Chapter 2
The Standard Model and Beyond

It is often said that the Standard Model (SM) is a theory of interactions.[1] That means, that it describes the laws of nature by assigning its pieces a susceptibility to certain forces. This is modelled as a charge with respect to a field, which in this respect is nothing more than a quantum of how strongly it couples to the force carriers of that field.

The most familiar of charges is probably electric charge. Consider how static electricity separates the straws of your hair—this happens when there are a lot of same-sign charges repelling each other, a large total charge.[2] It does not happen when there are only local fluctuations up and down in charge, as there normally is (they largely cancel). The same way, the magnitude of the charge on a fundamental particle determines how strongly it is coupled to the corresponding field.

But how do the straws of your hair know about the electric charge of their neighbours? Well, the charge is communicated by the exchange of a messenger: a field quantum. The field quantum of electromagnetism is the photon—a particle of light. In every interaction in the SM, a field quantum is exchanged. These are commonly called gauge bosons. The different forces of nature in the SM all correspond to their own field, and are communicated with each their own set of gauge bosons. For gravity to fit into this picture, it too should probably be mediated by a particle: the stipulated *graviton*, which remains to be observed. In fact, that it is not observed, and that mass (the coupling to gravity) is not quantised, indicates that gravity cannot yet be described as a *quantum field theory* like the other forces of nature.[3] From now on, we will not consider gravity further, and as a matter of fact, we can safely neglect

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[1]For a general introduction to the Standard Model, see for instance the review in [1], and references therein.

[2]A net charge arises as the hair is stripped of or receives *electrons*—fundamental particles with electric charge $-1e$. Unlike a compound object, a fundamental particle has an intrinsic, fixed charge.

[3]This could be an indication of a more fundamental theory than the SM.

L.K. Bryngemark, *Search for New Phenomena in Dijet Angular Distributions at $\sqrt{s}$ = 8 and 13 TeV*, Springer Theses, DOI 10.1007/978-3-319-67346-2_2

Table 2.1 The four fundamental interactions currently known, their strength relative to the strong interaction at their respective appropriate scale, and range in metres [2]

Force	Relative strength	Range (m)
Strong	1	10^{-15}
Electromagnetic	$\frac{1}{137}$	∞
Weak	10^{-6}	10^{-18}
Gravity	10^{-39}	∞

it, as it is many orders of magnitude weaker than the other three known forces of nature, which completely dominate particle interactions.

Moving from macroscopic compound objects like a straw of hair, the fundamental particles the SM deals with are *fermions* and *bosons*, with half-integer and integer (including zero) *spin*, respectively. Like charge, spin is a quantum number intrinsic to the particle, and it has a sign (is a directional quantity). In addition, a particle may carry charge under several fields, and thus interact with several forces. The combination of quantum numbers (spin type and charges) and mass[4] uniquely defines a fundamental particle. In total, the SM describes the interactions of 17 fundamental particles. The interactions and their range and relative strengths are listed in Table 2.1.

Although the table lists the properties of the fundamental interactions, let me immediately introduce a caveat. It so happens, that the strength of the interactions depends on the energy scale at which the interactions are probed. This is called "running of the coupling constants" and actually implies that at certain energies, forces can unite (unless they evolve exactly the same way). For instance at energy scales accessible to today's particle physics experiments, we often refer to electroweak[5] (EW) interactions.

As mentioned, the SM is a theory of interactions, and it is through the laws of interaction we can distinguish the particles. I will thus introduce the fundamental particles in the SM in terms of the interactions. It will become evident that some interactions and prediction techniques are more relevant to my work, as they will be described in greater detail, and will serve as a use-case for some of the general features of the SM formalism. Mathematically, the SM is also a theory of symmetries; from symmetries, interactions and conservation laws arise. Conservation laws have profound implications on the interpretation of the theory, but are also part of our experimental tool-box, as they allow us to deduce certain quantities that aren't directly observed.

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[4]Here it is again, the elusive, seemingly fundamental, concept of mass.

[5]*Electroweak* as in the unification of electromagnetic and weak interactions.

2.1 Electromagnetism: QED

Magnetism has been known by humanity for thousands of years, and even used (e.g. for navigation). Electricity was understood as a force much later, in the 19th century. The electron would be the first particle discovered which is still considered fundamental.

In the quantum world, electromagnetism is described by Quantum ElectroDynamics (QED). Its mediating gauge boson is the photon (often represented by a γ (gamma)). It is an infinite-range force, since the mediator is mass- and chargeless. This is the force which keeps atoms together, from the opposite electric charge sign of electrons and atomic nuclei. It also governs the electromagnetic waves we encounter in our everyday lives in form of radio (cell phone) signals, visible light or X-rays.

QED is one of the most tested theories we have—that is, we can both predict and measure quantities very precisely. The energy in an atomic energy level transition in hydrogen is often quoted as an example, as it is measured to 14 digits [3]! Yet, as we shall see, it is not a complete theory to all scales.

2.1.1 The Charged Leptons

Here we encounter our first matter particle type: the electrically charged leptons. One of these, the lightest, is the aforementioned electron (e). It partly makes up matter as we know it in our everyday life. However, it has heavier cousins: the muon, μ, and the tau lepton, τ. These cousins have different *flavour*, and different mass, but apart from that they are similar. Flavour is a quantum number that is conserved under the electromagnetic interaction. The charged leptons have unit electric charge.[6]

2.2 The Weak (Nuclear) Interaction

The weak interaction is suitably named, as it is substantially weaker than both the strong and electromagnetic interaction. It is mediated via massive vector bosons, the electrically charged W and the neutral Z boson, and unlike electromagnetism, it can transform particles into a cousin of different flavour. The masses of the gauge bosons make it a short range force. The weak interaction charge is called *weak isospin*,[7] and it is only carried by particles of *left-handed chirality*.

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[6] The electron being the first fundamental particle discovered, it set the standard for electric charge—as the name suggests.

[7] In the unified electroweak force, the charge is instead *weak hypercharge*, which takes both weak isospin and electric charge into account.

A particle of right-handed helicity is one where spin orientation and direction of motion coincides, while for a left-handed particle these two are opposite. This means that handedness depends on the reference frame of the observer. For massless particles, there is no choice of two frames with respect to which the massless particle can appear to move in opposite directions, since no observer can travel faster than the particle. Thus they are always of definite helicity, which coincides with its chirality. For massive particles, only chirality is invariant of choice of reference frame. This "handedness" or chirality is necessary to explain certain experimental observations, such as parity violation.[8]

2.2.1 The Neutral Leptons

Along with the weak interaction, the need for neutral leptons—*neutrinos*—arises. They are ordered in flavour doublets[9] together with the charged leptons as illustrated below, in order of increasing mass:

$$\begin{pmatrix} e \\ \nu_e \end{pmatrix}_L \quad \begin{pmatrix} \mu \\ \nu_\mu \end{pmatrix}_L \quad \begin{pmatrix} \tau \\ \nu_\tau \end{pmatrix}_L$$

As for the neutrino masses themselves, they are too small to have been directly measured yet. That neutrinos do have mass is however established through the phenomenon of neutrino oscillations: neutrinos produced in one flavour state can oscillate into another flavour state[10] as they travel. And travel they do! Since they only carry charge under the weak interaction, they rarely interact, and are very likely to travel straight through even large macroscopic objects like planets.

The weak interaction can convert an upper particle in a doublet to its lower counterpart. This is possible since there are charged weak bosons, $W^{\pm}$, which can carry the incoming charge such that it is overall conserved. For instance, in radioactive β decay, it is the weak interaction which is at play: $n \to p + e^- + \bar{\nu}_e$ involves the exchange of a W boson. But to understand that process, we first need to introduce a set of particles commonly associated with the last known fundamental force of nature.

[8]We won't need to discuss parity further in this work, but for a historical experiment, the interested reader is referred to Ref. [4].

[9]L denotes left-handed. The right-handed counterparts are flavour singlets, and thus stand alone: $e_R, \mu_R, \ldots$.

[10]Flavour oscillations are a quantum mechanical subtlety, relating to the flavour *eigenstate* not being the same as the mass eigenstate. Oh, yes, there it is again.

2.3 The Strong (Nuclear) Interaction: QCD

In our everyday lives, the main effect of the strong interaction is to keep the atomic nuclei together. This is not a small impact! The strong interaction is however a short-range force, limited to within the size of a nucleon, and only a smaller residual force is actually felt between the nucleons.

Colour charge is the quantum number making particles susceptible to the strong interaction or colour force, described by Quantum ChromoDynamics (QCD). The colour charges are, in an analogy to the components of white light, *red, green* and *blue*, expressed below in a colour triplet:

$$\psi_a = \begin{pmatrix} \psi_1 \\ \psi_2 \\ \psi_3 \end{pmatrix} \tag{2.1}$$

The gauge boson of the strong interaction is the *gluon*. Gluons carry colour charge themselves. Thus, in contrast to QED, where the photon does not carry electric charge, two gluons can interact. This in turn makes the range of the strong interaction finite even though gluons are massless.

The QCD Lagrangian, the equation of motion describing all of the workings of the theory, is formulated in a gauge invariant way as

$$\mathcal{L} = \mathcal{L}_q + \mathcal{L}_g = \bar{\psi}_a(i\gamma^\mu \partial_\mu \delta_{ab} - g_s \gamma^\mu t^C_{ab} \mathcal{A}^C_\mu - m\delta_{ab})\psi_b - \frac{1}{4} F^{\mu\nu}_A F^A_{\mu\nu}, \tag{2.2}$$

where Eq. 2.1 enters, and the field tensor

$$F^A_{\mu\nu} = \partial_\mu \mathcal{A}^A_\nu - \partial_\nu \mathcal{A}^A_\mu + g_s f^{ABC} \mathcal{A}^B_\mu A^C_\nu \tag{2.3}$$

makes up the kinetic term in the gauge field. The third term of Eq. 2.2 makes $\bar{\psi} i \not\partial \psi$ gauge invariant. Gauge invariance is a means for making local symmetries in a theory evident, and in practice it means that a given new choice of coordinate system must be accompanied by a choice of *covariant* derivatives (the ∂_μ for instance), such that there is no net change on the predictions of the theory. The physics doesn't change! But the choice of formalism can make it more or less obscure. Since local symmetries give rise to forces, this is a central point in the Lagrangian formulation. On a similar note, global symmetries correspond to conserved currents, or put more simply, conservation laws.

In Eqs. 2.2 and 2.3, the eight[11] gluons enter in the $\mathcal{A}^1_\mu, \ldots, \mathcal{A}^8_\mu$, accompanied by the eight generators t_{ab} and the structure constants f^{ABC}. The superscripts here are colour indices implicitly summed over. From the strong coupling strength, g_s, we

[11] $8 = 3^2 - 1$, QCD being an $SU(3)$ symmetry group.

define the strong coupling constant $\alpha_s = g_s^2/(4\pi)$. The last term in Eq. 2.3 is the self-interaction term due to the colour charge of the gluons.

2.3.1 The Quarks

The six quarks are fermions—building blocks of larger compounds of particles. They carry colour charge, meaning they belong in colour triplets, and non-integer[12] electric charge: *up* (u), *charm* (c), *top* (t) have $+\frac{2}{3}e$, while *down* (d), *strange* (s) and *bottom* (b) carry $-\frac{1}{3}e$. Note that gluons carry one colour and one anti-colour, giving them the possibility to change the colour state of for instance a quark in an interaction. None of the other fermions in the SM interact via the strong interaction—they are colourless, or colour singlets. Like the leptons, the quarks also come in three generations, ordered in flavour doublets as represented below, again ordering the doublets in increasing mass:

$$\begin{pmatrix} u \\ d \end{pmatrix}_L \quad \begin{pmatrix} c \\ s \end{pmatrix}_L \quad \begin{pmatrix} t \\ b \end{pmatrix}_L$$

From this structure, it should be clear that the quarks also carry weak isospin and take part in weak interactions. However, due to the much smaller weak interaction coupling strength, QCD processes are much more probable and thus happen more often.

2.4 The Brout–Englert–Higgs Mechanism and the Particle Masses

No thesis covering work done in ATLAS in recent years would be complete without mentioning the Brout–Englert–Higgs (BEH) mechanism, and the related H boson discovered by ATLAS and CMS in 2012. This mechanism gives masses to the fermions and weak gauge bosons via the mechanism of electroweak symmetry breaking, splitting the massless gauge bosons of the underlying symmetry into the massless photon and the massive W and Z bosons, thus splitting the electroweak theory into electromagnetic and weak interaction. Knowing at which energy we have unification, we could predict approximately what the mass of the H boson should be, even though mass is always a free parameter in the SM.

In the general picture of quantised coupling strengths, the H boson is a little special since the coupling to different particles is related to their mass. Or, conversely, the mass of a particle is a measure of—given by!—how strongly it couples to the BEH

[12]Had the history of discovery been different, the electric charge of the electron had likely been defined as $-3e$ instead.

Table 2.2 The masses of fundamental particles as experimentally measured, or in most quark cases, calculated [2]. Note that the light quark masses are current quark masses, as calculated in the $\overline{\mathrm{MS}}$ scheme at a scale of 2 GeV

	Particle	Symbol	Mass
Leptons	Neutrinos	ν_e, ν_μ, ν_τ	< 25 eV
	Electron	e	511 keV
	Muon	μ	105.6 MeV
	Tau lepton	τ	1776.2 ± 0.1 MeV
Quarks	Up	u	$2.3^{+0.7}_{-0.5}$ MeV
	Down	d	$4.8^{+0.5}_{-0.3}$ MeV
	Strange	s	95 ± 5 MeV
	Charm	c	1.275 ± 0.025 GeV
	Bottom	b	4.18 ± 0.03 GeV
	Top	t	$173.21 \pm 0.51 \pm 0.71$ GeV
Bosons	Photon	γ	0
	Gluon	g	0
	Charged weak	W	80.4 GeV
	Neutral weak	Z	91.2 GeV
	Higgs boson	H	125.7 ± 0.4 GeV

field. In relativity, mass governs how fast[13] something can travel at a given energy. Nothing travels faster than light in vacuum, precisely because photons are massless. And even though the BEH field permeates even the vacuum, photons don't interact with it and remain massless. Other particles can't travel as fast, as they are interrupted by having to interact with the medium. It is actually very similar to light in an atomic medium, such as glass. Here light travels more slowly than in vacuum, which gives glass its refractive index. At an atomic level, what happens is that the photon is constantly absorbed and re-emitted, slowing it down. On top of that, it is emitted in any random direction. From quantum mechanical effects, however, the sum of all possible paths introduces a lot of cancellations, and one direction of a light ray will be the final one. The final effect is that the light ray has refracted. In the process, the photons were moving more slowly, which can be thought of as acquiring an effective mass. Analogously, particles interacting with the BEH field acquire their masses too—the only difference being, that this medium exists **everywhere**. The masses of the fundamental particles as currently known are listed in Table 2.2.

For comparison, the proton and neutron weigh in at about 1 GeV. It is obvious that there are many fundamental particles which are heavier than these composite ones! Why the masses differ by up to five orders of magnitude between the fundamental particles is indeed a mystery in the present theoretical system.

[13]The relation between energy and velocity is given by $E^2 = m^2 + \vec{p}^2$.

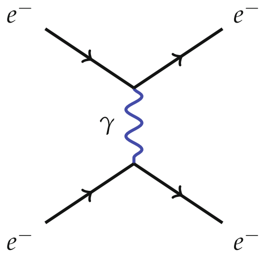

Fig. 2.1 Feynman diagram illustrating $e^-e^- \to e^-e^-$ scattering, under the exchange of a photon (γ)

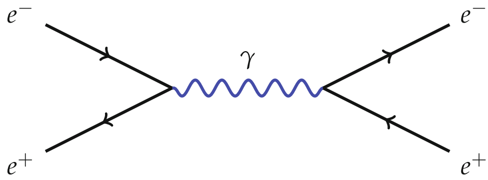

Fig. 2.2 Feynman diagram illustrating $e^+ - e^-$ annihilation into a photon (γ), and pair production back into an $e^+ - e^-$ pair

2.5 Antiparticles and Feynman Diagrams

For all of the fermions, there are also antiparticles, with the opposite sign on charges (charge conjugation). These are, for the electrically charged leptons, simply denoted with a $+$ instead of a $-$: the electron e^- has an anti-particle e^+. For neutrinos and quarks, antiparticles are denoted with a bar: u and $\bar{u}$.

The seemingly simple concept of antiparticles is still a crucial ingredient in charge conservation: only if the net charge is equal before and after the interaction, a transformation from energy in the form of one set of particles to another can occur. This is achieved in the annihilation or creation of particle-antiparticle pairs, where the net charge is 0 both before and after the interaction.

To guide intuition, there is the useful construct of a Feynman diagram. It has a profound interpretation in terms of probabilities of different processes, but let's focus on its illustrative strengths for now. In these diagrams, time flows from left to right, lines represent particles, and each vertex represents an interaction. Fermions are represented with solid straight lines, with arrows pointing right for particles and left[14] for antiparticles. Gauge bosons are represented with wavy or curly lines for electroweak bosons and gluons, respectively. Figure 2.1 is our first encounter: it illustrates how two electrons interact with (repel) each other under the exchange of a photon, the gauge boson of QED. As mentioned before, this gauge boson exchange is the model for how particles are affected by each other's presence.

Figure 2.1 shows a "space like" process. If we rotate the diagram by 90°, we get a "time like" process, as shown in Fig. 2.2.

Guided by the direction of the arrows, we realise that what is depicted in Fig. 2.2 is particle-antiparticle annihilation and pair production. The mass energy of the par-

[14]This convention goes back to considering antiparticles as particles moving backwards in time, as introduced in [5].

ticles is converted into photon energy. This is in turn converted back into a particle-antiparticle pair. As long as the available energy is large enough, a vertex like this can go in any direction (creation as well as annihilation). There is no requirement that the photon conserves flavour; it has no memory thereof as its flavour quantum number is zero (as is the combined positive and negative flavour quantum numbers of the electron and anti-electron[15]). As long as the other vertex conserves the flavour content, by for instance creating a muon-antimuon pair which taken together has zero flavour, all is well, and if the energy of the photon is large enough to create the mass of two muons, this can happen.

2.6 Hadron Case Study: The Proton

At this point, we have covered all the fundamental particles. But there is one more particle that is important to consider here: the proton, which we use for particle collisions. The proton is one example of a hadron[16]—a particle composed of quarks. Being composite, it is a suitable strong interaction case study, and we will use it to introduce some additional concepts. This is however a fairly complex topic, and we need to split it into pieces.

While quarks carry colour, hadrons as a whole are colourless. This can be accomplished in two ways: by a combination of colour-anticolour (e.g. a red-antired) as in *mesons*, or in a combination of all three (anti)colours red–green–blue, as in *baryons*. Hadrons thus consist of two or three (anti)quarks.[17] These are called valence quarks. In addition, there always occur quantum fluctuations[18] where a gluon splits into

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[15] Anti-electron: also known as *positron*.

[16] The concept of hadrons is older than the quark model, so, they must have certain unique characteristics, evident already before.

[17] Colourless combinations thereof, such as *pentaquarks*, have also been observed [6].

[18] Virtual particles can "borrow" additional energy from the vacuum, but only for a short time.

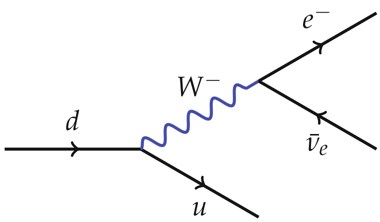


Fig. 2.3 Feynman diagram illustrating what nuclear β decay looks like at quark level, if one could resolve the W boson

a quark-antiquark pair which then annihilate back into a gluon. These fluctuation quarks are virtual, or *sea*, quarks.

Firstly, we establish that the proton is a baryon: it consists of three valence quarks, uud. This gives the proton a net electrical charge of $+1e$, and as mentioned before, no net colour charge. The other baryon making up ordinary matter, the neutron, has valence quarks udd, making it electrically neutral. The neutron is slightly heavier than the proton,[19] and an isolated neutron thus decays to a proton. At quark level, the transformation from d to u would imply weak decay involving a W boson, as illustrated in Fig. 2.3. There is in general not enough energy to create real W bosons when this happens, only virtual or *off-shell* W bosons that immediately produce a real lepton and neutrino. The comparatively long life-time of the isolated neutron, ~ 13 min, reflects all of this.

2.6.1 Parton Distribution Functions

Since the proton is a composite particle, if we accelerate the proton to carry a certain momentum, it is its constituents that carry this net momentum. The motion of constituents inside the proton is not restricted and can be both lateral and longitudinal, but the net effect has to be the overall proton momentum. We can thus stipulate

$$\sum_i \int x q_i(x) dx = 1, \tag{2.4}$$

where the x is the Bjorken x [7], which is the longitudinal momentum fraction carried by a parton, and the sum is over the quark indices i. We have already touched upon the concept of sea quarks, originating from quantum fluctuations inside the protons. By denoting proton as uud, we mean that we get a non-vanishing result

$$\int (u(x) - \bar{u}(x))\, dx = 2 \tag{2.5}$$

and

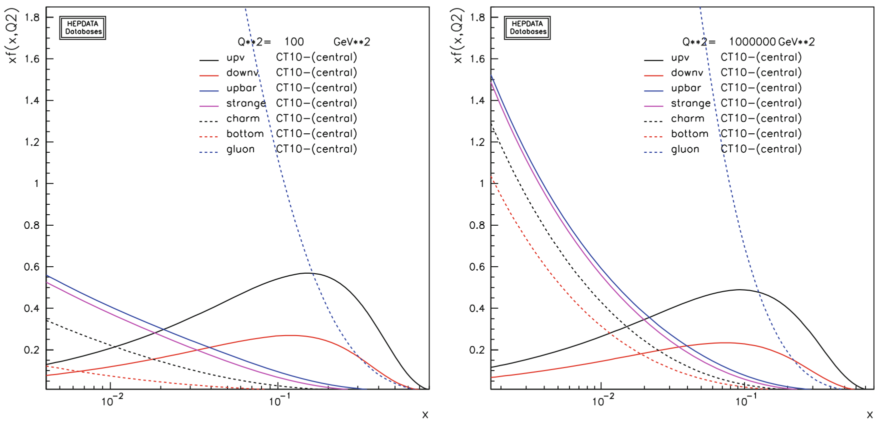
[19]More strictly speaking: $m_n > m_p + m_e + m_{\bar{\nu}_e}$.

$$\int \left(d(x) - \bar{d}(x)\right) dx = 1 \tag{2.6}$$

when we integrate over all the q and $\bar{q}$ content of the proton. The number of accessible sea quark flavours depends on the energy scale at which the proton is probed. This immediately means that the fraction of the proton momentum carried by gluons and sea quarks, respectively, depends on the energy transfer Q in the collision that probes the proton structure. In fact the fractions vary also for the valence quarks. Overall, the quarks and the gluons carry about half the momentum each. The fractions are given in the Parton Distribution Function (PDF). Two examples at different Q^2 are shown in Fig. 2.4, which shows that when the proton is probed at larger momentum transfer, the valence quarks become increasingly less dominant also at higher x. Albeit not theoretically known per se, the PDF evolution with Q^2 can be calculated from a given starting point using the Dokshitzer-Gribov-Lipatov-Altarelli-Parisi (DGLAP) equations. The starting point has to be an experimental measurement of the PDF at some Q^2. This can be data from for instance electron-proton or proton-proton collisions since the proton structure itself is universal and not dependent on the type of experiment. However, in the former case only one proton PDF is probed, making the extraction of information a little less involved.

2.6.2 *Perturbative QCD Calculations*

The logic of the Feynman diagrams, with a vertex for each interaction and mediating particles, easily lends itself to *perturbation theory*. Perturbative calculations split



(a) $Q^2 = 100$ GeV^2 (b) $Q^2 = 1$ TeV^2

Fig. 2.4 PDFs using NLO predictions including LHC data, for two values of Q^2: **a** 100 GeV^2 and **b** 1 TeV^2 [8, 9]

complicated calculations in pieces of increasing fine tuning, and start with the coarsest approximation. The method is to make an expansion[20] in increasing orders of your variable in a region where it is small, such that higher order contributions rapidly get smaller. In practice, a suitably truncated expansion is often good enough—luckily, since higher-order corrections are often not known, or computationally expensive, for a complicated expression.

Considering a process illustrated by a Feynman diagram, there is generally more than one way to draw it; there is more than one imaginable way to go from a given initial to final state, with more or less complicated steps in between. In quantum mechanics, we can't distinguish different possible histories—the intermediate steps in a process—leading up to a measured final state. But they all happen, with some probability! In a full calculation of the probability of an outcome, all of these possible paths need to be calculated, and summed correctly taking quantum mechanical interference into account. But in a Feynman diagram every vertex represents an interaction with a coupling strength, and all the vertices are multiplied to give the total probability, or *cross section*. This means that two different paths, with a different total number of vertices, are at different orders in coupling strength. If the coupling strength is small enough—which, as we shall see shortly, is the case for the small-distance, high energy transfer collisions explored in this thesis—the more complicated paths contribute increasingly little to the final result. In a perturbative calculation of the cross section of the process, we can thus truncate the expansion at some level of complexity without much loss of precision! Perturbation theory holds already for $Q > 1$ GeV, which is the proton mass and approximate confinement scale in QCD. Often the leading, or lowest, order (LO) result is a good approximation, but the next-to-leading order (NLO)corrections can be substantial.

2.6.3 Renormalisation

As mentioned, when applying the Feynman rules, all possibilities have to be integrated over, and they often come with momenta in the denominator. This gives rise to divergent (infinite) integrals, which would have to be cut off at some finite scale Λ to give finite results. Mathematically, this is not isolated to quantum field theories, even if it is a common feature of them.[21] Rather, it arises when one makes an expan-

[20]The idea is similar to the method of Taylor expansion.

[21]This discussion loosely follows Ref. [10], which gives an overview of the renormalisation idea that is worth a read!

sion of a dimensionless quantity (e.g., a probability) around a small dimensionless parameter (say, coupling strength) of a function that depends on a dimensional parameter (for instance momenta). To remain dimensionless, the calculated quantity has to depend on the dimensional parameter through the ratio with another parameter of the same dimension—a *regulator*, say, Λ. After choosing a regularisation scheme, one can redefine couplings, masses and other parameters to absorb the divergences. Typically the redefinition corresponds to a physically measured quantity (such as a coupling constant) at a given scale, which we call the renormalisation scale μ_R, with the dimensions of mass. In practice what happens is that the implicit dependence on Λ in the original expansion was removed. Only after this, we let $\Lambda \to \infty$ and get finite results. The price paid in this procedure is that the coefficients in the perturbative expansion only make sense in a given context of scale and corresponding coupling. In addition, we must abandon thinking of parameters as constant: when a quantity normalised at one scale is measured at a very different scale, the couplings and masses adjust. Also, the Λ introduced as an upper cut-off of the integrals to remove the divergence, can be thought of as the scale at which the physical theory no longer holds—a scale at which *new physics* enters.[22]

The Running of α_s

This immediately brings us to the question of the strong coupling constant. As indicated above, its value will depend on the scale at which we measure it. Experimentally, the value of α_s is given at the Z mass, and the world average is $\alpha_s(M_Z) = 0.1185(6)$ [2]. The scale dependence of α_s is controlled by the β function, which is precisely one of those parameters which do not depend on Λ:

$$\alpha^2 \frac{d\alpha_s}{d\alpha^2} = \beta(\alpha_s) = -(b_0\alpha_s^2 + b_1\alpha_s^3 + \mathcal{O}(\alpha_s^4)), \tag{2.7}$$

where $b_0 = (33 - 2n_f)/(12\pi)$, $b_1 = (153 - 19n_f)/(24\pi^2)$, and n_f is the number of accessible quark flavours. If we let $\alpha^2 = Q^2$, we can express the effective coupling strength as $\alpha_s(Q^2)$, where Q is the scale of the momentum transfer in the process at hand. Equation 2.7 shows a negative evolution of the coupling constant with the renormalisation scale μ_R. The implications are even more evident in the expression for α_s itself: from the β function, we obtain

$$\begin{aligned}\alpha_s(Q^2) = {} & \frac{4\pi}{b_0 \ln(Q^2/\Lambda_{QCD}^2)} \cdot \\ & \cdot \left[1 - \frac{2b_1}{b_0^2}\frac{\ln[\ln(Q^2/\Lambda_{QCD}^2)]}{\ln(Q^2/\Lambda_{QCD}^2)} + \mathcal{O}\left(\frac{1}{\ln^2(Q^2/\Lambda_{QCD}^2)}\right)\right]\end{aligned} \tag{2.8}$$

[22]For QED the physically meaningful upper cut-off is the scale of unification with the weak interaction.

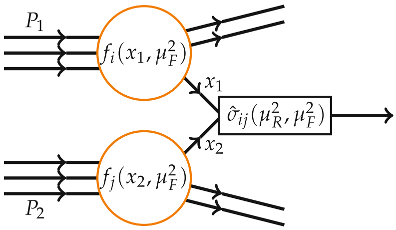


Fig. 2.5 Schematic illustration of the factorisable processes in a *pp* collision, where one parton from each proton undergoes a hard scattering

Here the reference scale $\Lambda_{QCD} \sim 200$ MeV is the confinement scale of QCD: this is the limit where α_s diverges and becomes strong. In this regime, the perturbative approach is no longer valid. In the limit $Q \to \infty, \alpha_s \to 0$. In between these regimes, α_s depends only logarithmically on Q. Furthermore, it is immediately clear that also the α_s value will depend on the order to which the perturbative expansion is carried out.

Confinement and Asymptotic Freedom

A possible way to think of a physical cause of the running coupling constant is in terms of (anti-)screening. Consider an electron. Just like a gluon fluctuates in and out of sea quark pairs, an electron constantly emits and reabsorbs field quanta, most likely photons. This can in turn create virtual loops of electron/positron pairs, which screen the charge seen farther from the electron. The net effect is a smaller effective charge of the electron, making the field around it weaker. Similarly, gluons are constantly emitted from and reabsorbed by the quark. These can in turn create virtual gluon loops, which enhance the field strength at a distance, but smear the quark colour charge as we look closely. So, the strong interaction coupling "constant" depends on the distance, or equivalently energy,[23] at which it is probed. At smaller distances (higher energies) α_s is smaller. In fact, at higher energies, more pair production becomes possible—this is one way of seeing why the classical (or leading order) approach breaks down: as we need to consider more possible paths, we need to introduce renormalisation.

The small coupling constant at high energies is called asymptotic freedom: at small distances, well inside the hadron, partons barely interact and are very loosely bound. As two quarks are increasingly separated, the potential binding energy increases. In fact the potential between them increases linearly—much like in a classical spring or rubber band, a picture exploited in the Lund string model [11], which we will summarise shortly. This theoretically requires a non-Abelian term, causing self-interactions.[24] Confinement means, that one can never observe a free quark.

[23]In the *natural units* commonly used in particle physics, where the speed of light in vacuum $c = 1$, distance has dimensions of 1/(energy).

[24]The electroweak theory is also non-Abelian, and W and Z bosons are self-interacting. Photons are not.

2.6.4 Factorisation Theorem

We concluded that we can use perturbative calculations for the high-energy processes that we are generally interested in. We have also seen, that the effective energy at which we are probing the proton, and as a result the rate of the process, depends on the PDFs. These are however not possible to calculate perturbatively, which mathematically manifests itself as divergent integrals. But luckily, the two regimes are independent—they are *factorisable*. This means that we can rely on the calculation of the DGLAP evolution for the non-perturbative PDF part, and do perturbative calculations of the hard scatter part, without loss of generality. Technically this introduces a *factorisation scale* μ_F, with $1\ \text{GeV}^2 \leq {\mu_F}^2 < Q^2$. For the regime below the factorisation scale, we use the non-perturbative proton quark distribution. The hard-scatter cross section $\hat{\sigma}_{i,j}$ is governed by short-distance processes and perturbatively calculable. We can then express the cross section for a hard scatter in a hadronic collision factorised as

$$\sigma(P_1, P_2) = \sum_{i,j} \int dx_1 dx_2 f_i(x_1, {\mu_F}^2) f_j(x_2, {\mu_F}^2) \hat{\sigma}_{i,j}(\alpha^2, {\mu_F}^2), \tag{2.9}$$

where the $P_{1,2}$ denote the incoming hadron momenta and the participating partons carry $p_1 = x_1 P_1$, $p_2 = x_2 P_2$. The $f_{i,j}(x, {\mu_F}^2)$ are the PDFs at some given Bjorken x, as given at the factorisation scale. This factorisation is schematically illustrated in Fig. 2.5.

2.6.5 Hadronisation

Since only colourless particles can travel macroscopic[25] distances, an outgoing parton from a hard scatter has to *hadronise*. This is a non-perturbative process, occurring at lower energy and correspondingly larger distances than the hard scatter, where α_s is large.

In the Lund string model, the force between two partons is pictured as a string. It has the properties of a classical string in the sense that the field contains a constant

[25] Macroscopic—or even outside the proton radius.

amount of field energy[26] per unit length, meaning that the potential increases linearly when the string is stretched [11]. If two quarks are pulled apart, in for instance a high energy collision, the binding energy becomes so large that it is energetically "cheaper" to create a real quark-antiquark pair between them, which breaks the string without resulting in free quarks (but in new strings between quarks and anti-quarks). This process is repeated as long as there is sufficient energy. The end result is a collimated hadron shower, called a *jet*, in the direction of the original quark. This jet essentially carries the energy, momentum and other properties of the original quark. Note that since hadronisation happens at longer time scales than the hard scatter process, it can't affect the partonic cross section of a process, or violate conservation laws. Measuring the jet properties is thus the way to access the properties of the original quark, even if it can't be isolated and measured itself. It is also a good way to measure their interactions.

2.6.6 Underlying Event

The remaining piece of our proton case study, is the remnants of the proton itself after a hard scatter involving one of its partons. In a violent high-energy collision, an outgoing parton produces jets due to confinement, as we have seen. Similarly, the proton remnants (illustrated in Fig. 2.5) acquire colour in the collision, and will undergo similar hadronisation. The remnants, however, often travel along the direction of the incident proton, and predominantly produce soft and diffuse radiation as measured in the transverse direction to the beam.

2.7 Monte Carlo Generators

In order to discern deviations from the expected SM behaviour in the processes studied, we need to make predictions of the SM. Our theoretical framework allows for perturbative calculations to finite orders, and non-perturbative processes such as hadronisation will remain. Using a Monte Carlo (MC) event generator, we can obtain a (pseudo-)random representation of the possible outcomes in for instance a proton collision, mimicking the stochastic processes by sampling a probability distribution. Complete generators will model both the hard-scatter process and parton showers (initial and final state radiation), hadronisation, multiple interactions and underlying event, providing a list of produced particles and their four-vectors at a given stage of the process. There are also incomplete generators calculating the hard-scatter cross sections only, which in turn may provide these calculations to higher orders.

[26] The colour field lines are not radial (as in electromagnetism) but compressed in a flux tube between the partons.

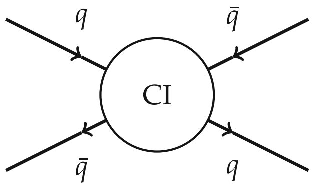


Fig. 2.6 Feynman diagram illustrating the unresolved interaction leading to $q\bar{q}$ production in a Contact Interaction (CI) approach

The underlying hypotheses for the non-perturbative processes giving these distributions can vary: the widely used complete MC generator Pythia [12] uses the Lund string model. This is the main generator used for the work described in this thesis.

2.8 Theories Beyond the Standard Model

There are numerous proposed extensions of the SM, intended to answer one or more of the outstanding questions posed by observations that seemingly have no fundamental explanation in the existing theoretical framework. Particle masses are, as I may have hinted before, a free parameter in the SM which still seems to be of some profound importance, especially if we want to unify all the known forces of nature. There are also numerous independent observations of phenomena that tell us that only about 5% of the total energy content in the universe is matter as we know it, and as all theories used in any field of science describe it. There is evidence that there is about five times as much *Dark Matter* as normal matter; the rest of the energy content in the universe is considered to be *Dark Energy* [13], the general properties of which are completely unknown. Finally, there is no a priori knowledge that the particles considered fundamental right now would not in fact have constituents—the history of particle physics actually points in the other direction. One could also argue that the mass hierarchy and generational structure points to fermion compositeness. All in all, the SM seems to be an effective theory holding up very well at the scales and the precision at which we have been probing it so far, but it may eventually have to yield to a more complete description of nature.

The measurements described in this thesis would be sensitive to many of the new phenomena predicted by such proposed extended theories. The strategy relies on simple yet powerful assumptions on what we can expect from SM processes, and the primary goal is to quantify the deviations in data from the SM prediction, rather than discover a specific hypothesised new phenomenon. Here I will focus on describing the so-called *benchmark models* used in the analysis: models making distinct predictions of observable distributions compared to the SM. When comparing these predictions to measured data, we can often make statements about the degree of compatibility with data, given certain parameter values in the model. Thus we learn something even from not discovering anything new: we learn how we *can't* describe nature.

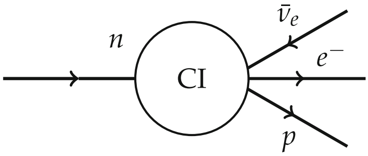


Fig. 2.7 Feynman diagram illustrating the effective field theory approach to β decay

2.8.1 Contact Interactions

One way to model fermion compositeness is to consider that at some energy scale, a new force of nature becomes manifest, as we resolve what is keeping the composite particles together. Well before that energy, however, there may be an effect on the probability and kinematic characteristics of a process, such as jet production. We can thus discern that there is something new before resolving the details of the process. This situation can be satisfactorily modelled with an *Effective Field Theory* (EFT), as depicted in Fig. 2.6. Actually, this approach is the same as in the four-point interaction of Fermi, describing nuclear β decay when there is not enough energy to resolve the W boson exchange. This is drawn in Fig. 2.7. In such a description, a scale Λ is introduced, dictating at which point we can resolve the processes hidden in the circle—the Contact Interaction (CI) [14–16]. It follows that as Λ grows, the signal strength gets weaker, if we keep the probe energy constant. The description chosen in this work is an additional effective Lagrangian:

$$\begin{aligned}\mathcal{L}_{CI}(\Lambda) = \frac{g^2}{2\Lambda^2}[&\eta_{LL}\left(\bar{q}_{iL}\gamma^{\mu}q_{iL}\right)\left(\bar{q}_{jL}\gamma_{\mu}q_{jL}\right)\\ &+\eta_{RL}\left(\bar{q}_{iR}\gamma^{\mu}q_{iR}\right)\left(\bar{q}_{jL}\gamma_{\mu}q_{jL}\right)\\ &+\eta_{RR}\left(\bar{q}_{iR}\gamma^{\mu}q_{iR}\right)\left(\bar{q}_{jR}\gamma_{\mu}q_{jR}\right)],\end{aligned} \tag{2.10}$$

where $i\,(j)$ is a flavour index, g denotes the strong coupling strength, and $\eta = 0, \pm 1$ represents the sign of the interference between CI and two-quark initial and final states of QCD: $+$ for destructive and—for constructive interference.[27] The CI is characterised by the compositeness scale Λ and its mode of interference with the QCD $q\bar{q} \to q\bar{q}$ process, where constructive interference is overall expected to lead to an enhanced signal, while for destructive interference, the effects of signal and interference compete. The CI modelling leads to non-resonant enhancement (or suppression) of jet production.

[27] Sign convention; confusing but true.

2.8.2 *Quantum Black Holes*

In a scenario where gravity propagates in more dimensions than the other fundamental forces, it would be diluted, causing it to appear much weaker than the others [17, 18]. This mechanism thus provides an explanation to the experimental observation that gravity is weaker than the other forces. The full set of space-time dimensions is commonly referred to as the *bulk*, while particles interacting under the SM[28] live on the *brane*, a 4D hypersurface in the $4 + n$ dimensional scenario. The number n of extra dimensions vary between realisations; typically $n = 1$ in a Randall–Sundrum (RS) scenario [17] and $n = 6$ in an Arkani-Hamed–Dimopoulos–Dvali (ADD) scenario [18]. This would in both cases lower the fundamental scale of gravity, M_D,[29] from the *Planck scale* $M_{Pl} \sim 10^{18}$ GeV to the vicinity of the electroweak scale $m_{EW} \sim 100$ GeV, which is clearly accessible at the LHC (see Chap. 3). This idea elegantly solves the so-called *hierarchy problem* in the SM, which is the question why these two seemingly fundamental scales are so widely separated, and it does so without introducing any new symmetries or interactions but by instead changing the space–time metric.

It does however introduce the possibility that microscopic or Quantum Black Holes (QBHs) are produced at the LHC. A TeV scale black hole created in a collision would decay to bulk and brane particles, giving an experimental possibility to detect it. If the black hole mass is larger than M_D, the black hole will thermalise and decay to high-multiplicity final states; however, there are many reasons to suspect that this is not the first mode of discovery, but rather 2-body final states are, as suggested in Ref. [19] and briefly summarised here.

Firstly, since they have not been discovered yet, it is unlikely that the energy threshold needed has been surpassed in previous experiments. Secondly, there is large suppression of Bjorken x through the PDFs, and energy loss from the initial parton-parton system, pushing the available black hole masses down for a given available centre-of-mass energy. In a regime below the production threshold energy, strong gravitational effects enhance the 2-body final state cross section through exchange of a mediating particle produced in strong gravity, even if the final state is not a black hole. Finally, even for cases with larger multiplicities, it may be seen as contrived to assume a complexity in which not also 2-body final states would be enhanced. Even so, a multi-body final state would still contribute to an analysis of 2-body final states which doesn't impose an upper limit on the number of final state objects.

[28] Note that since they don't interact in the SM, right-handed neutrinos are here not constrained to stay on the brane!

[29] The naming conventions and parameter choices vary between models. Here we choose the ADD representation.

2.8.3 Dark Matter

There is very little known about the properties of Dark Matter (DM). It interacts gravitationally, and makes up about 1/4 of the energy content in the universe—a factor 5 more than the normal matter (at least partly) described by the SM. DM particles remain to be detected.

A common approach in collider searches for DM is to assume that the DM particles produced in a collision escape detection. However, for them to be produced in the first place, there has to be a production mechanism involving coupling to partons, leaving a non-zero probability also of jet production. The production mechanism is often modelled in an EFT approach, where the scale of the phenomena is too high to be resolved using the available collision energy. This resembles the treatment of CI outlined above. Care must however be taken to avoid using an EFT approach in the regime where the available energy is larger than the scale of the new phenomenon. Here a simplified theory, assuming a mediator with some mass and a set of coupling strengths to fermions and dark matter, is a more suitable approach.

2.8.4 Excited Quarks

One consequence of quark compositeness would be the possibility of excited quark (q^*) states. Deexcitation proceeds through the emission of a gluon, making a resonant qg final state, since excitation energies would be discrete. Excited quark production and subsequent decay to quarks and gluons via gauge interactions has been used as a common benchmark for the dijet mass resonance search [20–24], and it is described in detail in Refs. [25, 26]. It is used in this thesis as a representative model for resonant dijet production.

References

1. G. Altarelli, The standard model of particle physics. CERN-PHTH-2005-206 (2005)
2. K.A. Olive et al., Particle data group. Chin. Phys. C **38**, 090001 (2014)
3. M. Niering et al., Measurement of the hydrogen 1S–2S transition frequency by phase coherent comparison with a microwave cesium fountain clock. Phys. Rev. Lett. **84**(24), 5496–5499 (2000)
4. C.S. Wu et al., Experimental test of parity conservation in beta decay. Phys. Rev. **105**(4), 1413–1415 (1957)
5. E.C.G. Stueckelberg, Remarks on the creation of pairs of particles in the theory of relativity. Helv. Phys. Acta **14**, 588–594 (1941)
6. LHCb Collaboration, Observation of J/ψp resonances consistent with pentaquark states in $\Lambda_b^0 \rightarrow J/\psi K^- p$ decays. Phys. Rev. Lett. **115**, 072001 (2015)
7. J.D. Bjorken, Asymptotic sum rules at infinite momentum. Phys. Rev. **179**, 1547–1553 (1969)
8. H.L. Lai et al., New parton distributions for collider physics. Phys. Rev. D **82**, 074024 (2010)
9. The Durham HepData Project. PDF Plotter (2016), http://hepdata.cedar.ac.uk/pdf/pdf3.html

10. B. Delamotte, A hint of renormalization. Am. J. Phys. **72**, 170–184 (2004)
11. B. Andersson et al., Parton fragmentation and string dynamics. Phys. Rep. **97**(2–3), 31–145 (1983). ISSN: 0370-1573
12. T. Sjostrand, S. Mrenna, P.Z. Skands, A brief introduction to PYTHIA 8.1. Comput. Phys. Commun. **178**, 852 (2008)
13. Planck Collaboration, Planck 2013 results. I. Overview of products and scientific results. Astron. Astrophys. **571**, A1 (2014)
14. E. Eichten et al., Supercollider physics. Rev. Mod. Phys. **56**, 579 (1984)
15. E. Eichten et al., Erratum: supercollider physics. Rev. Mod. Phys. **58**, 1065 (1986)
16. P. Chiappetta, M. Perrottet, Possible bounds on compositeness from inclusive one jet production in large hadron colliders. Phys. Lett. B **253**, 489 (1991)
17. L. Randall, R. Sundrum, A large mass hierarchy from a small extra dimension. Phys. Rev. Lett. **83**, 3370–3373 (1999)
18. N. Arkani-Hamed, S. Dimopoulos, G. Dvali, The hierarchy problem and new dimensions at a millimeter. Phys. Lett. B **429**(3–4), 263–272 (1998). ISSN: 0370-2693
19. P. Meade, L. Randall, Black holes and quantum gravity at the LHC. JHEP **05**, 003 (2008)
20. ATLAS Collaboration, Search for new particles in two-jet final states in 7 TeV proton-proton collisions with the ATLAS detector at the LHC. Phys. Rev. Lett. **105**, 161801 (2010)
21. ATLAS Collaboration, Search for quark contact interactions in dijet angular distributions in pp collisions at $\sqrt{s} = 7$ TeV measured with the ATLAS detector. Phys. Lett. B **694**, 327 (2011)
22. ATLAS Collaboration, Search for new physics in dijet mass and angular distributions in pp collisions at $\sqrt{s} = 7$ TeV measured with the ATLAS detector. New J. Phys. **13**, 053044 (2011)
23. ATLAS Collaboration, Search for new physics in the dijet mass distribution using 1 fb^{-1} of pp collision data at $\sqrt{s} = 7$ TeV collected by the ATLAS detector. Phys. Lett. B **708**, 37–54 (2012)
24. ATLAS Collaboration, Search for new phenomena in the dijet mass distribution using pp collision data at $\sqrt{s} = 8$ TeV with the ATLAS detector. Phys. Rev. D **91**, 052007 (2015)
25. U. Baur, I. Hinchliffe, D. Zeppenfeld, Excited quark production at hadron colliders. Int. J. Mod. Phys. A **2**, 1285 (1987)
26. U. Baur, M. Spira, P.M. Zerwas, Excited quark and lepton production at hadron colliders. Phys. Rev. D **42**, 815–825 (1990)

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