mathbb{Z}$ (the integers), with standard addition as the operation $\circ$. Then $e = 0$ and $g^{-1} = -g$. Alternatively we can restrict the group to $\mathbb{Z}_n$, where the operation is addition modulo n . In this group, $g_i^{-1} = n - g_i$ and the unit element is again $e = 0$.¹ Here, $\mathbb{Z}$ is an example of an **infinite** group, the set has an infinite number of members, while $\mathbb{Z}_n$ is **finite**, with **order** n , meaning n members. Both are **abelian** groups, meaning that the elements commute: $g_i \circ g_j = g_j \circ g_i$, because the standard addition commutes.

The simplest, non-trivial, of these $\mathbb{Z}_n$ -groups is $\mathbb{Z}_2$ which only has the members $e = 0$ and 1. The “multiplication” operation is completely defined by the three possibilities $0 + 0 = 0$, $0 + 1 = 1$ and $1 + 1 = 0$. Now, compare this to the set $G = \{-1, 1\}$ with the ordinary multiplication operation. Here, all the possible operations are $1 \cdot 1 = 1$, $1 \cdot (-1) = -1$ and $(-1) \cdot (-1) = 1$. This has exactly the same structure as $\mathbb{Z}_2$, only that the identity element is now 1. We say that these two groups are **isomorphic**, because there is a mapping between all the (two) elements $0 \rightarrow 1$ and $1 \rightarrow -1$, and the results of the multiplication operation is the same, and in fact we consider them as the *same group* despite the considerable apparent visual differences.² This notion of isomorphic (identical) groups is very important, and we will return to it in more detail in Sec. 1.3.

A somewhat more sophisticated example of a group can be found in the Taylor expansion of a function F , where

$$\begin{aligned}
 F(x+a) &= F(x) + aF'(x) + \frac{1}{2}a^2F''(x) + \dots \\
 &= \sum_{n=0}^{\infty} \frac{a^n}{n!} \frac{d^n}{dx^n} F(x) \\
 &= e^{a \frac{d}{dx}} F(x).
 \end{aligned}$$

The last equality uses the formal definition of the exponential series, but may drive some mathematicians crazy.³ The resulting operator $T_a = e^{a \frac{d}{dx}}$ is called the **translation operator**,

¹Note that we here use e for the identity in an abstract group, while we will later use I or 1 as the identity matrix in matrix representations of groups.

²This observation generalises to the set $G = \{e^{2\pi i k/n} | k = 1, \dots, n-1\}$, the **n-th roots of unity**, which, together with the standard multiplication operation, is isomorphic to $\mathbb{Z}_n$.

³We will not discuss this further, but there is a deep question here whether the operator formed by this exponentiation is well defined.

in this case in one dimension, since it shifts the coordinate x of the function f it is operating on by an amount a . Defining the (natural) multiplication operation $T_a \circ T_b = T_{a+b}$ it forms the **translational group** $T(1)$, where we can show that $T_a^{-1} = T_{-a}$.⁴ In n dimensions the group $T(n)$ has the elements $T_{\mathbf{a}} = e^{\mathbf{a} \cdot \nabla}$. Whereas we say that the groups $\mathbb{Z}$ and $\mathbb{Z}_n$ are **discrete groups**, since we can count the number of elements, T_a is a **continuous group** since the parameter a can be any real number.

1.2 Matrix groups

We next define some groups that are very important in physics and to the discussion in these notes. They have in common that they are defined in terms of square matrices.

1.2.1 General and special linear groups

The largest matrix group for a given matrix dimension is the general linear group.

Definition: The **general linear group** $GL(n)$ is defined by the set of invertible $n \times n$ matrices A under matrix multiplication. If we additionally require that $\det(A) = 1$, the matrices form the **special linear group** $SL(n)$.

The existence of the group identity is guaranteed by the identity matrix I being an invertible matrix (with I as the inverse). Since the existence of an inverse is also necessary in the group definition, we can not construct larger matrix groups. The general linear group also give us our first example of a **non-abelian group**, since matrix multiplication does not in general commute. For two matrices A and B , we may have $AB \neq BA$.

We usually take the matrices in matrix groups to be defined over the field of complex numbers $\mathbb{C}$. If we want to specify the field we may use the notation $GL(n, \mathbb{R})$, signifying that the group is defined over the real numbers. Defined over the complex numbers the $GL(n)$ groups have $2n^2$ free parameters since each of the n^2 elements of the matrices can be a complex number, needing two parameters. The $SL(n)$ group has $2n^2 - 2$ free parameters since the requirement on the determinant fixes both the real and imaginary part of the determinant.

1.2.2 Unitary and special unitary groups

We first remind you that the **Hermitian conjugate** or conjugate transpose of a matrix is given by transposing the matrix and taking the complex conjugate of its elements. Here, we will use the symbol $\dagger$ for this operation, so that for a matrix A , $A^\dagger = (A^T)^*$.

We now define the unitary groups.

Definition: The **unitary group** $U(n)$ is defined by the set of complex unitary $n \times n$ matrices U , *i.e.* matrices such that $U^\dagger U = I$ or $U^{-1} = U^\dagger$. If we additionally require that $\det(U) = 1$ the matrices form the **special unitary group** $SU(n)$.

⁴We could instead have defined the operation between two group elements to be ordinary multiplication and used that to show the relationship $T_a \circ T_b = T_{a+b}$. However, it is important to notice that showing this is not entirely trivial because ordinary arithmetic rules for exponentials fail for operators. In this particular case the proof is fairly simple, but this is in general not so.

Since for $U \in U(n)$,

$$\det(UU^\dagger) = \det(U) \det(U^\dagger) = \det(U) \det(U^T)^* = \det(U) \det(U)^* = \det(I) = 1,$$

we have that the determinant of these matrices must be complex numbers on the unit circle, *i.e.* $\det(U) = e^{i\theta}$. It can be shown (see exercises) that the $U(n)$ groups have n^2 independent parameters, while the $SU(n)$ groups have $n^2 - 1$.

It is these groups that form the gauge symmetry groups of the Standard Model: $SU(3)$, $SU(2)$ and $U(1)$. The group $U(1)$ makes perfect sense despite the odd matrix dimensions. This is simply the set of all complex numbers of unit length with ordinary multiplication, *i.e.* $U(1) = \{e^{i\alpha} | \alpha \in \mathbb{R}\}$, but notice that $SU(1)$ would be trivial since it contains only 1.

The unitary group has the important property that for $\forall \mathbf{x}, \mathbf{y} \in \mathbb{C}^n$ multiplication by a unitary matrix leaves scalar products unchanged. If $\mathbf{x}' = U\mathbf{x}$ and $\mathbf{y}' = U\mathbf{y}$, then

$$\begin{aligned} \mathbf{x}' \cdot \mathbf{y}' &\equiv \mathbf{x}'^\dagger \mathbf{y}' = (U\mathbf{x})^\dagger U\mathbf{y} \\ &= \mathbf{x}^\dagger U^\dagger U \mathbf{y} = \mathbf{x}^\dagger \mathbf{y} = \mathbf{x} \cdot \mathbf{y}. \end{aligned}$$

Thus, its members do not change the length of the vectors they act on.

Since we would like to let our group representations act on vectors that describe quantum mechanical states, the unitary groups then conserve probability for these states. For example, when acting on a complex number (a complex scalar), such as a wavefunction $\psi(x)$, the elements of $U(1)$ rotate the phase of ψ , however, the magnitude is conserved since $\psi' = e^{i\alpha}\psi$ gives $|\psi'|^2 = \psi^* e^{-i\alpha} e^{i\alpha} \psi = |\psi|^2$.

1.2.3 Orthogonal and special orthogonal groups

If we restrict the unitary matrices to be real, we get the orthogonal groups.

Definition: The **orthogonal group** $O(n)$ is the group of real $n \times n$ orthogonal matrices O , *i.e.* matrices where $O^T O = I$. If we additionally require that $\det(O) = 1$ the matrices form the **special orthogonal group** $SO(n)$.

It follows from the definition of the orthogonal group that the determinant of the members is either 1 or -1 , thus the special orthogonal group is simply one half of the members. For $\mathbf{x} \in \mathbb{R}^n$ the orthogonal group has the same property as the unitary group of leaving the length of vectors invariant.

Matrices in the $O(n)$ and $SO(n)$ groups have $n(n-1)/2$ independent parameters since an $n \times n$ matrix with real entries has n^2 elements, and there are $n(n+1)/2$ equations to be satisfied by the orthogonality condition.⁵

The special orthogonal groups $SO(2)$ and $SO(3)$ are much used because their elements represent rotations in two and three dimensions, respectively, while $SO(n)$ extends this to higher dimensions and represents the symmetries of a sphere in n dimensions. To see this we can start from the fact that rotations, by definition, conserve angles and distances (and orientation). This means that the original set of orthogonal axis – or orthogonal basis vectors if you wish – must transform into another orthogonal set of axis under the rotation. The matrix performing the rotation must then be orthogonal, and thus the collection of rotations must be $O(n)$. If we additionally require that orientation is preserved, this removes the matrices with negative determinant, leaving the $SO(n)$ group.

⁵Only the upper triangular part of $O^T O$ has independent equations since $O^T O$ is a symmetric matrix, $(O^T O)^T = O^T (O^T)^T = O^T O$.

$SO(2)$

Given that $SO(2)$ has only one parameter, we can write a general group member R as parameterised by θ

$$R(\theta) = \begin{bmatrix} \cos \theta & -\sin \theta \\ \sin \theta & \cos \theta \end{bmatrix}.$$

As expected, we can recognise this as the matrix of rotations of an angle θ around a point in the plane, and it represents the (orientation preserving) symmetries of a circle. Some tinkering with this representation will show that $SO(2)$ is in fact an abelian group, despite the matrix definition. From a physical viewpoint this should be expected: the order of rotations in the plane should not matter.

It is interesting to observe that the elements of $SO(2)$ rotate points in the plane, while the elements of $U(1)$ rotate complex numbers, which can be represented by points in the plane. Indeed, a one-to-one correspondence can be found between the members of the two groups so that the groups are indeed the same, or, as we say, isomorphic, $SO(2) \cong U(1)$.

 $SO(3)$

This group has three free parameters. Already at this point writing down the explicit form of a general group member is not very enlightening. There are also a number of different conventions in use, so proper care is advised when using results from the literature. In terms of a general rotation in three dimensions this can either be viewed as rotation angles around three fixed axis, or as the fixing of a rotation axis by two angles, with a third rotation angle around that axis.

One particular explicit form, where the angles correspond to the three **Euler angles** of rotation in three dimensions, α , β and γ , is

$$R(\alpha, \beta, \gamma) = \begin{bmatrix} \cos \alpha \cos \beta \cos \gamma - \sin \alpha \sin \gamma & -\cos \alpha \cos \beta \sin \gamma - \sin \alpha \cos \gamma & \cos \alpha \sin \beta \\ \sin \alpha \cos \beta \cos \gamma + \cos \alpha \sin \gamma & -\sin \alpha \cos \beta \sin \gamma + \cos \alpha \cos \gamma & \sin \alpha \sin \beta \\ -\sin \beta \cos \gamma & \sin \beta \sin \gamma & \cos \beta \end{bmatrix},$$

where $0 \leq \alpha, \gamma < 2\pi$ and $0 \leq \beta \leq \pi$. These rotations do not commute, so the group is non-abelian.

We have seen above that $SU(2)$ has three free parameters, just the same as $SO(3)$. You may at this point guess that $SO(3)$ is isomorphic to $SU(2)$. That would be a very good guess, however, it would also be wrong. We will return to this later, but this is one of those things in mathematics that turn out to be disappointingly only almost true.

1.3 Group properties

1.3.1 Subgroups

We now extend our vocabulary for groups by defining the **subgroup** of a group G .

Definition: A subset $H \subset G$ is a **subgroup** if and only if:^a

- i) $h_i \circ h_j \in H$ for $\forall h_i, h_j \in H$, (closure)
- ii) $h_i^{-1} \in H$ for $\forall h_i \in H$. (inverse)

H is a **proper** subgroup if and only if $H \neq G$ and $H \neq \{e\}$.

^aAn alternative, equivalent, and more compact way of writing these two requirements is the single requirement $h_i \circ h_j^{-1} \in H$ for $\forall h_i, h_j \in H$. This is often utilised in proofs.

We have already seen some examples of subgroups: the $SU(n)$ groups are subgroups of $U(n)$, and the $SO(n)$ groups are subgroups of the $O(n)$ groups. This can easily be shown using the properties of determinants.

There is a very important class of subgroup called the **normal** subgroup. The importance will become clear in a moment.

Definition: A subgroup H is a **normal** (invariant) subgroup, if and only if the **conjugation** of any element by any $g \in G$ is in H ,^a meaning

$$ghg^{-1} \in H \text{ for } \forall h \in H.$$

A **simple** group G has no proper normal subgroup. A **semi-simple** group G has no proper abelian normal subgroup.

^aAnother, pretty, but slightly abusive, way of writing the definition of a normal group is to say that $gHg^{-1} = H$. This implies (correctly), that the image of H under the conjugation operation is guaranteed to be the whole of H .

We can for example show that for $n > 1$, $SU(n)$ is a normal subgroup of $U(n)$ (see exercises).

1.3.2 Quotient groups

The normal subgroup can be seen as a factor in the original group that can be divided out to form a simpler group that only retains the structure that was not in the normal group. To be more precise we need the concept of **cosets**.

Definition: A **left coset** of a subgroup $H \subset G$ with respect to $g \in G$ is the set of members $\{gh|h \in H\}$, and a **right coset** of the subgroup is the set $\{hg|h \in H\}$. These are sometimes written gH and Hg , respectively.

For normal subgroups H it can be shown that the sets of left and right cosets coincide and form a group. This is called the **quotient** or **coset group** and denoted G/H .⁶ This has as its members all the *distinct sets* $\{gh|h \in H\}$, that can be generated by a $g \in G$, and has the binary operation $*$ with $\{gh|h \in H\} * \{g'h|h \in H\} = \{(g \circ g')h|h \in H\}$. To simplify notation this can be written $gH * g'H = (g \circ g')H$.

Let us briefly discuss an example of a quotient group. We already know that $SU(n)$ is a normal subgroup to $U(n)$. This means that $U(n)/SU(n)$ is a group. What sort of group is this? Notice that two matrices U_i and U_j live in the same coset of $SU(n)$ if and only if

⁶Sometimes also called the **factor group**. The notation is pronounced “G mod N”, where “mod” is short for modulo.

$\det(U_i U_j^{-1}) = 1$, which is true if and only if $\det(U_i) = \det(U_j)$. In other words, each coset constructed from $SU(n)$ is simply the set of all matrices with a given determinant, which we saw for the $U(n)$ group can be any unit complex number. Observe that this means that many of the cosets are the same set. Here, all cosets generated by members in $U(n)$ with the same determinant is the same coset. When these cosets act on each other with the group operation of $U(n)/SU(n)$ they form new cosets of matrices with a determinant that is the product of their individual determinants, *i.e.* with $U, U' \in U(n)$ we have the product of quotient group member $\{US|S \in SU(n)\} * \{U'S|S \in SU(n)\} = \{UU'S|S \in SU(n)\}$. Thus, the group behaves exactly as $U(1)$, and is in fact isomorphic to it.

1.3.3 Product groups

Now that we have introduced group division, we also need to introduce products of groups.

Definition: The **direct product** of groups G and H , $G \times H$, is defined as the ordered pairs (g, h) where $g \in G$ and $h \in H$, with component-wise operation $(g_i, h_i) \circ (g_j, h_j) = (g_i \circ g_j, h_i \circ h_j)$. $G \times H$ is then a group and G and H can be shown to be normal subgroups of $G \times H$.

We should note there that the subgroups are strictly $G \times \{e_H\}$ and $\{e_G\} \times H$, but these are isomorphic to G and H .

Because it has at least one important guest star appearance in this text we also need the definition of the **semi-direct product**.

Definition: The **semi-direct product** of groups G and H , $G \rtimes H$, where G is also a mapping $G : H \rightarrow H$, is defined by the ordered pairs (g, h) where $g \in G$ and $h \in H$, with component-wise operation $(g_i, h_i) \circ (g_j, h_j) = (g_i \circ g_j, h_i \circ g_i(h_j))$. Here H is not a normal subgroup of $G \rtimes H$, but G is.

Note how the semi-direct product is not symmetric between the factors.

The famous Standard Model gauge group $SU(3) \times SU(2) \times U(1)$ is an example of a direct product. Direct products are “trivial” structures because there is no “interaction” between the subgroups, the elements of each group act only on elements of the same group. This is not true for semi-direct products.

1.3.4 Isomorphic groups

We have already talked about how two groups are the same if they have a correspondence between their members and the same results for the multiplication operation, and we have called this an isometry. Let us now try to put this notion of when two groups are the same into a more formal language.

Definition: Two groups G and H are **homomorphic** if there exists a map between the elements of the groups $\rho : G \rightarrow H$, such that for $\forall g, g' \in G$, $\rho(g \circ g') = \rho(g) \circ \rho(g')$.

For homomorphic groups we say that the mapping conserves the structure of the group, or in other words, all the rules for the group operation/multiplication. This leads to our notion of group equality, namely **isomorphic** groups:

Definition: Two groups G and H are **isomorphic**, written $G \cong H$, if they are homomorphic and the relevant mapping is one-to-one.

The one-to-one mapping ensures that there is a one-to-one correspondence between the elements of the group, so that isomorphic groups effectively contain both the same members and have the same multiplication operation.

For matrix groups, a good way of checking the plausibility of isomorphism is to count the number of free parameters. The difference of the parameters for the two factors should be equal to the number of parameters for the quotient group. In our example $U(n)/SU(n)$, $U(n)$ has n^2 , while $SU(n)$ has $n^2 - 1$, and this gives $n^2 - (n^2 - 1) = 1$ parameters for the quotient group, which matches the one parameter of $U(1)$.

1.4 Representations

In some contrast to the treatment in most introductory group theory texts in mathematics, physicists are mostly interested in the properties of groups G where the elements of G act to transform some elements of a vector space $v \in V$, $g(v) = v' \in V$.⁷ Here, the members of V can for example be the state of a system, say a wave-function in quantum mechanics or a field in quantum field theory. To be useful in physics, we would like that the result of the group operation $g_i \circ g_j$ acts as $(g_i \circ g_j)(v) = g_i(g_j(v))$ and the group identity acts as $e(v) = v$.

We begin with the abstract definition of a representation that ensures these properties.

Definition: A **representation** of a group G on a vector space V is a map $\rho : G \rightarrow GL(V, K)$, where $GL(V, K)$ is the **general linear group** on V , *i.e.* the invertible matrices of the field K of V ,^a such that for $\forall g_i, g_j \in G$,

$$\rho(g_i \circ g_j) = \rho(g_i)\rho(g_j). \text{ (homomorphism)}$$

If this map is also isomorphic, we say that the representation is **faithful**.

^aTechnically, we can only be sure that we can write $GL(V, K)$ as matrices as long as V is a finite dimensional vector space. However, we shall do our best not to ass around with infinite dimensional representations.

From a physicist point of view, the underlying point here is that (the members of) our groups will be used on quantum mechanical states, or fields in field theory, which can be just complex numbers (functions) or multi-component vectors of such. They are thus members of a vector space, and the definition of representations force the transformation properties of the group to be written in terms of matrices. Furthermore, the mapping from the group, or, if you like, the concrete way of writing the abstract group elements, must be homomorphic (structure preserving), meaning that if we can write a group element as the product of two others, the matrix for that element must be the product of the two matrices for the individual group elements it can be written in terms of.

You may by now have realized that the matrix groups defined in Sec. 1.2 have the property that they are defined in terms of one of their representations. These are called the **fundamental** or **defining representations**. However, we will also have use for other representations,

⁷This part of group theory is called *representation theory*.

e.g. the **adjoint representation** which we will introduce later.

Let us now take a few examples that connect to our definition. We saw earlier that for $U(1)$ the group members can be written as the complex numbers on the unit circle $e^{i\theta}$, which can be used as phase transformations on wavefunctions $\psi(x)$ that form a one dimensional vector space over the complex numbers.

For $SU(2)$, a quick count of the number of free parameters in $SU(2)$ should convince us that we need three real numbers to parametrise the group elements, here the α_i : each matrix consists of four complex numbers, or eight real, the unitarity requirement removes four free parameters, while the requirement on the determinant one more. This means that that we should be able to write a general matrix in $SU(2)$ in terms of the linear combination of three unitary “basis” matrices. One common choice for this is the **Pauli matrices**.

$$\sigma_1 = \begin{bmatrix} 0 & 1 \\ 1 & 0 \end{bmatrix}, \quad \sigma_2 = \begin{bmatrix} 0 & -i \\ i & 0 \end{bmatrix}, \quad \sigma_3 = \begin{bmatrix} 1 & 0 \\ 0 & -1 \end{bmatrix}. \quad (1.1)$$

Since the three Pauli matrices are linearly independent, the sum $\alpha_i \sigma_i$, $\alpha_i \in \mathbb{R}$, should then (hopefully) span all of $SU(2)$.⁸ In the Standard Model this representation is applied to weak doublets of fields, *e.g.* the electron–neutrino doublet $\psi = (\nu_e, e)$ that form a two-dimensional vector space, as the $SU(2)_L$ gauge transformation. In the SM the group elements are written as the exponential $e^{i\alpha_i \sigma_i}$. We will return to why we want to write the group as an exponential in Sec. 1.7.

However, we can construct many more representations of a single group such as $SU(2)$. Using the three free parameters in $SU(2)$ it turns out that we can also write the elements in the group in terms of three 3×3 -matrices that act on vectors in a three-dimensional space. For example we can use the matrices

$$J_1 = \begin{bmatrix} 0 & 0 & 0 \\ 0 & 0 & -1 \\ 0 & 1 & 0 \end{bmatrix}, \quad J_2 = \begin{bmatrix} 0 & 0 & 1 \\ 0 & 0 & 0 \\ -1 & 0 & 0 \end{bmatrix}, \quad J_3 = \begin{bmatrix} 0 & -1 & 0 \\ 1 & 0 & 0 \\ 0 & 0 & 0 \end{bmatrix}. \quad (1.2)$$

This forms the adjoint representation of $SU(2)$ and is essentially the representation used in the SM for operations on the gauge fields that transform under $SU(2)_L$ which is a three-vector. The central point here is that the group structure is the same (isometric), even if the objects in the representation are different.

The existence of more representations necessitates a definition of when representations are actually equivalent or isomorphic. This should not be confused with whether groups are isomorphic, but removes differences in representations that are simply due to a change in basis for the vector space the group is acting on, or trivial changes in the dimension of the vector space.⁹

Definition: Two representations ρ and ρ' of G on V and V' are **equivalent** if and only if there exists a map $A : V \rightarrow V'$, that is one-to-one, such that for $\forall g \in G$, $A\rho(g)A^{-1} = \rho'(g)$.

⁸We will later see to what extent this is true.

⁹Imagine, for example, that you create a representation for $U(1)$ that consists of diagonal 2×2 -matrices with the unit complex numbers repeated twice on the diagonal. This is not essentially different from the one-dimensional representation, and should not be considered as such.

1.4.1 Irreducible representations

The building blocks of representations are so-called **irreducible representations**, also called **irreps**. These are the essential ingredients in representation theory, and are defined as follows:

Definition: An **irreducible representation** ρ is a representation where there is *no* proper subspace $W \subset V$ that is closed under the group, *i.e.* there is no $W \subset V$ such that for $\forall w \in W, \forall g \in G$ we have $\rho(g)w \in W$.^a

^aIn other words, we can not split the matrix representation of G in two parts that do not “mix”.

Let us take an example to try to clear up what a *reducible* representation means in contrast to an irreducible. Assume the representation $\rho(g)$ for $g \in G$ acts on a vector space V as matrices. If these matrices $\rho(g)$ can all be decomposed into $\rho_1(g)$, $\rho_2(g)$ and $\rho_{12}(g)$ such that for $\mathbf{v} = (\mathbf{v}_1, \mathbf{v}_2) \in V$

$$\rho(g)\mathbf{v} = \begin{bmatrix} \rho_1(g) & \rho_{12}(g) \\ 0 & \rho_2(g) \end{bmatrix} \begin{bmatrix} \mathbf{v}_1 \\ \mathbf{v}_2 \end{bmatrix},$$

then ρ is **reducible**. The subspace W of V spanned by $\mathbf{v}_1$ violates the irreducibility condition above.

If we also have $\rho_{12}(g) = 0$ we say that the representation is **completely reducible**. It can be shown that in most cases a reducible representation is also completely reducible. In fact, representations for which this is not true tend to be mathematical curiosities. As a result, there is a tendency in physics to use the term “reducible” where we should use the term “completely reducible”. In the case of a completely reducible representation we can split the vector space V into two vector spaces $V = W_1 \oplus W_2$, where $\mathbf{v}_1 \in W_1$ and $\mathbf{v}_2 \in W_2$, and define a representation of G on each of them using ρ_1 and ρ_2 , which in turn could either be reduced more, or would be irreducible.

We end this section with an important theorem that helps us decide whether a representation is irreducible, and ultimately gives a property identifying the representation. As many important theorems, it is called a lemma.

Theorem: (Schur’s Lemma [1])

If we have an irreducible representation ρ of a group G , all the matrices A that commute with $\rho(g)$ for $\forall g \in G$ are proportional to the identity, $A = \lambda I$. Here the λ are constants that label the representation.

1.5 Lie groups

In physics we are particularly interested in a special type of group, the **Lie group**, a class of continuous groups that we can parametrise and which are the basic tool we use to describe continuous symmetries. In order to define Lie groups we will need to use the technical term (smooth) manifold, meaning a mathematical object (formally a topological space) that

locally¹⁰ can be parametrised as a function of $\mathbb{R}^n$ or $\mathbb{C}^n$. We will describe a Lie group G in terms of a parameterisation of the members $g(\mathbf{a}) \in G$, where $\mathbf{a} \in \mathbb{R}^n$ (or $\mathbb{C}^n$). Additionally, in order to describe continuous symmetries these parameterisations need to be *smooth*, also in the technical sense of smooth, which means infinitely differentiable.

Definition: A Lie group G is a finite-dimensional **smooth manifold** where group multiplication and inversion are smooth functions, meaning that given elements $g(\mathbf{a}), g(\mathbf{a}') \in G$, $g(\mathbf{a}') \circ g(\mathbf{a}) = g(\mathbf{b})$ where $\mathbf{b}(\mathbf{a}, \mathbf{a}')$ is a smooth function of $\mathbf{a}$ and $\mathbf{a}'$, and $g^{-1}(\mathbf{a}) = g(\mathbf{b})$ where $\mathbf{b}(\mathbf{a})$ is a smooth function of $\mathbf{a}$.

The situation we are usually interested in is that a Lie group G acting on an m -dimensional vector space V through a representation. Here, it can be shown (for finite-dimensional representations) that we can locally write the map of the representation $G \times V \rightarrow V$ for $\mathbf{x} \in V$ in terms of an explicit function f called the **composition function**. We have $x_i \rightarrow x'_i = f_i(\mathbf{x}, \mathbf{a})$, where the composition function f_i is analytic¹¹ in x_i and a_i . Additionally f_i should have an inverse.

From our earlier example we immediately see that the translation group $T(1)$, given the parameterisation of the elements $g(a) = e^{a \frac{d}{dx}}$, is a Lie group since $g(a) \cdot g(a') = g(b) = g(a+a')$ and $b = a + a'$ is an analytic function of a and a' , and for the inverse $g^{-1}(a) = g(b) = g(-a)$ where $b = -a$ is a smooth function of a . Here, we can write the action of the group on the vector space $\mathbb{R}^1$ as $x' = f(x, a) = x + a$. The matrix groups are also Lie groups as they have $n \times n$ -matrix representations where the matrices A are parametrised by a certain number of parameters $\mathbf{a}$, so that $x'_i = f_i(\mathbf{x}, \mathbf{a}) = [A(\mathbf{a})\mathbf{x}]_i$.

By the analyticity of the explicit function f we can always construct the parameterisation so that the zero parameter corresponds to the identity element of the group, $g(0) = e$, which means that $f_i(\mathbf{x}, 0) = x_i$. By an infinitesimal change da_j of the j -th parameter we then get the following Taylor expansion¹²

$$\begin{aligned} x'_i &= x_i + dx_i = f_i(\mathbf{x}, d\mathbf{a}) \\ &= f_i(\mathbf{x}, 0) + \frac{\partial f_i}{\partial a_j} da_j + \dots \\ &= x_i + da_j \frac{\partial f_i}{\partial a_j}. \end{aligned}$$

This is the result of the transformation by the member of the group that in the parameterisation sits $d\mathbf{a}$ from the identity. For a matrix group the change in the vector space is

$$dx_i = da_j \frac{\partial f_i}{\partial a_j} = da_j \left[\frac{\partial A(\mathbf{a})}{\partial a_j} \mathbf{x} \right]_i = [da_j X_j \mathbf{x}]_i,$$

where

$$X_j \equiv \left. \frac{\partial A(\mathbf{a})}{\partial a_j} \right|_{\mathbf{a}=0},$$

¹⁰This insistence on local means that the parameterisation is not necessarily the same for the whole group.

¹¹Meaning infinitely differentiable and in possession of a convergent Taylor expansion. As a result analytic functions (on $\mathbb{R}$) are smooth, but the reverse does not hold.

¹²The fact that f_i is analytic means that this Taylor expansion must converge in some radius around $f_i(\mathbf{x}, 0)$.

are the n matrix **generators** X_j of the Lie group.

If we now let F be a function from the vector space V that we are interested in, to either the real $\mathbb{R}$ or complex numbers $\mathbb{C}$ – giving for example some interesting physical quantity – then the group transformation defined by da changes F by

$$\begin{aligned} dF &= \frac{\partial F}{\partial x_i} dx_i \\ &= \frac{\partial F}{\partial x_i} \frac{\partial f_i}{\partial a_j} da_j \\ &\equiv da_j X_j F \end{aligned}$$

where the operators defined by

$$X_j \equiv \frac{\partial f_i}{\partial a_j} \frac{\partial}{\partial x_i},$$

are called the n differential **generators** of the Lie group. We see here that the number of generators is the same as the dimension of the manifold, which is also the number of free parameters in the parameterisation of the group. It is these generators X that then define the effect of the Lie group members in a given representation (near the zero parameter), while the n a_j 's are mere parameters. We say that the generators determine the local structure of the group.

As an example of the above we can now go in the opposite direction and look at the two-parameter transformation *defined* by

$$x' = f(x) = a_1 x + a_2,$$

which gives

$$X_1 = \frac{\partial f}{\partial a_1} \frac{\partial}{\partial x} = x \frac{\partial}{\partial x},$$

which is the generator for **dilation** (scale change) in one dimension, and

$$X_2 = \frac{\partial f}{\partial a_2} \frac{\partial}{\partial x} = \frac{\partial}{\partial x},$$

which is the generator for translation, meaning the generator of the group that we have called $T(1)$. These are examples of representations that are not in terms of matrices, but rather **differential representations**. Notice that we can show the following relationship for these two generators: $[X_1, X_2] = X_1 X_2 - X_2 X_1 = -X_2$.

As we have seen, the group $SU(2)$ has three free parameters a_i , so it must have three generators X_i . We can show that the generators for $SU(2)$ in the two-dimensional representation are proportional to the Pauli matrices in (1.1), namely $X_i = \frac{1}{2} \sigma_i$, see exercises. By multiplying out we can also show the following commutation relationships between the generators

$$[X_i, X_j] = i \epsilon_{ijk} X_k. \quad (1.3)$$

These commutators should look familiar to us as they have the same structure as the commutators for the spin S_i and angular momentum operators L_i in quantum mechanics. This is no coincidence, the group that describes rotations is $SO(3)$, and one can show that it has the same generators as $SU(2)$.

In general, the commutator of the generators of a Lie group satisfy

$$[X_i, X_j] = iC_{ij}^k X_k,$$

where C_{ij}^k are called the **structure constants** of the group.¹³ We can easily see that these are antisymmetric in i and j , $C_{ij}^k = -C_{ji}^k$. Lie also showed [2] that there is a **Jacobi identity** among the generators,

$$[X_i, [X_j, X_k]] + [X_j, [X_k, X_i]] + [X_k, [X_i, X_j]] = 0. \quad (1.4)$$

This immediately leads to the following identity for the structure constants: $C_{ij}^k C_{kl}^m + C_{jl}^k C_{ki}^m + C_{li}^k C_{kj}^m = 0$.

We discussed the defining or fundamental representation of a matrix based group earlier. These representations usually have the lowest possible dimension. The **adjoint representation** consists of the n matrices M_i with elements:

$$(M_i)_{jk} = -iC_{ij}^k,$$

where C_{ij}^k are the structure constants. From the Jacobi identity it follows that $[M_i, M_j] = iC_{ij}^k M_k$, meaning that the matrices of the adjoint representation fulfils the same commutation relationship as the fundamental (generators). Note that the dimension of the adjoint representation for $SO(m)$ and $SU(m)$ is equal to the number of independent parameters, $m(m-1)/2$ and $m^2 - 1$, respectively.

We can now briefly return to the $SU(2)$ example. With structure constants $C_{ij}^k = \epsilon_{ijk}$ we get the elements of the adjoint matrices $(M_i)_{jk} = -i\epsilon_{ijk}$, which gives $M_i = iJ_i$, with the J_i being the matrices in (1.2).

1.6 Algebras

To further study the structure of groups we begin by defining an **algebra**. An algebra extends the familiar structure of vector spaces by adding a multiplication operation for the vectors which gives a new vector.

Definition: An **algebra** A over a field (say $\mathbb{R}$ or $\mathbb{C}$) is a linear vector space with a binary (multiplication) operation $\circ : A \times A \rightarrow A$.

It is important to remember here that the vector space part of the definition implies that there is field, so for example from $x \in A$ and $a \in \mathbb{R}$ (as the field) we can always form new members $ax \in A$.

As a very simple example, the vector space $\mathbb{R}^3$ together with the standard cross-product constitutes an algebra since the cross product results in a new vector in $\mathbb{R}^3$. Even more trivially perhaps is that $\mathbb{R}$ with ordinary multiplication as the binary operation fulfils the algebra requirements (more on this case below).

¹³There is an annoying difference in notation here between physics and mathematics, where the i is commonly dropped.

1.6.1 Normed division algebras*

As a slight aside to the main argument of the text, we want to give an interesting example of algebras. We start with a **division algebra**, which informally is an algebra where the binary operation of the algebra also has a meaningful (implicit) concept of division for all members (except division by zero).

Definition: An algebra D is a **division algebra** if for any element x in D and any non-zero element y in D there exists precisely one element z in D with $x = y \circ z$ and precisely one element z' in D such that $x = z' \circ y$. In this sense y is a divisor of x .

We see that division algebras have the addition and subtraction properties of “ordinary numbers” (reals) since they are vector spaces, and they have the multiplication (algebra) and division (division algebra) properties of the reals as well. So, in a sense, division algebras are structures close to the reals in terms of properties – and, of course, the reals are again an example of a division algebra.

We can now add to the division algebra the notion of a **norm**, or length, of the members $\|x\|$. This is a map from the algebra to the field, $\|\cdot\| : D \rightarrow \mathbb{R}$, so that we can discuss for example convergence of the members as we do for the reals. We have to require here that the norm is homomorphic, meaning that it preserves the structure of the algebra, so that for example the “product” of two objects with large norm has a large norm. We then get a **normed division algebra**.

Definition: If there exists a homomorphic norm for the division algebra, *i.e.* one where $\|x \circ y\| = \|x\|\|y\|$ for all $x, y \in D$, then the division algebra is a **normed division algebra**.

There is an important theorem by Hurwitz (1923) that demonstrates that only four of these real number “lookalikes” exist.

Theorem: Hurwitz’s theorem. There are only four normed division algebras over the reals (up to isomorphism), the reals themselves $\mathbb{R}$, the complex numbers $\mathbb{C}$, the quaternions $\mathbb{H}$, and the octonions $\mathbb{O}$.

In addition it is (relatively speaking) easy to show that $\mathbb{R} \subset \mathbb{C} \subset \mathbb{H} \subset \mathbb{O}$. One perspective on the relationship between these algebras is that the reals is the only ordered normed division algebra, *i.e.*, where we can compare uniquely the elements $a > b$. This is not possible for the complex numbers, but they keep the commutative and associate properties of the reals. The quaternions in turn break the commutativity of the complex numbers, while being associative, while the most unruly of the bunch, the octonions, are not even associative.

1.7 Lie algebras

We will now turn to the most crucial type of algebras for physics, namely **Lie algebras**. To distinguish these from the more general algebras that we have introduced above, we will use

the notation $[\cdot, \cdot]$ for the binary operation in Lie algebras.

Definition: A **Lie algebra** L is an algebra where the binary operator $[\cdot, \cdot]$, called the **Lie bracket**, has the properties that for $x, y, z \in L$ and $a, b \in \mathbb{R}$ (or $\mathbb{C}$):

i) (bilinearity)

$$[ax + by, z] = a[x, z] + b[y, z]$$

$$[z, ax + by] = a[z, x] + b[z, y]$$

ii) (anti-commutation)

$$[x, y] = -[y, x]$$

iii) (Jacobi identity)

$$[x, [y, z]] + [y, [z, x]] + [z, [x, y]] = 0$$

If the algebra is over $\mathbb{R}$ it is a real Lie algebra, and over $\mathbb{C}$ a complex Lie algebra.

Again the vectors of $\mathbb{R}^3$ with the Lie bracket defined in terms of the cross product, $[\mathbf{x}, \mathbf{y}] = \mathbf{x} \times \mathbf{y}$, is a simple example of a Lie algebra. However, we usually restrict ourselves to algebras of linear operators where the Lie bracket is the commutator $[x, y] = xy - yx$, where the defining properties follow automatically. Thus also explaining the notation that we have used for the binary operator.

From what we learnt in Section 1.5 the generators of an n -dimensional Lie group then span a unique n -dimensional Lie algebra where a general element X of the algebra can be written in terms of the generators X_i , as $X = a_i X_i$. However, the reverse is not true. There can be multiple Lie groups that have the same algebra. The often quoted example is $SO(3)$ and $SU(2)$. Here the generators are the same, however, while closely related, the groups are not isomorphic. We will denote the Lie algebra of a matrix group using lowercase letters, for example the algebra of $SU(n)$ is $su(n)$.¹⁴

1.7.1 Exponential map

As we also discussed in Section 1.5, the generators describe the local structure of the group. We can now finally look at how the group (and matrix representation) is reconstructed from the algebra. For this we use what is called the **exponential map**.

Definition: The **exponential map** from the Lie algebra L of a matrix group G to G is defined by $\exp : L \rightarrow G$, where for $X \in L$ we get the element $g \in G$ given by

$$g = \exp(iX) \equiv \sum_{n=0}^{\infty} \frac{(ia_i X_i)^n}{n!}. \quad (1.5)$$

The infinite sum in this definition is nothing more than the formal series definition of an exponential of a matrix. So, why use the exponential? One reason is that it gives us a

¹⁴In some literature so-called **fraktur** is used to represent the algebra, e.g. $\mathfrak{su}(2)$, as this is butt-ugly and unreadable we will seek to avoid it.

natural expansion of the group around the identity element. With the parameters $\mathbf{a} = 0$ the resulting matrix is the identity matrix, which means that the group element is the identity element $g = e$. Similarly, for an infinitesimal $d\mathbf{a}$ we have the group elements near the identity, $g = I + ida_i X_i$.

The question of whether the exponential map reaches all of the members of the group, *i.e.* that it is surjective on G , depends on the properties of the group. What we do know is that locally, meaning sufficiently close to the identity group element, the exponential map generates the group.

For groups that can not be written as matrices the exponential map can be generalised, however, this is somewhat beyond the scope of these notes.

Let us end here by returning to some of our examples. For $U(1)$ we saw that we could write a generic group member as $e^{i\theta}$. Comparing to the exponential map we see that the single generator must here simply be 1, while θ is the parameter. For $SU(2)$ the generators were the Pauli matrices σ_i , and the exponential map is thus, as we alluded to earlier, $e^{i\alpha_i \sigma_i}$, where α_i are the parameters.

For the translation group $T(1)$ we saw that a group member could be written as $e^{a \frac{d}{dx}} = e^{ia(-i \frac{d}{dx})}$. Thus the generator is the differential operator $P = -i \frac{d}{dx}$, which we should recognise from quantum mechanics as the momentum operator, suggesting a connection between translations and momentum. (And indeed, Noether's theorem tells us that it is symmetry under translations that leads to momentum conservation.)

1.8 Exercises

Exercise 1.1

Show that $T_a^{-1} = T_{-a}$ and that $T(1)$ is group.

Exercise 1.2

What are the elements of $O(1)$? Can you find a group that it is isomorphic to?

Exercise 1.3

Show that the $U(n)$ and $SU(n)$ groups have n^2 and $n^2 - 1$ independent (real) parameters, respectively. *Hint:* Consider the fact that $M = U^\dagger U$ is a hermitian matrix, *i.e.* $M^\dagger = M$.

Exercise 1.4

Show that for a subset $H \subset G$, if $h_i \circ h_j^{-1} \in H$ for $\forall h_i, h_j \in H$, then H is a subgroup of G .

Exercise 1.5

Show that if H is a subgroup of G , then $h_i \circ h_j^{-1} \in H$ for $\forall h_i, h_j \in H$.

Exercise 1.6

Show that $SU(n)$ is a proper subgroup of $U(n)$ and that $U(n)$ is not simple.

Exercise 1.7

If H is a normal subgroup of G , show that its left and right cosets are the same, and show that the set formed of the cosets is a group under an appropriate group operation.

Exercise 1.8

Show that the factors in a direct product of groups are normal groups to the product.

Exercise 1.9

Show that the group G and the group $(G \times H)/H$ are isomorphic.

Exercise 1.10

Show that $U(1) \cong \mathbb{R}/\mathbb{Z} \cong SO(2)$.

Exercise 1.11

Find the dimensions of the fundamental and adjoint representations of $SU(n)$.

Exercise 1.12

Find the fundamental representation for $SO(3)$ and the adjoint representation for $SU(2)$. What does this say about the groups and their algebras?

Exercise 1.13

Find an expression for the generators of $SU(2)$ and their commutation relationships. *Hint:* One answer uses a composition function but this approach has some dangers, try also to derive the properties of the generators from a member of the group an infinitesimal distance from the identity.

Exercise 1.14

What are the structure constants of $SU(2)$?

Exercise 1.15

Let A be an algebra based on a finite-dimensional vector space over a field F , with a basis $B = \{b_i | i = 1, \dots, n\}$. Show that the multiplication of elements in A is completely determined by the n^2 products $b_i b_j$ for each pair of basis vectors in B .

Exercise 1.16

Let V be a finite-dimensional vector space over a field F with a basis $B = \{b_i | i = 1, \dots, n\}$. Let $\{c_{rst} | r, s, t = 1, \dots, n\}$ be a collection of n^3 elements in F . Show that there exists one, and only one, multiplication operation on V so that V is an algebra over F under this multiplication and

$$b_r b_s = c_{rst} b_t,$$

for every pair of basis vectors in B . The elements c_{rst} are the structure constants of the algebra.

Exercise 1.17

Show that $\mathbb{R}^3$ with the binary operator $[\mathbf{x}, \mathbf{y}] = \mathbf{x} \times \mathbf{y}$ is a Lie algebra.

Exercise 1.18

Let A_i be the generators of the group G and B_j be the generators of group H . Explain in what sense the A_i and B_j are generators of the direct product group $G \times H$ and show that $[A_i, B_j] = 0$.

Exercise 1.19 Properties of the exponential map

Prove the following useful properties of the exponential map. A and B are matrices.

1. If A and B commute, $e^{A+B} = e^A e^B$.
2. If B has an inverse, $e^{BAB^{-1}} = B e^A B^{-1}$.
3. $e^{A^*} = (e^A)^*$
4. $e^{A^T} = (e^A)^T$
5. $e^{A^\dagger} = (e^A)^\dagger$
6. $e^{-A} = (e^A)^{-1}$
7. $\det e^A = e^{\text{Tr}[A]}$.

Chapter 2

The Poincaré algebra and its extensions

We now take a look at the symmetry groups behind Special Relativity (SR), the Lorentz and Poincaré groups. We will first see what sort of states transform properly under SR, which has surprising connections to already familiar physics. We will then look for ways to extend these external symmetries of the coordinates to internal symmetries of quantum fields, *i.e.* the symmetries of gauge groups.

2.1 The Lorentz Group

Einstein's requirement in Special Relativity was that the laws of physics should be invariant under rotations and/or boosts (changes of velocity) between different inertial reference frames. A point in the Minkowski space-time manifold $\mathbb{M}_4$ is given by a four-vector $x^\mu = (t, x, y, z)$. The resulting allowed transformations of the space-time coordinates are captured in the Lorentz group.

Definition: The **Lorentz group** is the group of linear transformations $x^\mu \rightarrow x'^\mu = \Lambda^\mu{}_\nu x^\nu$ such that $x^2 \equiv x_\mu x^\mu = x'_\mu x'^\mu$ is invariant. The **proper orthochronous** or **restricted Lorentz group** is a subgroup of the Lorentz group where $\det \Lambda = 1$ (**proper**) and $\Lambda^0{}_0 \geq 1$ (**orthochronous**).

The physical interpretation of the orthochronous property is that it keeps the direction (sign) of time of the four vector, while a proper group preserves orientation in rotations.

Since the definition of the Lorentz group effectively gives a composition function we can easily conclude that it is a Lie group. In fact, if we allow for a slight extension of the orthogonal group $O(n)$ to the **indefinite orthogonal group** $O(m, n)$, where instead of the orthogonality property for group members O , $O^{-1} = O^T$, we demand $O^{-1} = g^{-1} O^T g$ where

$$g = \text{diag}(\underbrace{1, \dots, 1}_n, \underbrace{-1, \dots, -1}_m),$$

is the “metric”,¹ then we can write the Lorentz group as $SO^+(1, 3)$, where the plus sign signi-

¹Indeed, we can recognise this matrix relationship as one of the defining (necessary) properties of Lorentz transformations $\Lambda^T g \Lambda = g$.

fies the orthochronous property. The counting of the free parameters of $SO(n, m)$ works just as for $SO(n)$, giving a total of six free parameters for $SO^+(1, 3)$. Physically, we can identify these with the three parameters needed to specify a general rotation in three dimensions, and the three parameters needed to specify a boost (the velocity components).

Since the rotation operations are closed, *i.e.* two rotations result in another rotation, one can prove that this forms a subgroup of $SO^+(1, 3)$. We have earlier claimed that the generators of $SO(3)$ (rotations in three dimensions) are identical to the generators of $SU(2)$. This now allows us to identify three of the generators of $SO^+(1, 3)$ as the J_i that fulfil the $su(2)$ algebra

$$[J_i, J_j] = i\epsilon_{ijk}J_k. \quad (2.1)$$

The boost operations are not closed, and one can show that their generators K_i (exercises) have the following relationships with the rotation generators

$$[K_j, J_i] = i\epsilon_{ijk}K_k, \quad (2.2)$$

$$[K_i, K_j] = -i\epsilon_{ijk}J_k. \quad (2.3)$$

where (2.1) and (2.3) then defines the complete algebra of $SO^+(1, 3)$.

To simplify notation these generators can further be structured into a matrix M given by

$$M = \begin{bmatrix} 0 & -K_1 & -K_2 & -K_3 \\ K_1 & 0 & J_3 & -J_2 \\ K_2 & -J_3 & 0 & J_1 \\ K_3 & J_2 & -J_1 & 0 \end{bmatrix}.$$

In terms of M the commutation relations (2.1) and (2.3) can be written:

$$[M_{\mu\nu}, M_{\rho\sigma}] = -i(g_{\mu\rho}M_{\nu\sigma} - g_{\mu\sigma}M_{\nu\rho} - g_{\nu\rho}M_{\mu\sigma} + g_{\nu\sigma}M_{\mu\rho}). \quad (2.4)$$

From the discussion in Sec. 1.7 and here, any $\Lambda \in SO^+(1, 3)$ can now be written as

$$\Lambda^\mu{}_\nu = \left[\exp \left(\frac{i}{2} \omega^{\rho\sigma} M_{\rho\sigma} \right) \right]^\mu{}_\nu, \quad (2.5)$$

where $\omega_{\rho\sigma} = -\omega_{\sigma\rho}$ are the six free parameters of the transformation and $M_{\rho\sigma}$ are the generators of the group $SO(1, 3)$ and form the basis of the Lie algebra for $O(1, 3)$. In fact, this is also the algebra of $O(1, 3)$ since the orthochronous and proper requirements do not change the number of free parameters, but rather restricts us to a subset of the matrices. This also nicely illustrates the local property of the exponential map: using these generators we can in fact not get outside of the $SO^+(1, 3)$ subgroup of $O(1, 3)$. The larger group $O(1, 3)$ can be seen as four connected components with $\det \Lambda = \pm 1$ and $|\Lambda^0{}_0| \geq 1$ that are joined by the time T and parity P inversion operators.

2.2 The Poincaré group

We can now extend $O(1, 3)$ by adding translation by a constant four-vector a^μ to the transformation of the Lorentz group: $x^\mu \rightarrow x'^\mu = \Lambda^\mu{}_\nu x^\nu + a^\mu$. This transformation leaves lengths $(x - y)^2$ invariant in $\mathbb{M}_4$.

Definition: The **Poincaré group** is the group of all transformations of the form

$$x^\mu \rightarrow x'^\mu = \Lambda^\mu_\nu x^\nu + a^\mu.$$

We can also construct the **restricted Poincaré group** by restricting the matrices Λ in the same way as in $SO^+(1, 3)$.

Writing a group member in terms of its parameters (Λ, a) , we can see from the explicit form of the transformation that the composition of two elements in this group is:

$$(\Lambda_1, a_1) \circ (\Lambda_2, a_2) = (\Lambda_1 \Lambda_2, \Lambda_1 a_2 + a_1).$$

This tells us that the Poincaré group is *not* a direct product of the Lorentz group and the translation group, but rather a *semi-direct product* of $O(1, 3)$ and the (indefinite) translation group $T(1, 3)$, $O(1, 3) \ltimes T(1, 3)$. The translation group is a normal subgroup, and while the Lorentz group is a subgroup, it is not normal. The restricted Poincaré group is written in the same way as $SO^+(1, 3) \ltimes T(1, 3)$.

The translation part of the Poincaré group adds four parameters to the six parameters of the rotations and boosts. This means that there are four more generators compared to the Lorentz group. Given our earlier discussion of the translation group in Sec. 1.7.1 we can convince ourselves that we can use the momentum operators $P_\mu = -i\partial_\mu$. These generators have a trivial commutation relationship:

$$[P_\mu, P_\nu] = 0. \quad (2.6)$$

Finally, one can show the following commutators with the generators of the Lorentz group:²

$$[M_{\mu\nu}, P_\rho] = -i(g_{\mu\rho}P_\nu - g_{\nu\rho}P_\mu). \quad (2.7)$$

Equations (2.4), (2.6) and (2.7) together form the **Poincaré algebra**, a Lie algebra. This allows us to write a general member g of the restricted Poincaré group by using the exponential map

$$g = \exp \left(\frac{i}{2} \omega^{\rho\sigma} M_{\rho\sigma} + i a^\mu P_\mu \right), \quad (2.8)$$

where a^μ are the additional parameters of the translation.

2.3 Irreducible representations of the Poincaré group

We would now like to ask the question: what sort of particles, or, if you like, quantum fields, can exist if we require that they are representations of the Poincaré group?³

To answer that question we will need to classify all the irreducible representations of the Poincaré group. This seems like a dramatically difficult task, however, we can now use Schur's lemma that we saw in Sec. 1.4.1. To do this we need to find the **Casimir operators** of the algebra.

²For a rigorous derivation of this see Chapter 1.2 of [3]. The proof is constructed by looking at the infinitesimal action of the generators.

³In the sense of being described by a vector space that the group representations act on.

Definition: The **Casimir operators** of a Lie algebra are the elements that commute with all other elements of the algebra

From Schur's lemma the Casimir operators should then be proportional to the identity for irreducible representations, and, most importantly, the constants of proportionality classify the irrep.

Let us take an example to demonstrate how this works. We saw earlier that $SO(3)$ and $SU(2)$ had the same algebra, with three members that we can generically write as J_i . We can show that $J^2 = J_1^2 + J_2^2 + J_3^2$ is a Casimir invariant of this algebra, meaning that in a given representation, we can write $J^2 = \lambda I$, where λ is this constant and I is the identity matrix. It may not surprise you to find out that the constant is $\lambda = \ell(\ell + 1)$, where ℓ can take half-integer values, meaning that we have the relationship $J^2 = \ell(\ell + 1)I$ that is familiar from quantum mechanics. We can now go back and test this, checking the relationship for the Pauli matrices and the J_i matrices in (1.2). We will find that if $J_i = \frac{1}{2}\sigma_i$, where σ_i are the Pauli matrices, $J^2 = \frac{3}{4}$, corresponding to $\ell = \frac{1}{2}$, and for the J_i matrices, $\ell = 1$. The point here is that ℓ labels the representation, here the spin- $\frac{1}{2}$ and spin-1 representations.

We can now use the constants of proportionality to classify the (irreducible) representations of our Lie algebra (and group). For the Poincaré algebra $P^2 = P_\mu P^\mu$ is a Casimir operator because the following holds:

$$[P_\mu, P^2] = 0, \quad (2.9)$$

$$[M_{\mu\nu}, P^2] = 0. \quad (2.10)$$

This allows us to label the irreducible representation of the Poincaré group with a quantum number that we will (randomly, or maybe not) name $m^2 \in \mathbb{R}$, writing a corresponding state in the vector space as $|m\rangle$, such that:⁴

$$P^2|m\rangle = m^2|m\rangle.$$

If we go to the rest frame of a particle the state has eigenvalues $(m, \vec{0})$ for the operator P_μ , where m is the mass (rest energy) of the particle.⁵ This demonstrates that the label m^2 can indeed be interpreted as the (square) of the mass,

The number of Casimir operators is equal to the **rank** of the algebra, *e.g.* the rank of $su(n)$ is $n - 1$. It turns out that the Poincaré algebra has rank 2, and thus two Casimir operators. To demonstrate this is rather involved, and we will not make an attempt here, but note that it can be shown that $SO^+(1, 3)$ is homomorphic to $SU(2) \times SU(2)$, because of the structure of the boost and rotation generators, where the algebra of each $SU(2)$ has rank 1.

So, what is the second Casimir of the Poincaré algebra?

Definition: We define the **Pauli-Ljubanski polarisation vector** by:

$$W_\mu \equiv \frac{1}{2} \epsilon_{\mu\nu\rho\sigma} P^\nu M^{\rho\sigma}. \quad (2.11)$$

where $\epsilon_{\mu\nu\rho\sigma}$ is the totally antisymmetric Levi-Civita tensor with $\epsilon_{0123} = 1$.

⁴Note that in general m^2 is not restricted to be larger than zero.

⁵This does not lose generality since physics should be independent of frame, however, this argument needs to be modified somewhat for massless particles.

We can show that this vector is translation invariant, *i.e.* that it commutes with the translation operator,

$$[P_\mu, W_\nu] = 0. \quad (2.12)$$

Then $W^2 = W_\mu W^\mu$ is a Casimir operator of the Poincaré algebra since we can show that

$$[P_\mu, W^2] = 0. \quad (2.13)$$

$$[M_{\mu\nu}, W^2] = 0, \quad (2.14)$$

Note that the second of these relationships is not trivial to demonstrate. See [3] for a complete proof.

If we again look at the situation in the rest frame we can write

$$W_i = \frac{1}{2} \epsilon_{i0jk} m M^{jk} = m S_i, \quad (2.15)$$

where $S_i = \frac{1}{2} \epsilon_{ijk} M^{jk}$ is the **spin operator**.⁶ By showing that $WP = 0$ we also have $W_0 = 0$ in this reference frame. This gives $W^2 = -\mathbf{W}^2 = -m^2 \mathbf{S}^2$. Since the spin operator acts on a state with spin s as $\mathbf{S}^2|s\rangle = s(s+1)|s\rangle$, we have that

$$W^2|m, s\rangle = -m^2 s(s+1)|m, s\rangle.$$

In addition this, also leads to the immediate conclusion that for each irrep with spin s there are $2s+1$ states with spin component $s_3 = 0, \frac{1}{2}, 1, \frac{3}{2}, \dots$

The conclusion of this subsection is that anything transforming under the Poincaré group, meaning the objects considered by special relativity, can be classified by two quantum numbers: mass m^2 and spin s .

What do these (irreducible) representations then look like? If we start with spin-0 representations, $s = 0$, we can write the corresponding states without any vector structure as $|m, 0\rangle \sim e^{\pm i p x}$, where p_μ is the four-momentum of the particle, since $P^2|m, 0\rangle = -\partial_\mu \partial^\mu |m, 0\rangle \sim p^2 |m, 0\rangle = m^2 |m, 0\rangle$. This exponential part of states can then always be used to take care of the eigenvalues of the P^2 -Casimir, and is often just implicitly implied in the states/fields.

We can also immediately write down the $s = 1$ vector representation of the Poincaré group, $|m, 1\rangle \sim \epsilon_\mu e^{i p x}$. We simply use a four-vector ϵ_μ that transforms under the fundamental (four-dimensional) representation of the Lorentz group $SO^+(1, 3)$. In order to fulfil the eigenvalue equation of the W^2 -Casimir, and describe the three spin states $s_3 = -1, 0, 1$, this vector (called the polarisation vector) needs to fulfil certain requirements which we do not detail here (see a course on quantum field theory).

However, to find a spin- $\frac{1}{2}$ representation we need to take some more care. In fact, we will find representations both in four and two dimensions. For those familiar with quantum field theory, these will as expected be the Dirac and Weyl spinor representations.

⁶Observe that this discussion is problematic for massless particles. However, it is possible to find a similar relation for massless particles, when we chose a frame where the momentum of the particle is in a single direction.

2.4 Weyl spinors

Interestingly, there exists a homomorphism between the groups $SO^+(1, 3)$ and $SL(2, \mathbb{C})$. This homomorphism, with $\Lambda^\mu{}_\nu \in SO^+(1, 3)$ and $M \in SL(2, \mathbb{C})$, can be explicitly given by:⁷

$$\Lambda^\mu{}_\nu(M) = \frac{1}{2} \text{Tr}[\bar{\sigma}^\mu M \sigma_\nu M^\dagger], \quad (2.16)$$

$$M(\Lambda^\mu{}_\nu) = \pm \frac{1}{\sqrt{\det(\Lambda^\mu{}_\nu \sigma_\mu \bar{\sigma}^\nu)}} \Lambda^\mu{}_\nu \sigma_\mu \bar{\sigma}^\nu, \quad (2.17)$$

where $\bar{\sigma}^\mu = (1, -\vec{\sigma})$ and $\sigma^\mu = (1, \vec{\sigma})$.

This two-to-one correspondence means that $SO^+(1, 3) \cong SL(2, \mathbb{C})/\mathbb{Z}_2$. Thus we can look at the representations of $SL(2, \mathbb{C})$ instead of the Lorentz group, when we describe particles, but what are those representations? It turns out that there exist two inequivalent fundamental representations ρ of $SL(2, \mathbb{C})$:

- i) The self-representation $\rho(M) = M$ acting on an element ψ of a representation vector space V :

$$\psi'_A = M_A{}^B \psi_B, \quad A, B = 1, 2.$$

- ii) The complex conjugate self-representation $\rho(M) = M^*$ working on a vector $\bar{\psi}$ in a space $\bar{V}$:

$$\bar{\psi}'_{\dot{A}} = (M^*)_{\dot{A}}{}^{\dot{B}} \bar{\psi}_{\dot{B}}, \quad \dot{A}, \dot{B} = 1, 2.$$

The vectors ψ and $\bar{\psi}$ in these representation spaces are called, respectively, **left- and right-handed Weyl spinors**.

The indices here follow the same summation rules as four-vectors. Indices can be lowered and raised with:

$$\epsilon_{AB} = \epsilon_{\dot{A}\dot{B}} = \begin{pmatrix} 0 & -1 \\ 1 & 0 \end{pmatrix}, \quad (2.18)$$

$$\epsilon^{AB} = \epsilon^{\dot{A}\dot{B}} = \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}. \quad (2.19)$$

The dots on the indices for the complex conjugate representation are there to help us remember which representation we are using and does not carry any additional importance. Since (2.17) gives M in terms of the Pauli matrices, their index structure must be $\bar{\sigma}^{\mu\dot{A}A}$ and $\sigma^\mu_{A\dot{A}}$. For a consistent index notation, the relationship between the vectors ψ and $\bar{\psi}$ can be expressed with:

$$\bar{\sigma}^{0\dot{A}A}(\psi_A)^* = \bar{\psi}^{\dot{A}}.$$

This may be seen as a bit of an overkill in indices as $\bar{\sigma}^{0\dot{A}A} = \delta^{\dot{A}A}$, and we will in the following often omit the matrix and simply write $(\psi_A)^* = \bar{\psi}^{\dot{A}}$. Note that from the above the following relationships hold for the hermitian conjugate:

$$(\psi_A)^\dagger = \bar{\psi}_{\dot{A}}$$

⁷The choice of sign in Eq. (2.17) is the reason that this is a homomorphism, instead of an isomorphism. Each element in $SO^+(1, 3)$ can be assigned to two in $SL(2, \mathbb{C})$.

$$(\bar{\psi}_{\dot{A}})^\dagger = \psi_A.$$

We further define contractions of Weyl spinors that are invariant under $SL(2, \mathbb{C})$ transformations – just as contractions of four-vectors are invariant under Lorentz transformations – as follows:

Definition: The contraction of two Weyl spinors is given by $\psi\chi \equiv \psi^A\chi_A$ and $\bar{\psi}\bar{\chi} \equiv \bar{\psi}_{\dot{A}}\bar{\chi}^{\dot{A}}$.

With this in hand we see that

$$\psi^2 \equiv \psi\psi = \psi^A\psi_A = \epsilon^{AB}\psi_B\psi_A = \epsilon^{12}\psi_2\psi_1 + \epsilon^{21}\psi_1\psi_2 = \psi_2\psi_1 - \psi_1\psi_2.$$

This quantity is zero if the Weyl spinors commute. In order to avoid this we make the following assumption which is consistent with how we treat fermions as anti-commuting operators:

Postulate: All Weyl spinors anticommute:^a $\{\psi_A, \psi_B\} = \{\bar{\psi}_{\dot{A}}, \bar{\psi}_{\dot{B}}\} = \{\psi_A, \bar{\psi}_{\dot{B}}\} = \{\bar{\psi}_{\dot{A}}, \psi_B\} = 0$.

^aThis means that Weyl spinors are so-called **Grassmann numbers**.

This means that the contraction evaluates as

$$\psi^2 \equiv \psi\psi = \psi^A\psi_A = -2\psi_1\psi_2.$$

Because of (2.17) we can now find how the Lorentz part of the Poincaré group acts on the Weyl spinors, and can use this to show that they fulfil the requirements of the spin- $\frac{1}{2}$ representation of the Poincaré group. In the rest frame of a particle, this is relatively straight forward since the spin operators S_i in (2.15) in the fundamental representation of $SL(2, \mathbb{C})$ can be written in terms of the Pauli matrices σ_i as $S_i = \frac{1}{2}\sigma_i$, and we already know that $\sigma^2 = 3I$, so $S^2 = \frac{3}{4}I$, which corresponds to $s = \frac{1}{2}$.

The Weyl spinors can in turn be used in a four-dimensional representation of the Poincaré group, stacking two Weyl spinors into a four-component Dirac spinor ψ_a :

$$\psi_a = \begin{pmatrix} \psi_A \\ \bar{\chi}^{\dot{A}} \end{pmatrix}.$$

Here, we have in general $(\psi_A)^* \neq \bar{\chi}^{\dot{A}}$. In order to describe a Dirac fermion, which has both particle and antiparticle states, with this Dirac spinor we need two distinct Weyl spinors with different handednesses. For Majorana fermions that are their own antiparticles we have:

$$\psi_a = \begin{pmatrix} \psi_A \\ \bar{\psi}^{\dot{A}} \end{pmatrix}.$$

2.4.1 Useful relationships for Weyl spinors

For Weyl spinors ψ , η , and ϕ we can prove the following relationships⁸

$$\eta\psi = \psi\eta, \quad (2.20)$$

$$\bar{\eta}\bar{\psi} = \bar{\psi}\bar{\eta}, \quad (2.21)$$

$$(\eta\psi)^\dagger = \bar{\psi}\bar{\eta}, \quad (2.22)$$

$$(\eta\psi)(\eta\phi) = \frac{1}{2}(\eta\eta)(\psi\phi) \quad (2.23)$$

$$\eta\sigma^\mu\bar{\psi} = -\bar{\psi}\bar{\sigma}^\mu\eta, \quad (2.24)$$

$$(\sigma^\mu\bar{\eta})_A(\eta\sigma^\nu\bar{\eta}) = \frac{1}{2}g^{\mu\nu}\eta_A(\bar{\eta}\bar{\eta}), \quad (2.25)$$

$$(\eta\sigma^\mu\bar{\eta})(\eta\sigma^\nu\bar{\eta}) = \frac{1}{2}g^{\mu\nu}(\eta\eta)(\bar{\eta}\bar{\eta}), \quad (2.26)$$

$$(\eta\sigma^\mu\partial_\mu\bar{\psi})(\eta\psi) = -\frac{1}{2}(\psi\sigma^\mu\partial_\mu\bar{\psi})(\eta\eta), \quad (2.27)$$

$$(\partial_\mu\psi\sigma^\mu\bar{\eta})(\bar{\eta}\bar{\psi}) = -\frac{1}{2}(\partial_\mu\psi\sigma^\mu\bar{\psi})(\bar{\eta}\bar{\eta}). \quad (2.28)$$

$$\eta\sigma^{\mu\nu}\psi = -\psi\sigma^{\mu\nu}\eta \quad (2.29)$$

Here $\sigma^{\mu\nu} = \frac{i}{4}(\sigma^\mu\bar{\sigma}^\nu - \sigma^\nu\bar{\sigma}^\mu)$.

2.5 The no-go theorem and graded Lie algebras

The Poincaré group contains the complete set of transformations for the symmetries of special relativity (invariance under rotations, translations and boosts), and we have seen that this implies certain properties for the particles, or rather fields, that want to live in representations of the Poincaré group. At the same time we know that the quantum fields have (internal) gauge symmetries. It would then be tempting to ask if these are somehow related and can be described in a larger symmetry.

Unfortunately, the answer to that question is ‘no’, at least as long as we keep to describing our symmetries using Lie algebras. In 1967 Coleman and Mandula [4] showed that under reasonable assumptions any extension of the restricted Poincaré group P to include gauge symmetries is isomorphic to $G_{\text{gauge}} \times P$. A direct product like this means that the generators of the two groups all commute, meaning that the generators B_i of the standard model gauge groups all have

$$[P_\mu, B_i] = [M_{\mu\nu}, B_i] = 0.$$

The result is that there can be no real interaction between the external and internal symmetries.

Not to be defeated by a simple mathematical proof, in 1975 Haag, Łopuszański and Sohnius (HLS) [5] showed that there is a way around Coleman and Mandula’s no-go theorem, if one introduces the concept of $\mathbb{Z}_2$ **graded Lie superalgebras**.⁹

⁸For clarity we have inserted parenthesis to show the different contractions.

⁹The definition of graded Lie algebras can be extended to $\mathbb{Z}_n$ by a direct sum over n vector spaces L_i , $L = \bigoplus_{i=0}^{n-1} L_i$, such that $x_i \circ x_j \in L_{i+j \bmod n}$, with the same requirements for supersymmetrization and Jacobi identity as for the $\mathbb{Z}_2$ graded algebra.

Definition: A **graded Lie superalgebra** is a vector space L that is a direct sum of two vector spaces L_0 and L_1 , $L = L_0 \oplus L_1$, with a binary operation $\circ : L \times L \rightarrow L$ such that for $\forall x_i \in L_i$

- i) $x_i \circ x_j \in L_{i+j \bmod 2}$ (grading)^a
- ii) $x_i \circ x_j = -(-1)^{ij} x_j \circ x_i$ (supersymmetrization)
- iii) $x_i \circ (x_j \circ x_k)(-1)^{ik} + x_j \circ (x_k \circ x_i)(-1)^{ji} + x_k \circ (x_i \circ x_j)(-1)^{kj} = 0$
(generalised Jacobi identity)

^aFor $x_0 \in L_0$ and $x_1 \in L_1$, this means that $x_0 \circ x_0 \in L_0$, $x_1 \circ x_1 \in L_0$ and $x_0 \circ x_1 \in L_1$.

The second requirement generalises the definition of a Lie algebra in Sec. 1.7 to allow for anti-commutators, $x \circ y = \{x, y\} \equiv xy + yx$, as the binary operation for elements in L_1 .

We can now start, following HLS, with the Poincaré Lie algebra ($L_0 = P$) and add a new vector space L_1 spanned by some generators Q_a . It can be shown that the superalgebra requirements are fulfilled if there are four generators, $a = 1, 2, 3, 4$, that together form a four-component Majorana spinor, also called the **supercharges**. The algebra is then

$$[Q_a, P_\mu] = 0 \quad (2.30)$$

$$[Q_a, M_{\mu\nu}] = (\sigma_{\mu\nu} Q)_a \quad (2.31)$$

$$\{Q_a, \bar{Q}_b\} = 2\mathcal{P}_{ab}, \quad (2.32)$$

where $\sigma_{\mu\nu}$ is given in terms of the γ -matrices, $\sigma_{\mu\nu} = \frac{i}{4}[\gamma_\mu, \gamma_\nu]$, and as usual $\mathcal{P} = P_\mu \gamma^\mu$ and $\bar{Q}_a = (Q^\dagger \gamma_0)_a$.¹⁰ Together with the commutators in (2.4), (2.6) and (2.7) this is called the **super-Poincaré algebra**.

Because of (2.31) this new algebra is a non-trivial extension of the Poincaré algebra that avoids the no-go theorem. However, in the Q_a we have introduced new operators that (disappointingly) do not correspond to the generators of the gauge groups (which should be related by commutators, not anti-commutators). The gauge group generators *can* appear in the algebra if we extend the algebra with a set of $N > 1$ new spinors Q_a^α , where $\alpha = 1, \dots, N$. This gives rise to so-called $N > 1$ supersymmetries.¹¹ However, this seems impossible to realise in nature at energies accessible in experiments due to an extensive number of extra particles.¹² This extension, including the potential for $N > 1$, can be proven, under some reasonable assumptions, to be the *largest possible* extension of the symmetries of special relativity.

We can also write the super-Poincaré algebra in terms of the Weyl spinors introduced in Sec. 2.4. With

$$Q_a = \begin{pmatrix} Q_A \\ \bar{Q}^{\dot{A}} \end{pmatrix}, \quad (2.33)$$

¹⁰Alternatively, (2.32) can be written as $\{Q_a, Q_b\} = -2(\gamma^\mu C)_{ab} P_\mu$.

¹¹Given a gauge group algebra $[B_i, B_j] = iC_{ij}^k B_k$, we can extend the $N = 1$ superalgebra by the non-trivial commutator $[Q_a^\alpha, B_i] = iS_i^{\alpha\beta} Q_a^\beta$, where S_i are matrix representations of the gauge symmetry group.

¹²Note that $N > 8$ would also include elementary particles with spin greater than 2, which seems to be in contradiction with quantum field theory.

for the Majorana spinor charges, we have instead

$$[Q_A, P_\mu] = [\bar{Q}_{\dot{A}}, P_\mu] = 0, \quad (2.34)$$

$$[Q_A, M^{\mu\nu}] = \sigma_A^{\mu\nu B} Q_B, \quad (2.35)$$

$$\{Q_A, Q_B\} = \{\bar{Q}_{\dot{A}}, \bar{Q}_{\dot{B}}\} = 0, \quad (2.36)$$

$$\{Q_A, \bar{Q}_{\dot{B}}\} = 2\sigma_{A\dot{B}}^\mu P_\mu, \quad (2.37)$$

where now the $\sigma^{\mu\nu}$ are given in terms of the Pauli matrices $\sigma^{\mu\nu} = \frac{i}{4}(\sigma^\mu \bar{\sigma}^\nu - \sigma^\nu \bar{\sigma}^\mu)$.

2.6 The Casimir operators of the super-Poincaré algebra

It is easy to see that P^2 is also a Casimir operator of the superalgebra. From Eq. (2.30) P_μ commutes with the Q s, so in turn P^2 must commute.¹³ However, W^2 is not a Casimir because of the following result:¹⁴

$$[W^2, Q_a] = W_\mu (\not{P} \gamma^\mu \gamma^5 Q)_a + \frac{3}{4} P^2 Q_a.$$

We want to find an extension of W that commutes with the Q s while retaining the commutators we already have. The construction

$$C_{\mu\nu} \equiv B_\mu P_\nu - B_\nu P_\mu,$$

where

$$B_\mu \equiv W_\mu + \frac{1}{4} X_\mu,$$

and with

$$X_\mu \equiv \frac{1}{2} \bar{Q} \gamma_\mu \gamma^5 Q,$$

has the required relation:

$$[C_{\mu\nu}, Q_a] = 0.$$

Note that we also have

$$[X_\mu, P_\nu] = 0. \quad (2.38)$$

We can show that C^2 then indeed commutes with all the generators in the algebra:

$$\begin{aligned} [C^2, Q_a] &= 0, & (\text{trivial}) \\ [C^2, P_\mu] &= 0, & (\text{proof by excessive algebra}) \\ [C^2, M_{\mu\nu}] &= 0. & (\text{because } C^2 \text{ is a Lorentz scalar}) \end{aligned}$$

Thus C^2 is a Casimir operator for the superalgebra.

¹³Although the fact that Eq. (2.30) holds crucially depends on Q_a being four-dimensional. P_μ and Q_a would not commute if there had been five Q s.

¹⁴Which, by the way, is really hard work!

As was the case for the original Poincaré group, states are labeled by m , where m^2 is the eigenvalue of P^2 . To find the possible eigenvalues of C^2 , let us again assume without loss of generality that we are in the rest frame (RF).¹⁵ For C^2 we have to do a bit of calculation:

$$\begin{aligned} C^2 &= 2B_\mu P_\nu B^\mu P^\nu - 2B_\mu P_\nu B^\nu P^\mu \\ &\stackrel{RF}{=} 2m^2 B_\mu B^\mu - 2m^2 B_0^2 \\ &= 2m^2 B_k B^k, \end{aligned}$$

where we used that $[B_\mu, P_\nu] = 0$, which we get from (2.12) and (2.38). From the definition of B_μ :

$$B_k = W_k + \frac{1}{4}X_k = mS_k + \frac{1}{8}\bar{Q}\gamma_k\gamma^5 Q \equiv mJ_k. \quad (2.39)$$

The operator we just defined, $J_k \equiv \frac{1}{m}B_k$, is an extension of the ordinary spin operator S_k . Just like the spin operator it can be shown to fulfil the angular momentum $su(2)$ algebra:

$$[J_i, J_j] = i\epsilon_{ijk}J_k,$$

and, additionally, it commutes with the Q s¹⁶

$$[J_k, Q_a] = 0. \quad (2.40)$$

This gives us, still in the rest frame,

$$C^2 = 2m^4 J_k J^k = -2m^4 J^2,$$

so that the eigenvalue equation is:

$$C^2|m, j\rangle = -m^4 j(j+1)|m, j\rangle,$$

for $j = 0, \frac{1}{2}, 1, \frac{3}{2}, \dots$, where the proof follows that for the angular momentum eigenvalues in quantum mechanics, because this relies only on the properties of the $su(2)$ algebra. In addition, for each irrep with a value of j , the angular momentum algebra allows us to show there are $2j+1$ distinct *states* with labels $j_3 = -j, -j+1, \dots, j-1, j$, so that we may further write

$$C^2|m, j, j_3\rangle = -m^4 j(j+1)|m, j, j_3\rangle,$$

labelling also the states of the irrep.¹⁷ So, in summary, the irreducible representations of the superalgebra can be labeled by (m, j) , and any given set of m and j will give us $2j+1$ states with different j_3 .¹⁸

2.7 Short interlude on γ -matrices

When we deal with four-component spinors we have a use for γ -matrices. These are defined as objects that fulfil a type of **Clifford algebra** given by¹⁹

$$\{\gamma_\mu, \gamma_\nu\} = 2g_{\mu\nu}. \quad (2.41)$$

¹⁵We can again carry out a similar argument in a different frame for massless particles.

¹⁶Again the proof is algebraically extensive, and the interested reader is suggested to pursue [3].

¹⁷Note that, as we shall see, unlike for spin this does not exhaust the number of states for the irrep.

¹⁸Make sure you remember that that j is *not* the spin, but a generalisation of spin.

¹⁹Be aware that the expression on the right hand side should be read as consisting of a rank-2 tensor *with each element* being a 4×4 identity matrix.

There exists different representations of this algebra, just as for the Lie algebras. In these notes we will use what is called the **Weyl-representation** where the γ -matrices are given in terms of the Pauli matrices as

$$\gamma^\mu = \begin{pmatrix} 0 & \sigma^\mu \\ \bar{\sigma}^\mu & 0 \end{pmatrix}. \quad (2.42)$$

The γ -matrices can be used to form a ‘fifth’ γ -matrix

$$\gamma^5 \equiv i\gamma^0\gamma^1\gamma^2\gamma^3 = \begin{pmatrix} -\sigma^0 & 0 \\ 0 & \sigma^0 \end{pmatrix}.$$

Using these expressions one can find the relationships between the four-component and two-component (Weyl spinor) notation for the supercharges.

2.8 Irreducible representations of the super-Poincaré group

What sort of particles transform under the super-Poincaré group? Or, in other words, what are the properties of the irreducible representations of the group? In the following we will construct all the states for a given representation labeled by the set (m, j) . To do this it is very useful to write the generators Q in terms of two-component Weyl spinors instead of four-component Dirac spinors, making explicit use of their Majorana nature, as we did in Section 2.4. We note that from Eq. (2.40) above

$$[J_k, Q_A] = [J_k, \bar{Q}_{\dot{B}}] = 0.$$

We begin by claiming that for any state with a given value of j_3 there must then exist a state $|\Omega\rangle$ – possibly the same state – that has the same value of j_3 and for which

$$Q_A|\Omega\rangle = 0. \quad (2.43)$$

This state is called the **Clifford vacuum**.²⁰

To show this, start with $|\beta\rangle$, a state with j_3 . Then the construction

$$|\Omega\rangle = Q_1 Q_2 |\beta\rangle,$$

has these properties. Using (2.36) we first we show that (2.43) holds:

$$Q_1 Q_1 Q_2 |\beta\rangle = -Q_1 Q_1 Q_2 |\beta\rangle = 0,$$

and

$$Q_2 Q_1 Q_2 |\beta\rangle = -Q_1 Q_2 Q_2 |\beta\rangle = Q_1 Q_2 Q_2 |\beta\rangle = 0.$$

For this state we also have:

$$J_3 |\Omega\rangle = J_3 Q_1 Q_2 |\beta\rangle = Q_1 Q_2 J_3 |\beta\rangle = j_3 |\Omega\rangle,$$

in other words, $|\Omega\rangle$ has the same value for j_3 as the $|\beta\rangle$ it was constructed from and the Clifford vacuum exists. This proof demonstrates a general feature of the supercharges, if

²⁰It is called the Clifford vacuum because the operators satisfy a Clifford algebra $\{Q_A, \bar{Q}_{\dot{B}}\} = 2m\sigma_{A\dot{B}}^0$. Do not confuse this with a vacuum state, it is only a name.

one supercharge with a particular index repeats in a term, then the term is zero by the anticommutation property of the supercharges.

We can now use the explicit expression for J_k in terms of the two-component supercharges

$$J_k = S_k - \frac{1}{4m} \bar{Q}_{\dot{B}} \bar{\sigma}_k^{\dot{B}A} Q_A, \quad (2.44)$$

in order to find the spin for this state. First we can see that

$$S_k |\Omega\rangle = J_k |\Omega\rangle = j_k |\Omega\rangle,$$

meaning that $s_3 = j_3$ is the eigenvalue of S_3 for the Clifford vacuum $|\Omega\rangle$. Further, since

$$S^2 |\Omega\rangle = J^2 |\Omega\rangle = j(j+1) |\Omega\rangle,$$

the eigenvalue of S^2 is $s(s+1) = j(j+1)$ for the Clifford vacuum.

We can construct three more states from the Clifford vacuum:²¹

$$\bar{Q}^{\dot{1}} |\Omega\rangle, \quad \bar{Q}^{\dot{2}} |\Omega\rangle, \quad \bar{Q}^{\dot{1}} \bar{Q}^{\dot{2}} |\Omega\rangle.$$

This means that there are four possible states that can be constructed out of any state with the quantum numbers m, j, j_3 . Taking a look at:

$$J_k \bar{Q}^{\dot{A}} |\Omega\rangle = \bar{Q}^{\dot{A}} J_k |\Omega\rangle = j_k \bar{Q}^{\dot{A}} |\Omega\rangle,$$

this means that all these states have the same j_3 (and j) quantum numbers.²² We can now find their eigenvalues for S_3 . From the superalgebra (2.35) we have:

$$[M^{ij}, \bar{Q}^{\dot{A}}] = -(\sigma^{ij})^{\dot{A}}_{\dot{B}} \bar{Q}^{\dot{B}},$$

so that:

$$S_3 \bar{Q}^{\dot{A}} |\Omega\rangle = \bar{Q}^{\dot{A}} S_3 |\Omega\rangle - \frac{1}{2} (\bar{\sigma}_3 \sigma^0)^{\dot{A}}_{\dot{B}} \bar{Q}^{\dot{B}} |\Omega\rangle = \left(j_3 \mp \frac{1}{2} \right) \bar{Q}^{\dot{A}} |\Omega\rangle,$$

where $-$ is for $\dot{A} = \dot{1}$ and $+$ is for $\dot{A} = \dot{2}$. We can similarly show that

$$S_3 \bar{Q}^{\dot{1}} \bar{Q}^{\dot{2}} |\Omega\rangle = j_3 \bar{Q}^{\dot{1}} \bar{Q}^{\dot{2}} |\Omega\rangle.$$

This means that for an irrep with labels m and j , there are $2j+1$ different values of j_3 , each giving two states with $s_3 = j_3$, and two with $s_3 = j_3 \pm \frac{1}{2}$, meaning two bosonic and two fermionic states with the same mass m , and in total $4(2j+1)$ states per irrep.

We should be careful to note here that we have only found the spin-components s_3 of these states, not their spins s . For the state $\bar{Q}^{\dot{1}} \bar{Q}^{\dot{2}} |\Omega\rangle$, s is the same as for the Clifford vacuum, *i.e.* $s = j$. This is because in the application of S_k from (2.44) to the state the terms with supercharges will all be zero since at least one of the $\bar{Q}^{\dot{A}}$ will repeat in each term, thus the eigenvalues of S^2 are the same as the eigenvalues of J^2 . For the other states we may need to combine states into definite spin states using Clebsch-Gordan coefficients.

²¹All other possible combinations of Q s and $|\Omega\rangle$ give either one of the other four states, or the zero state which is trivial and of no interest.

²²The same can easily be shown for $\bar{Q}^{\dot{1}} \bar{Q}^{\dot{2}} |\Omega\rangle$.

The above explains the much repeated statement that any supersymmetry theory has an equal number of bosons and fermions, which, incidentally, is not true. What is true, is that there must be an equal number of bosonic and fermionic states.

Theorem: For any representation of the superalgebra where P_μ is a one-to-one operator there is an equal number of boson and fermion states.

To show this, divide the representation into two sets of states, one with bosons and one with fermions. Let $\{Q_A, \bar{Q}_{\dot{B}}\}$ act on the members of the set of bosons. $\bar{Q}_{\dot{B}}$ transforms bosons to fermions and Q_A does the reverse mapping. If P_μ is one-to-one, then so is $\{Q_A, \bar{Q}_{\dot{B}}\} = 2\sigma^\mu_{A\dot{B}} P_\mu$. Thus there must be an equal number in both sets.

2.8.1 Examples of irreducible representations

Finally, let us briefly look at two examples of irreducible representations for a fixed non-zero m .

$j = 0$

For $j = 0$, we must have $j_3 = 0$ and as a result the Clifford vacuum $|\Omega\rangle$ has $s = 0$, $s_3 = 0$, and is a bosonic state. We can then create two states $\bar{Q}^{\dot{A}}|\Omega\rangle$ with $s_3 = \pm\frac{1}{2}$ and $s = \frac{1}{2}$, and one state $\bar{Q}^{\dot{1}}\bar{Q}^{\dot{2}}|\Omega\rangle$ with $s_3 = 0$ and $s = 0$. Note that we should really check the total spin s of each of the fermion states, which would involve some algebra. In total there are two scalar states and two spin- $\frac{1}{2}$ fermion states. We will later represent this set of states by the so-called **scalar superfield**.

We should use be a little careful about using the term particle about these states since what we have found are in fact Weyl spinor states. As we saw in Sec. 2.4 a Dirac fermion can then only be described by a $j = 0$ representation together with a different complex conjugate representation, thus consisting of four states. The complex conjugate representation of the first representation together with the self-representation of the second then form the anti-particle of the fermion, and provide an additional two scalars. So the particle count from the two irreducible representations is a fermion–anti-fermion pair, and four scalars. Note that all of the resulting particles have the same mass m .

For a Majorana fermion the situation is simpler, since we only need one self-representation and its complex conjugate representation.

$j = \frac{1}{2}$

For $j = \frac{1}{2}$ we have two Clifford vacua $|\Omega\rangle$ with $j_3 = \pm\frac{1}{2}$, and with $s = \frac{1}{2}$ and $s_3 = \pm\frac{1}{2}$, thus they are fermionic states. For the moment we label them as $|\Omega; \frac{1}{2}\rangle$ and $|\Omega; -\frac{1}{2}\rangle$. From each of these we can construct two further fermion states $\bar{Q}^{\dot{1}}\bar{Q}^{\dot{2}}|\Omega; \pm\frac{1}{2}\rangle$ where we know $s = \frac{1}{2}$ and $s_3 = \mp\frac{1}{2}$. Together these four states can form two fermions with $s = \frac{1}{2}$ and $s_3 = \pm\frac{1}{2}$.

In addition to this we have the two states $\bar{Q}^{\dot{1}}|\Omega; \frac{1}{2}\rangle$ and $\bar{Q}^{\dot{2}}|\Omega; -\frac{1}{2}\rangle$ with $s_3 = 0$, the state $\bar{Q}^{\dot{2}}|\Omega; \frac{1}{2}\rangle$ with $s_3 = 1$, and the state $\bar{Q}^{\dot{1}}|\Omega; -\frac{1}{2}\rangle$ which has $s_3 = -1$. From these we can create three states with $s = 1$, and $s_3 = 1, 0, -1$, and one state with $s = 0$ and $s_3 = 0$, representing

one massive vector particle and one scalar. We will later refer to this set of states as the **vector superfield**.

2.9 Exercises

Exercise 2.1

Show the commutation properties of the Lorentz group generators J_i and K_i .

Exercise 2.2

Show the commutation properties of the Poincaré group generators P_μ and $M_{\mu\nu}$.

Exercise 2.3

Show that

$$[P_\mu, P^2] = 0, \quad (2.45)$$

$$[M_{\mu\nu}, P^2] = 0. \quad (2.46)$$

Exercise 2.4

Show that $[P_\mu, W_\nu] = 0$.

Exercise 2.5

Show that

$$[P_\mu, W^2] = 0 \quad (2.50)$$

$$[M_{\mu\nu}, W^2] = 0. \quad (2.51)$$

Hint: You can use that²³

$$W^2 = -\frac{1}{2}M_{\mu\nu}M^{\mu\nu}P^2 + M^{\rho\sigma}M_{\nu\sigma}P_\rho P^\nu,$$

Exercise 2.6

Starting from the four-component form of the super-Poincaré algebra, derive the two-component (Weyl spinor) form.

Exercise 2.7

Show that $[X_\mu, P_\nu] = 0$.

Exercise 2.8

Show that $SO^+(1, 3)$ and $SL(2, \mathbb{C})$ are indeed homomorphic, *i.e.* that the mapping defined by (2.16) or (2.17) has the property that $\Lambda(M_1 M_2) = \Lambda(M_1)\Lambda(M_2)$ or $M(\Lambda_1 \Lambda_2) = M(\Lambda_1)M(\Lambda_2)$.

²³This is non-trivial to demonstrate, see Chapter 1.2 of [3].

Exercise 2.9

Show that the generalization of the spin operator, $J_k \equiv S_k + \frac{1}{8m} \bar{Q} \gamma_\mu \gamma^5 Q$, fulfils the algebra

$$[J_i, J_j] = i\epsilon_{ijk} J_k.$$

Exercise 2.10

What are the states for $j = 1$?

Chapter 3

Superspace

In this chapter we will introduce a very handy notation system for considering supersymmetry transformations effected by the Q elements of the superalgebra, or, more correctly, the elements of the super-Poincaré group and their representations. This notation uses a coordinate system called superspace, and allows us to define so-called superfields as a replacement of ordinary field theory fields. This mirrors the Lorentz invariance built into relativistic field theory by using four-vectors. In order to do this we first need to know a little more about the properties of Grassman numbers.

3.1 Superspace calculus

Grassmann numbers θ are numbers that anti-commute with each other but not with ordinary numbers. We will here use four such numbers and in addition we want to place them in Weyl spinors, indexed by A and $\dot{A}$.¹

$$\{\theta^A, \theta^B\} = \{\theta^A, \bar{\theta}^{\dot{B}}\} = \{\bar{\theta}^{\dot{A}}, \theta^B\} = \{\bar{\theta}^{\dot{A}}, \bar{\theta}^{\dot{B}}\} = 0.$$

From this we get the relationships:²

$$\theta_A^2 = \theta_A \theta_A = -\theta_A \theta_A = 0, \quad (3.1)$$

$$\theta^2 \equiv \theta \theta \equiv \theta^A \theta_A = -2\theta_1 \theta_2, \quad (3.2)$$

$$\bar{\theta}^2 \equiv \bar{\theta} \bar{\theta} \equiv \bar{\theta}_{\dot{A}} \bar{\theta}^{\dot{A}} = 2\bar{\theta}^{\dot{1}} \bar{\theta}^{\dot{2}}. \quad (3.3)$$

Notice that if we have a function f of a Grassman number, say θ_A , then the all-order expansion of that function in terms of θ_A , is

$$f(\theta_A) = a_0 + a_1 \theta_A, \quad (3.4)$$

as there are simply no more terms because of (3.1).

¹We can already see how this can be useful: if we consistently use $\theta^A Q_A$ and $\bar{\theta}_{\dot{A}} \bar{Q}^{\dot{A}}$ instead of only Q_A and $\bar{Q}^{\dot{A}}$ in Eqs. (2.36)–(2.35) we can actually rewrite the superalgebra as an ordinary Lie algebra, but with Grassman elements, because of these commutation properties.

²There is no summation implied in the first line. These are of course the same relations we already used for the Weyl spinors.

We now need to define differentiation and integration on these numbers in order to create a calculus for them.³

Definition: We define differentiation by:

$$\partial_A \theta^B \equiv \frac{\partial}{\partial \theta^A} \theta^B \equiv \delta_A^B,$$

with a product rule

$$\begin{aligned} \partial_A (\theta^{B_1} \theta^{B_2} \theta^{B_3} \dots \theta^{B_n}) &\equiv (\partial_A \theta^{B_1}) \theta^{B_2} \theta^{B_3} \dots \theta^{B_n} \\ &\quad - \theta^{B_1} (\partial_A \theta^{B_2}) \theta^{B_3} \dots \theta^{B_n} \\ &\quad + \dots + (-1)^{n-1} \theta^{B_1} \theta^{B_2} \dots (\partial_A \theta^{B_n}). \end{aligned} \quad (3.5)$$

This implies that the differential operator ∂_A is itself a Grassmann number and anti-commutes. We can show that the indices of these operators can be raised and lowered with the ϵ in Eqs. (2.18) and (2.19).

Definition: The integral of a function f of a Grassmann number θ_A is defined as a functional $I[f]$

$$I[f] = \int d\theta_A f(\theta_A),$$

which evaluates to a c -number. The evaluation is defined by the relations

$$\int d\theta_A \equiv 0, \quad \int d\theta_A \theta_A \equiv 1,$$

and similarly for $\bar{\theta}_A$, and we demand linearity:

$$\int d\theta_A [af(\theta_A) + bg(\theta_A)] \equiv a \int d\theta_A f(\theta_A) + b \int d\theta_A g(\theta_A).$$

In these definitions there is no implied summation over the index A .

This has one surprising property. If we take the integral of (3.4) we get:

$$\int d\theta_A f(\theta_A) = a_1 = \partial^A f(\theta_A),$$

meaning that differentiation and integration has the same effect on Grassman numbers.

To integrate over multiple Grassman numbers we define volume elements as

³These definitions have no infinitesimal interpretations, not is there any notion of definite integrals.

Definition:

$$\begin{aligned} d^2\theta &\equiv -\frac{1}{4}d\theta^A d\theta_A, \\ d^2\bar{\theta} &\equiv -\frac{1}{4}d\bar{\theta}_{\dot{A}} d\bar{\theta}^{\dot{A}}, \\ d^4\theta &\equiv d^2\theta d^2\bar{\theta}. \end{aligned}$$

and we demand that the integral operators anti-commute, just as the differential operators

$$\left\{ \int d\theta_A, \int d\theta_B \right\} = \left\{ \int d\theta_A, \theta_B \right\} = 0.$$

This specific volume element definition is made to normalise the following integrals

$$\begin{aligned} \int d^2\theta \theta\theta &= 1, \\ \int d^2\bar{\theta} \bar{\theta}\bar{\theta} &= 1, \\ \int d^4\theta (\theta\theta)(\bar{\theta}\bar{\theta}) &= 1. \end{aligned}$$

Delta functions of Grassmann variables are given by:

$$\begin{aligned} \delta(\theta_A) &= \theta_A, \\ \delta^2(\theta_A) &= \theta\theta, \\ \delta^2(\bar{\theta}^{\dot{A}}) &= \bar{\theta}\bar{\theta}, \end{aligned}$$

and we can easily show that these functions satisfy, just as the usual definition of delta functions,

$$\int d\theta_A f(\theta_A) \delta(\theta_A) = f(0).$$

3.2 Superspace definition

Superspace⁴ is a coordinate system where supersymmetry transformations are *manifest*, in other words, the action of elements in the super-Poincaré group (SP) are treated like Lorentz-transformations are in Minkowski space.

Definition: Superspace is an eight-dimension manifold that can be constructed from the **coset space** of the super-Poincaré group (SP) and the Lorentz group (L), SP/L , by giving coordinates $z^\pi = (x^\mu, \theta^A, \bar{\theta}_{\dot{A}})$, $\pi = \mu, A, \dot{A}$, where x^μ are the ordinary Minkowski coordinates, and where θ_A and $\bar{\theta}^{\dot{A}}$ are four Grassman (anti-commuting) numbers, being the parameters of the Q -operators in the algebra.

⁴First introduced by Salam & Strathdee [6].

To understand what all of this means we begin by writing a general element of SP , $g \in SP$, using the exponential map defined in Section 1.7:

$$g = \exp[ix^\mu P_\mu + i\theta^A Q_A + i\bar{\theta}_{\dot{A}} \bar{Q}^{\dot{A}} + \frac{i}{2}\omega_{\rho\nu} M^{\rho\nu}],$$

where x^μ , θ^A , $\bar{\theta}_{\dot{A}}$ and $\omega_{\rho\nu}$ constitute the parametrisation of the group, and P_μ , Q_A , $\bar{Q}^{\dot{A}}$ and $M_{\rho\nu}$ are the generators. We can now parametrise SP/L simply by setting $\omega_{\mu\nu} = 0$.⁵ The remaining parameters of SP/L are then the coordinates that span superspace.

As physicists we also want to know the dimensions of our new parameters θ . To do this we first look at Eq. (2.37):

$$\{Q_A, \bar{Q}_{\dot{B}}\} = 2\sigma^\mu_{A\dot{B}} P_\mu,$$

where we know that P_μ has mass dimension $[P_\mu] = M$. This means that $[Q^2] = M$ and $[Q] = M^{1/2}$. In the exponential, all terms must have mass dimension zero to make sense. This means that $[\theta Q] = 0$, and therefore $[\theta] = M^{-1/2}$.

We now want to find the effect of supersymmetry transformations (transformations by the super-Poincaré group) on the superspace coordinates, and we begin by noting that any SP transformation can effectively be written in the following way without the boosts and rotations of the Lorentz group:

$$g_0(a, \alpha) = \exp[ia^\mu P_\mu + i\alpha^A Q_A + i\bar{\alpha}_{\dot{A}} \bar{Q}^{\dot{A}}],$$

because one can show that⁶

$$\exp\left[\frac{i}{2}\omega_{\rho\nu} M^{\rho\nu}\right] g_0(a, \alpha) = g_0(\Lambda a, S(\Lambda)\alpha) \exp\left[\frac{i}{2}\omega_{\rho\nu} M^{\rho\nu}\right], \quad (3.6)$$

i.e. all that a Lorentz boost does is to transform spacetime coordinates by Λ and Weyl spinors by $S(\Lambda)$, which is a spinor representation of Λ . Thus, in more colloquial terms, for the supersymmetry transformation it does not matter which reference frame we are working in, we know how the transformation changes between the frames, given by Λa and $S(\Lambda)\alpha$.

We can now find the transformation of superspace coordinates under a supersymmetry transformation, just as we have seen the Minkowski coordinates transform under Lorentz transformations. It might be tempting to look at the effects of an element $g_0(a, \alpha)$ on a function on superspace coordinates, $F(z^\pi) = F(x^\mu, \theta^A, \bar{\theta}^{\dot{A}})$ just as we did for the translation group. However, the powers of the sum of generators in the infinite series will become very unwieldy and messy. Instead, we pull out the coordinate dependence from the function $F(z^\pi) = e^{iz^\pi K_\pi} f(0)$, where $K_\pi = (P_\mu, Q_A, \bar{Q}^{\dot{A}})$, and we look at the transformation $z^\pi \rightarrow z'^\pi$ given by

$$g_0 e^{iz^\pi K_\pi} = e^{iz'^\pi K_\pi}.$$

In moving the problem to the exponential, we can use the power of some results on the exponentials of non-commuting quantities, in particular the Campbell-Baker-Hausdorff expansion.

⁵ SP/L is *not* a coset group as defined previously, because L is not a normal subgroup of SP , but its parametrisation still forms a vector space (the coset space) which we call superspace.

⁶Fortunately we are not going to do this because it is messy, but it can be done using the algebra of the group and the series expansion of the exponential function. Note, however, that the proof rests on the P s and Q s forming a closed set, which we saw in the algebra Eqs. (2.36)–(2.35).

We have⁷

$$\begin{aligned} g_0 e^{iz^\pi K_\pi} &= \exp(ia^\nu P_\nu + i\alpha^B Q_B + i\bar{\alpha}_{\dot{B}} \bar{Q}^{\dot{B}}) \exp(iz^\pi K_\pi) \\ &= \exp(ia^\nu P_\nu + i\alpha^B Q_B + i\bar{\alpha}_{\dot{B}} \bar{Q}^{\dot{B}} + iz^\pi K_\pi \\ &\quad - \frac{1}{2}[ia^\nu P_\nu + i\alpha^B Q_B + i\bar{\alpha}_{\dot{B}} \bar{Q}^{\dot{B}}, iz^\pi K_\pi] +) \end{aligned}$$

Now we take a closer look at the commutator:⁸

$$\begin{aligned} [,] &= -[\alpha^B Q_B, \bar{\theta}_{\dot{A}} \bar{Q}^{\dot{A}}] - [\bar{\alpha}_{\dot{B}} \bar{Q}^{\dot{B}}, \theta^A Q_A] \\ &= \alpha^B \bar{\theta}_{\dot{A}} \epsilon^{\dot{A}\dot{C}} \{Q_B, \bar{Q}_{\dot{C}}\} + \bar{\alpha}_{\dot{B}} \theta^A \epsilon^{\dot{B}\dot{C}} \{\bar{Q}_{\dot{C}}, Q_A\} \\ &= 2\alpha^B \bar{\theta}_{\dot{A}} \epsilon^{\dot{A}\dot{C}} \sigma_{B\dot{C}}^\mu P_\mu + 2\bar{\alpha}_{\dot{B}} \theta^A \epsilon^{\dot{B}\dot{C}} \sigma_{A\dot{C}}^\mu P_\mu \\ &= 2(\alpha^B \bar{\theta}^{\dot{C}} \sigma_{B\dot{C}}^\mu + \bar{\alpha}^{\dot{C}} \theta^A \sigma_{A\dot{C}}^\mu) P_\mu \end{aligned}$$

We can relabel $B = A$ and $\dot{C} = \dot{A}$ which leads to

$$-\frac{1}{2}[,] = -(\alpha^A \sigma_{A\dot{A}}^\mu \bar{\theta}^{\dot{A}} - \theta^A \sigma_{A\dot{A}}^\mu \bar{\alpha}^{\dot{A}}) P_\mu.$$

The commutator is proportional with P_μ , and will therefore commute with all operators, in particular the higher terms in the Campbell-Baker-Hausdorff expansion, meaning that the series reduces to

$$g_0 e^{iz^\pi K_\pi} = \exp[i(x^\mu + a^\mu + i\alpha^A \sigma_{A\dot{A}}^\mu \bar{\theta}^{\dot{A}} - i\theta^A \sigma_{A\dot{A}}^\mu \bar{\alpha}^{\dot{A}}) P_\mu + i(\theta^A + \alpha^A) Q_A + i(\bar{\theta}_{\dot{A}} + \bar{\alpha}_{\dot{A}}) \bar{Q}^{\dot{A}}].$$

So superspace coordinates transform under supersymmetry transformations as:

$$(x^\mu, \theta^A, \bar{\theta}_{\dot{A}}) \rightarrow (x^\mu + a^\mu - i\alpha^A \sigma_{A\dot{A}}^\mu \bar{\theta}^{\dot{A}} + i\theta^A \sigma_{A\dot{A}}^\mu \bar{\alpha}^{\dot{A}}, \theta^A + \alpha^A, \bar{\theta}_{\dot{A}} + \bar{\alpha}_{\dot{A}}). \quad (3.7)$$

or given more explicitly as a composition function

$$f(x^\mu, \theta^A, \bar{\theta}_{\dot{A}}, a^\mu, \alpha^A, \bar{\alpha}_{\dot{A}}) = (x^\mu + a^\mu - i\alpha^A \sigma_{A\dot{A}}^\mu \bar{\theta}^{\dot{A}} + i\theta^A \sigma_{A\dot{A}}^\mu \bar{\alpha}^{\dot{A}}, \theta^A + \alpha^A, \bar{\theta}_{\dot{A}} + \bar{\alpha}_{\dot{A}}). \quad (3.8)$$

As a crucial by-product we can now write down a differential representation for the supersymmetry generators by applying the standard expression for the generators X_i of a Lie algebra, given the composition functions f_π :

$$X_j = \frac{\partial f_\pi}{\partial a_j} \frac{\partial}{\partial z_\pi},$$

which gives us:⁹

$$P_\mu = -i\partial_\mu, \quad (3.9)$$

$$iQ_A = -i(\sigma^\mu \bar{\theta})_A \partial_\mu + \partial_A, \quad (3.10)$$

$$i\bar{Q}^{\dot{A}} = -i(\bar{\sigma}^\mu \theta)^{\dot{A}} \partial_\mu + \partial^{\dot{A}}. \quad (3.11)$$

⁷Here we use Campbell-Baker-Hausdorff expansion $e^A e^B = e^{A+B-\frac{1}{2}[A,B]+\dots}$ where the terms that follow all contain commutators of the first commutator $[A,B]$ and the operators A and B .

⁸Using that P_μ commutes with all elements in the algebra, as well as $[\theta^A Q_A, \xi^B Q_B] = -\theta^A \xi^B \{Q_A, Q_B\} = 0$, and the same for $\bar{Q}^{\dot{B}}$.

⁹We define the generators X_i as iP_μ , iQ_A and iQ_B respectively. Again, this is in order to end up with hermitian instead of anti-hermitian generators.

3.3 Covariant derivatives

Similar to the properties of covariant derivatives for gauge transformations in gauge theories, it would be nice to have a derivative that is invariant under supersymmetry transformations, *i.e.* commutes with supersymmetry operators. Obviously $P_\mu = -i\partial_\mu$ does this, but more general covariant derivatives can be made.

Definition: The following **covariant derivatives** commute with supersymmetry transformations:

$$D_A \equiv \partial_A + i(\sigma^\mu \bar{\theta})_A \partial_\mu, \quad (3.12)$$

$$\bar{D}_{\dot{A}} \equiv -\partial_{\dot{A}} - i(\theta \sigma^\mu)_{\dot{A}} \partial_\mu. \quad (3.13)$$

These can be shown to satisfy the following relations that are useful in calculations:

$$\{D_A, D_B\} = \{\bar{D}_{\dot{A}}, \bar{D}_{\dot{B}}\} = 0, \quad (3.14)$$

$$\{D_A, \bar{D}_{\dot{B}}\} = -2\sigma^\mu_{A\dot{B}} P_\mu, \quad (3.15)$$

$$D^3 = \bar{D}^3 = 0, \quad (3.16)$$

$$D^A \bar{D}^2 D_A = \bar{D}_{\dot{A}} D^2 \bar{D}^{\dot{A}}. \quad (3.17)$$

Here, D^3 and $\bar{D}^3$ means the application of at least three of these covariant derivatives.

3.4 Superfields

Using the superspace coordinates we can now define functions of these to use in a field theory. Naturally we should call these objects superfields.

Definition: A **superfield** Φ is an operator valued function on superspace $\Phi(x, \theta, \bar{\theta})$.

Notice how we write, just as for ordinary fields depending on four-vector coordinates, the coordinates sans indices.

We can expand any such superfield $\Phi(x, \theta, \bar{\theta})$ as a power series in θ and $\bar{\theta}$. This gives, in general,

$$\begin{aligned} \Phi(x, \theta, \bar{\theta}) = & f(x) + \theta^A \phi_A(x) + \bar{\theta}_{\dot{A}} \bar{\chi}^{\dot{A}}(x) + \theta\theta m(x) + \bar{\theta}\bar{\theta} n(x) \\ & + \theta\sigma^\mu \bar{\theta} V_\mu(x) + \theta\theta \bar{\theta}_{\dot{A}} \bar{\lambda}^{\dot{A}}(x) + \bar{\theta}\bar{\theta} \theta^A \psi_A(x) + \theta\theta \bar{\theta}\bar{\theta} d(x), \end{aligned} \quad (3.18)$$

where the functions of space-time coordinates x are called the **component fields**. Any superfield without explicit spinor indices, such as the one above, commutes with any other superfield, because all the Grassmann numbers appear in contracted pairs.

The properties of the component fields can be deduced from the requirement that Φ must be an (operator valued) Lorentz scalar or pseudoscalar. This is shown in Table 3.1 along with the corresponding degrees of freedom each field has.

Component field	Type	d.o.f.
$f(x), m(x), n(x)$	Complex (pseudo) scalar	2
$\psi_A(x), \phi_A(x)$	Left-handed Weyl spinor	4
$\bar{\chi}^{\dot{A}}(x), \bar{\lambda}^{\dot{A}}(x)$	Right-handed Weyl spinor	4
$V_\mu(x)$	Lorentz 4-vector (complex)	8
$d(x)$	Complex scalar	2

Table 3.1: Field content of a general superfield.

One can show that under supersymmetry transformations these component fields transform linearly into each other, thus we say that superfields (with the differential form of the supersymmetry generators) are *representations* of the supersymmetry (super-Poincaré) group just as ordinary quantum fields are representations of the Poincaré group, albeit *highly reducible* representations!¹⁰

We can recover the known irreducible representations, see Section 2.8, by imposing some restrictions on the fields. To do this we define the following three types of superfields that we will discuss below:

$$\bar{D}_{\dot{A}}\Phi(x, \theta, \bar{\theta}) = 0 \quad (\text{left-handed scalar superfield}) \quad (3.19)$$

$$D_A\Phi^\dagger(x, \theta, \bar{\theta}) = 0 \quad (\text{right-handed scalar superfield}) \quad (3.20)$$

$$\Phi^\dagger(x, \theta, \bar{\theta}) = \Phi(x, \theta, \bar{\theta}) \quad (\text{vector superfield}) \quad (3.21)$$

Note that it is $\Phi^\dagger$ which is the right handed superfield in Eq. (3.20), not Φ . However, we can show that the hermitian conjugate of a left-handed scalar superfield fulfils the condition for a right-handed scalar superfield. Supersymmetry transformations can also be shown to transform left-handed superfields into left-handed superfields, right-handed superfields into right-handed superfields and vector superfields into vector superfields. This means that they can be considered separate representations.

Products of the same type of superfield is a superfield of the same type since for left-handed scalar superfields Φ_i and Φ_j ,

$$\bar{D}_{\dot{A}}(\Phi_i\Phi_j) = (\bar{D}_{\dot{A}}\Phi_i)\Phi_j + \Phi_i(\bar{D}_{\dot{A}}\Phi_j) = 0,$$

and similarly for a right-handed scalar superfields, and for vector superfields Φ_i and Φ_j ,

$$(\Phi_i\Phi_j)^\dagger = \Phi_j^\dagger\Phi_i^\dagger = \Phi_j\Phi_i = \Phi_i\Phi_j.$$

The product of a left-handed scalar superfield Φ and its hermitian conjugate $\Phi^\dagger$, $V = \Phi\Phi^\dagger$, is a vector superfield since

$$V^\dagger = (\Phi\Phi^\dagger)^\dagger = \Phi\Phi^\dagger = V.$$

The same is true for sums of superfields of the same type. These properties will be important when creating a superfield version of a Lagrangian.

¹⁰Indeed, they are linear representations since a sum of superfields is a superfield, and the differential supersymmetry operators act linearly.

3.4.1 Scalar superfields

What is the connection of the scalar superfields to the $j = 0$ irreducible representation? To see this we use a cute trick:¹¹ Change to the variable $y^\mu \equiv x^\mu + i\theta\sigma^\mu\bar{\theta}$. Then the covariant derivatives simplify to

$$D_A = \partial_A + 2i\sigma^\mu_{A\dot{A}}\bar{\theta}^{\dot{A}}\frac{\partial}{\partial y^\mu}, \quad (3.22)$$

$$\bar{D}_{\dot{A}} = -\partial_{\dot{A}}. \quad (3.23)$$

This means that a field fulfilling $\bar{D}_{\dot{A}}\Phi = 0$ in the new set of coordinates must be independent of the $\bar{\theta}$ coordinates. Thus we can write this field as:

$$\Phi(y, \theta) = A(y) + \sqrt{2}\theta\psi(y) + \theta\theta F(y), \quad (3.24)$$

and looking at the much restricted component field content we get the result in Table 3.2.

Component field	Type	d.o.f.
$A(x), F(x)$	Complex scalar	2
$\psi_A(x)$	Left-handed Weyl spinor	4

Table 3.2: Fields contained in a left-handed scalar superfield.

Since we wish to interpret the Weyl spinor here as a fermion quantum field with dimension $M^{3/2}$, and given that $[\theta] = M^{-1/2}$, the scalar superfield itself must have mass dimension $[\Phi] = 1$. This means that the scalar field A has the expected mass dimension M^1 of an ordinary scalar quantum field, however, the scalar F has the odd dimension $[F] = 2$.

We can undo the coordinate change in (3.24) by inserting for y and expanding in powers of θ and $\bar{\theta}$, giving

$$\Phi(x, \theta, \bar{\theta}) = A(x) + i(\theta\sigma^\mu\bar{\theta})\partial_\mu A(x) - \frac{1}{4}\theta\theta\bar{\theta}\bar{\theta}\square A(x) + \sqrt{2}\theta\psi(x) - \frac{i}{\sqrt{2}}\theta\theta\partial_\mu\psi(x)\sigma^\mu\bar{\theta} + \theta\theta F(x). \quad (3.25)$$

By using the transformation $y^\mu \equiv x^\mu - i\theta\sigma^\mu\bar{\theta}$ we can show a similar field content for the right-handed scalar superfield. The general form of a right handed scalar superfield is then:

$$\Phi^\dagger(x, \theta, \bar{\theta}) = A^*(x) - i(\theta\sigma^\mu\bar{\theta})\partial_\mu A^*(x) - \frac{1}{4}\theta\theta\bar{\theta}\bar{\theta}\square A^*(x) + \sqrt{2}\bar{\theta}\bar{\psi}(x) + \frac{i}{\sqrt{2}}\bar{\theta}\bar{\theta}\sigma^\mu\partial_\mu\bar{\psi}(x) + \bar{\theta}\bar{\theta}F^*(x). \quad (3.26)$$

We can now compare the above to the $j = 0$ representation with two scalar states and two fermionic states (d.o.f.). After applying the equations of motions (e.o.m.) the **auxiliary** field $F(x)$, with the strange mass dimension, can be completely eliminated as it does not have any derivatives. The e.o.m. also eliminates two of the fermion d.o.f. Thus, after the equations of motion, we are left with the same states as in the $j = 0$ representation.

However, the scalar superfields will not correspond directly to particle states for the known SM particles since, as we discussed in Sec. 2.8.1, a Weyl spinor on its own cannot describe a Dirac fermion. When we construct particle representations we will take one left-handed scalar superfield and one *different* right-handed scalar superfield. These will form a Dirac fermion and two scalars (and their anti-particles) after application of the e.o.m.

¹¹Here cute is used in the widest sense.

3.4.2 Vector superfields

If we take the general superfield in (3.18) and compare the Φ and $\Phi^\dagger$ expressions we can see that the following is the restricted component field structure of a vector superfield:

$$\begin{aligned}\Phi(x, \theta, \bar{\theta}) = & C(x) + \theta\phi(x) + \bar{\theta}\bar{\phi}(x) + \theta\theta M(x) + \bar{\theta}\bar{\theta}M^*(x) \\ & + \theta\sigma^\mu\bar{\theta}V_\mu(x) + \theta\theta\bar{\theta}\bar{\lambda}(x) + \bar{\theta}\bar{\theta}\theta\lambda(x) + \theta\theta\bar{\theta}\bar{\theta}D(x).\end{aligned}$$

The properties of the component fields are summarised in Table 3.3.

Component field	Type	d.o.f.
$C(x), D(x)$	Real scalar field	1
$\phi_A(x), \lambda_A(x)$	Weyl spinor	4
$M(x)$	Complex scalar field	2
$V_\mu(x)$	Real Lorentz 4-vector	4

Table 3.3: Field content of a general vector superfield.

With the large number of component fields, you may now be a little suspicious that this vector superfield does not correspond to the promised degrees of freedom in the $j = \frac{1}{2}$ representation of the superalgebra. However, gauge-freedom now comes to our rescue.

3.5 Supergauge

We first define what we will mean by an abelian (super)gauge transformation of a superfield.¹² Later we will see how it relates to the ordinary gauge transformations of quantum fields. We begin with the scalar superfields.

Definition: The abelian **supergauge transformation** (local or global) on a left handed scalar superfield Φ_i is defined as:

$$\Phi_i \rightarrow \Phi'_i = e^{-iq_i\Lambda}\Phi_i \quad (3.27)$$

where q_i is the charge of Φ_i under that gauge group and Λ (global), or $\Lambda(x)$ (local), is the parameter of the gauge transformation.

The definition is of course completely equivalent for right-handed scalar fields. For the definition to make sense the transformed field Φ'_i must be a left-handed scalar superfield itself, thus

$$\bar{D}_{\dot{A}}\Phi'_i = 0,$$

and this requires:

$$\bar{D}_{\dot{A}}\Phi'_i = \bar{D}_{\dot{A}}e^{-iq_i\Lambda}\Phi_i = e^{-iq_i\Lambda}\bar{D}_{\dot{A}}\Phi_i - iq_i(\bar{D}_{\dot{A}}\Lambda)e^{-iq_i\Lambda}\Phi_i = -iq_i(\bar{D}_{\dot{A}}\Lambda)\Phi'_i = 0.$$

Thus we must have $\bar{D}_{\dot{A}}\Lambda = 0$, which means that the parameter Λ is also a left-handed superfield.

¹²We promise that we will get back to the corresponding definition for non-abelian transformations.

Next, we move to the vector superfields.

Definition: Given a vector superfield $V(x, \theta, \bar{\theta})$, we define the **abelian supergauge transformation** as

$$V(x, \theta, \bar{\theta}) \rightarrow V'(x, \theta, \bar{\theta}) = V(x, \theta, \bar{\theta}) + i[\Lambda(x, \theta, \bar{\theta}) - \Lambda^\dagger(x, \theta, \bar{\theta})],$$

where the parameter of the transformation Λ is a scalar superfield.

With the expressions for scalar superfields in (3.25) and (3.26), using $\Phi = i\Lambda$, we can show that under supergauge transformations the vector superfield components transform as:

$$C(x) \rightarrow C'(x) = C(x) + A(x) + A^*(x) \quad (3.28)$$

$$\phi(x) \rightarrow \phi'(x) = \phi(x) + \sqrt{2}\psi(x) \quad (3.29)$$

$$M(x) \rightarrow M'(x) = M(x) + F(x) \quad (3.30)$$

$$V_\mu(x) \rightarrow V'_\mu(x) = V_\mu(x) + i\partial_\mu(A(x) - A^*(x)) \quad (3.31)$$

$$\lambda(x) \rightarrow \lambda'(x) = \lambda(x) \quad (3.32)$$

$$D(x) \rightarrow D'(x) = D(x). \quad (3.33)$$

If we look at the transformation this implies for the vector field, this is equivalent to the ordinary abelian gauge transformation for a vector field $V_\mu(x) \rightarrow V'_\mu(x) = V_\mu(x) + \partial_\mu g(x)$, with the gauge parameter given by the scalar component field of Λ , $g(x) = i[A(x) - A^*(x)] = -2\text{Im}(A(x))$. It then immediately follows that the standard field strength for a vector field, $F_{\mu\nu} = \partial_\mu V_\nu - \partial_\nu V_\mu$, is supergauge invariant.

If we demand that our theory is invariant under these gauge transformations we can choose the component fields of Λ in order to eliminate some the remaining reducibility in the representation.

Definition: The **Wess-Zumino (WZ) gauge** is a supergauge transformation of a vector superfield by a scalar superfield with

$$\psi(x) = -\frac{1}{\sqrt{2}}\phi(x), \quad (3.34)$$

$$F(x) = -M(x), \quad (3.35)$$

$$A(x) + A^*(x) = -C(x). \quad (3.36)$$

A vector superfield in the WZ-gauge can then be written:

$$V_{WZ}(x, \theta, \bar{\theta}) = (\theta\sigma^\mu\bar{\theta})[V_\mu(x) + i\partial_\mu(A(x) - A^*(x))] + \theta\theta\bar{\theta}\bar{\lambda}(x) + \bar{\theta}\bar{\theta}\theta\lambda(x) + \theta\theta\bar{\theta}\bar{\theta}D(x), \quad (3.37)$$

which, considered carefully, contains one real scalar field d.o.f., three gauge field d.o.f.¹³ and four fermion d.o.f., meaning it corresponds to the representation $j = \frac{1}{2}$. Unfortunately, supersymmetry transformations break this gauge, meaning that a vector superfield in the

¹³Hang on, where did that last d.o.f. go from $V(x)$? We have a remaining gauge freedom in the choice of the imaginary component of $A(x)$, which is the ordinary gauge freedom of a $U(1)$ field theory. This can be used to eliminate one d.o.f. from the vector field.

WZ-gauge will no longer be in the WZ-gauge after a supersymmetry transformation. Again, the equations of motion will eliminate the auxiliary D -field, as well as one d.o.f. from the gauge field (a massless on-shell gauge boson has two), and two d.o.f. from the (Majorana) fermion formed by λ . Repeating the arguments of mass dimension, if the V_μ is to represent a vector quantum field with $[V_\mu] = M^1$, then the dimension of the vector superfield must be $[V] = M^0$.

Notice that the WZ gauge is particularly convenient for calculations because:

$$V_{WZ}^2 = \frac{1}{2} \theta \theta \bar{\theta} \bar{\theta} [V_\mu(x) + i \partial_\mu (A(x) - A^*(x))] [V^\mu(x) + i \partial^\mu (A(x) - A^*(x))], \quad (3.38)$$

and, since multiplying in any θ or $\bar{\theta}$ into V_{WZ}^2 will then yield zero, we have

$$V_{WZ}^3 = 0,$$

so that

$$e^{V_{WZ}} = 1 + V_{WZ} + \frac{1}{2} V_{WZ}^2.$$

3.6 Exercises

Exercise 3.1

Show that

$$\int d^2\theta \theta \theta = 1.$$

Exercise 3.2

Check that the explicit differential forms of the generators in Eqs. (3.9)–(3.11) fulfil the superalgebra in Eqs. (2.36)–(2.34).

Exercise 3.3

Demonstrate the correctness of the general expression for the left-handed scalar superfield in Eq. (3.25). *Hint:* Prove the spinor identities

$$\begin{aligned} (\theta \sigma^\mu \bar{\theta})(\theta \sigma^\nu \bar{\theta}) &= \frac{1}{2} g^{\mu\nu} \theta \theta \bar{\theta} \bar{\theta}, \\ \theta \partial_\mu \psi \theta \sigma^\mu \bar{\theta} &= \frac{1}{2} \theta \theta \partial_\mu \psi \sigma^\mu \bar{\theta}. \end{aligned}$$

Exercise 3.4

Show the vector superfield supergauge transformation properties for the component fields. *Hint:* Use the field redefinitions:

$$\begin{aligned} \lambda(x) &\rightarrow \lambda(x) - \frac{i}{2} \sigma^\mu \partial_\mu \bar{\phi}(x), \\ D(x) &\rightarrow D(x) - \frac{1}{4} \square C(x). \end{aligned}$$

Exercise 3.5

Find the expression for V_{WZ}^2 .

Chapter 4

Construction of a low-energy supersymmetric Lagrangian

We would now like to construct models, in the form of a field theory Lagrangians, that are invariant under supersymmetry transformations, much in the same way that the Standard Model Lagrangian is invariant under Poincaré transformations. However, just as for ordinary quantum field theory Lagrangians we will need to be able to impose gauge invariance for a choice of gauge group, and we want to limit the models to models that are *renormalisable*, *i.e.* models where any infinities can be cancelled by a finite number of counter-terms. Along the way we will also find a way to deal with the problem of equal masses for the particles in an irreducible representation – in other words, where are all the supersymmetric partners of the Standard Model particles? – taking inspiration from the Standard Model Higgs mechanism using spontaneous symmetry breaking.

4.1 Supersymmetry invariant Lagrangians and actions

As should be well known the relativistic field theory **action**

$$S \equiv \int_R d^4x \mathcal{L}, \quad (4.1)$$

is invariant under supersymmetry transformations if this transforms the Lagrangian by a total derivative term $\mathcal{L} \rightarrow \mathcal{L}' = \mathcal{L} + \partial^\mu f(x)$, where $f(x) \rightarrow 0$ on the surface $S(R)$ of the integration region R . The question then becomes: how can we construct a Lagrangian from superfields with this property?

An infinitesimal **supersymmetry transformation** of a function on superspace $F(x, \theta, \bar{\theta})$ can be written as

$$F'(x, \theta, \bar{\theta}) = \exp(i\alpha Q + i\bar{\alpha}\bar{Q})F(x, \theta, \bar{\theta}) \simeq F(x, \theta, \bar{\theta}) + (i\alpha Q + i\bar{\alpha}\bar{Q})F(x, \theta, \bar{\theta}), \quad (4.2)$$

where α^A and $\bar{\alpha}_{\dot{A}}$ are the infinitesimal parameters of the transformation. If these parameters are constant then we say that this is a **global** supersymmetry transformation which is what we are (mostly) going to concern ourselves with in these notes. Replacing $\alpha \rightarrow \alpha(x)$ gives a **local** supersymmetry transformation.¹ Here we have ignored the action of the P_μ operators

¹Allowing for local supersymmetry transformations leads to a gauge theory of supersymmetry called supergravity, which can incorporate gravity in a natural way.

since they commute with all the Q s, so that this part of the group acts independently, and only results in a translation of the Minkowski coordinate. The total change in the function under the supersymmetry transformation is $\delta_s F = F' - F$, where

$$\delta_s = i(\alpha Q + \bar{\alpha} \bar{Q}). \quad (4.3)$$

We can show that the highest order component fields in θ and $\bar{\theta}$ of a superfield always transform as a total derivative under δ_s , *e.g.* for the general superfield the highest order component field $d(x)$ transforms as²

$$\delta_s d(x) = d'(x) - d(x) = \frac{i}{2}(\partial_\mu \psi(x) \sigma^\mu \bar{\alpha} - \partial_\mu \bar{\lambda}(x) \sigma^\mu \alpha), \quad (4.4)$$

and the same is naturally true for the D component-field of a vector superfield. For a scalar superfield it is the F -field which has this property:

$$\delta_s F(x) = -i\sqrt{2}\partial_\mu \psi(x) \sigma^\mu \bar{\alpha}. \quad (4.5)$$

These highest power θ -components can be isolated by using the projection property of integration in Grassmann calculus writing

$$S = \int_R d^4x \int d^4\theta \mathcal{L},$$

where the Lagrangian density $\mathcal{L}$, which is a function of superfields, is guaranteed to give a supersymmetry invariant action. Note that this constitutes a redefinition of what we mean by the Lagrangian $\mathcal{L}$ when working with superfields. We should in particular be careful when counting the dimension of terms. We have, see Section 3.2, $[\theta] = M^{-1/2}$, which, since $\int d\theta d\bar{\theta} = 1$, leads to $[\int d\theta] = M^{1/2}$. We then have $[\int d^4\theta] = M^2$. Since we must have $[\int d^4\theta \mathcal{L}] = M^4$ for the action to be dimensionless, we need $[\mathcal{L}] = M^2$, which is different from the standard field theory Lagrangian with dimension M^4 .

We can now write down the generic form for a supersymmetric Lagrangian consisting of scalar superfields, where the indices indicate the highest power of θ in the term:

$$\mathcal{L} = \mathcal{L}_{\theta\theta\bar{\theta}\bar{\theta}} + \theta\theta\mathcal{L}_{\bar{\theta}\bar{\theta}} + \bar{\theta}\bar{\theta}\mathcal{L}_{\theta\theta}.$$

Here $\mathcal{L}_{\theta\theta}$ ($\mathcal{L}_{\bar{\theta}\bar{\theta}}$) is a function of *only* left-handed (right-handed) scalar superfields where we project out the F -field by multiplying by $\bar{\theta}\bar{\theta}$ ($\theta\theta$) and integrating over all the θ . These terms form what we call the **superpotential**. Meanwhile, the $\mathcal{L}_{\theta\theta\bar{\theta}\bar{\theta}}$ term is a real valued function of the superfields where we project out the d -field, called the **Kähler potential**.

The requirement of renormalisability puts further restrictions on the fields in $\mathcal{L}$. We can at most have three factors of scalar superfields in each term of the superpotential, for details see *e.g.* Wess & Bagger [7].³ Since the action must be real, the (almost) most general supersymmetry Lagrangian that can be written in terms of scalar superfields is:

$$\mathcal{L} = \Phi_i^\dagger \Phi_i + \bar{\theta}\bar{\theta}W[\Phi] + \theta\theta W[\Phi^\dagger].$$

²A word of warning: we really have to be slightly more careful here with the supersymmetry transformation since the superfields are operator valued functions acting on a state-space. A unitary transformation $U|\psi\rangle$ of the state implies a unitary transformation $U^\dagger A U$ of an operator A on that state.

³The argument is similar to ordinary quantum field theory, terms with couplings with negative mass dimension is forbidden.

Here the first term is called the **kinetic term**,⁴ and W is the symbol for the **superpotential**, which by renormalisability is restricted to

$$W[\Phi] = g_i \Phi_i + m_{ij} \Phi_i \Phi_j + \lambda_{ijk} \Phi_i \Phi_j \Phi_k. \quad (4.6)$$

This means that to specify a supersymmetric Lagrangian we only need to specify the superfield content of the model and the form of the superpotential.

Dimension counting, starting from $[\mathcal{L}] = M^2$ and $[\Phi] = 1$, means that the superpotential terms must each be M^3 , for the couplings this gives $[g_i] = M^2$, $[m_{ij}] = M$ and $[\lambda_{ijk}] = 1$. Notice also that m_{ij} and λ_{ijk} must be symmetric since the superfields commute.

4.2 Abelian gauge theories

We will of course require not only a supersymmetry invariant Lagrangian, but also a gauge invariant Lagrangian, and we start with the requirements for an abelian gauge group.

Let us first look at the transformation of the superpotential W under the gauge transformation of the scalar superfields in (3.27):

$$W[\Phi] \rightarrow W[\Phi'] = g_i e^{-iq_i \Lambda} \Phi_i + m_{ij} e^{-i(q_i + q_j) \Lambda} \Phi_i \Phi_j + \lambda_{ijk} e^{-i(q_i + q_j + q_k) \Lambda} \Phi_i \Phi_j \Phi_k.$$

Requiring gauge invariance, $W[\Phi] = W[\Phi']$, we must have:

$$q_i = 0 \text{ or } g_i = 0, \quad (4.7)$$

$$q_i + q_j = 0 \text{ or } m_{ij} = 0, \quad (4.8)$$

$$q_i + q_j + q_k = 0 \text{ or } \lambda_{ijk} = 0. \quad (4.9)$$

This puts great restrictions on the form of the superpotential and the charge assignments of the superfields, as in ordinary gauge theories.

What then about the kinetic term? This transforms as

$$\Phi_i^\dagger \Phi_i \rightarrow \Phi_i^\dagger e^{iq_i \Lambda^\dagger} e^{-iq_i \Lambda} \Phi_i = e^{iq_i (\Lambda^\dagger - \Lambda)} \Phi_i^\dagger \Phi_i.$$

As in ordinary gauge theories we can introduce a **gauge compensating vector (super)field** V with the appropriate gauge transformation to make the kinetic term invariant under supersymmetry transformations. We write the kinetic term instead as $\Phi_i^\dagger e^{q_i V} \Phi_i$, which gives us an invariant term:⁵

$$\Phi_i^\dagger e^{q_i V} \Phi_i \rightarrow \Phi_i^\dagger e^{iq_i \Lambda^\dagger} e^{q_i (V + i\Lambda - i\Lambda^\dagger)} e^{-iq_i \Lambda} \Phi_i = \Phi_i^\dagger e^{q_i V} \Phi_i.$$

This definition of gauge transformation can be shown to recover the Standard Model **minimal coupling** for the component fields through the covariant derivative

$$D_\mu^i = \partial_\mu - \frac{i}{2} q_i V_\mu,$$

where V_μ is the vector component field of the vector superfield.

⁴The constant in front of this term can always be chosen to be one because we can rescale the whole Lagrangian. It is possible to start from more general Kähler potentials, but this is mostly beyond the scope of these notes, although we will return to modify this term to account for gauge invariance.

⁵In case you were worried: we can use the WZ gauge to show that the new kinetic term $\Phi_i^\dagger e^{q_i V} \Phi_i$ has no term with dimension higher than four, and is thus renormalizable.

4.3 Non-Abelian gauge theories

How do we extend the above to deal with the much more complicated non-abelian gauge theories? Let us take a gauge group G with the Lie algebra of group generators t_a in the fundamental representation

$$[t_a, t_b] = if_{ab}^c t_c, \quad (4.10)$$

where f_{ab}^c are the structure constants. We recap that for an element g in the group G we want to write down a unitary⁶ representation $U(g)$ that transforms a vector Ψ in the representation by $\Psi \rightarrow \Psi' = U(g)\Psi$.⁷ With an exponential map we can write this representation as $U(g) = e^{i\lambda^a t_a}$, where λ^a are the parameters, as you may perhaps have expected.

As in ordinary gauge theories, we simply copy the structure of the abelian supergauge transformation, and transform (a vector of) scalar superfields Φ as

$$\Phi \rightarrow \Phi' = e^{-iq\Lambda^a T_a} \Phi,$$

in a non-abelian supergauge transformation. Here q is the charge of Φ under G . At this point we can choose a representation different from the fundamental, reflected in a different choice for the generators which we write as T_a . Since for gauge groups we are almost exclusively interested in groups defined by a matrix representation $U(g)$ will be a matrix with dimension fixed by the dimension chosen for the representation. Again we can easily show that we must require that the Λ^a are left-handed scalar superfields for Φ to transform to a left-handed scalar superfield.

For the superpotential to be invariant we must now have:

$$g_i U_{ir} = g_r \text{ or } g_i = 0, \quad (4.11)$$

$$m_{ij} U_{ir} U_{js} = m_{rs} \text{ or } m_{ij} = 0, \quad (4.12)$$

$$\lambda_{ijk} U_{ir} U_{js} U_{kt} = \lambda_{rst} \text{ or } \lambda_{ijk} = 0, \quad (4.13)$$

where the indices on U are its matrix indices.

We also want a similar construction for the kinetic terms as for abelian gauge theories, $\Phi^\dagger e^{qV^a T_a} \Phi$, to be invariant under non-abelian gauge transformations. Now, using that we write the generators t_a as hermitian operators,

$$\Phi^\dagger e^{qV^a T_a} \Phi \rightarrow \Phi'^\dagger e^{qV'^a T_a} \Phi' = \Phi^\dagger e^{iq\Lambda^{a\dagger} T_a} e^{qV'^a T_a} e^{-iq\Lambda^a T_a} \Phi,$$

so we have to require that the vector superfield V transforms as:⁸

$$e^{qV'^a T_a} = e^{-iq\Lambda^{a\dagger} T_a} e^{qV^a T_a} e^{iq\Lambda^a T_a}. \quad (4.14)$$

When we look at this as an infinitesimal transformation in Λ we can show that

$$V'^a = V^a + i(\Lambda^a - \Lambda^{a\dagger}) - \frac{1}{2} q f_{bc}^a V^b (\Lambda^{c\dagger} + \Lambda^c) + \mathcal{O}(\Lambda^2),$$

⁶By unitary we mean, as usual, that $U^\dagger = U^{-1}$ so that $U^\dagger U = 1$.

⁷Of, course, you may ask, how do we even know that we can find a unitary representation for a particular Lie group? For the unitary group this is by definition, but is it possible in general? It turns out that this is always true for a subset of Lie groups, called **compact Lie groups**. These are the Lie groups where the parameters vary over a closed interval.

⁸This is independent of our choice of representation for the gauge group for the supergauge transformation.

which reduces to the definition for abelian groups in (3.28) since all the higher order terms contain the structure constant f_{bc}^a , which is zero for abelian groups. If we look at the component vector fields, V_μ^a , these transform just like in a standard non-abelian gauge theory:

$$V_\mu^a \rightarrow V_\mu'^a = V_\mu^a + i\partial_\mu(A^a - A^{a*}) - gf_{bc}^a V_\mu^b (A^c - A^{c*}),$$

in the adjoint representation of the gauge group.

The supergauge transformations of vector superfields can be written more efficiently in a representation independent way as

$$e^{V'} = e^{-i\Lambda^\dagger} e^V e^{i\Lambda},$$

and the inverse transformation is then given by

$$e^{-V'} = e^{-i\Lambda} e^{-V} e^{i\Lambda^\dagger},$$

where $\Lambda \equiv q\Lambda^a T_a$ and $V \equiv qV^a T_a$, such that $e^V e^{-V} = e^{V'} e^{-V'} = 1$.⁹

4.4 Supersymmetric field strength

There is one type of term missing from the supersymmetric Lagrangian we are constructing, namely field strength terms that make the gauge fields dynamical, *e.g.* terms to describe the electromagnetic field strength $-\frac{1}{4}F^{\mu\nu}F_{\mu\nu}$ in QED.

Definition: The **supersymmetric field strength** for a vector superfield V is defined by the spinor (matrix) scalar superfields W_A and $\bar{W}_{\dot{A}}$ given by

$$W_A \equiv -\frac{1}{4}\bar{D}\bar{D}e^{-V}D_A e^V,$$

and

$$\bar{W}_{\dot{A}} \equiv -\frac{1}{4}DDe^{-V}\bar{D}_{\dot{A}} e^V,$$

where $V = V^a T_a$.

For an abelian gauge field this definition reduces to

$$\begin{aligned} W_A &= -\frac{1}{4}\bar{D}\bar{D}D_A V, \\ \bar{W}_{\dot{A}} &= -\frac{1}{4}DD\bar{D}_{\dot{A}} V, \end{aligned}$$

where V is simply the vector superfield of the gauge group.

We can show that the components of W_A ($\bar{W}_{\dot{A}}$) are left-handed (right-handed) scalar superfields, and that both $\text{Tr}[W^A W_A]$ and $\text{Tr}[\bar{W}_{\dot{A}} \bar{W}^{\dot{A}}]$ are supergauge invariant constructions, and thus potential terms in the supersymmetry Lagrangian. Firstly

$$\bar{D}_{\dot{A}} W_A = -\frac{1}{4}\bar{D}_{\dot{A}} \bar{D}\bar{D}e^{-V}D_A e^V = 0,$$

⁹Notice that despite the non-commutative nature of the matrices involved, the identity $e^A e^{-A} = 1$ holds. See Exercise 1.19.

because from Eq. (3.16), $\bar{D}^3 = 0$, and similarly for $\bar{W}_A$, showing that they are both superfields, and as such can be used to form a supersymmetric Lagrangian.

Under a supergauge transformation we have:

$$\begin{aligned}
W_A \rightarrow W'_A &= -\frac{1}{4}\bar{D}\bar{D}e^{-i\Lambda}e^{-V}e^{i\Lambda^\dagger}D_Ae^{-i\Lambda^\dagger}e^Ve^{i\Lambda} \\
(\bar{D}_A\Lambda = 0) &= -\frac{1}{4}e^{-i\Lambda}\bar{D}\bar{D}e^{-V}e^{i\Lambda^\dagger}D_Ae^{-i\Lambda^\dagger}e^Ve^{i\Lambda} \\
(D_A\Lambda^\dagger = 0) &= -\frac{1}{4}e^{-i\Lambda}\bar{D}\bar{D}e^{-V}D_Ae^Ve^{i\Lambda} \\
&= -\frac{1}{4}e^{-i\Lambda}\bar{D}\bar{D}e^{-V}[(D_Ae^V)e^{i\Lambda} + e^V(D_Ae^{i\Lambda})] \\
&= e^{-i\Lambda}W_Ae^{i\Lambda} - \frac{1}{4}e^{-i\Lambda}\bar{D}\bar{D}D_Ae^{i\Lambda}.
\end{aligned} \tag{4.15}$$

We are free to add zero to (4.15) in the form of $-\frac{1}{4}e^{-i\Lambda}\bar{D}D_A\bar{D}e^{i\Lambda} = 0$,¹⁰ giving

$$\begin{aligned}
W'_A &= e^{-i\Lambda}W_Ae^{i\Lambda} - \frac{1}{4}e^{-i\Lambda}\bar{D}\{\bar{D}, D_A\}e^{i\Lambda} \\
&= e^{-i\Lambda}W_Ae^{i\Lambda} + \frac{1}{2}e^{-i\Lambda}\bar{D}_A\sigma^\mu_{AB}\epsilon^{\dot{A}\dot{B}}P_\mu e^{i\Lambda} \\
&= e^{-i\Lambda}W_Ae^{i\Lambda},
\end{aligned}$$

where we have used Eq. (3.15) to replace the anti-commutator. This means that the following trace is gauge invariant:

$$\begin{aligned}
\text{Tr}[W'^AW'_A] &= \text{Tr}[e^{-i\Lambda}W^Ae^{i\Lambda}e^{-i\Lambda}W_Ae^{i\Lambda}] \\
&= \text{Tr}[e^{i\Lambda}e^{-i\Lambda}W^AW_A] = \text{Tr}[W^AW_A],
\end{aligned}$$

and can be used in the Lagrangian.

If we expand W_A in the component fields we find, as we might have hoped, that it contains the ordinary field strength tensor:

$$F_{\mu\nu}^a = \partial_\mu V_\nu^a - \partial_\nu V_\mu^a + gf_{bc}^a V_\mu^b V_\nu^c$$

and that the trace indeed contains terms with $F_{\mu\nu}^a F^{\mu\nu a}$.

4.5 The (almost) complete supersymmetric Lagrangian

We can now compile all our results up to now to write down the complete Lagrangian for a supersymmetric theory with (possibly) non-abelian gauge symmetries:¹¹

$$\mathcal{L} = \Phi^\dagger e^V \Phi + \delta^2(\bar{\theta})W[\Phi] + \delta^2(\theta)W[\Phi^\dagger] + \frac{1}{2T(R)}\delta^2(\bar{\theta})\text{Tr}[W^AW_A]. \tag{4.16}$$

¹⁰Which is zero because Λ is a left-handed scalar superfield, $\bar{D}_A\Lambda = 0$.

¹¹There is no hermitian conjugate of the field strength term, and a slightly odd normalisation using 1/2 instead of 1/4. This is because the term can be proven to be real, although this is sometimes overlooked in the literature when authors instead use

$$\mathcal{L} \sim \frac{1}{4T(R)}\delta^2(\bar{\theta})\text{Tr}[W^AW_A] + \frac{1}{4T(R)}\delta^2(\theta)\text{Tr}[\bar{W}^A\bar{W}_A],$$

which looks more symmetric with respect to the superpotential part.

Here, $T(R)$ given by $T(R)\delta_{ab} = \text{Tr}[T_a T_b]$ is called the **Dynkin index** of the representation R of the gauge group using the generators T_a .

The Dynkin index appears in the field strength term to correctly normalise the energy density for the chosen representation R of the gauge group. To see that this factor cancels in a natural way, note that since W_A is spanned by T_a for a given representation, we can write $W_A = W_A^a T_a$, where W_A^a are some coefficients. Then

$$\text{Tr}[W^A W_A] = W^{aA} W_A^b \text{Tr}[T_a T_b] = W^{aA} W_A^b \delta_{ab} T(R) = T(R) W^{aA} W_A^a. \quad (4.17)$$

In addition to the above, it is also possible to add a pure vector superfield term as part of the Kähler potential, which is not constructed from scalar superfields, of the form $\mathcal{L}_{FI} \sim -kV$ where V is the vector superfield and k some constant. This kind of term is called a **Fayet–Iliopoulos term**. However, this is only possible for non-abelian gauge groups since a term $-kV^a$ could not be supergauge invariant.¹²

4.6 Finding the equations of motion

To find the equations of motion from our supersymmetric Lagrangian construction we can now perform the integration over superspace coordinates θ and $\bar{\theta}$ in the action and then apply the standard Euler–Lagrange field equations that minimise the action for the component fields ϕ_i :

$$\frac{\partial \mathcal{L}}{\partial \phi_i} - \partial_\mu \left(\frac{\partial \mathcal{L}}{\partial (\partial_\mu \phi_i)} \right) = 0. \quad (4.18)$$

We will now discuss some of the generic properties of this solution without specifying field content or gauge groups.

As already intimated the auxiliary component fields with unusual mass dimension, F_i and D^a , from scalar Φ_i and vector V^a superfields, respectively, will be eliminated.¹³ Starting with the F_i -fields these occur only in the kinetic term and the superpotential. From the kinetic term the contribution is $|F_i|^2$ as this is the only term with F_i that can survive the θ -integration. In the superpotential only the terms where a number of scalar fields A_j with no θ -coordinate multiply F_i can survive the integration. Their derivatives with respect to F_i can be expressed in terms of the superpotential functional as

$$W_i \equiv \frac{\partial W[A_1, \dots, A_n]}{\partial A_i}, \quad (4.19)$$

where the scalar superfields have been replaced by their scalar component fields. This means that the Euler–Lagrange equation for F_i is

$$\frac{\partial \mathcal{L}}{\partial F_i^*} = F_i + W_i^* = 0,$$

which is used to eliminate F_i .

¹²Why is the abelian case supergauge invariant then? Well, the superspace integration in the action will isolate the D -term from V , and this term is itself invariant under abelian supergauge transformations, see Eq. (3.33). For non-abelian supergauge transformations the various V^a fields mix.

¹³In the following we use the notation in Eqs. (3.25) and (3.37).

For a concrete example of how this can be worked out, see Exercise 4.1. Generalising this also leads to an expression for the action for any number of scalar superfields Φ_i , ignoring gauge interactions, which is:

$$S = \int d^4x \{ i \partial_\mu \bar{\psi}_i \sigma^\mu \psi_i - A_i^* \square A_i - \frac{1}{2} W_{ij} \psi_i \psi_j - \frac{1}{2} W_{ij}^* \bar{\psi}_i \bar{\psi}_j - |W_i|^2 \}$$

where W_{ij} , given by

$$W_{ij} \equiv \frac{\partial^2 W[A_1, \dots, A_n]}{\partial A_i \partial A_j}, \quad (4.20)$$

is called the **fermionic mass matrix**.

The D^a -fields occur in the kinetic terms and the field strength terms. In the kinetic term, in the Wess–Zumino gauge, it is relatively easy to see that D^a appears only once in the expansion of the exponential as $q A_i^* T_{ij}^a A_j D^a$ in the action. In the field strength term one can show with some more effort that the contribution is $D^a D^a$ in the Wess–Zumino gauge. For abelian gauge fields there is also a possible kD contribution from the Fayet–Iliopoulos term. This leads to the Euler–Lagrange equation

$$\frac{\partial \mathcal{L}}{\partial D^a} = q A_i^* T_{ij}^a A_j + 2 D^a = 0,$$

which can be used to eliminate D^a .

Note that the terms discussed here with F - and D -fields exhaust all the non-derivative terms that only have scalar fields since no terms with only A -fields can survive the θ -integration. Thus the complete non-derivative scalar field contribution to the Lagrangian, called the **scalar potential**, can be written as

$$V(A_i, A_i^*) = \sum_i |F_i|^2 + \sum_a D^a D^a = \sum_i \left| \frac{\partial W[A_1, \dots, A_n]}{\partial A_i} \right|^2 + \frac{1}{4} \sum_a q^2 (A_i^* T_{ij}^a A_j)^2. \quad (4.21)$$

4.7 The hierarchy problem

There is a fundamental problem with having scalar particles in a quantum field theory. Let us take the Standard Model Higgs boson h as an example, however, the following would be true for any scalar particle. If we naively calculate loop-corrections to its mass in self-energy diagrams like the ones shown in Fig. 4.1, where f is a fermion and s some other scalar, they both diverge due to their loop momenta integrals, meaning they are infinite. This then needs what is called **regularisation** in order to yield a finite answer.

There are different ways of achieving regularisation. Since we know that the Standard Model is an incomplete theory, at least when we go up to Planck scale energies where we need an unknown quantum theory of gravity, we can introduce a **cut-off regularisation** where we limit the integral in the loop-correction to momenta below a scale Λ_{UV} . Then the loop-correction to the Higgs mass is, at leading order in Λ_{UV} ,

$$\Delta m_h^2 = -\frac{|\lambda_f|^2}{8\pi^2} \Lambda_{UV}^2 + \frac{\lambda_s}{16\pi^2} \Lambda_{UV}^2 + \dots \quad (4.22)$$

where λ_f and λ_s are the couplings of f and s to the Higgs, respectively. Picking Λ_{UV} suggestively as the Planck scale, $\Lambda_{UV} = M_P = 2.44 \times 10^{18}$ GeV, meaning that the quantum

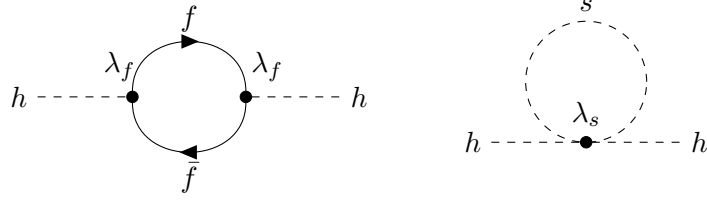


Figure 4.1: Higgs self-energy diagrams with fermions f (left) and scalars s (right).

corrections to the Higgs mass are some 15 orders of magnitude greater than its measured value.

We observe that the difference in sign between the fermions and bosons in (4.22) means that it could *in principle* be possible that these huge contributions cancel to keep $m_h \sim 125 \text{ GeV}$ as measured, however, as physicists we should worry why the Universe seems to be organised in such a strange way without any specific cause. This is known as the **Higgs fine tuning problem**, or the **hierarchy problem** because it is fundamentally a problem of why the electroweak scale at around 100 GeV where the Higgs lives, is so much lighter than the Planck scale of 10^{18} GeV , when the former should in principle be pulled up to the latter by the sensitivity of scalars to loop corrections.¹⁴

Enter supersymmetry to the rescue: with unbroken supersymmetry we find that we automatically have $|\lambda_f|^2 = \lambda_s$ and exactly twice as many scalar as fermion degrees of freedom running around in loops. This provides a magic cancellation of the quadratic divergence in Eq. (4.22). To see that this relation between the couplings holds, notice that it is the superpotential term $W \sim \lambda_{ijk} \Phi_i \Phi_j \Phi_k$ which is responsible for generating Lagrangian Yukawa terms of the form $\lambda_{ijk} \psi_i \psi_j A_k$ that couple two fermions to a scalar. At the same time, in the scalar potential, that same superpotential term is responsible for the terms

$$V(A, A^*) \sim \left| \frac{\partial W}{\partial A_i} \right|^2 = |\lambda_{ijk}|^2 A_j^* A_k^* A_j A_k. \quad (4.23)$$

When the scalar field A_k is the Higgs field, the loop fermion is represented by $\psi_i = \psi_j$ and the second scalar by A_j , these two terms are responsible for the two types of vertices in Fig. 4.1 with $\lambda_f = \lambda_{ijk}$ and $\lambda_s = |\lambda_{ijk}|^2$. The fact that there are loop contributions from two scalars for each fermion due to the same number of states (two scalar states per fermion state), means that unbroken supersymmetry predicts an exact cancellation in Eq. (4.22).¹⁵

4.8 Vacuum energy

This measured value of the dark energy in terms of an energy scale is $\Lambda_{DE} \simeq 2.4 \times 10^{-3} \text{ eV}$. This can be interpreted as vacuum energy, *i.e.* the energy predicted from the contribution of Feynman diagrams with *no* external particles. However, the predicted value of the vacuum

¹⁴With some background in renormalisation you may ask: What about choosing dimensional regularisation instead where there is no cut-off scale? That would in principle work, however, as soon as you introduce *any* new particle (significantly) heavier than the Higgs there is still a quadratic correction of the form of (4.22) with the new particle masses as the scale, meaning that we cannot complete the Standard Model at a higher scale without reintroducing the problem.

¹⁵To keep the argument simple we have avoided the contributions from vector bosons, however, we can show that these also cancel exactly against the contributions from their fermionic partners.

energy in the Standard Model is $\Lambda_{VE} \sim M_P$.¹⁶ If we want to compare energy densities, we should compare Λ_{DE}^4 to Λ_{VE}^4 ,¹⁷ which means that the prediction is off from the measurement by some 120 orders of magnitude, which is said to be a record for the greatest ever discrepancy between theory and experiment. This problem is the **hierarchy problem for vacuum energy**.

Can supersymmetry save us here as well? Sort of. For an unbroken global supersymmetry we can use the **non-renormalisation theorem**¹⁸ of Grisaru, Roach and Siegel (1979) [8] to show that the prediction in supersymmetry is exactly $\Lambda_{VE} = 0$.

Theorem: Non-renormalisation theorem

All higher order contributions to the effective supersymmetric action for a process S_{eff} can be written:

$$S_{\text{eff}} = \sum_n \int d^4x_i \dots d^4x_n d^4\theta F_1(x_1, \bar{\theta}, \theta) \times \dots \times F_n(x_1, \bar{\theta}, \theta) \times G(x_1, \dots, x_n), \quad (4.24)$$

where F_i are products of the n external superfields and their covariant derivatives, and G is a supersymmetry invariant function.

While this is a rather technical statement, that fact that vacuum diagrams do not have external superfields means that the higher order contributions are all zero, and since there are no tree-level contributions for vacuum diagrams – you must have loops to not have external legs – there can be no contribution at all.

This was a victory for supersymmetry before the discovery of dark energy, when the problem was instead to prove that $\Lambda = 0$, modulo the fact that the breaking supersymmetry, as we will be doing in Chapter 5, changes this prediction. As we shall see there, the scale of this contribution has to be the mass scale of the supersymmetric particles, so with for example $m_{SUSY} \simeq 2\text{TeV}$, we have $m_{SUSY}/\Lambda_{DE} \simeq 10^{15}$, some 15 orders of magnitude too large, which is better than the Standard Model prediction of $M_P/\Lambda_{DE} = 10^{30}$, but still a bit off the measured value.¹⁹

So now we are left with showing that the contribution is very small but non-zero, which is in general though to be a much harder problem than finding models where it is exactly zero. However, in supergravity something interesting happens. Introducing a local supersymmetry the scalar potential is not simply given by the superpotential derivatives in (4.21), but instead

¹⁶The origin of this is just the same as the quadratic divergence for the Higgs mass. It is the same type of diagrams contributing, only without external legs.

¹⁷In natural units the unit of energy density is $[\rho] = [E/V] = M/L^3 = M^4$.

¹⁸The name of this theorem is a bit funny. Why? Well, mainly because the result is not about not being able to renormalise the theory, but about not *needing* to renormalise certain parts of it. Another consequence of the theorem, which gives it its name, is that the couplings of the superpotential do not need separate renormalisation. The renormalisation of the superfields suffices.

¹⁹So, one would be tempted to say that supersymmetry has solved half the problem. On a more serious note, there is a significant difference in that the Standard Model prediction requires a cut-off to be finite, while supersymmetry predicts a finite but too large value.

is (ignoring the effects of gauge fields)

$$V(A, A^*) = e^{K/M_P} \left[K_{ij}(D_i W)(D_j W^*) - \frac{3}{M_P^2} |W|^2 \right], \quad (4.25)$$

where $K_{ij} = \partial_i \partial_j K[A, A^*]$ is the **Kähler metric**, D_i the **Kähler derivative** $D_i = \partial_i + \frac{1}{M_P^2}(\partial_i K)$, and all the derivatives are with respect to the scalar fields in the Kähler potential K and superpotential W . In the $M_P \rightarrow \infty$ limit, the low energy limit, we see that we recover the flat space result of Eq. (4.21). What is important to notice is that there is now a second *negative* term in the potential that can in principle cancel the positive supersymmetry breaking contribution, however, given the large size of the breaking contribution this will come at the price of fantastic fine-tuning unless some mechanism can be found where this is natural.

4.9 Exercises

Exercise 4.1

Write down the Lagrangian and find the action of the simplest possible supersymmetric field theory with a single scalar superfield, without gauge transformations, in terms of component fields, and show that it contains no kinetic terms for the F -field. Then show how the F -field can be eliminated by the equations of motion. *Hint:* The kinetic part of the action will turn out to be

$$S_{\text{kin}} = \int d^4x \left\{ -A^*(x) \square A(x) + |F(x)|^2 + i \partial_\mu \psi(x) \bar{\sigma}^\mu \psi(x) \right\}. \quad (4.26)$$

To show this you may have use of the identities in Sec. 2.4.1.

Exercise 4.2

Show that the supersymmetric field strength term for an abelian gauge field can be written in terms of component fields as

$$\mathcal{L} = -i\lambda(x)\sigma^\mu\partial_\mu\bar{\lambda}(x) + 2D^2(x) - \frac{1}{4}F_{\mu\nu}(x)F^{\mu\nu}(x) - \frac{i}{4}F_{\mu\nu}(x)\tilde{F}^{\mu\nu}(x). \quad (4.39)$$

Try to find an argument why the last term with the dual field tensor $\tilde{F}^{\mu\nu}(x)$ can be ignored when finding the equations of motion. *Hints:* To get started it may be productive to consider the coordinate change $y^\mu = x^\mu + i\theta\sigma^\mu\bar{\theta}$. You will also likely need the following algebraic relationship

$$(\sigma^{\mu\nu}\theta)^A F_{\mu\nu}(\sigma^{\rho\sigma}\theta)_A F_{\rho\sigma} = -\frac{1}{2}\theta\theta[F_{\mu\nu}F^{\mu\nu} + iF_{\mu\nu}\tilde{F}^{\mu\nu}]. \quad (4.40)$$

Exercise 4.3

Extend Exercise 1 to include a single abelian gauge group under which the scalar superfield has charge q . Simplify your answer using the covariant derivative $D_\mu \equiv \partial_\mu - i\frac{q}{2}V_\mu$. *Hint:* The spinor relationship

$$\theta\psi\bar{\theta}\bar{\psi}(\theta\sigma^\mu\bar{\theta}) = -\frac{1}{4}(\theta\theta)(\bar{\theta}\bar{\theta})\bar{\psi}\bar{\sigma}^\mu\psi, \quad (4.48)$$

may also come in handy.

Exercise 4.4

Find the Lagrangian for supersymmetric QED (SQED) in terms of component fields, eliminating any non-dynamical auxiliary fields. Supersymmetric QED is the smallest (in terms of field content) supersymmetric theory that has the particles and interactions in QED as a subset.

Chapter 5

Breaking supersymmetry

In Chapter 2 we saw that supersymmetry predicts scalar partner particles with the same mass as the known fermions, and new fermions with the same mass as the known vectors, living in the same irreducible representation. These, somewhat unfortunately, contradict experiment by not existing. In this chapter we will look at how we can break supersymmetry so that the new states become more massive, and avoid current experimental bounds. This discussion will be held at a general level, and concrete realisations of this breaking will have to wait until Chapter 6.

5.1 Spontaneous supersymmetry breaking

In the Standard Model we have already faced this problem: the vector bosons should remain massless under the gauge symmetry of the model. Yet, they are observed to be very massive. This is solved with the introduction of the Higgs mechanism and **spontaneous symmetry breaking** in the **scalar potential**.¹ The idea is that while there is a symmetry of the Lagrangian – in the Standard Model the gauge symmetry $SU(3)_c \times SU(2)_L \times U(1)_Y$ – this may not be a symmetry for the particular vacuum state of the theory (the lowest energy state), thereby allowing the properties of the vacuum to supply the masses. In the Standard Model this is achieved by the shape of the scalar potential having a minimum for a non-zero value of the Higgs field which means that the value of the field there – its vacuum expectation value – can contribute to the masses. Would it not be great if we could have a similar *spontaneous symmetry breaking* in order to break supersymmetry this way and boost the masses of supersymmetric particles beyond current limits?

To find the properties of the vacuum state for our supersymmetric models we start by pointing out that we can write the supersymmetric Hamiltonian as

$$H = \frac{1}{4}(Q_1 \bar{Q}_1 + \bar{Q}_1 Q_1 + Q_2 \bar{Q}_2 + \bar{Q}_2 Q_2).$$

To see this, consider first

$$\begin{aligned} \{Q_A, \bar{Q}_{\dot{B}}\} \bar{\sigma}^{\nu \dot{B} A} &= 2\sigma^\mu_{A\dot{B}} \bar{\sigma}^{\nu \dot{B} A} P_\mu \\ &= 2 \text{Tr}[\sigma^\mu \bar{\sigma}^\nu] P_\mu \\ &= 4g^{\mu\nu} P_\mu = 4P^\nu. \end{aligned}$$

¹The potential of the Lagrangian are those terms not containing derivatives of the fields (kinetic terms). The scalar potential are such terms that contain only scalar fields.

Now,

$$\begin{aligned} H &= P^0 = \frac{1}{4} \{Q_A, \bar{Q}_{\dot{B}}\} \bar{\sigma}^{0\dot{B}A} \\ &= \frac{1}{4} (Q_1 \bar{Q}_{\dot{1}} + \bar{Q}_{\dot{1}} Q_1 + Q_2 \bar{Q}_{\dot{2}} + \bar{Q}_{\dot{2}} Q_2). \end{aligned}$$

As discussed in Section 2.4 we have $Q_A^\dagger = \bar{Q}_{\dot{A}}$. Thus the Hamiltonian is **semipositive definite**, *i.e.* $\langle \Psi | H | \Psi \rangle \geq 0$ for any state $|\Psi\rangle$, so the energy of any state is always non-negative.

Imagine now that there exists some lowest lying state (or possibly degenerate states), the ground state(s) $|0\rangle$, that has vanishing energy $\langle 0 | H | 0 \rangle = 0$. These states are supersymmetric – meaning invariant under the supersymmetry transformation – since, to fulfil the energy assumption, we must have

$$Q_A |0\rangle = \bar{Q}_{\dot{A}} |0\rangle = 0 \quad \text{for } \forall A, \dot{A}, \quad (5.1)$$

and are thus invariant under the supersymmetry transformations given by (4.3)

$$\delta_S |0\rangle = (\alpha^A Q_A + \bar{\alpha}_{\dot{A}} \bar{Q}^{\dot{A}}) |0\rangle = 0. \quad (5.2)$$

The vanishing energy means that at this supersymmetric minimum of the potential the scalar potential must contribute zero, $V(A, A^*) = 0$, and thus, from Eq. (4.21),

$$\left\langle \frac{\partial W}{\partial A_i} \right\rangle = 0 \quad \text{and} \quad \langle A_i^* T_{ij}^a A_j \rangle = 0.$$

Conversely, if the scalar potential does contribute to the energy for the vacuum (ground state) $|0\rangle$, meaning that either

$$\left\langle \frac{\partial W}{\partial A_i} \right\rangle \neq 0 \quad \text{or} \quad \langle A_i^* T_{ij}^a A_j \rangle \neq 0,$$

in the minimum of the potential for some A_i , then *supersymmetry must be broken*. As in the Standard Model, the Lagrangian is still (super)symmetric, but $|0\rangle$ is not because (5.1) can no longer hold for all the Q s.

The **O’Raifeartaigh model** (1975) [9] is an example of a model that spontaneously breaks supersymmetry with three scalar superfields X , Y , Z , and the superpotential

$$W = \lambda Y Z + g X (Z^2 - m^2), \quad (5.3)$$

where λ , g and m are real non-zero parameters. The scalar potential is

$$\begin{aligned} V(A, A^*) &= \left| \frac{\partial W}{\partial A_X} \right|^2 + \left| \frac{\partial W}{\partial A_Y} \right|^2 + \left| \frac{\partial W}{\partial A_Z} \right|^2 \\ &= |g(A_Z^2 - m^2)|^2 + |\lambda A_Z|^2 + |\lambda A_Y + 2g A_X A_Z|^2, \end{aligned} \quad (5.4)$$

which can never be zero because setting $A_Z = 0$, which is needed for the second term, gives a non-zero contribution $g^2 m^4$ from the first term. Since the expectation value at the minimum that breaks supersymmetry is $\langle 0 | \frac{\partial W_i}{\partial A_i} | 0 \rangle$, and $F_i = -\frac{\partial W_i^*}{\partial A_i}$, the condition for spontaneous supersymmetry breaking with the O’Raifeartaigh mechanism can be written

$$\langle F_i \rangle \equiv \langle 0 | F_i(x) | 0 \rangle > 0, \quad (5.5)$$

hence it is often given the more generic name **F-term breaking** outside of the specific O’Raifeartaigh model. In general *F*-term breaking it is the **vacuum expectation value** (**vev**) of the auxiliary field of a scalar superfield that supplies the breaking.

In a gauge theory, a similar mechanism is found by adding the Fayet-Iliopolous term $\mathcal{L}_{FI} \sim 2kV$ where V is an abelian vector superfield. The vev of the $D(x)$ auxiliary field, $\langle D \rangle = \langle 0|D(x)|0 \rangle$, will then create a non-zero scalar potential.² This is called the **Fayet-Iliopolous model**, or **D-term breaking**.

5.2 Supertrace

Unfortunately, the above does not work in practice if all particles are at a low energy scale. The problem is that at *tree level* the **supertrace**, STr , of the mass matrix $\mathcal{M}$, meaning the weighted sum of eigenvalues of the mass matrix, can be shown to vanish, $\text{STr } \mathcal{M}^2 = 0$ even after spontaneous supersymmetry breaking.³

Definition: The **supertrace** is given by

$$\text{STr } (\mathcal{M}^2) \equiv \sum_s (-1)^{2s} (2s + 1) \text{Tr}[M_s^2], \quad (5.6)$$

where $\mathcal{M}$ is the complete mass matrix of the Lagrangian, s is the spin of particles and M_s is the mass matrix of all spin- s particles.

For a theory with only scalar superfields, with two fermionic and two bosonic degrees of freedom each, and with mass matrices $M_{1/2}$ and M_0 containing the masses of each particle on the diagonal, this means that $\text{STr } (\mathcal{M}^2) = \text{Tr}[M_0^2 - 2M_{1/2}^2] = 0$. While before spontaneous supersymmetry breaking each fermion and boson in the same superfield had equal masses, we now have that the sum of (the square of) the fermion and boson masses over all the superfields is the same. The crucial consequence is that not all the scalar partners can be heavier than our known fermions in order to balance out the supertrace relationship.

Adding vector superfields does not help because the fermions there have spin-1 vector bosons partners, and in the supertrace these contributions get opposite sign, and cancel each other out. This relationship is a tree level relationship, meaning it does not take into account quantum corrections. However, accounting for those and using strong couplings to move the supertrace value away from zero, the effect does not seem to be enough to build a consistent model.

The solution to the supertrace dilemma is to put the spontaneous supersymmetry breaking up at some higher energy scale $\sqrt{\langle F \rangle}$ – significantly higher than the electroweak scale – that we do not currently have experimental access to, so that there are also superfields with fermion and boson partners with much higher masses that fulfil the supertrace relationship, and the electroweak scale fermions and bosons become only a small perturbation on it.

²It is *always* the auxiliary fields’ fault!

³See Ferrara, Girardello and Palumbo (1979) [10].

5.3 Soft breaking

There are many alternatives in the literature of how we can break supersymmetry spontaneously at a high energy scale. The names of some popular examples that we will return to are:

- Planck-scale Mediated Symmetry Breaking (PMSB)
- Gauge Mediated Symmetry Breaking (GMSB)
- Anomaly Mediated Symmetry Breaking (AMSB)

Common to all of them is that the mass scale of the particles in the supersymmetry breaking fields is so high, that we can not observe their effects directly. What we can do instead is to parameterise our ignorance by adding effective explicit supersymmetry breaking terms to our low-energy Lagrangian.

However, we cannot simply add arbitrary terms to the Lagrangian. The terms we can add are so-called **soft terms** with couplings of mass dimension one or higher. The disallowed terms with smaller mass dimension are terms that can lead to divergences in loop contributions to scalar masses (such as the Higgs) that are quadratic or worse (because of the high dimensionality of the fields in the loops). This will re-introduce the hierarchy problem.

The allowed terms are in superfield notation as follows:

$$\begin{aligned}\mathcal{L}_{\text{soft}} = & -\frac{1}{4T(R)}M\theta\theta\bar{\theta}\bar{\theta}\text{Tr}\{W^AW_A\} - \frac{1}{6}a_{ijk}\theta\theta\bar{\theta}\bar{\theta}\Phi_i\Phi_j\Phi_k \\ & -\frac{1}{2}b_{ij}\theta\theta\bar{\theta}\bar{\theta}\Phi_i\Phi_j - t_i\theta\theta\bar{\theta}\bar{\theta}\Phi_i + \text{h.c.} \\ & -m_{ij}^2\theta\theta\bar{\theta}\bar{\theta}\Phi_i^\dagger\Phi_j,\end{aligned}\tag{5.7}$$

where $M, a_{ijk}, b_{ij}, t_i \in \mathbb{C}$ and $m_{ij}^2 \in \mathbb{R}$ are the couplings. Note that these terms are explicitly *not* supersymmetric. From the $\theta\theta\bar{\theta}\bar{\theta}$ -factors we see that only the lowest order in θ component fields of the superfields contribute.

There are also some terms that are called "maybe-soft" terms:

$$\mathcal{L}_{\text{maybe}} = -\frac{1}{2}\theta\theta\bar{\theta}\bar{\theta}c_{ijk}\Phi_i^\dagger\Phi_j\Phi_k + \text{h.c.}\tag{5.8}$$

with $c_{ijk} \in \mathbb{C}$. This last – oft ignored – type of term is soft as long as none of the scalar superfields is a singlet under all gauge symmetries. It is, however, quite difficult to get large values for c_{ijk} with spontaneous supersymmetry breaking. In the above terms we have not specified any gauge symmetry, which will, in the same way as it did for the superpotential, severely restrict the allowed terms. However, it turns out that soft-terms are responsible for most of the parameters in supersymmetric theories!

We can instead write the soft terms in terms of their component fields as

$$\begin{aligned}\mathcal{L}_{\text{soft}} = & -\frac{1}{2}M\lambda^A\lambda_A - \left(\frac{1}{6}a_{ijk}A_iA_jA_k + \frac{1}{2}b_{ij}A_iA_j + t_iA_i + \frac{1}{2}c_{ijk}A_i^*A_jA_k + \text{c.c.}\right) \\ & -m_{ij}^2A_i^*A_j.\end{aligned}\tag{5.9}$$

Here we clearly see that the m_{ij} couplings provide extra mass terms to the scalar partners A_i of the fermions in the scalar superfields, and M to the fermion partners λ_A of the vectors in

the vector superfields. Thus, they split the masses of the partners as we wanted to achieve with breaking supersymmetry.

Now, what happens to the hierarchy problem? Restricting ourselves to soft supersymmetry breaking terms guarantees that we end up with contributions to the Higgs mass of at most

$$\Delta m_h^2 = -\frac{\lambda_s}{16\pi^2} m_s^2 \ln \frac{\Lambda_{UV}^2}{m_s^2} + \dots, \quad (5.10)$$

at the leading order in Λ_{UV} , where m_s is the mass scale of the soft term. The correction to the Higgs mass is thus proportional to the masses of the supersymmetric particles.

This is the most important argument in favour of supersymmetry existing at low energy scales where we can detect it, because m_s can not be too large if we want the above corrections to be reasonably small. This is called the **little hierarchy problem**, and numerically, given that $16\pi^2 \sim 100$, and the couplings λ_s are not expected to be above unity, this means that we want $m_s \sim \mathcal{O}(1 \text{ TeV})$ in order to keep the cancellations reasonable and of the order of the measured Higgs mass, with little fine-tuning.

5.4 Exercises

Exercise 5.1

Show that adding a Fayet-Iliopolous term to the SQED constructed in Exercise 4.4 will break supersymmetry spontaneously.

Exercise 5.2

Write down the supertrace relationship for a spontaneously broken SQED.

Chapter 6

The Minimal Supersymmetric Standard Model

The Minimal Supersymmetric Standard Model (MSSM) is a minimal supersymmetric model in the sense that it has the smallest field (and gauge) content consistent with the known Standard Model fields. We will now construct this model on the basis of what we have learnt the previous chapters, and look at some of its consequences.

6.1 MSSM field content

To specify a supersymmetric model we need to specify the superfield content and the gauge symmetries. The gauge symmetry is already given as the Standard Model $SU(3)_C \times SU(2)_L \times U(1)_Y$. Previously we learnt that each (left-handed) scalar superfield S has a (left-handed) Weyl spinor ψ_A , a complex scalar A_s and an auxiliary complex scalar field F_s , since they are a $j = 0$ representation of the superalgebra. After an application of the equations of motion ψ_A and A_s have two fermionic and two bosonic degree of freedom remaining respectively, while the auxiliary field has been eliminated along with two fermionic d.o.f..

In order to construct a Dirac fermion, which are plentiful in the Standard Model, we need a right-handed Weyl spinor as well. We can acquire the needed right-handed Weyl spinor from the hermitian conjugate $\bar{T}^\dagger$ of a different scalar superfield $\bar{T}$ with the right-handed Weyl spinor $\bar{\varphi}_{\bar{A}}$.¹ With these four fermionic d.o.f. we can construct *two* Dirac fermions, a particle–anti-particle pair, and four scalars, two particle–anti-particle pairs.

We use these two superfield ingredients to construct all the known fermions:

- To represent the Standard Model leptons we introduce the superfields l_i and $\bar{E}_i$ for the charged leptons (i is the generation index) and ν_i for the neutrinos, and form $SU(2)_L$ doublet superfield vectors $L_i = (\nu_i, l_i)$, which contain the left-handed Weyl spinors for the particles. We do not introduce the $\bar{N}_i$ that would contain right-handed neutrino spinors needed for massive Dirac neutrinos.² This is a convention – the MSSM is older than the discovery of neutrino mass – and including $\bar{N}_i$ fields has some interesting

¹The bar here is used to (not) confuse us, it is part of the name of the superfields and does not denote any hermitian or complex conjugate. The bar signifies that this is the field where, when hermitian conjugated, we will pick the right-handed (bared) Weyl-spinor from in the Dirac fermion.

²The anti-neutrino contained in the superfield $\nu_i^\dagger$ is right-handed consistent with experiment.

consequences.³ The right-handed neutrinos themselves would not couple to anything since they are SM singlets,⁴ however, their supersymmetric scalar partners could be a potential dark matter candidate.

- For quarks the situation is similar. Up-type and down-type quarks get the superfields u_i , $\bar{U}_i$ and d_i , $\bar{D}_i$, forming the $SU(2)_L$ doublets $Q_i = (u_i, d_i)$.⁵

With all possible apologies, we now change notation for the component fields to what is conventional in phenomenology, as opposed to pure theory. The left-handed Weyl spinors in these superfields will now be named for example l_{iL} for the spinor in the superfield l_i , and $\bar{l}_{iL}$ for the one in $\bar{E}_i$, and similarly for the other superfields. This signifies that they are the left-handed Weyl spinors for the particle and anti-particle, respectively. For the right-handed Weyl spinors in the hermitian conjugate superfields we swap the bar and L/R labels, so that for example $l_i^\dagger$ contains $\bar{l}_{iR}$ and $\bar{E}_i^\dagger$ contains l_{iR} .

The scalar component fields are named after the fermions using the tilde notation, for example in the superfield l_i we have the scalar $\tilde{l}_{iL}$ as the supersymmetric partner particle, often just called **sparticle**, of the fermion l_{iL} . Similarly, $\bar{E}_i$ contains $\tilde{l}_{iL}^*$ which is the anti-particle of $\tilde{l}_{iL}$, $l_i^\dagger$ contains $\tilde{l}_{iR}^*$, and $\bar{E}_i^\dagger$ contains $\tilde{l}_{iR}$. Since scalars do not have any notion of handedness the L or R here is just part of the conventional name, we still call these particles for left-handed and right-handed scalar leptons though.

Additionally, we need vector superfields, which, after the e.o.m. have eliminated the auxiliary field, contain a *massless* vector boson component field with two scalar d.o.f. and two Weyl-spinors, one of each handedness λ and $\bar{\lambda}$, with two fermionic degrees of freedom. Together these form a $j = \frac{1}{2}$ representation of the superalgebra. If the vector superfield is neutral, the fermions can form a Majorana fermion, if not they can be combined with the Weyl-spinors from other fields to form Dirac fermions.

Looking at the construction $V \equiv qt^a V^a$ for the vector superfields in the supersymmetric Lagrangian we see that, as expected, we need one superfield V^a per generator t^a of the algebra, with the normal $SU(3)_C$, $SU(2)_L$ and $U(1)_Y$ vector bosons as component fields. We call these superfields C^a , W^a and B^0 .⁶ In order to be really confusing, we use the following symbols for the fermions constructed from the respective Weyl-spinors: $\tilde{g}$, $\tilde{W}^0$ and $\tilde{B}^0$. The tilde here is supposed to tell us – just as for the scalar component fields of the scalar superfields – that they are supersymmetric partners of the known Standard Model particles.

We also need superfields for the scalar Higgs boson. Now life gets interesting. The Higgs $SU(2)_L$ doublet scalar field H used in the Standard Model cannot give mass to all fermions because it relies on the construction $H^C \equiv -i(H^\dagger \sigma_2)^T$ to give masses to up-type quarks, and possibly neutrinos. The superfield version of this cannot appear in the superpotential because it would mix left- and right-handed superfields. The minimal Higgs content we can get away with are two Higgs superfield $SU(2)_L$ doublets, which we will call H_u and H_d , indexing the quarks they give mass to.⁷ These must have (more on the reason for that in Sec. 6.2) weak

³Note that component fields in the same superfield must have the same charge under all the gauge groups, *i.e.* the scalar partner of the electron has electric charge $-e$, so it cannot be a neutrino.

⁴They can not be colour-charged, they are right-handed singlets under $SU(2)_L$ thus they have zero weak isospin, but since they should also have zero electric charge the hypercharge must also be zero.

⁵Here we should really also include a color index a so that u_i^a is a component in an $SU(3)_C$ triplet superfield vector. We omit these for simplicity.

⁶And there we have another W .

⁷In some further insanity some authors prefer H_1 and H_2 so that you have no idea which is which.

hypercharge $Y = \pm 1$ for H_u and H_d respectively, so that we have the superfield doublets:

$$H_u = \begin{pmatrix} H_u^+ \\ H_u^0 \end{pmatrix}, \quad H_d = \begin{pmatrix} H_d^0 \\ H_d^- \end{pmatrix}, \quad (6.1)$$

where we have given the electric charges of the scalar superfield components of the superfield doublets based on the standard $Q = \frac{1}{2}Y + T_3$ relationship after electroweak symmetry breaking, where T_3 is the weak isospin. The fermion component fields of these superfields will be denoted $\tilde{H}_u^+$, $\tilde{H}_u^0$, $\tilde{H}_d^0$, and $\tilde{H}_d^-$. The scalar component fields of these fields, before their mixing following electroweak symmetry breaking, will have the same symbols as the superfields. (Yes, really!)

6.2 The kinetic terms

It is now straight forward to write down the kinetic terms of the MSSM Lagrangian giving the matter-gauge interaction terms

$$\begin{aligned} \mathcal{L}_{\text{kin}} = & L_i^\dagger e^{\frac{1}{2}g\sigma W - \frac{1}{2}g'B} L_i + Q_i^\dagger e^{\frac{1}{2}g_s\lambda C + \frac{1}{2}g\sigma W + \frac{1}{3}\cdot\frac{1}{2}g'B} Q_i \\ & + \bar{U}_i^\dagger e^{\frac{1}{2}g_s\lambda C - \frac{4}{3}\cdot\frac{1}{2}g'B} \bar{U}_i + \bar{D}_i^\dagger e^{\frac{1}{2}g_s\lambda C + \frac{2}{3}\cdot\frac{1}{2}g'B} \bar{D}_i \\ & + \bar{E}_i^\dagger e^{2\frac{1}{2}g'B} \bar{E}_i + H_u^\dagger e^{\frac{1}{2}g\sigma W + \frac{1}{2}g'B} H_u + H_d^\dagger e^{\frac{1}{2}g\sigma W - \frac{1}{2}g'B} H_d, \end{aligned} \quad (6.2)$$

where g' , g and g_s are the couplings constants (strengths) of $U(1)_Y$, $SU(2)_L$, and $SU(3)_C$, respectively, and $\frac{1}{2}\sigma_i$ and $\frac{1}{2}\lambda_a$ are the generators of $SU(2)_L$ and $SU(3)_C$. As a convention we assign the charges under $U(1)$, hypercharge, in units of $\frac{1}{2}g'$. All non-singlets of $SU(2)_L$ and $SU(3)_C$ have the same charge, the factor $\frac{1}{2}$ here is used to get by without accumulation of numerical factors since the algebras for the Pauli (σ_i) and Gell-Mann matrices (λ_a) are:

$$[\frac{1}{2}\sigma_i, \frac{1}{2}\sigma_j] = i\epsilon_{ijk}\frac{1}{2}\sigma_k,$$

and

$$[\frac{1}{2}\lambda_a, \frac{1}{2}\lambda_b] = if_{ab}^c\frac{1}{2}\lambda_c.$$

These conventions lead to the wanted Standard Model gauge transformations for the component fields and the familiar relations after electroweak symmetry breaking, $e = g \sin \theta_W = g' \cos \theta_W$.

We mentioned earlier that the two Higgs superfields have opposite hypercharge. This is needed for so-called **anomaly cancellation** in the MSSM. Gauge anomaly is the possibility that at loop level contributions to processes such as in Fig. 6.1 break the gauge invariance that we have established at the classical level in the Lagrangian, and ruins the predictability of the theory. This rather miraculously does not happen in the SM because it has the exactly field content it has, so that all such possible gauge anomalies exactly cancel – we do not know of a deeper reason for why it has exactly this field content. If we have only *one* Higgs doublet this cancellation does not happen for the MSSM. With two Higgs doublets, and with opposite hypercharge, it does.



Figure 6.1: The tree level coupling between three gauge bosons B (left), and the one-loop fermion contribution to the same process (right).

6.3 Gauge terms

The pure gauge terms with supersymmetric field strengths are also fairly easy to write down:

$$\mathcal{L}_V = \bar{\theta}\bar{\theta} \text{Tr}[C^A C_A] + \bar{\theta}\bar{\theta} \text{Tr}[W^A W_A] + \frac{1}{2} \bar{\theta}\bar{\theta} B^A B_A, \quad (6.3)$$

where we have used the Dynkin indices of the gauge group representations

$$T(R)_L = \text{Tr} \left[\frac{1}{2} \sigma_1 \cdot \frac{1}{2} \sigma_1 \right] = \frac{1}{2},$$

and

$$T(R)_C = \text{Tr} \left[\frac{1}{2} \lambda_1 \cdot \frac{1}{2} \lambda_1 \right] = \frac{1}{2},$$

in the normalization of the terms, and where the field strengths are given as:

$$C_A = -\frac{1}{4} \bar{D} \bar{D} e^{-C} D_A e^C, \quad C = \frac{1}{2} \lambda_a C^a, \quad (6.4)$$

$$W_A = -\frac{1}{4} \bar{D} \bar{D} e^{-W} D_A e^W, \quad W = \frac{1}{2} \sigma_i W^i, \quad (6.5)$$

$$B_A = -\frac{1}{4} \bar{D} \bar{D} D_A B, \quad B = \frac{1}{2} B^0. \quad (6.6)$$

6.4 The MSSM superpotential

With the gauge structure of the Standard Model in place we are ready to write down all possible terms in the superpotential. First, we notice that there can be no **tadpole terms** (terms with only one superfield), since there are no superfields that are singlets (zero charge) under all Standard Model gauge groups. The only alternative would be if we introduced right-handed neutrino superfields $\bar{N}_i$.

For the possible mass terms $m\Phi_i\Phi_j$ we first check the abelian gauge group $U(1)_Y$, where the requirement reduces to the relatively simple sum of hypercharges $Y_i + Y_j = 0$. In Table 6.1 we see that the only possible contributions from the superfields that contain the Standard Model fermions are particle–anti-particle combinations such as $l_{iL}\bar{l}_{iR}$, but these come from superfields with different handedness and cannot be used together.

The exception is a mass term with the two Higgs superfields that have opposite hypercharge $Y = \pm 1$. In order to also be invariant under $SU(2)_L$ we have to write this superpotential term as⁸

$$W_{\text{mass}} = \mu H_u^T i\sigma_2 H_d, \quad (6.7)$$

⁸This is immediately invariant under $SU(3)_C$ since the Higgs fields have no colour charge.

Superfield	L_i	$\bar{E}_i^\dagger$	Q_i	$\bar{U}_i^\dagger$	$\bar{D}_i^\dagger$	$L_i^\dagger$	$\bar{E}_i$	$Q_i^\dagger$	$\bar{U}_i$	$\bar{D}_i$
Fermion	ν_{iL}, l_{iL}	l_{iR}	u_{iL}, d_{iL}	u_{iR}	d_{iR}	$\bar{\nu}_{iR}, \bar{l}_{iR}$	$\bar{l}_{iL}$	$\bar{u}_{iR}, \bar{d}_{iR}$	$\bar{u}_{iL}$	$\bar{d}_{iL}$
Hypercharge	-1	-2	$\frac{1}{3}$	$\frac{4}{3}$	$-\frac{2}{3}$	1	2	$-\frac{1}{3}$	$-\frac{4}{3}$	$\frac{2}{3}$

Table 6.1: The MSSM superfields with Standard Model fermion content and their hypercharges Y .

where μ is the superpotential mass parameter. This is invariant under $SU(2)_L$ because, with the supergauge transformations $H_d \rightarrow e^{-ig\frac{1}{2}\sigma_k W^k} H_d$ and $H_u^T \rightarrow H_u^T e^{-ig\frac{1}{2}\sigma_k^T W^k}$, we get

$$\begin{aligned} H_u^T i\sigma_2 H_d &\rightarrow H_u^T e^{ig\frac{1}{2}\sigma_k^T W^k} i\sigma_2 e^{ig\frac{1}{2}\sigma_k W^k} H_d \\ &= H_u^T i\sigma_2 e^{i\frac{1}{2}g\sigma_k W^k} e^{-i\frac{1}{2}g\sigma_k W^k} H_d = H_u^T i\sigma_2 H_d, \end{aligned}$$

since $\sigma_k^T \sigma_2 = -\sigma_2 \sigma_k$. Usually we ignore the $SU(2)_L$ specific structure and write terms like this as $\mu H_u H_d$, confusing the hell out of anyone that is not used to this convention since we really do mean Eq. (6.7). Notice that if we write (6.7) in terms of component fields we get

$$W_{\text{mass}} = H_u^T i\sigma_2 H_d = H_u^+ H_d^- - H_u^0 H_d^0,$$

which we should have been able to guess because the Lagrangian must also conserve electric charge.

If you have paid very close attention to the argument above you may have noticed that there is one more possibility, namely

$$W_{\text{mass}} = \mu'_i L_i H_u \equiv \mu'_i L_i^T i\sigma_2 H_u = \mu'_i (\nu_i H_u^0 - l_i H_u^+),$$

where μ' is some other mass parameter in the superpotential. This is clearly an allowable term (and we will return to it below), however, it also raises a very interesting question: Could we have $L_i \equiv H_d$? Could the lepton superfields L_i play the rôle of Higgs superfields, thus reducing the field content needed to describe the SM particles in a supersymmetric theory? While not immediately forbidden, this suggestions unfortunately leads to problems with anomaly cancelation, processes with large lepton flavour violation (LFV) and much too massive neutrinos, and has been abandoned.

We have now found all possible mass terms in the superpotential. What about the Yukawa terms? The hypercharge requirement here is $Y_i + Y_j + Y_k = 0$. From our table of hypercharges only the following terms are found to be viable:

$$W_{\text{Yukawa}} = y_{ij}^e L_i H_d E_j + y_{ij}^u Q_i H_u \bar{U}_j + y_{ij}^d Q_i H_d \bar{D}_j + \lambda_{ijk} L_i L_j \bar{E}_k + \lambda'_{ijk} L_i Q_j \bar{D}_k + \lambda''_{ijk} \bar{U}_i \bar{D}_j \bar{D}_k,$$

where we have named and indexed the Yukawa couplings in a hopefully natural way.⁹ For all these terms we can simultaneously keep $SU(2)_L$ invariance with the $i\sigma_2$ construction implicitly inserted between any superfield doublets. Note that because of the $SU(2)_L$ invariance, we must have $i \neq j$ for λ_{ijk} , since $i = j$ gives $L_i L_i \bar{E}_k = (\nu_i l_i - l_i \nu_i) \bar{E}_k = 0$.

⁹For some particular opinion of what is natural.

For $SU(3)_C$ to be conserved, we need to have colour singlets. Some of these terms are colour singlets by construction since they do not contain any coloured fields. The terms with only two quark superfields contain left-handed Weyl spinors for quarks and anti-quarks, which are $SU(3)_C$ singlets if the superfields come in colour-anti-colour pairs. In representation language they are in the $\mathbf{3}$ and $\bar{\mathbf{3}}$ representations of $SU(3)_C$. Written with colour indices explicit we have *e.g.* $L_i Q_j \bar{D}_k = L_i \sigma_2 Q_j^a \bar{D}_k^a$, where a is the colour index. The final term $\bar{U}_i \bar{D}_j \bar{D}_k$ is a colour singlet once we demand that it is totally anti-symmetric in the colour indices: $\bar{U}_i \bar{D}_j \bar{D}_k \equiv \epsilon_{abc} \bar{U}_i^a \bar{D}_j^b \bar{D}_k^c$. The anti-symmetry property of the Levi-Civita tensor ϵ_{abc} means that we must have $j \neq k$ in λ''_{ijk} .

Our complete superpotential is then:

$$\begin{aligned} W = & \mu H_u H_d + \mu'_i L_i H_u + y_{ij}^e L_i H_d E_j + y_{ij}^u Q_i H_u \bar{U}_j + y_{ij}^d Q_i H_d \bar{D}_j \\ & + \lambda_{ijk} L_i L_j \bar{E}_k + \lambda'_{ijk} L_i Q_j \bar{D}_k + \lambda''_{ijk} \bar{U}_i \bar{D}_j \bar{D}_k. \end{aligned} \quad (6.8)$$

The parameter μ , potentially complex, is a brand new supersymmetric parameter appearing in the superpotential, with no corresponding parameter existing in the Standard Model Lagrangian. However, the Yukawa couplings y_{ij} are identical to the Standard Model Yukawa couplings since they will be required to give mass to the Standard Model fermions after electroweak symmetry breaking when the Higgs fields get a vev, see Section 6.9. The fate of the other parameters will be discussed in the next section.

6.5 R-parity

The superpotential terms LH_u , LLE and $LQ\bar{D}$ that we have written down in (6.8) all violate lepton number conservation, and $\bar{U}\bar{D}\bar{D}$ violates baryon number conservation. Such terms do not exist in the Standard Model, although there is no symmetry forbidding their existence there, instead a combination of what fields exist and the gauge symmetries that limit their interactions means that there are no such tree-level interactions. We call this an **accidental symmetry** of the Standard Model.

Allowing such terms in the MSSM would lead to, among other phenomenological problems, processes like proton decay, for example through $p \rightarrow e^+ \pi^0$ as shown in Fig. 6.2, which breaks both baryon number and lepton number. We can estimate the resulting proton life-time by noting that the exchange of a scalar particle (in this example a strange squark $\tilde{s}$) creates an effective dimension-6 four-fermion interaction Lagrangian term $\lambda \bar{u} \bar{d} e u$ with coupling

$$\lambda = \frac{\lambda'_{112} \lambda''_{112}}{m_{\tilde{s}}^2}, \quad (6.9)$$

where the sparticle mass $m_{\tilde{s}}$ comes from the scalar propagator in the diagram. The resulting matrix element for the total proton decay process must then be proportional to $|\lambda|^2$, which has mass dimension M^{-4} . Since decay width has mass dimension M , the phase space part of the calculation must provide something of mass dimension M^5 . The only mass scale involved in the problem is the proton mass m_p , thus we approximate the phase space integration part of the proton decay width by m_p^5 . We then have

$$\Gamma_{p \rightarrow e^+ \pi^0} \sim |\lambda|^2 m_p^5 = |\lambda'_{112} \lambda''_{112}|^2 \frac{m_p^5}{m_{\tilde{s}}^4}. \quad (6.10)$$

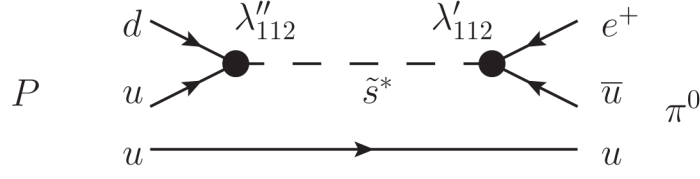


Figure 6.2: Possible Feynman diagram for proton decay with R-parity violating couplings λ''_{112} and λ'_{112} .

The current measured lower limit on the lifetime from watching a lot of protons not decay is $\tau_{p \rightarrow e^+ \pi^0} > 1.6 \cdot 10^{34}$ y [11], or, converting to seconds, $\tau_{p \rightarrow e^+ \pi^0} > \pi \cdot 10^7$ s/y $\times 1.6 \cdot 10^{34}$ y = $5.0 \cdot 10^{41}$ s, which gives a limit on the width

$$\Gamma_{p \rightarrow e^+ \pi^0} = \frac{\hbar}{\tau} < \frac{6.582 \cdot 10^{-25} \text{ GeV s}}{5.0 \cdot 10^{41} \text{ s}} \simeq 1.3 \cdot 10^{-66} \text{ GeV},$$

so that we have the following very strict limit on the combination of the two couplings

$$|\lambda'_{112} \lambda''_{112}| < 1.3 \cdot 10^{-27} \left(\frac{m_{\tilde{s}}}{1 \text{ TeV}} \right)^2. \quad (6.11)$$

Such strict limits can be found on most of the lepton and baryon number couplings using the measured properties of the Standard Model particles, with some exceptions for coupling involving the second and third generation fermions.

To avoid all such lepton and baryon number couplings Fayet (1975) [12] introduced the conservation of R-parity.

Definition: R-parity is a multiplicatively conserved quantum number given by

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

where s is a particle's spin, B its baryon number and L its lepton number.

For all Standard Model particles, including all the scalar Higgs bosons, this gives $R = 1$, while the superpartners all have $R = -1$. One usually *defines* the MSSM as conserving R-parity. For the MSSM this excludes the terms LH_u , $LL\bar{E}$, $LQ\bar{D}$ and $\bar{U}\bar{D}\bar{D}$ from the superpotential,¹⁰ leaving us with the R-parity conserving superpotential

$$W = \mu H_u H_d + y_{ij}^e L_i H_d E_j + y_{ij}^u Q_i H_u \bar{U}_j + y_{ij}^d Q_i H_d \bar{D}_j. \quad (6.12)$$

The consequence of this somewhat *ad hoc* definition is that in all interactions the total number of incoming and outgoing supersymmetric particles must be an even number $2n$, so that the total R-number is $(-1)^{2n} = 1$. This leads to the following very important phenomenological consequences:

¹⁰All the superpotential Yukawa terms lead to component field terms of the form $A_i \psi_j \psi_k$. If the scalar A_i here is *not* a Higgs boson, then it is a superpartner and if none of the fermions come from a Higgs superfield so that they are also a superpartner the term breaks R-parity conservation. The superpotential mass terms have component field terms of the form $\psi_i \psi_j$. If one of the fermions here comes from a Higgs superfield, then it is a superpartner, and if the other does not, the term breaks R-parity.

1. All sparticles must be produced in pairs (ignoring the very low probability of producing four or more).
2. Sparticles must annihilate in pairs.
3. The **lightest supersymmetric particle** (LSP) is absolutely stable, and every other sparticle must decay down to the LSP (possibly in multiple steps).

6.6 Supersymmetry breaking terms in the MSSM

We can directly apply our previous arguments on gauge invariance, that we used when discussing the superpotential, on the general soft-breaking terms in Eq. (5.7) in order to determine which supersymmetry breaking terms are allowed in the MSSM, keeping also in mind the requirement of R-parity conservation.

Mass terms of the form

$$-\frac{1}{4T(R)}M\theta\theta\bar{\theta}\bar{\theta}\text{Tr}[W^AW_A] + \text{h.c.},$$

are allowed because they have the same gauge structure as the supersymmetric field strength terms. In component fields only the fermions in the vector superfields survive, and are for the MSSM:

$$\mathcal{L}_{\text{soft}} = -\frac{1}{2}M_1\tilde{B}\tilde{B} - \frac{1}{2}M_2\tilde{W}^i\tilde{W}^i - \frac{1}{2}M_3\tilde{g}^a\tilde{g}^a + \text{c.c.},$$

where the M_i are potentially complex-valued. This gives six new parameters.

Yukawa terms

$$-\frac{1}{6}a_{ijk}\theta\theta\bar{\theta}\bar{\theta}\Phi_i\Phi_j\Phi_k + \text{h.c.},$$

are allowed when a corresponding term exist in the superpotential – meaning they are gauge invariant. In component fields only the scalar parts of the superfields survive, and the allowed terms are

$$\mathcal{L}_{\text{soft}} = -a_{ij}^e\tilde{L}_iH_d\tilde{l}_{jR}^* - a_{ij}^u\tilde{Q}_iH_u\tilde{u}_{jR}^* - a_{ij}^d\tilde{Q}_iH_d\tilde{d}_{jR}^* + \text{c.c.},$$

where we remind you that the H here refers to scalar parts of the Higgs superfields. We can see that all of these terms are R-parity conserving, since they consist of two sparticles and one (Higgs) particle. The couplings a_{ij} are all potentially complex valued, so this gives us 54 new parameters.

The mass terms

$$-\frac{1}{2}b_{ij}\theta\theta\bar{\theta}\bar{\theta}\Phi_i\Phi_j + \text{h.c.},$$

are again only allowed for corresponding terms in the superpotential, *i.e.*

$$\mathcal{L}_{\text{soft}} = -bH_uH_d + \text{c.c.},$$

where b is potentially complex valued, which gives us 2 new parameters.¹¹ Tadpole terms are not allowed, as there are no tadpoles in the superpotential.

Mass terms

$$-m_{ij}^2\theta\theta\bar{\theta}\bar{\theta}\Phi_i^\dagger\Phi_j,$$

¹¹The coupling b is sometimes written $B\mu$ where B is a unitless constant that indicates how different the coupling is from the corresponding coupling in the superpotential.

are allowed because they have the same gauge structure as the supersymmetric kinetic terms. In component fields again only the scalar fields survive, and in the MSSM they are:

$$\begin{aligned} \mathcal{L}_{\text{soft}} = & -(m_{ij}^L)^2 \tilde{L}_i^\dagger \tilde{L}_j - (m_{ij}^e)^2 \tilde{l}_{iR}^* \tilde{l}_{jR} - (m_{ij}^Q)^2 \tilde{Q}_i^\dagger \tilde{Q}_j - (m_{ij}^u)^2 \tilde{u}_{iR}^* \tilde{u}_{jR} - (m_{ij}^d)^2 \tilde{d}_{iR}^* \tilde{d}_{jR} \\ & - m_{H_u}^2 H_u^\dagger H_u - m_{H_d}^2 H_d^\dagger H_d, \end{aligned} \quad (6.13)$$

where the m_{ij}^2 are complex valued, however, also hermetic. This gives rise to 47 new parameters. Despite technically being allowed the MSSM ignores the “maybe-soft” terms in Eq. (5.8).

In total, after using our freedom to choose our basis for the fields wisely in order to remove what freedom we can, the MSSM has 105 new parameters compared to the Standard Model, 104 of these are soft-breaking terms and μ is the only new parameter in the superpotential.

6.7 Renormalisation group equations

Renormalisation, the removal of infinities from field theory predictions, introduces a fixed scale μ at which the parameters of the Lagrangian, the couplings, are defined. For example, the charge of the electron is not simply the bare charge e , but a charge at a given energy scale μ , $e(\mu)$, which is the scale at which the theory describes the electron, and which we can measure in an experiment at that scale. Scattering an electron at very high energy will require a different value of $e(\mu)$ than at a low energy. This is an experimentally well verified fact.¹²

However, since μ is not an observable *per se* but in principle a choice of how to write down the theory (at which energy to write down the Lagrangian), the action S should be invariant under a change of μ , which is expressed as:

$$\mu \frac{d}{d\mu} S(\Phi, \lambda, \mu) = 0, \quad (6.14)$$

where λ is a generic name for the couplings of the theory and Φ represents the (super)fields that have been renormalised.¹³ This equation can be re-written in terms of partial derivatives

$$\left(\mu \frac{\partial}{\partial \mu} + \mu \frac{\partial \lambda}{\partial \mu} \frac{\partial}{\partial \lambda} \right) S(\Phi, \lambda, \mu) = 0, \quad (6.15)$$

which is the **renormalisation group equation** (RGE).

We can look at the behaviour of a Lagrangian parameter λ as a function of the energy scale μ away from the value where it was defined, often denoted μ_0 . This is controlled by the **β -function**:

$$\beta_\lambda \equiv \mu \frac{\partial \lambda}{\partial \mu}. \quad (6.16)$$

These β -functions can be found from the so-called counterterm Z that renormalises the parameter.

¹²It is also impossible to avoid if we accept that the electron is a point particle. Since the potential has the form $V(r) \propto e/r$ an infinite energy would appear unless we were somehow to modify the charge at high energies, or equivalently, short distances.

¹³In Sec. 4.8 we mentioned how the non-renormalisation theorem implies that we do not need to renormalise the coupling constants of the superpotential separately.

As an example, take a **gauge coupling constant** g_0 defined (taken from measurement) at some scale μ_0 . This is given by (in $d = 4 - \epsilon$ dimensions):¹⁴

$$g_0 = Zg\mu^{-\epsilon/2}.$$

Then, differentiating both sides with respect to μ ,

$$\begin{aligned} 0 &= \frac{\partial Z}{\partial \mu} g\mu^{-\epsilon/2} + Z \frac{\partial g}{\partial \mu} \mu^{-\epsilon/2} - \frac{\epsilon}{2} Zg\mu^{-\epsilon/2-1} \\ \mu \frac{\partial g}{\partial \mu} &= \frac{\epsilon}{2} g - \frac{g\mu}{Z} \frac{\partial Z}{\partial \mu} \\ \mu \frac{\partial g}{\partial \mu} &= \frac{\epsilon}{2} g - g\mu \frac{\partial}{\partial \mu} \ln Z, \end{aligned}$$

and taking the limit $\epsilon \rightarrow 0$:

$$\beta_g = \mu \frac{\partial g}{\partial \mu} = -g\gamma_g,$$

where we have defined the **anomalous dimension** of g

$$\gamma_g = \mu \frac{\partial}{\partial \mu} \ln Z. \quad (6.17)$$

The reason for keeping around the factor of μ in the definition of the β -function is that it typically changes very slowly over large differences in energy scale, so it is practical to change variable to $t = \ln \frac{\mu}{\mu_0}$, so that $\mu \frac{\partial}{\partial \mu} = \frac{\partial}{\partial t}$ and $\beta_g = \frac{\partial g}{\partial t}$.

Z can now be calculated to the required loop-order to find the β -function to that order, and in turn the running of the coupling constant with μ . By evaluating to one-loop order we can find that for our particular example of a gauge coupling constant for a generic supersymmetric model

$$\gamma_g|_{1\text{-loop}} = -\frac{1}{16\pi^2} g^2 \left(\sum_R T(R) - 3C(A) \right), \quad (6.18)$$

where the sum of Dynkin indices is over all superfields that transform under a representation R of the gauge group in question, and $C(A)$ is the **quadratic Casimir invariant** of the representation A of the vector field under the gauge group

$$C(A)\delta_{ij} = (T^a T^a)_{ij}.$$

For the adjoint representation of $U(1)$ this is 0, and for $SU(N)$ this is N . The running of the coupling constants is particularly important since it will later lead us to the concept of gauge coupling unification.

As a second relevant example, for the soft-breaking parameters M_i in (5.7) we have the one-loop β -functions

$$\beta_{M_i} \equiv \frac{d}{dt} M_i = \frac{1}{16\pi^2} g_i^2 M_i \left(2 \sum_R T(R) - 6C(A) \right). \quad (6.19)$$

¹⁴The factor $\mu^{-\epsilon/2}$ is there to ensure that the scale of g_0 is correct.

6.8 Gauge coupling unification

The one-loop β -functions for gauge couplings in a generic supersymmetric model were given in Eq. (6.18). With the MSSM field content and the gauge couplings discussed in this chapter:¹⁵

$$g_1 = \sqrt{\frac{5}{3}}g', \quad g_2 = g, \quad g_3 = g_s,$$

we arrive at

$$\beta_{g_i}|_{1\text{-loop}} = \frac{1}{16\pi^2} b_i g_i^3, \quad (6.20)$$

with in the MSSM

$$b_i^{\text{MSSM}} = \left(\frac{33}{5}, 1, -3 \right).$$

For comparison, the same result in the Standard Model is

$$b_i^{\text{SM}} = \left(\frac{41}{10}, -\frac{19}{6}, -7 \right).$$

The values of b_i for the MSSM are found from the Casimir invariant and the Dynkin index of the gauge group representations

$$C(A)_{SU(3)} = 3, \quad C(A)_{SU(2)} = 2, \quad C(A)_{U(1)} = 0,$$

and

$$T(R)_{SU(3)} = \frac{1}{2}, \quad T(R)_{SU(2)} = \frac{1}{2}, \quad T(R)_{U(1)} = \frac{3}{5}Y^2,$$

where for example $b_3 = \frac{1}{2} \cdot 12 - 3 \cdot 3 = -3$ in the MSSM, because, after careful counting, we have twelve quark/squark scalar superfields transforming under $SU(3)_C$.

At one-loop order we can do a neat rewrite using $\alpha_i \equiv \frac{g_i^2}{4\pi}$. Since

$$\frac{d}{dt} \alpha_i^{-1} = -2 \frac{4\pi}{g_i^3} \frac{d}{dt} g_i,$$

we have:

$$\beta_{\alpha_i^{-1}} \equiv \frac{d}{dt} \alpha_i^{-1} = -\frac{8\pi}{g_i^3} \frac{1}{16\pi^2} g_i^3 b_i = -\frac{b_i}{2\pi}.$$

Thus α^{-1} runs linearly with t at one loop.

By running the couplings α_i^{-1} from their values measured at the electroweak scale to high energies it is observed that in the MSSM the coupling constants intersect at a single point, which they do not naturally do in the Standard Model. See Fig. 6.3, taken from Martin [13]. The common assumption made is then that a unified gauge group, *e.g.* $SU(5)$ or $SO(10)$, is broken at that scale, called the **grand unification theory** scale or GUT-scale, down to the Standard Model gauge group. This scale is $\mu_{\text{GUT}} \approx 2 \cdot 10^{16} \text{ GeV}$, about two orders of magnitude below the Planck scale.

¹⁵The normalisation choice for g_1 may seem a bit strange, however, this is the correct numerical factor when for example breaking a unified group such as $SU(5)$ or $SO(10)$ down to the Standard Model gauge group. This factor might be different with a different unified group.

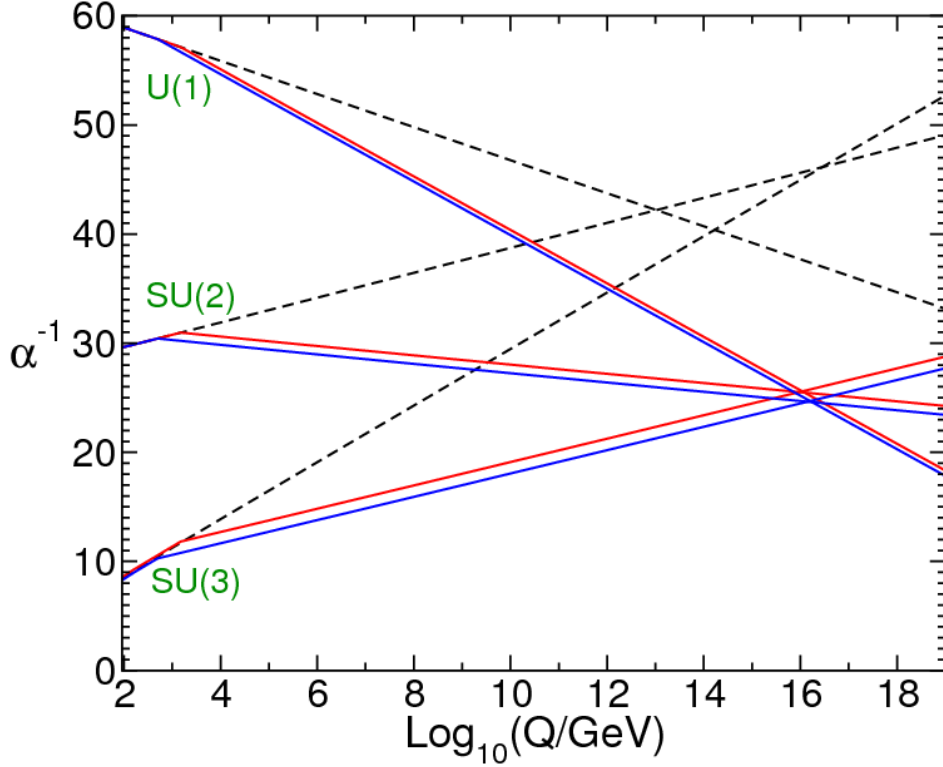


Figure 6.3: The RGE evolution of the inverse gauge couplings $\alpha_i^{-1}(Q)$ in the Standard Model (dashed lines) and the MSSM (solid lines). In the MSSM case, the sparticle mass thresholds are varied between 250 GeV and 1 TeV and $\alpha_3(m_Z)$ between 0.113 and 0.123 to create the bands shown by the red and blue lines. Two-loop effects are included.

Something funny happens to the gaugino soft-mass parameters M_i if we look at their running. From (6.19) the one-loop β -functions for the M_i in the MSSM are

$$\beta_{M_i}|_{1\text{-loop}} \equiv \frac{d}{dt} M_i = \frac{1}{8\pi^2} g_i^2 M_i b_i. \quad (6.21)$$

As a consequence all three ratios M_i/g_i^2 are scale independent at one loop. To see this let $R = M_i/g_i^2$, then

$$\beta_R \equiv \frac{dR}{dt} = \frac{\frac{d}{dt} M_i g_i^2 - M_i \frac{d}{dt} g_i^2}{g_i^4} = \frac{\frac{1}{8\pi^2} g_i^2 M_i b_i \cdot g_i^2 - M_i \cdot 2g_i \cdot \frac{1}{16\pi} g_i^3 b_i}{g_i^4} = 0. \quad (6.22)$$

In other words, R does not change with scale t .

If we now use that the coupling constants unify at the GUT scale to the coupling g_u , and assume that the soft-masses are the same at that scale $m_{1/2} = M_1(\mu_{\text{GUT}}) = M_2(\mu_{\text{GUT}}) = M_3(\mu_{\text{GUT}})$,¹⁶ it follows that

$$\frac{M_1}{g_1^2} = \frac{M_2}{g_2^2} = \frac{M_3}{g_3^2} = \frac{m_{1/2}}{g_u^2}, \quad (6.23)$$

at all scales!¹⁷ This is a very powerful and predictive assumption. Because of the relationship

¹⁶Again, not unreasonable if the spontaneous symmetry breaking mechanism acts uniformly for all the gauginos.

¹⁷At one-loop level.

between the electroweak couplings and the electric charge, $e = g' \cos \theta_W = g \sin \theta_W$, it leads to the following relation

$$M_3 = \frac{\alpha_s}{\alpha} \sin^2 \theta_W M_2 = \frac{3}{5} \frac{\alpha_s}{\alpha} \cos^2 \theta_W M_1, \quad (6.24)$$

which, inserting values for the fine structure constant, the strong coupling, and the Weinberg angle, numerically predicts

$$M_3 : M_2 : M_1 \simeq 6 : 2 : 1$$

at the electroweak scale. We will return to the implications of this when discussing the gauginos in Sec. 6.12.

In Fig. 6.4, again taken from Martin [13], we show the running of the gaugino mass parameters M_i (solid black), the Higgs mass parameters $m_{H_{d/u}}^2$ (dot-dashed green), the third generation sfermion soft terms m_{d_3} , m_{Q_3} , m_{u_3} , m_{L_3} and m_{e_3} (dashed red and blue, listed from top to bottom), and the corresponding first and second generation terms (solid lines). Here we assume unified soft-mass parameters $m_{1/2}$ for the gauginos and m_0 for the Higgs and sfermions at the GUT scale.

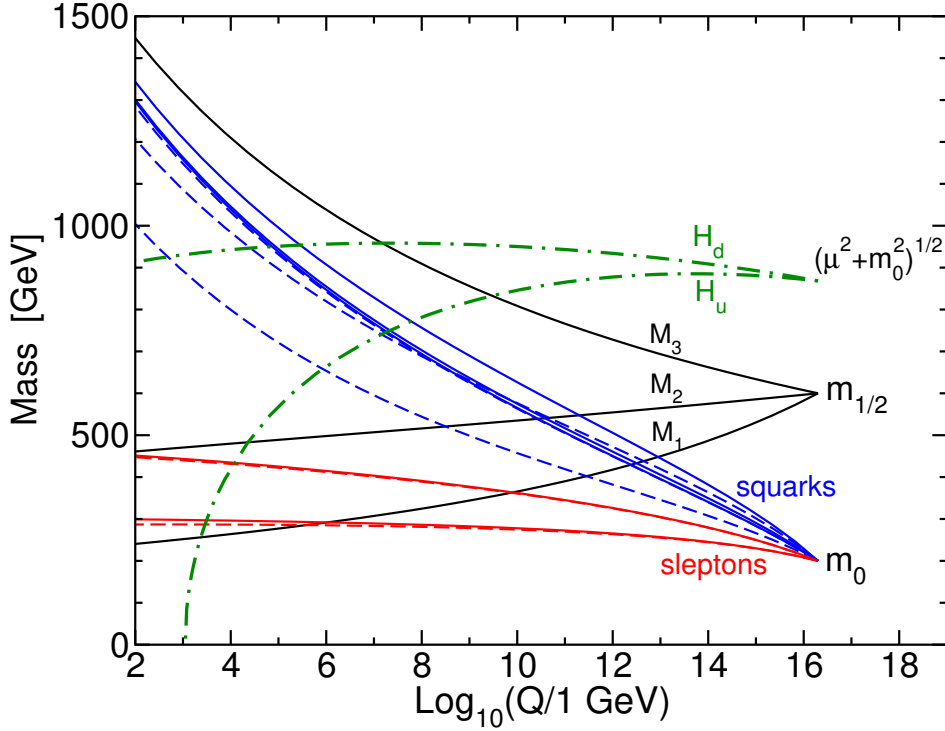


Figure 6.4: The RGE evolution of soft-mass parameters in the MSSM with typical minimal supergravity-inspired boundary conditions imposed at $2 \cdot 10^{16}$ GeV. The parameter values used for this illustration were $m_0 = 200$ GeV, $m_{1/2} = -A_0 = 600$ GeV, $\tan \beta = 10$, and $\text{sgn}(\mu) = +$.

6.9 Radiative Electroweak Symmetry Breaking

In the Standard Model both the vector bosons and the fermions are given mass spontaneously by electroweak symmetry breaking (EWSB), which is induced by the shape of the scalar

potential for the Higgs field H , which is an $SU(2)_L$ doublet of scalar fields:

$$V(H) = \mu^2 |H|^2 + \lambda |H|^4, \quad (6.25)$$

with $|H|^2 = H^\dagger H$.

Here, the requirement for successful EWSB is that $\lambda > 0$ and $\mu^2 < 0$.¹⁸ The first of these requirements ensures that the potential is **bounded from below**, *i.e.* that in the limit of large field values the potential does not turn to negative infinity. The second ensures that the minimum of the potential, the vacuum, is not given by zero field values, *i.e.* that the fields have **vacuum expectation values** (vevs).

In supersymmetry we found the following general scalar potential in Eq. (4.21) for unbroken supersymmetry,

$$V(A, A^*) = \sum_i \left| \frac{\partial W}{\partial A_i} \right|^2 + \frac{1}{2} \sum_a g^2 (A^* T^a A)^2 > 0, \quad (6.26)$$

where the first part is due to the elimination of the auxiliary F -fields in the scalar superfields, while the second part is due to the elimination of the auxiliary D -fields in the vector superfields. In addition we have to add all terms containing only relevant scalar fields from the soft breaking terms in Eq. 5.9.

For the scalar Higgs component fields in the MSSM this gives the potential

$$\begin{aligned} V(H_u, H_d) &= |\mu|^2 (|H_u^0|^2 + |H_u^+|^2 + |H_d^0|^2 + |H_d^-|^2) && \text{(from } F\text{-terms)} \\ &+ \frac{1}{8} (g^2 + g'^2) (|H_u^0|^2 + |H_u^+|^2 - |H_d^0|^2 - |H_d^-|^2)^2 && \text{(from } D\text{-terms)} \\ &+ \frac{1}{2} g^2 |H_u^+ H_d^{0*} + H_u^0 H_d^{-*}|^2 \\ &+ m_{H_u}^2 (|H_u^0|^2 + |H_u^+|^2) + m_{H_d}^2 (|H_d^0|^2 + |H_d^-|^2) && \text{(from soft breaking terms)} \\ &+ [b(H_u^+ H_d^- - H_u^0 H_d^0) + \text{c.c.}] && (6.27) \end{aligned}$$

This potential has 8 d.o.f. from 4 complex scalar fields H_u^+ , H_u^0 , H_d^0 and H_d^- . At the same time it has 6 parameters: the two Standard Model couplings g and g' , the magnitude of the supersymmetric parameter μ , and the three soft-breaking parameters b , $m_{H_u}^2$ and $m_{H_d}^2$.

We now want to do as in the Standard Model and break $SU(2)_L \times U(1)_Y \rightarrow U(1)_{\text{em}}$ in order to give masses to gauge bosons and SM fermions.¹⁹ To do this we need to show that (6.27) has: i) a minimum for finite, *i.e.* non-zero, field values, ii) that this minimum has a remaining $U(1)_{\text{em}}$ symmetry and iii) that the potential is bounded from below. We will here restrict our analysis to tree level, ignoring loop effects on the potential.

We start by using our $SU(2)_L$ gauge freedom, picking a gauge so that we rotate away any field value for H_u^+ at the minimum of the potential. So without loss of generality we can use that $H_u^+ = 0$ at the minimum in what follows. At the minimum we must have $\partial V / \partial H_u^+ = 0$, and by explicit differentiation of the potential one can show that $H_u^+ = 0$ then also leads to $H_d^- = 0$. This is good and proper since it guarantees our item ii), that $U(1)_{\text{em}}$ is a symmetry for the minimum of the potential, since the charged fields then have no vevs.

¹⁸This is called the **Mexican hat** or **wine bottle potential**, depending on preferences.

¹⁹You may ask why we can not use the soft-terms from the spontaneous breaking of *supersymmetry* to do this. However, the soft-terms are unable to effectively provide masses to vector bosons and fermions because they deal (mostly) with scalar fields.

We are now left with the following potential only in terms of the uncharged Higgs fields H_u^0 and H_d^0 (after gauge choice and at the minimum):

$$\begin{aligned} V(H_u^0, H_d^0) = & (|\mu|^2 + m_{H_u}^2)|H_u^0|^2 + (|\mu|^2 + m_{H_d}^2)|H_d^0|^2 \\ & + \frac{1}{8}(g^2 + g'^2)(|H_u^0|^2 - |H_d^0|^2)^2 - (bH_u^0H_d^0 + \text{c.c.}) \end{aligned} \quad (6.28)$$

We can now absorb a complex phase in H_u^0 or H_d^0 , in order to take b to be real and positive. This does not affect other terms because they are protected by absolute values. The minimum must also have the total phase of $H_u^0H_d^0$ real and positive, to get an as large as possible negative contribution from the b term. Thus the vevs $v_u = \langle H_u^0 \rangle$ and $v_d = \langle H_d^0 \rangle$ must have opposite phases. By the remaining $U(1)_Y$ gauge symmetry of the potential, which is effectively a phase rotation, and the fact that H_u^0 and H_d^0 have opposite hypercharge, we can transform v_u and v_d so that they are real and have the same sign. For the potential to have a negative mass term, and thus fulfill point i) above, we must then have

$$b^2 > (|\mu|^2 + m_{H_u}^2)(|\mu|^2 + m_{H_d}^2). \quad (6.29)$$

Since we have broken supersymmetry we must also check that the potential is actually bounded from below, our point iii), which was guaranteed for a supersymmetric vacuum. For large $|H_u^0|$ or $|H_d^0|$ the quartic gauge terms in (6.28) blows up to save the potential, except for $|H_u^0| = |H_d^0|$, the so-called D -flat directions. This means that we must also require

$$2b < 2|\mu|^2 + m_{H_u}^2 + m_{H_d}^2. \quad (6.30)$$

To summarise what we have learnt so far: at the minimum of the Higgs potential we know that the expectation values of the charged Higgs component fields are zero, $\langle H_u^+ \rangle = 0$ and $\langle H_d^- \rangle = 0$, and we fulfil the condition for the existence of an extremal point in the neutral Higgs component fields

$$\frac{\partial V}{\partial H_u^0} = \frac{\partial V}{\partial H_d^0} = 0. \quad (6.31)$$

In addition, for the minimum to have non-zero field values that break EWSB, and for the potential to be bounded from below, the parameters of the potential must simultaneously fulfil the inequalities

$$\begin{aligned} b^2 & > (|\mu|^2 + m_{H_u}^2)(|\mu|^2 + m_{H_d}^2), \\ 2b & < 2|\mu|^2 + m_{H_u}^2 + m_{H_d}^2. \end{aligned}$$

The resulting non-zero expectation values at the minimum for the neutral Higgs component fields are denoted v_u and v_d .

To satisfy (6.29) and (6.30), a negative value for $m_{H_u}^2$ (or $m_{H_d}^2$) can help, in particular if $|\mu|^2 + m_{H_u}^2 < 0$, as that automatically fulfils (6.29). Such a negative value is indeed perfectly allowed as a parameter in the Lagrangian. No negative particle masses will result.

In fact, if we assume that $m_{H_d} = m_{H_u}$ at some scale μ , for example the supersymmetry breaking scale, then (6.29) and (6.30) cannot be simultaneously be satisfied *at that scale*. However, to 1-loop the RGE running of these mass parameters is:

$$16\pi^2 \beta_{m_{H_u}^2} \equiv 16\pi^2 \frac{dm_{H_u}^2}{dt} = 6|y_t|^2(m_{H_u}^2 + m_{Q_3}^2 + m_{u_3}^2) + \dots$$

$$16\pi^2\beta m_{H_d}^2 \equiv 16\pi^2 \frac{dm_{H_d}^2}{dt} = 6|y_b|^2(m_{H_d}^2 + m_{Q_3}^2 + m_{d_3}^2) + \dots,$$

where y_t and y_b are the top and bottom quark Yukawa couplings, and $m_{Q_3} = m_{33}^Q$, $m_{u_3} = m_{33}^u$, and $m_{d_3} = m_{33}^d$, in a simplification of our previous notation for the soft-masses in Sec. 6.6. Because $y_t \gg y_b$, $m_{H_u}^2$ runs much faster with scale than $m_{H_d}^2$. If the parameters start out the same at some high scale, say from some universal supersymmetry breaking effect, as we go down to the electroweak scale $m_{H_u}^2$ becomes significantly smaller than $m_{H_d}^2$, and may become negative, fulfilling the EWSB criteria. For an illustration, see Fig. 6.4, where the running of $\mu^2 + m_{H_u}^2$ and $\mu^2 + m_{H_d}^2$ is shown. It is this property of starting out from some universal Higgs parameters at a high scale, which then by RGE effects break electroweak symmetry at lower scales, that is termed **radiative EWSB** (REWSB). Thus, in the MSSM with universal soft terms at a high scale there is an explanation why EWSB happens, it is not put in by hand in the potential as it is in the Standard Model.

Following EWSB, to get the familiar vector boson masses measured by experiment, the vevs need to satisfy the constraint from the electroweak parameters:

$$v_u^2 + v_d^2 \equiv v^2 = \frac{2m_Z^2}{g^2 + g'^2} \approx (174 \text{ GeV})^2. \quad (6.32)$$

Thus we have one free parameter coming from the two Higgs vevs in the MSSM. We can write this as

$$\tan \beta \equiv \frac{v_u}{v_d},$$

where by convention $0 < \beta < \pi/2$, so that $0 < \tan \beta < \infty$.

Using the condition for the existence of an extremal point in (6.31), the two non-SM parameters b and $|\mu|$ can be eliminated as free parameters from the model, however, not the sign of μ . Alternatively, we can choose to eliminate $m_{H_u}^2$ and $m_{H_d}^2$. You can look at this as giving away the freedom of these parameters to the vevs, and then fixing one vev by the electroweak constraint, and using $\tan \beta$ for the other.

Let us make a little remark here on the parameter μ . We have what is called the **μ problem**. The soft terms all get their scale from some common mechanism at some common high energy scale, it is assumed, however, μ is a mass term in the superpotential (the only one in fact) and could *a priori* take *any* value, even M_P . Why is μ then of the order of the soft terms, which is what allows us to achieve REWSB?²⁰

6.10 Higgs boson properties

Of the eight d.o.f. in the scalar potential for the Higgs component fields three are Goldstone bosons that get eaten by Z and $W^\pm$ to give them masses. The remaining five d.o.f. form two neutral scalars h , H , two charged scalars $H^\pm$ and one neutral pseudo-scalar (CP-odd) A .²¹

²⁰This problem can be solved in extensions of the MSSM such as the Next-to-Minimal Supersymmetric Standard Model (NMSSM).

²¹In addition to the scalars, we know that the Higgs supermultiplets contain four fermions, $\tilde{H}_u^0$, $\tilde{H}_d^0$, $\tilde{H}_u^+$ and $\tilde{H}_d^-$ (higgsinos). We will see later that these mix with the fermion partners of the gauge bosons (gauginos).

At tree level one can show that these have the masses:

$$m_A^2 = \frac{2b}{\sin 2\beta} = 2|\mu|^2 + m_{H_u}^2 + m_{H_d}^2, \quad (6.33)$$

$$m_{h,H}^2 = \frac{1}{2} \left(m_A^2 + m_Z^2 \mp \sqrt{(m_A^2 - m_Z^2)^2 + 4m_Z^2 m_A^2 \sin^2 2\beta} \right), \quad (6.34)$$

$$m_{H^\pm}^2 = m_A^2 + m_W^2. \quad (6.35)$$

As a consequence m_A and $\tan \beta$ can be used to parametrise the masses of the Higgs sector (at tree level), and H , $H^\pm$ and A are in principle unbounded in mass since they grow as $b/\sin 2\beta$. However, at tree level the lightest Higgs boson is restricted to

$$m_h < m_Z |\cos 2\beta| < 91.2 \text{ GeV}. \quad (6.36)$$

In contrast we have the current best measurement of the Higgs boson mass of $m_h = 125.10 \pm 0.14 \text{ GeV}$, combining results from the LHC [14].

Fortunately, there are large loop-corrections or the MSSM would have been excluded already.²² Because of the size of the Yukawa couplings the largest corrections to the mass of the lightest Higgs comes from loops with top quarks and its supersymmetric partners, the scalar top quarks, or stops, $\tilde{t}_L$ and $\tilde{t}_R$. See Fig. 4.1 for the relevant Feynman diagrams. In the limit where the mass of the stop quarks are larger than the top, $m_{\tilde{t}_R}, m_{\tilde{t}_L} \gg m_t$, and with stop mass eigenstates close to the chiral eigenstates (more on this later), we get the dominant loop correction

$$\Delta m_h^2 = \frac{3}{4\pi^2} \cos^2 \alpha y_t^2 m_t^2 \ln \left(\frac{m_{\tilde{t}_L} m_{\tilde{t}_R}}{m_t^2} \right), \quad (6.37)$$

where α is a mixing angle for h and H with respect to the superfield component fields H_u^0 and H_d^0 , given by

$$\frac{\sin \alpha}{\sin \beta} = -\frac{m_H^2 + m_h^2}{m_H^2 - m_h^2}, \quad (6.38)$$

at tree level.

With this and other corrections the upper bound on the lightest Higgs boson mass is weaker:

$$m_h \leq 135 \text{ GeV},$$

assuming a common sparticle mass scale of around $m_{\text{SUSY}} \leq 1 \text{ TeV}$. Higher values for the sparticle masses give large fine-tuning and weaken the bound very little because of the logarithm in Eq. (6.37). The bound can be further weakened by adding extra field content to the MSSM, *e.g.* as in the NMSSM, but there is an upper perturbative limit of $m_h \approx 150 \text{ GeV}$.

It is very interesting to discuss what the Higgs discovery actually implies for low-energy supersymmetry. As can be seen from the above numbers it requires rather large squark masses even in the favourable scenario with $\tan \beta \gg 1$ where the tree level mass is $m_h \sim 90 \text{ GeV}$. A naive estimate from Eq. (6.37) gives $m_{\tilde{t}} > 1 \text{ TeV}$. However, this does not take into account possible negative contributions to the Higgs mass from heavy gauginos (fermions in the vector superfields), and possible increases in the stop contribution due to tuning of the mixing of the chiral eigenstates $\tilde{t}_L$ and $\tilde{t}_R$, in the mass eigenstates $\tilde{t}_1$ and $\tilde{t}_2$.

²²It is worth pointing out here that the MSSM, despite its many parameters, is a falsifiable theory. For example, had the Higgs boson mass been $\sim 15 \text{ GeV}$ higher, which is perfectly allowed in the Standard Model, the MSSM would have been excluded.

Since the lightest stop quark is expected to be the lightest squark in scenarios with universal soft masses at some high scale – the reasoning here is the large downward RGE running of m_{33}^Q from a common squark soft mass at some high scale due to the large top Yukawa coupling – the expected sparticle spectrum lies mostly above 1 TeV, with the possible exception of gauginos/higgsinos. This points to so-called **Split-SUSY** scenarios with heavy scalars and light gauginos, and a relatively large degree of fine-tuning. If one can live with this little hierarchy problem, it will explain why no signs of supersymmetry have been seen yet at the LHC.

If you are willing to accept fine-tuning of the stop mixing instead, or come up with a good reason for why the mixing should be just-so to give a maximal Higgs mass, you can keep fairly light stop quarks. With the addition of light higgsinos and a light gluino the model is then technically natural, these scenarios are called **Natural SUSY** and could be within the current or near future reach of the LHC. The problem with these models, as we shall see, is that the higgsinos are degenerate, and thus difficult to detect.

To do calculations with the Higgs bosons in the MSSM we need the Feynman rules that result from the relevant Lagrangian terms. Since these have been listed elsewhere we will not repeat them here, but recommend in particular the PhD-thesis of Peter Richardson [15], where they can be found in Appendix A.6, including all interactions with fermions and sfermions. These can also be found, together with all gauge and self-interactions, in the classic paper by Gunion and Haber [16]. Note that in this paper a complex Higgs singlet appears which can safely be ignored.

6.11 The gluino

The fermion partner of the Standard Model gluon g is called the **gluino** $\tilde{g}$, and as the gluon it is a colour octet Majorana fermion. We usually talk about the gluino as being one particle, however, as an adjoint representation of $SU(3)$ there are actually eight (thus octet) distinct gluons, and we write $\tilde{g}^a$ when we want to make the distinction. As a colour octet it has nothing to mix with in the MSSM – this is still true even if we add R-parity violation – and at tree level the mass is given by the soft term M_3 . Since it lives in the same superfield as the massless gluon it would otherwise had zero mass.

The one complication for the gluino is that it is strongly interacting so $M_3(\mu)$ runs quickly with energy. It is useful to instead talk about the scale-independent **pole-mass**, *i.e.* the pole of the renormalised propagator, $m_{\tilde{g}}$. Including one loop effects due to gluon exchange and squark loops, see Fig. 6.5, in the $\overline{DR}$ renormalisation scheme we get:²³

$$m_{\tilde{g}} \simeq M_3(\mu) \left[1 + \frac{\alpha_s}{4\pi} \left(15 + 6 \ln \frac{\mu}{M_3} + \sum_{\text{all } \tilde{q}} A_{\tilde{q}} \right) \right],$$

where the squark loop contributions are

$$A_{\tilde{q}} = \int_0^1 dx x \ln \left(x \frac{m_{\tilde{q}}^2}{M_3^2} + (1-x) \frac{m_q^2}{M_3^2} - x(1-x) - i\epsilon \right).$$

Due to the 15-factor the correction can be significant (colour factor).

²³Note that the right-hand side here is μ dependent since the expression is only to finite order.

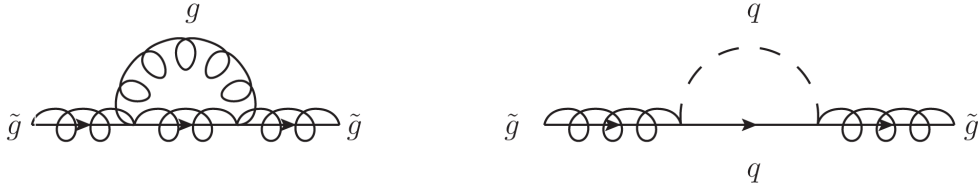


Figure 6.5: One loop contributions to the gluino mass.

Complete Feynman rules for gluinos can be found in Appendix C of the classic MSSM reference paper of Haber & Kane [17]. A more comprehensible alternative may be Appendix A.3 from the PhD-thesis of Bolz [18]. This thesis also provides a diagrammatic prescription of how to handle clashing fermion lines that can appear with Majorana fermions.

6.12 Neutralinos & Charginos

In the MSSM we have a bunch of fermion fields that can mix when electroweak symmetry is broken and we do not have to care about the $SU(2)_L \times U(1)_Y$ charges of the fields, only the charges under the remaining $U(1)_{\text{em}}$ symmetry matter. The candidates for mixing are the (Majorana) fermions from the $U(1)$ and $SU(2)$ vector superfields B^0 and W^a called the **gauginos**:

$$\tilde{B}^0 \text{ (bino)}, \quad \tilde{W}^0 \text{ (neutral wino)}, \quad \tilde{W}^\pm \text{ (charged wino)},$$

and the fermions from the Higgs superfields H_u and H_d , called **higgsinos**:

$$\tilde{H}_u^+, \quad \tilde{H}_u^0, \quad \tilde{H}_d^- \quad \text{and} \quad \tilde{H}_d^0.$$

Together these are called the **electroweakinos**.

In the Standard Model the neutral gauge field components B_μ^0 and W_μ^0 mix into the photon and the Z -boson. The neutral gauginos can mix the same way into the states

$$\tilde{\gamma} = N'_{11}\tilde{B}^0 + N'_{12}\tilde{W}^0 \quad \text{(photino)}, \quad (6.39)$$

$$\tilde{Z} = N'_{21}\tilde{B}^0 + N'_{22}\tilde{W}^0 \quad \text{(zino)}. \quad (6.40)$$

However, more generally, they also mix with the neutral higgsinos to form the four **neutralinos**:²⁴

$$\tilde{\chi}_i^0 = N_{i1}\tilde{B}^0 + N_{i2}\tilde{W}^0 + N_{i3}\tilde{H}_d^0 + N_{i4}\tilde{H}_u^0, \quad i = 1, 2, 3, 4, \quad (6.41)$$

where N_{ij} indicates size of the component of each of the fields in the **gauge eigenstate basis**

$$\tilde{\psi}^{0T} = (\tilde{B}^0, \tilde{W}^0, \tilde{H}_d^0, \tilde{H}_u^0). \quad (6.42)$$

²⁴The neutral higgsinos are also Majorana fermions despite coming from scalar superfields. Unlike the (s)fermion superfields the Higgs superfields have no $\tilde{H}$ chiral partners to supply the left-right Weyl spinor combinations required for Dirac fermions. Thus the neutralinos are Majorana fermions.

In the gauge eigenstate basis the neutralino mass term can be written as

$$\mathcal{L}_{\chi-\text{mass}} = -\frac{1}{2}\tilde{\psi}^{0T} M_{\tilde{\chi}} \tilde{\psi}^0 + \text{c.c.},$$

where the mass matrix $M_{\tilde{\chi}}$ is found from the bilinear terms in the Lagrangian with gauge eigenstates to be

$$M_{\tilde{\chi}} = \begin{bmatrix} M_1 & 0 & -\frac{1}{\sqrt{2}}g'v_d & \frac{1}{\sqrt{2}}g'v_u \\ 0 & M_2 & \frac{1}{\sqrt{2}}gv_d & -\frac{1}{\sqrt{2}}gv_u \\ -\frac{1}{\sqrt{2}}g'v_d & \frac{1}{\sqrt{2}}gv_d & 0 & -\mu \\ \frac{1}{\sqrt{2}}g'v_u & -\frac{1}{\sqrt{2}}gv_u & -\mu & 0 \end{bmatrix}. \quad (6.43)$$

In this matrix, the upper left diagonal part comes from the soft terms for the $\tilde{B}^0$ and the $\tilde{W}^0$, the lower right off diagonal matrix comes from the superpotential term $\mu H_u H_d$, while the remaining entries come from Higgs-higgsino-gaugino terms from the kinetic part of the Lagrangian, *e.g.* $H_u^\dagger e^{\frac{1}{2}g\sigma W + g'B} H_u$, which become mass terms when one of the neutral Higgs component fields acquires a vev. With the Z -mass condition on the vevs (6.32) we can also write

$$\frac{1}{\sqrt{2}}g'v_d = \cos\beta \sin\theta_W m_Z, \quad (6.44)$$

$$\frac{1}{\sqrt{2}}g'v_u = \sin\beta \sin\theta_W m_Z, \quad (6.45)$$

$$\frac{1}{\sqrt{2}}gv_d = \cos\beta \cos\theta_W m_Z, \quad (6.46)$$

$$\frac{1}{\sqrt{2}}gv_u = \sin\beta \cos\theta_W m_Z. \quad (6.47)$$

The mass matrix can now be diagonalised to find the $\tilde{\chi}_i^0$ masses. If N is a unitary diagonalisation matrix for $M_{\tilde{\chi}}$, we can write

$$\mathcal{L}_{\chi-\text{mass}} = -\frac{1}{2}\tilde{\psi}^{0T} N^T N^* M_{\tilde{\chi}} N^\dagger N \tilde{\psi}^0 + \text{c.c.}, = -\frac{1}{2}\tilde{\chi}^{0T} D \tilde{\chi}^0 + \text{c.c.},$$

where $D = N^* M_{\tilde{\chi}} N^\dagger$ is diagonal and contains the neutralino masses $m_{\tilde{\chi}_i^0}$, and we see that N gives the mixing of the gauge eigenstates $\tilde{\psi}^0$ into the mass eigenstates $\tilde{\chi}^0$. We number the neutralino eigenstates in (6.41), or, equivalently, sort the mass eigenvalues after diagonalisation, so that the neutralinos are numbered from lightest to heaviest.

The mass parameters of the neutralino mass matrix may in general be complex, leading to complex mass eigenvalues. Redefinition of fields can rotate away either the M_1 or M_2 phase, to make the parameter real and positive, but not both and not the μ -phase. This gives rise to problematic CP-violation that can easily be in contradiction with experiments. Therefore, these are often just assumed to be real in order not to violate experimental bounds. However, we are in principle perfectly happy with negative or even complex mass eigenvalues here, as this is just a phase for the corresponding mass eigenstate in (6.41).

One particularly interesting solution to the diagonalisation is in the limit where EWSB is a small effect, $m_Z \ll |\mu \pm M_1|, |\mu \pm M_2|$, and when $M_1 < M_2 \ll |\mu|$. Then the lightest neutralino

is bino-like, $\tilde{\chi}_1^0 \approx \tilde{B}^0$, the next-to-lightest is wino like, $\tilde{\chi}_2^0 \approx \tilde{W}^0$, and $\tilde{\chi}_{3,4}^0 \approx \frac{1}{\sqrt{2}}(\tilde{H}_d^0 \pm \tilde{H}_u^0)$, and the masses are to first order in $1/\mu$:

$$m_{\tilde{\chi}_1^0} = M_1 + \frac{m_Z^2 \sin^2 \theta_W \sin 2\beta}{\mu} + \dots \quad (6.48)$$

$$m_{\tilde{\chi}_2^0} = M_2 - \frac{m_W^2 \sin 2\beta}{\mu} + \dots \quad (6.49)$$

$$m_{\tilde{\chi}_{3,4}^0} = |\mu| + \frac{m_Z^2}{2\mu}(\text{sgn } \mu \mp \sin 2\beta) + \dots \quad (6.50)$$

Since the LSP is stable in R-parity conserving theories the lightest neutralino is an excellent candidate for dark matter. In particular since a neutralino with mass around 100 GeV has a natural relic density close to the measured dark matter density of the Universe. We will return to this issue in Chapter 8.

From the charged electroweakinos we can make **charginos** $\tilde{\chi}_i^\pm$ that are Dirac fermions with mass term

$$\mathcal{L}_{\chi^\pm\text{-mass}} = -\frac{1}{2}\tilde{\psi}^{\pm T} M_{\chi^\pm} \tilde{\psi}^\pm + \text{c.c.},$$

where the gauge eigenstate basis is $\tilde{\psi}^{\pm T} = (\tilde{W}^+, \tilde{H}_u^+, \tilde{W}^-, \tilde{H}_d^-)$, and the mass matrix is given by

$$M_{\tilde{\chi}^\pm} = \begin{bmatrix} 0 & 0 & M_2 & gv_d \\ 0 & 0 & gv_u & \mu \\ M_2 & gv_u & 0 & 0 \\ gv_d & \mu & 0 & 0 \end{bmatrix}.$$

Here the M_2 terms come from the soft terms for the charged winos $\tilde{W}^\pm$, the μ terms come from the superpotential as above, while the remaining terms come from the kinetic terms. We can here re-write

$$gv_d = \sqrt{2} \cos \beta m_W, \quad (6.51)$$

$$gv_u = \sqrt{2} \sin \beta m_W. \quad (6.52)$$

Diagonalising this mass matrix gives the mass eigenstates $\tilde{\chi}_i^\pm$, $i = 1, 2$. The eigenvalues are doubly degenerate, giving the same masses to the $\tilde{\chi}_i^+$ and $\tilde{\chi}_i^-$ particle and anti-particle pairs, and are explicitly given as:

$$m_{\tilde{\chi}_{1,2}^\pm} = \frac{1}{2} \left(|M_2|^2 + |\mu|^2 + 2m_W^2 \mp \sqrt{(|M_2|^2 + |\mu|^2 + 2m_W^2)^2 - 4|\mu M_2 - m_W^2 \sin^2 \beta|^2} \right).$$

In the same limit of small EWSB effects discussed above we have $\tilde{\chi}_1^\pm \approx \tilde{W}^\pm$ and $\tilde{\chi}_2^\pm \approx \tilde{H}_u^\pm / \tilde{H}_d^\pm$ with

$$m_{\tilde{\chi}_1^\pm} = M_2 - \frac{m_W^2}{\mu} \sin 2\beta + \dots, \quad (6.53)$$

$$m_{\tilde{\chi}_2^\pm} = |\mu| + \frac{m_W^2}{\mu} \text{sgn } \mu + \dots \quad (6.54)$$

Note that in this limit $m_{\tilde{\chi}_2^0} \simeq m_{\tilde{\chi}_1^\pm}$.

We saw earlier that the soft-mass ratio

$$M_3 : M_2 : M_1 \simeq 6 : 2 : 1,$$

appears at a scale of around 1 TeV if the same soft-masses unify at the GUT-scale. From our above discussion, as long as $|\mu| \gg M_1, M_2$, this gives the very predictive mass relationships $m_{\tilde{g}} \simeq 6m_{\tilde{\chi}_1^0}$, $m_{\tilde{\chi}_2^0} \simeq m_{\tilde{\chi}_1^\pm} \simeq 2m_{\tilde{\chi}_1^0}$. However, it is important to remember that this often used relationship is based on the *conjecture* of gauge coupling unification, and the unification of gaugino soft masses!

We should mention that some authors prefer other symbols for the neutralinos and charginos. Common examples are $\tilde{N}_i$ or $\tilde{Z}_i$ for neutralinos, and $\tilde{C}_i$ or $\tilde{W}_i$ (again!) for the charginos.

Feynman rules for charginos & neutralinos can again be found in Haber & Kane [17].

6.13 Sleptons & Squarks

The **sfermions**, the scalar partners of the Standard Model fermions, consists of the **squarks** and the **sleptons**. Reading of from the MSSM Lagrangian there are multiple tree-level contributions to the sfermion masses. In the following $\tilde{F}_i$ represents a $SU(2)_L$ doublet of sfermions with generation index i , for example $\tilde{Q}_i = (\tilde{u}_{iL}, \tilde{d}_{iL})$, while $\tilde{f}_{iR}$ represents a singlet. We make the following list:

- i) Under the reasonable assumption that soft masses are (close to) diagonal²⁵ the sfermions get contributions $-m_F^2 \tilde{F}_i^\dagger \tilde{F}_i$ and $-m_f^2 \tilde{f}_{iR}^* \tilde{f}_{iR}$ from the soft terms.
- ii) There are F -term contributions that come from Yukawa terms in the superpotential of the form $y_f F H \bar{K}$. From the contribution $\sum |W_i|^2$ to the scalar potential these give Lagrangian terms $y_f^2 H^{0*} H^0 \tilde{f}_{iL}^* \tilde{f}_{iL}$ and $y_f^2 H^{0*} H^0 \tilde{f}_{iR}^* \tilde{f}_{iR}$. After EWSB we then get the mass terms $m_f^2 \tilde{f}_{iL}^* \tilde{f}_{iL}$ and $m_f^2 \tilde{f}_{iR}^* \tilde{f}_{iR}$, where $m_f = v_{u/d} y_f$. These are only significant for large Yukawa coupling y_f , and give the same mass as their Standard Model fermion partner gets from the same Yukawa terms.
- iii) There are also so-called **hyperfine** terms that come from D -terms $\sum g^2 (A^* T^a A)^2$ in the scalar potential that give Lagrangian terms of the form $(\text{sfermion})^2 (\text{Higgs})^2$ when one of the scalar fields A is a neutral Higgs field, and the other is a sfermion. Under EWSB, when the Higgs field gets a vev these become mass terms. They contribute with a mass

$$\Delta_F = (T_{3F} g^2 - Y_F g'^2)(v_d^2 - v_u^2) = (T_{3F} - Q_F \sin^2 \theta_W) \cos 2\beta m_Z^2,$$

where the weak isospin, T_3 , hypercharge, Y , and electric charge, Q , are for the left-handed supermultiplet F to which the sfermion belongs. However, these contributions are usually quite small.

- iv) Furthermore, there are also F -terms that combine scalars from the $\mu H_u H_d$ term and Yukawa terms $y_f F H \bar{K}$ in the superpotential. These give Lagrangian terms $-\mu^* H^{0*} y_f \tilde{f}_L \tilde{f}_R^*$. With a Higgs vev this gives mass terms $-\mu^* v_{u/d} y_f \tilde{f}_R^* \tilde{f}_L + \text{c.c.}$

²⁵This assumption is of course made to avoid flavour changing neutral currents (FCNCs). However, it is also reasonable in that if the soft masses are diagonal, or even all the same, at a high scale, the RGE running will not create large off-diagonal terms.

- v) Finally, the soft Yukawa terms of the form $a_f \tilde{F} H \tilde{f}_R^*$ with a Higgs vev give mass terms $a_f v_{u/d} \tilde{f}_L \tilde{f}_R^* + \text{c.c.}$ ²⁶

For the first two generations of sfermions, terms of type ii), iv) and v) are small due to small Yukawa couplings. Then the sfermion masses are for example

$$m_{\tilde{u}_L}^2 = m_{Q_1}^2 + \Delta \tilde{u}_L, \quad (6.55)$$

$$m_{\tilde{d}_L}^2 = m_{Q_1}^2 + \Delta \tilde{d}_L, \quad (6.56)$$

$$m_{\tilde{u}_R}^2 = m_{u_1}^2 + \Delta \tilde{u}_R. \quad (6.57)$$

Mass splitting between same generation slepton/squark is then given by the hyperfine splitting

$$m_{\tilde{e}_L}^2 - m_{\tilde{\nu}_{eL}}^2 = m_{\tilde{d}_L}^2 - m_{\tilde{u}_L}^2 = -\frac{1}{2} g^2 (v_d^2 - v_u^2) = -\cos 2\beta m_W^2,$$

since they have the same hypercharge, see Table 6.1. For $\tan \beta > 1$ this gives the definite prediction $m_{\tilde{e}_L}^2 > m_{\tilde{\nu}_{eL}}^2$ and $m_{\tilde{d}_L}^2 > m_{\tilde{u}_L}^2$.

The **third generation sfermions** $\tilde{t}$, $\tilde{b}$ and $\tilde{\tau}$ have a more complicated mass matrix structure, *e.g.* in the gauge eigenstate basis $(\tilde{t}_L, \tilde{t}_R)$ for stop quarks the mass term is

$$\mathcal{L}_{\text{stop}} = -(\tilde{t}_L \quad \tilde{t}_R) m_{\tilde{t}}^2 \begin{pmatrix} \tilde{t}_L \\ \tilde{t}_R \end{pmatrix},$$

where the mass matrix is given by

$$m_{\tilde{t}}^2 = \begin{bmatrix} m_{Q_3}^2 + m_t^2 + \Delta \tilde{u}_L & v(a_t^* \sin \beta - \mu y_t \cos \beta) \\ v(a_t \sin \beta - \mu^* y_t \cos \beta) & m_{u_3}^2 + m_t^2 + \Delta \tilde{u}_R \end{bmatrix}. \quad (6.58)$$

Here the diagonal elements come from i), ii) and iii), while the off-diagonal elements come from iv) and v).

To find the particle masses, we must diagonalise this matrix, writing it in terms of the mass eigenstates $\tilde{t}_1$ and $\tilde{t}_2$, acquiring also here a unitary mixing matrix for the mass eigenstates in terms of the gauge eigenstates $\tilde{t}_L$ and $\tilde{t}_R$:

$$\begin{pmatrix} \tilde{t}_1 \\ \tilde{t}_2 \end{pmatrix} = \begin{bmatrix} c_{\tilde{t}} & -s_{\tilde{t}}^* \\ s_{\tilde{t}} & c_{\tilde{t}} \end{bmatrix} \begin{pmatrix} \tilde{t}_L \\ \tilde{t}_R \end{pmatrix}, \quad (6.59)$$

where the matrix entries are related by $|c_{\tilde{t}}|^2 + |s_{\tilde{t}}|^2 = 1$ and $m_{\tilde{t}_1}^2 < m_{\tilde{t}_2}^2$ are the eigenvalues of (6.58). The suggestive form of the mixing matrix indicates that if the off-diagonal elements of the original mass matrix has only real elements, this mixing matrix can be written as an element in $SO(2)$, using sine and cosine of a mixing angle $0 \leq \theta_{\tilde{t}} < \pi$, $c_{\tilde{t}} = \cos \theta_{\tilde{t}}$ and $s_{\tilde{t}} = \sin \theta_{\tilde{t}}$. The matrices for $\tilde{b}$ and $\tilde{t}$ have the same structure.

Since the third generation sneutrino $\tilde{\nu}_{eL}$ does not have a corresponding right-handed state in the MSSM, there is no mixing, and it has the same mass term as the first and second generation sneutrinos.

A good source for sfermion interaction Feynman rules is the PhD-thesis of Richardson [15].

²⁶We often assume that $a_f = A_0 y_f$ in order to further reduce the FCNC, meaning that there is a global constant A_0 with unit mass relating the Yukawa couplings and the trilinear A-term couplings.

6.14 Exercises

Exercise 6.1

Using the explicit form of the $SU(3)_C$ transformations with the Gell-Mann matrices, show that with our definition of the superpotential term $\bar{U}_i \bar{D}_j \bar{D}_k$ this is invariant under $SU(3)_C$.

Exercise 6.2

Show how you can eliminate the parameters $|\mu|$ and b by using the properties of the minimum of the potential in Eq. (6.28).

Exercise 6.3

Show Eqs. (6.44)–(6.47).

Chapter 7

Sparticle phenomenology

In this chapter we discuss the phenomenology of supersymmetric models and how to search for low-energy particle realisations of supersymmetry in experiments. We begin by returning to supersymmetry breaking in order to define some reasonable and (partially) motivated subsets of the 124 MSSM parameters which can be used to define more constrained models. We then discuss supersymmetry at lepton and hadron colliders, and finally look at precision measurements that are indirectly sensitive to the existence of sparticles.

7.1 Models for supersymmetry breaking

Let us begin by taking a closer look at the models we use to motivate supersymmetry breaking, and what their phenomenological consequences are. This is important to keep in mind as most searches for supersymmetry are interpreted under certain assumptions on the supersymmetry breaking mechanism.

Generically such models can be illustrated as shown in Fig. 7.1. There is one or more **hidden sector** (HS) scalar superfield X – by hidden we mean that it has no or very small direct couplings to the MSSM fields – that has an *effective (non-renormalizable) coupling* to the MSSM scalar fields Φ_i of the form

$$\mathcal{L}_{\text{HS}} = -\frac{1}{M}(\bar{\theta}\theta)X\Phi_i\Phi_j\Phi_k, \quad (7.1)$$

where M is some large scale, *e.g.* the Planck scale, that suppresses the interaction. Figure 7.2 shows interactions that can lead to such terms, where M is the mass scale of some **mediator** particle Y .

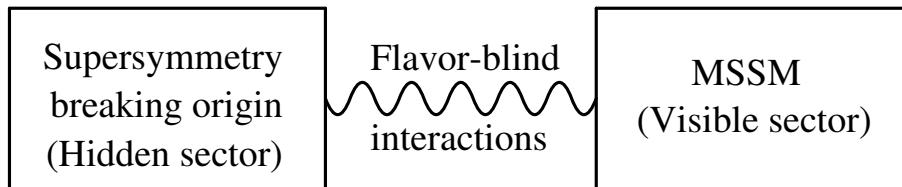


Figure 7.1: A generic illustration of how to generate soft breaking terms [13].

If the hidden sector is constructed so that X develops a vev for its auxillary F -component field, F_X ,

$$\langle X \rangle = \theta\theta\langle F_X \rangle, \quad (7.2)$$

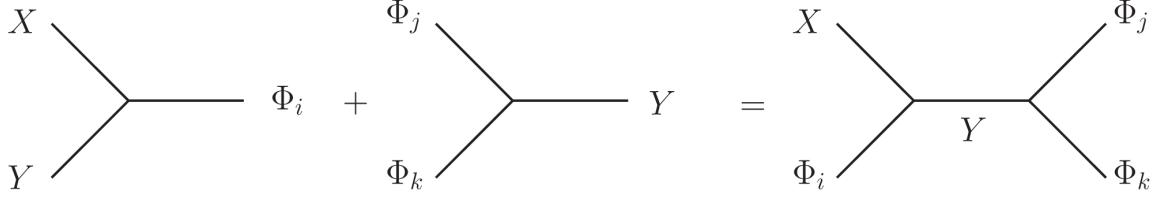


Figure 7.2: Interactions leading to effective 4-particle couplings in our example.

it breaks supersymmetry spontaneously, see the discussion leading up to Eq. (5.5). As a result, the interaction in (7.1) will produce a soft-term of the form of the second term in Eq. (5.9),

$$\mathcal{L}_{\text{soft}} = -\frac{\langle F_X \rangle}{M} A_i A_j A_k, \quad (7.3)$$

with the soft-mass parameter

$$m_{\text{soft}} = \frac{\langle F_X \rangle}{M}.$$

This has reasonable limits in that $m_{\text{soft}} \rightarrow 0$ as $\langle F_X \rangle \rightarrow 0$, which is the limit of no supersymmetry breaking, and $m_{\text{soft}} \rightarrow 0$ as $M \rightarrow \infty$, where the interaction with the hidden sector is decoupled because the mediating particle Y becomes too heavy to have any influence.

We will now look at two possible ways to construct such a hidden sector called Planck-scale Mediated Supersymmetry Breaking (PMSB) and Gauge Mediated Supersymmetry Breaking (GMSB).

7.1.1 Planck-scale Mediated Supersymmetry Breaking (PMSB)

In Planck-scale mediated supersymmetry breaking (PMSB) we blame some gravity mechanism for mediating the supersymmetry breaking from the hidden sector to the MSSM so that the scale of the breaking is $M = M_P = 2.4 \cdot 10^{18}$ GeV. Then we need to have $\sqrt{\langle F_X \rangle} \sim 10^{10} - 10^{11}$ GeV in order to get $m_{\text{soft}} \simeq 50 - 5000$ GeV, which is roughly of the right magnitude not to re-introduce the hierarchy problem. The use of $\sqrt{\langle F_X \rangle}$ is just a conventional shorthand notation for the magnitude of the vev of whichever F -term that breaks supersymmetry. This is called the **supersymmetry breaking scale**.

The complete soft terms for such a mechanism can then be shown to be

$$\begin{aligned} \mathcal{L}_{\text{soft}} = & -\frac{\langle F_X \rangle}{M_P} \left(\frac{1}{2} f_i \lambda_i^a \lambda_i^a + \frac{1}{6} y'_{ijk} A_i A_j A_k + \frac{1}{2} \mu'_{ij} A_i A_j + \frac{\langle F_X \rangle^*}{M_P^2} x_{ijk} A_i^* A_j A_k + \text{c.c.} \right) \\ & - \frac{|\langle F_X \rangle|^2}{M_P^2} k_{ij} A_i A_j^*. \end{aligned} \quad (7.4)$$

Incidentally, we can now see why we assumed the maybe-soft breaking terms to be unimportant, as in this model they are suppressed by $|\langle F_X \rangle|/M_P^2$ compared to the other masses.

If one assumes a minimal form for the parameters at the GUT scale, motivated by the wish for unification, *e.g.* $f_i = f$, $y'_{ijk} = \alpha y_{ijk}$, $\mu'_{ij} = \beta \mu$, and $k_{ij} = k \delta_{ij}$, then all the soft terms

are fixed by just four parameters

$$m_{1/2} = f \frac{\langle F_X \rangle}{M_P}, \quad m_0^2 = k \frac{|\langle F_X \rangle|^2}{M_P^2}, \quad A_0 = \alpha \frac{\langle F_X \rangle}{M_P}, \quad B_0 = \beta \frac{\langle F_X \rangle}{M_P}.$$

The resulting phenomenology is called **minimal supergravity**, or **mSUGRA/CMSSM**. This is minimal in the sense of the form of the parameters, and is the most studied, but perhaps not best motivated, version of the MSSM. Usually, B_0 and $|\mu|$ are exchanged for $\tan \beta$ at low scales using the EWSB condition in Eq. (6.31), so it is common to say that there are four and a half parameters in the model: $m_{1/2}$, m_0 , A_0 , $\tan \beta$ and $\text{sgn } \mu$.

7.1.2 Gauge Mediated Supersymmetry Breaking (GMSB)

An alternative to PMSB is gauge-mediated supersymmetry breaking where soft terms come from *loop diagrams* with **messenger** superfields that get their own mass by coupling to the hidden sector supersymmetry breaking vev, and that have Standard Model gauge interactions. By dimensional analysis we must have

$$m_{\text{soft}} = \frac{\alpha_i}{4\pi} \frac{\langle F \rangle}{M}.$$

If now the supersymmetry breaking scale $\sqrt{\langle F \rangle}$ and the messenger mass M are roughly comparable in size, which is reasonable given where the messenger mass comes from, then $\sqrt{\langle F \rangle} \simeq 100 \text{ TeV}$ can give a viable sparticle spectrum. Notice that there is now a lot less RGE running for the parameters in the GMSB compared to the PMSB since the soft masses are created at a rather low scale.

One way of thinking about how these mass terms appear is that the messenger field(s) get masses from hidden sector vevs and contribute to for example gaugino mass terms through diagrams such as the one in Fig. 7.3, where messenger scalars and fermions run in the loop, and their masses from the hidden sector vevs are symbolised by the mass insertions. Note that scalars can only get mass contributions like this at two-loop order since the messenger interaction is a gauge interaction, involving gauge bosons or gauginos in the MSSM. In order not to spoil GUT unification messengers are often assumed to have small mass splittings and come in N_5 complete $\mathbf{5} + \bar{\mathbf{5}}$ (fundamental) representations of $SU(5)$.

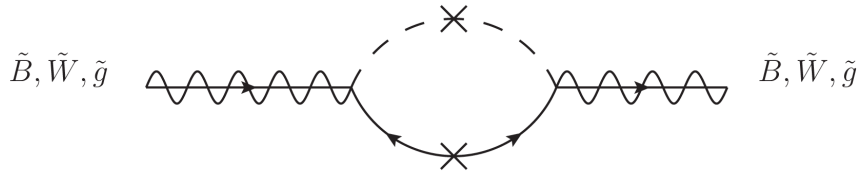


Figure 7.3: Diagram for GMSB giving masses to the gauginos. The messenger scalars and fermions run in the loop.

The minimal parametrisation of GMSB models is in terms of the scale $\Lambda = \frac{\langle F \rangle}{M}$, the messenger mass M , the number of representations N_5 , and $\tan \beta$ for the EWSB criterion

(instead of μ and b). This gives the soft masses

$$M_i = \frac{\alpha_i}{4\pi} \Lambda N_5, \quad (\text{gaugino masses}) \quad (7.5)$$

$$m_j^2 = 2\Lambda^2 N_5 \sum C(R)_i \left(\frac{\alpha_i}{4\pi} \right)^2, \quad (\text{scalar masses}) \quad (7.6)$$

where the sum is over the quadratic Casimir invariant for the scalar superfield Φ_j that the scalar field belongs to. We clearly see that the scalar soft-masses are a two-loop effect as discussed above.

While this parameterisation looks independent of M , the messenger scale sets the starting point of the RGE running of the sparticle masses, and thus influences their magnitude. For example, the tri-linear soft-term couplings a_{ijk} are expected to be very small at the messenger scale, and are effectively set to zero, however, due to RGE running they are small, but non-zero at the electroweak scale. Since the scalar masses m_j scale as $\sqrt{N_5}$ compared to N_5 for the gauginos, we in general expect the scalars to be lighter in GMSB models. One should also notice that this parameterisation gives the same hierarchy of gaugino masses as in mSUGRA, $M_3 > M_2 > M_1$, since (7.5) is ordered in terms of the strength of the gauge couplings α_i . The origin of the hierarchy is different since in mSUGRA it comes from the running of the parameters down from the GUT scale.

7.2 Supersymmetry at lepton colliders

High-energy lepton colliders have traditionally been circular e^+e^- -colliders. These have the advantage of a well known centre-of-mass (CoM) energy given by the total energy of the electron-positron pair, clean final states, and periodic application of the accelerating gradient. The challenge is to reach high energies and high luminosities (collision rates). Since the electrons are light they radiate a lot of bremsstrahlung photons when bent in orbits. The highest energy so-far at an e^+e^- -collider was the 209 GeV CoM-energy at LEP2 in 2000. Plans are being made both for linear e^+e^- -colliders, and a muon collider where there is less bremsstrahlung because of the higher muon mass, meaning that higher energies can be reached. However, there are significant technical challenges ahead for both. The linear colliders need very long installation tunnels and a very high accelerating gradient, while the muon collider must be able to produce and store the unstable muons.

Most supersymmetry searches at lepton colliders rely on the pair production of oppositely charged sparticles with electroweak couplings from $e^+e^- \rightarrow \gamma^*/Z^* \rightarrow S\bar{S}$, where S symbolises a generic sparticle. Due to R-parity conservation these sparticles both decay to lighter sparticles and Standard Model decay products, until they, potentially after several successive decays known as a **cascade decay**, leave only the stable LSP. This LSP must be electrically (and most likely colour) neutral due to experimental constraints on massive long-lived charged particles that would bind to atomic nuclei. The neutrality of the LSP means that it escapes detectors unseen.

The search for supersymmetry thus focuses on events with the Standard Model decay products of sparticles and an imbalance in momentum conservation due to the two missing LSPs. By the measured sum of the momenta of all the visible decay products the sum of the total momenta for the invisible particles can be inferred as going in the opposite direction. This is known as a **missing energy** signature.

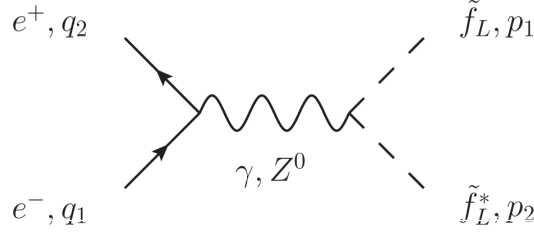


Figure 7.4: Feynman diagram for the pair production of left-handed sfermions in the s-channel at an e^+e^- -collider.

In practice such missing energy measurements are challenging experimentally due to the absence of detectors near the incoming beams in the longitudinal direction, where particles that are in principle visible may escape undetected and create an artificial momentum imbalance in the longitudinal direction. At high energies this is exacerbated by the increase of collinear bremsstrahlung from the incoming electrons at the interaction point due to interaction of the beams. This will be a particularly difficult for a future 0.5 – 3.0 TeV CoM International Linear Collider (ILC) or the Compact LInear Collider (CLIC) project.¹

We can now discuss more specifically the search for sfermions, neutralinos & charginos, and Higgs bosons at lepton colliders.

7.2.1 Sfermions

We can estimate the leading order amplitude of the s -channel sfermion pair production process shown in Fig. 7.4. We begin by writing down the matrix element with an intermediary γ as:

$$\mathcal{M} = \bar{v}ie\gamma^\mu u \frac{-ig_{\mu\nu}}{k^2 + i\epsilon} [-ie \cdot e_f(p_1 - p_2)^\nu], \quad (7.7)$$

which gives a squared matrix element of, assuming that the CoM s is much greater than m_Z and taking into account both the photon and the Z ,

$$|\mathcal{M}|^2 \simeq \frac{g^4 e_f^2}{8 \cos \theta_W} \frac{st + (m_{\tilde{f}}^2 - t)^2}{s^2} \times (1 + (4 \sin^2 \theta_W - 1)^2). \quad (7.8)$$

Here, we safely take $(1 + (4 \sin^2 \theta_W - 1)^2) \simeq 1$. The complete differential cross-section is then:

$$\frac{d\sigma}{dt} = \frac{1}{32\pi} \frac{1}{s^2} |\mathcal{M}|^2. \quad (7.9)$$

This cross section is relatively small due to the electroweak coupling factor g^4 and sfermion mass suppression, so sfermion production events will be rare. We show an example of the slepton pair production cross section including the Z -resonance, and the t -channel contributions from neutralinos (see below) in Fig. 7.5.

After production sfermions will decay, typically to a lighter neutralino or chargino and a Standard Model fermion f , in processes of the type $\tilde{f} \rightarrow f\tilde{\chi}_i^0$ or $\tilde{f} \rightarrow f'\tilde{\chi}_i^\pm$. Since the energy

¹For more information on these projects see the websites for the International Linear Collider <http://www.linearcollider.org/> and the Compact LInear Collider <https://clic.cern/>

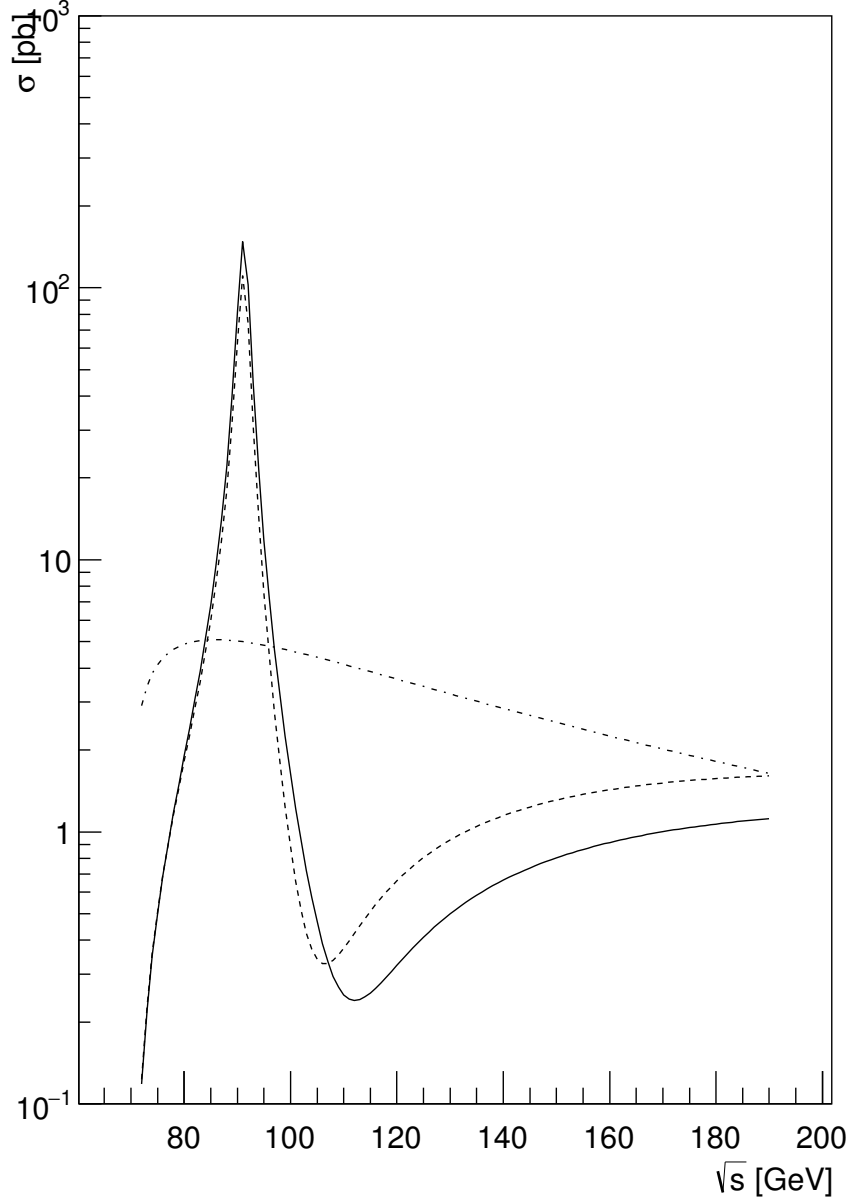


Figure 7.5: Cross sections for selectron pair production as a function of CoM energy. Cross sections for $\tilde{e}_L^* \tilde{e}_L$ (solid line), $\tilde{e}_R^* \tilde{e}_R$ (dashed line), and $\tilde{e}_L^* \tilde{e}_R$ (dashed dotted line) are shown separately. The particular model point has a common slepton mass of $m_{\tilde{e}_{L/R}} = 35$ GeV.

reach is limited, only the very lightest sfermions are likely to be producible, which are usually the sleptons, and these are likely to be near (but above) the mass of the lightest neutralino. See Fig. 6.4 for some expectations of the masses in an mSUGRA context. This means that

the most expected decay is $\tilde{\ell} \rightarrow \ell \tilde{\chi}_1^0$, and the total process is then $e^+e^- \rightarrow \tilde{\ell}^+\tilde{\ell}^- \rightarrow \ell^+\ell^- \tilde{\chi}_1^0\tilde{\chi}_1^0$, giving two oppositely charged Standard Model leptons and missing energy as the experimental signature.

Such a (weak) signal then needs to be discriminated from Standard Model backgrounds from for example W^+W^- pair production, which can lead to the final state $\ell^+\nu_\ell\ell^-\bar{\nu}_\ell$, featuring the same leptons and missing energy from the neutrinos. Since backgrounds are usually well under control at lepton colliders, the limit to searches is the ability to produce the sparticles at all. With a CoM energy $\sqrt{s}$ sparticles up to $\sqrt{s}/2$ are energetically possible, and a rule of thumb is that the reach approaches $\sqrt{s}/2$ from below given sufficient data, never quite getting to $\sqrt{s}/2$.

Should some excess be discovered, we need a smoking duck in order to confirm that this is indeed supersymmetry. We would like to identify and measure the masses of as many new particles as possible, and hopefully also their spin. The properties of the sparticles can be measured through the inferred cross section, and the kinematical distribution of the final state products. The ability of a lepton collider to easily change the CoM energy $\sqrt{s}$, allows for a so-called **threshold scan** of the cross section where the cross section is measured as a function of $\sqrt{s}$ around where it becomes zero. This in turn allows for a very precise measurement of the mass of the pair produced particle. As an example of a kinematic distribution, for the process discussed here, it can be shown that the energy distribution for the final state leptons is a uniform distribution between $E_{\min}$ and $E_{\max}$ where

$$E_{\max/\min} = \frac{\sqrt{s}}{4} \left(1 - \frac{m_{\tilde{\chi}_1^0}^2}{m_{\tilde{\ell}}^2} \right) \left(1 \pm \left(1 - \frac{4m_{\tilde{\ell}}^2}{s} \right)^{1/2} \right), \quad (7.10)$$

which, with a known slepton mass, also gives a handle on the LSP mass $m_{\tilde{\chi}_1^0}$ even though it is undetected.

7.2.2 Neutralinos & charginos

For charginos and neutralinos the production cross section depends on their wino, bino and higgsino components. You would be forgiven to think that pair production of the lightest neutralino $e^+e^- \rightarrow \tilde{\chi}_1^0\tilde{\chi}_1^0$ would be the natural sparticle to search for, however, this has some significant problems. Since it is usually the LSP it does not decay, and there is nothing in the event that can actually be measured. We cannot use the missing energy as that requires an *imbalance* in momentum. Given sufficiently hard (energetic) radiation from either the initial electron or positron, a single photon recoiling against missing energy could potentially be measured, and this, so-called **mono-photon search**, was indeed a search channel for dark matter production at LEP. However, for neutralino dark matter this does not work all that well for other reasons. The $Z\tilde{\chi}_i^0\tilde{\chi}_j^0$ vertex shown in Fig. 7.6 has the Feynman rule

$$\frac{ig}{2\cos\theta_W}\gamma^\mu \left[(N_{i3}N_{j3}^* - N_{i4}N_{j4}^*) P_L - (N_{i3}^*N_{j3} - N_{i4}^*N_{j4}) P_R \right], \quad (7.11)$$

which depends only on the higgsino components of the neutralinos, N_{i3} and N_{i4} . This can be understood from the fact that there are no ZZZ or $Z\gamma\gamma$ vertices in the Standard Model that can be supersymmetrised, only a Zhh vertex. For the photon there is no tree level coupling to the neutralinos at all since there are no direct couplings between the Higgs and the photon in the Standard Model. Thus, only neutralinos with significant higgsino components can be

produced this way. To top it off, a light higgsino with a mass dominated by the μ parameter would have very similar values of N_{i3} and N_{i4} , thus canceling the coupling.

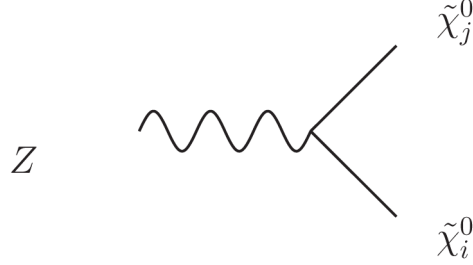


Figure 7.6: Vertex for $Z\tilde{\chi}_i^0\tilde{\chi}_j^0$.

The selectron and electron sneutrino have a special rôle for e^+e^- colliders due to the resulting t -channel diagrams. Figure 7.7 shows the t -channel diagrams that are important in pair production at a e^+e^- collider. Neutralino pair production with t -channel selectron exchange does not suffer from the same problems as neutralino pair production in the s -channel. However, the process depends on the selectron mass as $m_{\tilde{e}}^{-4}$ for large mass values and is rapidly suppressed.

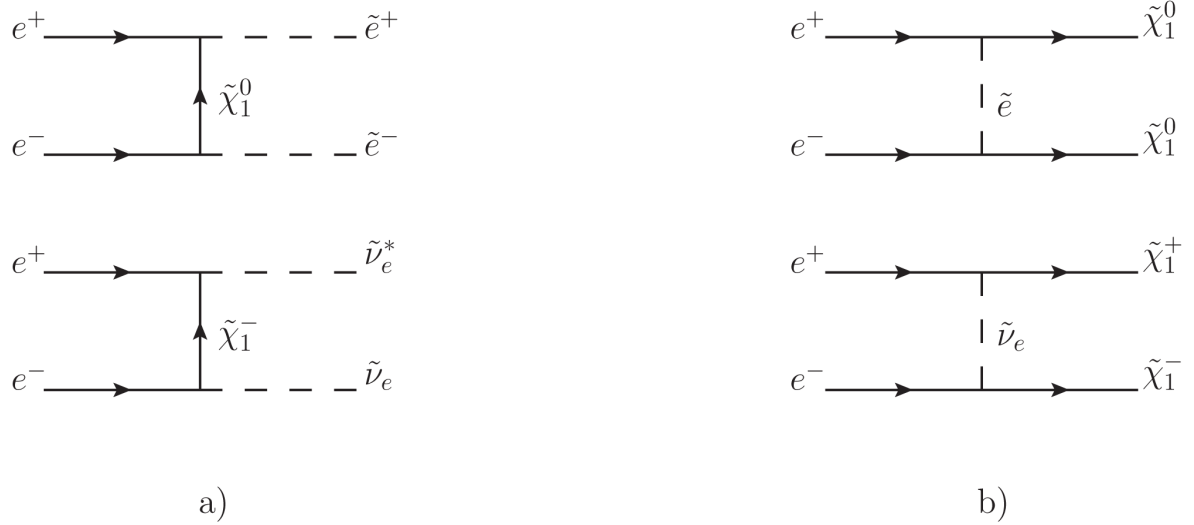


Figure 7.7: The t -channel diagrams for pair production of a) selectrons and electron sneutrinos, and b) neutralinos and charginos.

The consequences of the above discussion is that it is the production of the second lightest neutralino and the lightest chargino that is typically searched for in the combinations $\tilde{\chi}_2^0\tilde{\chi}_1^0$, $\tilde{\chi}_2^0\tilde{\chi}_1^0$, and $\tilde{\chi}_1^+\tilde{\chi}_1^-$. These decay – unless there is a slepton with a mass intermediary to the produced particle and the LSP – to the LSP and a possibly off-shell vector boson, $\tilde{\chi}_2^0 \rightarrow Z^{(*)}\tilde{\chi}_1^0$ and $\tilde{\chi}_1^\pm \rightarrow W^{(*)}\tilde{\chi}_1^0$. Thus the expected experimental signature is the fermionic decay products of the vector bosons with additional missing energy from the LSPs.

If the mass relevant differences between the neutralinos and charginos are small then the

resulting decay products of the off-shell vector bosons will have little energy – the technical term for this is that they are **soft**, as opposed to very energetic particles that are **hard**. Soft particles are difficult to reconstruct in detectors, and events dominated by soft particle production are easily missed, meaning that degenerate scenarios are difficult to discover, and sparticles mass bounds become very poor in this case. This is particularly unfortunate since a scenario where $\mu \ll M_1, M_2$ is well motivated on theoretical grounds, and with μ controlling two neutralino masses and one chargino, the lightest two neutralinos and the lightest chargino indeed become degenerate.

7.2.3 Higgs bosons

Because of the very small Yukawa coupling of electrons, Higgs bosons are not expected to have significant direct production cross sections at a e^+e^- collider. The most realistic production channels are Higgs-strahlung where the Higgs boson is attached to a vector boson: $e^+e^- \rightarrow Zh$, $e^+e^- \rightarrow ZH$, $e^+e^- \rightarrow Ah$ and $e^+e^- \rightarrow AH$, charged Higgs pair production $e^+e^- \rightarrow H^+H^-$, and W^+W^- vector boson fusion into Higgs bosons $e^+e^- \rightarrow \nu_e \bar{\nu}_e h(H)$.

7.2.4 Current bounds at lepton colliders

The below bounds are all from the LEP (Large Electron Positron) collider, running from 1989 until 2000, which outdated all previous bounds with a top energy of $\sqrt{s} = 209$ GeV, recording an integrated luminosity (amount of data) of 233 pb^{-1} above 204 GeV. Results exist from all four LEP experiments ALEPH, DELPHI, L3 and OPAL.² The numbers below are all taken from the 2014 PDG (Particle Data Group) review [19], but the conclusions from the LEP data has changed little since then.

While these bounds often come from pair-production of the relevant sparticles, and thus are less model dependent than the hadron collider bounds presented in the next section, there remains some model dependence in many results, which, unfortunately, is sometimes ignored in the literature. Complicating matters is a reliance by the LEP experiments on theoretical assumptions such as GUT-scale coupling and gaugino mass unification in many searches.

- Selectron: The strongest limit given is $m_{\tilde{e}_L} > 107$ GeV and $m_{\tilde{e}_R} > 73$ GeV (ALEPH 2002) in searches for acoplanar di-electrons.³ However, the limit is the result of a scan over MSSM parameter space simplified by assuming a common m_0 and $m_{1/2}$ at GUT scale. Interpreted in the even more constrained mSUGRA with $A_0 = 0$ the bounds are even stronger, 152 GeV and 95 GeV, respectively. In contrast, due to strict limits on the precisely measured Z -width, there is a fully model independent limit of $m_{\tilde{e}_{L/R}} > 40$ GeV.⁴
- Smuon: $m_{\tilde{\mu}_R} > 94$ GeV (DELPHI 2003). The limit is obtained as in the MSSM scenario for the selectron.
- Stau: $m_{\tilde{\tau}_1} > 81.9$ GeV (DELPHI 2003) assuming exclusive $\tilde{\tau}_1 \rightarrow \tau \tilde{\chi}_1^0$ decays and $m_{\tilde{\tau}_1} - m_{\tilde{\chi}_1^0} > 15$ GeV.

²Most of which are silly acronyms of course.

³The observant reader will notice that two electrons are always in the same plane, however, when experimentalists say acoplanar, they mean not in one plane with the beam axis.

⁴Similar model independent limits around half the Z -mass exists for all sparticles that couple to the Z .

- Sneutrinos: From LEP we have an indirect limit of $m_{\tilde{\nu}} > 94$ GeV (DELPHI 2003) in neutralino & slepton searches. This assumes $m_{\tilde{e}_R} - m_{\tilde{\chi}_1^0} > 10$ GeV. From the Z-width we can obtain the model independent limit $m_{\tilde{\nu}} > 44.7$ GeV.
- Neutralino: The limit $m_{\tilde{\chi}_1^0} > 46$ GeV (DELPHI 2003) has been derived from the direct searches for $\tilde{\chi}_1^0 \tilde{\chi}_2^0$ and $\tilde{\chi}_2^0 \tilde{\chi}_2^0$. This assumes gauge coupling unification and a common gaugino mass $m_{1/2}$ at GUT scale. Even in the Z-decays, the contribution to the width depends on the higgsino component of the lightest neutralino, so $m_{\tilde{\chi}_1^0} \simeq 0$ GeV is still in principle allowed [20].
- Here we have $m_{\tilde{\chi}_1^\pm} \geq 94$ GeV (DELPHI 2003), assuming GUT scale universality of m_0 and $m_{1/2}$ and using multiple direct search channels from production of charginos, neutralinos and sleptons. It also assumes either no third generation sfermion mixing or $m_{\tilde{\chi}_1^\pm} - m_{\tilde{\chi}_1^0} > 6$ GeV. From the Z-width we can extract a strict limit of $m_{\tilde{\chi}_1^\pm} \geq 45$ GeV.

7.3 Supersymmetry at hadron colliders

Since the particles collided at hadron colliders (protons and other nuclei) are heavier they are not as susceptible to losing energy through bremsstrahlung as the light leptons. Thus they can more readily be accelerated in a circular collider with periodic acceleration. However, this mass both means that linear colliders are inappropriate since the acceleration is smaller for the same applied fields, and that it is more difficult to steer the hadrons in a circular orbit. Thus the energy available at a hadron collider is limited by the strength of the bending magnets.

Up to 2009 the highest energy hadron collider was the Tevatron based at Fermilab outside of Chicago, that collided protons and anti-protons at $\sqrt{s} = 1.96$ TeV. The current record is held by the Large Hadron Collider (LHC) at CERN, colliding protons at a top energy of $\sqrt{s} = 13$ TeV.⁵

Since a proton-proton collider mostly collides partons, the quarks and gluons inside the protons, and since the strong coupling is large, this means that we can get large cross sections and potentially many events for QCD charged sparticles, *i.e.* squarks and gluinos, provided their masses are low enough. Cross sections for particles that have only electroweak charges (sleptons, electroweakinos, Higgs bosons) are expected to be considerably smaller, however, these sparticles are also expected to be lighter. The balance between these factors means that both types of sparticles are searched for.

As discussed earlier, with R-parity conservation sparticles are produced in pairs and both decay to the LSP. In Fig. 7.8 we show the sparticle mass spectrum for a particular mSUGRA benchmark point called SPS1a that has $m_0 = 100$ GeV, $m_{1/2} = 250$ GeV, $A_0 = -100$ GeV, $\tan \beta = 10$ and $\text{sgn } \mu = +$ [21].⁶ By producing the squarks or gluinos near the top of the mass spectrum in a collision we see that we can get many different quite long decay chains, called cascades. We illustrate such an event in Fig. 7.9.

In hadron collisions the momenta of the incoming partons is unknown, so no complete momentum balance can be made. However, the sum of the momenta in the direction transverse

⁵The LHC was designed to operate at $\sqrt{s} = 14$ TeV, but has had problems fully reaching that energy. In the next round of the LHC to start in 2022, called Run III, the plan is to go to $\sqrt{s} = 13.5$ TeV.

⁶This benchmark point is now quite old and clearly excluded by searches, but it still serves as a nice illustration of the kind of sparticle spectrum expected in mSUGRA models.

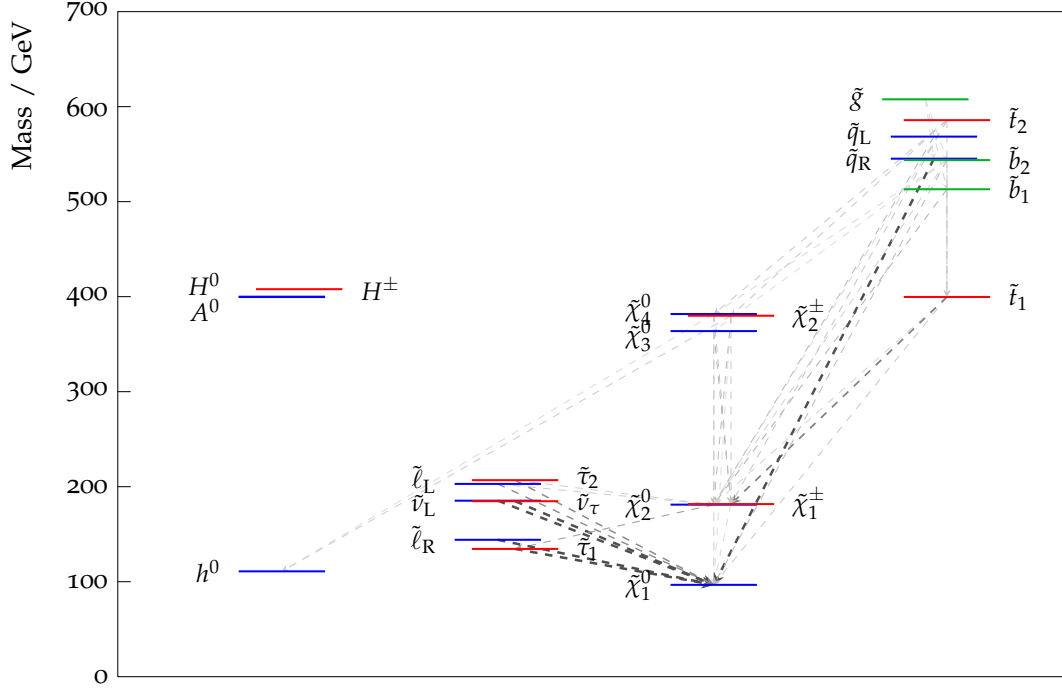


Figure 7.8: Supersymmetric particle mass spectrum (coloured solid lines) and possible decay channels (dashed lines) for the SPS1a benchmark point. Only decays with branching ratios above 5% are shown. The line opacity indicates relative branching ratios. This plot was generated using PySLHA 3.0.1 [22].

to the beam direction is zero. Thus with escaping LSPs we can observe **missing transverse energy**, $\cancel{E}_T$ or E_T^{miss} , *i.e.* an imbalance in the directional sum of all energy deposits transverse to the beam direction. The rest of the signal depends on which particles are searched for. In broad terms this means that if searching for gluino or squark production then at some point the cascade decay must shed colour charge to arrive at the LSP, resulting in Standard Model quarks or gluons that hadronise and produce jets in the detector (hadronic showers). If instead searching for electroweakly produced particles or Higgs bosons, one usually assumes an absence of significant jet activity, and instead looks for leptons from the decays of the sleptons or electroweakinos, possibly through vector bosons.

At the LHC the Standard Model backgrounds are much more significant and problematic. For example, at design luminosity the LHC produces about 10 pairs of top quarks per second. Due to the top quark decay $t \rightarrow bW^+$, resulting in charged leptons and potential missing energy from neutrinos in W decays, as well as jets produced in hadronic W decays and from the b -quarks, this alone is a significant background to many supersymmetry searches. Figure 7.10 shows the expected backgrounds and signals produced in different channels at the $\sqrt{s} = 14$ TeV LHC for different assumed particle masses. One can see that even the largest supersymmetric cross sections that we get for squark and gluino production are orders of magnitude below the Standard Model backgrounds for masses beyond 500 GeV.

To cope with this challenge, the searches define kinematical variables that are designed to separate supersymmetry events from Standard Model events. For example, searching for the production of squarks and/or gluinos decaying to the LSP, one expects jets and missing

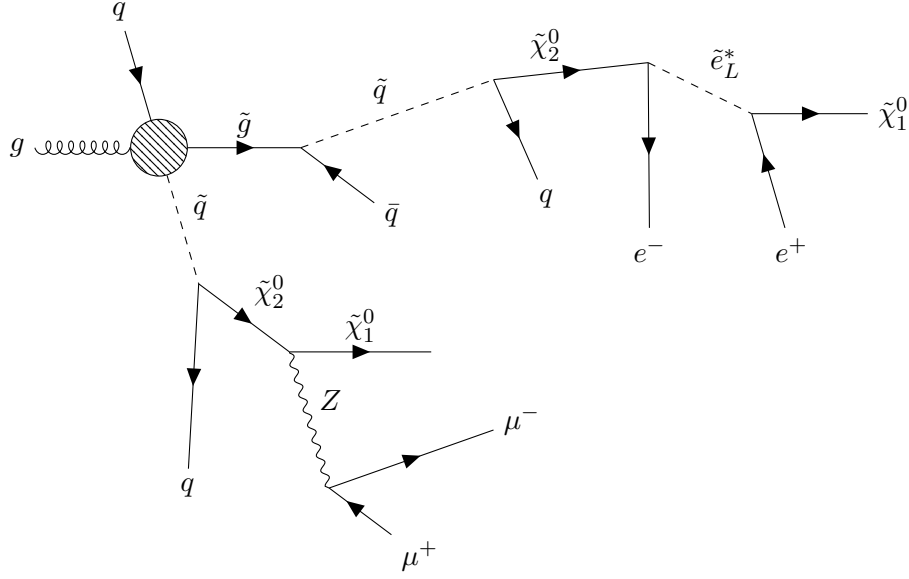


Figure 7.9: Cascade decays starting with the production of a gluino and squark pair from a gluon–quark interaction.

energy from two LSPs. One can then define the **effective mass**

$$M_{\text{eff}} = \sum p_T^{\text{jet}} + \cancel{E}_T, \quad (7.12)$$

from the sum of the transverse momentum of the jets in the event, p_T , and the missing energy. The expectation is then that this is more significant in squark and gluino production events than in Standard Model events, so that a high requirement on M_{eff} reduces the number of background events without significantly removing events with sparticles. This helps for example to combat $t\bar{t}$ production because the jet p_T s in those events are limited by the energy in the top quark mass, while the more massive squarks and gluinos will have higher jet p_T s.

However, there are also models where this kind of approach is ineffective. Imagine a scenario where only the lightest stop $\tilde{t}_1$ is light enough to be copiously produced. If $m_{\tilde{t}_1} - m_{\tilde{\chi}_1^0} < m_W$ then the stop will dominantly decay as $\tilde{t}_1 \rightarrow c\tilde{\chi}_1^0$ or $\tilde{t}_1 \rightarrow b\nu\tilde{\chi}_1^0$, where all final state particles have rather low energy (p_T), so-called **soft particles**. As for the degenerate neutralino scenario discussed above, this is very difficult to discover with standard techniques.

Another example of the use of kinematics is the distribution of invariant masses. In Fig. 7.11 we show an example of two sequential two-body decays. Even if particle A here is invisible, for example the LSP, we can use two visible decay products a and b to form the invariant mass m_{ab} . One can show that the distribution for m_{ab} has a triangular shape with a sharp endpoint at the maximum

$$(m_{ab}^{\text{max}})^2 = \frac{(m_C^2 - m_B^2)(m_B^2 - m_A^2)}{m_B^2}, \quad (7.13)$$

where we have assumed that a and b are massless.⁷ This can be used both to select events that potentially have supersymmetric particles in them depending on the value of the invariant

⁷A more complicated expression covers the massive case.

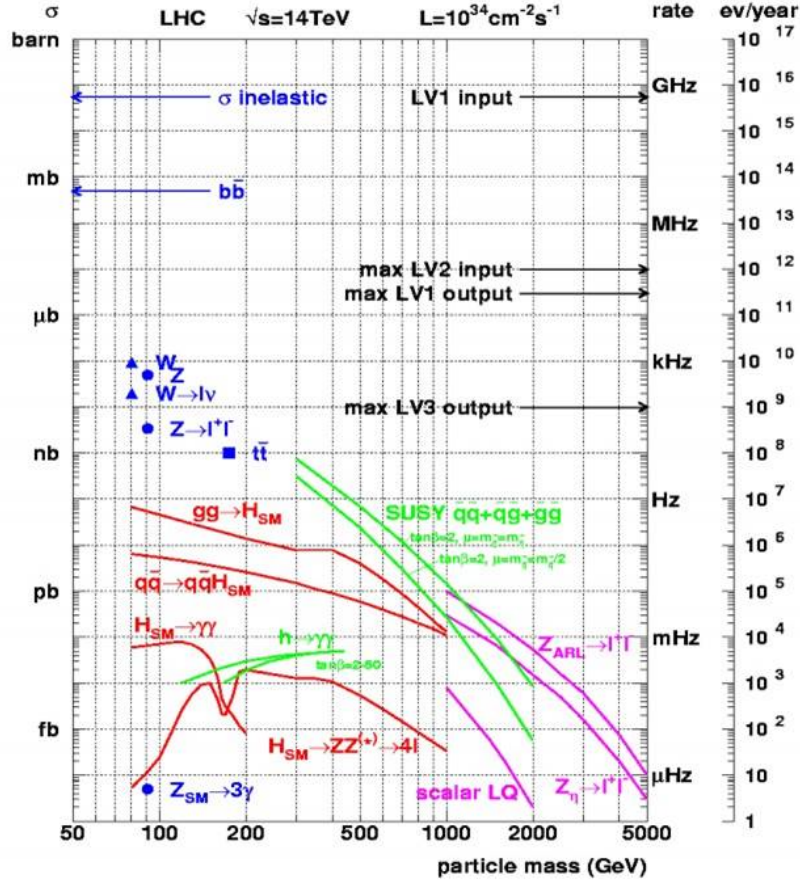


Figure 7.10: Plot of the expected cross sections and rates for various processes at the $\sqrt{s} = 14$ TeV LHC plotted against the mass of the particles, assuming the design final luminosity of the LHC (rate of data taking). The current Run II of the LHC has collected around 140 fb^{-1} of data, so the lowest point on the cross section axis indicates that $\mathcal{O}(1)$ events are expected to have been produced by now.

mass, and to determine relationships between the particle masses. In Fig. 7.12 we show a simulation of the invariant mass distribution of two opposite-sign same-flavour (OSSF) leptons $m_{\ell\ell}$ from the production of $\tilde{\chi}_2^0$ and its decay chain $\tilde{\chi}_2^0 \rightarrow \ell^\pm \tilde{\ell}_R^\mp \rightarrow \ell^\pm \ell^\mp \tilde{\chi}_1^0$.

As alternatives to these standard searches for pair produced particles with missing energy there are ongoing searches for decaying LSPs when R-parity is violated, or the production of single particles.⁸ There is also the possibility of **massive metastable charged particles** (MMCPs), typically in scenarios with a gravitino LSP, where the next-to-lightest supersymmetric particle (NLSP) is charged and long-lived because the decay to the gravitino is via a very weak gravitational coupling. The latter also includes so-called **R-hadrons** if the NLSP has colour charge, which means that it will hadronise after production and be a short-lived but very massive meson or baryon. We should also mention the searches for the extra Higgs

⁸Single particle production at the LHC requires rather large R-parity violating couplings for the $LQ\bar{D}$ or $\bar{U}\bar{D}\bar{D}$ operators, of the order of $\lambda > 10^{-2}$.

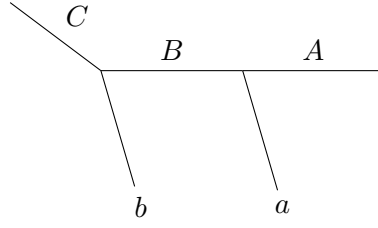


Figure 7.11: A generic illustration of two successive two-body decays $C \rightarrow bB$ and $B \rightarrow aA$.

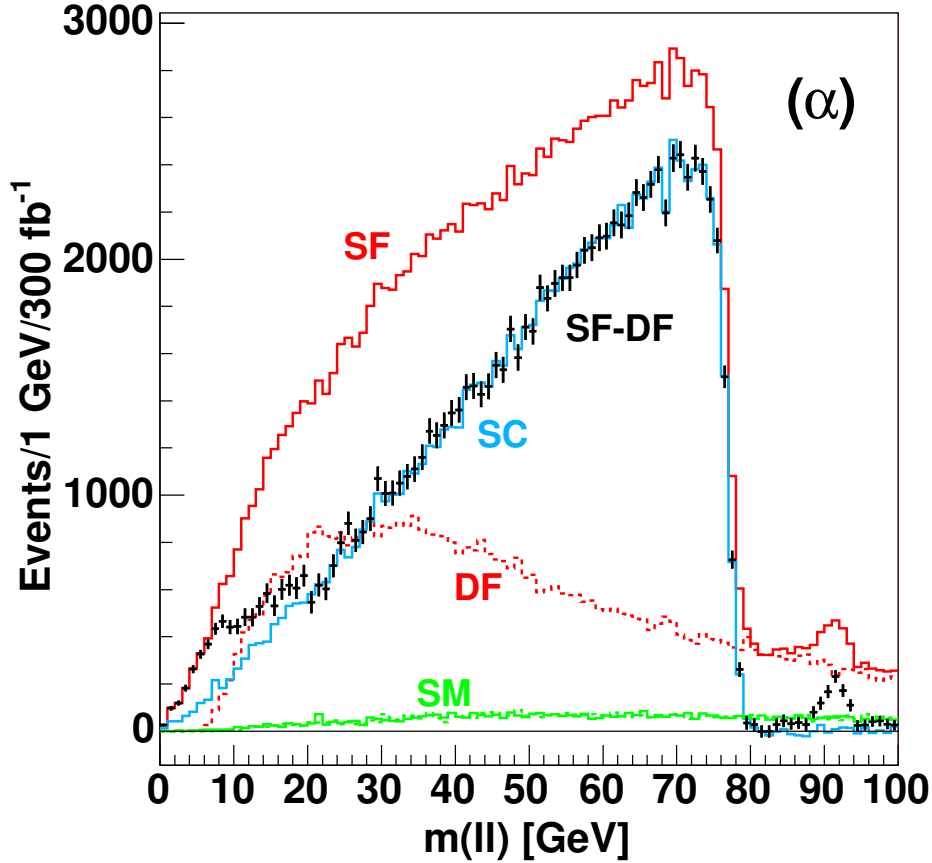


Figure 7.12: Invariant mass distribution of opposite sign same flavour (OSSF) dileptons for the mSUGRA benchmark model point SPS1a [23].

states predicted in the MSSM.⁹

7.4 Current bounds on sparticle masses

The LHC has finished its Run II and collected a total of around 140 fb⁻¹ of data per experiment at $\sqrt{s} = 13$ TeV of energy. It is currently preparing for Run III to start in 2022. Direct bounds from the LHC experiments ATLAS and CMS now supersede bounds from other colliders (Tevatron and LEP) in almost all channels, with some exceptions for models

⁹But we don't really have time.

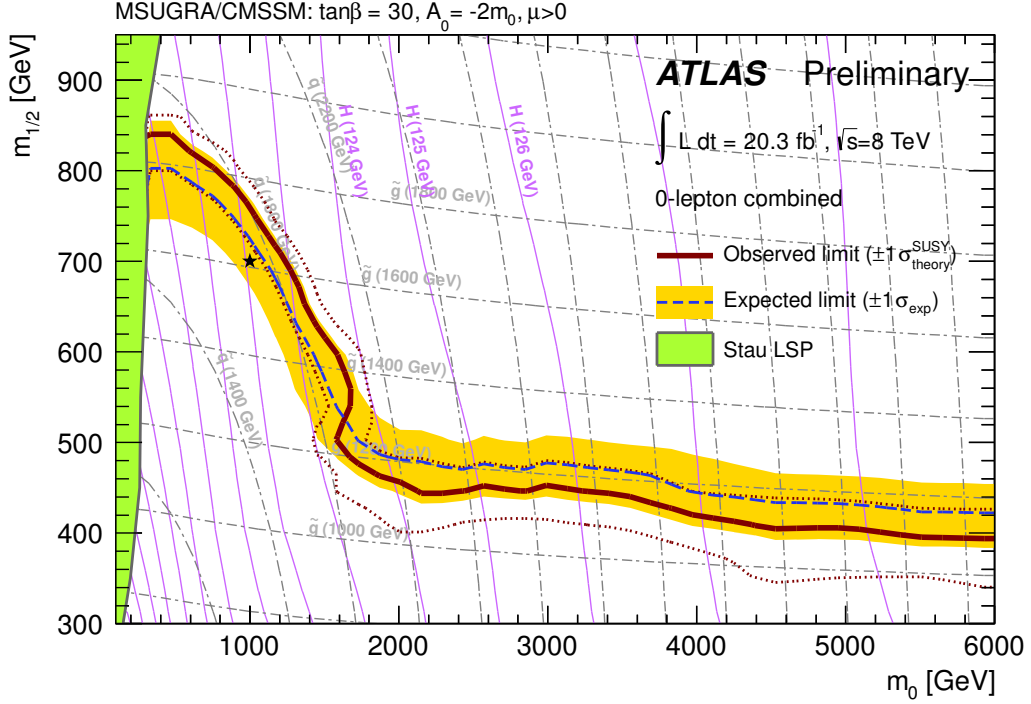


Figure 7.13: Plot of the excluded area in the $m_{1/2}$ - m_0 plane of the mSUGRA parameter space for $\tan\beta = 30$, $A_0 = -2m_0$ and $\mu > 0$ for searches using missing energy and jets. The limit is the red line. The green area is theoretically forbidden because it has a charged LSP (the stau) [24].

with degenerate masses. The below limits are mostly limits with the full Run II dataset and represent the current state-of-art in sparticle searches. We mostly use examples from ATLAS, the corresponding plots from CMS are very similar.

The strongest current limits in terms of mass are on the gluino and squarks simply because of the large production cross sections. Significant bounds on electroweakinos and sleptons exist, but these are either model dependent (depend on squark/gluino mass assumptions and cascade decays), or weaker if they rely only on electroweak production.

7.4.1 Squarks and gluinos

In Fig. 7.13 we show the limits from the ATLAS experiment on the mSUGRA model using searches for jets plus missing energy with all the data collected at 8 TeV. The mSUGRA parameters $\tan\beta$ and A_0 have been chosen in order to give relatively large Higgs masses for small values of $m_{1/2}$ and m_0 . The figure also shows the corresponding first and second generation squark masses, the gluino mass (both dot-dashed lines), and the Higgs mass (purple) for these parameter values. We then have the following approximate bounds in mSUGRA: $m_{\tilde{q}} > 1600$ GeV and $m_{\tilde{g}} > 1100$ GeV. Bounds on the mSUGRA space directly have not been updated since this plot.

Notice that in the figure the squark mass bound is more or less equivalent to the mass required for a sufficiently heavy Higgs, thus the direct search does not constrain the squarks

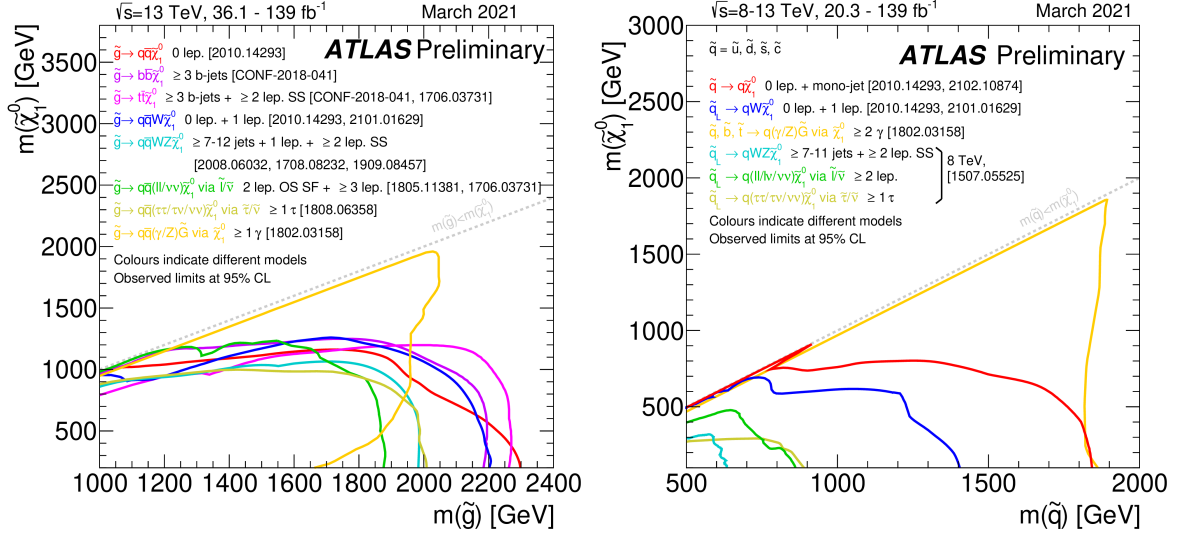


Figure 7.14: Excluded regions in the $(m_{\tilde{g}}, m_{\tilde{\chi}_1^0})$ (left) and $(m_{\tilde{q}}, m_{\tilde{\chi}_1^0})$ (right) mass planes under different assumptions on the decays of the gluino and squarks [25]. References for the individual analysis given in the figure.

masses significantly more here than the indirect constraint from the Higgs mass.

An important question is how these bounds change as we move away from the mSUGRA assumptions. In mSUGRA the resulting gluino and squark decays are very much constrained by the model. For more general models the limits depend on what decay chains are dominant. In Fig. 7.14 we show current limits in the gluino–lightest neutralino and squark–lightest neutralino mass planes under various assumptions on the decay chain shown in different colours.

For the gluino the most minimal assumption we can make is that the gluino sheds its adjoint colour charge in the decay to two quarks of the first or second generation and a neutralino, $\tilde{g} \rightarrow q\bar{q}\tilde{\chi}_1^0$, either by a direct three-body decay or via an intermediary squark (red line). We then get limits up to $m_{\tilde{g}} > 2300$ GeV under the assumption of a very light neutralino. This limit gradually weakens as the neutralino mass increases, and disappears beyond $m_{\tilde{\chi}_1^0} > 1000$ GeV for $m_{\tilde{g}} > 1000$ GeV. For $m_{\tilde{g}} < 1000$ GeV (not visible on the plot) most of the parameter space is excluded, with the exception that a small sliver remains when $m_{\tilde{g}} \simeq m_{\tilde{\chi}_1^0}$ in the degenerate scenario.

For the first two generations of squarks the most minimal decay is $\tilde{q} \rightarrow q\tilde{\chi}_1^0$. Here the bound (red) reaches to around $m_{\tilde{q}} > 1800$ GeV for a light neutralino, while for heavier neutralinos the bound disappears beyond $m_{\tilde{\chi}_1^0} > 700 - 800$ GeV for $m_{\tilde{q}} > 800$ GeV. Below a squark mass of around 800 GeV most of the parameter space is excluded, but we again lose sensitivity when there is degeneracy between the squark and the neutralino. Here a search for mono-jets has some impact in closing the gap, which shows as a spike in the excluded region between 800 and 900 GeV, however, despite poor visibility in the plot, a very degenerate squark–neutralino pair is still allowed below 800 GeV.

An additional assumption made in this plot is that all the eight squarks of the first two-generations are degenerate in mass and adding to the cross section. Should one squark

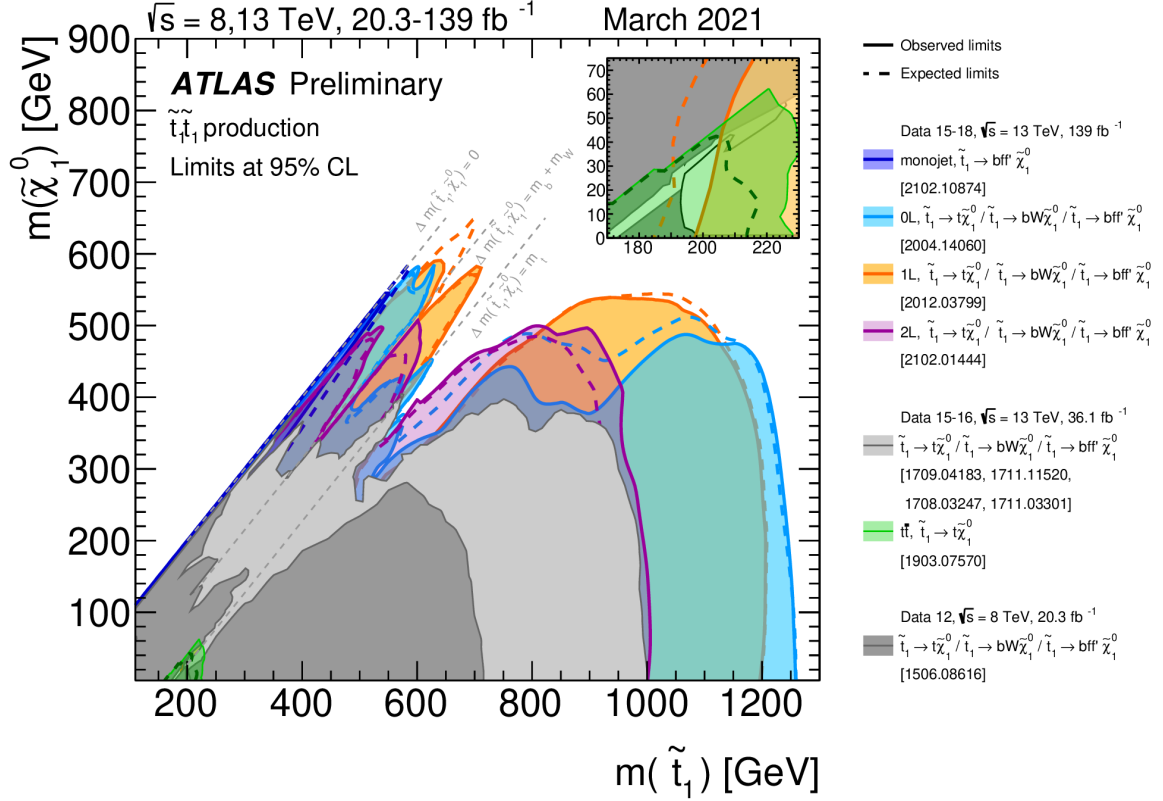


Figure 7.15: Plot of the excluded area in the $(m_{\tilde{t}_1}, m_{\tilde{\chi}_1^0})$ mass plane [25]. References for the individual analysis given in the figure.

generation or flavour be significantly lighter than the others this means a further reduction in the production cross section and thus a weaker bound. It is also fairly clear that removing R-parity, meaning that the LSP decays, also weakens the above conclusions due to the possible absence of significant missing energy.

7.4.2 Stop

For the stop there are many possible competing decay channels, meaning that limits set are rather model dependent. The two main decay categories for the lightest stop are decays via the chargino, $\tilde{t}_1 \rightarrow b\tilde{\chi}_1^+$, a supersymmetrised version of the Standard Model $t \rightarrow bW^+$, and decays directly to the neutralino $\tilde{t}_1 \rightarrow t\tilde{\chi}_1^0/bW\tilde{\chi}_1^0/bff\tilde{\chi}_1^0$, where ff' represents the fermions in a W decay, and where the dominant decay mode depends on the stop–neutralino mass difference. Since the chargino decay is also typically $\tilde{\chi}_1^+ \rightarrow ff'\tilde{\chi}_1^0$, the experimental signature of stop production is $bff'\tilde{\chi}_1^0$, a mixture of missing energy, fermions that may either be leptons or result in jets, and a b -quark that forms a jet with distinct properties, called a b -jet.

A summary of (the many) ATLAS searches for the stop is found in Fig. 7.15 showing limits in the $(m_{\tilde{t}_1}, m_{\tilde{\chi}_1^0})$ mass plane. The dashed grey lines show the regions where the different stop decays are kinematically possible. We observe that for light neutralinos the limit goes all the

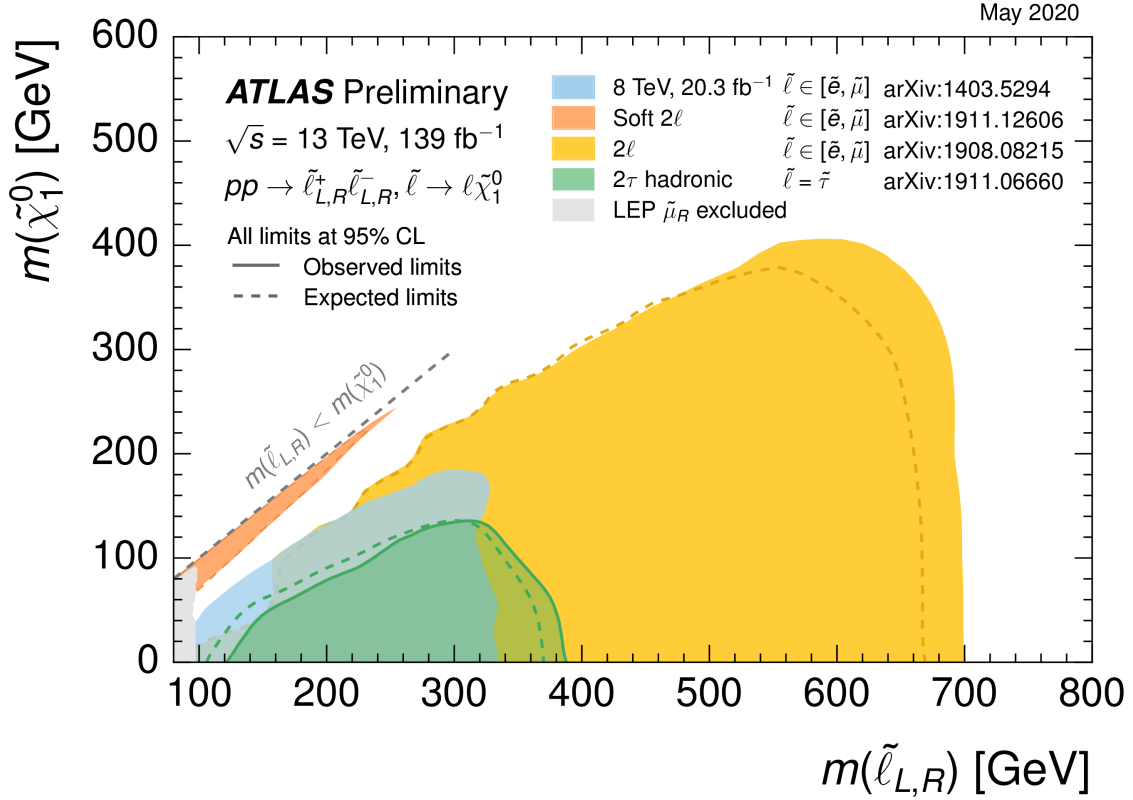


Figure 7.16: Plot of the excluded area in the $(m_{\tilde{\ell}_{L,R}}, m_{\tilde{\chi}_1^0})$ plane for mass degenerate right- and left-handed sleptons of different flavours [25].

way up to $m_{\tilde{t}_1} > 1250 \text{ GeV}$, while there are essentially no limits for $m_{\tilde{\chi}_1^0} > 580 \text{ GeV}$. Again degenerate scenarios have weaker bounds. When the stop approaches the neutralino plus top mass stop masses down to 600 GeV are allowed. And, again not very visible, for a degenerate stop-neutralino the limit is virtually non-existent. As an additional complication, if the stop-neutralino mass difference is below the bottom quark mass, the standard stop decays no longer work and the flavour changing neutral current decay $\tilde{t}_1 \rightarrow c \tilde{\chi}_1^0$ becomes dominant. This decay has multiple still open theoretical questions, both about the potentially long lifetime of the stop in this scenario, and exactly where in parameter space this transition occurs.

7.4.3 Sleptons

The mass bounds on sleptons will be very dependent on the assumed production mechanism. If the sleptons are produced indirectly in cascade decays from heavier squarks or gluinos they could have large cross sections, however, the most model independent bounds come from assuming only direct electroweak pair production through a virtual photon or Z .

The result for degenerate right- and left-handed charged sleptons from electroweak production, assuming decays with 100% branching ratio to the lightest neutralino, are shown in Fig. 7.16. These limits separate between selectron and smuon production, and stau production. The former are assumed degenerate, and have stricter bounds than the stau, which is harder to reconstruct due to the many different possible tau decays involving extra neutri-

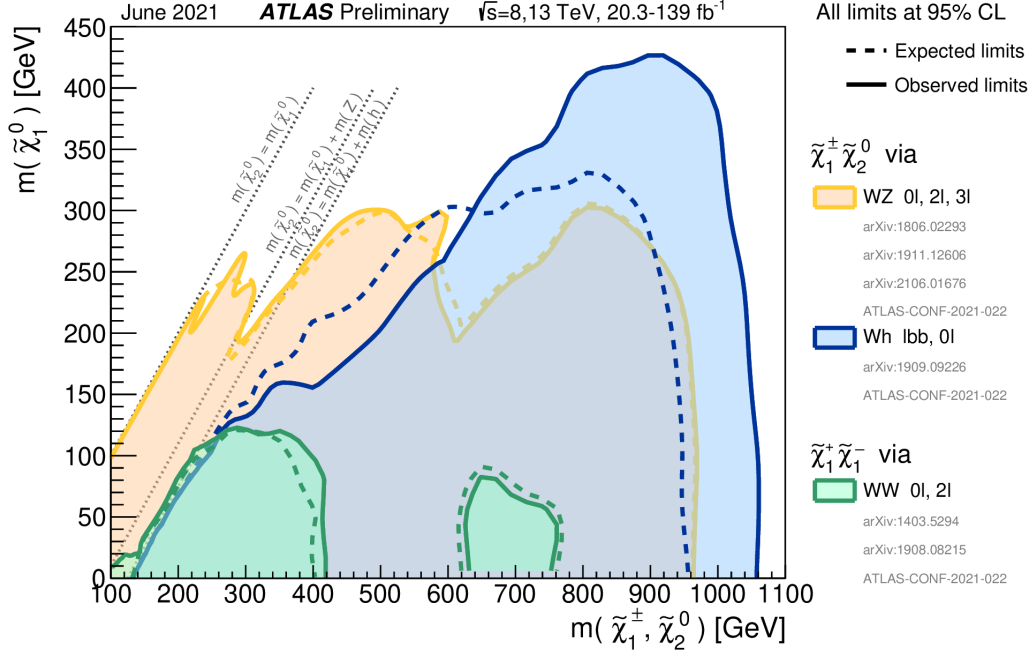


Figure 7.17: Plot of the excluded area in the $(m_{\tilde{\chi}_1^\pm, \tilde{\chi}_2^0}, m_{\tilde{\chi}_1^0})$ mass plane for different search signatures [25]. This plot assumes a degenerate wino-like $\tilde{\chi}_2^0$ and $\tilde{\chi}_1^\pm$ and a 100% branching ratio to the given decay channels.

nos. For all slepton flavours we see that there is a gap in the mass plane down to very low masses where no sleptons are excluded, all the way down to masses below 100 GeV where the LEP bound applies. This is yet another example of the problems with degeneracy. If the slepton–neutralino mass difference is around 50 GeV and lower, it becomes difficult to reliably reconstruct the soft leptons from the slepton decay and to separate them from the ordinary Standard Model pair production of lepton pairs.

7.4.4 Charginos and neutralinos

As for the sleptons, bounds are dependent on the production process assumed. The search for direct electroweak pair production of the lightest neutralino, $\tilde{\chi}_1^0 \tilde{\chi}_1^0$, has the same problem as at LEP, the coupling is vanishingly small for a non-higgsino $\tilde{\chi}_1^0$. This means that it is the heavier neutralinos and charginos that are typically searched for. Here some considerations on the possible masses hierarchies come into play. If the lightest neutralino is dominantly bino and the next-to-lightest dominantly wino ($M_1 < M_2 < |\mu|$) then the $\tilde{\chi}_2^0$ and $\tilde{\chi}_1^\pm$ states are degenerate in mass and the most important searches will be for $\tilde{\chi}_2^0 \tilde{\chi}_1^\pm$ and $\tilde{\chi}_1^\pm \tilde{\chi}_1^\mp$ production.¹⁰ These then decay as $\tilde{\chi}_2^0 \rightarrow Z \tilde{\chi}_1^0$ or $\tilde{\chi}_2^0 \rightarrow h \tilde{\chi}_1^0$, and $\tilde{\chi}_1^\pm \rightarrow W \tilde{\chi}_1^0$, where the bosons may be off-shell if the mass difference is small. This is the scenario usually considered by the experiments, in part because winos give the largest electroweak production cross sections.

¹⁰Again because of the coupling in Eq. (7.11) the production of a pair of wino-like $\tilde{\chi}_2^0 \tilde{\chi}_2^0$ is suppressed and the cross section is small.

We show the current limits assuming degenerate wino-like states $\tilde{\chi}_2^0$ and $\tilde{\chi}_1^\pm$ in Fig. 7.17. We can see that limits up to 1050 GeV can be set for $\tilde{\chi}_2^0$ and $\tilde{\chi}_1^\pm$ when the lightest neutralino is below around 200 GeV. When the bosons in the decay go off-shell we see that the sensitivity of the searches go down significantly. If we move away from the wino-production scenario the production cross sections drop, and so does the reach of the experiments in the mass plane.

An additional important complication appears if the lightest neutralino is wino ($M_2 < M_1, |\mu|$) or higgsino ($|\mu| < M_1, M_2$). Then the lightest neutralino and chargino, $\tilde{\chi}_1^0$ and $\tilde{\chi}_1^\pm$ (for higgsinos also $\tilde{\chi}_2^0$), are degenerate, and the crucial question becomes how degenerate they are. If the mass difference is small but above around 300 MeV, then the particles are very difficult to discover at all, despite potentially having very large cross sections if they are light. The decays of the produced $\tilde{\chi}_1^\pm$ (and for higgsinos $\tilde{\chi}_2^0$) into $\tilde{\chi}_1^0$ lead to very soft decay products, either leptons in $\tilde{\chi}_1^\pm \rightarrow \ell^\pm \nu_\ell \tilde{\chi}_1^0$ (and $\tilde{\chi}_2^0 \rightarrow \ell^+ \ell^- \tilde{\chi}_1^0$), or decays to pions (the lightest hadronic states) $\tilde{\chi}_1^\pm \rightarrow \pi^\pm \tilde{\chi}_1^0$ (and $\tilde{\chi}_2^0 \rightarrow \pi^0 \tilde{\chi}_1^0$), that are unobservable above the backgrounds of a hadron collider.

If instead the mass difference becomes of the order off or less than the pion mass the width of the chargino can become so small that the state is sufficiently long lived to decay somewhere inside the detector, and not, effectively, at the interaction point. This would lead to charged tracks from the charginos that disappear somewhere inside the detector. Such disappearing track signatures are readily observable, and in fact the limits on very degenerate neutralinos and charginos are stronger than the limits on only somewhat degenerate ones.

The consequences of this complication is that it is very difficult to exclude near massless wino- or higgsino-like charginos and neutralinos at the LHC. However, since the chargino mass is constrained by the Z -width measured at LEP, there is an absolute lower limit of 45 GeV for both in this scenario. To improve significantly on this limit one may have to wait for a new lepton collider.

7.5 Precision observables

A different way to search for the signs of supersymmetric particles is their indirect effect on very accurately measured Standard Model particle properties and processes, so-called **precision observables**, mostly through loop diagrams with sparticles. We will here discuss four different types of probes: the electroweak precision observables, the value of the anomalous magnetic moment of the muon $(g-2)_\mu$, the flavour changing neutral current (FCNC) process $b \rightarrow s\gamma$ and the very rare (and FCNC) process $B_s \rightarrow \mu\mu$.

7.5.1 Electroweak precision observables

When we talk about **electroweak precision observables**, we study particle properties such as M_W (or M_Z), Γ_W , Γ_Z , m_t and $\sin\theta_W$, as well as the Higgs mass m_h and the properties of the Higgs such as its couplings to all the other particles (gauge and Yukawa couplings) and its self-coupling.

Up to 2012 we could study all of these as *functions* of the unknown Higgs mass, looking for deviations that could be a sign of for example supersymmetry. We show a fit of the Higgs mass to all available electroweak data and direct exclusion bounds in Fig. 7.18, made by the Gfitter Collaboration in 2011 just before the LHC started taking significant amounts of data [26], a fit pretty much indicating that the most probable Standard Model Higgs mass

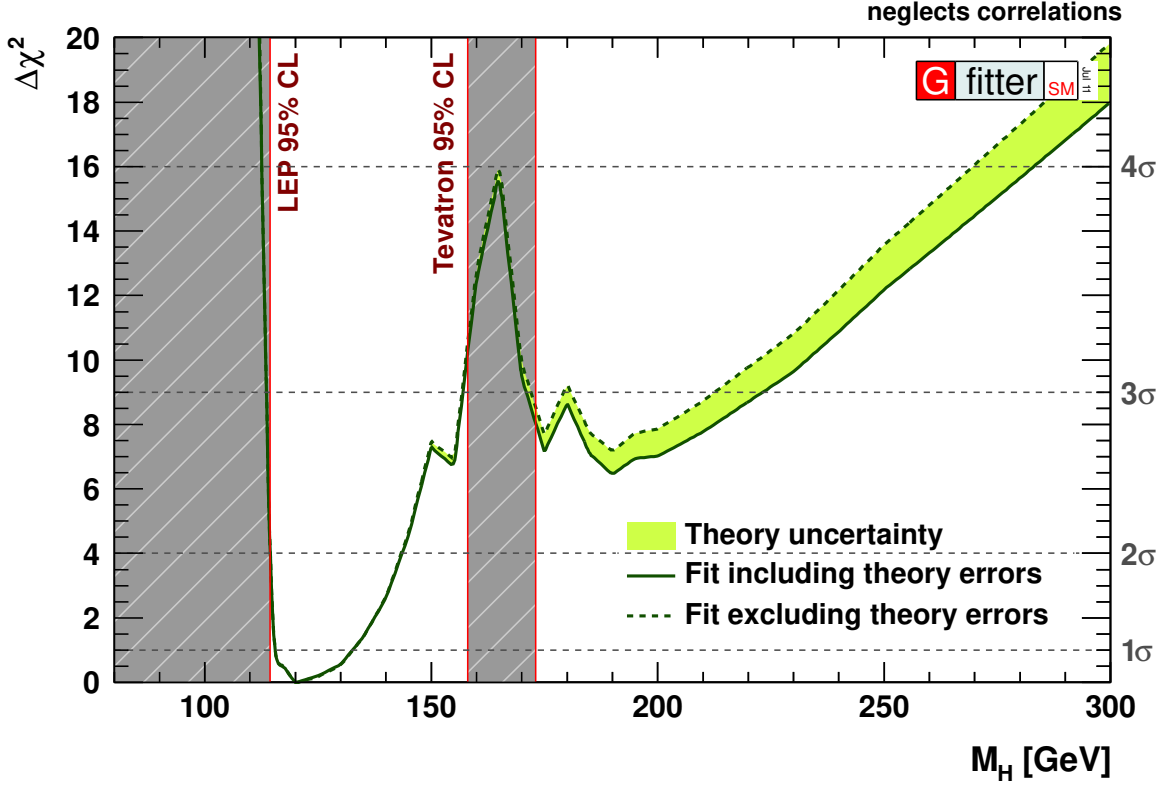


Figure 7.18: Plot of the total $\Delta\chi^2$ from all precision observable measurements and the direct exclusions bounds for the Standard Model Higgs set by LEP and the Tevatron, as a function of the Higgs mass [26].

was around 120 GeV. Here,

$$\Delta\chi^2 = \chi^2(m_h) - \min_{m_h} \chi^2(m_h),$$

is the difference in the χ^2 from the best fit point as a function of the Higgs mass, which has a standard interpretation in terms of σ s of significance for a one dimensional fit.

Figure 7.19 shows a similar plot for mSUGRA. At that time the absolute minimum of the fit, even taking into account the different number of parameters, gave a better fit for mSUGRA, $\min \chi^2_{\text{mSUGRA}} < \min \chi^2_{\text{SM}}$, but this changed quickly when the Higgs was found because of the position of the two minima.

Now all the parameters of the Standard Model – neutrinos excepted – have been determined to some precision. Thus the Standard Model is a completely constrained system. If we now do a electroweak fit the situation looks like that in Fig. 7.20, where we show the global fit to all measurements *except* the top and W masses, compared to the measured values of the W and top masses. Clearly what we are seeing here is (still?) consistent with the Standard Model.

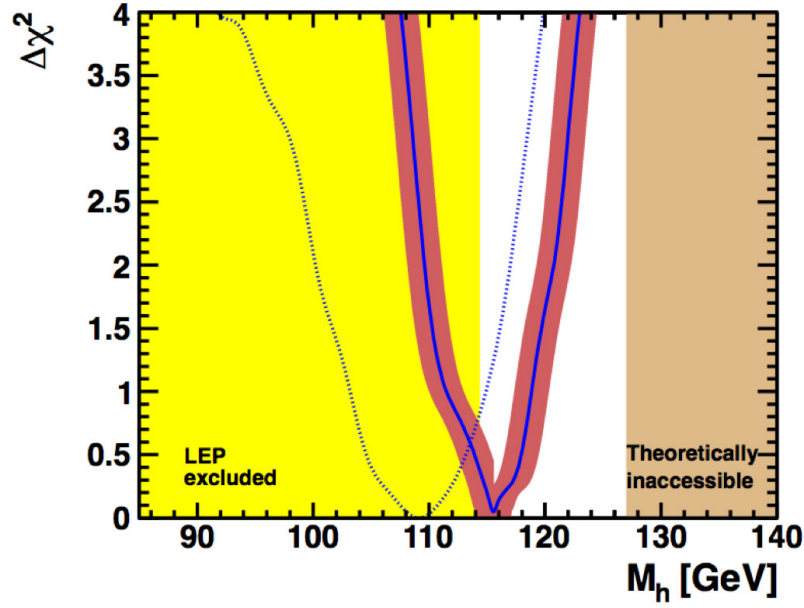


Figure 7.19: Plot of the $\Delta\chi^2$ from all precision observable measurements for mSUGRA as a function of the Higgs mass. The yellow area shows the area excluded by LEP searches (*not* included in the fit), while the brown shows the theoretically inaccessible area.

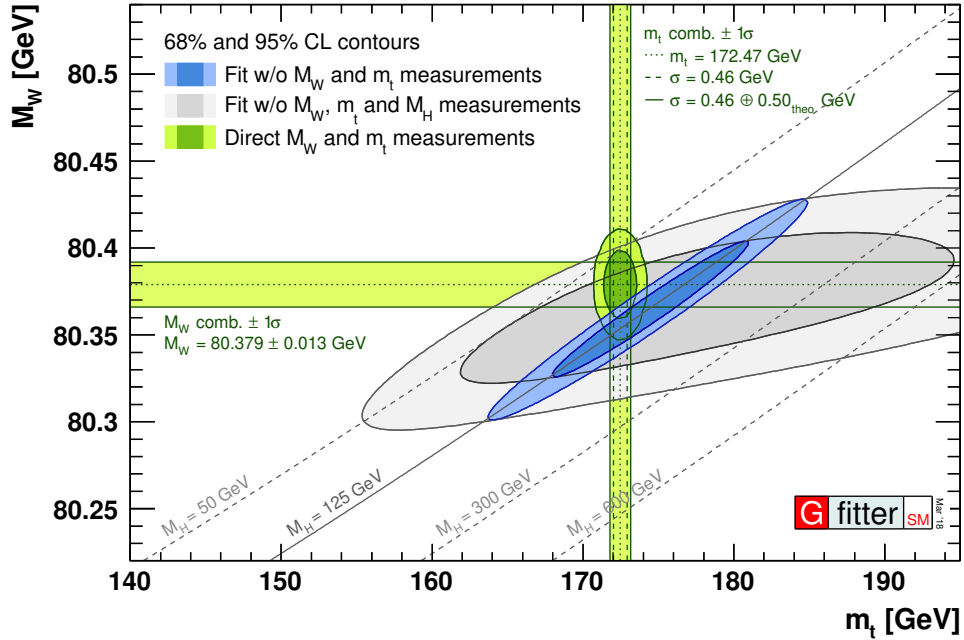


Figure 7.20: Electroweak fit excluding M_W and m_t (blue), compared to their measured values (green) [27].

7.5.2 $(g - 2)_\mu$

The **anomalous magnetic moment** of the muon, g_μ , in effect the interaction of a muon with an external electromagnetic field shown in Fig. 7.21 a), was calculated by Dirac in relativistic quantum mechanics as $g_\mu = 2$. However, in quantum field theory we have higher order corrections from loop-process such as shown in Fig. 7.21 b), that makes g_μ deviate from two.

It was very precisely measured by the E821 experiment at BNL [28] to be:

$$g_\mu^{\text{exp}} = 2.00116592089(63),$$

or, in terms of the deviation $a_\mu = g_\mu - 2$,

$$a_\mu^{\text{exp}} = 11659208.9(6.3) \cdot 10^{-10},$$

where the parenthesis indicates the uncertainty on the last digits.

In the Standard Model we find the prediction

$$a_\mu^{\text{SM}} = 11659183.0(5.1) \cdot 10^{-10},$$

giving a difference with respect to the experimental value of

$$\Delta a_\mu \equiv a_\mu^{\text{exp}} - a_\mu^{\text{SM}} = (25.9 \pm 8.1) \cdot 10^{-10},$$

a value which is 3.2σ away from zero. Recently, the Muon g-2 Collaboration at Fermilab made a new measurement [29] confirming this discrepancy with the value

$$\Delta a_\mu = 25.1 \pm 5.9 \cdot 10^{-10},$$

corresponding to 4.2σ . The Muon g-2 Collaboration is set to make further improvements on this number, expecting to reduce the uncertainty by a factor of 2 – 4.

However, we should be aware that some of the Standard Model contributions involve hadronic loops, *e.g.* the so-called **hadronic vacuum polarisation** shown in Figure 7.21 b), where one has to rely on experimental information from low energy $e^+e^- \rightarrow \gamma^* \rightarrow \text{hadrons}$ in order to estimate a contribution of $a_\mu^{\text{HVP}} = 10.5(2.6) \cdot 10^{-10}$, which is of the same order of magnitude as the discrepancy, and may be prone to errors in the interpretation.

In the MSSM we can have contributions to a_μ at the one-loop level, either by the exchange of one of the new Higgs bosons across the muon line, or by loops containing a smuon $\tilde{\mu}$ or muon-sneutrino $\tilde{\nu}_\mu$, together with a neutralino or chargino. The diagrams for the latter two processes are shown in Figure 7.21 c) and d). These contribute opposite sign terms $a_\mu(\tilde{\chi}^0)$ and $a_\mu(\tilde{\chi}^-)$. A thorough analysis shows that we need $\mu > 0$ in order to give a positive contribution that will close the gap between the experimental value and the prediction. In order to get a sufficiently large contribution the loop masses must be less than 500 – 600 GeV for $\tan \beta = 40 - 50$ and 200 – 300 GeV for $\tan \beta \simeq 10$. However, as we saw in Sec. 7.4 this is not implausible.

7.5.3 $b \rightarrow s\gamma$

The quark level process $b \rightarrow s\gamma$ is a Flavour Changing Neutral Current (FCNC) process which must proceed through loops in the Standard Model since there are no tree-level FCNC

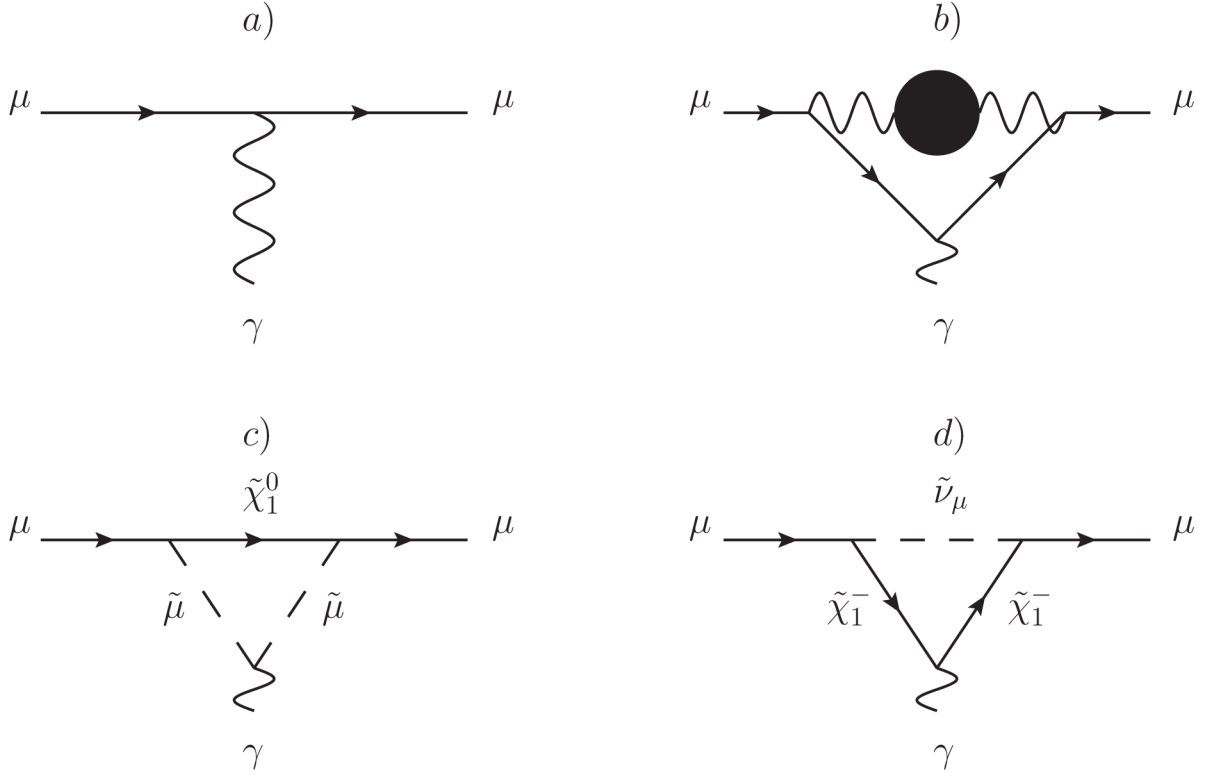


Figure 7.21: Diagrams for muon interaction with an electromagnetic field. Loop corrections to the tree level diagram a) give the value of a_μ . Diagram b) shows hadronic vacuum polarization where the blob contains QCD fields. Diagrams c) and d) show the lowest order MSSM contributions to a_μ .

interactions there. At meson level it leads to measurable decays of the type $B \rightarrow X_s \gamma$, where X_s is some meson with a strange quark, *e.g.* the decay $B \rightarrow K \gamma$. Figure 7.22 a) shows the one-loop Standard Model contribution with a virtual W and up-type quarks. This contribution is suppressed by the smallness of the CKM entries in the W -vertices which favours diagrams with a top quark in the loop, and the large masses, M_W and m_t , in the loop.

The process has been calculated at next-to-next-to-leading order (NNLO) in the Standard Model to be $\text{Br}(B \rightarrow X_s \gamma)_{\text{SM}} = (3.36 \pm 0.23) \cdot 10^{-4}$ for $E_\gamma \geq 1.6 \text{ GeV}$ [30, 31], inclusively summing over all meson final states.¹¹

Supersymmetric particles may contribute to this process, for example with off diagonal soft mass terms in diagrams such as Fig. 7.22 b) and d), where in b) a $m_{bs}^2 \tilde{b}^* \tilde{s}$ soft-breaking off-diagonal Lagrangian mass term changes a $\tilde{b}_1$ squark to a $\tilde{s}$ squark, and in d) where a $m_{tc}^2 \tilde{t}^* \tilde{c}$ term changes a $\tilde{t}_1$ to a $\tilde{c}$. Charged Higgs bosons with flavour non-diagonal couplings can also give important contributions in diagrams such as Fig. 7.22 c).

The main MSSM contributions are expected to come from the chargino–stop¹² and charged Higgs–top loops, as shown in Figs. 7.22 c) and d), respectively. However, there is little room

¹¹For the process $b \rightarrow d \gamma$ the Standard Model calculation yields $\text{Br}(B \rightarrow X_d \gamma) = 1.73_{-0.22}^{+0.12} \cdot 10^{-5}$.

¹²We usually expect higher generation off-diagonal soft terms to be larger due to RGE running controlled by Yukawa couplings.

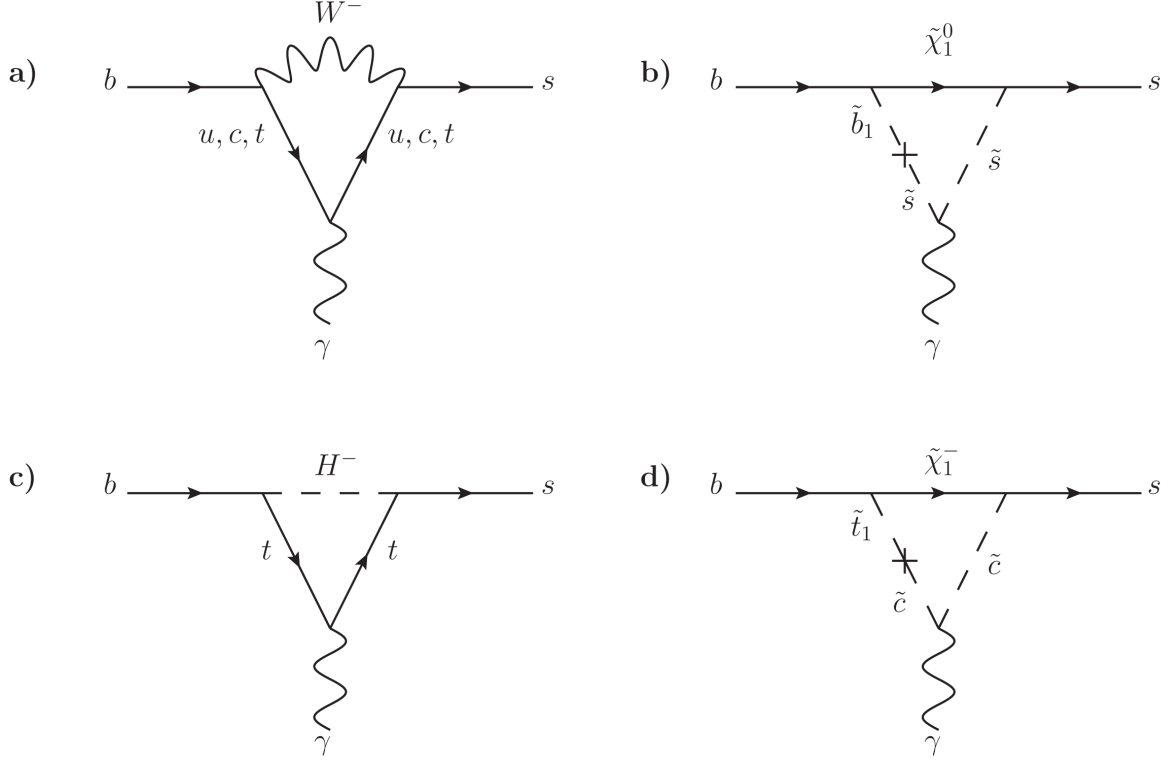


Figure 7.22: Diagrams for the process $b \rightarrow s \gamma$. a) shows the SM diagram while b), c) and d) show MSSM contributions.

for effects from supersymmetry since the current experimental world average is $\text{Br}(B \rightarrow X_s \gamma) = (3.32 \pm 0.15) \cdot 10^{-4}$ [32]. This means that either the charged Higgs is heavy enough and the stop-scharm soft mass term small enough not to contribute significantly, or that there are cancellations between the two contributions. This is both an example of a process that puts significant constraints on off-diagonal soft terms, and a process that is important in constraining the possible values of the charged Higgs mass.

7.5.4 $B_s \rightarrow \mu^+ \mu^-$

The process $B_s \rightarrow \mu^+ \mu^-$ is another FCNC process as either the bottom or the strange quark must change flavour in order to get a coupling to the muons. The Standard Model process is shown in Fig. 7.23 a), where the CKM factor in the W vertex allows one of the quarks to change flavour, but also suppresses the decay rate. For strange historical reasons involving darts and illegal substances this process is known as a **penguin diagram** [33].

On top of the small CKM factor, the process also suffers from what is called **helicity suppression** in the Standard Model. The Z -boson is spin-1, while the starting point meson B_s is spin-0 (pseudoscalar), meaning that the spins of the quarks in the meson are opposite. At some point in the diagram the helicity (chirality) must “flip”. This introduces an extra suppression proportional to $m_\mu^2/M_{B_s}^2$, making the expected rate extremely small and sensitive to supersymmetry contributions. We get a similarly suppressed process for B_d with a $\bar{d}$ -quark instead of the $\bar{s}$ in the initial state.

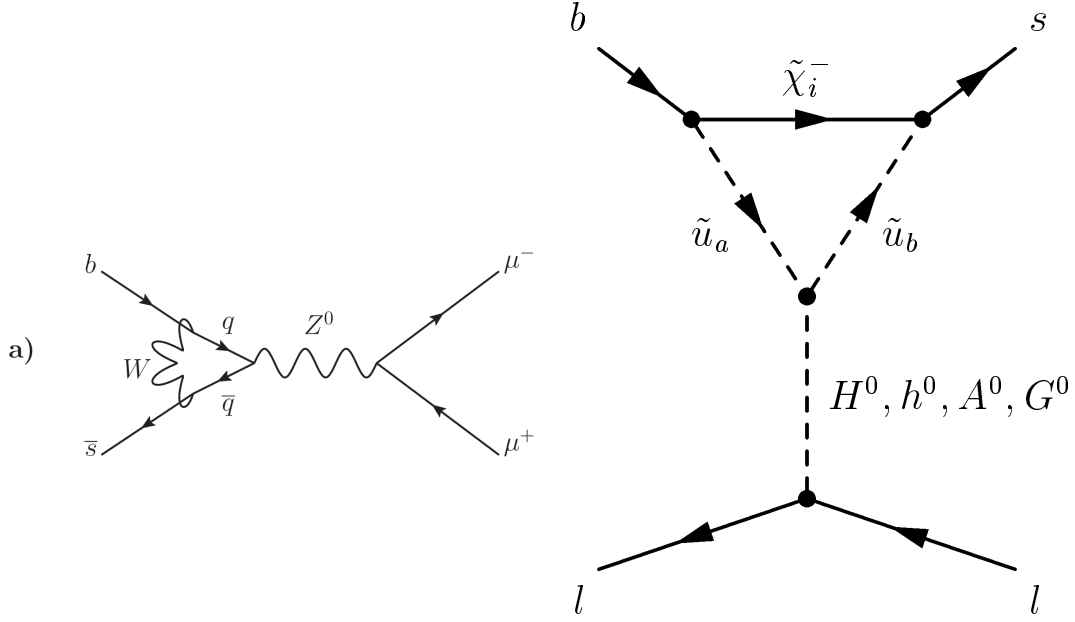


Figure 7.23: Diagrams for the process $B_s \rightarrow \mu^+ \mu^-$. Diagram a) shows one of the leading SM contributions, while b) shows one contribution from the MSSM taken from [34].

The predicted SM branching ratios for these processes are [35]:

$$\text{Br}(B_s \rightarrow \mu^+ \mu^-) = (3.65 \pm 0.23) \cdot 10^{-9}, \quad (7.14)$$

$$\text{Br}(B_d \rightarrow \mu^+ \mu^-) = (1.06 \pm 0.09) \cdot 10^{-10}. \quad (7.15)$$

First evidence for the $B_s \rightarrow \mu^+ \mu^-$ decay was shown by the LHCb Collaboration in 2012. The current value for the branching ratio from combined measurements at LHCb, ATLAS and CMS is [36]

$$\text{Br}(B_s \rightarrow \mu^+ \mu^-) = 2.69_{-0.35}^{+0.37} \cdot 10^{-9}, \quad (7.16)$$

$$\text{Br}(B_d \rightarrow \mu^+ \mu^-) < 1.6 \cdot 10^{-10}, \quad (7.17)$$

where the value for B_d is an upper limit at the 90% confidence level. It is interesting to observe that the measured value is lower than the Standard Model prediction at around the 2σ level, this has lead to speculations of destructive interference from new physics.

If we limit ourselves to scenarios with only diagonal soft terms the dominant contribution in the MSSM is from process such as shown in Fig. 7.23 b), with a chargino and up-type squark loop. These contributions are proportional to $\tan^6 \beta$, which makes the decay process highly sensitive to scenarios with large $\tan \beta$. To see this dependence, notice that μ couples to the mediating heavy higgses H/A^0 through the Yukawa term $y_{22}^l L_2 H_d \bar{E}_2$ in the superpotential, and the Yukawa constant in this term, $y_{22}^l = y_\mu$, is connected to the fermion mass through $m_\mu = y_\mu v \cos \beta$. Thus this vertex is proportional to $1/\cos \beta$ or $\tan \beta$, giving a factor $\tan^2 \beta$

in the amplitude squared.¹³

Furthermore, a chargino(higgsino)–stop loop can couple the strange and bottom quarks to the Higgs. These couplings are proportional to the bottom Yukawa coupling y_b , from the superpotential terms $y_{33}^d Q_3 H_d \bar{D}_3$, which appears in the stop–chargino–bottom vertex, and the $y_{32}^d Q_3 H_d \bar{D}_2$, which appears in the strange–chargino–stop vertex. Both these Yukawa couplings are proportional to y_b and thus to $1/\cos\beta$, giving a further factor of $\tan^4\beta$ in the amplitude squared. This $\tan\beta$ dependence makes $B_s \rightarrow \mu^+\mu^-$ an excellent channel for discovering supersymmetry, and puts very stringent bounds on the sparticle masses in large $\tan\beta$ scenarios.

7.6 Exercises

Exercise 7.1

From relativistic kinematics, show Eq. (7.13). *Hint:* the choice of rest frame is very important in order to simplify the calculation.

Exercise 7.2

Find the total cross section for the process $q\bar{q} \rightarrow \tilde{q}\tilde{q}^*$ via an s-channel gluon shown in Fig. 7.24.

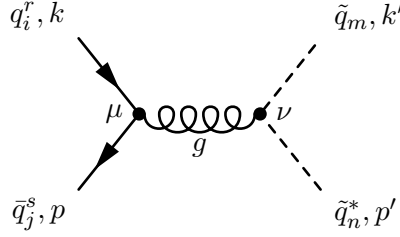


Figure 7.24: Strong production of a squark–anti-squark pair through a gluon.

¹³Remember that in the limit of large $\tan\beta$

$$\cos\beta = \pm \frac{1}{\sqrt{1 + \tan^2\beta}} = \pm \frac{1}{\tan\beta \sqrt{1 + \frac{1}{\tan^2\beta}}} \simeq \pm \frac{1}{\tan\beta}. \quad (7.18)$$

Chapter 8

Supersymmetric dark matter

We have a standard model also for cosmology, the Lambda Cold Dark Matter model (Λ CDM). This models the observed universe on three main ingredients: the known Standard Model matter (dominantly the baryonic matter of the atoms, radiation and neutrinos), dark energy (Λ) and cold dark matter (CDM), where the cold signifies that this ingredient is non-relativistic. In this chapter we will look closer at candidates for dark matter (DM) in supersymmetric models.

8.1 Evidence for dark matter

The idea of dark matter goes back quite a long way. Today we have evidence for the existence of dark matter through several effects where we observe its gravitational influence on ordinary matter. We list the evidence below:

- 1) Kinematics (Zwicky 1933 [38]): The motion of galaxies (velocity dispersion) cannot be explained by the visible matter. This has also been observed on the scales of galaxies in their rotation curves (Rubin 1970 [39]).
- 2) Gravitational lensing (Tyson 1996 [40]). First observed in galactic clusters. Clusters show evidence of lensing not explained by luminous matter. Dark matter dynamics (non-interacting) are demonstrated by the Bullet cluster (Clowe 2006 [41]).
- 3) Large scale structures (clusters, superclusters, filaments and voids): The structures observed in the 2dFGRS (2-degree Field Galaxy Redshift survey Colles 2001 [42]) and the SDSS (Sloan Digital Sky Survey Tegmark 2004 [43]) imply a relative matter density of $\Omega_m \equiv \frac{\rho_m}{\rho_c} \simeq 0.29$ where $\rho_c = 1.05 \cdot 10^{-5} h^2 \text{ GeV/cm}^3$ is the critical energy density for a flat universe.¹ They also imply that the majority of DM must be **cold** (non-relativistic), because warm DM would suppress clustering.
- 4) Big-Bang Nucleosynthesis (BBN): The formation of light elements in the period $t = 1 - 1000 \text{ s}$ after the Big Bang. Measurements of Early Universe abundance of light elements, mainly D and He, points to a baryonic matter density of $\Omega_b \approx 0.04$. This gives $\Omega_{\text{leftover}} = \Omega_{\text{DM}} \approx 0.25$.

¹ h is defined through the Hubble constant H_0 as $H_0 = 100h \text{ km/Mpc/s}$.

- 5) Supernovae (Riess 1998 [44] and Perlmutter 1999 [45]): Measurements of type Ia supernovae (SNe Ia) were used as standard candles to show an accelerated expansion of the Universe. This fixes $\Omega_\Lambda - \Omega_m \simeq k$ where Ω_Λ is the energy density of dark energy/cosmological constant.
- 6) Cosmic Microwave Background (CMB) (Penzias & Wilson 1965 [46]): The temperature variation of the CMB over the sky of the order of 0.0002 K is sensitive to all cosmological parameters, and gives $\Omega_\Lambda + \Omega_m \simeq k$, where k is some constant.

The evidence above can be used to constrain the Λ CDM concordance model of cosmology, which has just a handful of ingredients such as baryonic and dark matter, radiation (photons) and dark energy. In Fig. 8.1 we show the effects of the SNe, CMB and large scale structure data (BAO) on this model in the plane of matter density Ω_m and dark energy density Ω_Λ .

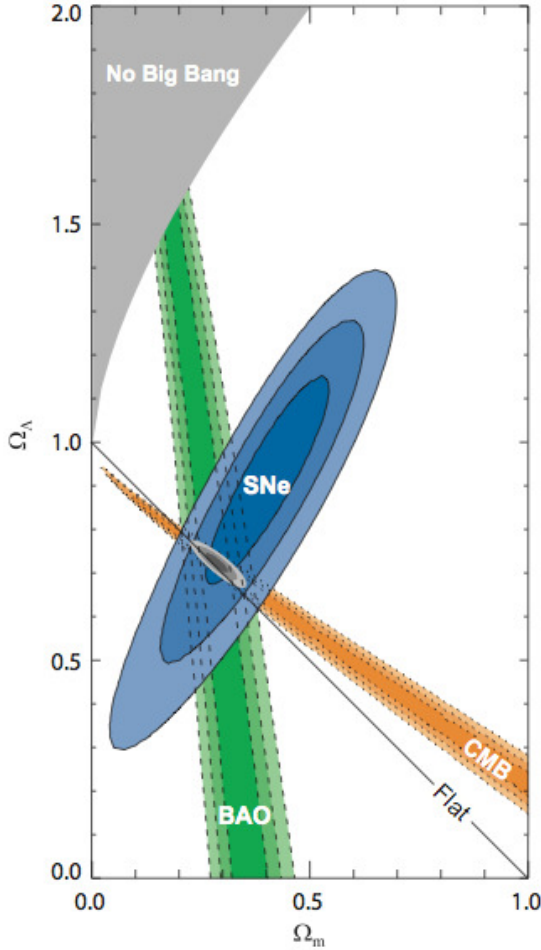


Figure 8.1: Limits from different observational sources on the dark energy density, Ω_Λ , compared to the total mass density in the universe, Ω_m .

A maximum likelihood fit to a selected subset of the measurements by the Planck Collaboration [47] gives the parameters for the model shown in Table 8.1. This means that the non-baryonic matter constitutes some 85% of the matter in the Universe, and that its relative energy density is quite well determined, at the percent level.

Parameter	Ω_Λ	$\Omega_m h^2$	$\Omega_b h^2$	H_0 [km/Mpc/s]	t_0 [Gy]
Value	$0.685^{+0.018}_{-0.016}$	0.1426 ± 0.0025	0.02205 ± 0.00028	67.3 ± 1.2	13.817 ± 0.048

Table 8.1: Measured values for cosmological parameters [47].

8.2 WIMP magic

The very existence of a stable Weakly² Interacting Massive Particle (WIMP) χ automatically gives an additional component to the total energy density of the Universe. WIMPs are found in a number of theories, for example the lightest neutralino of the MSSM, the lightest Kaluza-Klein particle of a theory with extra dimensions or an inert (no vev) Higgs boson.

This is due to the in equilibrium **thermal production** of the WIMP through the process $SM SM \rightarrow \chi\chi$, where SM are some Standard Model particles, and the reverse annihilation process $\chi\chi \rightarrow SM SM$, in the early hot Universe ($T \gg m_\chi$). As the temperature decreases to $T < m_\chi$ and there is not enough energy in an average collision for the production of χ to occur, only the reverse process can take place, and the comoving density³ falls with the temperature of the Universe.

The WIMPs then experience what is called a **chemical decoupling**, or loss of chemical equilibrium, due to the expansion of the Universe. This is when the WIMPs become so dilute, because of the expansion, that they in effect no longer interact inelastically, and this roughly happens when the expansion rate becomes larger than the rate of annihilation. The WIMPs then get a constant (comoving) density, we say that they experience a **freeze-out** at this temperature T_c . With weak-scale masses and couplings the freeze-out happens at $T_c \approx 0.05m_\chi$, however, this is before or at the same time as **kinetic decoupling** where the WIMPs effectively lose elastic interactions with the other matter, meaning that χ freezes-out with non-relativistic velocities and become cold dark matter.

The exact time (temperature) of freeze-out is controlled by the annihilation cross section of χ , larger cross sections keep chemical equilibrium for longer, in turn resulting in lower dark matter relic abundance. This abundance, in number density, can be found from the Boltzmann equation

$$\frac{dn_\chi}{dt} = -3Hn_\chi - \langle\sigma v\rangle(n_\chi^2 - n_\chi^{eq2}), \quad (8.1)$$

where n_χ^{eq} and n_χ are the chemical equilibrium and actual comoving number densities, H is Hubble's constant for the expansion rate, and $\langle\sigma v\rangle$ the velocity averaged annihilation cross section for $\chi\chi \rightarrow SM SM$. Solutions to this equation for increasing $\langle\sigma v\rangle$ are shown in Fig. 8.2.

In practice one must also often take into account co-annihilation with other particles with mass within 10%–20% of the χ , and numerical codes such as **DarkSUSY** [48] or **MicrOMEGAs** [49, 50] are typically used to solve the Boltzmann equation. For weak scale particles a rough approximation to the resulting dark matter density is

$$\Omega_\chi h^2 = 0.1 \times \frac{3 \times 10^{-26} \text{ cm}^3 \text{ s}^{-1}}{\langle\sigma v\rangle},$$

²Weak as in electro-weak, meaning on the same scale as the weak force.

³Taking the expansion of the universe into account by looking at the number of particles in a volume expanding at the same rate.

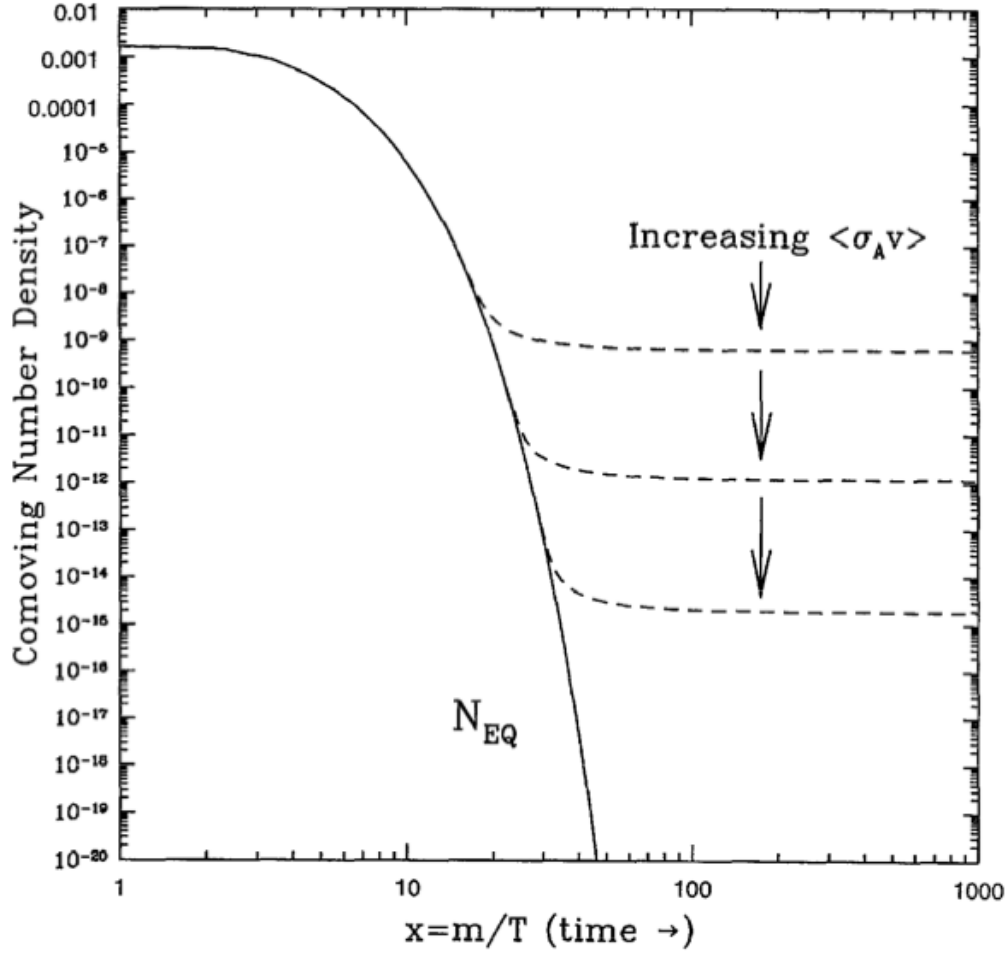


Figure 8.2: Illustration of the freeze-out of the comoving number density of a WIMP as a function of time, where the black line represents a model without chemical decoupling, and the dotted lines represent different freeze out temperatures for different velocity averaged annihilation cross sections.

and since the annihilation cross section can be shown to be

$$\langle\sigma v\rangle \approx \frac{\alpha_{weak}^2}{m_{weak}^2} \approx 10^{-25} \text{ cm}^3 \text{ s}^{-1}, \quad (8.2)$$

the predicted dark matter density is

$$\Omega_\chi h^2 \approx 0.1 \times \left(\frac{g_{weak}}{g_\chi} \right)^4 \left(\frac{m_\chi}{m_{weak}} \right)^2.$$

When compared to the values in Table 8.1 this is called the **WIMP-miracle** since it supplies just about the correct missing energy density for the WIMP as long as it has a weak scale mass and a weak scale coupling to the Standard Model. The strong sensitivity of the predicted dark

matter density to both the mass and the coupling, and the precisely measured experimental value, means that the WIMP scenario is very predictive and testable.

For a more detailed discussion of the WIMP miracle, see the standard cosmology book by Kolb and Turner [51].

8.3 Dark matter candidates in supersymmetry

With R-parity conservation in place we have seen that the LSP is stable and any neutral LSP can then in principle constitute all or part of the dark matter. Without R-parity only super-weakly coupling particles like gravitinos and axinos are candidates. Below we briefly discuss the various possibilities.

8.3.1 Neutralino

As soon as you have a stable neutralino LSP, you usually get into trouble trying to explain why there is so little dark matter. For example, the standard mSUGRA bino-like $\tilde{\chi}_1^0$ scenario prefers LSP masses around 100 GeV or lower not to overproduce dark matter. Due to current lower bounds on the $\tilde{\chi}_1^0$ mass and the measured Higgs mass this tends to be difficult to realise in practise. This scenario with a low mass bino-like neutralino is often called the **bulk region** scenario, which can be seen in Fig. 8.3 in the lower left corner.

Alternatives to the bulk region scenario use co-annihilation or resonant annihilation to increase $\langle\sigma v\rangle$ and thus decrease the dark matter density. The **stau-coannihilation region**, where $\tilde{\tau}_1 \tilde{\chi}_1^0 \rightarrow SM SM$ is an efficient annihilation process reducing the dark matter density, exists for $m_{\tilde{\tau}_1} - m_{\tilde{\chi}_1^0} \leq 10$ GeV. This makes the scenario difficult to discover at collider experiments due to the production of soft (low-energy) taus. In mSUGRA the stau-coannihilation region can be found for small m_0 , and is shown as the lower strip in Fig. 8.3 which follows the lower theoretical bound (brown) where the stau becomes the LSP. Outside of mSUGRA similar regions can also be found for smuons and selectrons, with large mass degeneracy between the slepton and the LSP.

The **stop-coannihilation region**, where $\tilde{t}_1 \tilde{\chi}_1^0 \rightarrow SM SM$ is efficient typically has $m_{\tilde{t}_1} - m_{\tilde{\chi}_1^0} \leq 25$ GeV. In mSUGRA this exists for large values of $|A_0|$, small m_0 and $m_{1/2}$. Again, this is difficult to discover because of the soft decay products of the stop.

The **Higgs funnel region** is found for $2m_{\tilde{\chi}_1^0} \simeq m_{A,H}$ and large $\tan\beta$, where the neutralino has resonant annihilation through a heavy Higgs boson. For mSUGRA, this is shown in Figure 8.3 as the two diagonal structures roughly in the middle of the plot, rising as a funnel upwards.

The **chargino-coannihilation** region can be found when we have higgsino or wino LSPs, where a chargino is automatically degenerate with the LSP resulting in a lower dark matter density typically below experimental bounds. In mSUGRA this region is called the **focus point region**, and is found for large m_0 and low μ , leading to so-called split-SUSY, as the sfermion masses need to be pressed up quite a bit, and there is a large mass difference between gauginos and sfermions. The focus point region can be seen in Fig. 8.3 following the upper theoretical bound where EWSB breaks down.

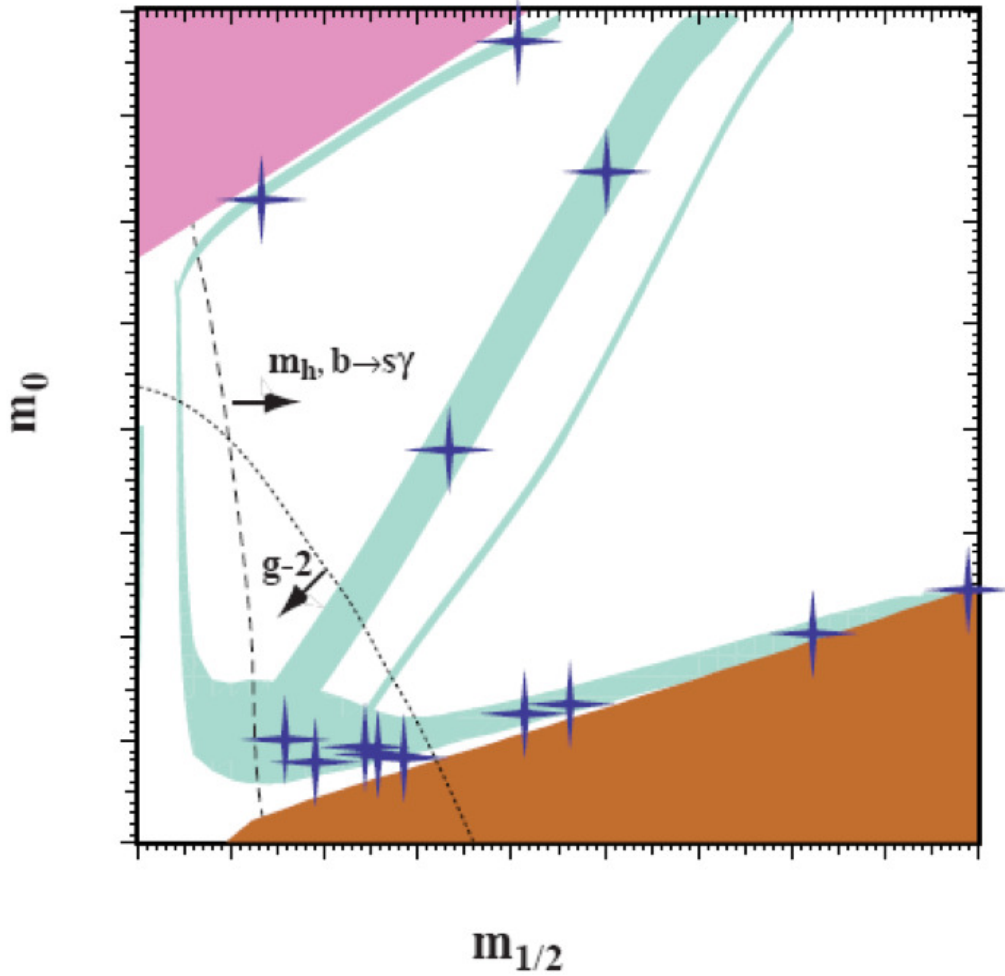


Figure 8.3: Generic illustration of the allowed neutralino dark matter regions (puke green) in the $(m_0, m_{1/2})$ -plane for mSUGRA. Except for the low m_0 and $m_{1/2}$ regions the area outside of the allowed green region gives too much dark matter. The dashed line shows the Higgs mass limit which pushes towards larger values of $m_{1/2}$, while the dotted line represents the limit from the anomalous magnetic moment of the muon.

8.3.2 Sneutrinos

The left handed sneutrino $\tilde{\nu}_L$ is happily excluded as a potential dark matter candidate due to the large cross section for $\tilde{\nu}_L q \rightarrow \tilde{\nu}_L q$ via Z -exchange.⁴ The large cross means that it should already have been seen by direct detection experiments, see Sec 8.4. It is also problematic to get $m_{\tilde{\nu}_L} < m_{\tilde{l}_L}$ due to hyperfine-splitting. However, the right-handed sneutrino $\tilde{\nu}_R$ couples very weakly and is still a viable candidate. This would necessitate adding right-handed

⁴For the neutralinos this problem only exists for a higgsino $\tilde{\chi}_1^0$ LSP, as the wino and bino do not couple to the Z .

neutrinos and sneutrino superfields to the MSSM.

8.3.3 Gravitino

The **gravitino** is the massive spin-3/2 supersymmetric partner of the graviton in models with supergravity. With only gravitational strength interactions, the gravitino is *not* a WIMP as it is never in chemical equilibrium. If it is the LSP it can be created from the decays of the next-to-lightest supersymmetric particle (NLSP), giving, in R-parity conserving scenarios,

$$\Omega_{\tilde{G}} = \frac{m_{\tilde{G}}}{m_{\text{NLSP}}} \Omega_{\text{NLSP}}.$$

However, these scenarios are problematic because the NLSP is long-lived due to the extremely weak interaction and creates potential trouble in BBN by injecting energy that changes the production of light elements.

Alternatively, the gravitino can be created in non-thermal production processes in the period of reheating after inflation. One potential process $gg \rightarrow \tilde{g}\tilde{G}$ is shown in Fig. 8.4. The reverse process $\tilde{g}\tilde{G} \rightarrow gg$ is not efficient as the density of gravitinos is never high enough given the small cross section. This type of dark matter creation process is often called **freeze-in**. For the gravitino this gives a new magic formula:

$$\Omega_{\tilde{G}} h^2 \approx 0.5 \cdot \left(\frac{T_R}{10^{10} \text{ GeV}} \right) \left(\frac{100 \text{ GeV}}{m_{\tilde{G}}} \right) \left(\frac{m_{\tilde{g}}}{1 \text{ TeV}} \right)^2, \quad (8.3)$$

where T_R is the unknown and very weakly constrained reheating temperature. This production mechanism is also valid for models with R-parity violation. There the small gravitino coupling $\propto \frac{1}{M_P}$ makes the gravitinos very long-lived, with lifetimes longer than the age of the Universe, but not absolutely stable because of the R-parity violating operators.

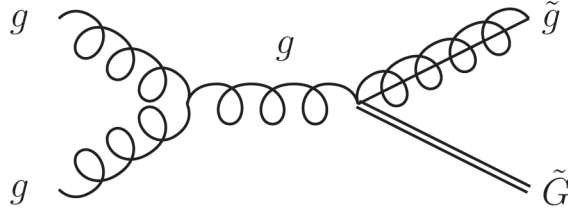


Figure 8.4: One possible diagram for the non-thermal production of gravitinos from the scattering of gluons.

8.3.4 Others

One could even imagine colour charged supersymmetric particles as dark matter, in particular the gluino, which, if stable, after hadronisation with Standard Model quarks and/or gluons would form so-called **R-hadrons**. These have very strict limits from searches, but these limits are somewhat obfuscated by complications in modelling R-hadron scattering.

8.4 Direct detection

In addition to the direct production of dark matter at colliders and the corresponding searches for missing energy, there are two other main ways to search for dark matter, direct and indirect detection. Here we briefly discuss **direct detection**.

Direct detection seeks to make weak and thus rare dark matter interactions with standard model matter visible by very low background searches in large volumes, using local galactic halo dark matter interacting with ordinary matter on Earth. The feasibility of direct detection is very dependent on the χ scattering cross section on nucleons (quarks), which can in principle be calculated in a given model,⁵ but also on the dark matter halo density distribution and velocity distribution, which have large uncertainties.

The rate of interaction between dark matter and a particular type of nucleons in a sample can be expressed in the differential scattering rate with respect to the recoil energy E_r of a scattered nucleon with mass M as

$$\frac{dN}{dE_r} = \frac{\sigma \rho_{DM}}{2\mu^2 m_\chi} |F(q)|^2 \int_{v_{min}}^{v_{esc}} \frac{f(\vec{v})}{v} d^3v, \quad (8.4)$$

where σ is the dark matter scattering cross section off the nucleus in question, ρ_{DM} is the dark matter halo density at Earth, $\mu = m_\chi M/(m_\chi + M)$ is the dark matter and nucleus reduced mass, $F(q)$ is a nuclear form factor dependent on the scattering momentum transfer $q = \sqrt{2ME_r}$ and the nucleon type, $f(\vec{v})$ is the velocity distribution of the dark matter in the halo, $v_{min} = \sqrt{ME_r/2\mu^2}$ is the minimal velocity that gives a recoil energy E_r and v_{esc} is the escape velocity from the halo.

In general direct detection relies heavily on suppressing Standard Model backgrounds which would be the overwhelmingly dominant cause of observed recoils. This is done by moving the experiments deep under ground where they are shielded from cosmic rays, and by specifically shielding the detector samples against local radioactivity. Beyond this, there are two main tactics followed in order to try to directly detect dark matter:

- Suppress (almost) all backgrounds and monitor as large a volume as possible. This used in experiments such as XENON, PANDAS and CDMS.
- Look for an annular modulation in the event rate – due to the Earth’s movement in the galactic rest frame and thus in the dark matter halo, changing the velocity integral in Eq. (8.4) – to observe a small dark matter signal on top of a constant background, such as used in the DAMA and CoGeNT experiments.

Figure 8.5 shows results from the current most important direct detection experiments as limits on the dark matter–nucleon cross section as a function of dark matter mass. Observe that while DAMA claims a detection (closed regions), these are already excluded by many other experiments. For reference, a neutralino WIMP could reasonably be expected to have a scattering cross section in the range $10^{-45} - 10^{-50} \text{ cm}^2$ for masses in the range 100–1000 GeV [52]. The dashed yellow line shows what is called the **neutrino floor**. This is the cross section that can be reached before the experiments would be saturated with recoils due to neutrinos scattering off the experiment sample. Going below the neutrino floor would need

⁵However, while the quark–dark matter cross section is relatively easy to calculate, the fact that the quark is bound in a nucleon, and the nucleon bound to a nucleus makes this more difficult in practice, involving also estimates from nuclear physics.

new experimental techniques, fundamentally, the direction of the particle creating the recoil would need to be reconstructed.

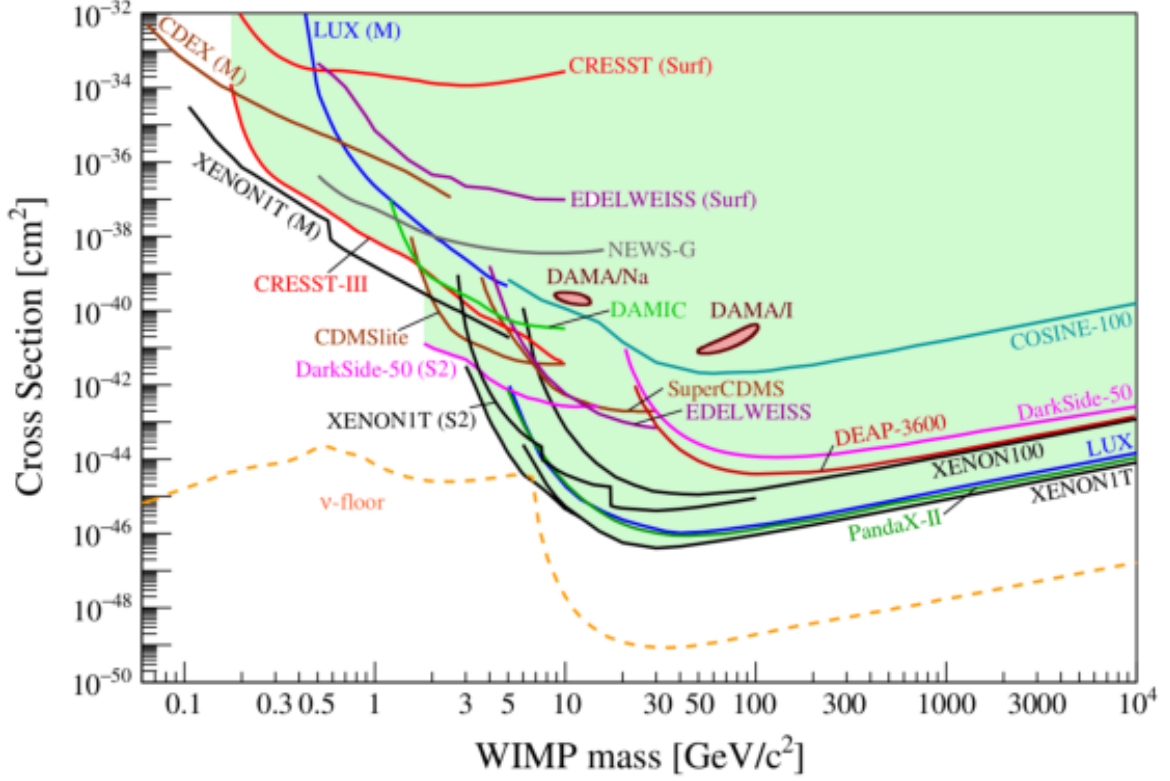


Figure 8.5: Plot of different exclusion and detection results for direct detection of dark matter in the WIMP mass versus WIMP–nucleon cross section plane [53]. The dashed yellow line shows the neutrino floor.

8.5 Indirect detection

In **indirect detection** we look for the annihilation, $\chi\chi \rightarrow SM SM$, or decay products of dark matter in multiple final (messenger) states in cosmic rays. The viable search channels must be stable Standard Model particles, so that they can reach the Earth (or satellites in orbit). The messengers should also have as low backgrounds from ordinary astrophysical processes as possible, this makes searches with electrons and protons extra challenging. The remaining candidates are photons, neutrinos, positrons, antiprotons and antideuterons.⁶ We now discuss the properties of each of these in a little more detail.

Photons

These can either come from direct production processes such as the annihilation channels $\chi\chi \rightarrow \gamma\gamma, Z\gamma$, which is relatively easy to detect because the spectrum is a sharp line at

⁶Potentially, even heavier antinuclei than antideuterons could be used, but they would be even rarer.

exactly the mass of the dark matter, or they could be γ from bremsstrahlung or pion decays in the annihilation decay products, which creates a broad spectrum. This is harder to detect because of the large background of bog-standard astrophysical photons, but is expected to make up the majority of photons from dark matter. Photons from dark matter have the advantage that they point to the source, so we can focus our searches on areas in our galaxy or nearby galaxies with large ρ_{DM} , where annihilation or decay is more probable, and thus reducing potential backgrounds relative to the signal. We can also look for photons that are extragalactic in origin, but then we have to account for the red-shifting of the spectrum.

Dark matter annihilating in our own galaxy into photons should result in a differential flux at Earth in terms of energy E and solid angle Ω given by

$$\frac{d\Phi}{dEd\Omega} = \frac{1}{8\pi m_\chi^2} \frac{dN_\gamma}{dE} \langle\sigma v\rangle \int_{\text{l.o.s.}} \rho_{DM}^2(l) dl, \quad (8.5)$$

where $N_\gamma(E)$ gives the number of photons with energy E in a single annihilation event and the integral is over the dark matter density in the line of sight. We see that the flux depends on the square of the dark matter density since annihilation requires two particles to be present. For decaying dark matter the corresponding expression is proportional to ρ_{DM} .

There have been some indications of an excess of photons above expected backgrounds from the galactic centre (a.k.a. the Hooperon [54]), however, no unambiguous dark matter signal has been confirmed. Current limits from the Fermi-LAT experiment seems to rule out most possible models for a dark matter explanation for this excess, and, more importantly, sets a cross section limit for dark matter annihilation close, and for some masses beyond, to the canonical limit of

$$\langle\sigma v\rangle = 3 \times 10^{-26} \text{ cm}^3\text{s}^{-1},$$

which would give the correct dark matter density, see Fig. 8.6.⁷ This limit is some of the strongest evidence today against WIMP dark matter.

Neutrinos

These also point to their source since since they are electrically neutral, and can be extragalactic in origin just like the photons, still reaching the Earth. The same flux calculation can be used, starting from the neutrino spectrum from dark matter annihilation. While the astrophysical background is much smaller, the neutrino signal is difficult to detect due to the weak matter interaction. The current leading experiment in detecting neutrinos from dark matter annihilation or decay is the IceCube experiment at the South Pole that has instrumented a cubic kilometre of the South Pole ice with photodetectors a kilometre below the surface.⁸

One interesting alternative possibility is that dark matter scatters on ordinary matter sufficiently strongly that the dark matter accumulates at the centre of the Sun (or possibly the Earth), because it scatters off the Sun's atoms, loses energy and becomes gravitationally bound, ultimately being further downscattered in energy in multiple interactions, settling at the Sun's centre. When these DM particles annihilate the only decay products that can escape the Sun's interior are neutrinos. These could then potentially be detected, and the fact that

⁷Keep in mind though, that the velocity distribution of any dark matter today would not be the same as during freeze-out at a much higher temperature, so the cross section numbers are not exactly translatable.

⁸This should, on galactic scales, get some sort of prize for the geniusly nutscase idea.

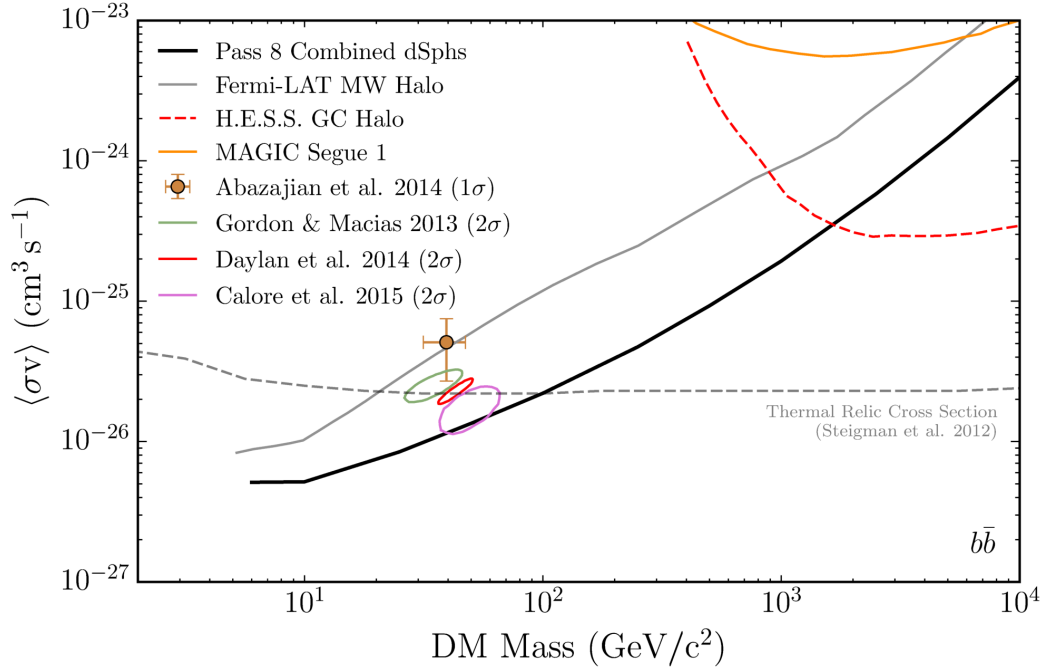


Figure 8.6: Results from Fermi-LAT indirect gamma-ray searches in the $\chi\chi \rightarrow b\bar{b}$ channel. Grey line shows limit from Milky Way halo search, black line from Milky Way dwarf spheroidal galaxy search with six years of data [55].

high energy neutrinos would be coming directly from the Sun would reduce astrophysical backgrounds significantly.

Positrons

Charged particles propagate in a complicated way through the galactic magnetic field, and they are therefore impossible to track back to the source. At the same time sources outside of our own Galaxy cannot contribute significantly to the flux at Earth. Positrons also have relatively large astrophysical backgrounds, for example pulsars are expected to produce a significant amount of lower energy positrons, so experiments search for small excesses on large backgrounds, mostly at high energies. Some potential excess has been seen in the past by Fermi-LAT and PAMELA, but this has not been conclusive.

Antiprotons

Just as the positrons these propagate in a complicated manner, but the backgrounds are under better control due to the lack of significant astrophysical antiproton production, except from the scattering of very high energy cosmic rays on the interstellar medium. The current best limits here come from the AMS-02 experiment.

Antideuterons

These have very low backgrounds because the production of antideuterons in astrophysical processes is extremely rare, in particular at low energies, due to the conservation of baryon number and the required high energy of cosmic rays to produce two antibaryons (at least two baryons must be produced in the same interaction). This means that the detection of an antideuteron should be a smoking duck signal of dark matter annihilation or decay. However, the physics of the formation of antideuterons from dark matter is quite complicated and hard to reliably calculate, bringing in extra uncertainty. Current searches set upper bounds on the flux of antideuterons at the Earth. It is hoped that AMS-02 will provide new better data on antideuterons soon.⁹

8.6 Exercises

Exercise 8.1

Show that $\chi\chi \rightarrow Z \rightarrow f\bar{f}$ gives

$$\sigma v \approx \frac{g^4 E_\chi^2}{128\pi m_Z^2}, \quad (8.6)$$

which in the low-velocity limit can be shown to be

$$\langle\sigma v\rangle_0 \approx 10^{-25} \text{ cm}^3\text{s}^{-1}.$$

⁹This has been hoped for a very long time.

Appendix A

Solutions to exercises

Exercise 1.1

Given the group multiplication definition the identity element must be $e = T_0 = 1$ since $T_a \circ T_0 = T_{a+0} = T_a$. Let $T_a^{-1} = T_b$. Since $T_a \circ T_b = T_{a+b} = 1$, we find $b = -a$. (This does not show that the inverse is unique, but this is not a requirement.) We have now demonstrated the required existence of an identity element (iii) and an inverse (iv) for $T(1)$ as required by the group definition (the order of operations can obviously be reversed). The closure of the multiplication operation (i) is true by its definition. The associativity (ii) can be demonstrated as follows $(T_a \circ T_b) \circ T_c = T_{a+b} \circ T_c = T_{a+b+c} = T_a \circ T_{b+c} = T_a \circ (T_b \circ T_c)$.

Exercise 1.3

An $n \times n$ matrix U over $\mathbb{C}$ has n^2 complex or $2n^2$ real parameters. The hermitian matrix $M = U^\dagger U$ has n real elements on the diagonal and $(n^2 - n)/2$ complex elements above and below the diagonal. Because $M^\dagger = M$ all the $(n^2 - n)/2$ elements below the diagonal are given by the complex conjugate the corresponding elements above the diagonal. Since U is unitary $M = I$. That all the n real diagonal elements of M is equal to one gives n restrictions (n real equations) for the elements of U . Further, all $(n^2 - n)/2$ complex elements above the diagonal are zero, which gives $(n^2 - n)/2$ complex equations, which means $n^2 - n$ real equations. For the terms below the diagonal we do not obtain any new equations. As a result the free parameters for the $U(n)$ matrices are $2n^2 - n - (n^2 - n) = n^2$. For $SU(n)$ we must use the additional requirement $\det U = 1$. Because $1 = \det M = \det(U^\dagger U) = \det U^\dagger \det U = |\det U|^2$, we see that $\det U$ must be a phase factor, and that the requirement $\det U = 1$ only gives one new condition (equation). Thus there are $n^2 - 1$ independent real parameters in $SU(n)$.

Exercise 1.5

Since H is a subgroup of G then by point ii) in the definition since $h_j \in H$ we must also have $h_j^{-1} \in H$. With $h_i, h_j^{-1} \in H$ by point i) in definition $h_i \circ h_j^{-1} \in H$.

Exercise 1.6

Let $U_i, U_j \in SU(n)$, then

$$\det(U_i U_j^{-1}) = \det(U_i) \det(U_j^{-1}) = 1.$$

This means that $U_i U_j^{-1} \in SU(n)$. In other words, $SU(n)$ is a proper subgroup of $U(n)$. Let $V \in U(n)$ and $U \in SU(n)$, then $VUV^{-1} \in SU(n)$ because:

$$\det(VUV^{-1}) = \det(V) \det(U) \det(V^{-1}) = \frac{\det(V)}{\det(V)} \det(U) = 1.$$

In other words, $SU(n)$ is also a normal subgroup of $U(n)$, thus $U(n)$ is not simple.

Exercise 1.13

A member U of $SU(2)$ fulfils $UU^\dagger = I$ and has $\det U = 1$ and can be written as

$$U = \begin{bmatrix} a + ib & c + id \\ -c + id & a - ib \end{bmatrix},$$

where $a, b, c, d \in \mathbb{R}$ and $a^2 + b^2 + c^2 + d^2 = 1$. Using a composition function here it is tempting to solve for $a = \sqrt{1 - b^2 - c^2 - d^2}$ and let a, b, c be the three free parameters of $SU(2)$. Note that solving for b, c or d is not really an option as a requirement for deriving the generators from the composition function is that the zero parameters, here $b = c = d = 0$, give the identity element of the group, which is the identity matrix. However, there is a potential issue here that choosing the sign on the square root restricts the parameterisation to apply for only group members with positive real components in the upper left element, meaning that this composition function only works for part of the group. This is another example of a local description of the group. Other parameterisations exist, however, they may have other problems such as a non-unique zero element (meaning no inverse of the composition function exists).

Fortunately, this parameterisation is enough to give all of the generators of the whole group:

$$\begin{aligned} X_1 &= \left. \frac{\partial U}{\partial d} \right|_{b=c=d=0} = \begin{bmatrix} 0 & i \\ i & 0 \end{bmatrix}, \\ X_2 &= \left. \frac{\partial U}{\partial c} \right|_{b=c=d=0} = \begin{bmatrix} 0 & 1 \\ -1 & 0 \end{bmatrix}, \\ X_3 &= \left. \frac{\partial U}{\partial b} \right|_{b=c=d=0} = \begin{bmatrix} i & 0 \\ 0 & -i \end{bmatrix}. \end{aligned}$$

While these are not the Pauli matrices we might be expecting, they are related to the Pauli matrices by a simple multiplicative factor $\sigma_1 = -iX_1$, $\sigma_2 = -iX_2$, and $\sigma_3 = -iX_3$. In fact the matrices we have found are typically used in mathematical texts as the generators of $SU(2)$, but then the exponential map has no factor of i . Notice that these matrices are anti-hermitian, while the Pauli matrices are hermitian.

The infinitesimal approach writes the group member as $U = e^{ida_i X_i} \simeq I + ida_i X_i$, and applying the unitarity requirement

$$UU^\dagger \simeq (I + ida_i X_i)(I + ida_i X_i)^\dagger = I + ida_i(X_i - X_i^\dagger) + \dots,$$

means that a necessary condition on the generators X_i is that they are hermitian $X_i = X_i^\dagger$. The general form of the generators is thus restricted to

$$X = \begin{bmatrix} \alpha & \gamma \\ \gamma^* & \beta \end{bmatrix},$$

where $\alpha, \beta \in \mathbb{R}$ and $\gamma \in \mathbb{C}$. The property of the determinant can be found as follows (see Exercise ??):

$$\det U = \det e^{ida_i X_i} = e^{\text{Tr}[ida_i X_i]} = e^{ida_i \text{Tr}[X_i]}.$$

This means that for the determinant to be one, each trace must be zero, meaning that the generators are traceless so that the general form of the generators is restricted to

$$X = \begin{bmatrix} \alpha & \gamma \\ \gamma^* & -\alpha \end{bmatrix}.$$

From the requirement $\alpha^2 + |\gamma|^2 = 1$ we can recreate the three Pauli matrices by the three choices: 1) $\alpha = 1, \gamma = 0$, 2) $\alpha = 0, \gamma = 1$, and 3) $\alpha = 0, \gamma = -i$. Naturally, there is a continuum of equivalent expressions for the generators from other compatible choices.

This approach emphasises two important properties of the generators. Firstly that the generators are hermitian. Since the generators will function as operators in a QM or QFT setting, this is very desirable. Second, the generators are traceless. This will often be a calculational advantage.

Exercise 2.3

The first relation follows trivially from the commutation of P_μ with P_ν . To show the second we first use that

$$[M_{\mu\nu}, P_\rho P^\rho] = [M_{\mu\nu}, P_\rho] P^\rho + P_\rho [M_{\mu\nu}, P^\rho], \quad (2.47)$$

and Eq. (2.7) to get:

$$[M_{\mu\nu}, P_\rho P^\rho] = -i(g_{\mu\rho} P_\nu - g_{\nu\rho} P_\mu) P^\rho - iP_\rho (g_\mu{}^\rho P_\nu - g_\nu{}^\rho P_\mu), \quad (2.48)$$

thus

$$[M_{\mu\nu}, P_\rho P^\rho] = -2i[P_\mu, P_\nu] = 0. \quad (2.49)$$

Exercise 2.4

From the definition of W_ν we have

$$\begin{aligned} [P_\mu, W_\nu] &= \frac{1}{2} \epsilon_{\nu\rho\sigma\tau} [P_\mu, P^\rho M^{\sigma\tau}] \\ &= \frac{1}{2} \epsilon_{\nu\rho\sigma\tau} g_{\mu\gamma} ([P^\gamma, P^\rho] M^{\sigma\tau} + P^\rho [P^\gamma, M^{\sigma\tau}]) \\ &= \frac{i}{2} \epsilon_{\nu\rho\sigma\tau} g_{\mu\gamma} P^\rho (g^{\sigma\gamma} P^\tau - g^{\tau\gamma} P^\sigma) \\ &= \frac{i}{2} \epsilon_{\nu\rho\sigma\tau} g_{\mu\gamma} P^\rho g^{\sigma\gamma} P^\tau - \frac{i}{2} \epsilon_{\nu\rho\sigma\tau} g_{\mu\gamma} P^\rho g^{\tau\gamma} P^\sigma \\ &= \frac{i}{2} \epsilon_{\nu\rho\mu\tau} P^\rho P^\tau - \frac{i}{2} \epsilon_{\nu\rho\sigma\mu} P^\rho P^\sigma \\ &= 0, \end{aligned}$$

using the Poincaré algebra properties from (2.7).

Exercise 2.7

$$[X_\mu, P_\nu] = \frac{1}{2}[\bar{Q}\gamma_\mu\gamma^5 Q, P_\nu] = \frac{1}{2}(\bar{Q}\gamma_\mu\gamma^5)_a[Q_a, P_\nu] + \frac{1}{2}[\bar{Q}_a, P_\nu](\gamma_\mu\gamma^5 Q)_a = 0$$

Exercise 3.1

$$\begin{aligned} \int d^2\theta \theta\theta &= -\frac{1}{4} \int d\theta^A d\theta_A \theta^A \theta_A \\ &= \frac{1}{2} \int d\theta^A d\theta_A \theta_1 \theta_2 = - \int d\theta_1 d\theta_2 \theta_1 \theta_2 = \int d\theta_1 \theta_1 \int d\theta_2 \theta_2 = 1 \end{aligned}$$

Exercise 3.3

We start from

$$\Phi(y, \theta) = A(y) + \sqrt{2}\theta\psi(y) + \theta\theta F(y),$$

and $y^\mu \equiv x^\mu + i\theta\sigma^\mu\bar{\theta}$. This gives

$$\begin{aligned} \Phi(x, \theta, \bar{\theta}) &= A(x + i\theta\sigma\bar{\theta}) + \sqrt{2}\theta\psi(x + i\theta\sigma\bar{\theta}) + \theta\theta F(x + i\theta\sigma\bar{\theta}) \\ &= A(x) + \partial_\mu A(x)i\theta\sigma^\mu\bar{\theta} + \frac{1}{2}\partial_\mu\partial_\nu A(x)(i\theta\sigma^\mu\bar{\theta})(i\theta\sigma^\nu\bar{\theta}) + \sqrt{2}\theta\psi(x) \\ &\quad + \sqrt{2}\partial_\mu\theta\psi(x)i\theta\sigma^\mu\bar{\theta} + \theta\theta F(x) \\ &= A(x) + i\theta\sigma^\mu\bar{\theta}\partial_\mu A(x) - \frac{1}{2}\partial_\mu\partial_\nu A(x)\frac{1}{2}g^{\mu\nu}\theta\theta\bar{\theta}\bar{\theta} + \sqrt{2}\theta\psi(x) \\ &\quad + i\sqrt{2}\theta\partial_\mu\psi(x)\theta\sigma^\mu\bar{\theta} + \theta\theta F(x) \\ &= A(x) + i\theta\sigma^\mu\bar{\theta}\partial_\mu A(x) - \frac{1}{4}\square A(x)\theta\theta\bar{\theta}\bar{\theta} + \sqrt{2}\theta\psi(x) + \frac{i}{\sqrt{2}}\theta\theta\partial_\mu\psi\sigma^\mu\bar{\theta} + \theta\theta F(x), \end{aligned}$$

where we have used that $(\theta\sigma^\mu\bar{\theta})(\theta\sigma^\nu\bar{\theta}) = \frac{1}{2}g^{\mu\nu}\theta\theta\bar{\theta}\bar{\theta}$ and $\theta\partial_\mu\psi\theta\sigma^\mu\bar{\theta} = \frac{1}{2}\theta\theta\partial_\mu\psi\sigma^\mu\bar{\theta}$.

Exercise 4.1

The Lagrangian for a single scalar superfield Φ is

$$\mathcal{L} = \Phi^\dagger\Phi + \delta^2(\bar{\theta})(g\Phi + m\Phi\Phi + \lambda\Phi\Phi\Phi) + \delta^2(\theta)(g\Phi^\dagger + m\Phi^\dagger\Phi^\dagger + \lambda\Phi^\dagger\Phi^\dagger\Phi^\dagger). \quad (4.27)$$

Using the following expressions for the scalar superfield in terms on component fields

$$\begin{aligned} \Phi(x, \theta, \bar{\theta}) &= A(x) + i(\theta\sigma^\mu\bar{\theta})\partial_\mu A(x) - \frac{1}{4}\theta\theta\bar{\theta}\bar{\theta}\square A(x) + \sqrt{2}\theta\psi(x) - \frac{i}{\sqrt{2}}\theta\theta\partial_\mu\psi(x)\sigma^\mu\bar{\theta} \\ &\quad + \theta\theta F(x). \\ \Phi^\dagger(x, \theta, \bar{\theta}) &= A^*(x) - i(\theta\sigma^\mu\bar{\theta})\partial_\mu A^*(x) - \frac{1}{4}\theta\theta\bar{\theta}\bar{\theta}\square A^*(x) + \sqrt{2}\bar{\theta}\bar{\psi}(x) + \frac{i}{\sqrt{2}}\bar{\theta}\bar{\theta}\theta\sigma^\mu\partial_\mu\bar{\psi}(x) \\ &\quad + \bar{\theta}\bar{\theta}F^*(x). \end{aligned}$$

we have for the kinetic term,

$$\begin{aligned} \int d^4\theta \Phi^\dagger \Phi &= -\frac{1}{4}A\Box A^* - \frac{1}{4}A^*\Box A + |F|^2 \\ &\quad + \int d^4\theta \{(\theta\sigma^\mu\bar{\theta})(\theta\sigma^\nu\bar{\theta})\partial_\mu A^*(\theta\sigma^\nu\bar{\theta})\partial_\nu A + i\bar{\theta}\bar{\theta}\theta\sigma^\mu\partial_\mu\bar{\psi}\theta\psi - i\bar{\theta}r\psi\theta\theta\partial_\mu\psi\sigma^\mu\bar{\theta}\}, \end{aligned}$$

where we have removed the factor $\theta\theta\bar{\theta}\bar{\theta}$ by integration. Using the identities

$$(\theta\sigma^\mu\bar{\theta})(\theta\sigma^\nu\bar{\theta}) = \frac{1}{2}g^{\mu\nu}\theta\theta\bar{\theta}\bar{\theta}, \quad (4.28)$$

$$\theta\sigma^\mu\partial_\mu\bar{\psi}\theta\psi = -\frac{1}{2}\psi\sigma^\mu\partial_\mu\bar{\psi}\theta\theta, \quad (4.29)$$

$$\partial_\mu\psi\sigma^\mu\bar{\theta}\bar{\theta}\bar{\psi} = -\frac{1}{2}\partial_\mu\psi\sigma^\mu\bar{\psi}\bar{\theta}\bar{\theta}, \quad (4.30)$$

gives

$$\begin{aligned} \int d^4\theta \Phi^\dagger \Phi &= -\frac{1}{4}A\Box A^* - \frac{1}{4}A^*\Box A + \frac{1}{2}\partial^\mu A\partial_\mu A^* + |F|^2 \\ &\quad - \frac{i}{2}\psi\sigma^\mu\partial_\mu\bar{\psi} + \frac{i}{2}\partial_\mu\psi\sigma^\mu\bar{\psi}. \end{aligned}$$

Now, since

$$\frac{1}{2}\partial^\mu A\partial_\mu A^* = \frac{1}{2}\partial^\mu(A\partial_\mu A^*) - \frac{1}{2}A\Box A^*, \quad (4.31)$$

and

$$\begin{aligned} -\frac{1}{4}A\Box A^* &= -\frac{1}{4}\partial_\mu(A\partial^\mu A^*) + \frac{1}{4}\partial_\mu A\partial^\mu A^* \\ &= -\frac{1}{4}\partial_\mu(A\partial^\mu A^*) + \frac{1}{4}\partial^\mu((\partial_\mu A)A^*) - \frac{1}{4}A^*\Box A, \end{aligned}$$

we can write

$$-\frac{1}{4}A\Box A^* - \frac{1}{4}A^*\Box A + \frac{1}{2}\partial^\mu A\partial_\mu A^* = -A^*\Box A + \text{total derivatives}. \quad (4.32)$$

Using (2.24) we can similarly write

$$\begin{aligned} -\frac{i}{2}\psi\sigma^\mu\partial_\mu\bar{\psi} + \frac{i}{2}\partial_\mu\psi\sigma^\mu\bar{\psi} &= \frac{i}{2}\partial_\mu\bar{\psi}\bar{\sigma}^\mu\psi - \frac{i}{2}\bar{\psi}\bar{\sigma}^\mu\partial_\mu\psi \\ &= i\partial_\mu\bar{\psi}\bar{\sigma}^\mu\psi - \frac{i}{2}\partial_\mu(\bar{\psi}\bar{\sigma}^\mu\psi). \end{aligned} \quad (4.33)$$

Removing the terms with total derivatives the kinetic term is

$$\int d^4\theta \Phi^\dagger \Phi = \int d^4x \{-A^*(x)\Box A(x) + |F(x)|^2 + i\partial_\mu\psi(x)\bar{\sigma}^\mu\psi(x)\}. \quad (4.34)$$

The (left-handed) superpotential terms are

$$\begin{aligned} \int d^4\theta \bar{\theta}\bar{\theta}\Phi &= F, \\ \int d^4\theta \bar{\theta}\bar{\theta}\Phi\Phi &= 2AF + \psi\psi, \\ \int d^4\theta \bar{\theta}\bar{\theta}\Phi\Phi\Phi &= 3A^2F + 3A\psi\psi. \end{aligned}$$

where we have used (2.23) to rewrite the terms with ψ . The total action is then

$$S = \int d^4x \left\{ -A^*(x) \square A(x) + |F(x)|^2 + i\partial_\mu \bar{\psi}(x) \bar{\sigma}^\mu \psi(x) + gF(x) + 2mA(x)F(x) + m\psi(x)\psi(x) + 3\lambda A(x)^2 F(x) + 3\lambda A(x)\psi(x)\psi(x) + h.c. \right\} \quad (4.35)$$

The equation of motion for F is then

$$\frac{\partial \mathcal{L}}{\partial F^*} = F + g + 2mA^* + 3\lambda A^{*2} = 0, \quad (4.36)$$

which we can solve for F . With $\frac{\partial W}{\partial A} = g + 2mA + 3\lambda A^2$, $F = -\frac{\partial W^*}{\partial A^*}$ and the action is

$$S = \int d^4x \left\{ -A^*(x) \square A(x) + i\partial_\mu \psi(x) \bar{\sigma}^\mu \psi(x) + m\psi(x)\psi(x) + 3\lambda A(x)\psi(x)\psi(x) + m\bar{\psi}(x)\bar{\psi}(x) + 3\lambda A^*(x)\bar{\psi}(x)\bar{\psi}(x) - \left| \frac{\partial W}{\partial A} \right|^2 \right\}. \quad (4.37)$$

Using the fermionic mass matrices in (4.20) this can be written as

$$S = \int d^4x \left\{ -A^*(x) \square A(x) + i\partial_\mu \bar{\psi}(x) \bar{\sigma}^\mu \psi(x) + \frac{1}{2} \frac{\partial^2 W}{\partial A^2} \psi(x)\psi(x) + \frac{1}{2} \frac{\partial^2 W^*}{\partial A^{*2}} \bar{\psi}(x)\bar{\psi}(x) - \left| \frac{\partial W}{\partial A} \right|^2 \right\}. \quad (4.38)$$

Exercise 4.2

To calculate the field strength in Wess–Zumino gauge we start from the vector superfield itself

$$V_{WZ}(x, \theta, \bar{\theta}) = \theta \sigma^\mu \bar{\theta} V_\mu(x) + \theta \theta \bar{\theta} \bar{\lambda}(x) + \bar{\theta} \bar{\theta} \theta \lambda(x) + \theta \theta \bar{\theta} \bar{\theta} D(x), \quad (4.41)$$

and make the coordinate change $y^\mu = x^\mu + i\theta \sigma^\mu \bar{\theta}$. This gives

$$\begin{aligned} V_{WZ}(x, \theta, \bar{\theta}) &= \theta \sigma^\mu \bar{\theta} V_\mu(y) - i\theta \sigma^\nu \bar{\theta} \partial_\nu \theta \sigma^\mu \bar{\theta} V_\mu(y) + \theta \theta \bar{\theta} \bar{\lambda}(y) + \bar{\theta} \bar{\theta} \theta \lambda(y) + \theta \theta \bar{\theta} \bar{\theta} D(y) \\ &= \theta \sigma^\mu \bar{\theta} V_\mu(y) + \theta \theta \bar{\theta} \bar{\lambda}(y) + \bar{\theta} \bar{\theta} \theta \lambda(y) + \theta \theta \bar{\theta} \bar{\theta} [D(y) - \frac{i}{2} \partial^\mu V_\mu(y)], \end{aligned} \quad (4.42)$$

where we have used (2.26).

Since, in this coordinate system, $D_A = \partial_A + 2i(\sigma^\mu \bar{\theta})_A \partial_\mu$ and $\bar{D}_{\dot{A}} = -\partial_{\dot{A}}$, so that $\bar{D} \bar{D} = \partial_{\dot{A}} \partial^{\dot{A}}$, we can then find the field strength as

$$\begin{aligned} W_A &= -\frac{1}{4} \bar{D} \bar{D} D_A V_{WZ}(y, \theta, \bar{\theta}) \\ &= -\frac{1}{4} \partial_{\dot{A}} \partial^{\dot{A}} [(\sigma^\mu \bar{\theta})_A V_\mu(y) + i\theta_A \bar{\theta} \bar{\theta} \partial^\mu V_\mu(y) + 2\theta_A \bar{\theta} \bar{\lambda}(y) + \bar{\theta} \bar{\theta} \lambda_A(y) + 2\theta_A \bar{\theta} \bar{\theta} D(y) \\ &\quad + 2i(\sigma^\nu \bar{\theta})_A \partial_\nu (\theta \sigma^\mu \bar{\theta} V_\mu(y) + \theta \theta \bar{\theta} \bar{\lambda}(y))] \\ &= -\frac{1}{4} \partial_{\dot{A}} \partial^{\dot{A}} [(\sigma^\mu \bar{\theta})_A V_\mu(y) + i\theta_A \bar{\theta} \bar{\theta} \partial^\mu V_\mu(y) + 2\theta_A \bar{\theta} \bar{\lambda}(y) + \bar{\theta} \bar{\theta} \lambda_A(y) + 2\theta_A \bar{\theta} \bar{\theta} D(y) \\ &\quad + 2i(\sigma^\nu \bar{\theta})_A \theta \sigma^\mu \bar{\theta} \partial_\nu V_\mu(y) - i\theta \theta \bar{\theta} \bar{\theta} (\sigma^\nu \partial_\nu \bar{\lambda}(y))_A] \\ &= i\theta_A \partial^\mu V_\mu(y) + \lambda_A(y) + 2\theta_A D(y) \\ &\quad + 2i\partial_{\dot{A}} \partial^{\dot{A}} (\sigma^\nu \bar{\theta})_A \theta \sigma^\mu \bar{\theta} \partial_\nu V_\mu(y) - i\theta \theta (\sigma^\nu \partial_\nu \bar{\lambda}(y))_A \\ &= \lambda_A(y) + 2\theta_A D(y) + (\sigma^{\mu\nu} \theta)_A F_{\mu\nu}(y) - i\theta \theta (\sigma^\mu \partial_\mu \bar{\lambda}(y))_A \end{aligned} \quad (4.43)$$

where we have used $\partial_A \partial^A (\bar{\theta}\bar{\theta}) = 4$ and $2i(\sigma^\nu \bar{\theta})_A \theta \bar{\theta} \partial_\nu \bar{\lambda}(y) = -i\theta \bar{\theta} \bar{\theta} (\sigma^\nu \partial_\nu \bar{\lambda}(y))_A$, and where $F_{\mu\nu} = \partial_\mu V_\nu - \partial_\nu V_\mu$ is the field strength for the field V_μ .

Then the terms with a $\theta\theta$ factor in $W^A W_A$ are

$$\begin{aligned}
W^A W_A|_{\theta\theta} &= -i\theta\theta\lambda(y)\sigma^\mu\partial_\mu\bar{\lambda}(y) - i\theta\theta(\sigma^\mu\partial_\mu\bar{\lambda}(y))^A\lambda_A(y) + 4\theta\theta D^2(y) \\
&\quad + 2D(y)\theta\sigma^{\mu\nu}\theta F_{\mu\nu}(y) + 2(\sigma^{\mu\nu}\theta)^A F_{\mu\nu}(y)\theta_A D(y) + (\sigma^{\mu\nu}\theta)^A F_{\mu\nu}(y)(\sigma^{\rho\sigma}\theta)_A F_{\rho\sigma}(y) \\
&= -2i\theta\theta\lambda(y)\sigma^\mu\partial_\mu\bar{\lambda}(y) + 4\theta\theta D^2(y) \\
&\quad + 4\theta\sigma^{\mu\nu}\theta D(y)F_{\mu\nu}(y) + (\sigma^{\mu\nu}\theta)^A F_{\mu\nu}(y)(\sigma^{\rho\sigma}\theta)_A F_{\rho\sigma}(y) \\
&= -2i\theta\theta\lambda(y)\sigma^\mu\partial_\mu\bar{\lambda}(y) + 4\theta\theta D^2(y) - \theta\theta\frac{1}{2}F_{\mu\nu}(y)F^{\mu\nu}(y) - \theta\theta\frac{i}{2}F_{\mu\nu}(y)\tilde{F}^{\mu\nu}(y) \quad (4.44)
\end{aligned}$$

where we have used that from (2.29) $\theta\sigma^{\mu\nu}\theta = 0$ to remove a term, and rewritten

$$(\sigma^{\mu\nu}\theta)^A F_{\mu\nu}(y)(\sigma^{\rho\sigma}\theta)_A F_{\rho\sigma}(y) = -\theta\theta\frac{1}{2}F_{\mu\nu}(y)F^{\mu\nu}(y) - \theta\theta\frac{i}{2}F_{\mu\nu}(y)\tilde{F}^{\mu\nu}(y), \quad (4.45)$$

where $\tilde{F}^{\mu\nu} = \frac{1}{2}\epsilon^{\mu\nu\rho\sigma}F_{\rho\sigma}$ is the dual field strength tensor.

Since these terms are invariant under the change of coordinates $y^\mu = x^\mu + i\theta\sigma^\mu\bar{\theta}$ – because all terms beyond the first order disappear in the expansion due to too many θ -factors – and since no other terms can survive the $\int d^4\theta$ -integration in the Lagrangian, we are left with

$$\int d^4\theta \bar{\theta}\bar{\theta} W^A W_A = -2i\lambda(x)\sigma^\mu\partial_\mu\bar{\lambda}(x) + 4D^2(x) - \frac{1}{2}F_{\mu\nu}(x)F^{\mu\nu}(x) - \frac{i}{2}F_{\mu\nu}(x)\tilde{F}^{\mu\nu}(x). \quad (4.46)$$

We see that the component field Lagrangian, as expected and as promised in Sec. 4.4, contains a field strength term for the vector component field. At this point we may worry about the appearance of the dual field strength tensor which we do not have in the Standard Model. However, we can show that the term in question can be written as a total derivative,

$$\begin{aligned}
F_{\mu\nu}\tilde{F}^{\mu\nu} &= \frac{1}{2}(\partial_\mu A_\nu - \partial_\nu A_\mu)\epsilon^{\mu\nu\rho\sigma}(\partial_\rho A_\sigma - \partial_\sigma A_\rho) \\
&= 2\epsilon^{\mu\nu\rho\sigma}(\partial_\mu A_\nu)(\partial_\rho A_\sigma) \\
&= 2\epsilon^{\mu\nu\rho\sigma}(\partial_\rho\partial_\mu A_\nu)A_\sigma - 2\epsilon^{\mu\nu\rho\sigma}\partial_\rho(\partial_\mu A_\nu A_\sigma) \quad (4.47)
\end{aligned}$$

where the first term disappears due to the asymmetry of the Levi-Civita symbol. That means that the term with the dual field strength tensor disappears if the fields fall off rapidly enough towards infinity. This will always be true for abelian field theories, but is not necessarily so for non-abelian theories, where we can have so-called instanton effects.

Exercise 4.3

We should note that in this case $W = 0$ as we can not construct gauge invariant terms in the superpotential with only one superfield that is charged under the gauge group. This also implies that $F(x) \equiv 0$.

For an abelian gauge group the kinetic term in Wess–Zumino gauge is

$$\mathcal{L}_{\text{kin}} = \int d^4\theta \Phi^\dagger e^{qV_{WZ}} \Phi = \int d^4\theta \left[\Phi^\dagger \Phi + q\Phi^\dagger \Phi V_{WZ} + \frac{1}{2}q^2 \Phi^\dagger \Phi V_{WZ}^2 \right]. \quad (4.49)$$

The first term in the integral we have calculated already in Exercise 1. Inserting Eqs. (3.25), (3.26) and (3.37) the second term is

$$\begin{aligned}
\int d^4\theta \Phi^\dagger \Phi V_{WZ} &= \int d^4\theta \left\{ [iA^*(\theta\sigma^\nu\bar{\theta})\partial_\nu A - iA(\theta\sigma^\nu\bar{\theta})\partial_\nu A^* + 2\bar{\theta}\bar{\psi}\theta\psi](\theta\sigma^\mu\bar{\theta})V_\mu \right. \\
&\quad \left. + \sqrt{2}A\bar{\theta}\bar{\psi}\theta\bar{\theta}\bar{\lambda} + \sqrt{2}A^*\theta\psi\bar{\theta}\bar{\theta}\theta\lambda + \theta\bar{\theta}\bar{\theta}\bar{\theta}|A|^2 D \right\} \\
&= \int d^4\theta \left\{ \frac{i}{2}A^*\theta\bar{\theta}\bar{\theta}\bar{\theta}\partial^\mu AV_\mu - \frac{i}{2}A\theta\bar{\theta}\bar{\theta}\bar{\theta}\partial^\mu A^*V_\mu - \frac{1}{2}(\theta\bar{\theta})(\bar{\theta}\bar{\theta})\bar{\psi}\bar{\sigma}^\mu\psi V_\mu \right. \\
&\quad \left. + \frac{1}{\sqrt{2}}A\theta\bar{\theta}\bar{\theta}\bar{\psi}\bar{\lambda} + \frac{1}{\sqrt{2}}A^*\theta\bar{\theta}\bar{\theta}\bar{\psi}\lambda + \theta\bar{\theta}\bar{\theta}\bar{\theta}|A|^2 D \right\} \\
&= \frac{i}{2}(A^*\partial^\mu A - A\partial^\mu A^*)V_\mu - \frac{1}{2}\bar{\psi}\bar{\sigma}^\mu\psi V_\mu \\
&\quad + \frac{1}{\sqrt{2}}A\bar{\psi}\bar{\lambda} + \frac{1}{\sqrt{2}}A^*\psi\lambda + |A|^2 D,
\end{aligned} \tag{4.50}$$

where we have used Eqs. (2.23), (2.26), and (4.48) for the second equality.

The third term is found from Eqs. (3.25), (3.26) and (3.38), with only one possible surviving term

$$\int d^4\theta \frac{1}{2}\Phi^\dagger \Phi V_{WZ}^2 = \int d^4\theta \frac{1}{4}|A|^2\theta\bar{\theta}\bar{\theta}\bar{\theta}V^\mu V_\mu = \frac{1}{4}|A|^2 V^\mu V_\mu. \tag{4.51}$$

Summing up these three terms the total contribution from the kinetic terms is then

$$\begin{aligned}
\mathcal{L}_{\text{kin}} &= -A^*\square A + i\partial_\mu\bar{\psi}\bar{\sigma}^\mu\psi + \frac{i}{2}q(A^*\partial^\mu A - A\partial^\mu A^*)V_\mu - \frac{1}{2}q\bar{\psi}\bar{\sigma}^\mu\psi V_\mu \\
&\quad + \frac{1}{\sqrt{2}}qA\bar{\psi}\bar{\lambda} + \frac{1}{\sqrt{2}}qA^*\psi\lambda + q|A|^2 D + \frac{1}{4}q^2|A|^2 V^\mu V_\mu.
\end{aligned} \tag{4.52}$$

We can simplify this expression by defining a covariant derivative from the vector component field $D_\mu \equiv \partial_\mu - i\frac{q}{2}V_\mu$.²⁰ We immediately see that

$$i\partial_\mu\bar{\psi}\bar{\sigma}^\mu\psi - \frac{1}{2}q\bar{\psi}\bar{\sigma}^\mu\psi V_\mu = iD_\mu^*\bar{\psi}\bar{\sigma}^\mu\psi, \tag{4.53}$$

and

$$\begin{aligned}
|D_\mu A|^2 &\equiv D_\mu^* A^* D^\mu A = (\partial_\mu A^* + i\frac{q}{2}V_\mu A^*)(\partial^\mu A - i\frac{q}{2}V^\mu A) \\
&= |\partial_\mu A|^2 + \frac{q^2}{4}V_\mu V^\mu |A|^2 + i\frac{q}{2}(A^*\partial^\mu A - \partial^\mu A^* A)V_\mu,
\end{aligned} \tag{4.54}$$

so that we can write, using $-A^*\square A = |\partial_\mu A|^2$,

$$\mathcal{L}_{\text{kin}} = \frac{1}{\sqrt{2}}qA\bar{\psi}\bar{\lambda} + \frac{1}{\sqrt{2}}qA^*\psi\lambda + q|A|^2 D + iD_\mu^*\bar{\psi}\bar{\sigma}^\mu\psi + |D_\mu A|^2$$

²⁰The unusual factor of 2 here compared to the usual abelian gauge covariant derivative is due to our earlier choices in how to write down the vector superfield in terms of component fields. If we want we can absorb this by redefining V_μ .

Adding the field strength terms from Exercise 2 the complete Lagrangian is

$$\begin{aligned}\mathcal{L} = & iD_\mu^* \bar{\psi}(x) \bar{\sigma}^\mu \psi(x) + |D_\mu A(x)|^2 - i\lambda(x) \sigma^\mu \partial_\mu \bar{\lambda}(x) \\ & + \frac{1}{\sqrt{2}} q A(x) \bar{\psi}(x) \bar{\lambda}(x) + \frac{1}{\sqrt{2}} q A^*(x) \psi(x) \lambda(x) + q |A(x)|^2 D(x) \\ & + 2D^2(x) - \frac{1}{4} F_{\mu\nu}(x) F^{\mu\nu}(x).\end{aligned}\quad (4.55)$$

Exercise 4.4

In QED we have three particles, the electron, positron and the photon. As a Dirac fermion the electron needs two scalar superfields to provide the two different left- and right-handed Weyl-spinor components. However, then the positron is already given by the hermitian conjugates of these spinors. For the photon we need an abelian vector superfield for the $U(1)_{\text{em}}$ gauge group. Thus the minimum field content consists of two scalar superfields L and $\bar{E}$, and a vector superfield V .

The supergauge transformations for the scalar superfields are $L \rightarrow L' = \exp(-ie\Lambda)L$ and $\bar{E} \rightarrow \bar{E}' = \exp(ie\Lambda)\bar{E}$, where e is the elementary electrical charge, and the difference in sign signifies the opposite charges of the electron and positron. The general form of the superpotential must then be $W = mLE$ with a single (mass) parameter m . No tadpole terms in L and $\bar{E}$ can survive since these are not gauge singlets, and only the LE mass term survives since the two fields in the mass term must have opposite charges under the gauge group. Similarly, no Yukawa term can fulfil the gauge invariance criterium.

Solving for the auxiliary F -fields we get

$$\begin{aligned}F_L^* &= -W_L = -\frac{\partial W[A_L, A_E]}{\partial A_L} = -mA_E, \\ F_E^* &= -W_E = -\frac{\partial W[A_L, A_E]}{\partial A_E} = -mA_L.\end{aligned}$$

Generalising the discussion in Exercise 1, the contribution to the Lagrangian from the superpotential in terms of component fields is

$$\mathcal{L}_{\text{superpot}} = m\psi_L\psi_E + m^*\bar{\psi}_L\bar{\psi}_E - |m|^2(|A_L|^2 + |A_E|^2), \quad (4.56)$$

where we must reasonably identify m with the electron mass, see also that, as expected the scalar fields A_E and A_L obtain the same mass.

From the field content we know that the kinetic terms take the form $L^\dagger e^{eV} L$ and $\bar{E}^\dagger e^{-eV} \bar{E}$ where the sign ensures the gauge invariance under the supergauge transformations of the scalar and vector superfields together. Taking results from Exercise 3 the Lagrangian contribution from the kinetic terms is

$$\begin{aligned}\mathcal{L}_{\text{kin}} = & \frac{1}{\sqrt{2}} e A_L \bar{\psi}_L \bar{\lambda} + \frac{1}{\sqrt{2}} e A_L^* \psi_L \lambda + e |A_L|^2 D + i D_\mu^* \bar{\psi}_L \bar{\sigma}^\mu \psi_L + |D_\mu A_L|^2 \\ & - \frac{1}{\sqrt{2}} e A_E \bar{\psi}_E \bar{\lambda} - \frac{1}{\sqrt{2}} e A_E^* \psi_E \lambda - e |A_E|^2 D + i D_\mu^* \bar{\psi}_E \bar{\sigma}^\mu \psi_E + |D_\mu A_E|^2,\end{aligned}$$

where λ is the Weyl fermion from the vector superfield V , D is the auxiliary field from V , and the covariant gauge derivative is $D_\mu = \partial_\mu - \frac{ie}{2} A_\mu$, where A_μ is the electromagnetic gauge field.

In Exercise 2 we saw that the contribution from the abelian field strength term is

$$\mathcal{L}_{\text{field strength}} = -i\lambda\sigma^\mu\partial_\mu\bar{\lambda} + 2D^2 - \frac{1}{4}F_{\mu\nu}F^{\mu\nu}, \quad (4.57)$$

where $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$.

Solving for the auxiliary D -field we get

$$D = \frac{e}{4}(|A_E|^2 - |A_L|^2). \quad (4.58)$$

This gives a total Lagrangian of

$$\begin{aligned} \mathcal{L} = & iD_\mu^*\bar{\psi}_L\bar{\sigma}^\mu\psi_L + |D_\mu A_L|^2 + iD_\mu^*\bar{\psi}_E\bar{\sigma}^\mu\psi_E + |D_\mu A_E|^2 - i\lambda\sigma^\mu\partial_\mu\bar{\lambda} - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} \\ & + m\psi_L\psi_E + m^*\bar{\psi}_L\bar{\psi}_E - |m|^2(|A_L|^2 + |A_E|^2) \\ & + \frac{1}{\sqrt{2}}eA_L\bar{\psi}_L\bar{\lambda} + \frac{1}{\sqrt{2}}eA_L^*\psi_L\lambda - \frac{1}{\sqrt{2}}eA_E\bar{\psi}_E\bar{\lambda} - \frac{1}{\sqrt{2}}eA_E^*\psi_E\lambda - \frac{e^2}{8}(|A_E|^2 - |A_L|^2)^2. \end{aligned}$$

Exercise 7.2

We will in the following take all outgoing momenta to go out of the vertex. The indices $abcd$ are gluon indices $(1, \dots, 8)$, rs are spin indices $(1, 2)$, and $ijmn$ are colour indices $(1, 2, 3)$. The relevant Feynman rules are as follows:

- Incoming quark: $u^r(k)$.
- Incoming antiquark: $\bar{v}^s(p)$.
- Gluon propagator: $-i\frac{g_{\mu\nu}}{s}\delta^{ab}$.
- Vertex $q\bar{q}g$: $-it_{ij}^a\gamma^\mu g_s$.
- Vertex $\tilde{q}\tilde{q}^*g$: $-it_{mn}^b(k' - p')^\nu g_s$.

We will assume the SM particles to have negligible mass compared to the squarks.

The matrix element is then given as

$$\mathcal{M} = -\frac{g_s^2}{s}t_{ij}^a t_{mn}^b \delta^{ab} \bar{v}^s \gamma^\mu u^r (k' - p')_\mu.$$

In the squared amplitude we average over all incoming spin and colour, and sum over the outgoing:

$$\begin{aligned} |\bar{\mathcal{M}}|^2 &= \frac{1}{4} \cdot \frac{1}{9} \frac{g_s^4}{s^2} \sum_{ab} \sum_{cd} \sum_{rs} \sum_{ijmn} t_{ij}^a t_{mn}^b t_{ij}^c t_{mn}^d \delta^{ab} \delta^{cd} \bar{v}^s \gamma^\mu u^r \bar{u}^r \gamma^\nu v^s (k' - p')_\mu (k' - p')_\nu \\ &= \frac{1}{4} \cdot \frac{1}{9} \frac{g_s^4}{s^2} \sum_{ijmn} \underbrace{\left(\sum_a t_{ij}^a t_{mn}^a \sum_c t_{ji}^c t_{nm}^c \right)}_{\equiv C_f} \sum_{rs} \bar{v}^s \gamma^\mu u^r \bar{u}^r \gamma^\nu v^s (k' - p')_\mu (k' - p')_\nu \\ &= \frac{1}{4} \cdot \frac{1}{9} \frac{g_s^4}{s^2} C_f p_\alpha k_\beta \text{TR}(\gamma^\alpha \gamma^\mu \gamma^\beta \gamma^\nu) (k' - p')_\mu (k' - p')_\nu \\ &= \frac{1}{9} \frac{g_s^4}{s^2} C_f p_\alpha k_\beta (p^\mu k^\nu - p^\beta k_\beta \eta^{\mu\nu} + p^\nu k^\mu) (k' - p')_\mu (k' - p')_\nu \\ &= \frac{1}{9} \frac{g_s^4}{s^2} C_f (2p \cdot (k' - p') k \cdot (k' - p') - p \cdot k (k' - p') \cdot (k' - p')), \end{aligned}$$

where we have isolated the colour factors into the coefficient C_f .

In the centre of mass frame, we have $p = (E, \vec{p})$, $k = (E, -\vec{p})$, $p' = (E, \vec{p}')$ and $k' = (E, -\vec{p}')$. From this we find

$$\begin{aligned} 2p \cdot (k' - p') k \cdot (k' - p') - p \cdot k (k' - p') \cdot (k' - p') &= 2 \cdot 4|\vec{p}|^2 |\vec{p}'|^2 (1 - \cos^2 \theta) \\ &= 2s |\vec{p}'|^2 (1 - \cos^2 \theta), \end{aligned}$$

where θ is the acute angle between the incoming quarks, and $s = (p + k)^2 = 4E^2 = 4|\vec{p}|^2$. This gives the squared averaged amplitude

$$|\bar{\mathcal{M}}|^2 = \frac{2}{9} \frac{g_s^4}{s} C_f \cdot |\vec{p}'|^2 (1 - \cos^2 \theta),$$

and the differential cross section

$$\begin{aligned} \frac{d\sigma}{d\Omega} &= \frac{|\vec{p}'|^2}{32\pi^2 \sqrt{ss}} |\bar{\mathcal{M}}|^2 \\ &= \frac{1}{144} \frac{g_s^4}{\pi^2 \sqrt{ss^2}} C_f |\vec{p}'|^3 (1 - \cos^2 \theta). \end{aligned}$$

Integrating over the solid angle then gives the total cross section

$$\begin{aligned} \sigma &= \int \left(\frac{d\sigma}{d\Omega} \right) d\Omega \\ &= 2\pi \frac{1}{144} \frac{g_s^4}{\pi^2 \sqrt{ss^2}} C_f |\vec{p}'|^3 \int_{-1}^1 (1 - \cos^2 \theta) d(\cos \theta) \\ &= \frac{4}{3} 2\pi \frac{1}{144} \frac{g_s^4}{\pi^2 \sqrt{ss^2}} C_f |\vec{p}'|^3 \\ &= \frac{1}{54\pi} \frac{g_s^4}{\sqrt{ss^2}} C_f |\vec{p}'|^3. \end{aligned}$$

We can rewrite $|\vec{p}'|^3$ by noticing that

$$|\vec{p}'| = \sqrt{E^2 - m^2} = \frac{1}{2} \sqrt{4E^2 - 4m^2} = \frac{\sqrt{s}}{2} \sqrt{1 - \frac{4m^2}{s}}.$$

The colour factor is calculated below and found to be $C_f = 2$, so the total cross section is

$$\sigma = \frac{g_s^4}{216\pi s} \sqrt{\left(1 - \frac{4m^2}{s}\right)^3}.$$

Using $\alpha_s = \frac{g_s^2}{4\pi}$, and assuming that both the left-, and right-handed squarks have the same mass, we arrive at the final expression

$$\sigma = \frac{4}{27} \frac{\pi \alpha_s^2}{s} \sqrt{\left(1 - \frac{4m^2}{s}\right)^3}.$$

To calculate the colour factor C_f we use that the sum over the generators t is given (see for example [37]) as:

$$\sum_a t_{ij}^a t_{mn}^a = \frac{1}{2} \left(\delta_{in} \delta_{jm} - \frac{1}{N_C} \delta_{ij} \delta_{mn} \right),$$

and using that $\delta_{ij} = \delta_{ji}$, we have

$$\begin{aligned} C_f &= \frac{1}{4} \sum_{ijmn} \left(\delta_{in} \delta_{jm} - \frac{1}{N_C} \delta_{ij} \delta_{mn} \right) \left(\delta_{jm} \delta_{in} - \frac{1}{N_C} \delta_{ji} \delta_{nm} \right) \\ &= \frac{1}{4} \sum_{ijmn} \left(\delta_{in} \delta_{jm} \delta_{jm} \delta_{in} - \frac{2}{N_C} \delta_{ij} \delta_{mn} \delta_{jm} \delta_{in} + \frac{1}{N_C^2} \delta_{ij} \delta_{mn} \delta_{ji} \delta_{nm} \right) \\ &= \frac{1}{4} \sum_{ijmn} \left(\delta_{in} \delta_{in} \delta_{jm} \delta_{jm} - \frac{2}{N_C} \delta_{ij} \delta_{mn} \delta_{jm} \delta_{in} + \frac{1}{N_C^2} \delta_{ij} \delta_{ij} \delta_{mn} \delta_{mn} \right) \\ &= \frac{1}{4} \left(N_C^2 - \frac{2}{N_C} N_C + \frac{1}{N_C^2} N_C^2 \right) \\ &= \frac{1}{4} (N_C^2 - 1) = 2. \end{aligned}$$

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