

Study of Higgs bosons in the WW final state and development of a fast calorimeter simulation for the ATLAS experiment

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Contents

Introduction	1
1 The Standard Model of elementary particle physics	3
1.1 Introduction to free fields and gauge interactions	3
1.2 Particle content of the Standard Model	5
1.3 Gauge groups and their interactions	6
1.3.1 The subgroup $SU(2)_L \otimes U(1)_Y$ of electroweak interactions	7
1.3.2 The group $SU(3)_C$ of strong interaction	8
1.4 Spontaneous symmetry breaking and the Higgs mechanism	9
1.4.1 The mass of the W and Z boson	11
1.4.2 The mass of the Higgs boson	11
1.4.3 The mass of fermions	12
2 Phenomenology at proton proton colliders	14
2.1 Partonic cross section	14
2.2 Hadronic cross section	14
2.3 Higher order corrections	16
2.4 Luminosity	18
2.5 Event generation for high energy hadron colliders	19
2.5.1 Hard scattering	21
2.5.2 Particle decays	21
2.5.3 Parton showers	22
2.5.4 Hadronisation	22
2.5.5 Underlying event and multiple interactions	22
2.5.6 Pile-up events	23
2.5.7 Summary	23
3 Higgs boson physics at the LHC	25
3.1 Standard Model Higgs boson production at the LHC	25
3.2 Decay modes of the Higgs boson	25
3.3 Higgs boson discovery potential	26
4 Determination of Higgs boson properties	31
4.1 Measuring the Higgs boson mass	31
4.2 Determining spin and CP quantum numbers	32

4.3	Determination of the Higgs boson coupling parameters	32
4.3.1	Limitations of the LHC	32
4.3.2	Measurement of ratios of coupling parameters	33
4.3.3	Determination of absolute coupling parameters	37
5	The Large Hadron Collider and the detector system of the ATLAS experiment	42
5.1	The Large Hadron Collider	43
5.2	The ATLAS experiment	45
5.2.1	Inner detector	46
5.2.2	Calorimeter system	50
5.2.3	Muon spectrometer	56
5.2.4	Trigger system	57
6	Development of a fast calorimeter simulation for the ATLAS detector	61
6.1	Introduction	62
6.1.1	The ATLAS full simulation	62
6.1.2	The ATLAS fast simulation Atlfast	63
6.2	Improving the fast simulation	63
6.2.1	Simplifications for FastCaloSim	64
6.2.2	Input to the parametrisation	67
6.3	Energy parametrisation	68
6.3.1	Optimal calibration	68
6.3.2	Reconstructed shower depth	69
6.3.3	Correlated energy fractions	74
6.3.4	Summary of simulation steps for the energy parametrisation	79
6.3.5	Consistency test of the energy simulation	80
6.4	Shape parametrisation	82
6.4.1	Input samples for the shape parametrisation	82
6.4.2	Parameters of the shape parametrisation	83
6.4.3	Fitting of the shower shape	89
6.4.4	Summary of the simulation steps for the shape parametrisation	92
6.5	Implementation in Athena	97
6.5.1	Important c++ class names	99
6.5.2	Execution steps of FastCaloSim	99
6.5.3	Python configuration	100
6.6	Possible improvements to FastCaloSim	100
6.6.1	Overall simulation	100
6.6.2	Energy parametrisation and simulation	101
6.6.3	Shower shape parametrisation and simulation	103

7	Validation of the fast calorimeter simulation	104
7.1	Validation against the input samples used for the parametrisation	104
7.1.1	Photon response and shape variables	105
7.1.2	Charged pion response and shape variables	108
7.1.3	Validation of jet response and efficiency	113
7.1.4	Summary	117
7.2	Performance of the fast simulation Atlfast-II	119
7.2.1	Performance of electrons	119
7.2.2	Performance of muons	122
7.2.3	Performance of jets	125
7.2.4	Performance of the missing transverse momentum	129
7.3	Summary	129
8	Study of the inclusive Higgs boson production in the $H \rightarrow WW$ final state	130
8.1	Simple cut based analysis	131
8.1.1	Selection of the leptonic final state	132
8.1.2	Event selection to suppress reducible backgrounds	133
8.1.3	Kinematic suppression of the irreducible WW background	138
8.1.4	Summary of the cut based analysis	144
8.2	Signal significance and systematic uncertainties	144
8.2.1	The influence of systematic uncertainties	146
8.3	Systematic uncertainties on the shape of the $q\bar{q} \rightarrow WW$ background	148
8.3.1	Uncertainty on $\Delta\phi_{ll}$ from the missing transverse momentum	148
8.3.2	Generator uncertainties on the shape of the $\Delta\phi_{ll}$ and m_T distribution	151
8.4	Systematic uncertainties from top backgrounds	151
8.5	Systematic uncertainties from the $gg \rightarrow WW$ background	155
8.6	The challenge of the background normalisation	158
9	Measurement of the shape and rate of the dominant dilepton backgrounds	161
9.1	Separation into exclusively longitudinal and exclusively transverse variables	161
9.1.1	The rapidity of the WW system	162
9.1.2	The rapidity difference of the W bosons	163
9.1.3	Observable longitudinal variables	165
9.1.4	Correlation between the observable longitudinal and transverse variables	168
9.2	Using uncorrelated longitudinal variables to measure transverse distributions	172
9.2.1	Measuring the distribution of a signal process	172
9.2.2	Measuring the unbiased distribution for two process classes	173

9.2.3	Measurement of all distributions and normalisation factors	173
9.3	Measuring the $q\bar{q} \rightarrow WW$ background	174
9.4	Measuring the backgrounds to the $gg \rightarrow H \rightarrow WW$ signal process	177
9.4.1	Measuring the background extrapolation ratio	179
9.4.2	Impact of the $q\bar{q} \rightarrow Z^{(*)} \rightarrow WW$ amplitude on the $q\bar{q} \rightarrow WW$ background	183
9.5	Combining both background normalisation methods	185
9.5.1	Optimising statistical versus systematic uncertainties	186
10	Analysis results based on $\sqrt{s}=10$ TeV simulations	187
10.1	Event selection cuts	187
10.2	Expected signal and background rate	188
10.2.1	Additional dilepton backgrounds	188
10.2.2	The W +jet background	190
10.2.3	The influence of detector effects on the measurement of the transverse variables	191
10.2.4	The transverse mass distribution for the signal and background processes	193
10.2.5	The accumulation of statistical uncertainties	203
10.3	Systematic background uncertainties	205
10.3.1	Systematic uncertainties on the shape of the $q\bar{q} \rightarrow WW$ background	205
10.3.2	Systematic uncertainties due to variations of the size of the top and $gg \rightarrow WW$ backgrounds	206
10.3.3	Total systematic uncertainty	209
10.3.4	Further reduction of systematic uncertainties	209
10.4	Discovery potential	209
10.5	Increasing the discovery potential	214
	Conclusion	217
A	Bijection transformation of histograms into Gaussian distributions	I
B	Simulation and properties of the signal and background processes	III
B.1	The Higgs boson signal process $gg \rightarrow H \rightarrow WW$	III
B.1.1	Treatment of spin correlations in the MC@NLO Higgs signal process	IV
B.2	Non-resonant dilepton backgrounds	IV
B.2.1	$q\bar{q} \rightarrow WW$	IV
B.2.2	$q\bar{q} \rightarrow WZ$	VI
B.2.3	$gg \rightarrow WW$	VI
B.2.4	$t\bar{t}$ and Wt	VI
B.3	Resonant dilepton backgrounds	X

B.3.1	Drell Yan $\gamma^*/Z \rightarrow ee/\mu\mu$ and $Z \rightarrow ee/\mu\mu$	X
B.3.2	$Z \rightarrow \tau\tau$	XI
B.4	Lepton reconstruction and the rejection of the W +jets background	XI
B.4.1	Electron reconstruction	XII
B.4.2	Muon reconstruction	XIII
B.4.3	Lepton isolation	XIII
B.4.4	Displaced leptons	XVII
B.4.5	Lepton selection requirements	XVII
B.4.6	Kinematic distribution	XVIII
C	The definition of the transverse mass	XX
C.1	Variants of the transverse mass	XX
C.1.1	The classical transverse mass	XX
C.1.2	The transverse mass assuming a WW system at rest . . .	XX
C.1.3	The true transverse mass	XX
C.1.4	The exclusive transverse mass	XXI
C.2	Properties of the transverse mass definitions	XXI
C.3	Advantages and disadvantages of the different transverse mass definitions	XXI
D	Event distributions	XXVII
	List of Tables	XXXII
	List of Figures	XXXIII
	Bibliography	XLIII

Introduction

The Standard Model of particle physics, which combines three of the four known fundamental interactions in a comprehensive quantum field theory, has been successfully developed during the past 50 years. These three interactions are the strong, the weak and the electromagnetic force, with the latter two being unified to the electroweak force. The bosonic force fields of the Standard Model interact with the elementary matter particles, fermions, which are grouped into leptons and quarks. While the electroweak force acts upon all fermions, the strong force couples only to quarks. Fermions occur in three generations of particles that are identical apart from their mass. For the leptons these are e , μ , τ and ν_e , ν_μ , ν_τ and for the quarks u , c , t and d , s , b .

During the last 40 years, all fermions predicted by the Standard Model were eventually discovered and also all bosons, the mediators of the three forces, were observed unambiguously. Despite increasingly accurate predictions and increasingly elaborate and precise experiments, no significant deviation between the Standard Model and experiments could be established.

With the imminent start of the next generation of experiments located at the LHC¹, the Standard Model will be probed at so far unreached energies. The LHC is designed to explore the energy range beyond the electroweak unification scale, where the consistency of the Standard Model requires the existence of another so far unobserved particle, the Higgs boson. The Higgs boson plays a central role in the Standard Model because its interactions generate the masses of all other particles. In case the Higgs boson as predicted by the Standard Model does not exist, the theory requires either a modification of or an extension to the fundamental forces and particles of the Standard Model within the energy range accessible by the LHC. Hence the LHC is expected to give access to so far unobserved fundamental properties of nature, whatever they might be.

This thesis was carried out within the collaboration of the ATLAS experiment², one of the four LHC experiments, and addresses central topics in the area of Higgs boson physics and detector simulation. It is structured as follows:

The first three chapters of this thesis give an overview of the theory of the Standard Model, its application to particle phenomenology at proton colliders and Higgs boson physics at the LHC.

As a first major part of this thesis, a study was performed how well the pa-

¹**LHC:** Large **H**adron **C**ollider

²**ATLAS:** A Toroidal **L**H**C** Apparatu**S**

rameters of a Higgs boson can be determined once it is discovered at the LHC. Methods to measure in particular the absolute coupling strength to fermions and bosons have been developed. The corresponding results and strategies are presented in Chapter 4.

Signals from so far undiscovered physics processes, like Higgs boson events, are very rare and require complex collider designs, like the LHC, that produce highly energetic proton collisions at very high luminosities. Due to the high collision energy and frequency, the LHC detectors need to be equally complex to measure and record interesting collisions. An overview of the LHC and of the ATLAS detector is given in Chapter 5.

Rare events like the Higgs boson production are concealed in an overwhelmingly large background of established Standard Model processes. In many cases the isolation of a signal is only possible with guidance from Monte Carlo simulations of the physics processes and the response of the detector. Due to the complexity of the detector, a detailed computer simulation requires several minutes per event. Due to the large Standard Model background many analyses require millions of simulated Monte Carlo events. Even with today's massive computing resources, it is difficult to meet both requirements.

This problem is addressed in Chapter 6, where a fast calorimeter simulation for the ATLAS detector is developed that reduces the event simulation time by a large factor without compromising the simulation quality significantly. Chapter 7 demonstrates that this fast simulation is suitable to be used in the analysis of many physics processes.

The last three chapters of this thesis study in detail one of the most promising channels for the discovery of the Higgs boson: the decay of the Higgs boson into two W bosons with a subsequent decay into leptons.

Chapter 8 gives an introduction to the analysis and shows that the discovery potential is high, provided the Standard Model background is known precisely. However, the Higgs signature is not very clear and the standard analyses of this channel have to rely heavily on Monte Carlo predictions of the backgrounds to reach this discovery potential. Especially the gluon induced WW background, that was studied for the first time in this thesis, can introduce large systematic uncertainties.

In Chapter 9 an alternative analysis strategy is developed that uses uncorrelated variables to isolate the Higgs boson signal and the dominant Standard Model background processes without using Monte Carlo predictions.

Finally, in Chapter 10 the analysis strategy for the isolation of the Higgs boson signal is applied to Monte Carlo events simulated with the full and fast detector simulation. It is shown that a discovery of the Higgs boson in the WW final state is possible using only information that can be extracted from recorded data and without the use of Monte Carlo predictions for the backgrounds.

The appendix contains details for some of the analysis and simulation topics. Throughout this thesis, natural units of $\hbar = 1$ and $c = 1$ are used.

1 The Standard Model of elementary particle physics

The developments in theoretical physics of the last 50 years have led to a consistent combination of quantum mechanics and special relativity into the quantum field theory of the Standard Model [1] of elementary particle physics. With a few elementary assumptions this theory is able to describe the properties of particles and the interactions between them. A very important outcome of the Standard Model is the large number of observable properties and processes that can be calculated and thus comparisons to existing measurements and predictions for new experiments are possible. The success of this theory is based on the remarkable agreement that was found between increasingly elaborate theoretical predictions and experimental observations.

The following chapter gives a brief introduction into the mathematical description of particles and their interactions. The essential properties of the Standard Model will be explained and the necessity of the Higgs field and the accompanying Higgs boson will be motivated.

A far more detailed introduction to the theory can be found in Reference [2].

1.1 Introduction to free fields and gauge interactions

A fermionic field f of mass m is described by the Lagrange density

$$\mathcal{L}_f = \bar{f}(x)(i\gamma^\mu\partial_\mu - m)f(x) \quad (1.1)$$

with $\bar{f} = f^\dagger\gamma_0$ (all formulas in natural units : $\hbar = 1, c = 1$). Applying the Euler-Lagrange-formalism to $\mathcal{L}_f$ gives the well known Dirac equation for the field f .

$$(i\gamma^\mu\partial_\mu - m)f(x) = 0 \quad . \quad (1.2)$$

Equation (1.1) is the Lagrange density and Equation (1.2) the equation of motion for the free movement of fermions of type f without interaction. One way to describe interactions between particle fields is the principle of invariance of fields under local gauge transformations: the action $S = \int \mathcal{L}_f d^4x$ is required to be

changed by not more than a total derivative by the local gauge transformation and the equation of motion has to stay unchanged (no effect on observable quantities). A simple example of such a gauge invariance is the invariance under a local $U(1)$ transformation

$$U_1(x) = e^{i\alpha(x)} \quad . \quad (1.3)$$

The effect on a field f is given by

$$f(x) \rightarrow f'(x) = U_1(x) f(x) = e^{i\alpha(x)} f(x) \quad . \quad (1.4)$$

All terms of the form

$$\bar{f}(x)f(x) \rightarrow \bar{f}'(x)f'(x) = \bar{f}(x)U_1^{-1}(x)U_1(x)f(x) = \bar{f}(x)f(x) \quad (1.5)$$

stay invariant and hence fulfil the gauge principle. In contrast, the terms containing derivatives of $f(x)$ in the Lagrange density (1.1) do not transform invariant:

$$\begin{aligned} \bar{f}(x)\partial_\mu f(x) &\rightarrow \bar{f}'(x)\partial_\mu f'(x) \\ &= [\bar{f}(x)U_1^{-1}(x)] \partial_\mu [U_1(x)f(x)] \\ &= \bar{f}(x)\partial_\mu f(x) + i\bar{f}(x)[\partial_\mu\alpha(x)]f(x) \quad . \end{aligned} \quad (1.6)$$

Therefore, the demanded gauge invariance of the Lagrange density is not fulfilled. To achieve this nevertheless, the usual derivative ∂_μ in Equation (1.1) is replaced by the covariant derivative $\mathcal{D}_\mu$:

$$\partial_\mu \rightarrow \mathcal{D}_\mu = \partial_\mu + igA_\mu(x) \quad . \quad (1.7)$$

The following transformation rule under the $U(1)$ transformation described by Equation (1.3) is defined for the field A_μ which is contained in the covariant derivative:

$$A_\mu(x) \rightarrow A'_\mu(x) = A_\mu(x) - \frac{1}{g}\partial_\mu\alpha(x) \quad . \quad (1.8)$$

The invariance of terms of the form $\bar{f}(x)\mathcal{D}_\mu f(x)$ then follows from Equations (1.7) and (1.8):

$$\begin{aligned} \bar{f}(x)\mathcal{D}_\mu f(x) &\rightarrow \bar{f}'(x)\mathcal{D}'_\mu f'(x) \\ &= [\bar{f}(x)U_1^{-1}(x)] \left[\partial_\mu + ig \left(A_\mu(x) - \frac{1}{g}\partial_\mu\alpha(x) \right) \right] [U_1(x)f(x)] \\ &= \bar{f}(x)\partial_\mu f(x) + ig\bar{f}(x)A_\mu(x)f(x) + \\ &\quad i\bar{f}(x)[\partial_\mu\alpha(x)]f(x) - i\bar{f}(x)[\partial_\mu\alpha(x)]f(x) \\ &= \bar{f}(x)\mathcal{D}_\mu f(x) \quad . \end{aligned} \quad (1.9)$$

Therefore, the redefined Lagrange density

$$\mathcal{L}_{f'} \equiv \mathcal{L}_f = \bar{f}(x)(i\gamma^\mu \mathcal{D}_\mu - m)f(x) = \bar{f}(x)(i\gamma^\mu \partial_\mu - g\gamma^\mu A_\mu(x) - m)f(x) \quad (1.10)$$

is invariant under local $U(1)$ transformations.

The field A_μ , that emerged from the demand of gauge invariance, generates an extra term

$$-g\bar{f}(x)\gamma^\mu A_\mu(x)f(x) \quad (1.11)$$

in Equation (1.10). This term is the well known QED interaction term between a fermion f with charge g and an external electromagnetic potential $A_\mu(x)$. The demand of local gauge invariance gives a well defined scheme to generate interaction terms between forces and particle fields.

Instead of being an external potential, $A_\mu(x)$ can be turned into a free photon field by adding the appropriate free Lagrangian

$$\mathcal{L}_A = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} \quad , \quad (1.12)$$

with $F_{\mu\nu} \equiv \partial_\mu A_\nu - \partial_\nu A_\mu$, to the overall Lagrange density. $A_\mu(x)$ itself is then treated as free field in the Lagrange formalism.

1.2 Particle content of the Standard Model

Table 1.1 shows an overview of the fundamental particles of the Standard Model and their most important quantum numbers.

All fermions (quarks and leptons) appear in three families, that differ only in their proper masses, with otherwise identical quantum numbers. The index L , respectively R refers to the left or right handed helicity state (weak eigenstate) of the fermion, given by:

$$f_L = \frac{1 - \gamma^5}{2} f \quad , \quad f_R = \frac{1 + \gamma^5}{2} f \quad . \quad (1.13)$$

The left handed helicity states of fermions with identical weak hypercharge Y are combined to a left handed doublet that expresses the two fermions as the two weak isospin components $T^3 = \pm\frac{1}{2}$ of an isospin state $T = \frac{1}{2}$. The relation between hypercharge Y , electric charge Q and the 3rd component T^3 of the weak isospin T is $Q = T^3 + \frac{Y}{2}$. Quarks carry an additional colour quantum number (with $C = 3$ independent eigenstates) and appear as colour triplets. All other fundamental particles are colour singlets.

Right handed neutrinos have vanishing quantum numbers for all known interactions and are consequently not directly observable. Under the assumption that neutrinos are Dirac fermions, their existence can be inferred from the observation of oscillations between neutrino families, which can only be explained by a

Leptons				s	T	T^3	Q	Y	
$L_f =$	$\begin{pmatrix} \nu_L^e \\ e_L^- \end{pmatrix}$	$\begin{pmatrix} \nu_L^\mu \\ \mu_L^- \end{pmatrix}$	$\begin{pmatrix} \nu_L^\tau \\ \tau_L^- \end{pmatrix}$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	0	-1	
				$\frac{1}{2}$	$\frac{1}{2}$	$-\frac{1}{2}$	-1	-1	
	ν_R^e	ν_R^μ	ν_R^τ	$\frac{1}{2}$	0	0	0	0	
	e_R^-	μ_R^-	τ_R^-	$\frac{1}{2}$	0	0	-1	-2	
Quarks				s	T	T^3	Q	Y	C
$Q_f =$	$\begin{pmatrix} u_L \\ d_L' \end{pmatrix}$	$\begin{pmatrix} c_L \\ s_L' \end{pmatrix}$	$\begin{pmatrix} t_L \\ b_L' \end{pmatrix}$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{2}{3}$	$\frac{1}{3}$	3
				$\frac{1}{2}$	$\frac{1}{2}$	$-\frac{1}{2}$	$-\frac{1}{3}$	$\frac{1}{3}$	3
	u_R	c_R	t_R	$\frac{1}{2}$	0	0	$\frac{2}{3}$	$\frac{4}{3}$	3
	d_R	s_R	b_R	$\frac{1}{2}$	0	0	$-\frac{1}{3}$	$-\frac{2}{3}$	3
Higgs				s	T	T^3	Q	Y	
$\phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} = \sqrt{\frac{1}{2}} \begin{pmatrix} \phi_1 + i\phi_2 \\ \phi_3 + i\phi_4 \end{pmatrix}$				0	$\frac{1}{2}$	$\frac{1}{2}$	1	1	
				0	$\frac{1}{2}$	$-\frac{1}{2}$	0	1	

Table 1.1: Table of the fundamental particles of the Standard Model. The quantum numbers are: s =spin, T =weak isospin, T^3 =3rd component of T , Q =electric charge, Y =hypercharge, C =number of QCD colour states. The states d_L' , s_L' and b_L' are linear combinations of the mass eigenstates d_L , s_L and b_L (see Section 1.3.2).

non-vanishing proper mass. This in turn makes the existence of a right helicity component necessary, since $\bar{f}f = \bar{f}_L f_R + \bar{f}_R f_L$.

Following the introduction in Section 1.1, the Lagrange density of the Standard Model contains one term of the form $\mathcal{L}_f = \bar{f}(x)(i\gamma^\mu \mathcal{D}_\mu - m)f(x)$ for every fermion in Table 1.1. The bosonic doublet ϕ of complex scalar fields (the Higgs field) will be treated in detail in Section 1.4.

1.3 Gauge groups and their interactions

The complete gauge group of the Standard Model is $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$ [1]. The interactions between fermions are derived from the invariance of the Lagrange density under transformations of this gauge group.

1.3.1 The subgroup $SU(2)_L \otimes U(1)_Y$ of electroweak interactions (Weinberg Salam model)

The subgroup $SU(2)_L \otimes U(1)_Y$ acts on matter fields with the gauge transformation

$$U(x) = e^{i \frac{Y}{2} \alpha(x) + i T \vec{\sigma} \cdot \vec{\beta}(x)} \quad , \quad (1.14)$$

where the specific values for the hypercharge Y and the weak isospin T for each fermion are taken from Table 1.1. The covariant derivative that belongs to this gauge transformation is:

$$\mathcal{D}_\mu = \partial_\mu + i g' \frac{Y}{2} B_\mu(x) + i g T \vec{\sigma} \cdot \vec{W}_\mu(x) \quad . \quad (1.15)$$

In Equations (1.14) and (1.15) $\vec{\beta}(x)$ is an arbitrary three-dimensional function, $\vec{\sigma}$ the vector of the three Pauli matrices, $B_\mu(x)$ the field associated to the group $U(1)_Y$ and $\vec{W}_\mu$ the vector of the gauge fields of the group $SU(2)_L$.

For left handed fermions the gauge transformation $U(x)$ and the covariant derivative $\mathcal{D}_\mu$ acts on the doublets $L_f/Q_f = \begin{pmatrix} \cdot \\ \cdot \end{pmatrix}_L$ with quantum number $T = \frac{1}{2}$ in Table 1.1. In order to preserve the gauge symmetry under $SU(2)_L$ transformations, all left handed fermion fields have to be grouped into such a doublet in the Lagrange density.

However, the arrangement of left handed fermions into doublets is not possible for the mass terms $-m \bar{f}(x) f(x)$ in the free fermion Lagrange density (1.1), since the mass of both isospin components would need to be identical. Furthermore, the Lagrange density is scalar, so the left handed doublet needs to be turned into a scalar quantity. Mass terms for fermions will be discussed in Section 1.4.3 in the context of the Higgs mechanism.

One should add that the gauge fields B_μ and $\vec{W}_\mu$ are not the physical fields but mix to the physical fields A (photon), Z (Z boson) and $W^\pm$ (W boson):

$$\begin{aligned} W_\mu^\pm &= \frac{1}{\sqrt{2}} (W_\mu^1 \mp i W_\mu^2) \\ A_\mu &= B_\mu \cos \theta_W + W_\mu^3 \sin \theta_W \\ Z_\mu &= -B_\mu \sin \theta_W + W_\mu^3 \cos \theta_W \quad . \end{aligned} \quad (1.16)$$

θ_W is the weak mixing angle, an experimentally determined parameter of the Standard Model, that connects the coupling strengths g and g' to the electric charge:

$$e = g \sin \theta_W = g' \cos \theta_W \quad . \quad (1.17)$$

The free Lagrange density for B_μ and $\vec{W}_\mu$ is

$$\mathcal{L}_{\text{free}(B, \vec{W})} = -\frac{1}{4} B^{\mu\nu} B_{\mu\nu} - \frac{1}{4} \sum_{a=1}^3 W^{a\mu\nu} W_{\mu\nu}^a \quad , \quad (1.18)$$

with $B_{\mu\nu} \equiv \partial_\mu B_\nu - \partial_\nu B_\mu$ and $W_{\mu\nu}^a = \partial_\mu W_\nu^a - \partial_\nu W_\mu^a - g \sum_{b,c=1}^3 f^{abc} W_\mu^b W_\nu^c$. The constants f^{abc} are the structure constants of the group $SU(2)$.

For a complete description of the physical W and Z bosons, mass terms of the form $-m_W^2 W_\mu^+ W^{-\mu} - m_Z^2 Z_\mu Z^\mu$ are still missing in the Lagrange density. However, a direct introduction of such mass terms is not possible, as they violate the gauge invariance of the Lagrange density. The most convincing way of implementing mass terms for the gauge bosons seems to be the concept of spontaneous broken symmetries as it will be described in Section 1.4.

1.3.2 The group $SU(3)_C$ of strong interaction

The group $SU(3)_C$ acts only on quarks and on itself due to the self interaction of the gluon fields. Every quark in Table 1.1 is a triplet in colour space that is affected by the $SU(3)_C$ transformation. All other particles are singlets with regard to $SU(3)_C$ transformations and stay unchanged by actions of this symmetry group. The full $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$ gauge transformation of the quark sector is:

$$U(x) = e^{i \frac{Y}{2} \alpha(x) + i T \vec{\sigma} \cdot \vec{\beta}(x) + i \sum_{a=1}^8 \gamma_a(x) \frac{\lambda_a}{2}} \quad (1.19)$$

$$\mathcal{D}_\mu = \partial_\mu + i g' \frac{Y}{2} B_\mu(x) + i g T \vec{\sigma} \cdot \vec{W}_\mu(x) + i g_s \sum_{a=1}^8 \frac{\lambda_a}{2} G_\mu^a(x) \quad (1.20)$$

The eight fields G_μ^a are the gluon fields of the QCD, the matrices λ_a are the Gell-Mann matrices. The free Lagrange density of the gluon fields is

$$\mathcal{L}_{\text{free}(G)} = -\frac{1}{4} \sum_{a=1}^8 G_a^{\mu\nu} G_{\mu\nu}^a, \quad (1.21)$$

with $G_{\mu\nu}^a \equiv \partial_\mu G_\nu^a - \partial_\nu G_\mu^a - g \sum_{b,c=1}^8 g^{abc} G_\mu^b G_\nu^c$. The constants g^{abc} are the structure constants of the group $SU(3)$.

The left handed eigenstates with respect to the group $SU(2)_L \otimes U(1)_Y$ are not the same as the eigenstates of the group $SU(3)_C$ and the mass eigenstates. The $SU(2)_L \otimes U(1)_Y$ eigenstates are related to the mass and $SU(3)_C$ eigenstates by an unitary rotation with the 3×3 CKM¹ matrix:

$$\begin{pmatrix} d'_L \\ s'_L \\ b'_L \end{pmatrix} = U_{\text{CKM}} \begin{pmatrix} d_L \\ s_L \\ b_L \end{pmatrix}. \quad (1.22)$$

¹Cabibbo-Kobayashi-Maskawa

1.4 Spontaneous symmetry breaking and the Higgs mechanism

In the previous sections, mass terms for fermions and bosons were shown to break the invariance of the Lagrange density under gauge transformation. To solve this problem, the Higgs mechanism is introduced. As already indicated in Table 1.1, a complex scalar Higgs doublet $\phi^T = (\phi_1 + i\phi_2, \phi_3 + i\phi_4)$ is postulated with the following free Lagrange density:

$$\mathcal{L}_\phi = |\mathcal{D}_\mu \phi(x)|^2 - V(\phi) \quad . \quad (1.23)$$

The expression $|\dots|^2$ is a short hand form for $(\dots)^\dagger (\dots)$. The gauge transformation for ϕ is the same as for colour neutral fermions, which is given by Equation (1.14) and the covariant derivative by Equation (1.15). The Higgs field is the only field in the Standard Model that has an external potential $V(\phi)$ in the Lagrange density. This potential $V(\phi)$ is postulated as:

$$V(\phi) = \mu^2 |\phi(x)|^2 + \lambda |\phi(x)|^4 \quad . \quad (1.24)$$

Under the special choice of parameters $\mu^2 < 0$ and $\lambda > 0$, the minimum of the external potential (1.24) is not at $|\phi|^2 = 0$ but can be found on the hypersurface of fields ϕ that fulfil

$$|\phi(x)|^2 = \phi^\dagger(x)\phi(x) = \frac{1}{2} (\phi_1^2(x) + \phi_2^2(x) + \phi_3^2(x) + \phi_4^2(x)) = -\frac{\mu^2}{2\lambda} \equiv v^2 \quad , \quad (1.25)$$

where v is defined as the vacuum expectation value of the Higgs field. An illustration of the shape of the Higgs potential can be seen in Figure 1.1.

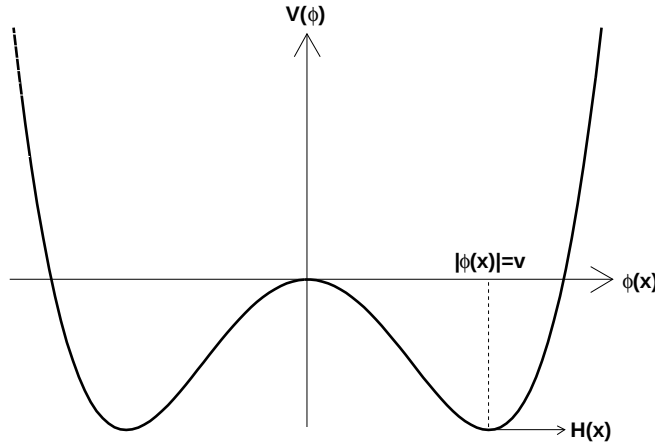


Figure 1.1: Illustration of the Higgs potential $V(\phi)$.

In the presence of such a potential with a non zero minimum, the field ϕ does not vanish in the absence of any excitation. This also implies that interactions of the Higgs field with other fields do not vanish. These non-zero ground state interactions will finally lead to the desired mass terms.

Unfortunately, it is not possible to calculate the complete solution of the Euler-Lagrange equations of motion of all fields of the Standard Model Lagrange density including interactions. As an approximation, most calculations use perturbation theory on the free non-interacting particle fields and converge only, if the necessary series expansion is performed with respect to the minimum (the ground state) of all active fields. Thus we rewrite the Higgs field as a function of fields that approach the total minimum, if all fields converge to zero. The usual minimum chosen for the Higgs field is $\phi_1(x) = \phi_2(x) = \phi_4(x) = 0$ and $\phi_3(x) = v$.

$$\phi_0(x) = \sqrt{\frac{1}{2}} \begin{pmatrix} 0 \\ v \end{pmatrix} \quad (1.26)$$

This choice is completely arbitrary and can be transformed by Equation (1.14) into any other solution of Equation (1.25). With this choice the representation of the Higgs field ϕ is

$$\phi(x) = \sqrt{\frac{1}{2}} \begin{pmatrix} \phi_1(x) + i\phi_2(x) \\ v + H(x) + i\phi_4(x) \end{pmatrix} , \quad (1.27)$$

where $\phi_3(x)$ was replaced by $v + H(x)$. This can be transformed into the usual representation of the Higgs doublet in the unitary gauge,

$$\phi(x) = \sqrt{\frac{1}{2}} \begin{pmatrix} 0 \\ v + H(x) \end{pmatrix} , \quad (1.28)$$

by local gauge transformations as given by Equation (1.14). The representation of Equation (1.28) contains only the physical Higgs field $H(x)$ and the spontaneously broken symmetry is directly visible in the form of the vacuum expectation value v . However, only gauge transformations and re-parametrisations of the fundamental fields were used to arrive at Equation (1.28) and the result is still completely equivalent to the original representation in Table 1.1. The degrees of freedom of the fields $\phi_1(x)$, $\phi_2(x)$ and $\phi_4(x)$ are completely contained in the choice of local gauge transformations that was used to obtain Equation (1.28).

Using Equation (1.15) and (1.28) in the Higgs Lagrange density (1.23) gives the Lagrange density after spontaneous symmetry breaking:

$$\begin{aligned} \mathcal{L}_H &= \frac{1}{4} (2 \partial_\mu \partial^\mu + \lambda v^2 - 2 \lambda v H(x) - \lambda H^2(x)) (v + H(x))^2 \\ &+ \frac{g^2}{8} \left(2 W_\mu^- W^{+\mu} + \frac{Z_\mu Z^\mu}{\cos^2 \theta_W} \right) (v + H(x))^2 . \end{aligned} \quad (1.29)$$

1.4.1 The mass of the W and Z boson

By extracting only those terms that contain the fields $W_\mu^\pm$ and Z_μ , but not the field $H(x)$,

$$\frac{1}{2} \left(\frac{vg}{2} \right)^2 \left(2 W_\mu^- W^{+\mu} + \frac{Z_\mu Z^\mu}{\cos^2 \theta_W} \right) , \quad (1.30)$$

from the Lagrange density (1.29), one finds the mass generating terms for the W and Z bosons, which give a mass of $m_W = \frac{vg}{2}$ and $m_Z = \frac{vg}{2 \cos \theta_W} = \frac{m_W}{\cos \theta_W}$. Since the fields of the W and Z bosons are now massive vector fields, each of them needs one longitudinal degree of freedom. These are provided by the three fields $\phi_1(x)$, $\phi_2(x)$ and $\phi_4(x)$ which were absorbed by the local gauge transformations in the transition from Equation (1.27) to Equation (1.28) [3].

Terms that contain the fields $W_\mu^\pm$ or Z_μ and also the field $H(x)$ describe the coupling of the Higgs boson to the W and Z bosons. One finds:

$$\mathcal{L}_{WW \rightarrow H(H)} = \frac{1}{2} \left(\frac{m_W}{v} \right)^2 (2v H(x) + H^2(x)) (2 W_\mu^- W^{+\mu}) \quad (1.31)$$

$$\mathcal{L}_{ZZ \rightarrow H(H)} = \frac{1}{2} \left(\frac{m_Z}{v} \right)^2 (2v H(x) + H^2(x)) (Z_\mu Z^\mu) , \quad (1.32)$$

thus the coupling strength of the Higgs boson to the W and Z bosons is proportional to m_W^2 , respectively m_Z^2 . However, this coupling is a gauge coupling and the parameters m_W , v and g are closely related to each other.

1.4.2 The mass of the Higgs boson

By separating all terms that are quadratic in the Higgs field $H(x)$ from the Lagrange density (1.29), one finds:

$$-\frac{1}{2} (2\lambda v^2) H^2(x) . \quad (1.33)$$

This is a mass term for the Higgs field itself, which describes a Higgs boson with a mass of $m_H = \sqrt{2\lambda}v$. The value of the parameter v can be determined experimentally from the known masses of the W and Z bosons to be $v \approx 246$ GeV. The value of the parameter λ is unknown and no prediction of the Higgs boson mass can be made from current measurements.

The direct searches at the LEP experiments give a lower bound on the mass of the Standard Model Higgs boson of $m_H > 114.4$ GeV [4]. Further limits on the mass of the Higgs boson can be extracted from a combined fit to the electroweak precision observables [5]: Figure 1.2(a) shows the individual bounds on the Higgs boson mass from the different observables and Figure 1.2(b) shows the $\Delta\chi^2$ distribution of the combined fit to all observables. The latter indicates

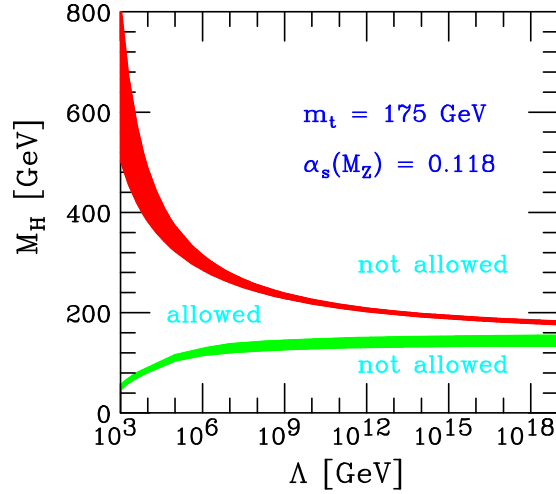


Figure 1.3: The triviality (upper) bound and the vacuum stability (lower) bound on the Higgs boson mass as a function of the new physics or cut-off scale Λ [6] for a top quark mass $m_t = (175 \pm 6)$ GeV and $\alpha_s(M_Z) = 0.118 \pm 0.002$. The allowed region lies between the bands and the coloured bands illustrate the impact of various uncertainties.

is defined. Substituting the Higgs field after symmetry breaking (1.28) gives:

$$\mathcal{L}_{m_f(\text{downtype})} = -\frac{G_f v}{\sqrt{2}} (\bar{f}_L f_R + \bar{f}_R f_L) - \frac{G_f}{\sqrt{2}} (\bar{f}_L f_R + \bar{f}_R f_L) H(x) \quad . \quad (1.35)$$

By defining the free coupling parameter $G_f = \frac{m_f \sqrt{2}}{v}$, one gets

$$\mathcal{L}_{m_f(\text{downtype})} = -m_f \bar{f}(x) f(x) - \frac{m_f}{v} \bar{f}(x) f(x) H(x) \quad , \quad (1.36)$$

using the equality $\bar{f}_L f_R + \bar{f}_R f_L = \bar{f} f$. Due to the Yukawa interaction, the Lagrange density after the symmetry breaking contains one mass term for each fermion f and one coupling term of the fermion f to the Higgs field $H(x)$, which is proportional to the fermion mass.

A similar approach is made for the upper component of a left handed doublet, only in this case the Higgs doublet ϕ is replaced by $\phi_c = -i \sigma_2 \phi^*$:

$$\mathcal{L}_{m_f(\text{uptype})} = -G_f (\bar{L}_f \phi_c f_R + \bar{f}_R \bar{\phi}_c L_f) \quad (1.37)$$

Using the same definition of G_f as above, one gets:

$$\mathcal{L}_{m_f(\text{uptype})} = -m_f \bar{f}(x) f(x) - \frac{m_f}{v} \bar{f}(x) f(x) H(x) \quad (1.38)$$

and thus the same type of mass and coupling terms to the Higgs field as for the downtype fermions in Equation (1.36).

2 Phenomenology at proton proton colliders

The previous chapter contained a very brief overview of the Standard Model and the Lagrange density for the various components of the theory. However, this alone is not sufficient to make calculations and predictions for existing or upcoming experimental measurements and a large number of perturbative as well as non-perturbative effects have to be taken into account to describe the phenomenology of particle production at proton-proton (pp) colliders. In most cases, complex Monte Carlo simulation techniques are needed to obtain reliable predictions. The following chapter gives an overview of the various elements involved in these Monte Carlo simulations.

2.1 Partonic cross section

The first step in obtaining a prediction for some particle production process is the calculation of the partonic cross section. The partonic cross section can be seen as the probability for the transition from a completely given initial state to a completely given final state. One commonly followed approach to obtain it, is to derive the invariant matrix element $\mathcal{M}$ from the interaction Lagrangian of the theory. With this matrix element, the partonic cross sections can be easily calculated from

$$d\sigma(a + b \rightarrow c) = \frac{1}{2\hat{s}} \frac{1}{(2\pi)^2} \int \frac{d^3p_c}{2E_c} \delta^4(p_a + p_b - p_c) \cdot F_{\text{color}} \cdot F_{\text{spin}} \sum_{\text{spin,color}} |\mathcal{M}|^2 \quad (2.1)$$

where the sum runs over all possible initial and final spin and colour states and F_{color} and F_{spin} are factors which stem from the averaging over the initial colour and spin states.

2.2 Hadronic cross section

In the case of hadron colliders, especially the initial state particles are often quarks or gluons, which do not exist as free particles but are contained in stable hadrons. Thus, for the calculation of measurable cross sections at the LHC,

hadronic interactions between protons have to be considered. The main interest at hadron colliders lies in the production of heavy particles (like electroweak gauge bosons or the Higgs boson), which usually occurs in interactions with high four momentum transfer Q^2 . The production cross section can be calculated in the framework of the parton model; for more details see e.g. Reference [7].

In the following, the parton a is considered as a constituent of hadron A and parton b as constituent of hadron B . To every hadronic interaction of the form

$$A + B \rightarrow c + X \quad (2.2)$$

where c could be a heavy particle (e.g. a Higgs boson) and X denotes hadronic debris left over after the interaction, the partonic subprocess

$$a + b \rightarrow c \quad (2.3)$$

is associated, whose cross section is given in Equation (2.1). To translate the partonic cross section $\sigma_{a+b \rightarrow c}$ back to the hadronic cross section $\sigma_{A+B \rightarrow c+X}$ the probability to encounter the parton a with a momentum fraction x_a within the hadron A is needed and likewise for b . The probability density $f_{a/A}(x_a, Q^2)$ is called “parton distribution function” (PDF). The products of the parton momentum fraction x and the quark or gluon PDFs, as obtained by the CTEQ group [8], are shown in Figure 2.1 for two example Q values of 2 GeV and 100 GeV.

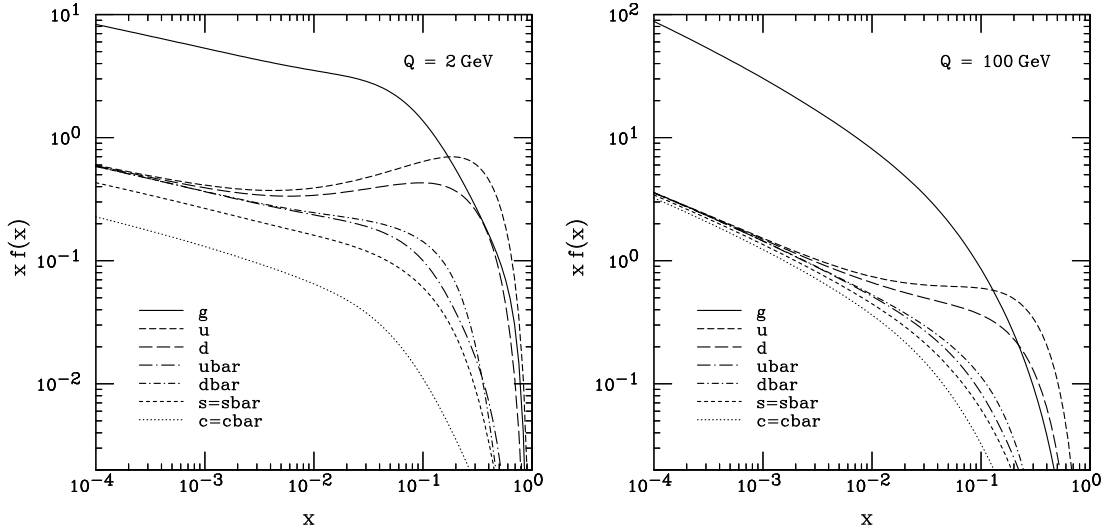


Figure 2.1: Product of the parton momentum fraction x and the quark or gluon PDFs $f(x, Q)$ as obtained by the CTEQ group [8] for Q values of 2 GeV and 100 GeV. The distributions are generated from PDF version CTEQ6M.

Finally, the hadronic cross section is calculated from the partonic cross section by a convolution with the PDFs for both initial partons a and b :

$$d\sigma(A+B \rightarrow c+X) = \sum_{a,b} \int_0^1 dx_a \int_0^1 dx_b f_{a/A}(x_a, Q^2) f_{b/B}(x_b, Q^2) d\sigma(a+b \rightarrow c) \quad (2.4)$$

where the sum extends over all possible initial partons a and b whose collisions can produce the final state c . A schematic view of this is shown in Figure 2.2.

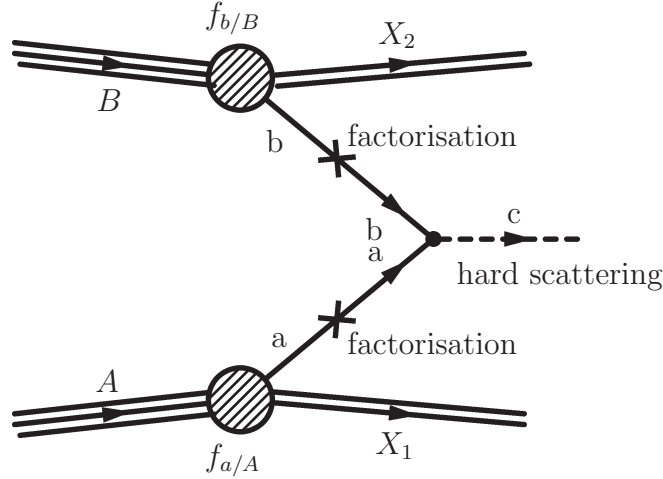


Figure 2.2: Schematic view of the factorisation of a high p_T process into a soft part absorbed into PDFs ($f_{a/A}$ and $f_{b/B}$) and a hard scattering.

The splitting into what is called “hard scattering” process with a partonic cross section and a “soft part” (low Q^2), hidden in the parton distribution functions, is called “factorisation”. Factorization is possible, since the PDFs are universal and hence independent of the hard process and can therefore be measured experimentally. The scale Q^2 , which is assumed to separate the hard process and the soft contributions, is called the “factorisation scale” μ_F . When cross sections are obtained at a fixed order in perturbation theory, the result depends on the unphysical factorisation scale. This dependency becomes weaker with increasing orders of calculation and would vanish if one could carry out the calculations to all orders of perturbation theory.

2.3 Higher order corrections

For the calculation of the partonic cross section, perturbation theory has to be used. Results obtained only at leading order (LO) suffer from large uncertainties. The variation of unphysical scales involved in the calculations, usually the factorisation scale μ_F and renormalisation scale μ_R , can give a handle on the theoretical uncertainty. If it were possible to calculate quantities to all orders,

dependencies on these unphysical scales would vanish. Therefore, it is indispensable to calculate, whenever practically possible, as many higher order corrections as possible to reduce the μ_R and μ_F dependency.

In Figure 2.3(a) the leading order contribution to the production of a Higgs boson in gluon fusion is illustrated. The higher order contributions consist of virtual corrections (see Figures 2.3(b) - 2.3(d)), propagator corrections (see Figures 2.3(e) - 2.3(f)) and emissions of real particles (see Figures 2.3(g) - 2.3(i)). Due to the large number of diagrams, not all possible next-to-leading order (NLO) Feynman graphs are shown. In a full NLO calculation, the soft and collinear singularities of the real emission cancel the singularities from the interference terms of the virtual corrections with the leading order contribution, such that the cross section receives a finite correction. This cancellation occurs in each order of perturbation theory.

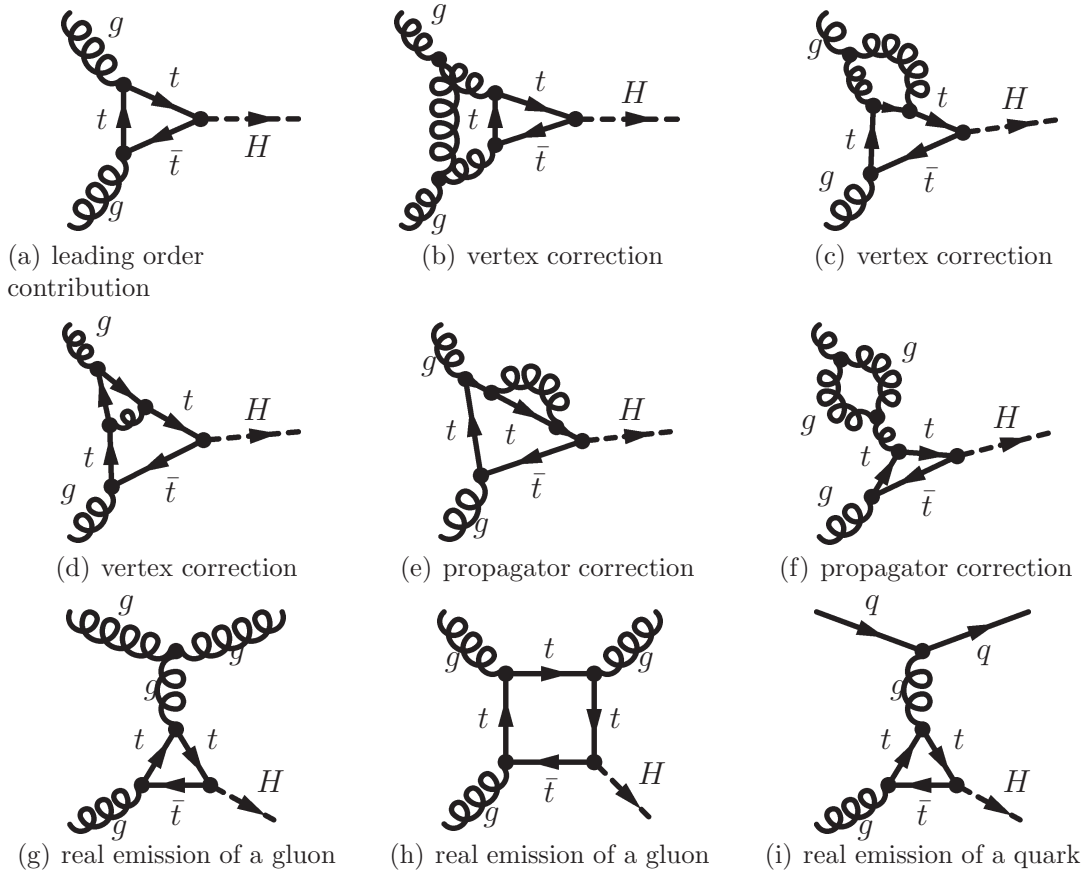


Figure 2.3: Tree level and some higher order QCD correction diagrams to the Higgs production in gluon fusion.

As can be seen from Figure 2.3, not only the inclusive cross sections, but also final states and kinematics are influenced by higher order corrections. From Figures 2.3(g) - 2.3(i) one obtains an additional quark or gluon in the final state,

which cannot be produced in a leading order calculation.

In addition to these fixed order higher order corrections, there are soft corrections that sum up large leading logarithms that appear at small transverse momenta. The current best calculation of the gluon fusion cross section, including all known effects from NNLO fixed order calculations, NNLL soft gluon logarithmic resummation and electroweak corrections (see Reference [9]) is shown in Figure 2.4 for LHC energies of 10 TeV and 14 TeV as a function of the Higgs boson mass.

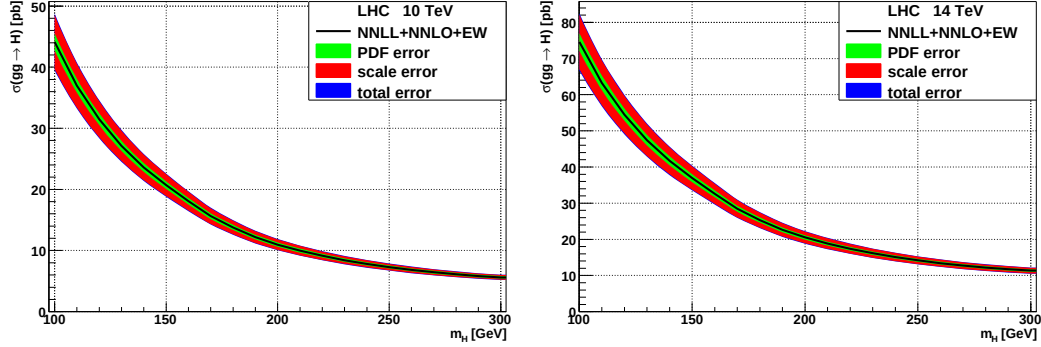


Figure 2.4: Gluon fusion cross sections at the LHC ($\mu_F = \mu_R = m_H$) with $\sqrt{s} = 10$ TeV (left) and $\sqrt{s} = 14$ TeV (right) using the MSTW2008 [10] parton densities. The cross sections and uncertainties are taken from the tables in the appendix of Reference [9].

2.4 Luminosity

Once the hadronic cross section is calculated, the expected event rate at a hadron collider can be predicted if the luminosity of the collider is known. As long as the geometry of the colliding beams is sufficiently simple, the instantaneous luminosity $\mathcal{L}$ can be calculated as

$$\mathcal{L} = f \frac{n_1 n_2}{A_{eff}} \quad (2.5)$$

where n_1, n_2 are the number of particles contained in the proton bunches of the beams 1 and 2, which are brought to collision with the frequency f in the effective transverse collision area A_{eff} . If the beam profile is perfectly Gaussian in both directions (x and y) transverse to the beam direction, the effective area is given by

$$A_{eff} = 4\pi \sigma_x \sigma_y \quad (2.6)$$

The rate of events, $R = dN/dt$ is the number of events of a certain physical process observed per time. It is given as the product of the hadronic cross section σ and the instantaneous luminosity $\mathcal{L}$

$$R = \sigma \cdot \mathcal{L} \quad . \quad (2.7)$$

Therefore the number of events N with a cross section σ is given by the product

$$N = \sigma \cdot L = \sigma \cdot \int \mathcal{L} dt \quad , \quad (2.8)$$

where L is the integrated luminosity usually given in units of “barn” (or fractions like pb, fb) where $1 \text{ barn} = 10^{28} \text{ cm}^2$.

2.5 Event generation for high energy hadron colliders

The theory outlined in the previous chapters can be used to calculate cross sections for certain processes and decay widths for particles. The final state particles of these calculations are leptons, quarks, gluons and electroweak bosons. Most of these particles are not experimentally observable, since a quark or gluon will always hadronise before observation and heavy bosons decay before detection.

The hadronisation of quarks and gluons cannot be treated with perturbation theory, since the strong coupling is of order 1 in the relevant energy regime of a few GeV and hence too big to be treated as a perturbation. Therefore, hadronisation has to be described using phenomenological models that are based on experimental input¹.

Coloured or electrically charged particles in the initial or final state might undergo radiation and emit additional particles as initial state radiation (ISR) and final state radiation (FSR).

The proton remnant, since it is not colour neutral, has to undergo additional interactions to produce colour neutral final state hadrons. This leads to an underlying event (UE). In addition, the incoming particle beams themselves can become so dense that more than one scattering event per bunch crossing is probable.

All these complex effects are taken into account in computer programs called “Monte Carlo Event Generators” which produce “events” for a desired process. These events consist of stable particles with specified creation points in space and time, four momenta and particle type. One possible event for the Higgs boson production in gluon fusion is sketched in Figure 2.5. For a comprehensive summary of all steps involved, see References [11] and [12].

¹Usually these models are derived from the current understanding of the underlying physics and contain a number of parameters, that need to be adjusted (tuned) to give the best possible agreement with experimental measurements.

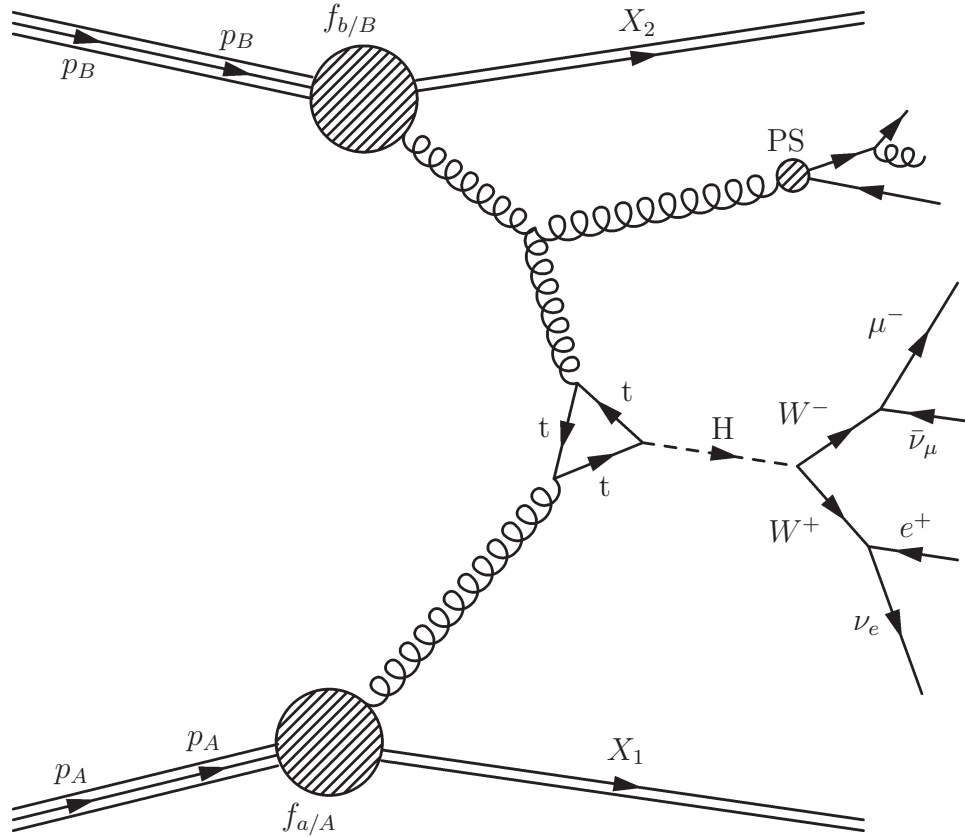


Figure 2.5: Generation of an event for a proton-proton collider. Shown is the reaction $p_A + p_B \rightarrow H + X$. Each proton emits a hard gluon. The two gluons fuse to give a Higgs boson (H) via a top-quark loop. The scalar Higgs boson decays into a pair of W bosons, which decay into $e^+ + \nu_e$ and $\mu^- + \bar{\nu}_\mu$. The upper gluon emits another gluon which fragments further in a parton shower step before the end products hadronize. Both protons also leave a beam remnant X_1 and X_2 which, again, has to hadronize.

An additionally necessary step in obtaining a prediction of what would be measured in a given experiment is the simulation of the detector. Detector simulation computer programs simulate the behaviour of a particle physics detector in response to (pseudo-) stable particles. It aims at predicting the detector signals produced by an event. An example of such a detector simulation is described in Chapter 6.

The various steps involved in the generation of an event will be explained in more detail in the following.

2.5.1 Hard scattering

The central part of an event generator program is given by the hard scattering and the convolution with the PDFs as it was described in Section 2.2. The calculations are usually performed at lowest order in perturbation theory, so that the hard scattering is given by a $2 \rightarrow n$, $n \lesssim 6$ scattering process. It should be noted that in the last few years the Monte Carlo event generation program MC@NLO [13] was developed which contains matrix elements and parton showers at NLO. Where possible and appropriate, NLO Monte Carlo event generation was used to model signal and background processes in Chapters 8–10. Some programs already contain libraries for a large number of matrix elements to perform the generation of events, but, in general, at least in leading order, automatic matrix element generators are used.

Once the total cross sections are evaluated for all allowed, contributing subprocesses, an event in a given kinematic configuration is selected probabilistically. At this stage the event still contains particles like quarks and gluons, and has a fixed number of outgoing particles, which might be unstable and decay rapidly.

This “bare” event has subsequently to be “dressed up” by the following steps.

2.5.2 Particle decays

Particle decay processes need to be simulated for unstable particles produced in the hard scattering or in later steps of the event generation in relative proportion to all relevant particle decay widths. The partial decay width can either be calculated with the knowledge of the couplings or is known from measurements in previous experiments. In an event generator program one of the available decay modes is usually selected on a probabilistic basis. If a precise kinematic description of the decay configuration is necessary, decay matrix elements and/or spin correlations need to be included in the calculation.

In the case of the Higgs boson analyses presented in Chapters 8 and 10, such spin correlations in the decay are important and necessary for all $W \rightarrow l\nu$ decays for the signal and the dominant background processes.

2.5.3 Parton showers

A realistic description of a high energy event at a hadron collider is obtained by taking QCD radiation effects into account. The use of matrix elements including more and more real emissions is impractical due to the rapidly increasing complexity and also due to singularities from soft and collinear emissions.

In the limit of soft or collinear configurations, the emission of quarks and gluons factorises from the matrix element and can be treated universally on a probabilistic basis. Several methods exist and allow for some tuning to measurements in such a way, that also non-collinear and hard emissions are produced with limited predictive power (although large uncertainties remain). This approximate inclusion of multiple emissions in an event generation is called “parton shower” (PS).

Also a combination of LO matrix elements and parton showers can be used if the overlap is treated correctly [11]. These methods promise to give an accurate description of hard emissions by using appropriate matrix elements with additional quarks and gluons in the final state and of soft emissions by using the parton shower.

A parton shower algorithm can also be applied to initial state partons, implementing a backward shower algorithm. Therefore initial state radiation is treated on a similar footing as final state radiation.

2.5.4 Hadronisation

After particle decays and parton showering have taken place, colour neutral particles have to be formed from the remaining quarks and gluons by hadronisation into mesons and baryons.

As stated before, the hadronisation is a non-perturbative process and must be described by phenomenologically driven models tuned to measurements in experiments. Several such models exist and are implemented in computer event generation programs. Since the configurations of quarks and gluons before hadronisation will depend on the type and parameters of the parton shower, the tuning of a hadronisation model will depend on the Parton Shower itself.

2.5.5 Underlying event and multiple interactions

Section 2.2 describes how a parton from each incoming proton is selected to participate in a hard scattering (also depicted in Figure 2.2). Usually this leaves a coloured beam remnant that will undergo fragmentation and hadronisation to produce stable, colour neutral particles that form the underlying event of the hard scattering. Due to the very low transverse momenta of these particles, the underlying event is produced by a non-perturbative process that can currently only be described by phenomenological models tuned to experimental measurements.

Since the fragments of the proton remnants are almost collinear with the incoming proton beams, they cannot be fully detected by any detector and the total energy and momentum of the collision is not completely measurable. Therefore, only the total transverse momentum is a well defined quantity at a hadron collider.

Also the treatment of multiple interactions is still a widely and controversially discussed field, closely coupled to the discussion of the underlying event. Since protons are composite objects, consisting of quarks and gluons, it can be assumed that each of these components has a certain probability to interact independently and give rise to additional scatterings.

Current models of the underlying event and multiple interactions are tuned to previous measurements at hadron colliders. But since the theoretical knowledge of this topic is still limited, the event activity coming from both effects at LHC energies is currently not very well predicted.

2.5.6 Pile-up events

As a consequence of the large density of protons in the LHC beams, there is a high probability that more than one proton per beam will undergo an inelastic interaction. The superposition of several inelastic proton-proton collisions is called “pile-up”.

Usually the additional interactions accompanying a hard scattering process are not hard scattering events due to the much larger cross section of low p_T inelastic interactions. These low p_T events are commonly called “minimum bias events” (MB).

Usually the effect of pile-up events is included in the event generation by overlaying the hard scattering event with a number of MB events that varies with luminosity. For the first three years of LHC running (low luminosity phase with $L = 2 \cdot 10^{33} \text{ cm}^{-2} \text{ s}^{-1}$) it is expected that on average 4.6 MB events are superimposed on each hard proton-proton collision. For the following years (high luminosity running, $L = 10^{34} \text{ cm}^{-2} \text{ s}^{-1}$) it is expected that the average rate of superimposed minimum bias events increases to 23.

Sometimes the type of pile-up described here is called in-time pile-up, to differentiate it from the effect that under certain circumstances the detector cannot distinguish events from different bunch crossings.

2.5.7 Summary

The elements described in this chapter allow for the generation of high energy physics events. Despite the relative success of the programs implementing these generation steps, it should always be kept in mind that they can only be an approximation of nature. Apart from the uncertainties involved in the individual steps, it is most probable that the separation of the various steps from the initial

protons to the final stable hadrons is less distinct in nature as it is in the current calculations and Monte Carlo simulation programs. Even if the final hadron state was correctly described in all details, any intermediate step might not be physical as it cannot be measured directly in any way².

Nevertheless, these programs are indispensable tools to analyse the data of high energy experiments, and also to develop strategies of how to extract the maximal information from data that will be recorded in the future.

²This is especially relevant when considering parton level simulations that are directly based on matrix element calculations. While results for the kinematical configuration of the hard scattering itself are reliable to some degree, all QCD aspects of the event are most likely not well described or not present at all.

3 Higgs boson physics at the LHC

In the subsequent chapter the phenomenology of Higgs boson production and decay in the Standard Model will be shortly discussed, followed by a brief overview of the current estimate of the discovery potential at the LHC. The Standard Model predicts all properties of the Higgs boson, except for the mass: it is a neutral, spin zero, CP-even particle, with gauge couplings to the W and Z bosons and Yukawa couplings to the fermions.

3.1 Standard Model Higgs boson production at the LHC

A result of the Higgs mechanism in the Standard Model as discussed in Chapters 1.4.1 and 1.4.3 is, that the coupling of the Higgs boson to fermions and gauge bosons is proportional to the mass of the particles. Thus it can be expected that Higgs bosons will be produced dominantly in processes involving the heaviest particles in the Standard Model: the W and Z boson, as well as the top quark, giving rise to the production channels shown in Figure 3.1.

The loop induced gluon fusion process is the dominant production mode at the LHC due to the huge gluon luminosity (see Figure 2.1). The Higgs boson production cross sections as function of the unknown Higgs boson mass are shown in Figure 3.2 for $\sqrt{s}=14$ TeV.

3.2 Decay modes of the Higgs boson

Similar to the Higgs boson production, also the decay of the Higgs boson is dominated by the coupling to heavy particles. Figure 3.3 shows the Higgs boson decay modes that are experimentally accessible at the LHC, Figure 3.4 shows the branching ratios of these decay modes as a function of the hypothetical Higgs boson mass.

Basically only final states with sizeable branching ratios of several percent are observable at the LHC. The one exception is the decay of $H \rightarrow \gamma\gamma$, which could be observable due to its very clear signature of two highly energetic photons showing a resonant peak at the Higgs boson mass.

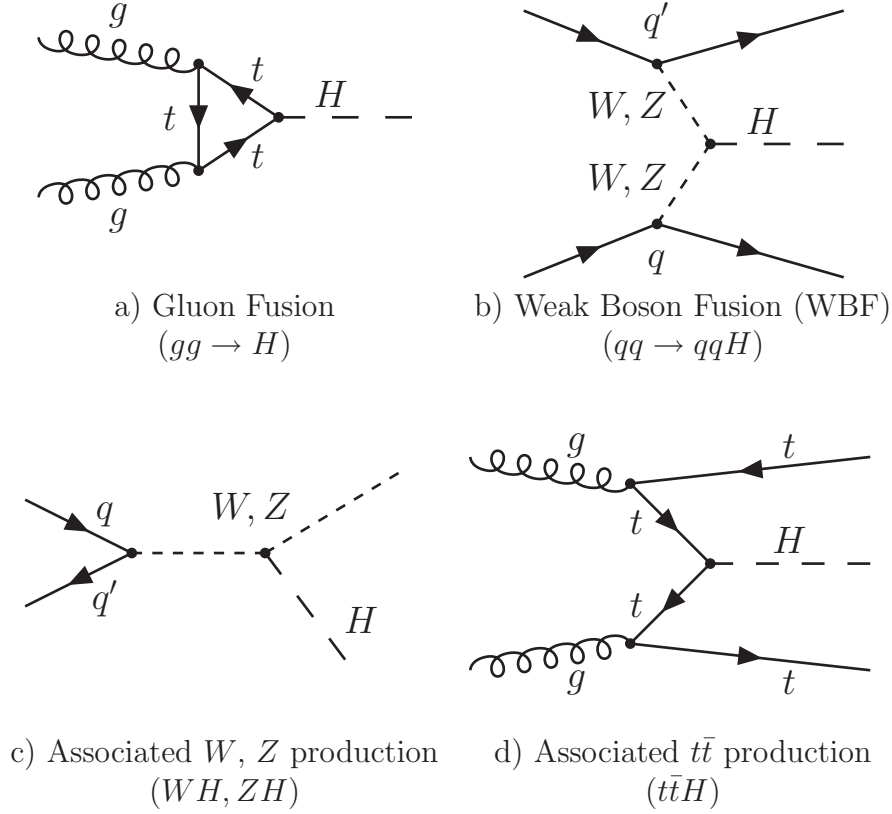


Figure 3.1: Dominant Standard Model Higgs boson production modes at the LHC. The WBF production process is sometimes also called Vector Boson Fusion (VBF).

3.3 Higgs boson discovery potential

During the last years the collaborations of both multi purpose LHC detector communities, ATLAS and CMS (see Chapter 5), have invested a huge amount of work to study all promising discovery modes and to establish the discovery potential for a Standard Model Higgs boson as a function of its hypothetical mass.

An observed signal is generally accepted as having been discovered, if the probability for the observation due to a fluctuation of the background in absence of the signal is less than $2.9 \cdot 10^{-7}$. This is the one-sided probability for a fluctuation in a Gaussian distribution with width σ to be observed at least 5σ away from the mean of the distribution. For convenience, the discovery potential of a channel is usually given in units of Gaussian σ -deviations that correspond to the probability to observe a background fluctuation of the size of the expected signal.

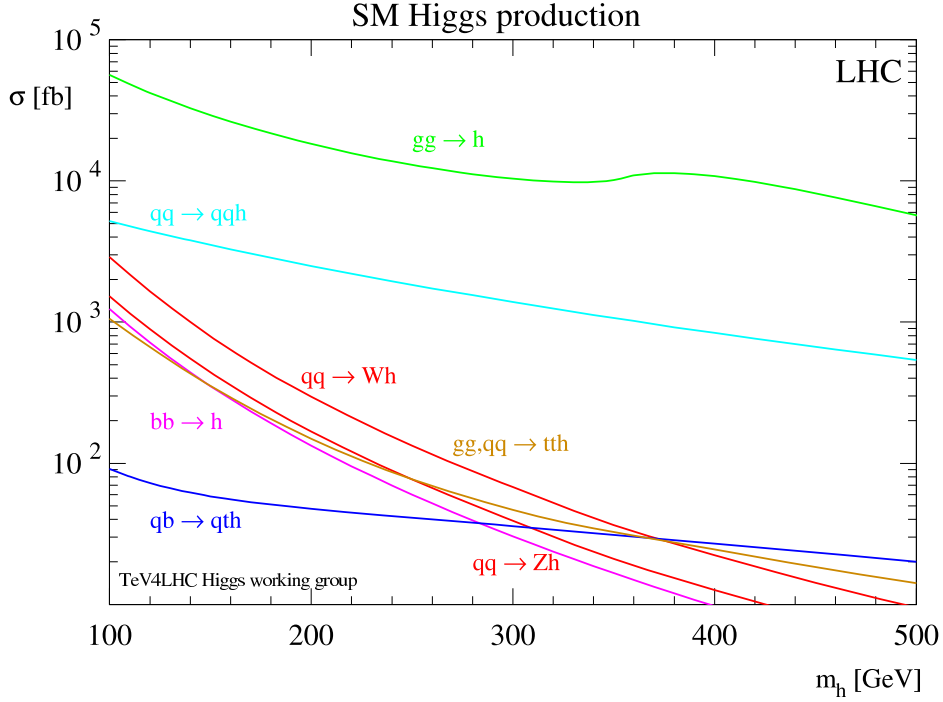


Figure 3.2: Standard Model Higgs boson production cross sections (fb) at the LHC ($\sqrt{s} = 14$ TeV) for the most relevant production mechanisms as a function of the Higgs boson mass [14]. Results for $gg \rightarrow h$, $q\bar{q} \rightarrow Vh$, $b\bar{b} \rightarrow h$ are calculated at NNLO in the QCD expansion. Weak boson fusion ($qq \rightarrow qqh$) and $t\bar{t}$ associated production are at NLO accuracy. Single-top associated production ($qb \rightarrow qth$) is at LO.

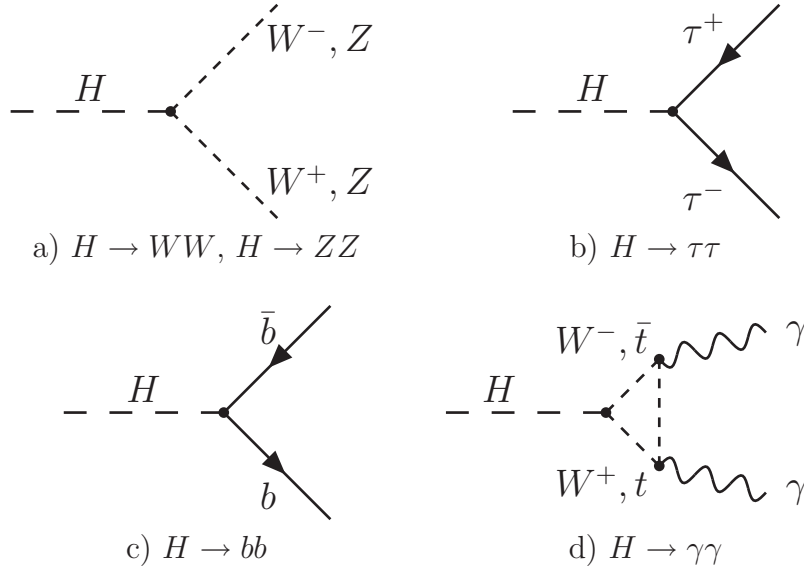


Figure 3.3: Experimentally accessible Standard Model Higgs boson decay modes at the LHC.

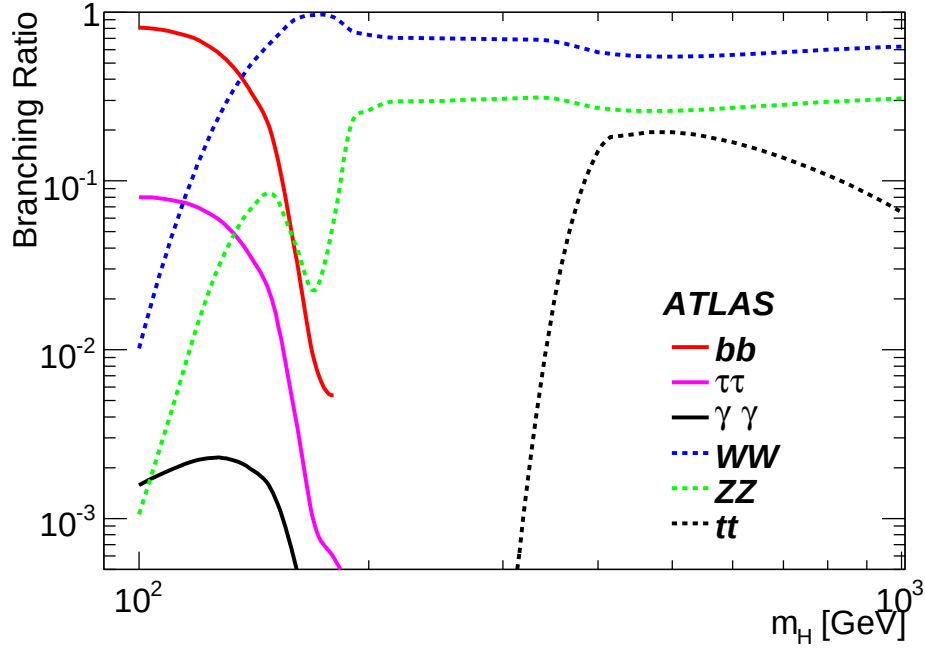


Figure 3.4: Calculated Standard Model Higgs boson branching ratios as function of the Higgs boson mass [15].

The discovery potential of the ATLAS experiment for a Standard Model Higgs boson is shown in Figure 3.5, using an integrated luminosity of 10 fb^{-1} . The discovery potential is calculated as the statistical combination of the five most significant Higgs boson channels. For the $H \rightarrow WW$ analyses only the dominant contribution with one electron and one muon in the final state is included. Using the most recent theoretical and experimental predictions, it is expected that a combined 5σ discovery potential can be reached by the ATLAS experiment in the complete mass range of approximately 125 GeV–400 GeV.

For a Higgs boson mass close to the WW threshold, a discovery using just $1\text{--}2 \text{ fb}^{-1}$ seems feasible, as can be seen in Figure 3.6. However, it should be noted that this mass range is already excluded at 95% CL by the combination of the two Tevatron experiments [16].

The current Tevatron Higgs boson sensitivity is shown in Figure 3.7. The Higgs boson exclusion limit is shown as ratio between the excluded cross sections and the Standard Model cross section and calculated from the combination of dedicated searches of both Tevatron experiments in different channels. In the narrow mass range of $160 \text{ GeV} < m_H < 170 \text{ GeV}$ the observed cross section limit is below the Standard Model expectation and the existence of a Standard Model Higgs boson can hence be excluded with a 95% confidence limit in this range.

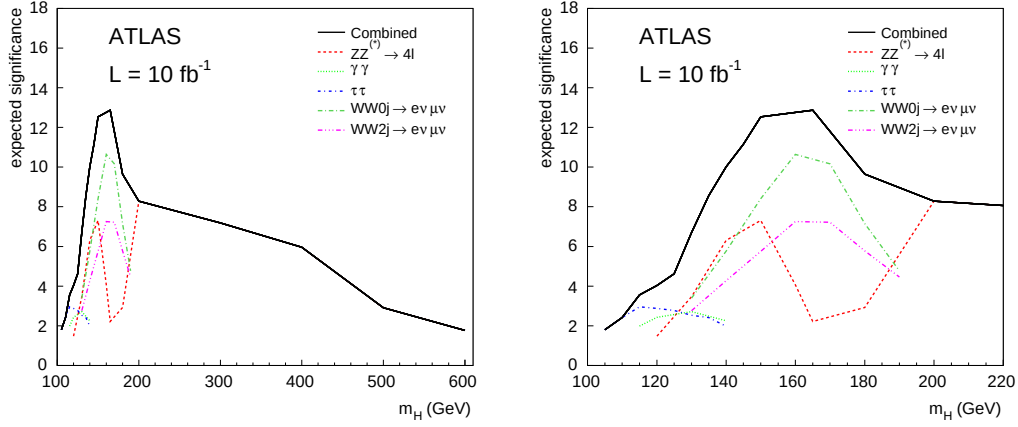


Figure 3.5: Standard Model Higgs boson discovery potential of the ATLAS experiment with an integrated luminosity of 10 fb^{-1} [15]. The left figure shows the complete mass range of 110 GeV–600 GeV, the right figure shows the same in the restricted mass range up to 220 GeV.

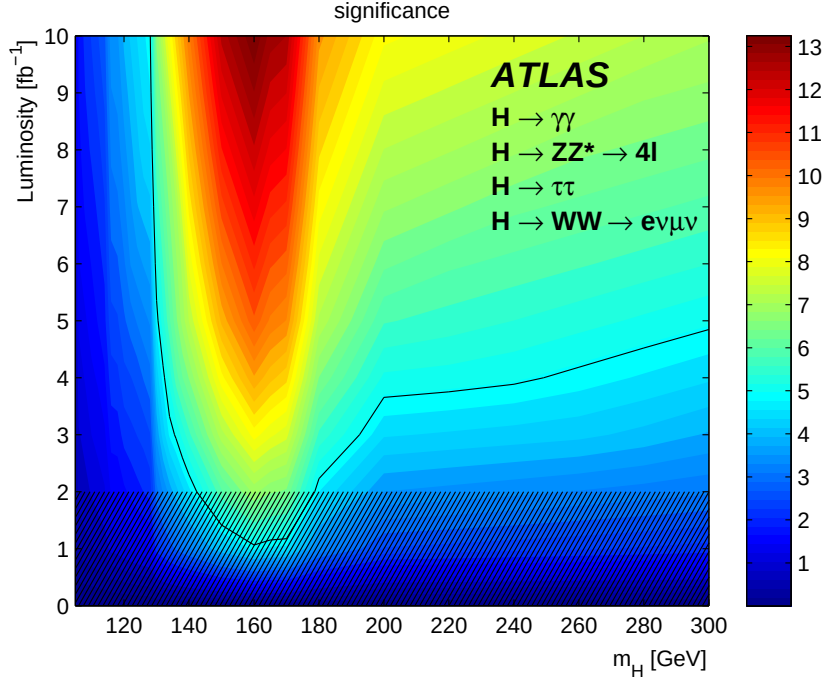


Figure 3.6: Standard Model Higgs boson discovery potential of the ATLAS experiment as function of the integrated luminosity and the Higgs boson mass [15]. The black line is the contour level, at which a combined 5σ discovery potential is expected.

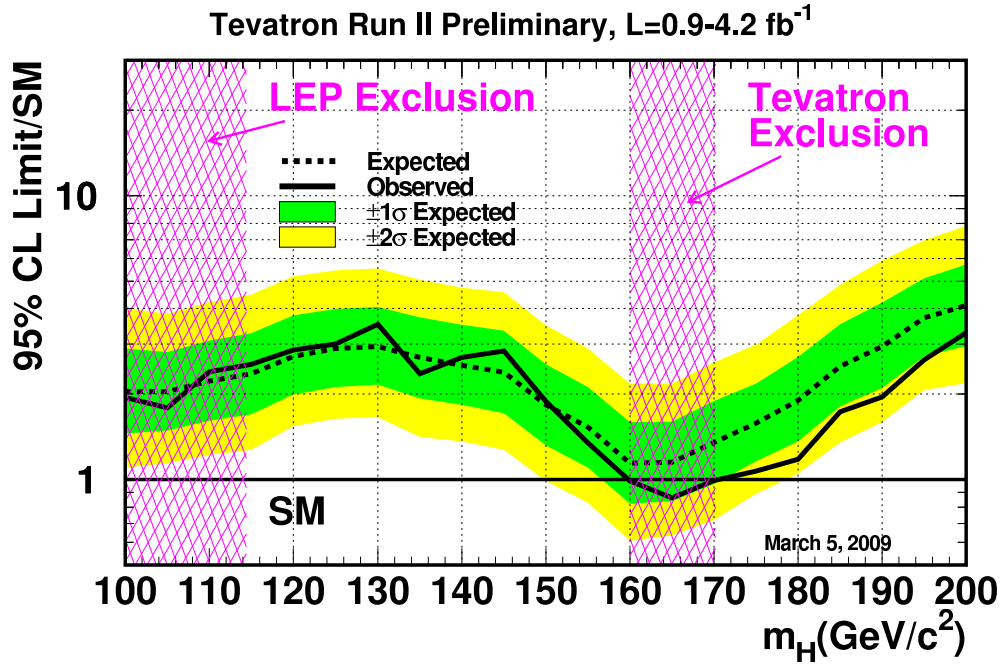


Figure 3.7: Observed and expected (median, for the background-only hypothesis) 95% CL upper limits on the ratios to the Standard Model cross section, as functions of the Higgs boson mass for the combined CDF and D0 analyses [16]. The limits are expressed as multiples of the Standard Model prediction for test masses (every 5 GeV) for which both experiments have performed dedicated searches in different channels. The points are joined by straight lines for better readability. The bands indicate the 68% and 95% probability regions, where the limits can fluctuate in the absence of signal.

4 Determination of Higgs boson properties

One of the most challenging tasks after an observation of an excess of events in one of the Higgs boson discovery channels will be to establish that the observed excess does really come from a (Standard Model) Higgs boson: one massive, CP even, spin 0 particle with couplings to other particles proportional to their mass.

4.1 Measuring the Higgs boson mass

One of the first parameters known after a discovery will be the mass of the possible Higgs boson. All discovery channels have some sensitivity to the mass, but using the narrow mass resonances observable in the $H \rightarrow \gamma\gamma$ and $H \rightarrow ZZ^{(*)}$ final states, a precision measurement with a final accuracy better than 1% should be possible. Figure 4.1 shows the resolution with which the ATLAS experiment is predicted to measure the Higgs boson mass resonance for a 120 GeV Higgs boson in the $H \rightarrow \gamma\gamma$ final state and for a 130 GeV Higgs boson in the $H \rightarrow ZZ^{(*)} \rightarrow 4\mu$ decay mode.

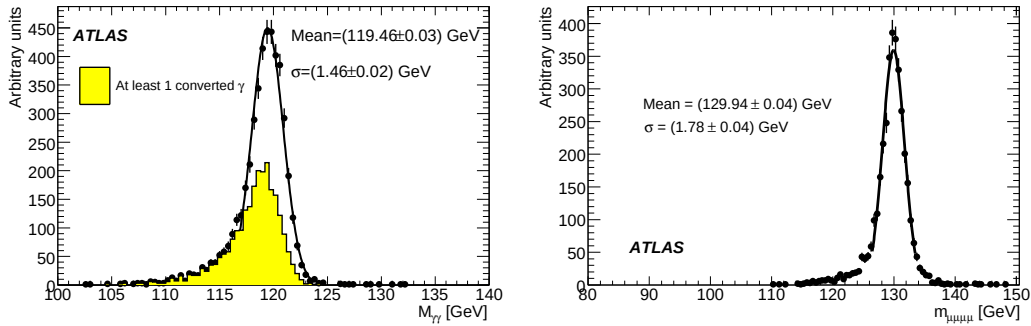


Figure 4.1: Reconstructed Higgs boson mass resonance and mass resolution, as it is expected for the ATLAS experiment [15], for a 120 GeV Higgs boson in the $H \rightarrow \gamma\gamma$ final state and for a 130 GeV Higgs boson in the $H \rightarrow ZZ^{(*)} \rightarrow 4\mu$ decay mode.

4.2 Determining spin and CP quantum numbers

The determination of spin and CP quantum numbers depends on the precise observation of angular correlations in the production and/or decay of the Higgs boson itself and subsequent decays into tertiary particles. However, it should be kept in mind that all non Standard Model configurations of spin and CP also alter the production rate of the (Higgs) boson in a way that is not predictable in general. Thus most studies of this topic are case studies that assume modified angular production/decay correlations while the total production and decay rates are kept at the Standard Model level.

As these studies are rather specific to limited decay modes and Higgs boson mass ranges, the topic is not further elaborated here. An overview of the possibilities to measure the spin and CP quantum numbers at the LHC can be found in Reference [17].

4.3 Determination of the Higgs boson coupling parameters

The Higgs boson coupling parameters can be extracted at the LHC from a combination of all observable (and confirmed as non-observable) Higgs boson production and decay modes [18, 19]. The present coupling analyses combine the results of the studies for the channels shown in Table 4.1 in a global maximum likelihood fit, including correlated theoretical and experimental systematic uncertainties. Details on the method of the combination and on the channels used in the combination can be found in Reference [18].

4.3.1 Limitations of the LHC

The observables at the LHC are measurements of partial cross sections $\sigma(XX \rightarrow H \rightarrow YY)$ for specific Higgs boson production and decay channels. Provided that there is only one sufficiently narrow Higgs boson resonance, this cross section can be expressed by the partial Higgs boson decay widths:

$$\sigma(XX \rightarrow H \rightarrow YY) = c_{XX} \cdot \frac{\Gamma_X \cdot \Gamma_Y}{\Gamma_H} \quad , \quad (4.1)$$

using some channel specific constant c_{XX} that can be calculated from theory.

However, at the LHC it is in principle impossible to observe either Γ_X or the branching ratio $\frac{\Gamma_Y}{\Gamma_H}$ directly. In order to measure Γ_X or $\frac{\Gamma_Y}{\Gamma_H}$, one would need to be able to select events from some Higgs boson production mode $\sigma(XX \rightarrow H)$ for at least one X , without including any signature from the decay in the analysis. Unfortunately, such an unbiased selection of Higgs boson events is not possible at the LHC.

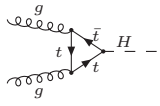
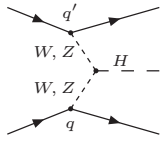
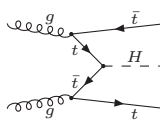
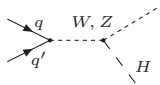
Production	Decay	Mass range
 GF: Gluon Fusion ($gg \rightarrow H$)	$H \rightarrow ZZ^{(*)} \rightarrow 4l$	110 GeV - 200 GeV
	$H \rightarrow WW^{(*)} \rightarrow l\nu l\nu$	110 GeV - 200 GeV
	$H \rightarrow \gamma\gamma$	110 GeV - 150 GeV
 WBF: Weak Boson Fusion ($qq \rightarrow H$)	$H \rightarrow ZZ^{(*)} \rightarrow 4l$	110 GeV - 200 GeV
	$H \rightarrow WW^{(*)} \rightarrow l\nu l\nu$	110 GeV - 190 GeV
	$H \rightarrow \tau\tau \rightarrow l\nu l\nu$	110 GeV - 150 GeV
	$H \rightarrow \tau\tau \rightarrow l\nu l \text{ had} \nu$	110 GeV - 150 GeV
	$H \rightarrow \gamma\gamma$	110 GeV - 150 GeV
 $t\bar{t}H$	$H \rightarrow WW^{(*)} \rightarrow l\nu l\nu (l\nu)$	120 GeV - 200 GeV
	$H \rightarrow b\bar{b}$	110 GeV - 140 GeV
	$H \rightarrow \gamma\gamma$	110 GeV - 120 GeV
 WH	$H \rightarrow WW^{(*)} \rightarrow l\nu l\nu (l\nu)$	150 GeV - 190 GeV
	$H \rightarrow \gamma\gamma$	110 GeV - 120 GeV
ZH	$H \rightarrow \gamma\gamma$	110 GeV - 120 GeV

Table 4.1: List of all ATLAS studies used as input for the combined fit of the Higgs boson coupling parameters [18]. The mass range is the range of the Higgs boson masses considered in the studies.

Also the total width Γ_H of the Higgs boson is not directly measurable from the reconstructed Higgs boson mass distributions in $H \rightarrow \gamma\gamma$ and $H \rightarrow ZZ \rightarrow 4l$ for a Standard Model Higgs boson mass below ≈ 200 GeV, because the experimental resolution is ≈ 1 GeV (see Figure 4.1) and therefore too imprecise to resolve the expected width of the Higgs boson.

Without being able to independently measure at least one of the components in Equation (4.1), it is impossible to extract any other component. Thus it is impossible to make absolute measurements of Higgs boson decay widths (or the total width) at the LHC without further assumptions. Consequently, it is also impossible to get absolute measurements of the Higgs boson coupling parameters.

4.3.2 Measurement of ratios of coupling parameters

However, instead of absolute measurements it is possible to measure ratios of coupling parameters. While measurements of ratios are not sufficient to establish that an observed particle is the Standard Model Higgs boson, they can be used to either find discrepancies from the Standard Model predictions or to rule out models beyond the Standard Model.

Measurement of ratios of Higgs boson decay width

A ratio of Higgs boson decay widths (equivalent to a ratio of branching ratios) can be obtained by combining two (or more) measurements in different production or decay channels. The easiest way of combining two channels is to build the ratio

$$\frac{\Gamma_Y}{\Gamma_{Y'}} = \frac{\sigma(XX \rightarrow H \rightarrow YY)}{\sigma(XX \rightarrow H \rightarrow Y'Y')} \quad . \quad (4.2)$$

When combining a large set of channels, it is advantageous to choose a common denominator Y' . Current studies show that for a Higgs boson mass below ≈ 200 GeV the $H \rightarrow WW$ channels are the best measurable channels over the complete mass range (see Figure 3.5), so $Y' = W$ is a natural choice.

The result of a combination of all channels in Table 4.1 in one global fit of coupling ratios $\frac{\Gamma_Y}{\Gamma_W}$, assuming in integrated luminosity of 30 fb^{-1} , is shown in Figure 4.2(a).

At best precisions of the order of $\approx 20\%$ can be reached for the ratio $\frac{\Gamma_Z}{\Gamma_W}$, while the measurement of $\frac{\Gamma_\gamma}{\Gamma_W}$ and $\frac{\Gamma_\tau}{\Gamma_W}$ is possible with precisions of 30% to 60%.

Measurement of ratios of Higgs boson couplings

In the combined fit of ratios of Higgs partial decay widths only the coupling information from the Higgs boson decay is used. This can be improved by including also the coupling information accessible through the Higgs boson production mode.

To do this, it is assumed that the tree level loop diagrams in the gluon fusion Higgs boson production and the $H \rightarrow \gamma\gamma$ decay are mediated only via Standard Model particles. Under this assumption it is possible to express all Higgs boson production channels and partial decay widths by the fundamental Higgs boson couplings to the W and Z bosons as well as to the b , t and τ fermions.

In general, the Higgs boson production cross sections are a function of the square of a Higgs boson coupling multiplied by a constant that can be calculated perturbatively:

$$\begin{aligned} \sigma_{ggH} &= \alpha_{ggH} \cdot g_t^2 \\ \sigma_{WBF} &= \alpha_{WF} \cdot g_W^2 + \alpha_{ZF} \cdot g_Z^2 \\ \sigma_{t\bar{t}H} &= \alpha_{t\bar{t}H} \cdot g_t^2 \\ \sigma_{WH} &= \alpha_{WH} \cdot g_W^2 \\ \sigma_{ZH} &= \alpha_{ZH} \cdot g_Z^2 \quad . \end{aligned} \quad (4.3)$$

Similarly, all Higgs boson branching ratios can be expressed in terms of the

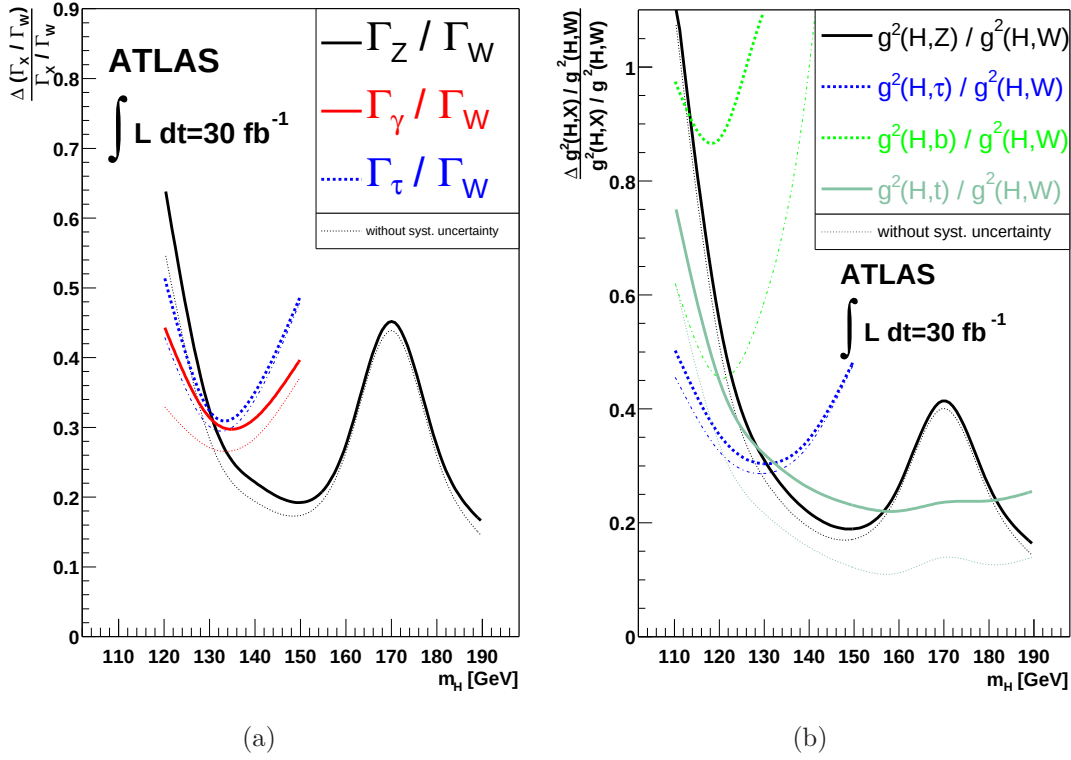


Figure 4.2: Expected accuracy of the ATLAS experiment [18] for the determination of ratios of Higgs boson coupling parameters using an integrated luminosity of 30 fb⁻¹. Results without systematic uncertainties are shown in dashed lines.

(a): Precision of measurements of ratios of Higgs boson decay widths.
 (b): Precision of measurements of ratios of Higgs boson coupling squares.

coupling squares and the total Higgs boson width:

$$\begin{aligned}
 \text{BR}(H \rightarrow WW) &= \beta_W \frac{g_W^2}{\Gamma_H} \\
 \text{BR}(H \rightarrow ZZ) &= \beta_Z \frac{g_Z^2}{\Gamma_H} \\
 \text{BR}(H \rightarrow \gamma\gamma) &= \frac{(\beta_{\gamma(W)} \cdot g_W - \beta_{\gamma(t)} \cdot g_t)^2}{\Gamma_H} \\
 \text{BR}(H \rightarrow \tau\tau) &= \beta_\tau \frac{g_\tau^2}{\Gamma_H} \\
 \text{BR}(H \rightarrow b\bar{b}) &= \beta_b \frac{g_b^2}{\Gamma_H} .
 \end{aligned} \tag{4.4}$$

The constants α_i and β_j in Equations 4.4 and 4.5 can be calculated within the Standard Model without further knowledge of the Higgs boson parameters. The expected theoretical uncertainties in the calculation of α_i and β_j are listed in Table 4.2.

Production		Decay	
α_{ggH}	20%	β_W	1%
α_{WF}	4%	β_Z	1%
α_{ZF}	4%	$\beta_{\gamma(W)}$	1%
$\alpha_{t\bar{t}H}$	15%	$\beta_{\gamma(t)}$	1%
α_{WH}	7%	β_τ	1%
α_{ZH}	7%	β_b	1%

Table 4.2: Overview of the systematic uncertainties assumed for the fit of relative Higgs boson couplings [18]. The uncertainties on W - and Z -fusion (α_{WF} and α_{ZF}) are treated as 100% correlated.

Apart from interference terms in the Higgs boson loop processes only squares of Higgs boson couplings appear in the expressions for the production cross sections and partial decay widths. Therefore squares of coupling ratios

$$\frac{g_Z^2}{g_W^2} \quad , \quad \frac{g_t^2}{g_W^2} \quad , \quad \frac{g_b^2}{g_W^2} \quad \text{and} \quad \frac{g_\tau^2}{g_W^2} \quad (4.5)$$

are extracted from the global fit together with the free parameter

$$\frac{g_W^2}{\sqrt{\Gamma_H}} \quad , \quad (4.6)$$

which cannot be resolved into its components due to the limitations of the LHC.

Figure 4.2(b) shows the expected accuracy of the ATLAS experiment in measuring the squares of ratios of fundamental Higgs boson couplings assuming an integrated luminosity of 30 fb^{-1} . Especially the measurement of ratios involving the top quark Yukawa coupling is of high interest, as the top quark plays a special role in many theories beyond the Standard Model. Due to the large contribution of the top quark to the well measurable gluon fusion production process, precisions of 20% to 30% are reachable for sufficiently high Higgs boson masses. However, one should note that this accuracy relies on the assumption that no unknown particles contribute to the gluon fusion process. Without this assumption, the top quark coupling can only be extracted from the less precisely measurable $t\bar{t}H$ channels.

4.3.3 Determination of absolute coupling parameters

In Section 4.3.1 it was elaborated that an absolute measurement of Higgs boson coupling parameters is not possible at the LHC without some assumption that imposes an absolute limit on at least one coupling parameter.

Already without any assumption one can get a lower limit on the total width of the Higgs boson, as it has to be larger than the sum of all observable decay widths:

$$\Gamma_H > \sum_Y \Gamma_Y \quad . \quad (4.7)$$

This implies that also all partial decay widths have a lower limit, as can be seen from the following simple example:

- The partial cross section for the WBF process $qq H \rightarrow qq WW$ (or similarly the process $WH(H \rightarrow WW)$ could be taken) is measurable and can be taken as a constant for the following argument:

$$\sigma(qq H \rightarrow qq WW) = c_{\text{WBF}} \cdot \frac{\Gamma_W \cdot \Gamma_W}{\Gamma_H} = \text{const.} \quad (4.8)$$

In order to simplify the argumentation, the relatively small contribution from Z boson fusion to the WBF production cross section is ignored.

- From Equation (4.7) follows:

$$\Gamma_H = \Gamma_W + x \quad , \quad \text{with } x > 0 \quad . \quad (4.9)$$

- Combining these two equations and solving for Γ_W gives:

$$\Gamma_W = \frac{\text{const.}}{c_{\text{WBF}}} \cdot \left(1 + \frac{x}{\Gamma_W} \right) \quad , \quad (4.10)$$

where all involved quantities are positive. Thus Γ_W has a non zero lower limit and a similar argument can be given for all other Higgs boson partial widths and hence also for the total width.

In Figure 4.3 the lower limit on the total Higgs boson width extracted from all observable channels is shown together with the theoretical prediction for the Standard Model width and the direct experimental resolution for the width from the $H \rightarrow \gamma\gamma$ and $H \rightarrow ZZ$ final states. Especially for low Higgs boson masses, the limit is one order of magnitude below the actual prediction and thus not very precise.

No similar argument can be made for an upper limit on any of the Higgs boson coupling parameters. However, the Higgs boson serves a second important role in the theory of the Standard Model, besides generating gauge invariant mass terms. Without the Higgs boson, the cross section for the scattering process $WW \rightarrow WW$ diverges with increasing centre-of-mass energy. The presence of the Higgs boson preserves unitarity through destructive interference:

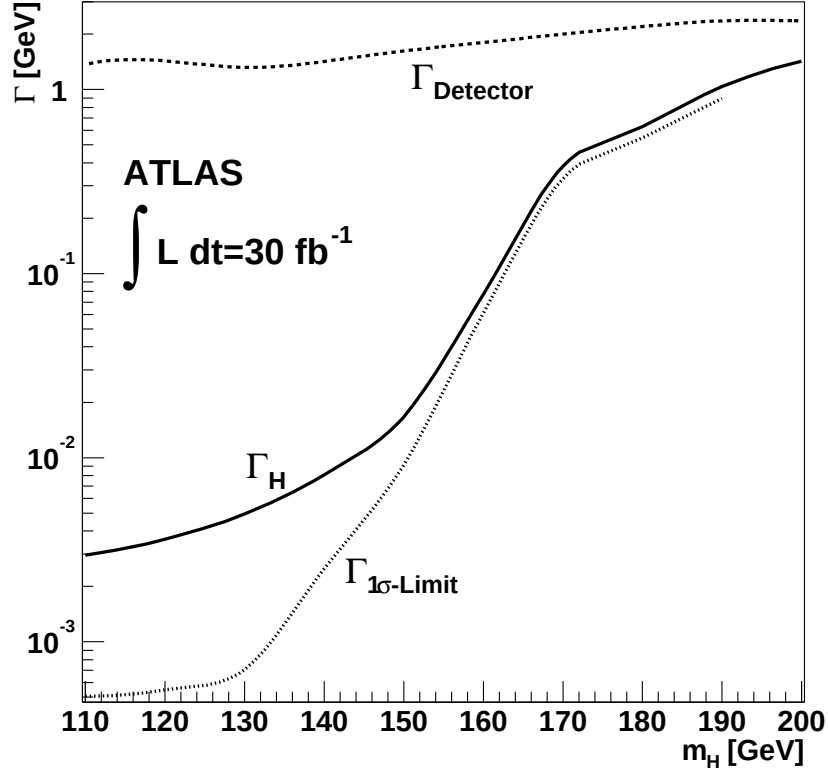


Figure 4.3: Lower 1σ limit on the total Higgs boson width for the ATLAS experiment [18] shown together with the theoretical prediction Γ_H for the Standard Model width and the direct experimental resolution Γ_{Detector} .

$$\left[\begin{array}{c} \text{Diagram 1: } W \text{ and } W^* \text{ exchange via } \gamma, Z \\ \text{Diagram 2: } W \text{ and } W^* \text{ exchange via } \gamma, Z \\ \text{Diagram 3: } W \text{ and } W^* \text{ exchange via } \gamma, Z \\ \text{Diagram 4: } W \text{ and } W^* \text{ exchange via } H \\ \text{Diagram 5: } W \text{ and } W^* \text{ exchange via } H \end{array} \right]^2 < \infty .$$

This cancellation is only possible, if the couplings to the W and Z boson of all Higgs like particles H_i in the theory sum up to exactly the Standard Model value for this coupling:

$$\sum_{H_i} g^2(H_i, W) = g^2(H_{\text{SM}}, W) \quad . \quad (4.11)$$

Since all squares of couplings are positive, Equation (4.11) gives an absolute upper limit on the strength of the coupling to W and Z for every Higgs boson

considered in the sum:

$$\begin{aligned} g^2(H, W) &< 1.05 \cdot g^2(H_{\text{SM}}, W) \\ g^2(H, Z) &< 1.05 \cdot g^2(H_{\text{SM}}, Z) \end{aligned} \quad (4.12)$$

The extra 5% margin allows for theoretical uncertainties in the translation between squares of couplings and partial widths, and also for small admixtures of exotic Higgs states, like SU(2) triplets. Even without the extra margin, Equation (4.12) is valid for models with an arbitrary number of Higgs singlets and doublets. Further details can be found in Reference [19].

Due to the additional constraints from Equations (4.7) and (4.12) it is possible to fit the numerator and denominator of Equation (4.6) separately and hence the absolute Higgs boson couplings

$$g_W^2, \quad g_Z^2, \quad g_t^2, \quad g_b^2 \quad \text{and} \quad g_\tau^2 \quad (4.13)$$

can be extracted from the global fit together with the total Higgs boson width Γ_H and the Higgs boson decay width Γ_{inv} to invisible and/or undetectable Higgs boson decay modes.

Figure 4.4 shows the precision with which absolute Higgs boson coupling squares and a limit on the width Γ_{inv} can be determined from the combination of two LHC experiments, assuming an integrated luminosity of 30 fb^{-1} per experiment.

Below a Higgs boson mass of $m_H \approx 160 \text{ GeV}$, all uncertainties on the coupling measurements are fairly large at 40% to 100% relative precision due to the large uncertainty on the measurement of the b -quark partial width Γ_b that dominates the total Higgs width up to a mass of $m_H \approx 135 \text{ GeV}$. Above 160 GeV, where the b contributions become negligible, relatively high precisions of 10% to 40% on the squares of absolute couplings are reachable. Depending on the Higgs boson mass, limits of 20% to 50% can be set on the width Γ_{inv} to invisible and/or undetectable Higgs boson final states.

The assumption from Section 4.3.2 that only Standard Model particles contribute to the Higgs boson loop processes, is not essential for the fit. Therefore the global fit can be extended to include the additional parameters

- $\Gamma_{\gamma(\text{new})}$ for the $H \rightarrow \gamma\gamma$ decay mode
- $\Gamma_g(\text{new})$ for the gluon fusion production mode

to account for unknown particles contributing to the gluon fusion production process and the $H \rightarrow \gamma\gamma$ decay mode.

Figure 4.5 shows the precision with which the absolute Higgs boson couplings and the additional partial widths due to unknown particles in the model can be determined from the combination of two LHC experiments, assuming an integrated luminosity of 30 fb^{-1} per experiment.

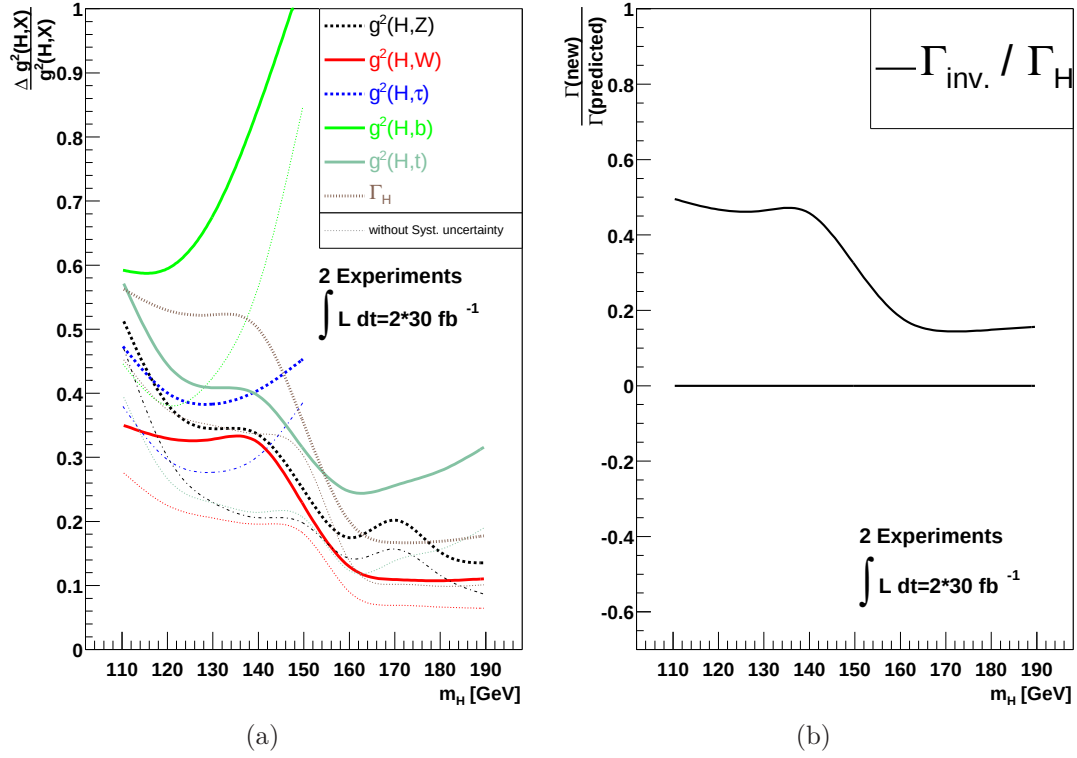


Figure 4.4: Expected accuracy of two LHC experiments [19] in the determination of absolute Higgs boson coupling parameters using an integrated luminosity of $2 \cdot 30 \text{ fb}^{-1}$. It is assumed that only Standard Model particles contribute to the Higgs boson loop processes. Further theoretical assumptions are needed for the parameter extraction, details are given in the text.

(a): Precision on measurements of absolute Higgs boson coupling squares.

(b): Limit on additional partial decay widths due to unknown particles.

Due to the additional fit parameters for unknown particles in Higgs boson loops, the precision of the coupling determination deteriorates slightly with respect to Figure 4.4. However, the overall picture is not modified significantly and only very loose limits can be set on additional contributions from new particles.

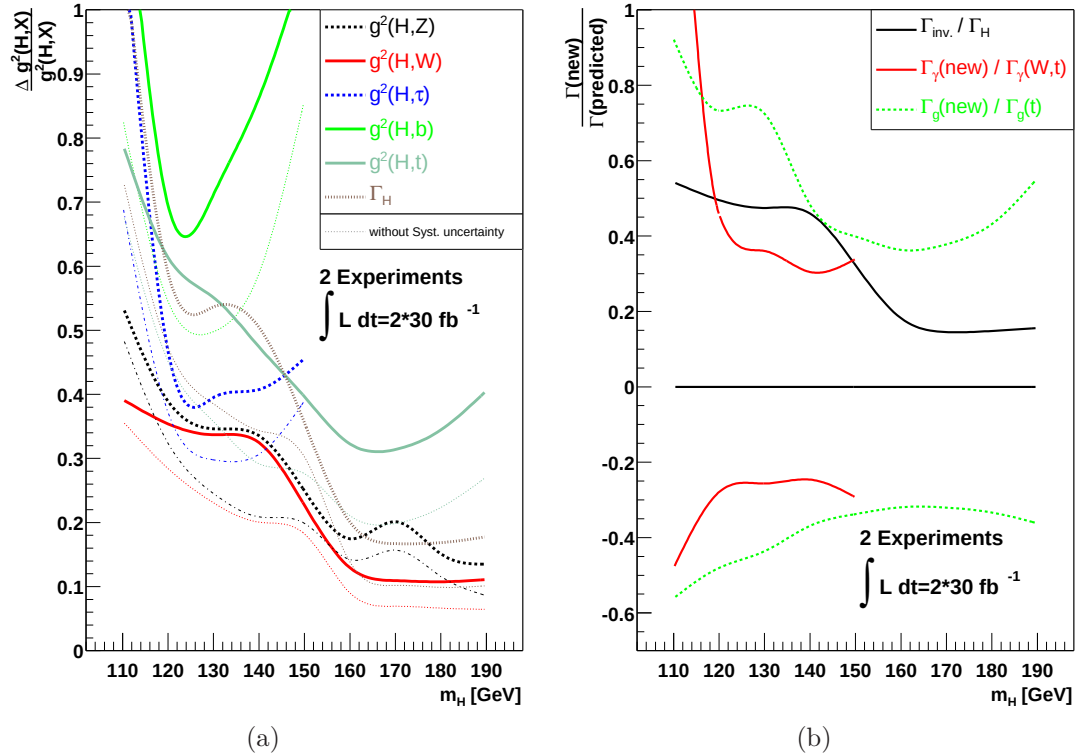


Figure 4.5: The same as Figure 4.4, but loop contributions from unknown particles to the gluon fusion production process and the $H \rightarrow \gamma\gamma$ decay mode are allowed.

5 The Large Hadron Collider and the detector system of the ATLAS experiment

During the last ≈ 10 years the LHC¹ was installed into the already existing tunnel of the LEP² accelerator ring at the European centre for particle physics CERN³ close to the Swiss town of Geneva. Together with the accelerator infrastructure four new experiments were built into the LHC tunnel at four interaction regions. A schematic view of the LHC ring and the layout of the CERN accelerator complex is shown in Figure 5.1.

The two major multi purpose detectors ATLAS⁴ [21] and CMS⁵ [22] are intended for the search of new phenomena in proton proton collisions at so far unreached energies and luminosities. The focus of the LHCb⁶ experiment is the understanding of the properties of b -quarks and CP violating decays of hadrons. Finally, ALICE⁷ is built to measure the collisions of heavy ions to analyse the behaviour of strongly interacting particles at extreme high densities and temperatures.

The LHC accelerator and experiments were already completely built and commissioned until September 2008 and the first beams were successfully circulated on 10th September 2008. Unfortunately on 19th September a serious fault damaged a number of superconducting magnets which resulted in more than one year of repair and re-engineering work. The new startup of the LHC is currently planned for November 2009.

In the following chapter a short overview of the LHC accelerator and the ATLAS detector with all its sub-detector systems will be given.

¹**LHC**: Large Hadron Collider

²**LEP**: Large Electron-Positron collider

³**CERN**: Conseil Européen pour la Recherche Nucléaire

⁴**ATLAS**: A Toroidal LHC ApparatuS

⁵**CMS**: Compact Muon Spectrometer

⁶**LHCb**: Large Hadron Collider beauty

⁷**ALICE**: A Large Ion Collider Experiment

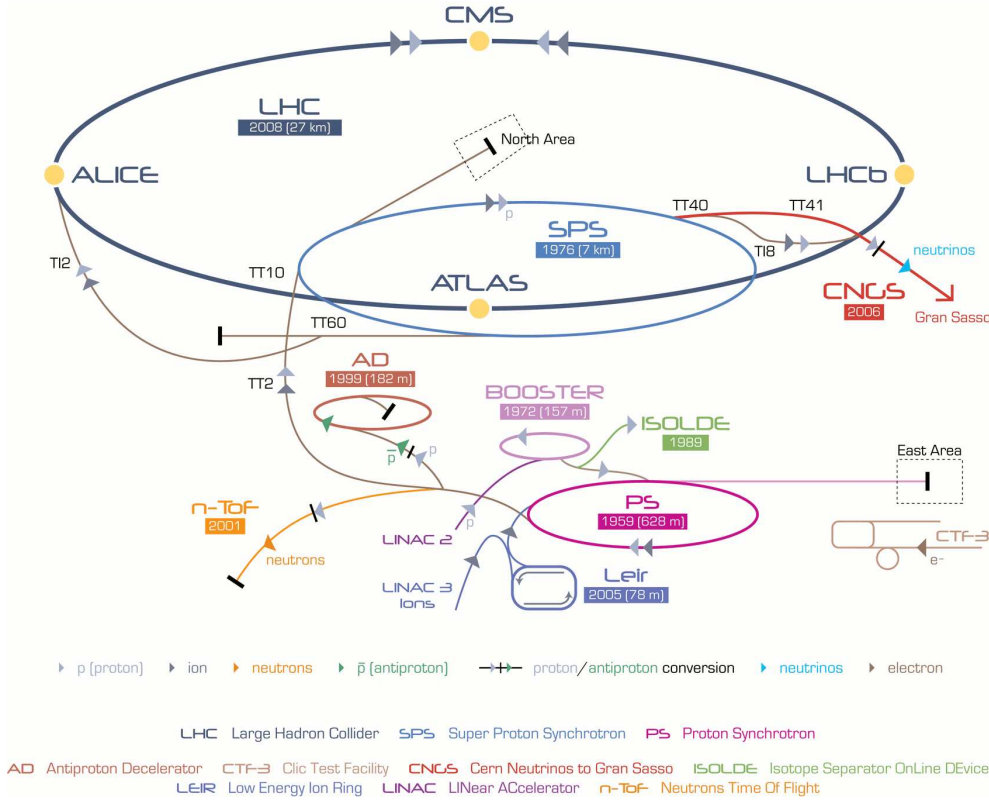


Figure 5.1: Schematic view of the position of the four LHC experiments and the layout of the CERN accelerator complex [20].

5.1 The Large Hadron Collider

The LHC is a proton proton collider with a centre-of-mass energy of 14 TeV and a design luminosity of up to 10^{34} particles/cm²/s. Besides the proton collision mode, on which this introduction will focus, the accelerator is also suited as heavy ion collider.

The LHC tunnel lies ≈ 100 m underground and has a circumference of 27 km. Two contrariwise circulating proton beams with a foreseen energy of 7 TeV each are brought to collision in the four interaction points positioned around the LHC ring (see Figure 5.1). Each beam will consist of 2808 so called bunches with $\sim 10^{11}$ protons per bunch. Each bunch has a length of 7.5 cm $\simeq 0.25$ ns and a diameter of 7 μ m. The current per beam will be 0.53 A.

To keep the high energy protons on track in the ring, super conducting dipole magnets are used that produce a magnetic field of 8.33 Tesla at a temperature of 1.9 K. The acceleration and the loss of energy due to synchrotron radiation once the protons have reached their design energy is coupled into the beam through 8

super conducting cavities with a field gradient of 5 MV/m.

Protons are accelerated to an injection energy of 450 GeV by the existing SPS⁸. The acceleration from 450 GeV to 7 TeV will take roughly 20 minutes. After this injection phase the half-life of the beam will be approximately 10 hours. The most important technical parameters of the LHC are given in Table 5.1.

centre-of-mass energy	14 TeV
injection energy	450 GeV
luminosity	$1 \cdot 10^{34} \frac{1}{\text{cm}^2\text{s}}$
Relativistic gamma	7461
beam current	0.582 A
bunch spacing	7.48 m
bunch spacing	24.95 ns
number of protons per bunch	$1.1 \cdot 10^{11}$
number of bunches	2808
beam-beam angle at the interaction point	$300 \mu \text{ rad}$
half-life-period of the luminosity	10 h
energy loss per cycle	7 keV
synchrotron radiation per beam	3.8 kW
stored energy per beam	362 MJ
number of dipole magnets	1232
dipole field strength at 7 TeV per beam	8.33 Tesla

Table 5.1: Design parameters of the LHC.

The unprecedented challenges in the design of the LHC detectors can be seen from the cross section and event rate overview shown in Figure 5.2. At the design luminosity, the LHC will produce 10^9 collisions per second. W and Z bosons, the particles that were deeply analysed in the previous LEP experiments, are produced at a rate of 10^3 - 10^4 Hz. Events from new physics are expected to be suppressed by a factor of $\approx 10^{10}$ or more with respect to the total proton-proton interaction rate. These rare events need to be filtered out, measured in detail and recorded by the experiments.

The LHC parameters given so far are the design parameters. Especially in the startup phase, lower centre-of-mass energies and luminosities will be used. The current plan foresees to start with injection energy and then to accelerate to $\sqrt{s} = 7$ TeV collision energy at luminosities of $\approx 10^{31} \frac{1}{\text{cm}^2\text{s}}$. The Higgs analysis presented in the Chapters 8 to 10 is built on the assumption of an initial collision energy of $\sqrt{s} = 10$ TeV, as this was the estimated startup energy during the last year.

⁸SPS: Super Proton Synchrotron

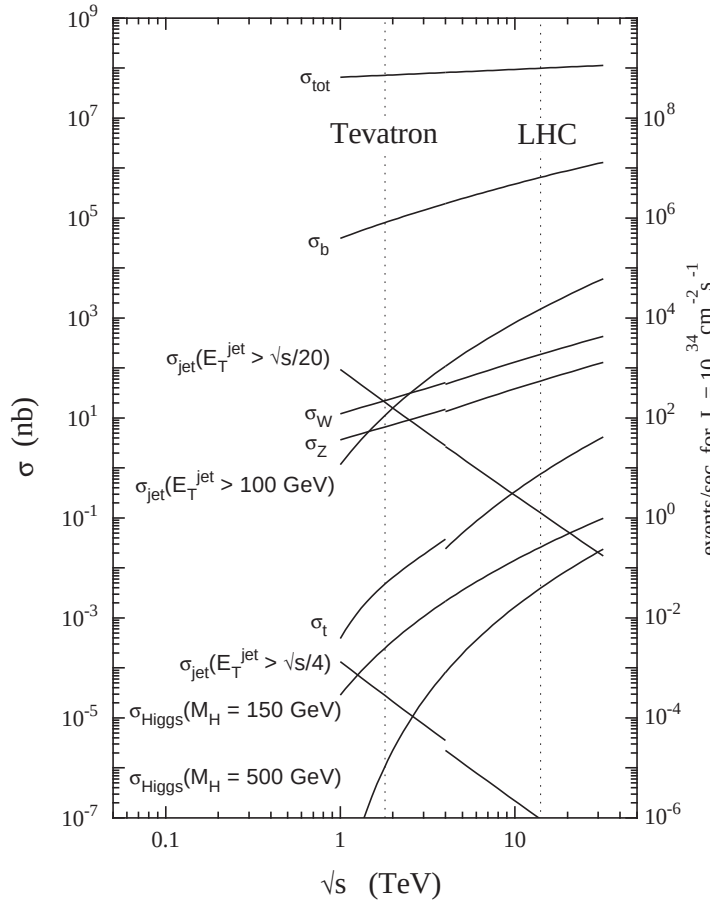


Figure 5.2: Cross sections and production rates for important processes at the LHC [23]; The left axis shows the cross section of a process, the right axis the expected rate. Values are given for both the Tevatron ($\sqrt{s} = 1.96 \text{ TeV}$) and the LHC ($\sqrt{s} = 14 \text{ TeV}$). The jump in the cross sections arises due to the fact that the Tevatron is a $p\bar{p}$ machine, while the LHC is a pp collider.

5.2 The ATLAS experiment

The ATLAS experiment is designed as a multi purpose high energy physics detector to exploit the full discovery potential of the LHC. The LHC will offer a large range of physics opportunities, ranging from Standard Model physics to potential black hole production. A major focus of interest will be the origin of mass at the electroweak scale, as it was outlined in the Chapters 1 to 3.

ATLAS is a large scale project that has around 1800 members from 35 countries. The proposal for the ATLAS experiment [21] was published in 1994 and the detector construction was finalised in 2008, in time for the first LHC beam.

As discussed in Chapter 3, one of the possible manifestations of the spontaneous electroweak symmetry-breaking mechanism could be the existence of one or several Higgs bosons. The wish to ensure the discovery of these Higgs particles, if realised in nature, has guided the design of the ATLAS detector alongside many studies to ensure that the detector concept is also suited for a large range of other physics scenarios, e.g. general searches for Supersymmetry, heavy additional vector bosons and other exotic particles.

These searches for existing Standard Model physics and new phenomena require the detector to be able to measure several final state objects like electrons, muons, photons, jets (from τ leptons, b quarks and light flavours) and missing transverse energy over a wide kinematical range. From these the design considerations for the ATLAS detector can be summarised as follows:

- good energy resolution and coverage for electrons and photons in the central electronic calorimeter (EM-Barrel) and in the **electromagnetic endcap calorimeter (EMEC)**;
- good energy resolution and coverage for hadronic showers in the central hadronic calorimeter (Tile-Calorimeter) and in the **hadronic endcap calorimeter (HEC)**;
- large coverage of the calorimeter system for a precise measurement of the missing transverse momentum;
- precise momentum measurement for leptons in the inner detector;
- precise momentum measurement for muons in the stand alone muon system;
- flexible and fast multi level trigger system to select events of physics interest.

The ATLAS detector consists of several sub-detector systems arranged radially around the interaction point, as shown in Figure 5.3. The calorimeters reach a coverage in pseudorapidity⁹ of $|\eta| \leq 4.9$, while the inner detector and the muon system provide track reconstruction for $|\eta| \leq 2.5$ (can be up to $|\eta| = 2.7$ for the muon system). In the following section the sub-detector systems will be described shortly. A detailed description of the whole detector, including technical aspects, can be found in Reference [24].

5.2.1 Inner detector

The inner detector (ID) is the sub-detector closest to the interaction point as can be seen in Figure 5.4. It is ≈ 7 m long, has an outer radius of 1.15 m and is placed in a 2 Tesla strong magnetic field, produced by a solenoid magnet situated between the ID and the calorimeter.

⁹Pseudorapidity: $\eta = -\ln(\tan \frac{\theta}{2})$, $\theta \equiv$ Polar angle

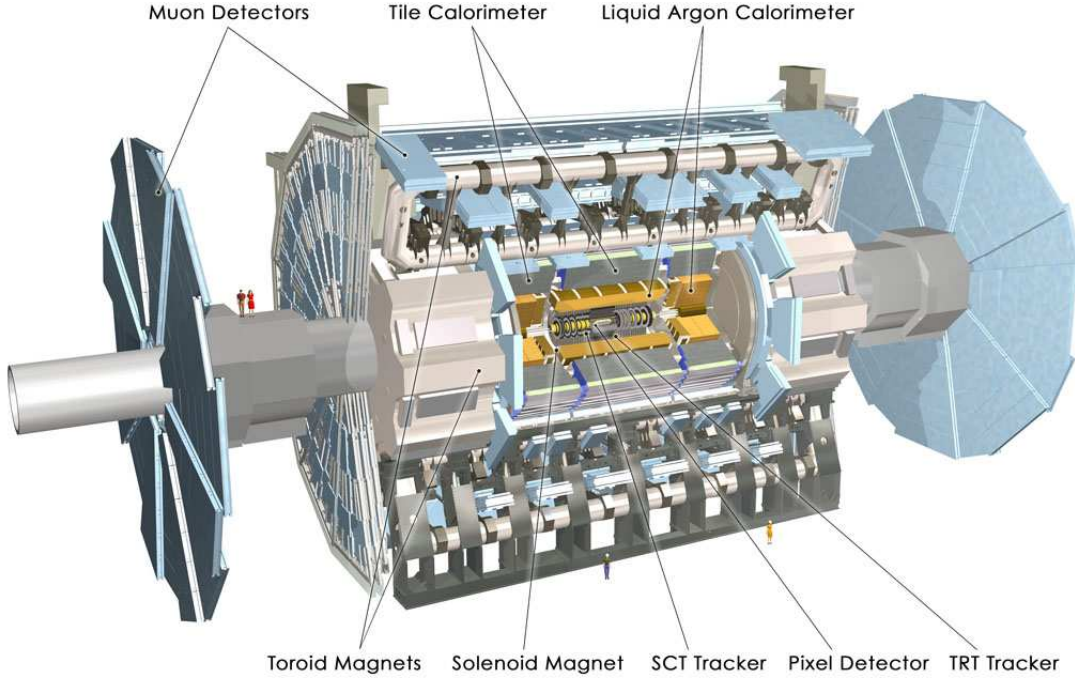


Figure 5.3: Cut-away view of the ATLAS detector [24]. The dimensions of the detector are 25 m in height and 44 m in length. The overall weight of the detector is approximately 7000 tons.

The ID is composed of three different detector sub systems, decreasing in granularity when going away from the interaction point. Closest to the beam pipe are three layers of silicon pixel detectors (PIXEL) which provide three space points per charged particle track. This is followed in the barrel by four stereo layers of silicon strip detectors (SCT) providing four space points and 9 SCT discs in the endcap. The PIXEL and SCT are surrounded by a transition radiation tracker (TRT) which usually provides 36 hits and additional electron - pion separation capabilities.

The inner detector is divided along the beam direction into a barrel and two end cap sections. In the barrel region, layers are arranged in concentric cylinders around the beam pipe while in the end cap region they are mounted as discs perpendicular to the beam axis. In this way a coverage of $|\eta| \leq 2.5$ is achieved. Figure 5.5 and Table 5.2 give an overview of the positioning of the active silicon elements in the inner detector.

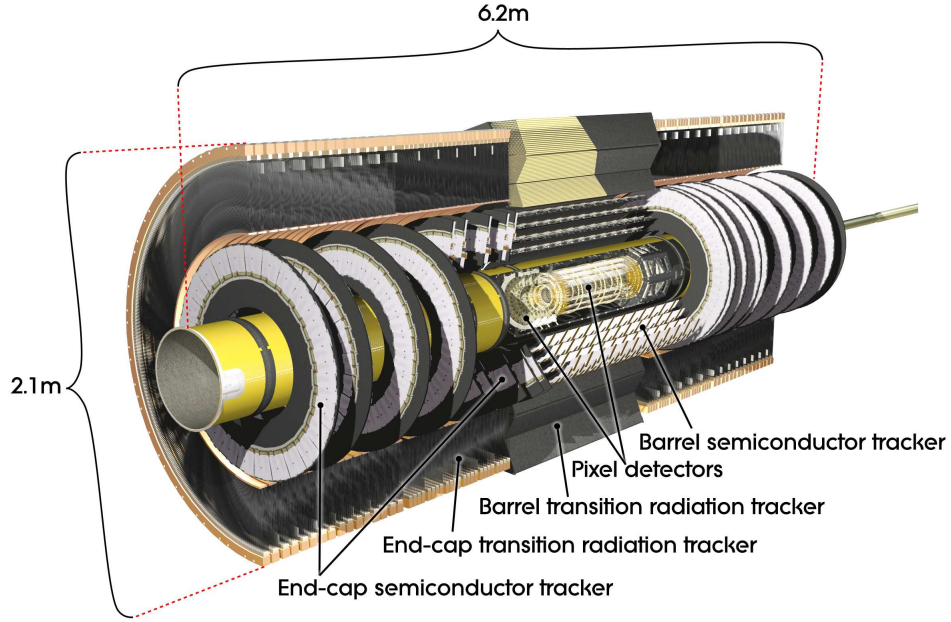


Figure 5.4: Cut-away view of the ATLAS inner detector [24].

PIXEL detector

The PIXEL detector consists of barrel layers and discs containing modules. Each module has 46080 elements which amounts to a total of 80 million pixels with a pitch of $50\ \mu\text{m}$ in the $R\text{-}\phi$ direction and $400\ \mu\text{m}$ in the z direction. The segmentation in two dimensions gives space point measurements without the usual ambiguities encountered when using projective geometries. However, this leads to complex electronics and it requires a very large number of read out channels.

SCT detector

The SCT is built of silicon microstrip detectors, with a fixed pitch of $80\ \mu\text{m}$ in the barrel and a variable pitch of $70\ \mu\text{m}$ to $95\ \mu\text{m}$ in the endcap. A small stereo angle of $40\ \text{mrad}$ between two sandwiched modules provides a z measurement. Altogether, the SCT consists of 4088 modules and provides 6.2 million read out channels.

TRT detector

The Transition Radiation Tracker (TRT) is composed of straw detectors, which can operate at the expected high rates due to their small diameter and the isolation of the sense wires within individual gas volumes. The capability to identify

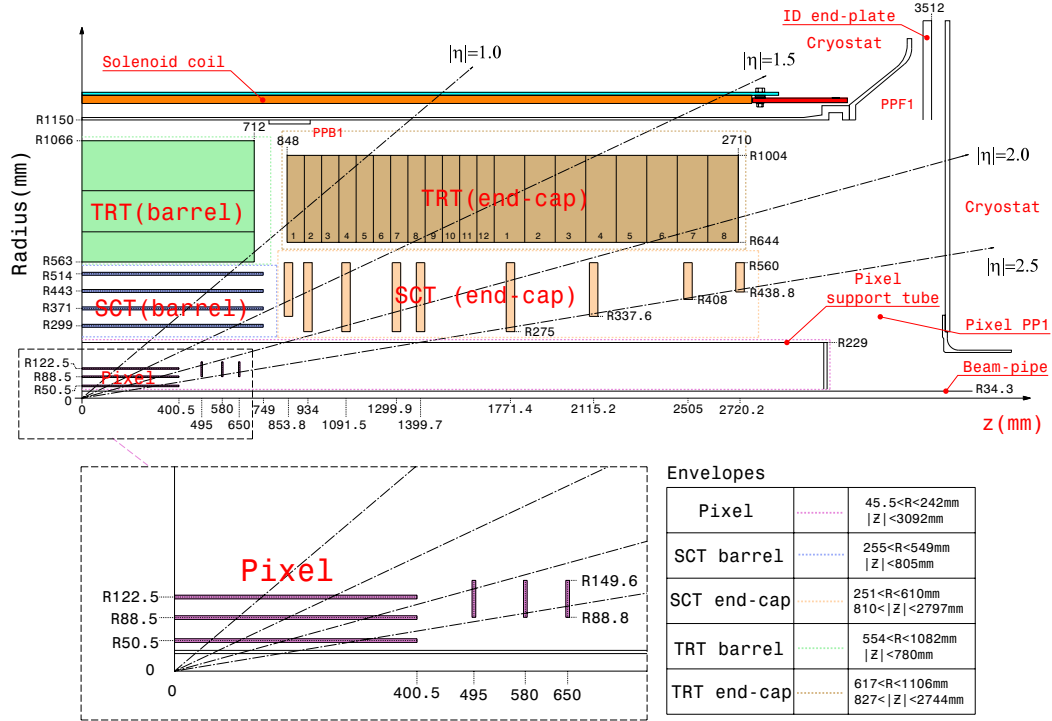


Figure 5.5: Schematic view of a quarter-section of the ATLAS inner detector showing each of the major detector elements with its active dimensions and envelopes [24].

Barrel		
System	Radii [mm]	η coverage
Pixel	50.5	± 2.5
	88.5, 122.5	± 1.7
SCT	299, 371, 443, 514	± 1.4
Endcap		
System	z [cm]	η coverage
Pixel	49.5, 58.0, 65.0	$\pm 1.7 - 2.5$
SCT	83.5, 92.5, 107.5	$\pm 1.4 - 2.5$
	126.0, 145.0, 169.5	$\pm 1.4 - 2.5$
	213.5, 252.8, 272.8	$\pm 1.4 - 2.5$

Table 5.2: Position and coverage of inner detector silicon sub systems.

electrons is added by employing Xenon gas to detect transition radiation photons created in a radiator between the straws. Each straw is 4 mm in diameter and equipped with a 30 μm diameter gold-plated W-Re wire. The straws are divided into two parts at the centre and read out at both ends, to reduce the occupancy. The end-caps contain 320 000 radial straws, with the read out at the outer radius. Each of the total of 420 000 channels provides a drift time measurement, giving a spatial resolution of 170 μm per straw, and two independent thresholds. These allow the detector to discriminate between tracking hits, which pass the lower threshold, and transition radiation hits, which pass the higher threshold.

Summary

The inner detector will provide the measurement of tracks with their direction, momentum and impact parameters for charged particles. These tracks can then be combined into primary or secondary vertices. The momentum is measured through the curvature of charged particle tracks in the magnetic field. A detailed description of the vertex and track reconstruction is given in Reference [15]. It is expected that the primary vertex can be reconstructed with a resolution of $\Delta x \cdot \Delta y \cdot \Delta z = 11 \mu\text{m} \cdot 11 \mu\text{m} \cdot 40 \mu\text{m}$ and that track momenta can be measured with a resolution of $\approx 5 \%$ at $p_T \sim 100 \text{ GeV}$ [15].

The material budget of the inner detector is important for the behaviour of particles before they reach the calorimeter. This material budget vs. $|\eta|$ is given in Figure 5.6. It shows the contributions from the PIXEL, SCT, TRT and service infrastructure separately. In the central barrel, the inner detector system contributes only half a radiation length X_0 . However, this increases to ≈ 2 radiation lengths around $|\eta| \approx 1.5$ (in the so called “crack region”), where most of the service infrastructure is located.

5.2.2 Calorimeter system

The calorimeter system of ATLAS is composed of several components that cover different solid angles and longitudinal depths. An overview of the whole calorimeter system can be seen in Figure 5.7. The main geometrical parameters are summarised in Table 5.3.

Electromagnetic calorimeter

The first calorimeter system that particles enter after the exit from the inner detector system is the electromagnetic (EM) calorimeter. This calorimeter is responsible for the measurement of electromagnetically interacting particles: electrons and photons. It is divided into the EM-Barrel in the central region and EMEC system in the more forward region of the detector. The EM-Barrel covers a range in pseudorapidity of $|\eta| < 1.475$, while the EMEC covers $1.375 < |\eta| < 3.2$.

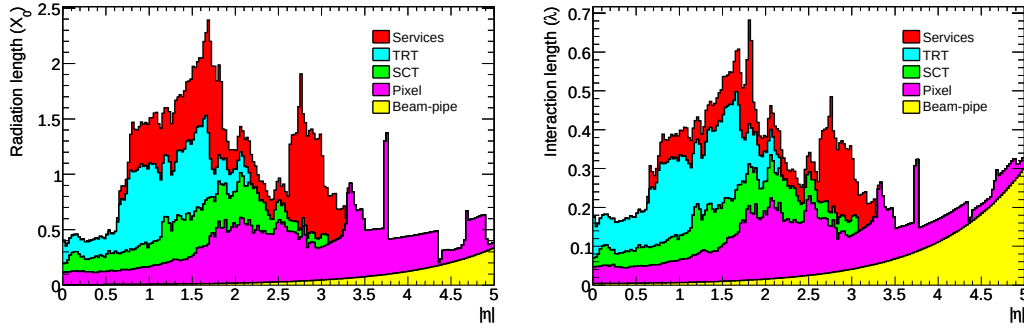


Figure 5.6: Material distribution in units of radiation length X_0 (left) and hadronic interaction length λ (right) at the exit of the inner detector envelope, including the services and thermal enclosures [24]. The distribution is shown as a function of $|\eta|$ and averaged over ϕ . The breakdown indicates the contributions of external services and of individual sub-detectors, including services in their active volume.

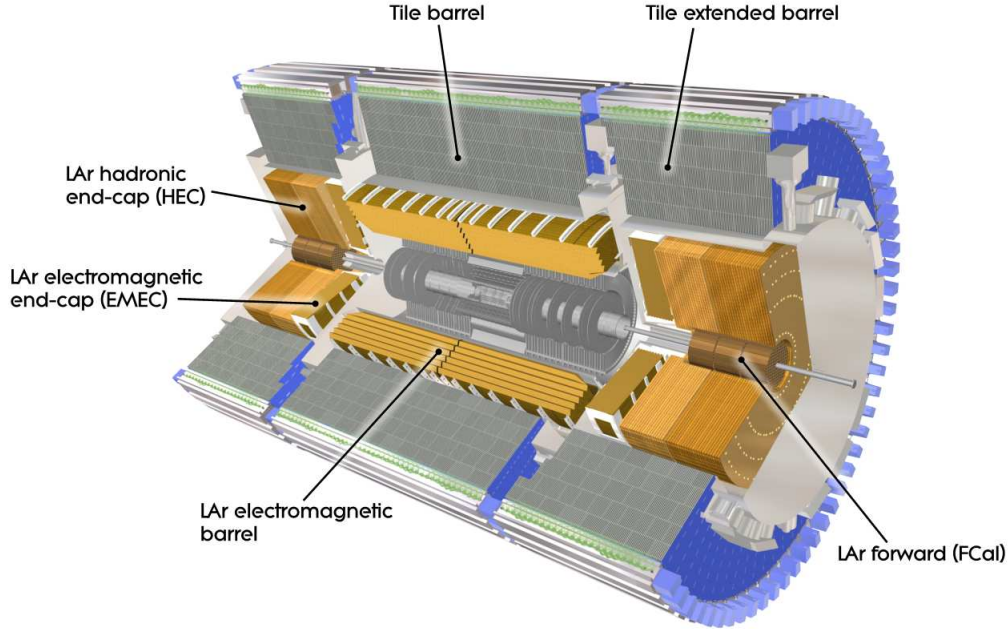


Figure 5.7: Cut-away view of the ATLAS calorimeter system [24].

	Barrel		End-cap	
EM calorimeter				
Number of layers and $ \eta $ coverage				
Presampler	1	$ \eta < 1.52$	1	$1.5 < \eta < 1.8$
Calorimeter	3	$ \eta < 1.35$	2	$1.375 < \eta < 1.5$
	2	$1.35 < \eta < 1.475$	3	$1.5 < \eta < 2.5$
			2	$2.5 < \eta < 3.2$
Granularity $\Delta\eta \times \Delta\phi$ versus $ \eta $				
Presampler	0.025×0.1	$ \eta < 1.52$	0.025×0.1	$1.5 < \eta < 1.8$
Calorimeter 1st layer	$0.025/8 \times 0.1$	$ \eta < 1.40$	0.050×0.1	$1.375 < \eta < 1.425$
	0.025×0.025	$1.40 < \eta < 1.475$	0.025×0.1	$1.425 < \eta < 1.5$
			$0.025/8 \times 0.1$	$1.5 < \eta < 1.8$
			$0.025/6 \times 0.1$	$1.8 < \eta < 2.0$
			$0.025/4 \times 0.1$	$2.0 < \eta < 2.4$
			0.025×0.1	$2.4 < \eta < 2.5$
			0.1×0.1	$2.5 < \eta < 3.2$
Calorimeter 2nd layer	0.025×0.025	$ \eta < 1.40$	0.050×0.025	$1.375 < \eta < 1.425$
	0.075×0.025	$1.40 < \eta < 1.475$	0.025×0.025	$1.425 < \eta < 2.5$
			0.1×0.1	$2.5 < \eta < 3.2$
Calorimeter 3rd layer	0.050×0.025	$ \eta < 1.35$	0.050×0.025	$1.5 < \eta < 2.5$
Number of readout channels				
Presampler	7808		1536 (both sides)	
Calorimeter	101760		62208 (both sides)	
LAr hadronic end-cap				
$ \eta $ coverage			$1.5 < \eta < 3.2$	
Number of layers			4	
Granularity $\Delta\eta \times \Delta\phi$			0.1×0.1	$1.5 < \eta < 2.5$
			0.2×0.2	$2.5 < \eta < 3.2$
Readout channels			5632 (both sides)	
LAr forward calorimeter				
$ \eta $ coverage			$3.1 < \eta < 4.9$	
Number of layers			3	
Granularity $\Delta x \times \Delta y$ (cm)			FCal1: 3.0×2.6	$3.15 < \eta < 4.30$
			FCal1: $\sim$ four times finer	$3.10 < \eta < 3.15,$ $4.30 < \eta < 4.83$
			FCal2: 3.3×4.2	$3.24 < \eta < 4.50$
			FCal2: $\sim$ four times finer	$3.20 < \eta < 3.24,$ $4.50 < \eta < 4.81$
			FCal3: 5.4×4.7	$3.32 < \eta < 4.60$
			FCal3: $\sim$ four times finer	$3.29 < \eta < 3.32,$ $4.60 < \eta < 4.75$
Readout channels			3524 (both sides)	
Scintillator tile calorimeter				
	Barrel		Extended barrel	
$ \eta $ coverage	$ \eta < 1.0$		$0.8 < \eta < 1.7$	
Number of layers	3		3	
Granularity $\Delta\eta \times \Delta\phi$	0.1×0.1		0.1×0.1	
	Last layer 0.2×0.1		0.2×0.1	
Readout channels	5760		4092 (both sides)	

Table 5.3: Main parameters of the ATLAS calorimeter system [24].

Both calorimeters use lead-liquid-argon layers for detection. The EM calorimeters have four layers in longitudinal direction. The first layer, called “presampler”, is used to recover for energy lost in dead material upstreams, like e.g. the solenoid magnet.

The next layer is finely granulated in η , therefore called η -strip layer. It is used for a very sensitive measurement of the fine structure of the electromagnetic shower shape and thus able not only to separate electrons and photons from hadronic showers but also to resolve the two close-by photons from a π^0 decay.

Then follows the EM-Middle layer which is the thickest part of the EM calorimeter with $\approx 16 X_0$ and offers a good granularity both in η and ϕ . The EM-Back layer completes the electromagnetic calorimeter system.

All EM layers together offer 24 units of radiation length X_0 in a geometrical depth of 47 cm. Figure 5.8 shows the radiation length of the calorimeter as function of the pseudorapidity.

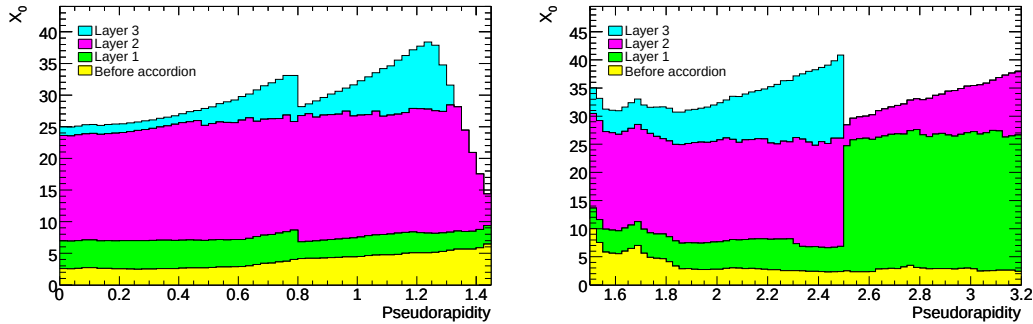


Figure 5.8: Cumulative amounts of material, in units of radiation length X_0 and as a function of $|\eta|$, in front of and in the electromagnetic calorimeters [24]. Shown is the thicknesses of each accordion layer as well as the amount of material in front of the accordion for the barrel (left) and endcap (right).

The EM calorimeter is built in an accordion geometry in the ϕ direction. This offers uniform detection of electromagnetic particles without any dead or less active calorimeter regions at the price of a more complicated design, slightly reduced resolution in ϕ and more involved shower reconstruction and energy calibration. A schematic drawing of a piece of the EM-Barrel is shown in Figure 5.9.

The EM calorimeters are sampling calorimeters with liquid argon as the active medium. The electrodes consist of capton plated copper plates which are segmented into strips that constitute the read out cells. The passive material is steel plated lead. The gap between electrodes and absorber plates is filled with liquid argon. Liquid argon was chosen because of its intrinsic radiation hardness and good energy resolution.

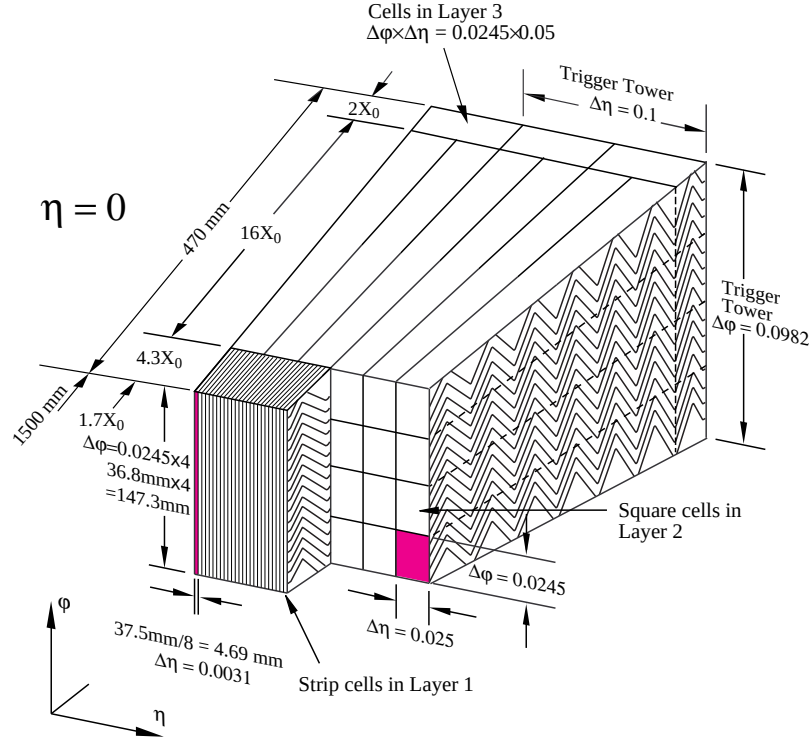


Figure 5.9: Sketch of a electromagnetic barrel module where the different layers are clearly visible with the ganging of electrodes in ϕ [24]. The granularity in η and ϕ of the cells of each of the three layers and of the trigger towers is also shown.

The aspired energy resolution is $\frac{\sigma_E}{E} = \frac{10\%}{\sqrt{E}} \oplus 0.7\%$ with a linearity of $< 0.5\%$ over the entire energy range. An angular resolution of 40 mrad requires a fine granularity that is given especially in the η -strip layer.

Hadronic calorimeter

Both in the barrel and in the endcap region the EM calorimeters are surrounded by hadronic calorimeters.

In the barrel part a steel-scintillator-sampling calorimeter (called TileCal) is used. It uses plastic scintillators with photo multiplier read out as the active medium and steel as passive medium. It consists of three parts, the barrel with a coverage of $|\eta| < 1.0$ and two extended barrels with a coverage of $0.8 < |\eta| < 1.7$. A schematic drawing of the Tile calorimeter system is shown in Figure 5.10. The read out elements follow a complicated structure to allow for a pseudoprojective arrangement and geometry of the individual calorimeter cells. The gap between the barrel and the extended barrel is filled with the electrical and mechanical

detector infrastructure of the inner detector and the EM calorimeter.

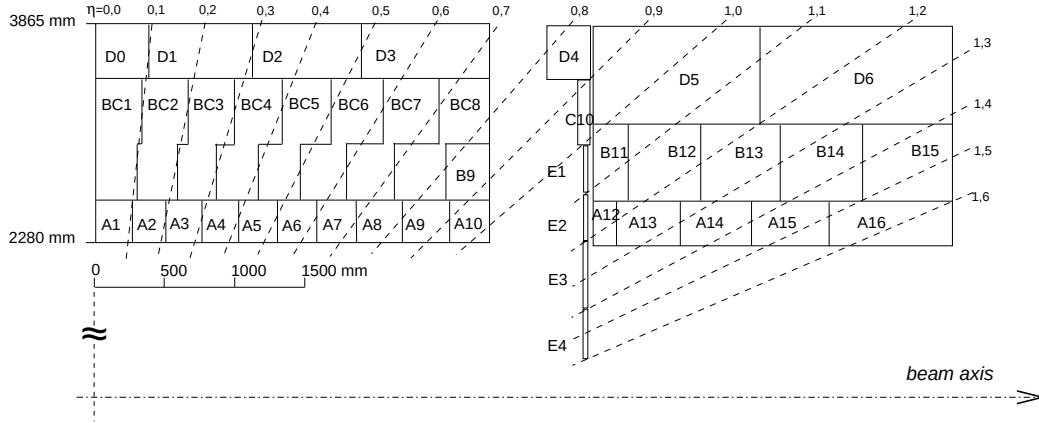


Figure 5.10: Schematic drawing of the segmentation in depth and η of the tile-calorimeter modules in the central (left) and extended (right) barrels. The bottom of the picture corresponds to the inner radius of the tile calorimeter. The tile calorimeter is symmetric about the interaction point at the origin [24].

The hadronic endcap calorimeter (HEC) is built mostly analogous to the EMEC, with the difference that copper is used as passive material and a planar geometry was chosen. It shares the cryostat with the EMEC and covers the range of $1.5 < |\eta| < 3.2$. In Figure 5.11 a schematic drawing of the HEC calorimeter is shown with η projective lines as guidance for the eye. Similar as with the Tile calorimeter, the read out cells are arranged in a pseudoprojective grid.

The design energy resolution is $\frac{\sigma_E}{E} = \frac{50\%}{\sqrt{E}} \oplus 3\%$ for both hadronic calorimeters.

Forward Calorimeter

An important contribution to the measurement of transverse missing energy gives the **forward calorimeter (FCAL)** at $3.1 < |\eta| < 4.9$. Due to its relatively fine segmentation, it also allows for the reconstruction of jets which is needed for physics signatures involving forward jets, e.g. Higgs boson production in Weak Boson Fusion. The FCAL is also constructed as a liquid argon sampling calorimeter and uses copper (for the first layer) and tungsten (for layer two and three) as absorber material.

Summary

The ATLAS calorimeter system provides at least 25 electromagnetic radiation lengths up to $|\eta| < 3.2$ (see Figure 5.8) and at least 10 hadronic interaction

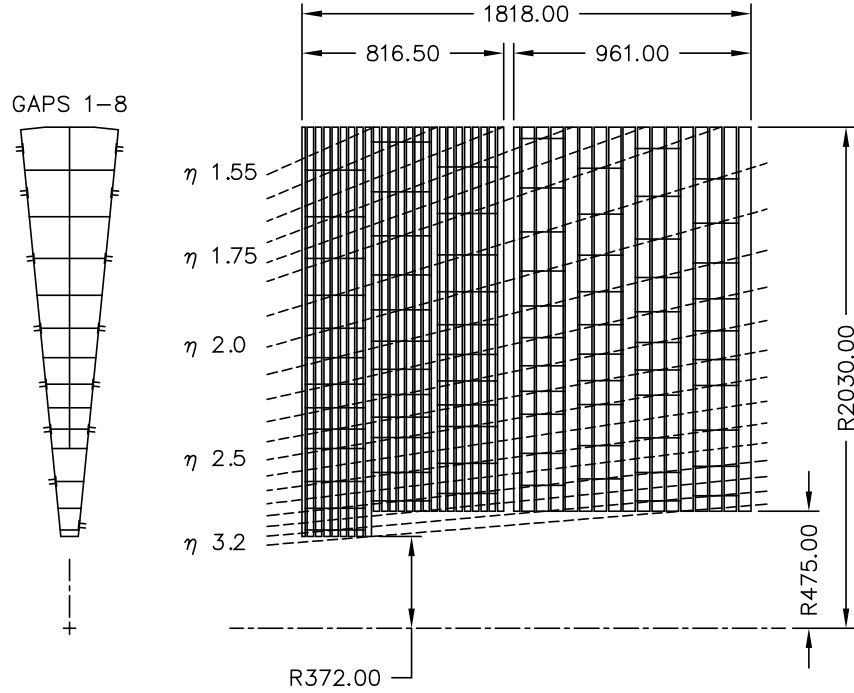


Figure 5.11: Schematic R - ϕ (left) and R - z (right) views of the hadronic end-cap calorimeter. The semi-pointing layout of the read out electrodes is indicated by the dashed lines. Dimensions are in mm. [24].

lengths up to $|\eta| < 4.9$. The profile of the absorption length vs. $|\eta|$ for all calorimeter components is shown in Figure 5.12.

This design allows for very good detection of electromagnetic particles in the most relevant central region of the calorimeter and hadronic activity up to very small angles to the beam pipe, which provides for an hermetic measurement of the missing transverse momentum.

5.2.3 Muon spectrometer

Final states with high energetic muons are among the most important signatures for interesting processes at the LHC. Therefore, an independent muon detector system was built, which surrounds the other ATLAS sub-detectors hermetically. It consists of a magnet system which produces a toroidal magnetic field of 0.3 Tesla over a large volume, using air-core super conducting magnets and several tracking elements. Figure 5.13 shows a schematic drawing of the muon system.

The magnetic field forces muons on a helix in the r - z -plane, where hits for

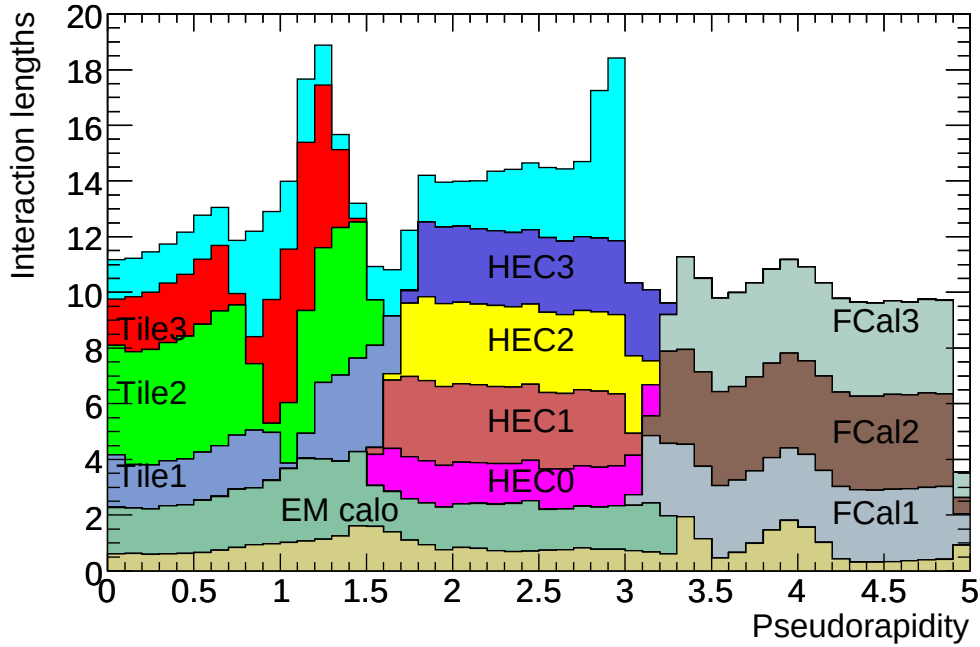


Figure 5.12: Cumulative amount of material, in units of hadronic interaction length λ , as a function of $|\eta|$, in front of the electromagnetic calorimeters, in the electromagnetic calorimeters themselves, in each hadronic compartment, and the total amount at the end of the active calorimetry. Also shown for completeness is the total amount of material in front of the first active layer of the muon spectrometer (up to $|\eta| < 3.0$) [24].

track reconstruction are measured by two different systems. These are drift tubes (MDT) in the barrel and cathode strip chambers (CSC) in the forward region. Both provide for a total momentum resolution of $\Delta p/p < 10\%$ at $p = 1$ TeV.

Since few particles, besides muons, penetrate the very thick calorimeter system, muons are also clean signatures to trigger on. In order to provide a fast muon trigger the muon spectrometer is equipped, in addition to the precision chambers, with resistive plate chambers (RPC) and thin gap chambers (TGC) which provide a less precise but very fast measurement of hits.

5.2.4 Trigger system

For a hadron collider experiment a potent trigger and data acquisition system (DAQ) is indispensable. The trigger has to filter interesting events out of an overwhelming background to reduce the rate by factor of $\sim 400,000$ and the

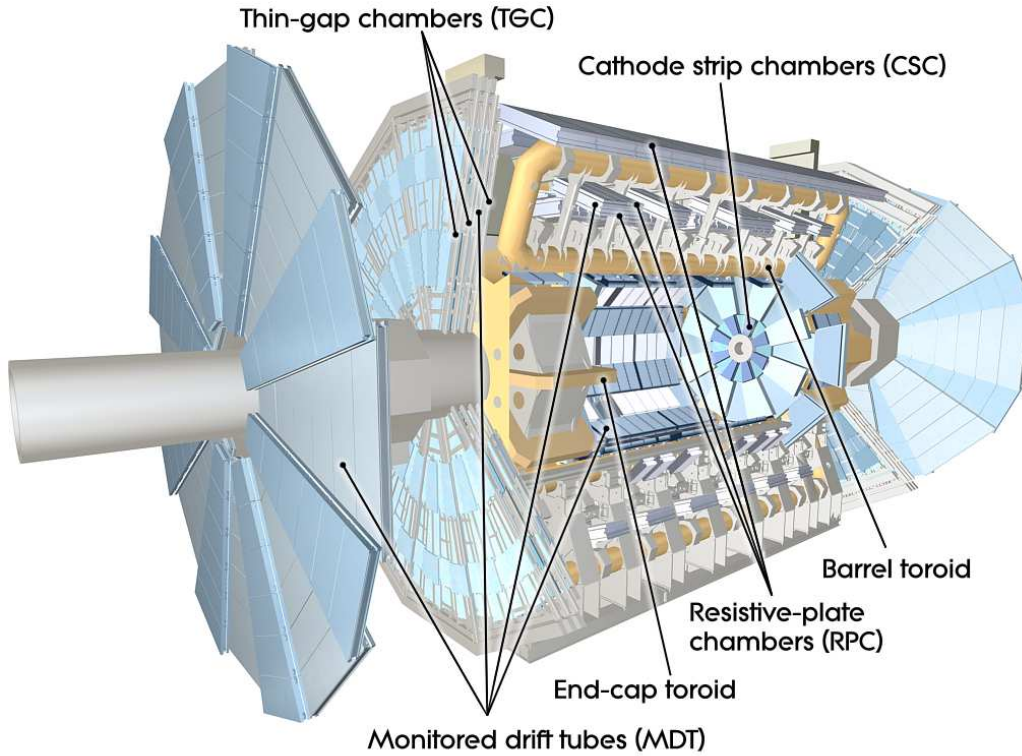


Figure 5.13: Cut-away view of the ATLAS muon system [24].

DAQ has to distribute and store the very large amount of collected data. A schematic representation of the ATLAS Trigger and data acquisition system is shown in Figure 5.14.

To achieve a data reduction of $\sim 400,000$, the trigger system of ATLAS is composed of three levels which have access to different levels of detector information, as is shown in the schematic overview in Figure 5.15.

Level One trigger

The Level One Trigger is based on custom built electronics to provide a latency of less than $2 \mu\text{s}$. It uses information from the calorimeters and the muon system only. The calorimeter information is read out in so called trigger towers which have a separate electronics path and correspond to a reduced granularity of $\Delta\eta \times \Delta\phi = 0.1 \times 0.1$. On this granularity Level One objects are built, which are electromagnetic clusters, τ -clusters, jets and missing transverse energy. The muon system provides muon trigger candidates. Based on these objects, with certain energy/momentum threshold, a trigger signal is then issued. The area covered by these objects also serves as a region of interest (ROI) for the Level Two Trigger. A maximum of 100 kHz is accepted by Level One and handed over to the Level Two Trigger.

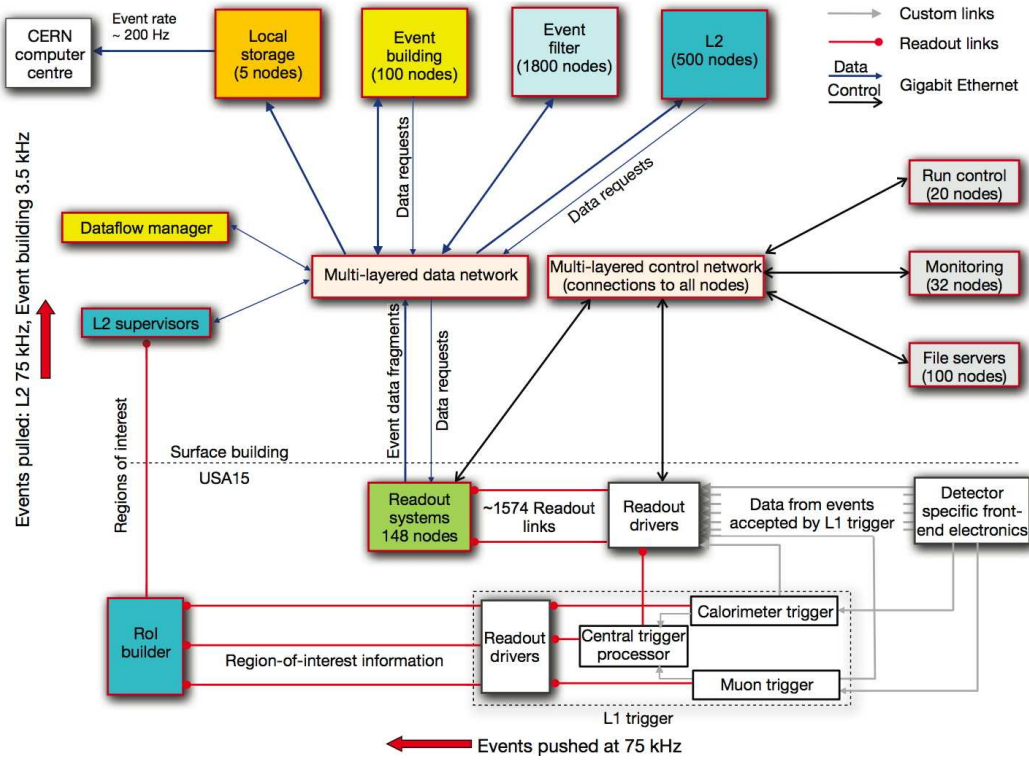


Figure 5.14: Block diagram of the ATLAS trigger and data acquisition systems [24].

Level Two trigger

The Level Two Trigger has access to the full calorimeter granularity, but only in the ROIs given by the Level One Trigger. In addition, tracks from the inner detector become available at this stage of the trigger chain. The higher granularity, track information and better reconstructed objects are then used to refine the decision of the Level One Trigger and to reduce the rate to ~ 1 kHz.

Event Filter

Events which are accepted by Level Two are reprocessed completely by the Event Filter with the full detector information available from all solid angles (not only from the ROIs). If such fully reconstructed events are accepted by the last trigger stage, they are written to storage with ~ 100 Hz. This corresponds to a data rate of 10-100 MB/s.

One copy of the raw data of accepted events is then stored in the CERN computer centre and a second copy directly distributed to one of the 10 ATLAS Tier One grid computing centres worldwide.

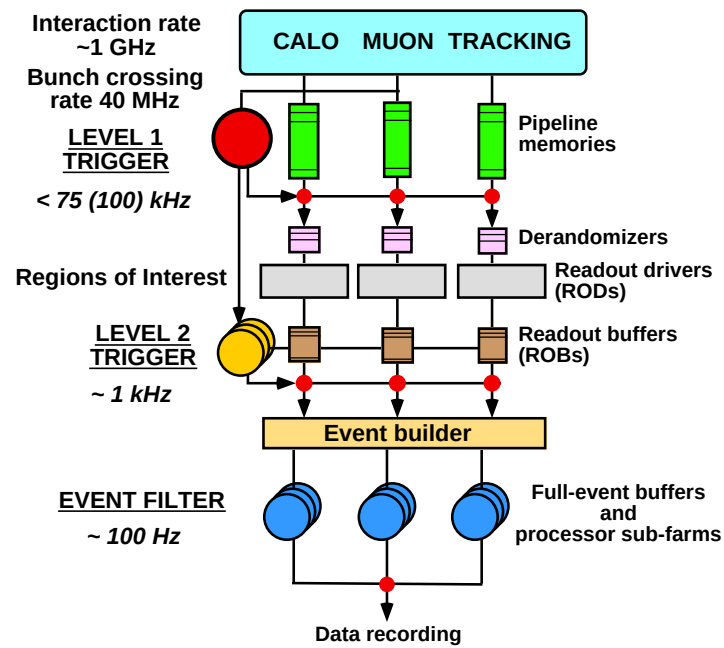


Figure 5.15: Schematic diagram of the ATLAS trigger system.

6 Development of a fast calorimeter simulation for the ATLAS detector

An understanding of events recorded with the detector and the physics processes that most likely produced these events is only possible at the hadron level. Apart from a few very clear final states, the recorded events contain hundreds of different final state objects. With a perfect detector, one could hope to identify every single of these objects and thus reconstruct the event at the stable hadron level. Once this is done, one could try to go back the chain to unstable resonances and finally to a reconstruction of the parton level event and an interpretation in terms of a physics process.

Despite the intricate and complex design of current particle physics detectors, the detectors and detector reconstruction algorithms are not able to

- reconstruct all objects.
- identify the objects unambiguously.
- measure the identified objects sufficiently precise to reconstruct all unstable resonances.

Thus going back to the parton level is almost impossible on a per event basis and even on a statistical basis it is a very complicated step that involves a deep understanding of the detector and the physics processes involved. So instead of going back to the parton level, the procedure is reversed and parton level events are generated and simulated all the way down to the detector level.

The basic aim of any detector simulation in particle physics is to allow a comparison between the data measured in the real detector and the expected measurements produced by known (and some unknown, but desired) physical phenomena. Due to imprecision of the detector in all three points mentioned above, it is unfortunately not sufficient to simulate only the events of interest, but one needs also a huge number of additional background events that might be falsely seen as signal events. Furthermore, during the construction phase of the detector, the simulation is also crucial to evaluate the impact of the detector design and to prepare software and physics analyses for the data taking phase. A large number of simulated events is also needed for this purpose.

The combination of a need for a large number of simulated events and the time consuming simulation of a complicated detector system makes it necessary to search for acceptable optimisations in the simulation. Thus, in the following chapter a fast simulation method for the ATLAS calorimeter system is introduced.

6.1 Introduction

In the case of the ATLAS calorimeter, the measured quantities in the detector are energies deposited inside calorimeter cells. These measured energies are used during the reconstruction step to calculate the total energy and particle identification variables from the relative distribution of energies in the cells (the “shape” of the energy distribution) for an object in the calorimeter. Once a particle enters the calorimeter, it produces a hadronic and/or electromagnetic shower that leads to energy depositions in several calorimeter cells. Even for identical particles these showers will always follow an individual history, leaving no two showers ever identical. Furthermore, the calorimeter technology will give an intrinsic uncertainty on any measurement of the deposited energy. Last, the read-out system will add electronic noise to the measurements.

Due to this intrinsic randomness there is no sense in comparing one measured event to one simulated event. Even if the underlying physics process and kinematic properties were completely known for the simulation for this particular event¹, the measured event in the detector and the simulated event would never be identical. A comparison is only sensible on a statistical basis, comparing significant distributions for certain particle types. Producing and analysing such significant distributions is an essential part of the detector reconstruction step.

Therefore any measure of quality of a simulation is based on reconstructed variables that were produced with the simulation as input for reconstruction. Any differences that do not show up after reconstruction will most likely not be found (and are most likely also unimportant).

There are several possible approaches to the detector simulation to reach a good agreement after reconstruction, which will be briefly discussed now.

6.1.1 The ATLAS full simulation

The most detailed approach is to simulate the propagation and interactions of all particles through a detailed computer model of the detector, including a simulation of the energy measurement process and the electronics used for read-out. For ATLAS, this is implemented using a Geant 4 simulation model [25–27], which will

¹The input to the detector simulation are all particles of some primary physics event simulation that are sufficiently stable to reach the active elements of the detector (hadron level). For a complete comparison of a single event, this hadron level input would need to be identical to the real event.

be called “full simulation” in the following. Provided that the computer model of the detector and the physics model of particle interactions in the detector material is sufficiently accurate, such a simulation should give very good results. The drawback is that the simulation time for an average ATLAS event is in the order of several minutes, most of which is spent for the calorimeter simulation.

6.1.2 The ATLAS fast simulation *Atlfast*

The simulation package *Atlfast* [28] was developed as a very fast alternative to the full ATLAS detector simulation and event reconstruction and follows a very different strategy. Instead of trying to simulate the detector, it tries to simulate the final physics objects after all simulation, reconstruction and identification steps. This is done using the truth information provided by the primary physics simulation. This truth information is smeared using parametrisations of resolution functions and is corrected for reconstruction and identification efficiencies.

The drawback of this fast approach is that efficiencies and rejection factors can show a strong dependency on the event type² that can not be easily implemented into a parametrisation. Also any analysis that depends on underlying information from the reconstruction (e.g. E/p of an electron, some tau discriminant variables, etc.) cannot be done, since this underlying information is not defined within the *Atlfast* simulation approach.

Since almost no simulation and reconstruction is done with *Atlfast*, each event takes less than a tenth of a second to be produced.

6.2 Improving the fast simulation

The main deficiency of *Atlfast* is that the reconstruction and identification step is not done by some reconstruction algorithm but emulated by the use of true information of the primary event generator without the use of simulated detector information. To improve this for the calorimeter, a more detailed simulation and reconstruction code using this simulated calorimeter is needed.

The aim of *FastCaloSim* is to provide a sufficiently good calorimeter simulation to run the ATLAS standard reconstruction and identification code that is also used on full simulation and will be used for data taking. “Sufficiently good” means that key features of the reconstructed properties are similar to the full simulation. Differences in details are assumed to be independent of the overall event activity and can therefore be absorbed in a parametrisation of residual differences.

For the simulation of the inner detector and the muon system three simulation options exist:

²For example, the electron or τ efficiencies and rejections in “quiet” events like $q\bar{q} \rightarrow Z$ can be very different to efficiencies and rejections in “busy” events like $gg \rightarrow t\bar{t}$.

1. The inner detector is simulated with the full simulation and the detailed parametrisation of Atlfast is used for muons. Since the muon system is to a large degree decoupled from the rest of the detector and the overall event activity, doing this is a good first approximation.
2. It is possible to combine the full simulation of the inner detector with a full simulation of muons in the calorimeter and muon system with a small increase in overall CPU time due to the muons in the outer detector systems.
3. The full simulation of the inner detector is replaced by the fast inner detector simulation FATRAS [29], that does also a fast simulation of muons in the calorimeter and muon system.

In the following chapter, only option 2 is considered to avoid unwanted side effects from another fast simulation.

6.2.1 Simplifications for FastCaloSim

It is clear that any significant improvement in speed must stem from simplifications in the simulation model. The design baseline was to make the simulation algorithm as fast as possible in all areas that cost a large amount of time.

- As described in Section 6.1, particle showers have fluctuations that make every shower different. Since any reconstruction algorithm should work with high efficiency on the particle type it is aimed for, such an algorithm is most sensitive to particle shower properties that are shared by most shower histories and (hopefully) independent from shower fluctuations. It is a different issue with misidentified objects, which are often caused by these fluctuations and thus sensitive to all effects that can cause fluctuations in the shower shape.

The simulation method chosen in FastCaloSim is to try to simulate the longitudinal shower properties, including fluctuations and correlations, but only average lateral shower properties and uncorrelated lateral energy fluctuations (fluctuations due to the calorimeter cell energy resolution³ and electronic noise). Lateral energy fluctuations from the individual shower development are ignored. This allows to deposit a particle in one simulation step into the calorimeter without tracking the particle propagation, interactions and shower development as in full simulation. From this, one would expect that particle identification efficiencies are much better described than misidentification rates.

³In case of the ATLAS LAr calorimeter, the intrinsic energy resolution of the technology is $\approx 10\%/\sqrt{E[\text{GeV}]}$.

The particle shower is described by two parametrisations, one providing the total energy deposited in each calorimeter layer (see Section 6.3) and the other describing the energy distribution within the specific calorimeter layer (see Section 6.4).

- Only two types of particles are parametrized and used for simulation: photons and charged pions. The photon parametrisation is used for photons and electrons. The pion parametrisation is used for all hadrons (neutral and charged).

Inside the calorimeter, the main difference between electrons and photons is that electrons tend to start the electromagnetic shower earlier. This gives a small difference in the reconstructed energy in the calorimeter and slightly broader shower shapes. This difference is visible but not crucial for the reconstruction quality and can be corrected by a later parametrisation or correction step.

Most of the hadronic activity in QCD events is dominated by charged pions, so the simplification of using the pion parametrisation for all hadrons is sufficient. Since hadrons are almost never isolated, the individual shower properties cannot be easily resolved in the reconstruction and a lot of the individual shower development of the different hadrons remains hidden⁴.

In principle, dedicated electron and hadron specific parametrisations could be created, but this would need a lot of additional parametrisation effort and memory during simulation. One important result from the detailed validation of FastCaloSim within the ATLAS collaboration [30,31] is that this effort seems justified in the case of electrons. Therefore a dedicated electron parametrisation will be included in the updated version of FastCaloSim.

- The simulation uses the reconstruction geometry of the calorimeter that describes calorimeter cells as cuboids in the r/z , η , ϕ space (the FCal cells are cuboids in x,y,z). This is an accurate assumption for the homogeneous regions of the EM calorimeter (but not for the hadronic Tile and HEC calorimeter or cells close to the calorimeter boundaries, where the real detector geometry is more complicated, see Section 5.2.2). Figure 6.1 shows the calorimeter reconstruction geometry in the r - z projection. For the purpose of the fast simulation, this simplified geometry is absolutely sufficient, although some particle shape distortions due to the simplified geometry close to the calorimeter cracks can be expected.

Parametrisations are done within the continuous calorimeter regions away from calorimeter boundaries and do not follow all the fine structure. Espe-

⁴This might cause differences in special cases like $\tau^\pm \rightarrow K^\pm + X$ decays, where the heavier hadron is sometimes visible as isolated particle. Another case, although hard to isolate, would be neutrons that are also treated like charged pions.

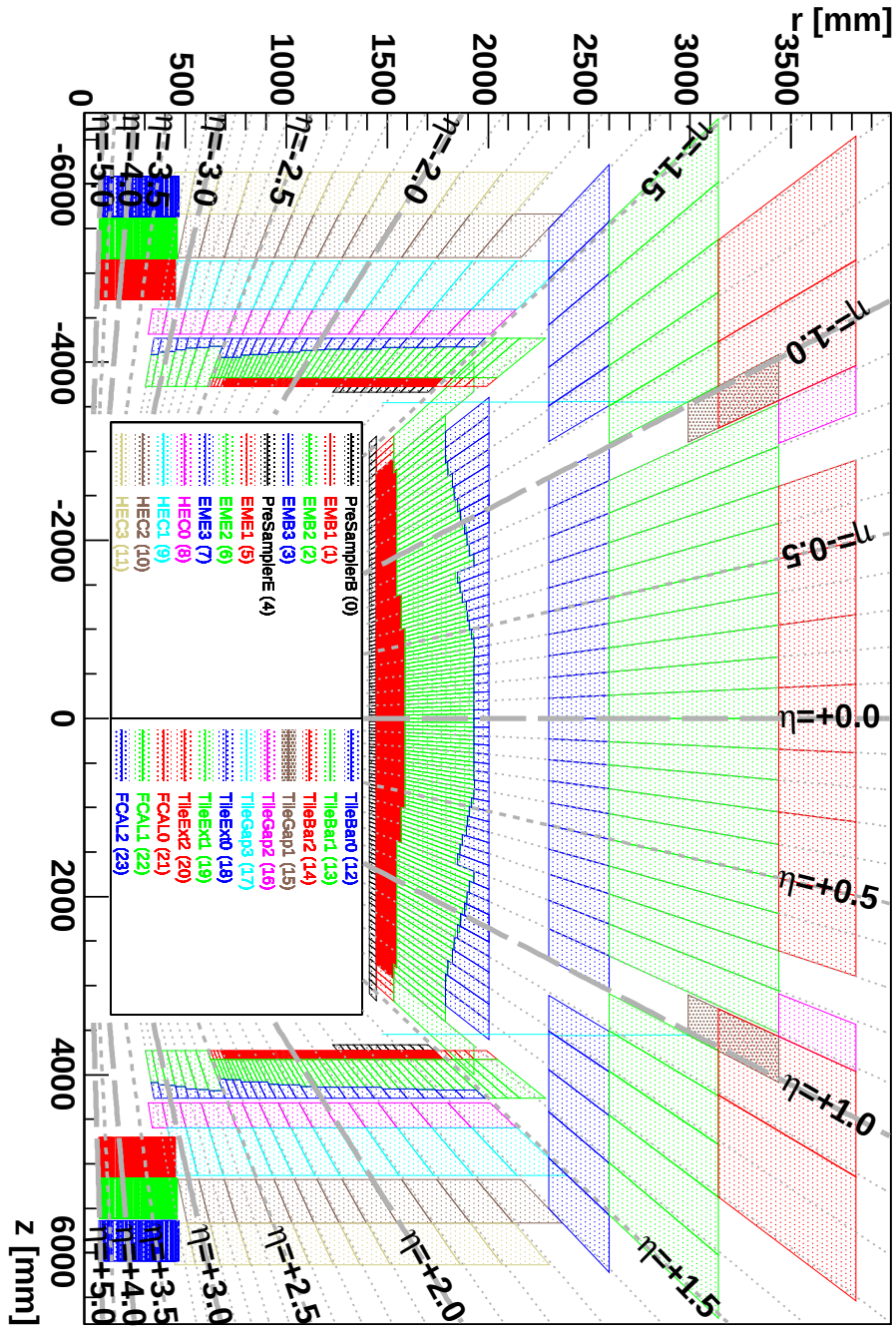


Figure 6.1: Visualisation of the reconstruction geometry of the calorimeter system in the r - z plane. The individual layers of the ATLAS calorimeter are colour coded and listed in the legend.

cially close to the boundaries, differences to the full simulation are anticipated.

- Muons are currently not simulated (see the discussion in the introduction to Section 6.2) and therefore no energy is deposited in the calorimeter by FastCaloSim itself. Only if FastCaloSim is run together with the FATRAS fast muon system simulation, the muon energy loss in the calorimeter simulated by FATRAS is passed on to FastCaloSim and deposited along the direction of flight of the muon in the calorimeter. If full simulated muons are used, the muon energy loss in the calorimeter is already included in the full simulation.

6.2.2 Input to the parametrisation

The parametrisations are based on full simulated single particle events created with Athena version 10.0.1. The generated particle energies and pseudorapidities are listed in Table 6.1. For each combination of particle type, energy and pseudorapidity 10,000 events were generated with the full simulation (in total $\approx 10\text{M}$ photons and $\approx 11\text{M}$ pions).

Particle	Constant energies [GeV]														
Photon	0.2	0.4	0.6	0.8	1	2	5	10	20	50	100	200	500		
Pion			0.6	0.8	1	2	5	7.5	10	15	20	50	100	200	500

	Pseudorapidity	
η -Range	0.05 - 3.35	3.50 - 4.90
Bins	66	8
Bin η -size	0.05	0.05 (gap of 0.15 to the next bin)

Table 6.1: Generated particle energies and pseudorapidities used as input for the FastCaloSim parametrisation. Particles were generated flat in ϕ in the range $[-\pi, +\pi]$. The pion raw energy response shows some small non-linear behaviour around $E=10$ GeV, so two extra energy points were added. In the FCal region ($\eta > 3.2$) cells have an average size in η of $\Delta\eta \approx 0.2$. To save CPU time, less bins were produced there.

The particles are generated at the centre of the detector with a displacement corresponding to the expected beam spread of the ATLAS interaction region. Only particles without interaction in the inner detector (unconverted photons and non-interacting pions) are used for the parametrisation. These interaction effects are added separately by the inner detector simulation and should not enter into the calorimeter parametrisation.

The Geant 4 simulation and detector geometry as of Athena version 10 is outdated by now, but the necessary size of the full simulation input samples and the time to produce a parametrisation do not allow frequent updates. An update of the parametrisation is done with a close to final Geant 4 simulation version and the settings and geometry used for the first data taking period, but using the same method as for the parametrisation described here.

6.3 Energy parametrisation

The first parametrisation step is to describe the total energy deposition of a particle and how this energy is distributed between the various calorimeter layers. The layers of the ATLAS calorimeter system are shown in Figure 6.1 together with the names that are used to identify the individual calorimeter layers.

FastCaloSim is designed to provide a parametrisation of the raw calorimeter energies as seen during the reconstruction step of a full simulation. After calorimeter object identification, the energy of the objects is calibrated using object type dependent calorimeter calibration factors. Such a calibration can only work with the raw calorimeter energies of the fast simulation, if the most important correlations between the energies in the various layers are preserved.

6.3.1 Optimal calibration

The optimal calibration is defined as the calibration function that gives as average calorimeter response to a specific particle an energy distribution peaking at the true particle energy with a minimal spread. The calibration needed to achieve this goal can be different for different physics objects, and all reconstruction algorithms will try to work with a specific calibration that is as close as possible to optimal for the intended object. The calorimeter raw response that is being parametrised for FastCaloSim can in principle be very different from this optimal calibration. During the reconstruction step (which is also applied on the events simulated by FastCaloSim) calibrations are applied to this raw response to reconstruct the “true” energy.

Preserving all correlations in a direct simulation of the raw energy response in such a way that a later calibration step during reconstruction works for all objects with FastCaloSim as with the full simulation is very challenging. The correlations are preserved much easier if the parametrisation is based on calorimeter energies after (optimal) object calibration for the considered particle. In this way, the most important correlations that are exploited in a calibration step during reconstruction are already taken into account. The raw energy response can then be calculated from the simulation of the calibrated energy response by applying the inverse calibration function as last step in the FastCaloSim energy simulation.

A sampling based calibration without constant offset

$$E_{\text{cal.}} = \sum_{i \in \text{layers}} c_i \cdot E_i \quad (6.1)$$

is used for the energy parametrisation. For every particle type, energy and η -position a different set of calorimeter calibration constants c_i is fitted for the energies E_i in the calorimeter layers. Figure 6.2 shows the calorimeter response to photons and charged pions before and after calibration.

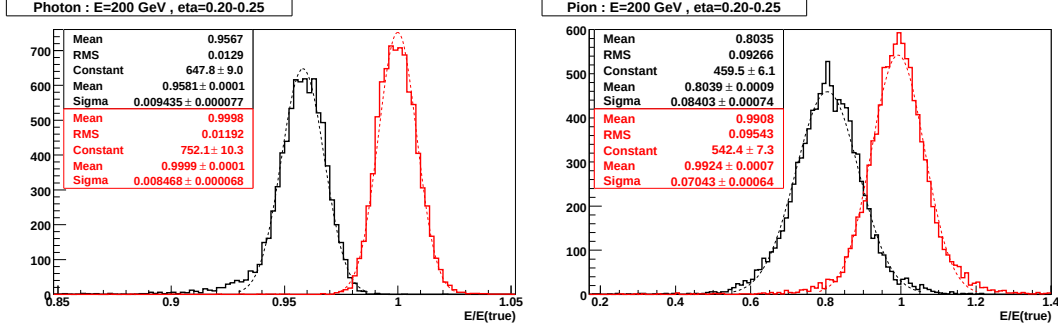


Figure 6.2: Total energy response of the calorimeter to photons (left) and charged pions (right) of 200 GeV energy in the range $0.20 < \eta < 0.25$. The black distribution is the raw response of the calorimeter, the red distribution the response after calibration. Overlaid is a Gaussian fit within 2σ of the mean as thin dashed line.

Figures 6.3–6.6 show the fitted calibration constants for photons and pions as function of the particle energy and pseudorapidity. The calibration constants have a smooth energy dependency and follow in η the distribution and thickness of the active calorimeter material. Close to or beyond the edges of calorimeter layers some calibration constants show a non-continuous behaviour. This is caused by high calibration factors on small energy depositions from tails of the shower development. The fluctuations observed for photons in the hadronic calorimeter layers are also due to uncertainties of this kind.

6.3.2 Reconstructed shower depth

One of the simplifications introduced in Section 6.2.1 was to describe showers only by their mean shape without taking into account fluctuations. Depending on the longitudinal depth in the calorimeter at which a shower starts, the mean shape measured inside one calorimeter layer can be very different. Since the point of the first interaction in the calorimeter was not recorded for the input sample, the reconstructed longitudinal shower centre was taken to model this key

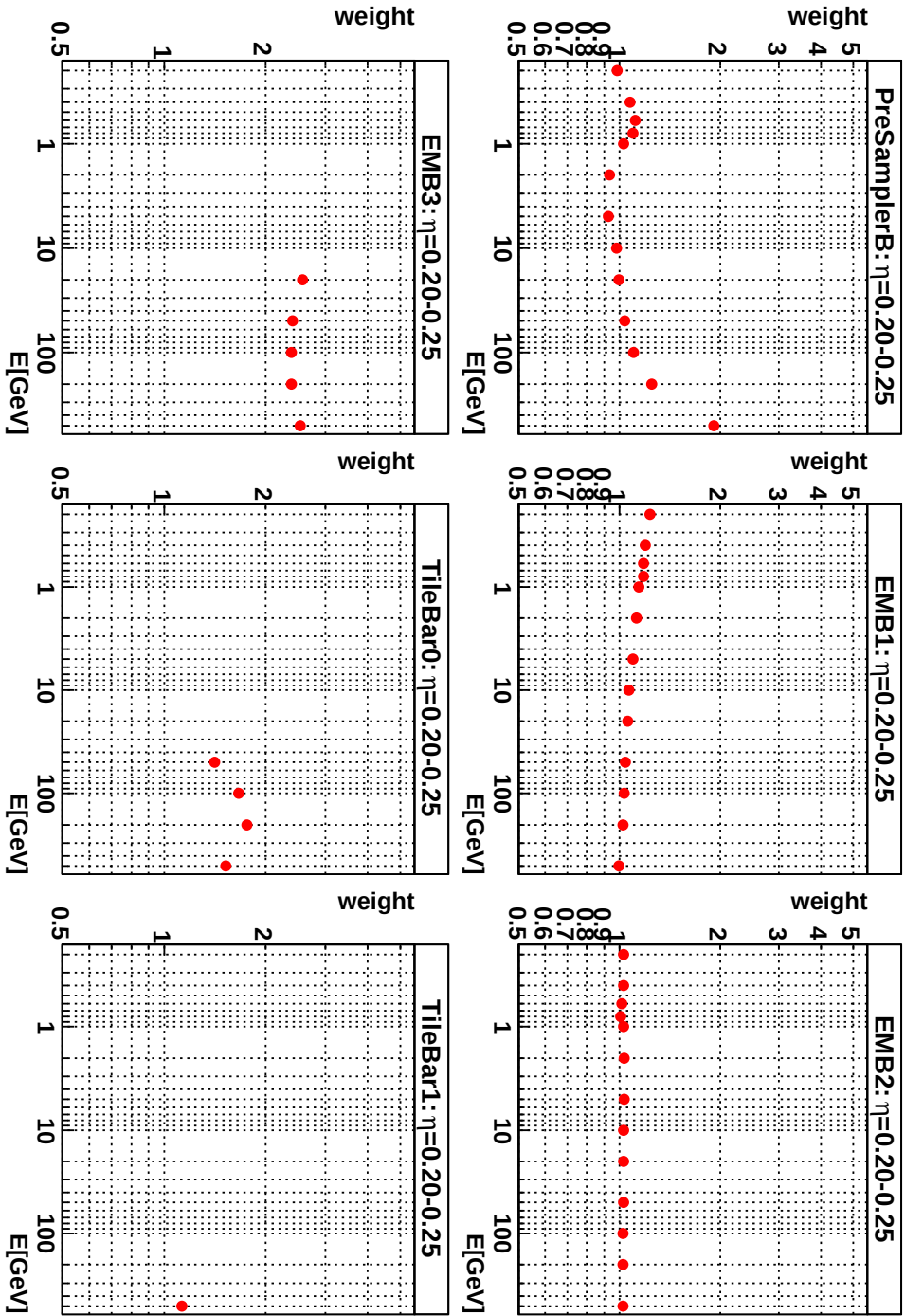


Figure 6.3: Fitted calorimeter calibration constants as function of the particle energy for photons in the range $0.20 < \eta < 0.25$.

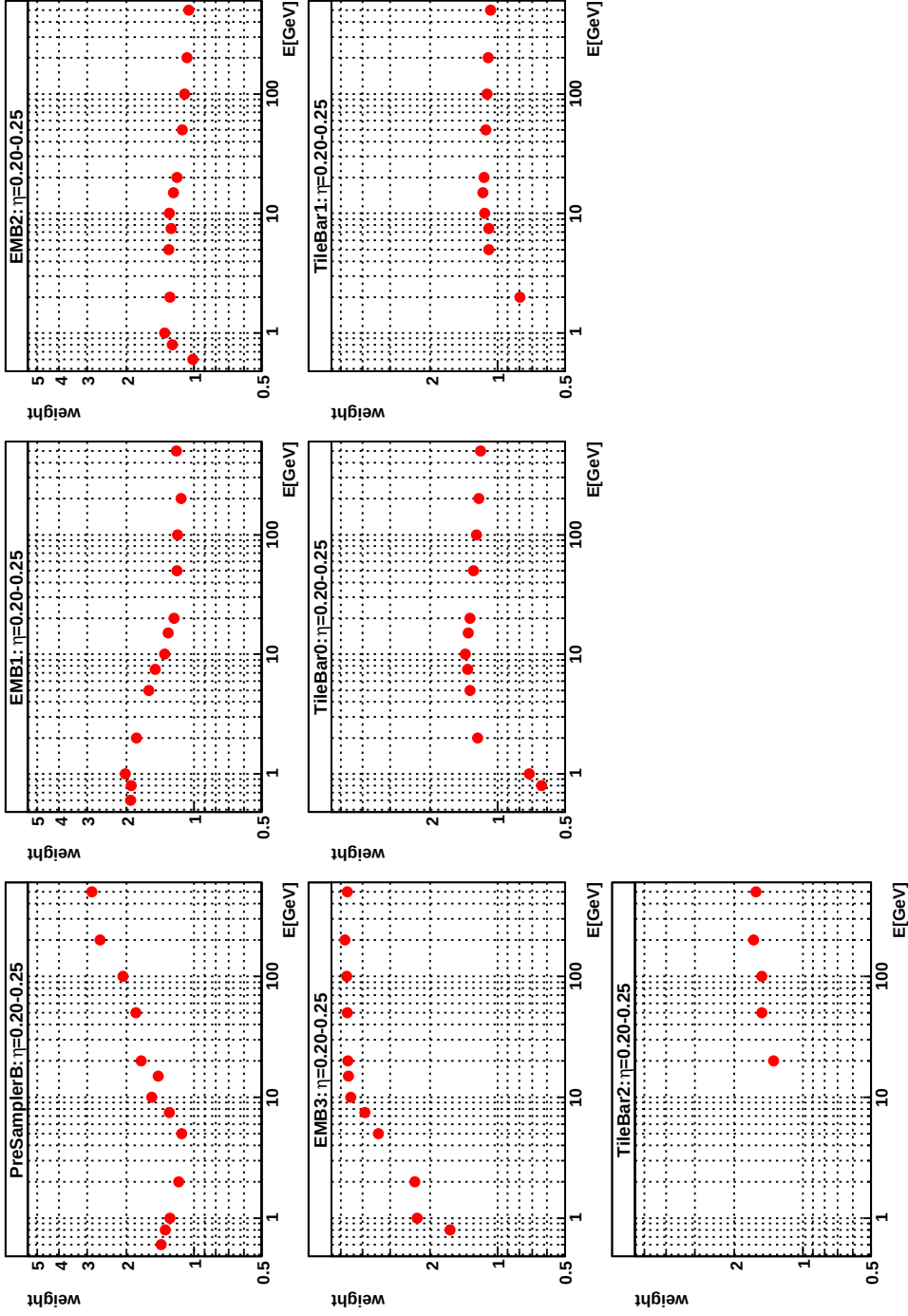


Figure 6.4: The same as Figure 6.3 for pions.

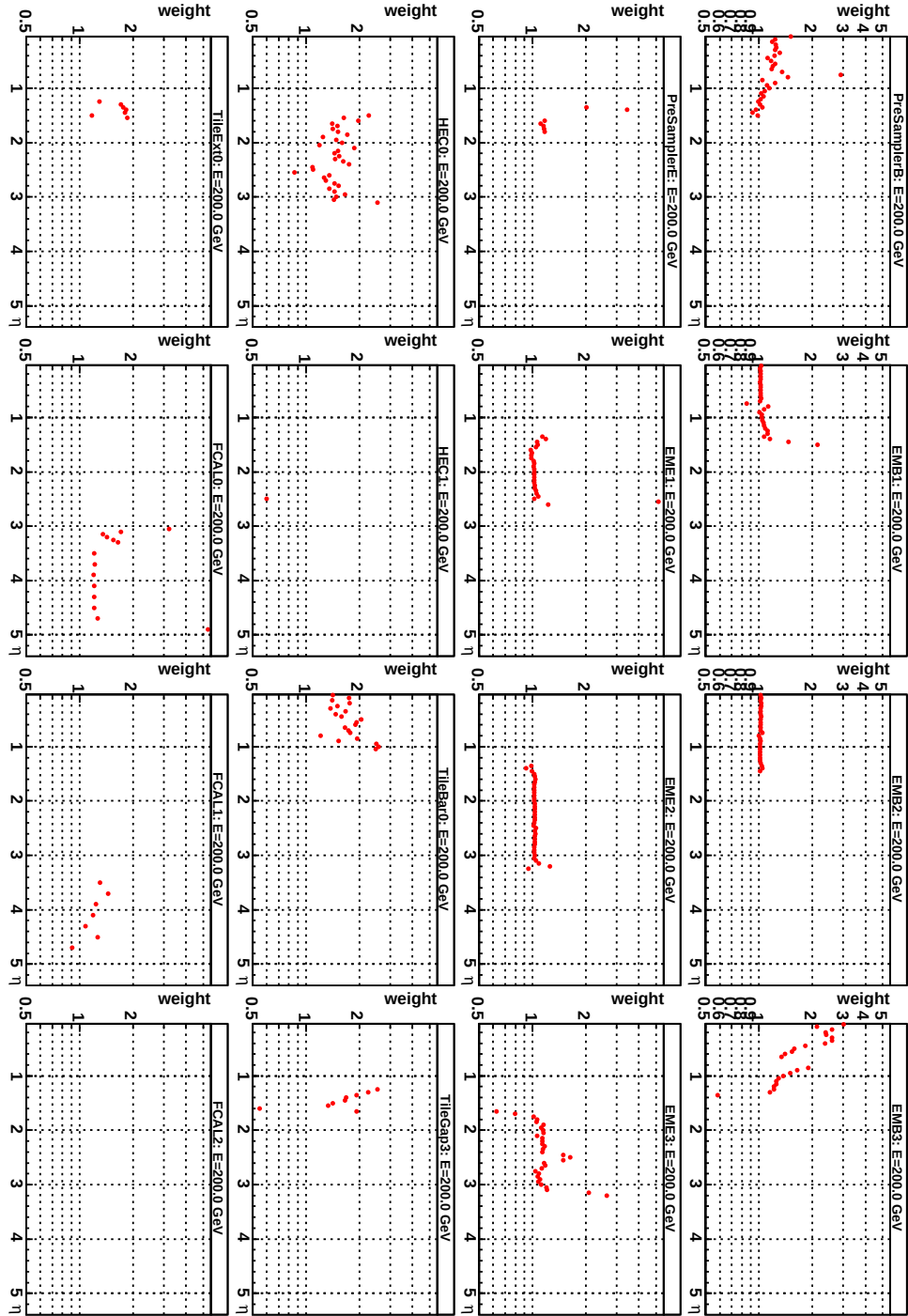


Figure 6.5: Fitted calorimeter calibration constants as function of the pseudorapidity for photons of 200 GeV energy.

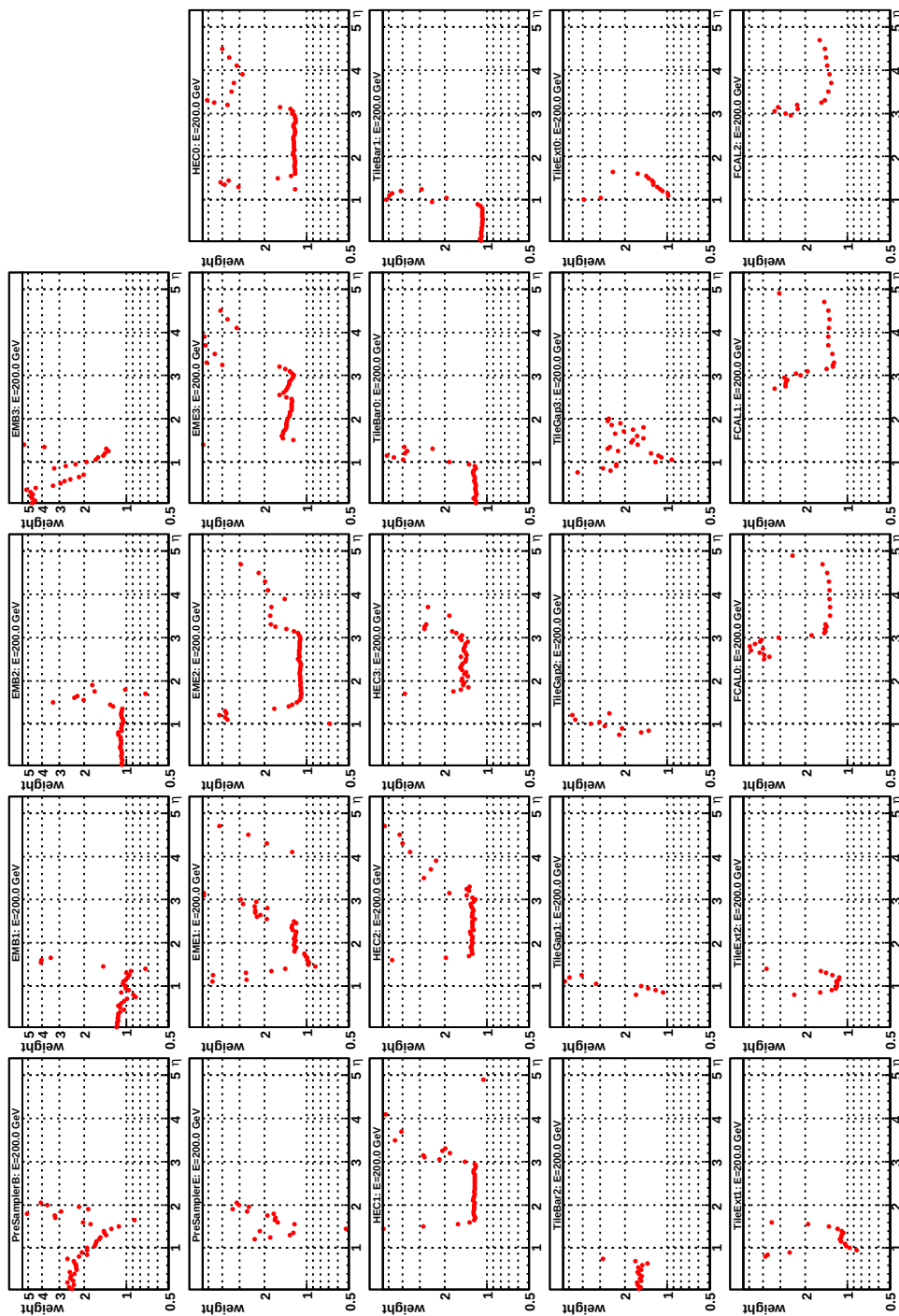


Figure 6.6: The same as Figure 6.5 for pions.

dependency of the shower development. This has the additional advantage of being a quantity that could be measured in the detector and thus allows to produce a parametrisation from measured data in the future. Furthermore, shower fluctuations that change the longitudinal shower development, get at least partly absorbed in the dependency of the longitudinal shower centre.

The reconstructed longitudinal shower centre is calculated from the energy weighted mean of the longitudinal centre positions d_{layer} of all calorimeter layers:

$$d = \frac{1}{E} \cdot \sum_{i \in \text{layers}} E_i \cdot d_i \quad . \quad (6.2)$$

In the following this distance d is called “shower depth”.

Figures 6.7 and 6.8 show on the left the shower depth for photons and pions. For the parametrisation and simulation, this distribution is subdivided in 10 bins of approximately equal number of events per bin (red histogram). Extreme events that have a very early or late shower (the first and last 1% of events in the histogram) are not included in the parametrisation. All following parametrisation steps are performed independently for each of these bins in the shower depth. On the right, the correlation between the calibrated energy and the shower depth is shown.

For photons (Figure 6.7, right side), which deposit their energy completely within the homogeneous LAr EM calorimeter, no dependency of the calibrated energy on the shower depth is visible. Pions (Figure 6.8, right side) show a small dip in the energy response, if part of the shower energy is lost in the cryostat region between EM and hadronic calorimeter. Also an improved energy resolution for pions which deposit the dominant part of their energy only in the EM calorimeter (first bin) is visible.

The simulation takes 2D histograms as in Figures 6.7 and 6.8 on the right side as input PDFs of the energy parametrisation and thus chooses pairs of reconstructed calibrated energy and shower depth with the correct correlation preserved. For intermediate points in energy and pseudorapidity between those in Table 6.1 the best matching parametrisation point is selected.

The fraction of the total energy deposited in the individual calorimeter layers is shown in Figures 6.9 and 6.10 as function of the shower depth. For photons a clear anti-correlation in the shower depth between EMB1 (strip layer) and EMB2 (EM bulk) is visible, all other layers contain only small amounts of energy. Pions show an anti-correlation between the total energy in the EM and Tile calorimeter.

6.3.3 Correlated energy fractions

If the simulation uses histograms as in Figures 6.9 and 6.10 as input (together with the already determined energy and shower depth), it is able to reproduce these distributions and also the dependency of the energy fractions on the shower

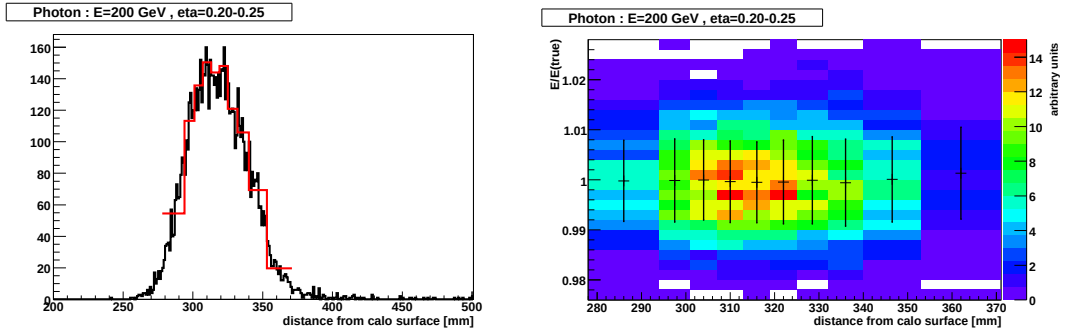


Figure 6.7: **Left:** The black histograms show the reconstructed longitudinal distance d of the shower centre from the calorimeter surface (shower depth) for photons of 200 GeV energy in the range $0.20 < \eta < 0.25$. In red, a rebinned version of this histogram that is used for the simulation is shown (see text).

Right: Correlation of the shower depth and the total calibrated energy response. The black bars indicate the mean and sigma of a Gaussian fit in each bin of the shower depth.

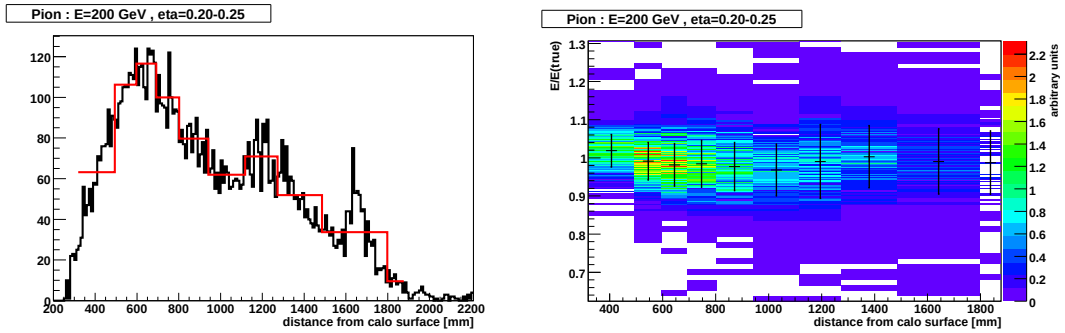


Figure 6.8: The same as Figure 6.7 for pions.

depth. However, the energy fractions deposited in each calorimeter layer should add up to the total energy. This is not ensured by drawing uncorrelated random numbers from the individual histograms. Furthermore, after dicing the energy fractions, the shower depth can be recalculated from these energy fractions. The result should be consistent with the shower depth initially determined in the simulation, which is not guaranteed either.

In order to solve this problem, it is necessary to take into account at least the dominant correlations between the energy and shower depth on the one hand and the energy fractions in each layer (and among themselves) on the other hand. One way of storing the leading correlations between all variables is to build their correlation matrix.

A correlations matrix is only well defined for Gaussian distributions and hence each vertical slice of each histogram in Figures 6.9 and 6.10 is transformed into

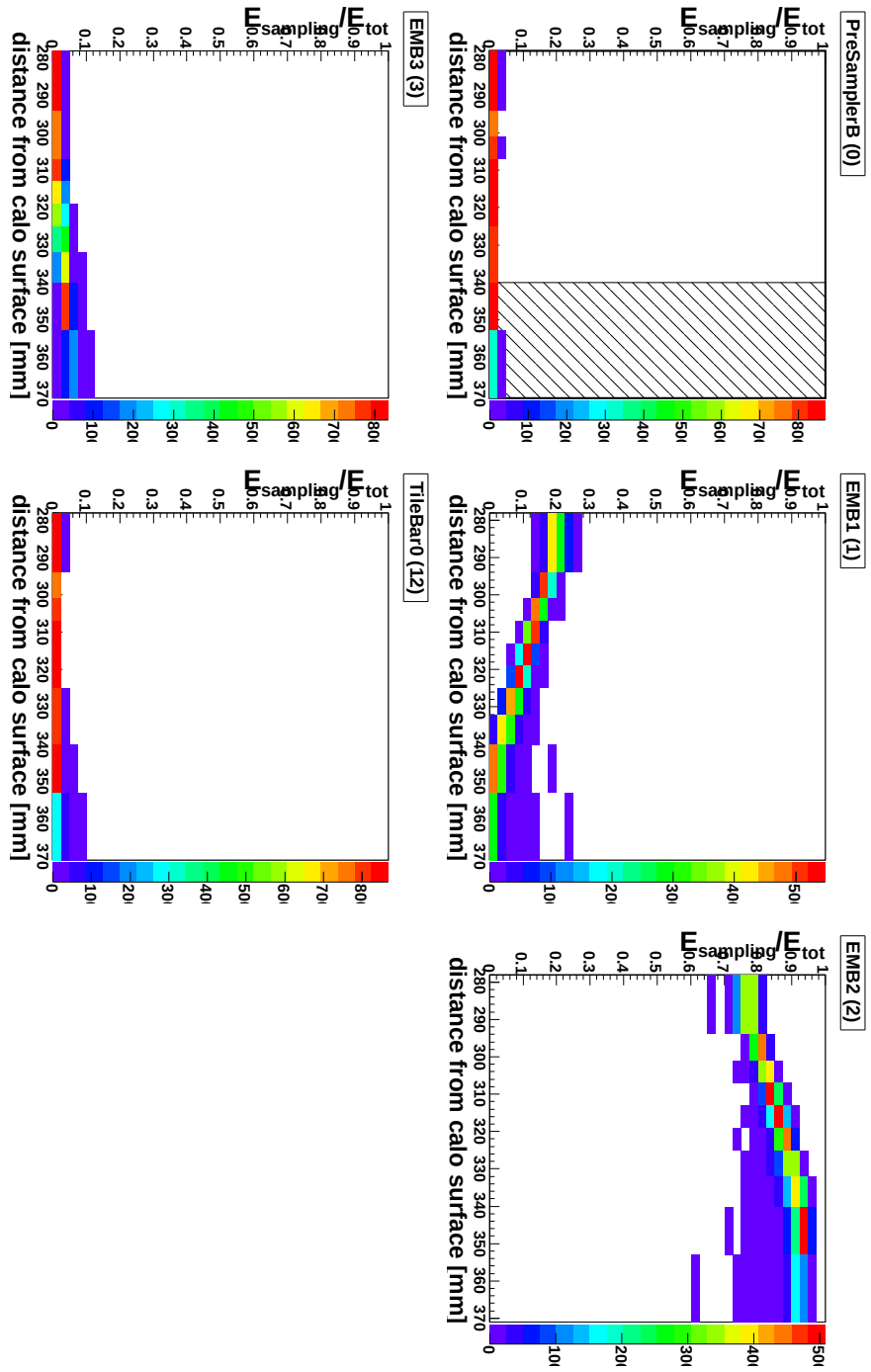


Figure 6.9: Energy response of the individual calorimeter layers to photons of 200 GeV energy in the range $0.20 < \eta < 0.25$. Shown is the fraction of the total energy in each layer as function of the reconstructed longitudinal shower centre. The hatched areas indicate regions that are ignored in the simulation due to low energy. For the naming convention of the calorimeter layers, see Figure 6.1.

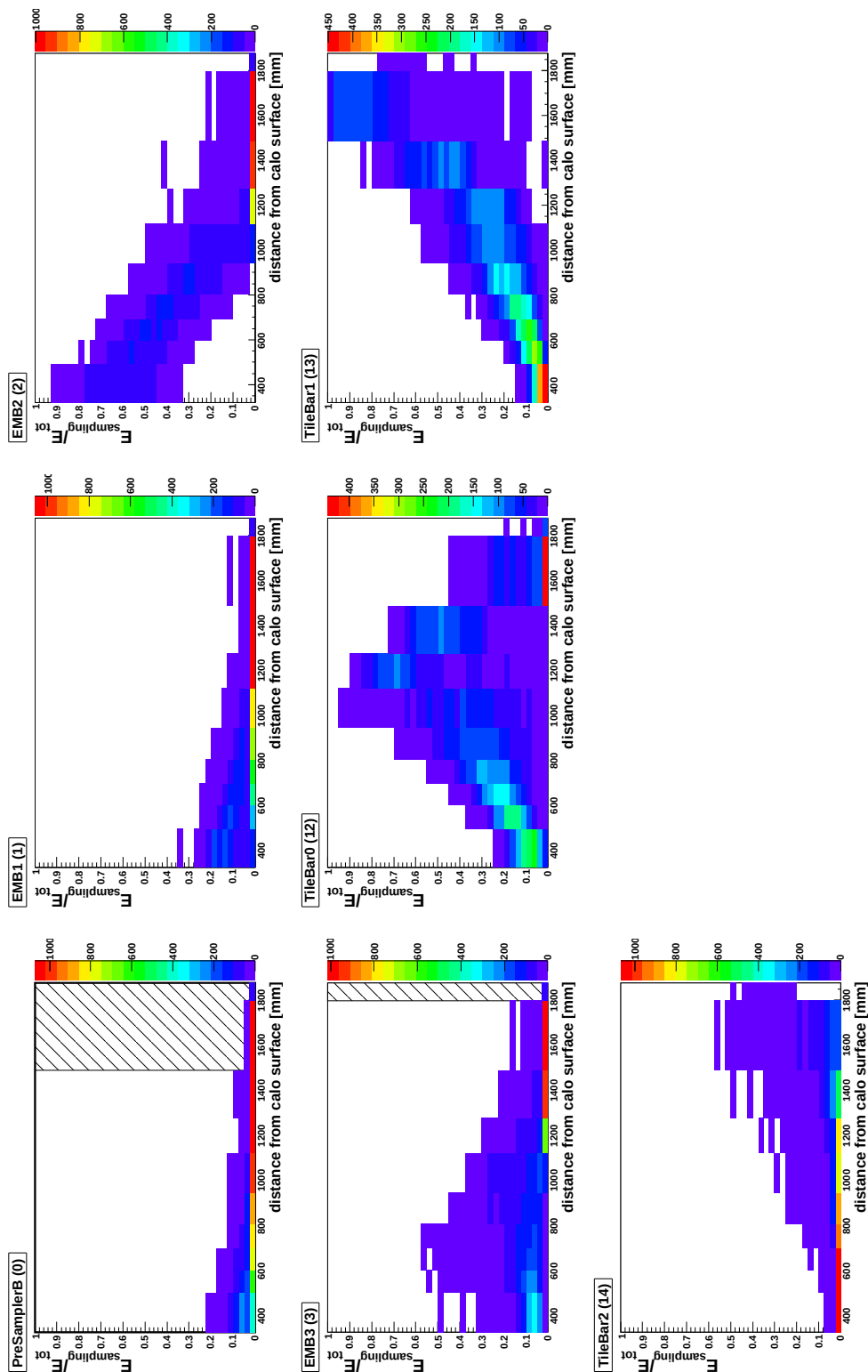


Figure 6.10: The same as Figure 6.9 for pions.

a Gaussian distribution. The details of the transformation scheme are given in Appendix A. The transformation works in such a way that probabilities are conserved. For the mapping $x \rightarrow x'$ the integral of the normalised histogram in the interval $[-\infty, x]$ is equal to the normalised Gaussian integral in $[-\infty, x']$.

Furthermore, the transformation scheme is reversible and can assign a value in an arbitrary source histogram bijectively to a value in a Gaussian distribution. The same transformation scheme is applied to all distributions of energy fractions. Therefore, at this stage, all energy fractions in the calorimeter layers are described by values following normalised Gaussian distributions.

As a last step a correlation matrix is built from these Gaussian-transformed energy fractions, together with the total energy in all calorimeter layers and the shower depth. For each bin in the shower depth distribution, this matrix contains the dominant correlations between the total energy, the shower depth distribution within the bin itself and the energy fractions in the individual calorimeter layers. Figure 6.11 shows such correlation matrices for photons and pions of 200 GeV energy in the range $0.20 < \eta < 0.25$ whose showers were found in the 4th bin of the shower depth distribution.

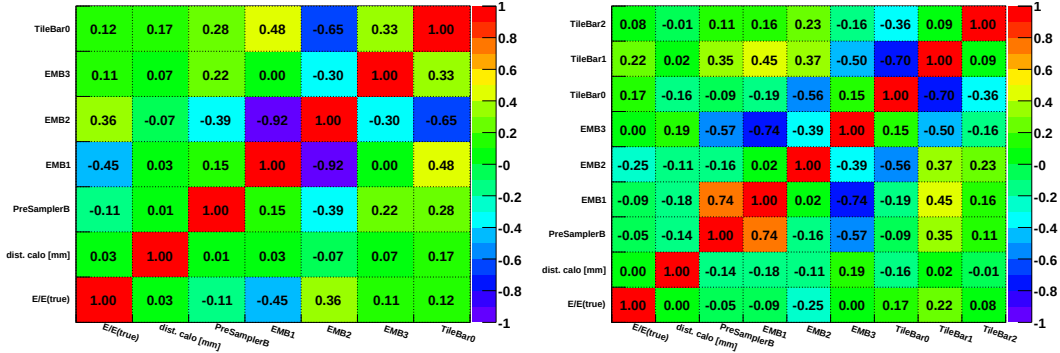


Figure 6.11: Correlation matrix of the calorimeter energy fractions for photons (left) and charged pions (right) of 200 GeV energy in the range $0.20 < \eta < 0.25$ for showers found in the 4th bin of the shower depth distribution (correlations of the distributions in the 4th vertical slice of Figures 6.9 and 6.10).

For the simulation the procedure is inverted. Using the a priori known values of the total energy and shower depth (1st and 2nd row and column in Figure 6.11), correlated Gaussian random numbers for the energy fractions in each layer are dived using the correlation matrix. For each layer the corresponding Gaussian random number is transformed back into the corresponding value of the original distribution that describes the energy fraction in this layer. Due to intrinsic uncorrelated energy fluctuations and because only linear correlations are preserved,

the energy fractions will still not add up to 1 (but the sum is usually close to 1). Therefore, all fractions are rescaled to give a sum of 1 to keep the total energy.

6.3.4 Summary of simulation steps for the energy parametrisation

The following gives a short summary of the simulation steps for the energy parametrisation:

1. Depending on the particle type, it is decided whether the photon or pion parametrisation is applied (see Section 6.2.1).
2. Depending on the particle energy and the η position on the calorimeter surface, one of the energy and η bins of the parametrisation as listed in Table 6.1 is selected.
3. From 2D histograms as in Figures 6.7 and 6.8, correlated values for the calibrated total energy and the reconstructed shower depth are diced.
4. Using the correlation matrix described in Section 6.3.3, correlated Gaussian random numbers are diced for the energy fractions in each calorimeter layer. These Gaussian random numbers are transformed back into energy fractions using the inverted transformation scheme of Appendix A.
5. The calibrated energy in each layer is the simple product of the total calibrated energy from step 3 and the energy fraction from step 4. If the energy fractions do not add up to 1, all fractions are scaled to give a sum of 1.
6. The raw energy in each layer is calculated by dividing the calibrated energy in the layer (step 5) by the calibration constant for this layer (see Section 6.3.1).
7. This raw energy is then added to the calorimeter, using the shape functions explained in the following section.

For each particle type, energy and η position the following parametrisation information must be stored:

- The calibration constants for all calorimeter layers.
- The 2D histogram of calibrated energy versus shower depth.
- The inverse Gaussian transformation for the energy fractions in each layer that needs simulation (only if the mean energy in the layer is sufficiently large, otherwise the whole calorimeter layer is skipped).
- The energy fraction correlation matrix for all layers that need simulation.

6.3.5 Consistency test of the energy simulation

In the following, a few basic consistency tests of the energy parametrisation and simulation are performed.

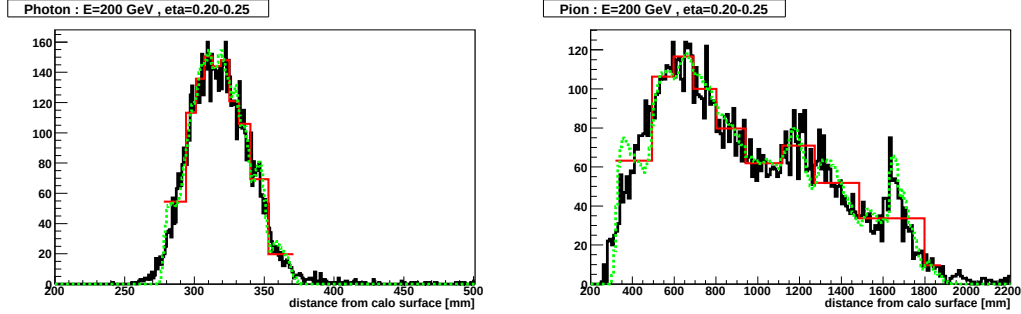


Figure 6.12: Consistency test for the shower depth for photons (left) and charged pions (right) of 200 GeV energy in the range $0.20 < \eta < 0.25$. The original input distribution is shown in black, the rebinned version used as input for the parametrisation and simulation is in red. In green is the distribution of the shower depth for simulated events.

Figure 6.12 shows the shower depth for the original input sample and for the simulated events. The red distribution is the rebinned distribution of the shower depth that is actually used as input for the parametrisation and simulation. One can observe that the simulated green distribution follows the original black distribution better as could be expected from the red input distribution to the simulation. There are two reasons for this: Due to the correlations in the energy parametrisation some information on the substructure within one bin in the shower depth is retained. More important, the shower depth is calculated from the longitudinal position of the calorimeter layers and will always show a bias towards the longitudinal position of the layer with the dominant energy fraction. This explains the reproduction of the structure of the bumps in the pion distribution.

Figures 6.13 and 6.14 show the performance of the energy parametrisation. On the left a comparison of the calibrated calorimeter energy response between the original input sample and simulated events is shown. The plot in the centre shows the raw calorimeter response of the input events and the simulated events after a reversal of the calibration step. On the right, the sum of all energy fractions diced in the simulation is shown. The blue distribution is produced using the correlated random numbers from the correlation matrix, the green distribution would be seen if uncorrelated random numbers were used.

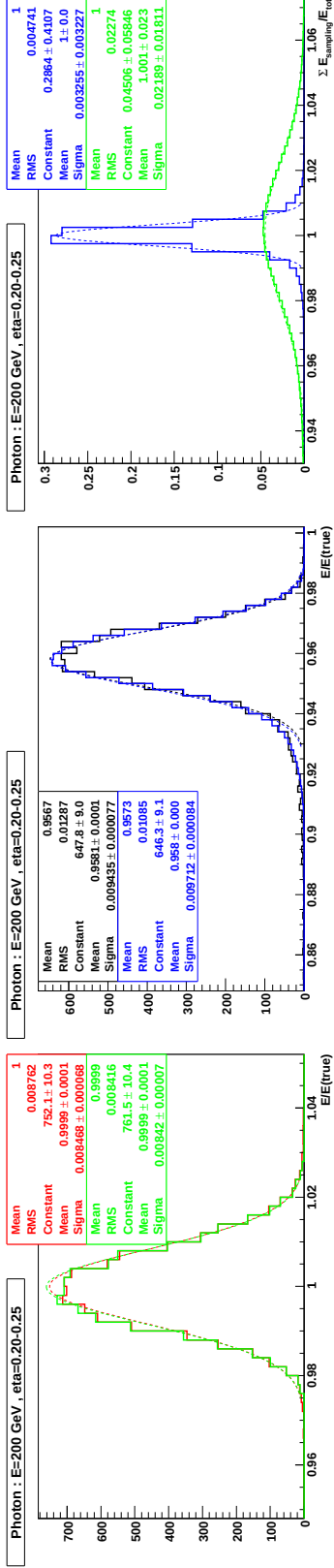


Figure 6.13: Consistency test of the energy parametrisation for photons with 200 GeV energy in the range $0.20 < \eta < 0.25$.

All distributions are overlaid with a Gaussian fit in the $\pm 3\sigma$ range (dashed line).

Left: Distribution of the calibrated energy for the input events (red) and the simulation (green).

Centre: Distribution of the raw energy for the input events (black) and the simulation (blue).

Right: In blue, the distribution of the sum of all energy fractions diced in step 4 of Section 6.3.4 is shown.

The green distribution shows the same sum if the energy fractions were diced uncorrelated.

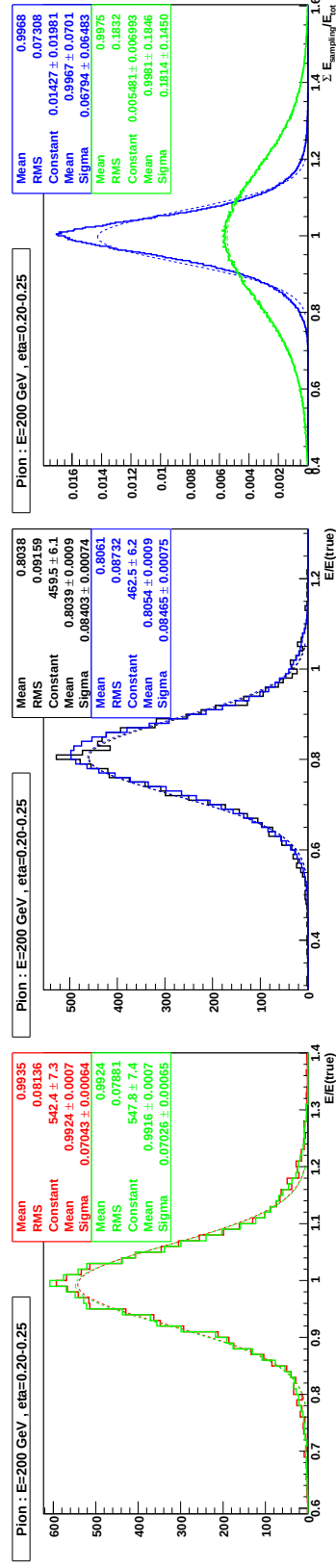


Figure 6.14: The same as Figure 6.13 for pions.

6.4 Shape parametrisation

The parametrisation of the average particle shape is based on the assumption that the energy distribution perpendicular to the direction of flight is radially symmetric. This is a very good assumption, if

1. the effect of the magnetic field is small
2. the calorimeter is homogeneous in all directions perpendicular to the direction of flight
3. the calorimeter layers are perpendicular to the direction of flight

At least items 2. and 3. are not fulfilled for all calorimeter layers, but can be treated as modifications of the radial symmetric shape.

Various methods of parametrisation of the radial particle shape can be found in the literature (for example see Reference [32]), but most are either specific to certain particles or to some calorimeter technology. They are also not directly suitable to produce a parametrisation that allows a deposition of the energy shape into a calorimeter layer in one simulation step (with the risk of some loss in accuracy compared to more specialised parametrisations, of course). Since the particle energies are already parametrised, the aim is to describe the shape of the energy deposition inside one calorimeter layer only⁵.

So instead of trying to find a general analytic function that is able to describe the radial shower profile in various stages of the shower development inside different calorimeter technologies, the shower profile is fitted with generic third order polynomial spline functions. Currently, these fits are performed for all but the forward calorimeters. For the forward calorimeters (FCAL) the complete energy is deposited in the cell hit by the particle trajectory.

6.4.1 Input samples for the shape parametrisation

Since no information on the underlying physical processes is contained in such an empiric parametrisation approach, the shape parametrisations have to be created for each calorimeter layer in a fine grid of particle energy, pseudorapidity and longitudinal shower depth:

- Particle type : Photon and Pion
- Constant particle energy : 5, 10, 20, 50, 100, 200, 500 GeV

⁵Most physics motivated methods of shower description (either in Geant 4 or a parametrisation) produce a three dimensional shower development. For these methods, the subdivision into calorimeter layers and cells is a partitioning of the 3D shape performed in the end. The parametrisation of the shape used here is only concerned with the shape within one layer, the longitudinal shower development is factorised out, only a radial symmetric 2D shape needs to be parametrised.

- Particle pseudorapidity : regions of size 0.05 and spaced by 0.2 in pseudorapidity, well away (if possible) from the edges of a calorimeter layer
- Calorimeter layer (only those where on average sufficient energy is deposited)
- Shower depth bin : same subdivisions as for the energy parametrisation in Section 6.3.2, on average divided into 10 bins.

As an example, the pseudorapidity regions parametrised for 200 GeV photons and pions are shown in Figures 6.15 and 6.16. As input samples, the same samples as discussed in Section 6.3 are used.

6.4.2 Parameters of the shape parametrisation

To fit a particle shape, the true particle trajectory is extrapolated to the longitudinal centre of the calorimeter layer. Figures 6.17 and 6.18 show examples of the displacement in η and ϕ between the energy weighted centre of the shower and the extrapolated position of the true particle trajectory. In most cases this displacement is close to zero (as in Figures 6.17 and 6.18), but since the whole shower will not propagate through the detector and magnetic field in the same way as a single particle, non-negligible offsets are needed farther in the calorimeter and at high η . The expected point of impact in the calorimeter layer is the sum in η and ϕ of the extrapolated true particle trajectory and the average shower displacement.

Figures 6.19 and 6.20 show an example of the energy distribution in the calorimeter around the expected impact point. Shown is the average energy of cells at a given distance from the impact point. The distance is calculated from the centre of the cell. Since no information on the energy distribution within one calorimeter cell is available, these plots show a convolution of the true average particle shape and the cell dimensions. For the example shown in Figures 6.19 and 6.20, the cells have a size of approximately 40 mm in η and ϕ direction. As can be seen, this is a distance scale at which cell energies fall off by approximately one order of magnitude, so the effect of the cell size cannot be neglected.

An independent shape function is fitted to all cells of all events within one bin of the shower depth, and this function is described by several parameters:

- An offset in η and ϕ direction which gives the average displacement of the shower centre with respect to the extrapolated particle position in the layer.
- A 3rd order polynomial spline that gives the average energy $E(r)$ found in a shower at a distance r (measured in mm) from the expected impact point. The energy in a shower falls off very rapidly with r , so the actually fitted spline is chosen to describe $\ln (E(r) / E(0))$.

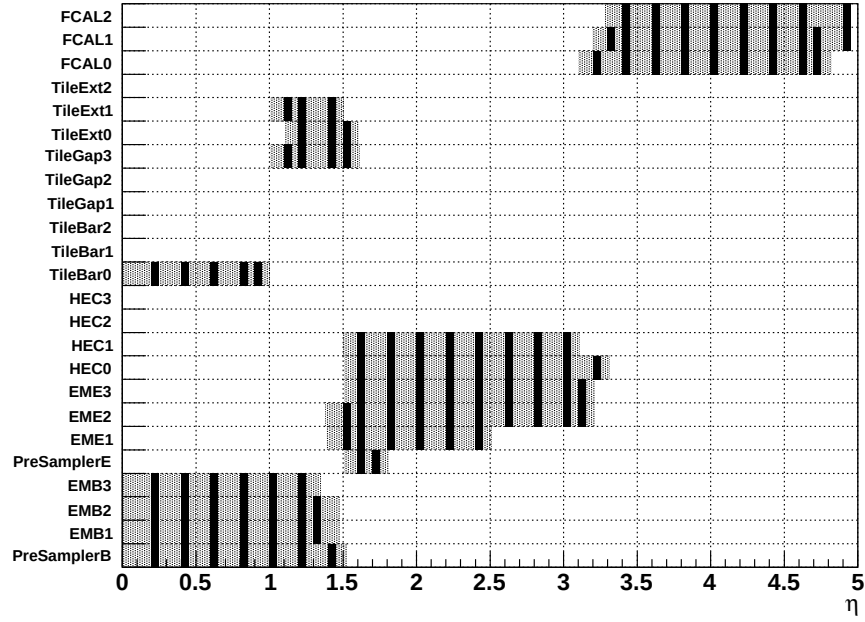


Figure 6.15: The pseudorapidity intervals within each calorimeter layer for which a parametrisation of the radial shower profile was prepared are shown as black marked regions for photons of 200 GeV energy. The hatched areas indicate the pseudorapidity ranges covered by the calorimeter layers.

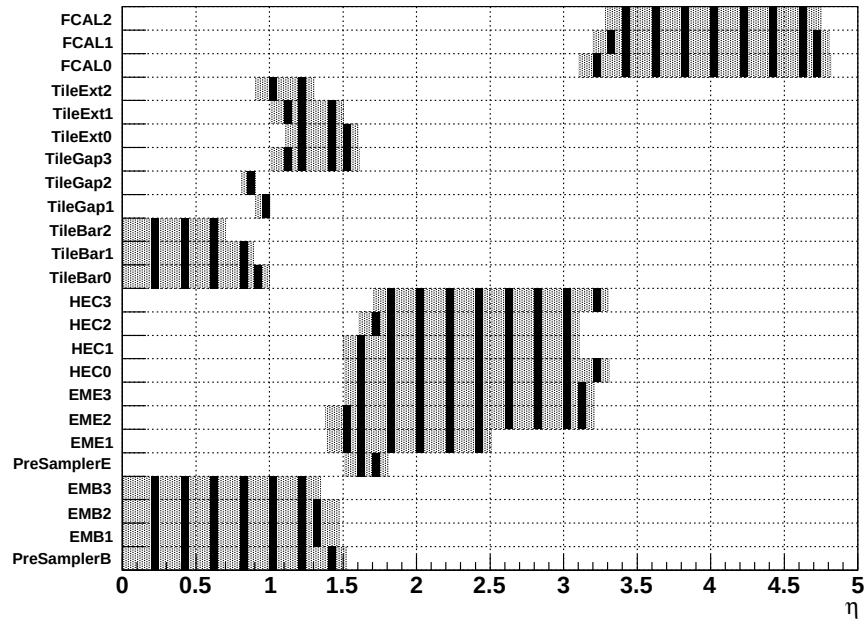


Figure 6.16: The same as Figure 6.15 for pions.

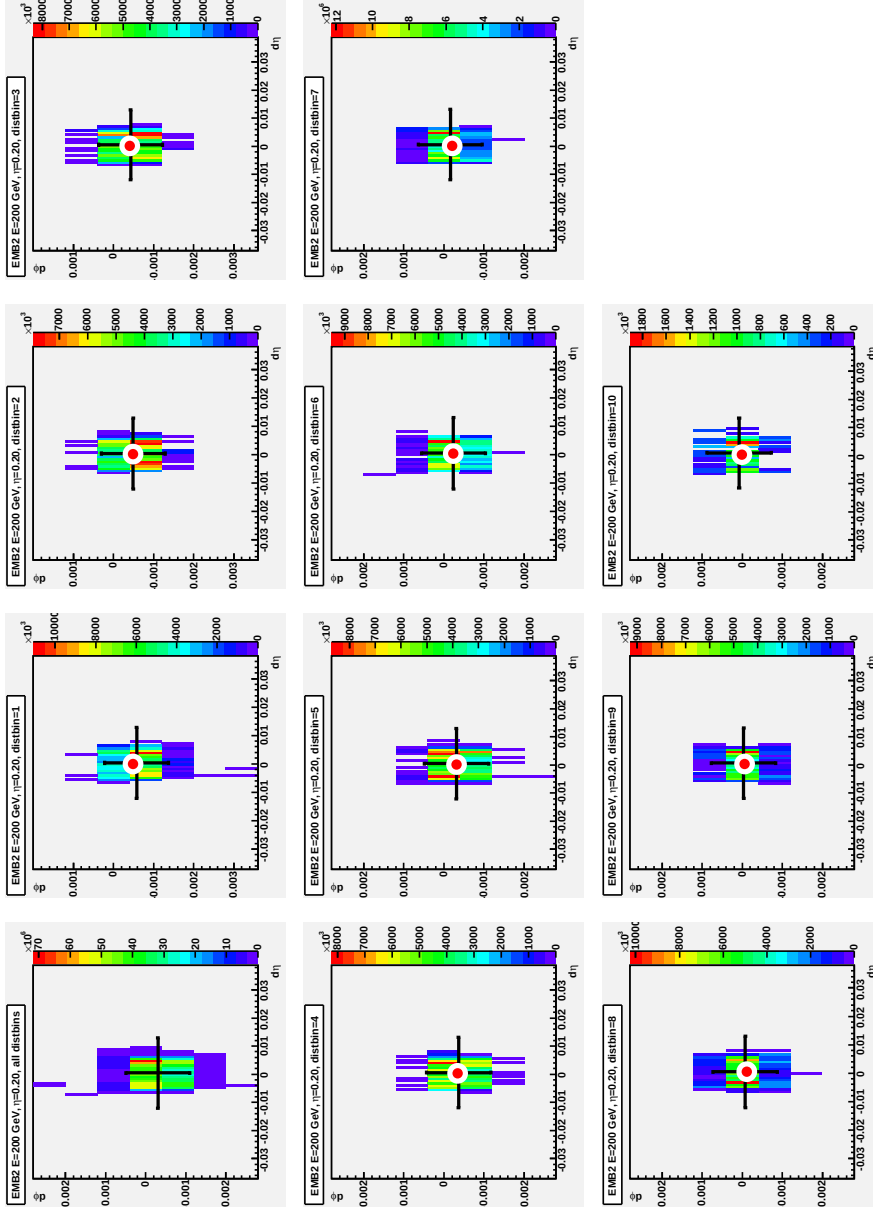


Figure 6.17: Displacement in η and ϕ between the energy weighted centre of the shower in the 2nd EM barrel layer and the extrapolated position of the true particle trajectory for photons of 200 GeV energy in the range $0.20 < \eta < 0.25$. The upper left plot shows the average displacement for all events. The other plots show the displacement in each bin of the shower depth distribution (bins defined as in the red histograms in the left of Figure 6.7). The black error bars give the position and RMS of the distributions, the red dot in the white circle the average displacement found by the fit.

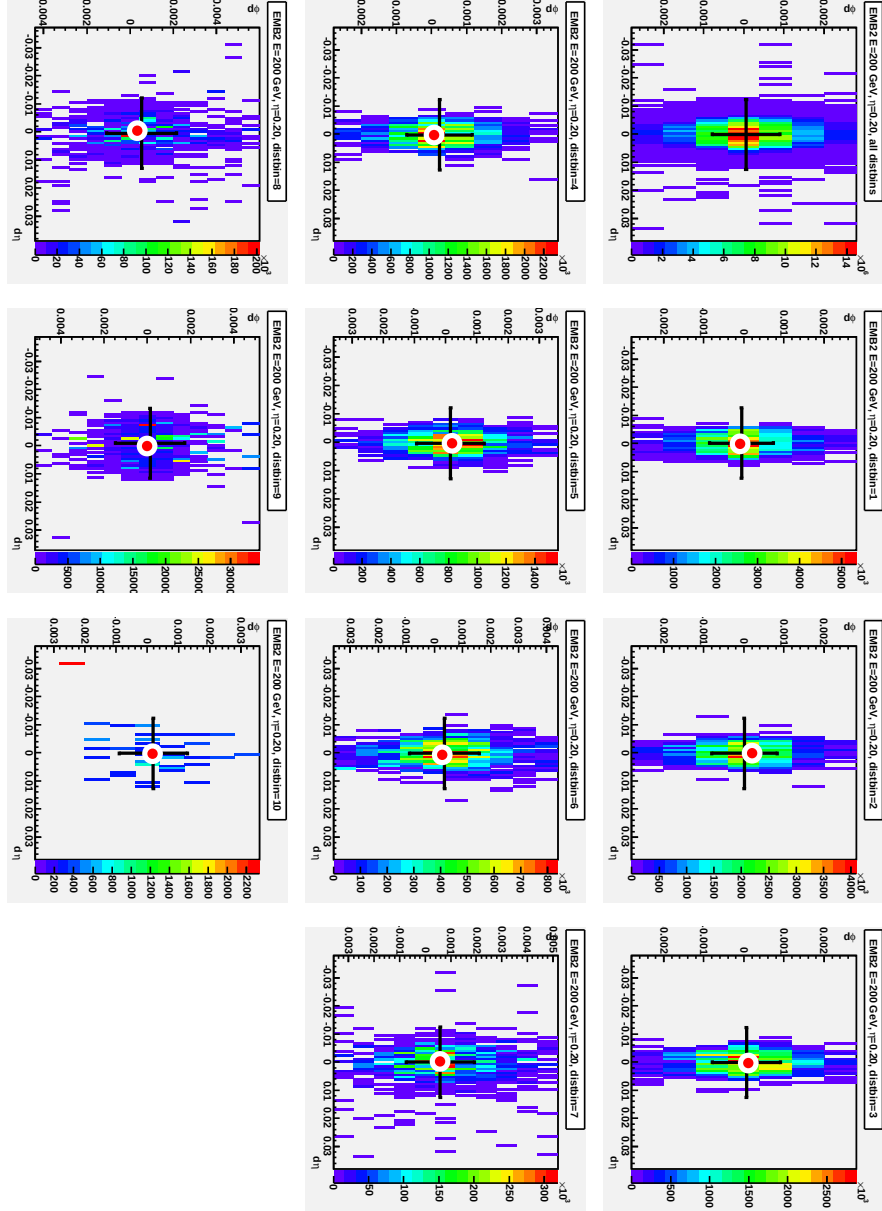


Figure 6.18: The same as Figure 6.17 for pions.

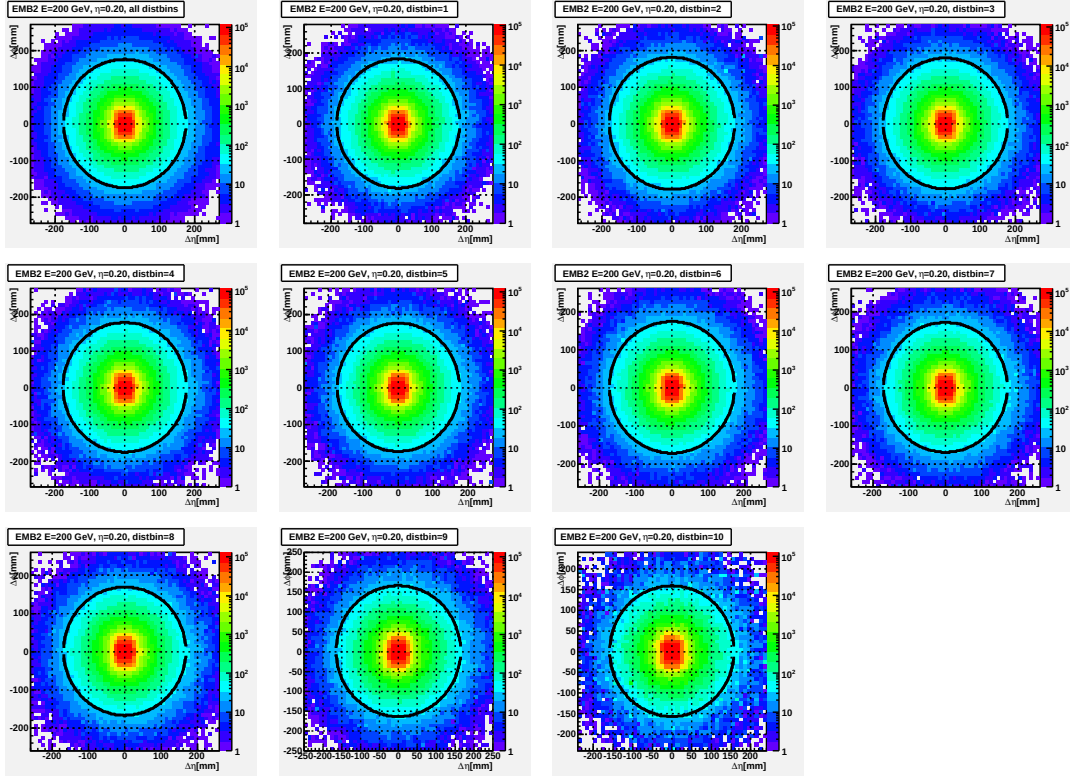


Figure 6.19: Energy distribution in the 2nd EM barrel calorimeter around the expected point of impact for photons of 200 GeV energy in the range $0.20 < \eta < 0.25$. The upper left plot shows the average energy distribution for all events. The other plots show the average energy distribution in each bin of the shower depth distribution (bins defined as in the red histograms in the left of Figure 6.7). Each colour coded entry in the histogram gives the average energy of cells whose centre position is at a distance from the impact point given by the bin position. Cell dimensions are approximately 40 mm in the η and ϕ direction (for $\eta \approx 0.2$ in this calorimeter layer), so this distribution can be seen as the average true distribution smeared by the cell dimensions. The black circle gives the maximum radius r_{max} up to which a parametrisation of the shape is attempted. Outside this region, less than 0.5% of the average particle energy in the layer can be found.

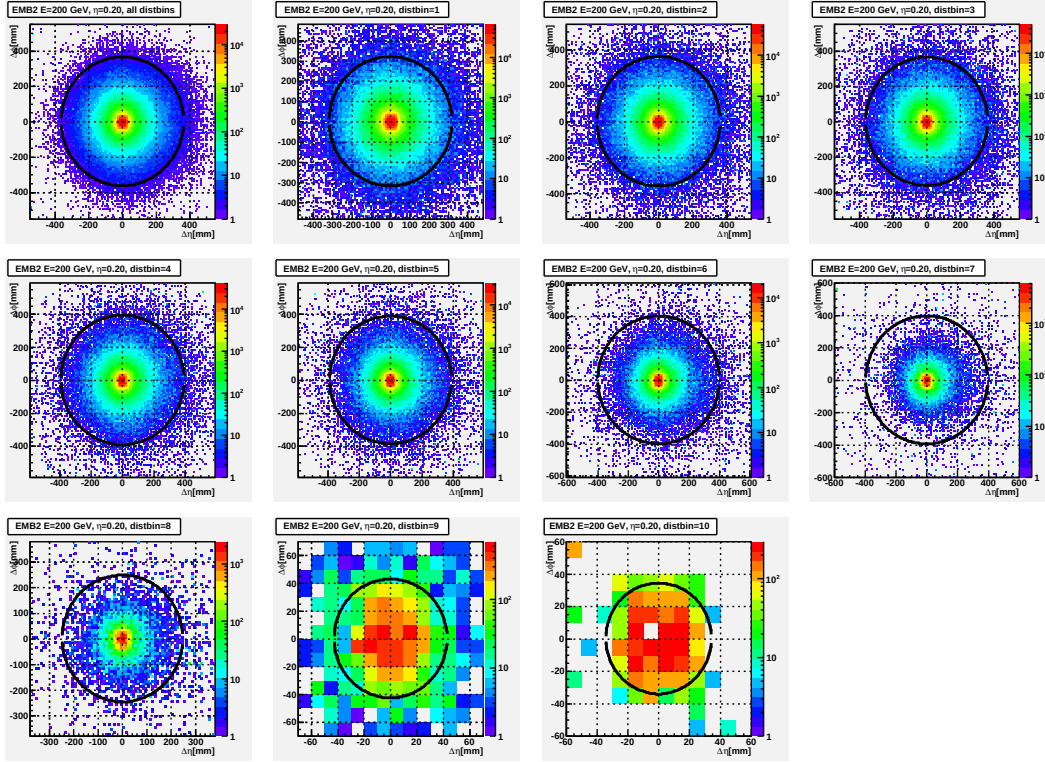


Figure 6.20: Same as Figure 6.19 for pions of 200 GeV energy in the range $0.20 < \eta < 0.25$ in the 2nd EM barrel calorimeter.

- The longitudinal boundaries of the calorimeter (see Figure 6.1) are not perpendicular to the axis of the shower development, which normally follows η projective lines. At small angles between the shower axis and the calorimeter surface, the component of the shower seen inside one calorimeter layer will show an asymmetry in the η direction (see illustration in Figure 6.21).

This asymmetry is described by a second spline that gives, depending on the η distance from the shower centre, a broadening factor of the shower shape in η direction. Only for the 1st EM calorimeter sampling (EMB1 and EME1), which has a very fine granularity in the η direction (see Figure 6.22 on the left), this asymmetry is sufficiently large to be visible after reconstruction and thus needs to be simulated. In all other calorimeter layers it can be neglected.

- Additional simulation parameters: a maximum radius $r_{\max}$ of the shower, chosen to contain 99.5% of the layer energy; an integration granularity Δ_I , chosen to be 1/6 of the average radial distance where 10% of the core energy is reached; a global scale factor $f_{\text{E-scale}}$ applied to all cell energies used for normalisation.

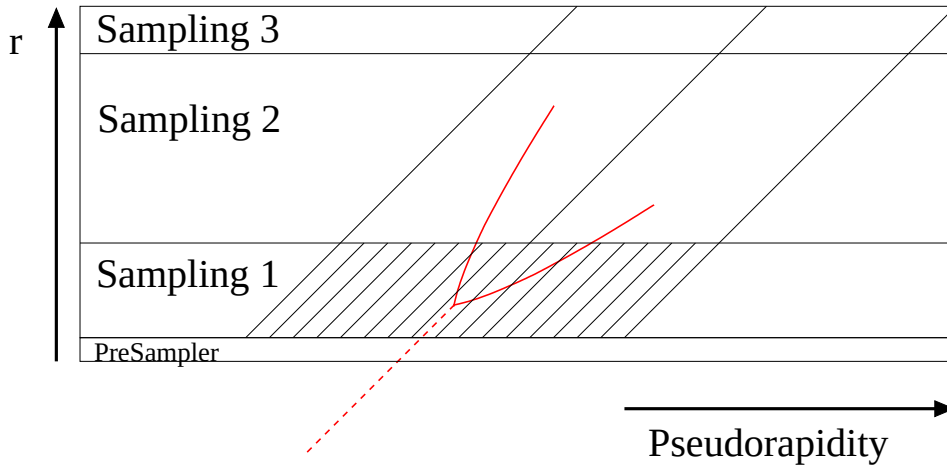


Figure 6.21: Illustration of the η asymmetry seen in the shower shape of a particle inside one calorimeter layer. This is a sketch of a photon entering the calorimeter at a large angle. The photon starts its shower somewhere in the 1st sampling of the calorimeter. The further the shower penetrates the calorimeter, the broader the shower gets. In the 1st sampling, cells located to the right of the shower centre see a larger part of the shower than cells to the left of the centre, so the shower shape appears broader to the right of the centre than to the left.

6.4.3 Fitting of the shower shape

The shape for one calorimeter layer is extracted from a χ^2 fit to all particles that belong to the same event category (consisting of particle type, energy, η , calorimeter layer and shower depth). The parametrisation of the total energy inside one calorimeter layer was already described in Section 6.3, so only the parametrisation of the relative distribution of energies in this calorimeter layer remains.

Especially pion shapes have large fluctuations which make it necessary to run several fit iterations before a sufficiently good shape description is found. After the end of one fit iteration, outliers are discarded for the next iteration, errors recalculated and fit settings are optimised. For each single fit step in one of these fit iterations, the shape function to be fitted is evaluated and a global χ^2 is calculated. The parameters of the shape function are modified in each fit step until a minimal χ^2 for this fit iteration is reached.

1. For each event i and for every calorimeter cell c in the layer, the expected energy content $\hat{E}_{i,c}$ is calculated using the shape function. To do this, the shape function is integrated in the η - φ -plane over the area of the cell. The integration is performed by equidistant sampling of the shape function

values (using the integration granularity parameter Δ_I as point distance) in the cell area. Only cells that are within the maximum shower radius $r_{\max}$ are considered. Cells that overlap with $r_{\max}$ get a fractional weight w_c that corresponds to the cell area fraction inside $r_{\max}$.

The EM calorimeter has an accordion shape in φ direction (see Figure 6.22 on the left). The effect of this shape is that a cell has a linearly increasing response function of the size of the accordion amplitude and thus a point-like energy deposition can end up in 2 cells. Up to $|\eta| < 2.5$ the accordion shape amplitude is $\Delta\varphi = 2\pi/256 \approx 0.025$, for $|\eta| > 2.5$ it is $\Delta\varphi = 2\pi/64 \approx 0.1$.

To take this into account, the actual cell integral includes the linear turn on function at the φ -edges of the cell and covers a sufficiently large area to sample the whole integral. In the ATLAS test beam, this linear energy turn on effect can be seen clearly for single muons (this is the best approximation of a point like energy deposition) [33]. The right side of Figure 6.22 shows the φ -dependency of the cell energy on the impact point of the muon in the cell.

2. All cell integrals are rescaled to give $\sum_c \hat{E}_{i,c} = f_{\text{E-scale}} \cdot \sum_c E_{i,c}$, where $E_{i,c}$ is the energy of cell c .
3. An average cell χ_i^2 for the event i is defined as

$$\chi_i^2 = \frac{1}{\sum_c w_c} \cdot \sum_c w_c \cdot \left(\frac{E_{i,c} - \hat{E}_{i,c}}{\Delta \hat{E}_{i,c}} \right)^2$$

where $\Delta \hat{E}_{i,c}$ is the error on the energy in cell c . The energy error $\Delta \hat{E}_{i,c}$ is taken to scale with $\sqrt{E_{i,c}}$: $\Delta \hat{E}_{i,c} = p_{\text{layer}} \cdot \sqrt{\hat{E}_{i,c}}$, where p_{layer} is 10% for the EM barrel and endcap, 50% for Tile and HEC and 100% for FCal (see Section 5.2.2). For the first fit iteration, $\Delta \hat{E}_{i,c}$ is calculated directly from the energy in the cell, once the fit is stable it is calculated from the estimated cell energies $\hat{E}_{i,c}$ to get the expected error on the cell energy⁶.

If the cell energy $\hat{E}_{i,c}$ is very small, also the error decreases as $\sqrt{\hat{E}_{i,c}}$ and approaches $\Delta \hat{E}_{i,c} \rightarrow 0$ in the limit of $\hat{E}_{i,c} \rightarrow 0$. Fitting cells in a χ^2 fit that have an error close to 0 (or, even worse $\Delta \hat{E}_{i,c} = 0$) will give these cells an unnatural large weight especially in case of fluctuations in the tails. If these

⁶If the error $\Delta \hat{E}_{i,c}$ were calculated from $\hat{E}_{i,c}$ in every fit step, the global χ^2 minimum of the fit would favour large estimates of $\hat{E}_{i,c}$, since they would give slightly bigger errors $\Delta \hat{E}_{i,c}$ as well and hence a smaller χ^2 . To avoid this bias, $\Delta \hat{E}_{i,c}$ is calculated only once in the beginning of each fit iteration with the best parameter values of the last fit iteration. During the fit the error stays constant. For a stable result, the fit needs to be iterated a few times.

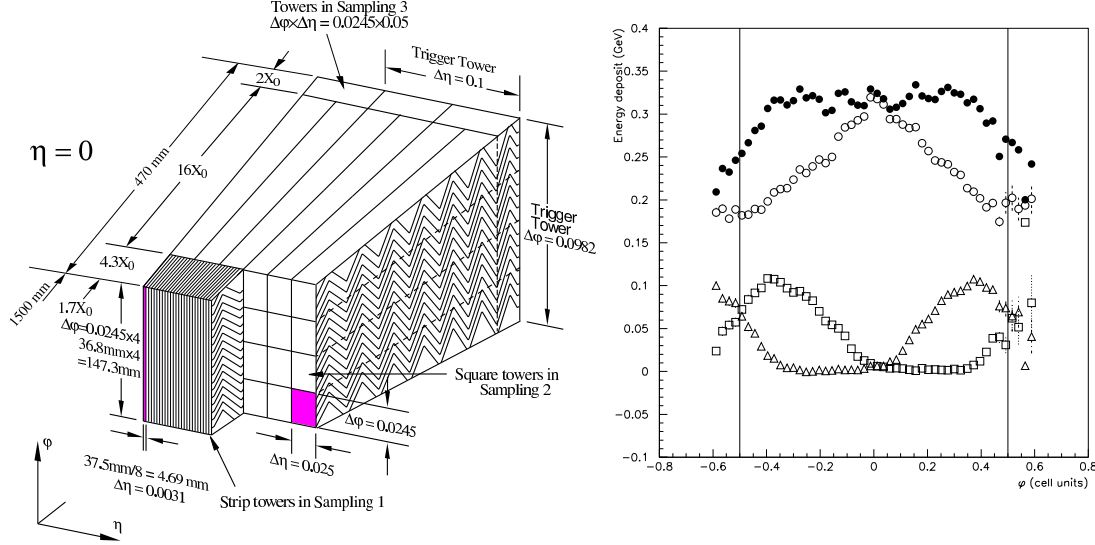


Figure 6.22: **Left:** Illustration of the EM calorimeter (see Section 5.2.2), showing the subdivision into longitudinal layers and cells in η and ϕ . One real cell in ϕ consists of 4 accordion shaped readout elements (in the EM2 layer) and actually covers twice the area in ϕ as given by the granularity.

Right: Figure taken from Reference [33]: Energy deposition of muons in the EM calorimeter as function of the position of the muon impact point in ϕ . Open circles : energy deposition in the cell closest to the impact point (central cell).

Open squares : energy deposition in the neighbouring cell in ϕ , in the direction of decreasing ϕ .

Open triangles : energy deposition in the neighbouring cell in the direction of increasing ϕ .

Filled circles : Sum of energy deposition in the central cell and in the adjacent ϕ cell closest to the muon impact point.

cells were completely skipped, the fit would produce a positive bias for cell energies in the tail of the energy distribution. This is solved by introducing a minimum cell energy error that is determined by the energy and error for which $\Delta\hat{E}_{i,c} = p_{\text{layer}} \cdot \sqrt{\hat{E}_{i,c}} = \hat{E}_{i,c}$. As soon as a cell energy is sufficiently small to give an error of 100%, all smaller energies are considered indistinguishable and the cells get the minimum error assigned. This approach will also give some bias, but it seems to be sufficiently small to be acceptable.

4. A global χ^2 is calculated as the average of the per event χ_i^2 values. Since a few extreme outliers (showers with very unusual shower development)

can produce a large bias in the fit, the 5% of events with the highest χ_i^2 values during the last fit iteration are excluded from the χ^2 calculation. The absolute value of this χ^2 is not meaningful, as only the intrinsic calorimeter resolution is included in the error definition of the cell energies but not the intrinsic fluctuations of the shower development itself.

5. The fit steps are continued until shape function parameters which give a stable minimal χ^2 are found. Once the stable minimum is reached, one fit iteration ends.

This fit to minimise the global χ^2 is performed iteratively until a stable shape is determined. Since the conditions of the fit (cell errors, outliers) depend on the final shape of the previous fit iteration, several iterations are needed to converge.

Figures 6.23 and 6.24 show the average ratio between the cell energies and the fitted energy shape as function of the cell distance from the impact point in η and ϕ direction.

For photons, which have only small fluctuations in the shower shape, the fits describe the energy distributions very well. Also pions with a small shower depth and therefore a large energy content in the 2nd EM layer (as shown in the 3rd plot in the top row of Figure 6.10) are quite well predicted, especially in the core region. With increasing shower depth, the energy content in the 2nd EM layer decreases and fluctuations, especially in the tail of the shape, get more and more pronounced. A lot of these fluctuations come from the intrinsic shower development of the pion and can lead to large energy depositions at a large distance from the shower centre. Such fluctuations cannot be described by an average shape function and dominate the tails of the ratio plots in Figure 6.24.

In the last bins of the shower depth distribution, the pion passes the 2nd EM calorimeter essentially as a minimal ionising particle and only a very narrow energy deposition is found.

The η - ϕ -distributions in Figures 6.23 and 6.24 also demonstrate, that the quality of the shape description is essentially radially symmetric. To visualise the fluctuations, Figures 6.25 and 6.26 show the distribution of the ratios between the cell energies and the fitted energy shape as function of the radial cell distance from the impact point. The fluctuations close to the impact point are very small, since the integral of the fitted function was normalised to the observed energy in each event. The total energy is deposited to a large degree in the cells close to the impact point, so most fluctuations in these cells are absorbed in the normalisation.

6.4.4 Summary of the simulation steps for the shape parametrisation

The following gives a short summary of the simulation steps for the shape parametrisation:

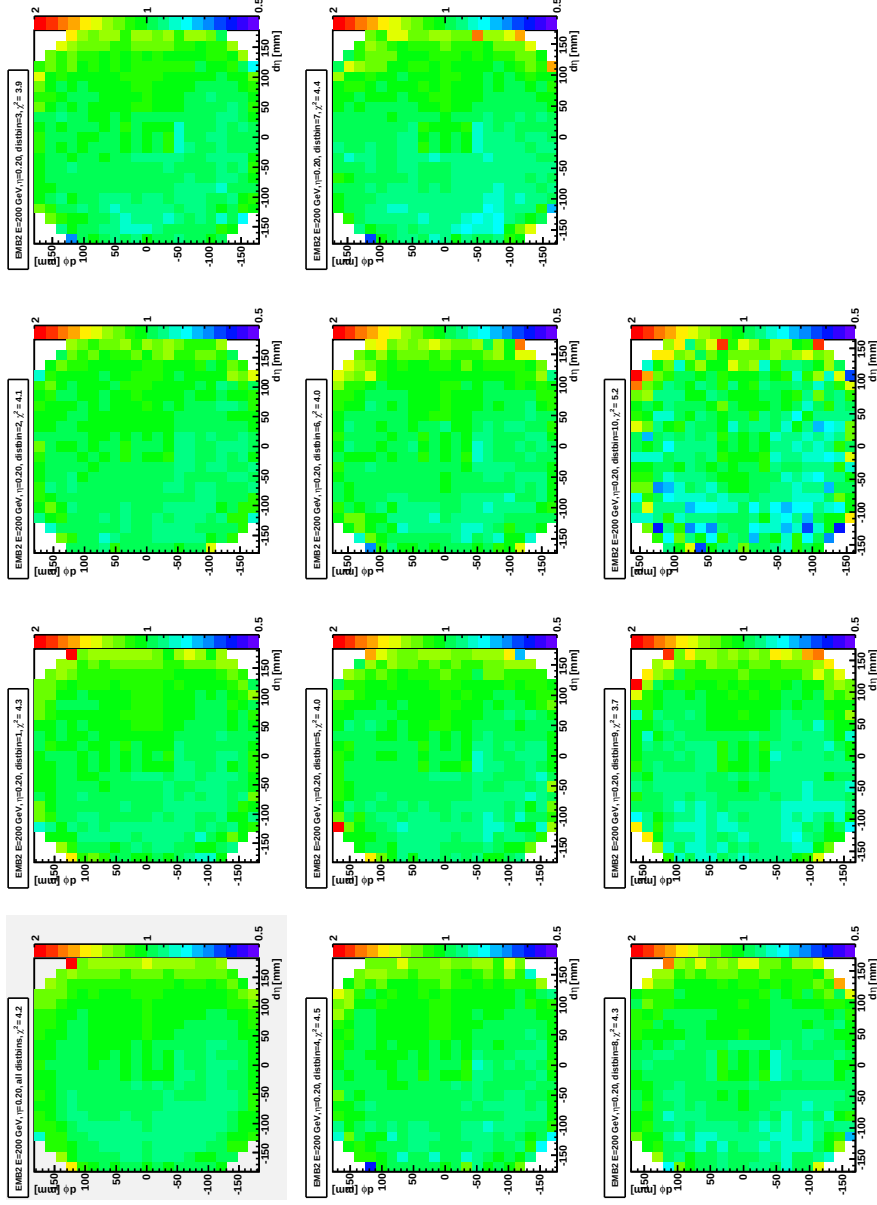


Figure 6.23: Average ratio between the cell energies and the fitted energy shape as function of the cell distance from the impact point in η and ϕ direction for photons of 200 GeV energy in the range $0.20 < \eta < 0.25$ in the 2nd EM barrel calorimeter. The upper left plot shows the average ratio for all events. The other plots show the average ratio in each bin of the shower depth distribution (bins defined as in the red histograms in the left of Figure 6.7).

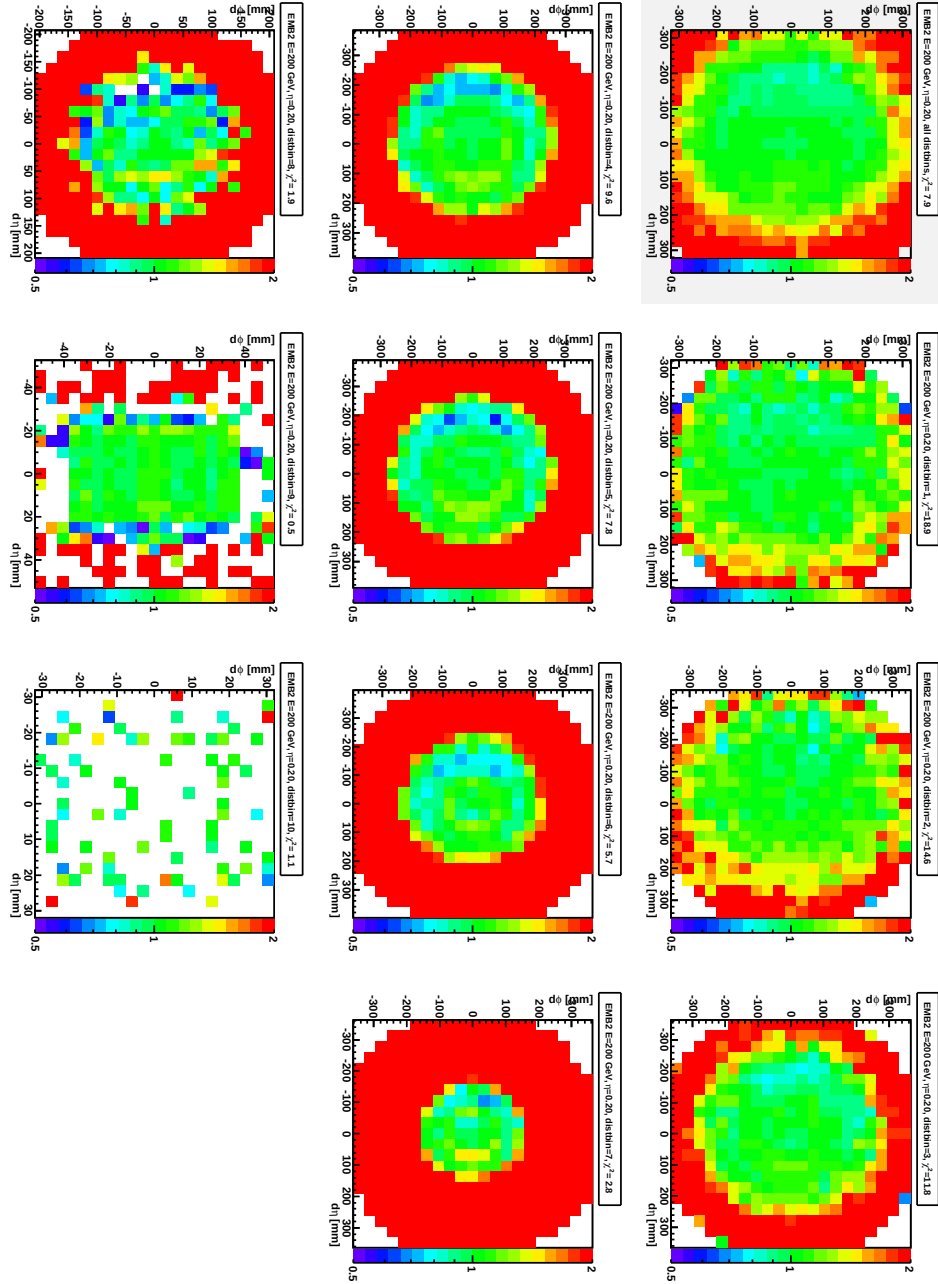


Figure 6.24: The same as Figure 6.23 for pions.

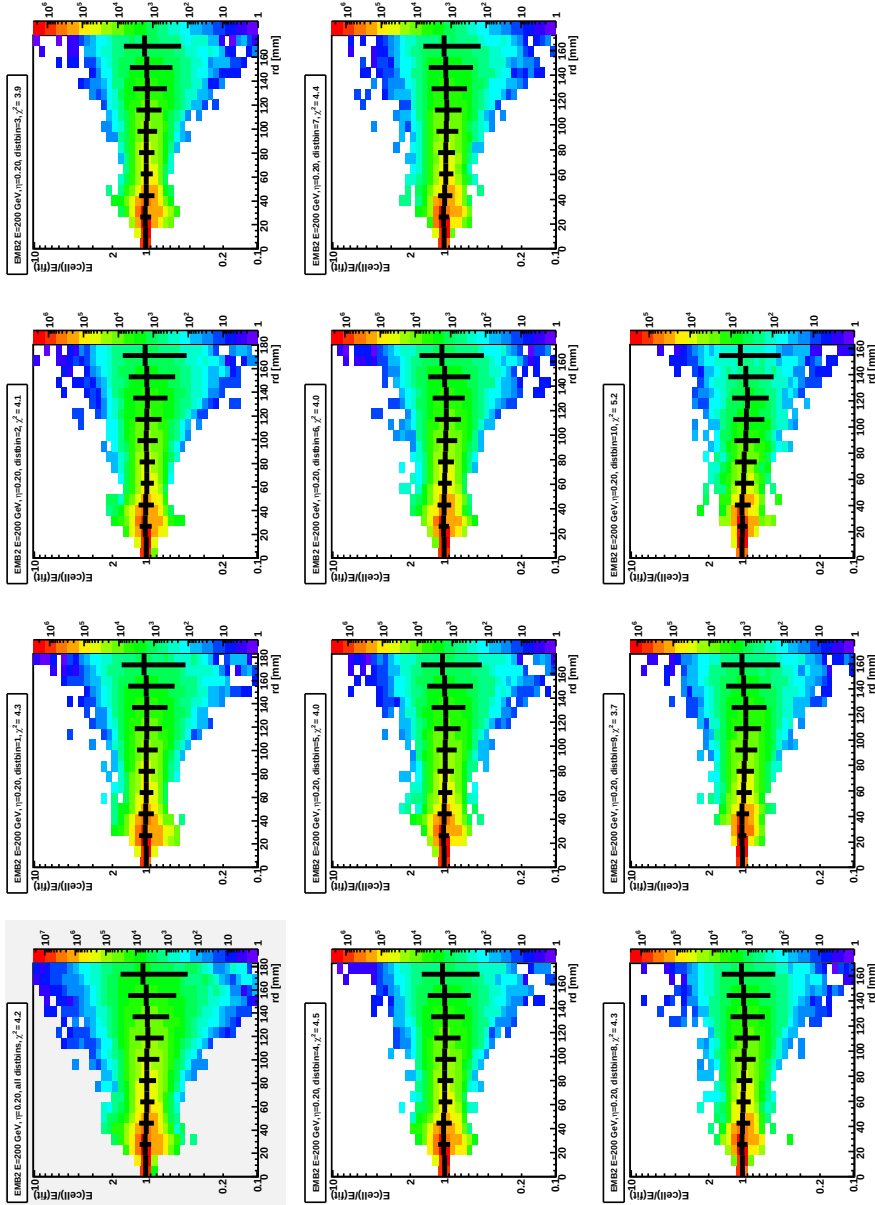


Figure 6.25: Ratios between the cell energies and the fitted energy shape as function of the radial cell distance from the impact point for photons of 200 GeV energy in the range $0.20 < \eta < 0.25$ in the 2nd EM barrel calorimeter. The upper left plot shows the distribution of ratios for all events. The other plots show this ratio in each bin of the shower depth distribution (bins defined as in the red histograms in the left of Figure 6.7). The black error bars give the mean and RMS of the ratio distribution.

1. Depending on the particle type, it is decided whether the photon or pion parametrisation should be applied.
2. The particle trajectory is extrapolated to the longitudinal centre of the calorimeter layer.
3. Depending on the calorimeter layer, the particle energy, the extrapolated η position and the shower depth dived in the energy simulation step, the most suitable shape parametrisation is selected. Parametrisations are provided for all energy and η positions as summarised in Section 6.4.1. Using the shape position offset (see Section 6.4.2), the impact point of the particle is calculated.
4. Particles with impact points beyond the edges of a calorimeter layer may still deposit energy in the layer, since showers are extended objects. In these cases, one cannot expect that the shapes fitted for the homogeneous calorimeter regions will be appropriate. However, at least the energy can be deposited with the available shapes, even if the shape is not well described.
5. For all cells that are within a radius of $r_{\max}$ of the impact point, the integral of the shape function is calculated in the same way as described for the fitting of the shower shape in Section 6.4.3. All integrals are rescaled to give in the sum the same energy as was determined for this calorimeter layer in the energy simulation step.
6. All cell energies are modified randomly (intrinsic calorimeter fluctuations) according to the intrinsic energy resolution of the calorimeter technology. The same fluctuation values as for the cell energy errors (step 3 of Section 6.4.3) are used.
7. The total effect of the calorimeter fluctuations are already included in the simulation of the energy fractions in the calorimeter layer. To avoid taking this into account twice, the sum of all cell energies is again rescaled to give the same energy as was determined in the energy simulation step.

6.5 Implementation in Athena

The fast calorimeter simulation FastCaloSim is integrated directly into the ATLAS reconstruction instead of being run independently before, as it is done for the full simulation. At the point where during a normal reconstruction run the calorimeter information for an event is read from the (simulated) detector output, FastCaloSim is executed instead. The overall execution procedure for the ATLAS reconstruction and the integration of FastCaloSim is schematically illustrated in Figure 6.27.

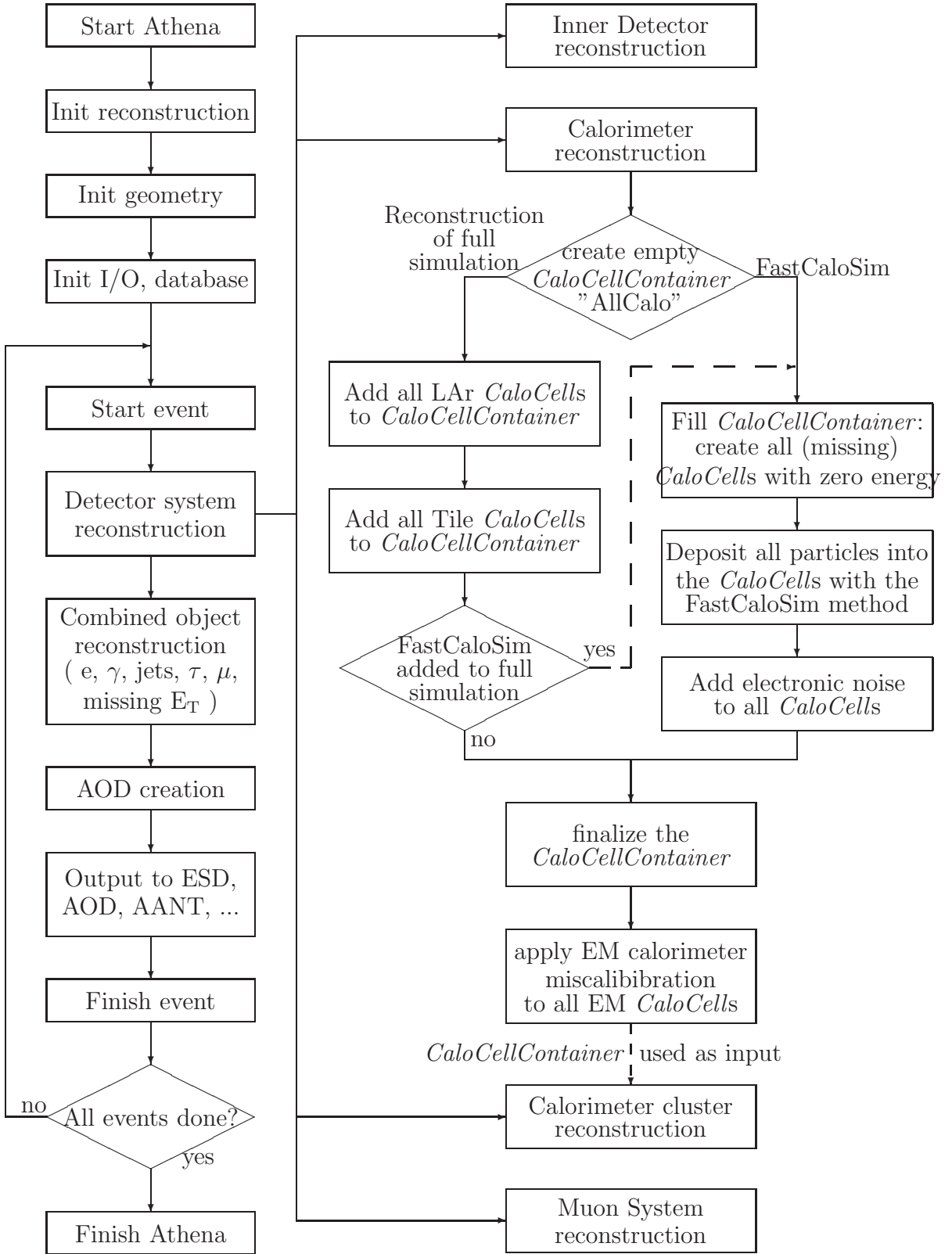


Figure 6.27: Schematic view of the implementation of the fast calorimeter simulation FastCaloSim in the ATLAS reconstruction software.

6.5.1 Important c++ class names

CaloCell is the c++ class name of the object that contains the energy of a calorimeter cell and its type and geometric position information. There is one instance of the *CaloCell* object for each real calorimeter cell in the ATLAS detector. The c++ object *CaloCellContainer* contains the collection of all calorimeter cells and represents the total ATLAS calorimeter. For each event one instance of *CaloCellContainer* with the name “AllCalo” is created and used for the further reconstruction steps.

6.5.2 Execution steps of FastCaloSim

The calorimeter reconstruction is controlled by the Athena *CaloCellMaker* algorithm. This in turn calls several tools that do the actual calorimeter processing. *CaloCellMaker* starts by creating an empty *CaloCellContainer*. For FastCaloSim, this container is then processed in several steps by the FastCaloSim tools:

1. The empty container is filled with all calorimeter cells (*CaloCell*) that are contained in the ATLAS detector by the *EmptyCellBuilderTool*. The energies of all cells are set to zero. Altogether, there are 173312 LAr EM cells, 5184 hadronic Tile cells, 5632 LAr HEC cells and 3524 LAr FCAL cells.
2. For each particle found in the truth information of the event, the energy and shape simulation, as described in Sections 6.3 and 6.4, is done by the *FastShowerCellBuilderTool*. The resulting cell energies are added to the corresponding cells in the *CaloCellContainer*.
3. Depending on the energy of the cells in the *CaloCellContainer*, the gain factor of the electronic read-out system and the RMS of the electronic noise, with which this energy would have been measured, is determined by *AddNoiseCellBuilderTool*. A random energy with an RMS of the electronic noise of the cell is then added to the cell energy.
4. The container *CaloCellContainer* is finalised by the *CaloCellContainerFinalizerTool*.
5. An EM scale miscalibration, as it is expected to be present for the real detector, is applied by the *LArCellEmMiscalib* tool.

At this stage, the produced *CaloCellContainer* is from the software point of view indistinguishable from the same container produced by the normal reconstruction used with full simulation and also with real data. The only information that is currently not filled is the cell time information.

It is also possible to combine FastCaloSim with the full simulation. A use case is to run a full simulation of all particles in the inner detector. After exiting

the volume of the inner detector, the full simulation is continued only for muons, which then also deposit energy in the calorimeter and get tracked through the muon system. In the reconstruction, the fully simulated calorimeter with the energies only from the muons is first processed and then FastCaloSim is executed as described above to add the simulation of the calorimeter energies for the remaining particles in the event.

6.5.3 Python configuration

The configuration is done on the python side by *CaloCellGetter.py*. The execution of FastCaloSim can be steered by flags in the *CaloCellFlags* jobproperties.

- *jobproperties.CaloCellFlags.doFastCaloSim* activates all FastCaloSim tools.
- *jobproperties.CaloCellFlags.doFastCaloSimNoise* switches on the *AddNoiseCellBuilderTool*.
- *jobproperties.CaloCellFlags.doFastCaloSimAddCells* controls that the full simulation calorimeter is reconstructed first and then the FastCaloSim simulation is added.

6.6 Possible improvements to FastCaloSim

The current version of FastCaloSim is the first implementation of the described method and is meant for testing if such a fast calorimeter simulation can be used to give satisfactory results. There are still a large number of improvements possible for different aspects.

6.6.1 Overall simulation

- The input samples used for the parametrisation (Section 6.2.2) were originally intended for an older fast simulation and were produced by a currently outdated version of the full simulation. Especially for the energy parametrisation an update is necessary.
- The particle simulation was started at the centre of the detector. It should be investigated, if a simulation starting at the end of the inner detector is better suited to separate the inner detector from the calorimeter simulation.
- The main difference visible between electrons and photons is a shift in the shower depth. Adding a parametrisation of this shift might give a significant improvement to the electron simulation.

- All hadrons, charged or neutral, are treated inside the calorimeter like charged pions. Adding more hadronic particles, e.g. charged kaons, protons or neutrals like K_L or neutrons, could improve the simulation but would also increase the parametrisation effort and the memory requirement by a large factor. For some hadron species, a simple energy rescaling factor might already give an improvement at very small run time costs.
- Especially for pions it might be possible to improve the simulation by subdividing the parametrisation sample and the simulation step into several categories that correspond to different categories of shower development. A possible choice of subcategories for pions could be to subdivide into the number of topoclusters found in the shower and then derive a parametrisation for these individual topoclusters instead of the whole pion.
- Particle parametrisations were performed for constant energy in a range from 200 MeV to 500 GeV (see Table 6.1). The upper limit of 500 GeV was determined by CPU time limitations for the full simulation of the input samples at high energies. At $\eta = 2.5$ an energy of $E = 500$ GeV corresponds to $p_T \approx 80$ GeV, at $\eta = 4.0$ it goes down to $p_T \approx 20$ GeV. Since ATLAS is in a hadron collider environment, very high particle energies can appear also in standard QCD events at large pseudorapidities and an extension of the energy parametrisation to higher energies should be done.

6.6.2 Energy parametrisation and simulation

- Muons are currently not simulated. Their energy deposit in the calorimeter is small (about 3 GeV), but plays a role in muon calorimeter isolation quantities (and possibly in the missing transverse momentum calculation, but no effect was seen up to now). For the muon shape it should be sufficient to deposit all energy in the hit cell. Concerning the energy, one could either produce a dedicated parametrisation or use the parametrisation of the estimated muon energy loss used in the muon track reconstruction.
- The energies in the individual calorimeter layers are currently parametrised as fractions of the total energy seen in all calorimeters (see Section 6.3.3). Alternatively, the energies in the individual calorimeter layers could be parametrised as fractions of the true particle energy. Both approaches have advantages and disadvantages, concerning how correlations are preserved:
 1. If the energies in the individual calorimeter layers are parametrised as fractions of the total energy seen in all calorimeters, an uncorrelated energy fluctuation in one calorimeter layer will also contribute to the total energy and change the energy fractions in all other layers.

For example, the correlation matrix for the energy fractions in Figure 6.11 (left) shows this as a strong anti-correlation between EMB1 and EMB2 for photons. On the other hand, correlated energy fluctuations from the shower development that affect the energies in several layers in a similar way automatically cancel in the fraction to the total energy.

2. One could avoid the propagation of uncorrelated fluctuations by parametrising the energies in the individual calorimeter layers as fractions of the true particle energy. In this case, an uncorrelated energy fluctuation in one layer would not influence the other layers, but still contribute to the total energy sum. In the end, it is necessary to force this energy sum to agree with the originally diced total energy, so the problem of item 1 reenters in the energy normalisation step and not via the correlation matrix. On the other hand, in this approach correlated energy fluctuations of all layers need to be described fully by the correlation matrix.

It was not tried to analyse in detail which approach is more accurate. Since uncorrelated fluctuations are mostly Gaussian, they can be well described by the correlation matrix. Correlated fluctuations come from the intrinsic shower development and are more complicated, so absorbing them to some degree in the fraction of energy seemed to be the better approach. The remaining part needs to be handled by the correlation matrix together with the uncorrelated fluctuations.

For a future parametrisation it would be of advantage if the calorimeter layer energy resolution, which contributes to a large part to the uncorrelated energy fluctuations, could be distinguished from the correlated fluctuations induced by the shower development (if the true energy deposited in a calorimeter was known, not only the measured energy). This would allow a parametrisation of the correlated true energy depositions to which the uncorrelated fluctuations could simply be added.

- Currently, the calibration constants that belong to one energy parametrisation are also used during the simulation to go back to the calorimeter raw response. If the energy of the particle to be simulated is somewhere between the energies for which the parametrisation were produced, an interpolation of the calibration constants might give an improved energy behaviour. As can be seen in Figures 6.3 and 6.4, a simple logarithmic interpolation of the calibration constants should be sufficient, but the interpolation needs to be handled with care if the constants have large uncertainties.

6.6.3 Shower shape parametrisation and simulation

- For a given particle energy and direction, shower shapes are parametrised in each calorimeter layer as average shapes only depending on the shower depth. For photons this seems to be a well suited variable. For pions it should be investigated if the point of the first hadronic interaction could be used as additional or alternative parameter.
- The impact point is calculated with respect to the longitudinal centre of each calorimeter layer. Especially for non-pointing particles, the shower position might be more accurate if the energy weighted average longitudinal depth in each layer would be used for the extrapolation.
- Uncorrelated energy fluctuations are included and they modify the average shower shape. The size of these fluctuations was taken as one energy dependent function for each calorimeter technology (see step 3 in Section 6.4.3). The actual uncorrelated fluctuations are most likely larger and extracting them from the fully simulated samples might give improved results.
- Correlated shape fluctuations are not included in the current model and it will not be easy to do a parametrisation for that.
 1. One possibility could be to assume that correlated fluctuations arise from subshowers that deposit a lot of energy in localised areas. One could try to simulate at least the first few steps of the shower history, similarly as it is done in the full simulation. The drawbacks of this approach are that particle multiplicities will increase by large factors and energy/shape parametrisations would be needed for particles produced deep inside the calorimeter.
 2. An empiric approach could be to try to build localised correlation matrices in bins of the distance to the impact point that describe the correlations of fluctuations in cells, depending on their distance to each other.

The first approach will slow the simulation down and will require physics input similar to Geant 4 or directly using Geant 4. The second approach could be almost as fast as the current model, but it needs to be shown that it can give a visible improvement to the current model after reconstruction.

- No shower shapes were fitted for the forward calorimeter (FCAL) and the complete energy is deposited in the hit cells. The reason is that the cells in the forward calorimeters have a more complicated x-y-z geometry instead of the simply η - ϕ -r/z projective geometry of the cells in other calorimeter layers. It should be possible to add shower shapes for the FCAL with some modifications to the cell integration.

7 Validation of the fast calorimeter simulation

In the previous chapter the method of the fast calorimeter parametrisation was explained and some comparisons of the internal consistency of the parametrisation with the full simulation were shown. However, these comparisons focused only on the variables and distributions that are included in the parametrisation.

During the object reconstruction in the ATLAS calorimeter, a lot of shape and energy distribution variables are calculated and used to identify objects and reconstruct their energy. In order to be universally usable, the fast calorimeter simulation needs to be able to reproduce also distributions that are not directly included in the parametrisation. The validation to ensure that the fast calorimeter parametrisation and simulation is usable for the simulation of physics processes is performed in two steps:

- The first step is to compare the fast calorimeter simulation to the full simulated single photon and charged pion samples from Athena release 10, that were used as input to the parametrisation, and to di-jet events simulated with the same Athena version. This validation step was carried out within the ATLAS collaboration and is documented in Reference [30].
- In the second step the fast detector simulation Atlfast-II is used to validate the performance of physics objects (electrons, muons, jets, missing E_T , etc.) against the full simulation of Athena release 14. The Atlfast-II simulation package is composed of the full simulation of the inner detector and the fast simulation FastCaloSim for the calorimeter system. For muons, the full simulation is also used in the calorimeter and the muon system. This validation step was also carried out within the ATLAS collaboration and is documented in Reference [31].

7.1 Validation against the input samples used for the parametrisation

The parametrisation of FastCaloSim is based on a rather old version of the ATLAS full simulation as included in Athena release 10.0.1. The first validation step is to do an internal validation by comparing photons and charged pions between

FastCaloSim and the input samples that were used to produce the FastCaloSim parametrisation.

The fast and full simulation samples were privately produced without electronic noise. Furthermore only events without interactions in the inner detector are used. This allows to compare the fully simulated input events from Athena version 10.0.1 to the parametrised fast simulation to study intrinsic differences of the parametrisation model. In order to identify which differences seen between FastCaloSim and the current full simulation are due to improved geometry and Geant 4 physics lists descriptions, a comparison to the full detector simulation of single photons in Athena release 13 without electronic noise is also shown. As far as this comparison is concerned, there are no differences between the full simulation in Athena release 13 and 14.

7.1.1 Photon response and shape variables

For the first validation step, two sets of photon samples with a photon energy of 20 GeV within $0.2 < \eta < 0.25$ and an energy of 100 GeV within $2.00 < \eta < 2.05$ are used.

Typical variables used to identify photons [15] are shown in Figures 7.1 and 7.2 for the 20 GeV and the 100 GeV photons, respectively. All distributions contain only photons without conversions in the inner detector to avoid a bias due to changes in the description of the material in the inner detector between Athena release version 10 and 13.

- The distribution of the total uncalibrated energy in Figures 7.1(a) and 7.2(a) shows a fair agreement between FastCaloSim and the old version of the ATLAS full simulation. There is a shift to higher reconstructed energies for the new version of the ATLAS full simulation. This shift can be traced back to updated sampling factors in the new full simulation.
- The energy fraction F1 in the first layer of the electromagnetic calorimeter (see Figures 7.1(b) and 7.2(b)) is well described by the fast simulation.
- However, differences are visible for the energy fraction F3 in the third calorimeter layer (see Figures 7.1(c) and 7.2(c)). FastCaloSim deposits no energy in a calorimeter layer if the average energy in this layer is below 100 MeV. Thus photons with very small energy fractions in the third layer are simulated with F3=0.
- Figures 7.1(d) and 7.2(d) show the shower width $\omega_{3 \text{ strips}}$ in the η direction. $\omega_{3 \text{ strips}}$ is calculated from the three central cells in the photon cluster in the first layer of the electromagnetic calorimeter. The distribution of the fast simulation is in good agreement to the old full simulation, but a bit narrower

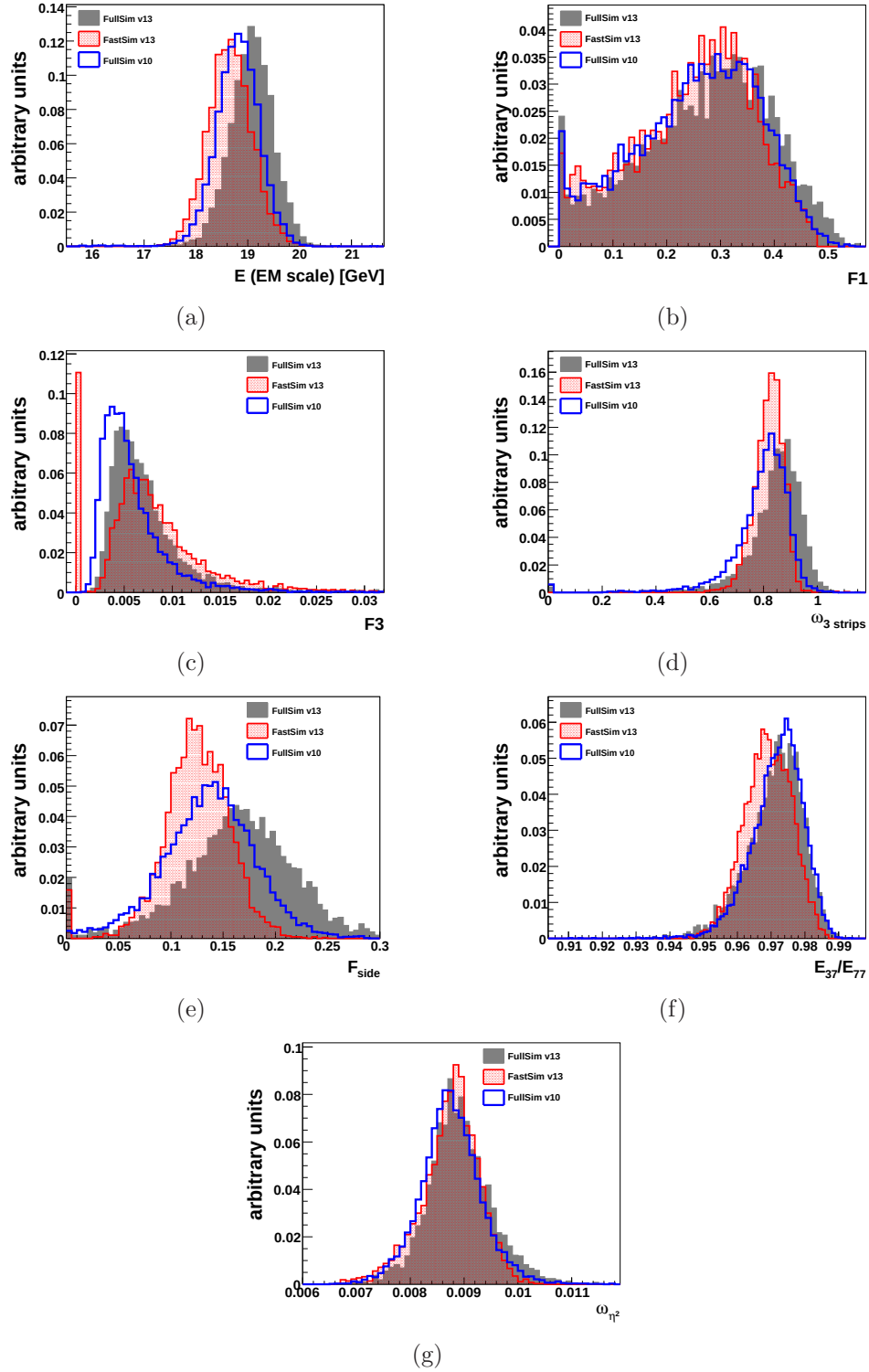


Figure 7.1: Comparison between full (Athena release v10 and v13) and fast simulation for unconverted photons with an energy of 20 GeV and $0.2 < \eta < 0.25$ [30]. See the text for a discussion of the variables.

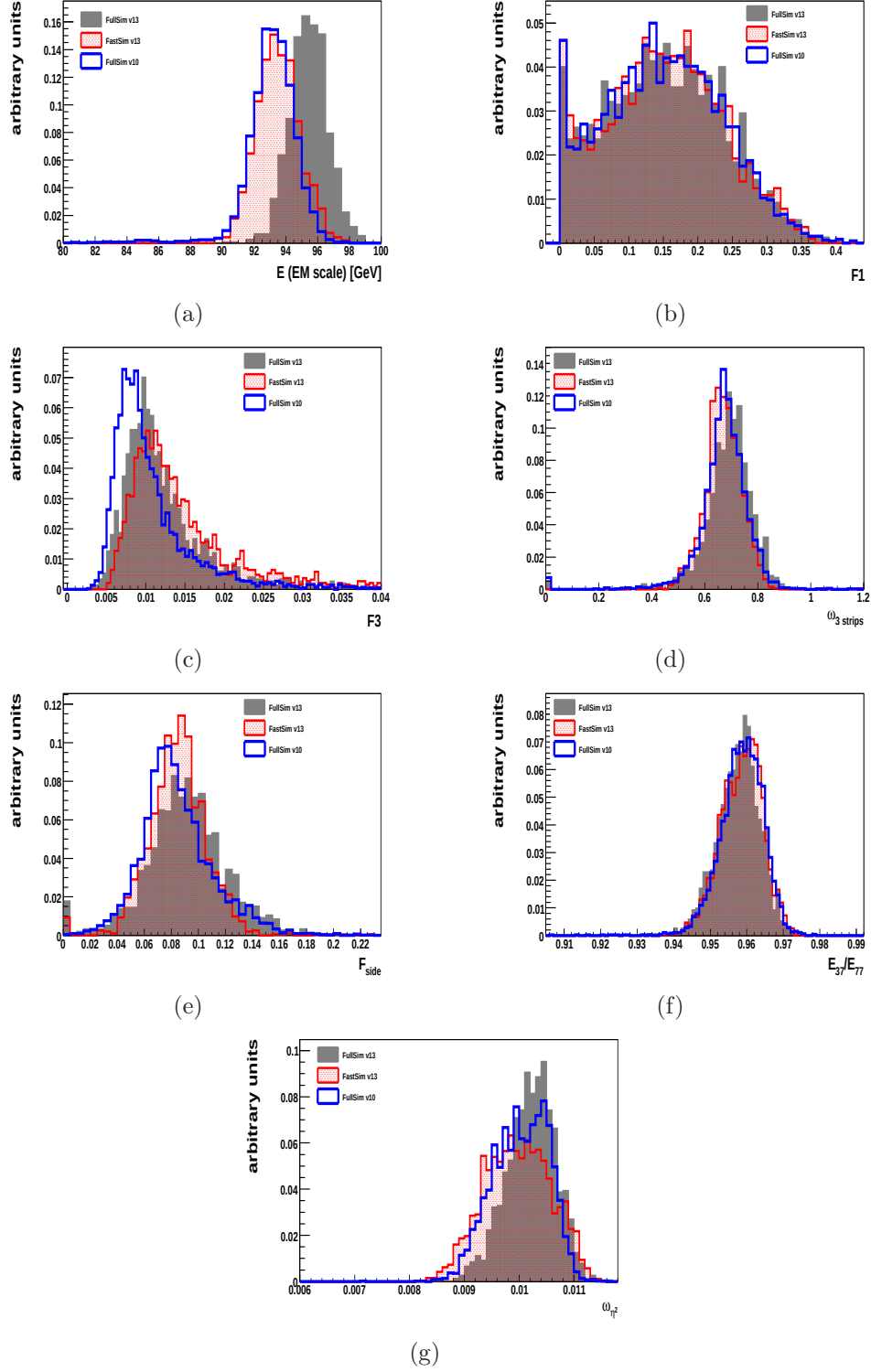


Figure 7.2: The same as Figure 7.1 for photons with an energy of 100 GeV and $2.0 < \eta < 2.05$ [30].

for 20 GeV photons due to missing intrinsic shower shape fluctuations in the fast simulation.

- The related variable F_{side} , calculated from the energy fraction outside the central strip cells in the first calorimeter layer, is shown in Figures 7.1(e) and 7.2(e). Due to the missing fluctuations in the fast simulation shower shape, FastCaloSim has a lower energy fraction outside the central cells. However, the agreement to the old full simulation is still acceptable, while the new full simulation is shifted towards high energy fractions.
- The lateral shower shape and width in the second layer of the electromagnetic calorimeter are described by the variables $R_\eta(37) = E_{37}/E_{77}$, the ratio of the energy in $3 \cdot 7$ cells in $\Delta\eta \cdot \Delta\phi$ to the energy in $7 \cdot 7$ cells (see Figures 7.1(f) and 7.2(f)), and ω_{η^2} , the width in the η direction in the 2nd EM calorimeter (see Figures 7.1(g) and 7.2(g)). The fast simulation shows a more symmetric shower shape distribution and a peak shifted to lower values compared to both full simulation versions. However, the overall agreement is good within the total width of the distribution.

In general the fast simulation shows reasonable agreement to the old version of the full simulation. In the cases where differences are seen, the distributions of the fast simulation are shifted towards more photon like values and hence it can be expected that the fast simulation efficiencies are overestimated. The remaining efficiency overestimation needs to be corrected probabilistically.

The differences are caused by the simplified description of shower shapes in FastCaloSim, especially due to the missing lateral shower shape fluctuations. Currently no parametrisation model exists that is able to implement these fluctuations without a severe increase in CPU time during the simulation. However, the comparison to the new full simulation shows that the more up to date Geant 4 version and the improved detector geometry description is causing differences of the same size as the simplifications in the fast simulation.

7.1.2 Charged pion response and shape variables

The largest fraction of charged pions in physics events appears inside jets and is not isolated, hence jet distributions are not very sensitive to the distribution of pions. However, isolated high p_T pions appear in hadronic τ -lepton decays and thus the hadronic τ identification variables [15] are most suited to compare the charged pion simulation between full and fast simulation. For the validation, samples of charged pions simulated with an energy of 20 GeV within $0.2 < \eta < 0.25$ and an energy of 100 GeV within $2.00 < \eta < 2.05$ are used.

Figures 7.3 and 7.4 show calorimeter specific τ identification variables. All distributions contain only pions without hard interactions in the inner detector

to avoid a bias due to changes in the description of the material in the inner detector between Athena release 10 and 13.

- The reconstructed transverse momentum p_T of the τ shows a good agreement to the old version of the full simulation.
- The transverse energy weighted width $\Delta\eta$ in the η direction in the first calorimeter layer is narrower for the fast simulation. However, this distribution is very sensitive to calorimeter noise and the differences disappear once noise is added.
- Due to the missing fluctuations in the lateral pion shape the energy weighted shower radius in the EM calorimeter R_{em} is narrower in the fast simulation.
- For the same reason the isolation energy ΔE_T^{12} in a hollow cone¹ $0.1 < \Delta R < 0.2$ around the centre of the τ is lower in the fast simulation than in both full simulations.

The last three variables show small difference between the full and fast simulations due to the missing shower fluctuations. However, similarly as it was for the photon identification, these differences cause a shift of the τ identification variables towards more τ like values and hence the τ identification efficiency is expected to be higher in fast simulation than in full simulation.

The photon identification variables from the previous section are intended to separate electrons and photons on the one hand from jets on the other hand. Jets that are candidates for a misidentification as electrons are often dominated by a single charged hadron. Therefore a comparison between the fast and full simulation of the electron and photon identification variables calculated for charged pions can give an idea how well the rejection of jets is described in the fast simulation.

In Figures 7.5(a)-(f) and 7.6(a)-(f) the comparison between the fast and full simulation is shown using the photon validation from the previous section.

- The fraction of energies in the first (see Figures 7.5(a) and 7.6(a)) and third (see Figures 7.5(b) and 7.6(b)) layer of the electromagnetic calorimeter are well described by FastCaloSim.
- The lateral shower variables (see Figures 7.5(c)–7.5(f) and Figures 7.6(c)–7.6(f)) show large differences between the fast and full simulations. While pion showers have large fluctuations from the hadronic shower development, the fast simulation deposits the energy using an average lateral shower profile.

¹The cone is defined in the boost invariant radius $\Delta R = \sqrt{(\Delta\eta)^2 + (\Delta\phi)^2}$

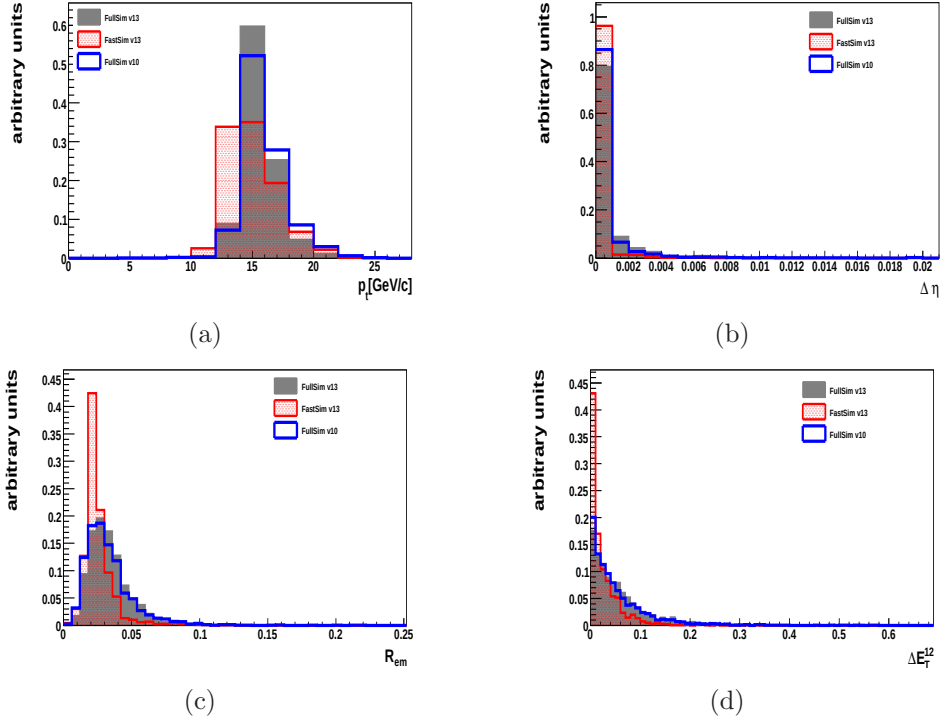


Figure 7.3: Comparison between full (Athena release v10 and v13) and fast simulation for charged pions with an energy of 20 GeV and $0.2 < \eta < 0.25$ [30]. See the text for a discussion of the variables.

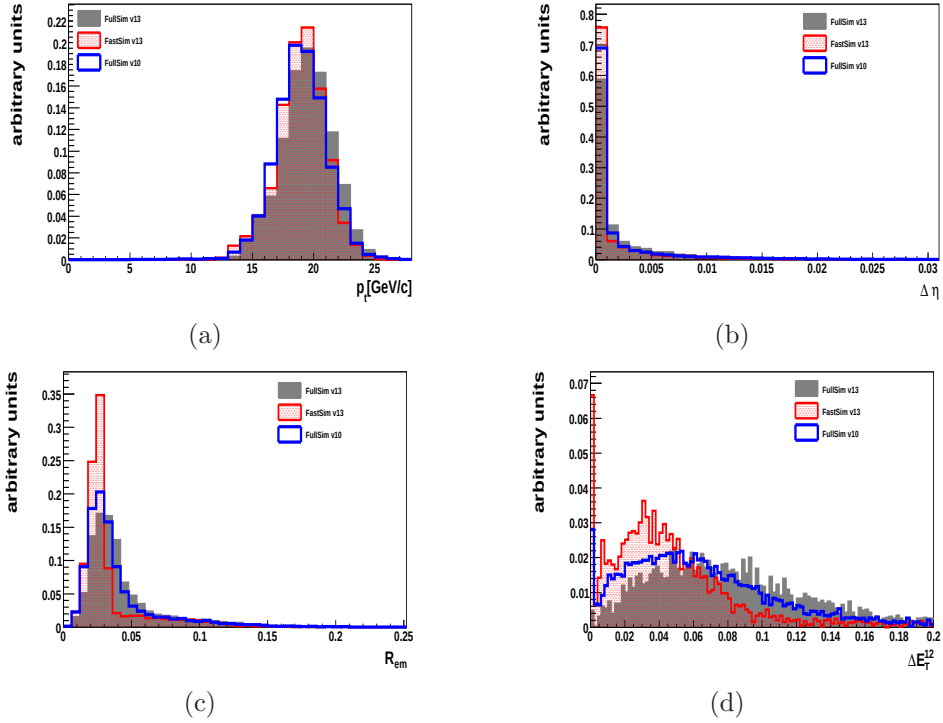


Figure 7.4: The same as Figure 7.3 for charged pions with an energy of 100 GeV and $2.0 < \eta < 2.05$ [30].

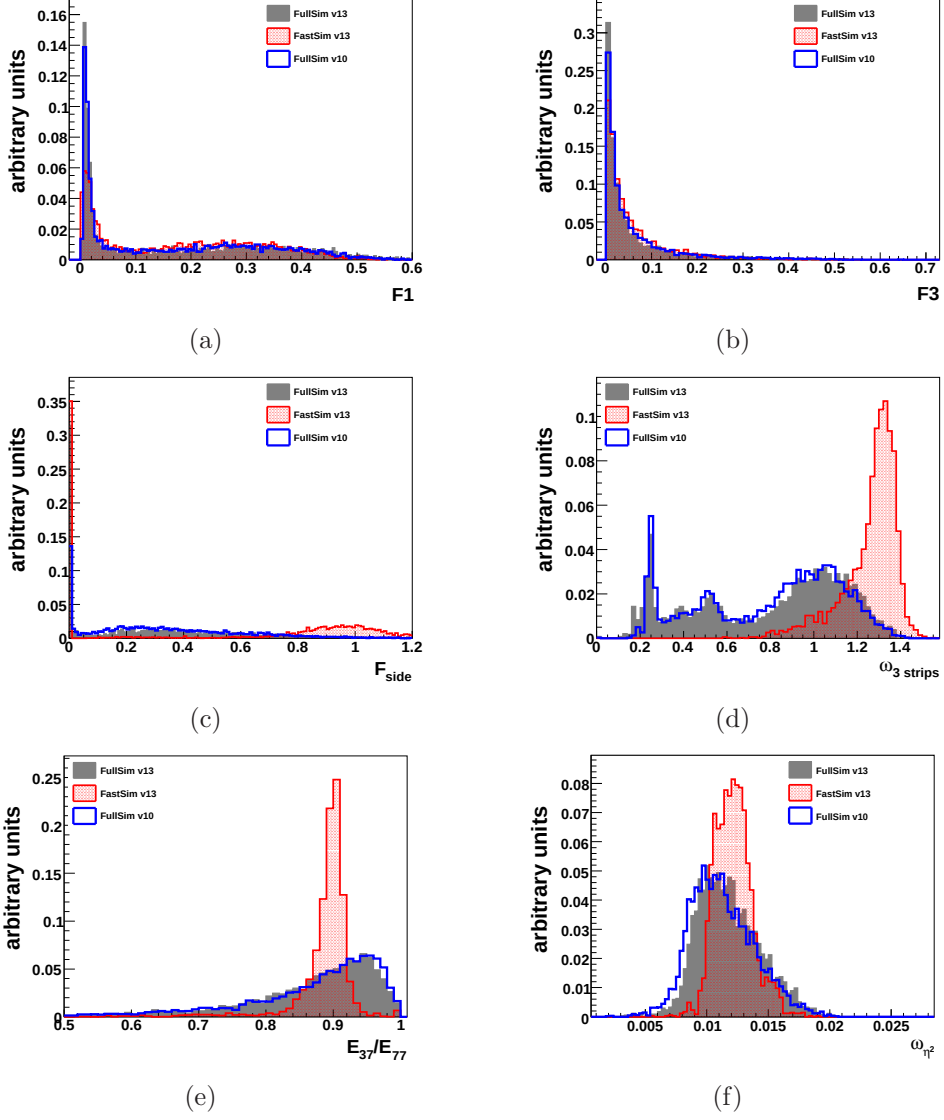


Figure 7.5: The same as Figure 7.1 for charged pions with an energy of 20 GeV and $0.2 < \eta < 0.25$ [30].

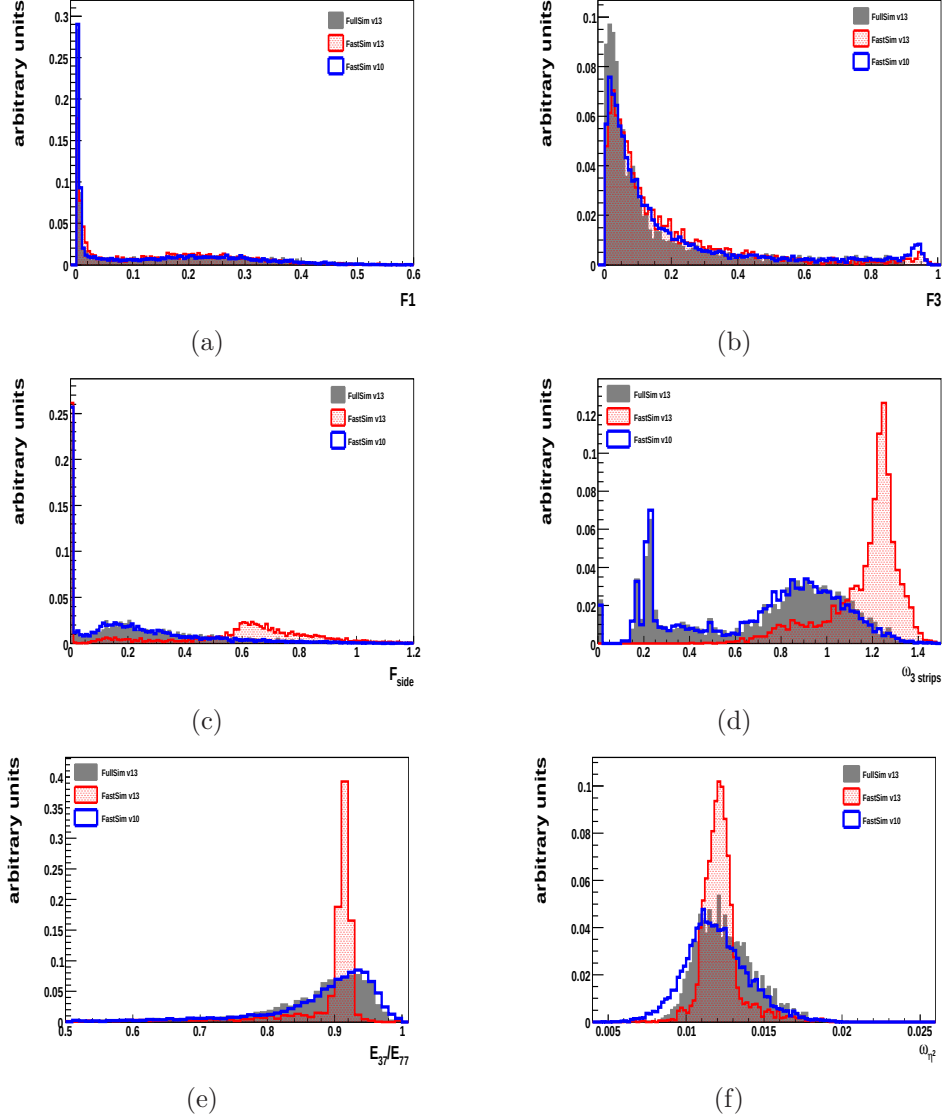


Figure 7.6: The same as Figure 7.1 for charged pions with an energy of 100 GeV and $2.0 < \eta < 2.05$ [30].

For $R_\eta(37) = E_{37}/E_{77}$ and ω_{η^2} the average shape gives a good description of the mean, but without the large spread due to the fluctuations. However, in the 1st calorimeter layer the energy is distributed broadly over many cells by the fast simulation, while a narrow fluctuation in the pion shower is usually picked up in the reconstruction of the full simulation and used to calculate the variables of the 1st layer. This leads to a peak of the $\omega_{3 \text{ strips}}$ distribution at $\sqrt{2}$ in the fast simulation and more energy being deposited outside the central region in the case of F_{Side} .

The fast calorimeter simulation is designed to give a good description of average particle properties. If reconstructed variables are sensitive to fluctuations, as it is the case for the misidentification of a charged pion as an electron or photon, the simulation cannot perform well. However, for most variables at least the mean of the distribution is described well. Hence it can be expected that the fast simulation is able to describe correctly the rejection of the largest fraction of jets, which are composed of multiple neutral and charged particles, and fails only for the remaining small fraction of jets, that are dominated by a single charged hadron.

7.1.3 Validation of jet response and efficiency

The validation of jets has been done considering mainly the following three quantities:

- The transverse energy response.
- The transverse energy resolution.
- The reconstruction efficiency.

Jets in FastCaloSim are reconstructed using the same algorithms and software implementation used for the full simulation. The most ATLAS-wide used jet algorithm (iterative cone with a cone size $\Delta R = 0.4$) has been used for the comparison. The input to the jet reconstruction has been chosen to be calorimeter towers.

To evaluate the performance of the jet reconstruction, the reconstructed jets are associated to particle jets (that is, to jets reconstructed from the truth final state particles using the same jet reconstruction algorithm) based on the angular distance ΔR . If a reconstructed jet has a particle jet within a distance of $\Delta R < 0.2$, the two are associated together.

Jet Response at the EM scale

If jets obtained from the FastCaloSim parametrisation are compared to consistent full simulation events, they are expected to have an energy response (EM energy in the following) very close to that of the full simulation. However, the jet EM energy is expected to be non-linear and different from the energy of the originating particle jet, which is due mainly to the fact that the ATLAS calorimeters are not compensating and also because of the presence of dead material. The jet energy is then corrected to obtain the particle jet energy.

The first validation step is to compare the jet EM energy for FastCaloSim and the full simulation. The comparison is done with datasets produced with Athena release 10.0.4. For this study QCD di-jet events are used. Events are generated in bins of the partonic p_T in the range $17 \text{ GeV} < p_T \lesssim 1 \text{ TeV}$.

The particle jet phase space is divided into bins of the transverse energy and broad pseudorapidity bins. For each bin, a histogram of $E_{T\text{rec}}/E_{T\text{mc}}$ is made, where $E_{T\text{mc}}$ is the reconstructed energy of the particle jet and $E_{T\text{rec}}$ is the transverse energy of the associated reconstructed jet. Non-associated jets are not considered. Each histogram is fitted with a Gaussian function once to obtain a pair $(\mu_{\text{raw}}, \sigma_{\text{raw}})$ for the mean value and the Gaussian sigma of the jet energy response function. Then it is fitted a second time (again with a Gaussian function) in the interval $[\mu_{\text{raw}} - 2\sigma_{\text{raw}}, \mu_{\text{raw}} + 2\sigma_{\text{raw}}]$, to obtain the final values for the energy response μ and resolution σ that are used in the comparison.

Figure 7.7 shows a summary of the μ values as a function of the particle jet E_T in four different pseudorapidity bins ($0 < |\eta| < 0.7$, $0.7 < |\eta| < 1.5$, $1.5 < |\eta| < 2.5$ and $2.5 < |\eta| < 3.5$). The agreement between FastCaloSim and the full simulation is very good. Even in the most critical calorimeter crack region ($2.5 < |\eta| < 3.5$), the maximum disagreement is found for the lowest E_T bin, and it is about 5%. Hence, FastCaloSim well reproduces the jet EM energy.

Figure 7.8 shows a summary of the transverse energy resolution σ/μ as a function of the particle jet E_T in the same four pseudorapidity bins. Clearly FastCaloSim is underestimating the jet energy fluctuations in the calorimeter. In the central region, the values of the resolution are $\sim 10\%$ lower for FastCaloSim than in full simulation. The disagreement increases at higher pseudorapidity.

Jet Reconstruction Efficiency

After the association procedure defined above, the jet reconstruction efficiency in full simulation and in FastCaloSim can be compared. Again, the particle jet phase space is divided in bins of transverse energy and pseudorapidity. For each bin, the jet reconstruction efficiency is defined as the number of associated particle jets divided by the total number of particle jets.

The efficiency as a function of the particle jet E_T are shown in Figure 7.9 for the full simulation and for FastCaloSim for different pseudorapidity bins. The

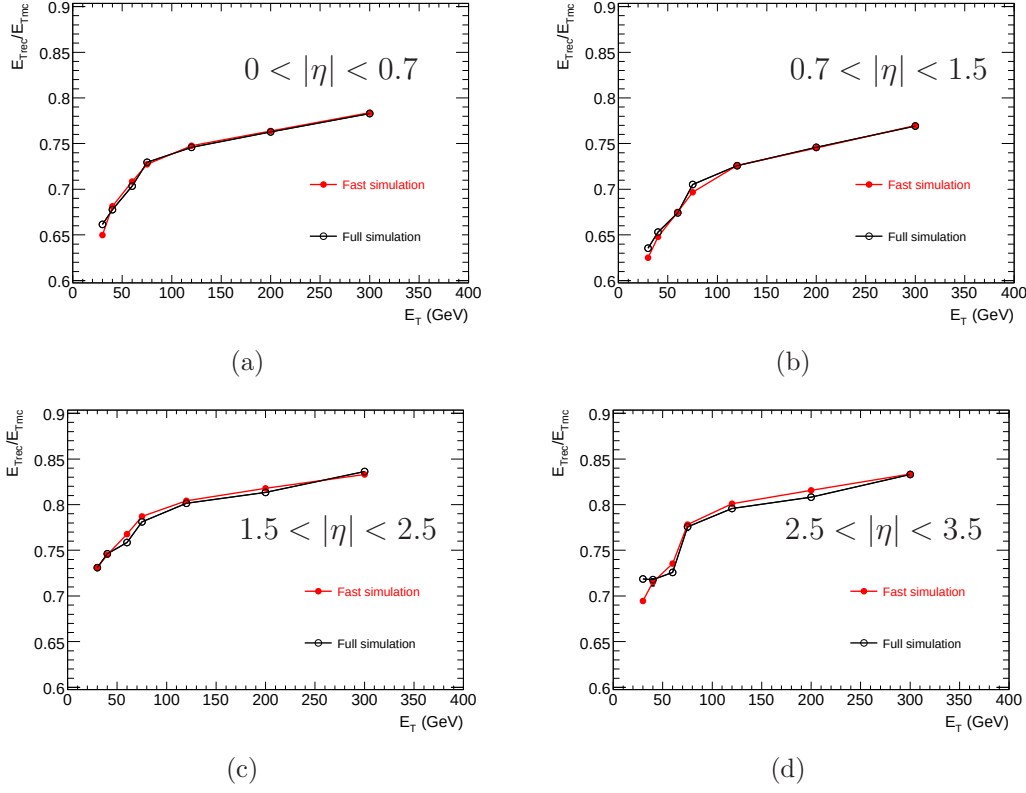


Figure 7.7: Comparison of the jet EM energy scale between full simulation and FastCaloSim in four different pseudorapidity regions [30].

fact that the efficiency does not fully saturate to 1 at moderate E_T in the central region is a consequence of the association cut $\Delta R < 0.2$. The angular resolution for the jet reconstruction depends on the jet energy (the higher the jet energy, the more precise the angular direction). In the forward calorimeter regions, for equivalent E_T , the jet energy is higher by a factor $\cosh \eta$ and hence the efficiency approaches 1 at high $|\eta|$.

The overall agreement of the reconstruction efficiency is good. The most relevant discrepancies are found at low E_T in the region $1.5 < |\eta| < 2.5$ (where the FastCaloSim efficiency is about 30% higher with respect to the full simulation) and at relatively low E_T in the $2.5 < |\eta| < 3.5$ calorimeter crack region, where the FastCaloSim is 15% lower than the full simulation efficiency.

Jet Response at the HAD scale

As mentioned previously, jets in full simulation are corrected mainly for calorimeter non-compensation and dead material using a cell weighting procedure à la H1, described in details in Reference [15]. The main principle is that a number of cell weights are determined minimizing the resolution of the reconstructed jet energy

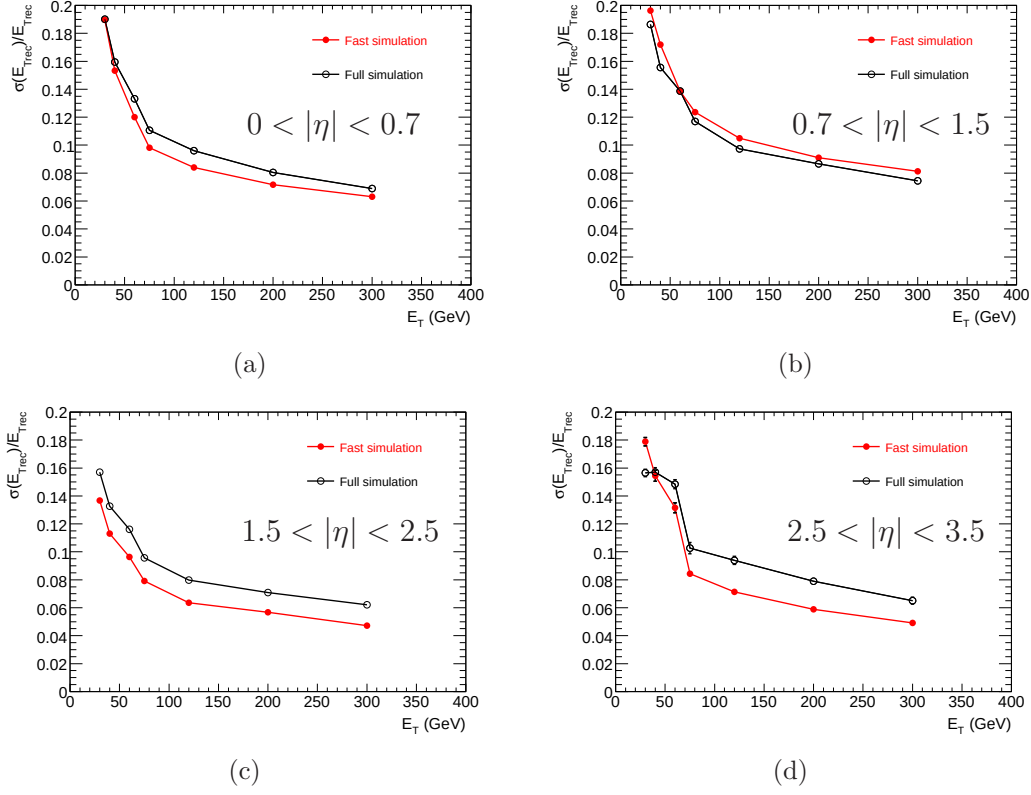


Figure 7.8: Comparison of the jet EM transverse energy resolution between full simulation and FastCaloSim in four different pseudorapidity regions [30].

with respect to the corresponding particle jet energy.

The jet corrections for full simulation are recomputed each time there is a major change in the detector description or in the Geant 4 physics lists, while the FastCaloSim parametrisation is based on the old Athena version 10. Therefore, dedicated jet constants have been computed for FastCaloSim [30], following exactly the same procedure that is used in full simulation to decouple the FastCaloSim jet corrections from the full simulation ones.

The jet linearity and resolution for FastCaloSim as a function of the jet transverse energy after the application of the global H1 corrections are shown in Figure 7.10 for three different pseudorapidity regions. The linearity is recovered with the largest deviation of the order of 2% at low E_T . Even in the most problematic calorimeter crack region, the linearity is within 2% on a large E_T range.

The resolution is improved with respect to jets at the EM scale. The resolution curves are fitted with a function of the form

$$\frac{\sigma(E_T)}{E_T} = \frac{a}{\sqrt{E_T}} \oplus b \oplus \frac{c}{E_T} \quad (7.1)$$

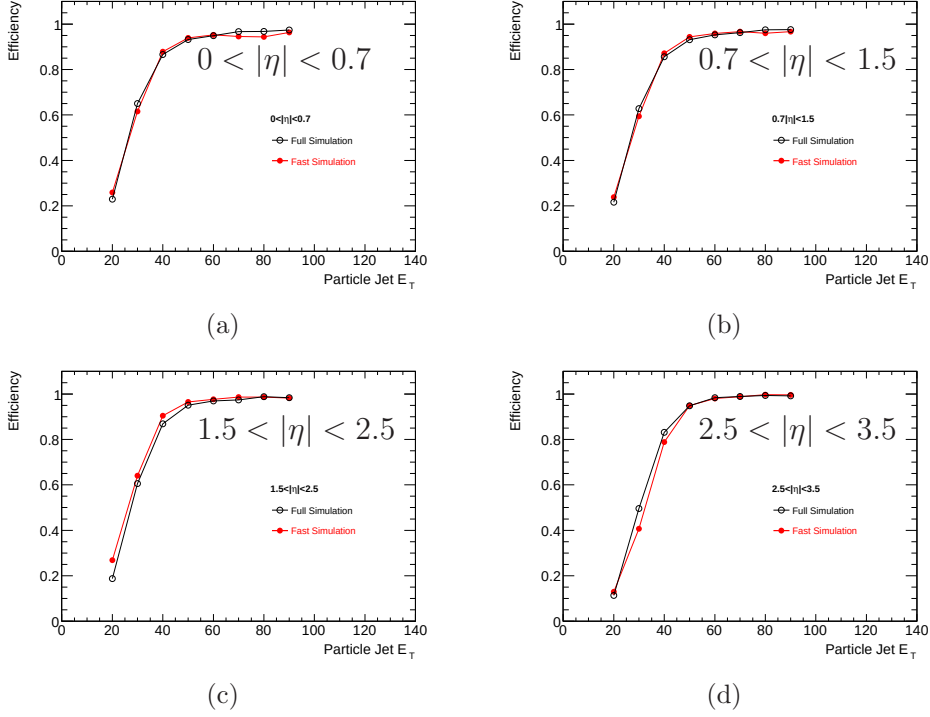


Figure 7.9: Comparison of the jet reconstruction efficiency between full simulation and FastCaloSim in four different pseudorapidity regions [30].

where the transverse energy is measured in GeV. The results of the fits in three pseudorapidity regions are shown in the right plot of Figure 7.10. Figure 7.11 (left) shows the jet linearity as a function of the pseudorapidity for three energy bins. Again, the maximum deviation from linearity is of the order of 2%. The right plot of Figure 7.11 shows the jet E_T resolution as a function of the pseudorapidity.

For Athena release 10 the H1 calibration of jets in full simulation events was not performing well. Therefore, a comparison of the jet response and resolution after the H1 calibration is shown in Section 7.2.3 for Athena release 14.

7.1.4 Summary

In summary, the comparison of FastCaloSim to the old full simulation in Athena release 10 shows that the reconstructed energy response is described well. Differences appear for lateral shape variables due to the missing lateral shower shape fluctuations. However, the direction of these differences are all such, that object identification efficiencies are expected to be higher in the fast simulation than in the full simulation.

The comparison to the new full simulation shows that the energy response

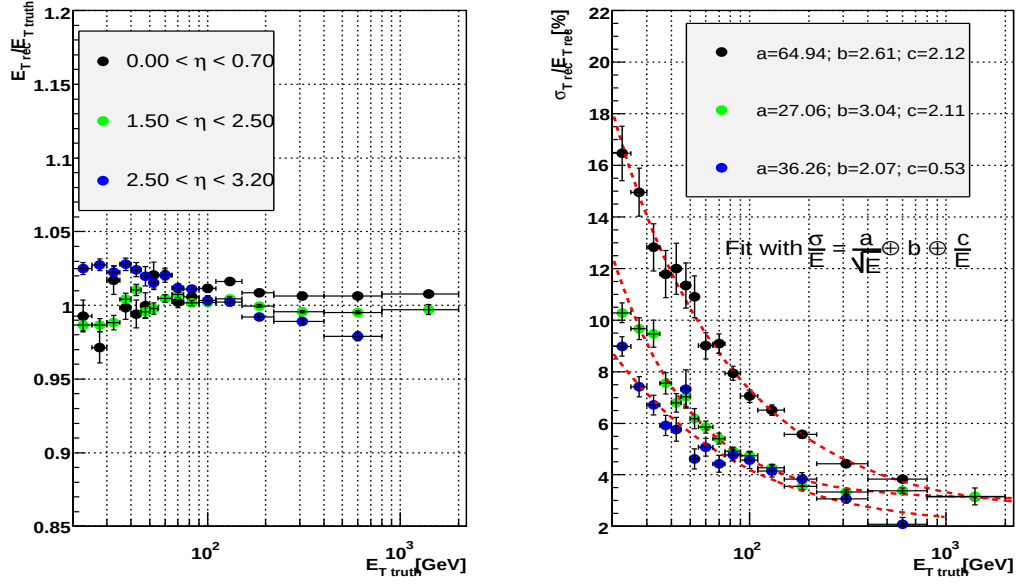


Figure 7.10: Jet linearity (left) and jet resolution (right) as a function of the jet transverse energy in three pseudorapidity bins [30].

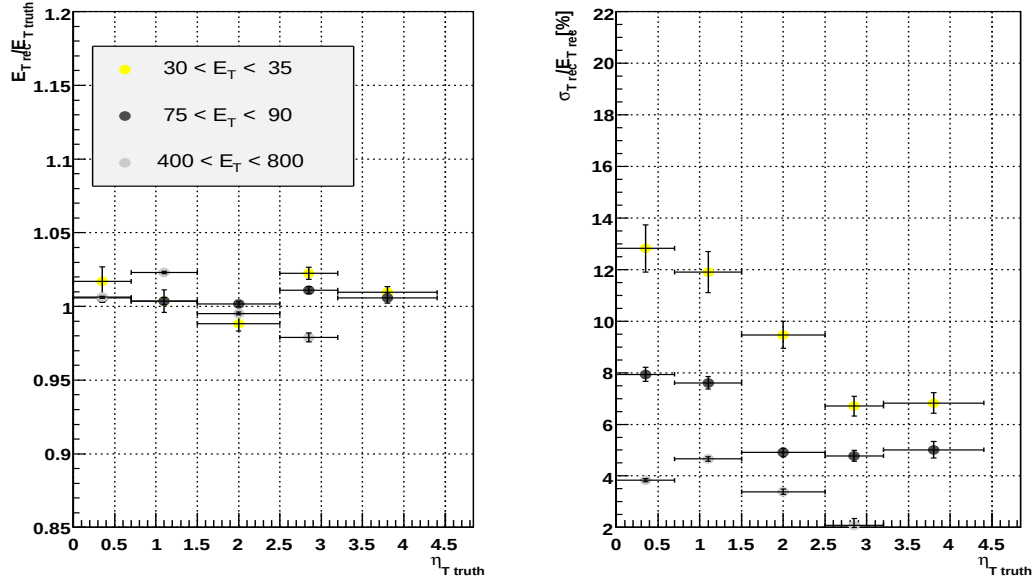


Figure 7.11: Jet linearity (left) and jet resolution (right) as a function of the jet pseudorapidity for three transverse energy bins [30].

function for photons and charged pions is changed due to intrinsic changes in the full simulation since Athena release 10. Also the lateral shower shape distributions have changed and therefore increased reconstruction and identification efficiency differences can be expected between FastCaloSim and new versions of the ATLAS full simulation.

7.2 Performance of the fast simulation Atlfast-II

The fast simulation FastCaloSim is used for the simulation of the calorimeter system in the Atlfast-II detector simulation package in combination with a full simulation of the inner detector and a full simulation of the calorimeter and muon system for muons. The simulation of the complicated ATLAS calorimeter system is by far the most time consuming element in the full simulation and therefore the use of Atlfast-II can reduce the total event simulation time by a factor of 10–20 [27].

For the validation of Atlfast-II, full and fast simulation Monte Carlo samples generated with Athena version 14 are used. In the previous section it was shown that some differences between the fast and full simulation are expected, on the one hand because of intrinsic limitations of the fast calorimeter parametrisation and on the other hand because of differences between the full simulation of Athena version 14 and the full simulation, that was used to produce the parametrisation.

The expected differences in the calorimeter energy response of the fast simulation are corrected by using dedicated electron, photon and jet energy calibration constants. The differences in efficiencies are corrected by using object specific reconstruction and identification efficiency correction factors as function of the object p_T and η . In the following comparisons, this corrected version of Atlfast-II is also shown wherever it is appropriate.

In the Chapters 8–10 Atlfast-II Monte Carlo samples are used in the context of the $gg \rightarrow H \rightarrow W^-W^+ \rightarrow l^-\bar{\nu} l^+\nu$ analysis. Therefore only the detector objects that are relevant for this analysis are discussed in this chapter. The validation of Atlfast-II for all detector objects and several physics benchmark analysis is discussed in Reference [31].

7.2.1 Performance of electrons

Electrons are reconstructed in the ATLAS detector from a combination of calorimeter and inner detector track measurements [15]. Three sets of standard ATLAS electron identification selection cuts are defined in the “isEM” electron selection: “loose”, “medium” and “tight”. These selection cuts give an increasing level of rejection against jets misidentified as electrons at the price of a reduced efficiency.

The electron selection efficiency is shown in Figure 7.12 as function of the true electron transverse momentum p_T and pseudorapidity η for loose, medium and

tight electrons. Without corrections, the medium and tight electron identification efficiency is overestimated by $\approx 5\%$ for electrons in the low p_T and $1 \lesssim |\eta| \lesssim 1.7$ regions. After the efficiency correction, the agreement between full and fast simulation is within 1–2% for all η and p_T regions.

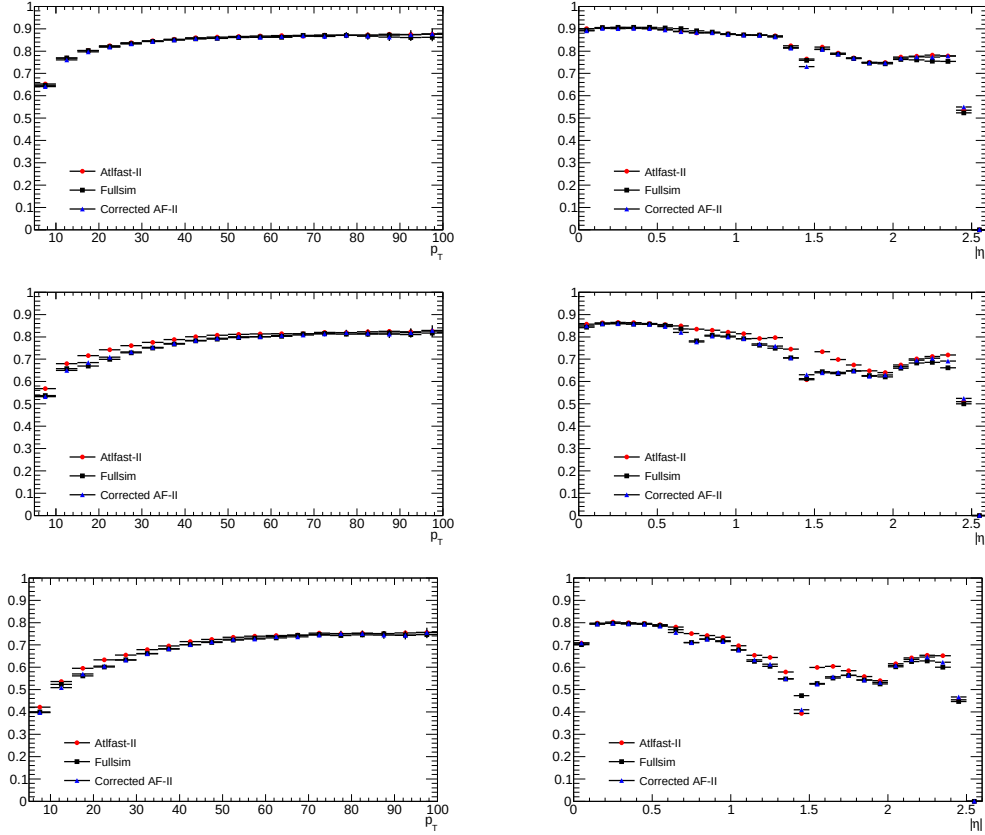


Figure 7.12: The electron identification efficiencies as a function of p_T (left) and $|\eta|$ (right) [31]. The standard loose (top), medium (middle) and tight (bottom) *isEM* selections are applied to identify electrons from a $Z \rightarrow ee$ sample. For the plots showing the η dependency a p_T cut of 20 GeV has been applied. The results obtained from full simulation (black), Atlfast-II (red) and corrected Atlfast-II (blue) are superimposed.

The electron energy response and resolution is shown in Figures 7.13 and 7.14 for medium electrons. Due to the dedicated electron calibration, the energy response and resolution is in very good agreement to the full simulation. For true electrons, the *isEM* electron selection cuts have almost no influence on the

energy reconstruction and hence the comparison of the response and resolution for loose and tight electrons shows the same results.

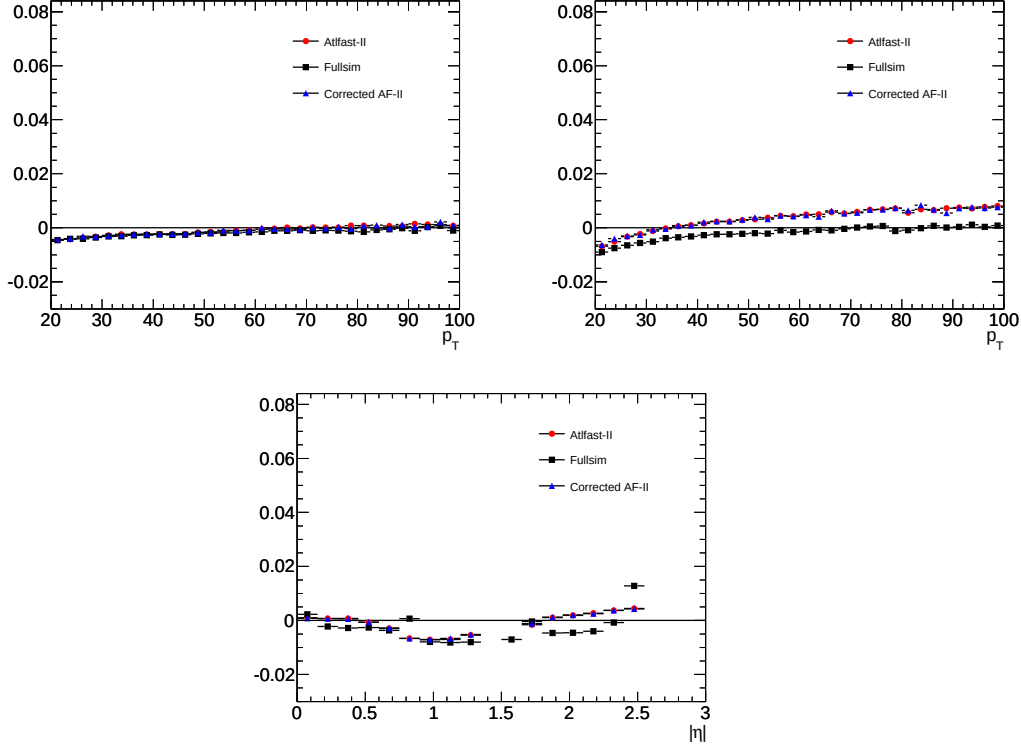


Figure 7.13: The linearity $(E_T^{\text{rec}} - E_T^{\text{true}})/E_T^{\text{true}}$ for electrons from $Z \rightarrow ee$ events as function of p_T in the barrel region ($|\eta| < 1.3$) (upper left) and endcap region (upper right) and as a function of $|\eta|$ (lower plot) [31]. The standard medium *isEM* selection is applied to identify electrons from a $Z \rightarrow ee$ sample. The results obtained from full simulation (black), Atlfast-II (red) and corrected Atlfast-II (blue) are superimposed.

The misidentification efficiency for a jet being identified as an electron is shown in Figure 7.15 as function of the electron candidate p_T and $|\eta|$ for loose, medium and tight selection cuts. The loose selection does not apply cuts on the lateral shower shape in the first electromagnetic layer, where the biggest differences between the fast and full simulation was seen in Section 7.1.1. Hence a good agreement is visible between the full and fast simulation.

For medium and tight selection cuts the cluster shape in the first layer is used for the electron identification and the misidentification efficiency is therefore underestimated by a large factor. However, it is possible to correct for this under-

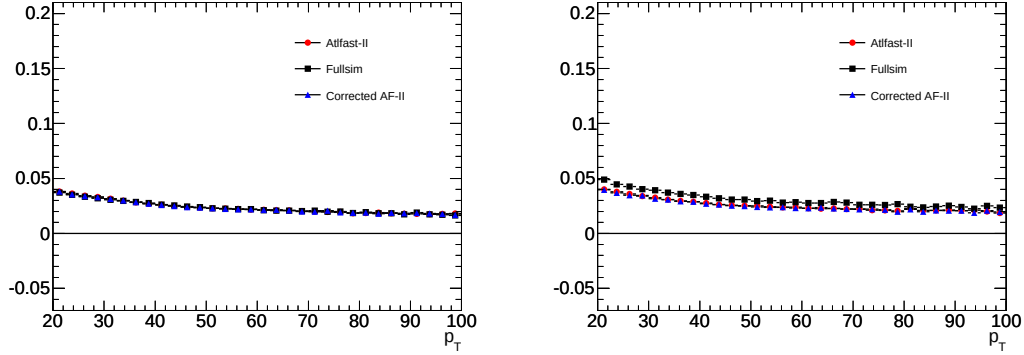


Figure 7.14: The energy resolution as function of p_T for electrons from $Z \rightarrow ee$ events for the barrel (left) and endcap (right) [31]. The standard medium $isEM$ selection is applied to identify electrons from a $Z \rightarrow ee$ sample. The results obtained from full simulation (black), Atlfast-II (red) and corrected Atlfast-II (blue) are superimposed.

estimation of the misidentification efficiency. Although comparisons to different samples indicate that this correction of the misidentification efficiency is rather universal, differences in the order of 20-50% between the full and fast simulation after corrections can still be expected.

7.2.2 Performance of muons

True muons are simulated in Atlfast-II with the full simulation. Therefore, no differences between the full and fast simulation are expected in the muon reconstruction using the “STACO” algorithm [15], that combines statistically the track measurement in the inner detector with the independent track measurement in the muon system.

Figure 7.16 shows that the muon reconstruction efficiency is indeed identical for both simulations as function of the true muon p_T and η . Also for the transverse momentum resolution of the reconstructed muon no differences between both simulations are visible, as shown in Figure 7.17.

While true muons show a very good agreement to the full simulation, basically because they are simulated with the full simulation of all detector components, misidentified muons can come from several sources. The p_T distribution of misidentified muons is shown in Figure 7.18. The misidentified muons can be sorted into two categories, depending on where and how they are produced:

- A large fraction of misidentified muons are true muons originating from

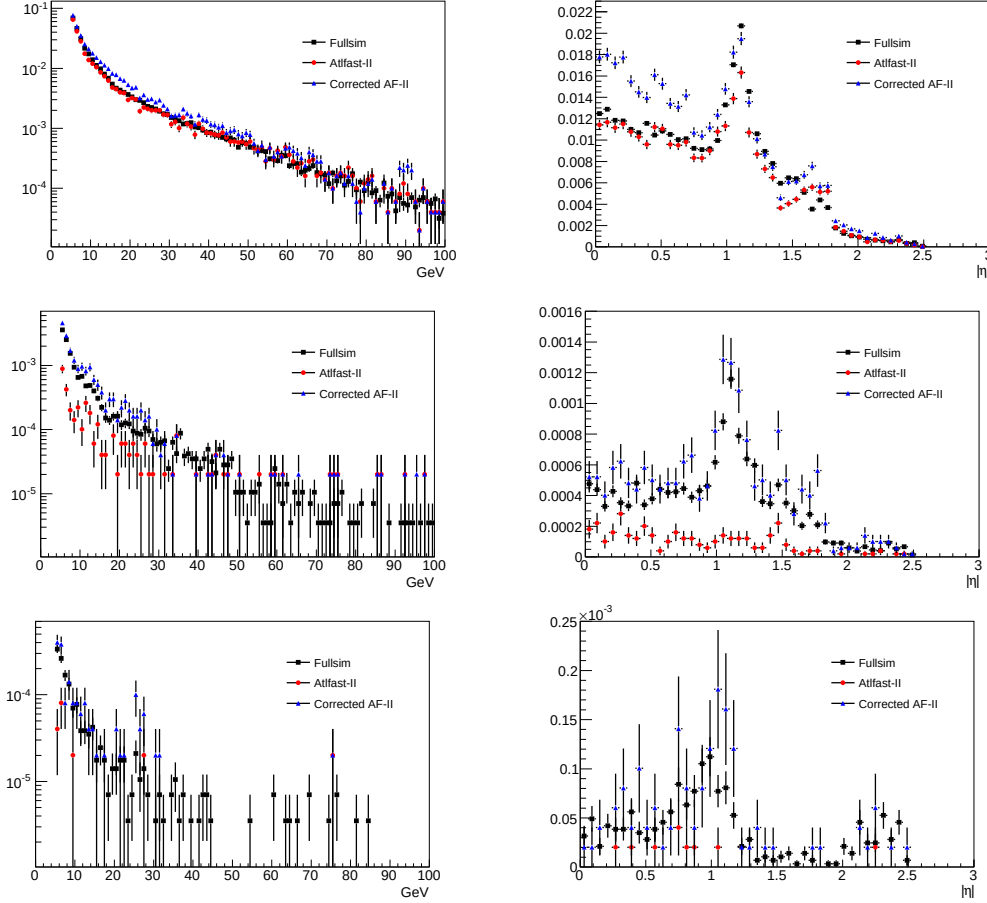


Figure 7.15: Probability for a jet to fake an electron as a function of p_T and $|\eta|$ [31]. The standard loose (top), medium (middle) and tight (bottom) $isEM$ selections are applied to jets from $t\bar{t}$ events. The results obtained from full simulation (black), Atlfast-II (red) and corrected Atlfast-II (blue) are superimposed.

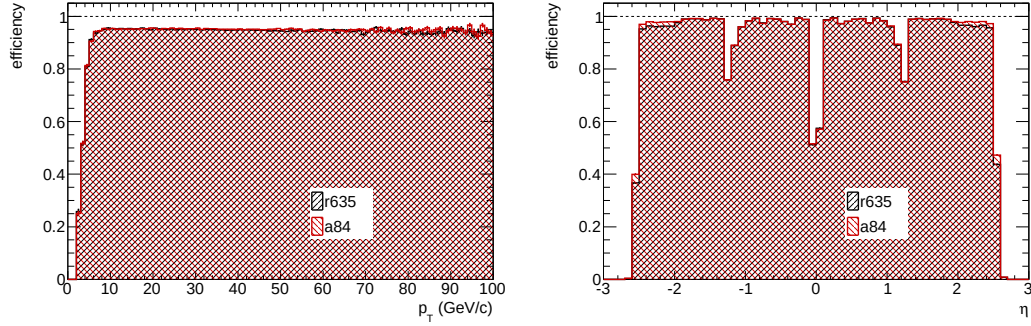


Figure 7.16: The identification efficiencies for muons as a function of p_T for $|\eta| < 2.5$ (left), and as a function of η for $p_T > 10$ GeV (right) [31]. The standard combined muon selection (STACO) is applied. The results are shown for full simulation (black) and Atlfast-II (red).

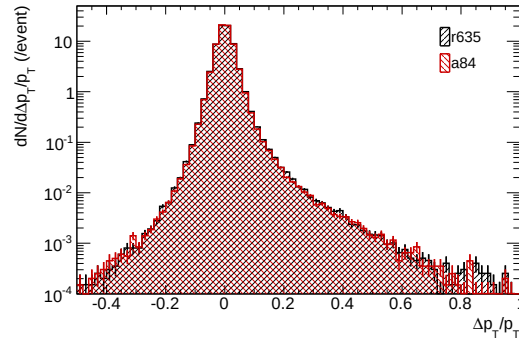


Figure 7.17: The relative muon momentum resolution $\Delta p_T/p_T$ for full simulation (black) and Atlfast-II (red) for the standard STACO combined muon selection for muons with $p_T > 10$ GeV from $Z \rightarrow \mu\mu$ decays.

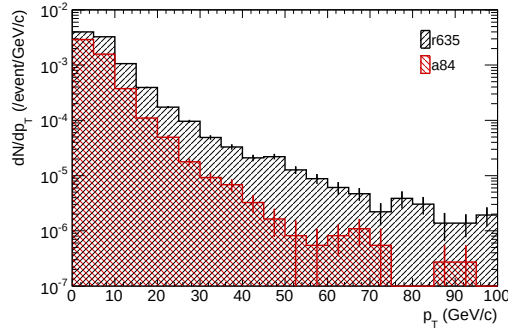


Figure 7.18: The p_T distributions of misidentified muons in $t\bar{t}$ events for full simulation (black) and Atlfast-II (red) [31].

pion and kaon decays in the inner detector. These muons are treated like true muons in the Atlfast-II simulation and hence are simulated correctly.

- However, the remaining fraction of misidentified muons originates from hadronic showers that are not contained in the calorimeter and leak into the muon system, decays of hadrons inside the calorimeter that produce secondary muons and wrongly reconstructed muons due to pattern recognition problems. This second fraction of misidentified muons is not simulated in Atlfast-II.

A correction of the muon misidentification efficiency, as it was done for electrons, is not easily possible because the second type of misidentified muons is completely missing and hence a reweighting of this component by a correction factor is not possible. However, it is planned to add a parametrisation of misidentified muons of the second type to Atlfast-II in the future.

7.2.3 Performance of jets

The performance of jets was already compared in Section 7.1.3 to the full simulation in Athena release 10 and showed a good agreement in all variables. In this section the comparison is repeated against full simulated jets from Athena release 14.

The energy response of the calorimeter in the full simulation was changed between the Athena release 10, that was used for the FastCaloSim parametrisation, and release 14 and this change in the energy response is therefore also visible between the fast simulation and the full simulation in release 14. In order to compare jet distributions from both simulations, independent jet calibrations are used (see Section 7.1.3) and only calibrated jets are used for the comparison.

Figure 7.19 shows the jet reconstruction efficiency as function of p_T for several bins in the pseudorapidity. Apart from the low p_T region, where small differences are visible, the jet efficiency in fast simulation agrees very well to the full simulation. The increased efficiency for very low p_T jets is most likely caused by the missing lateral shape fluctuations in the pion simulation that cause more localised energy depositions which are more efficient to be seen as jet.

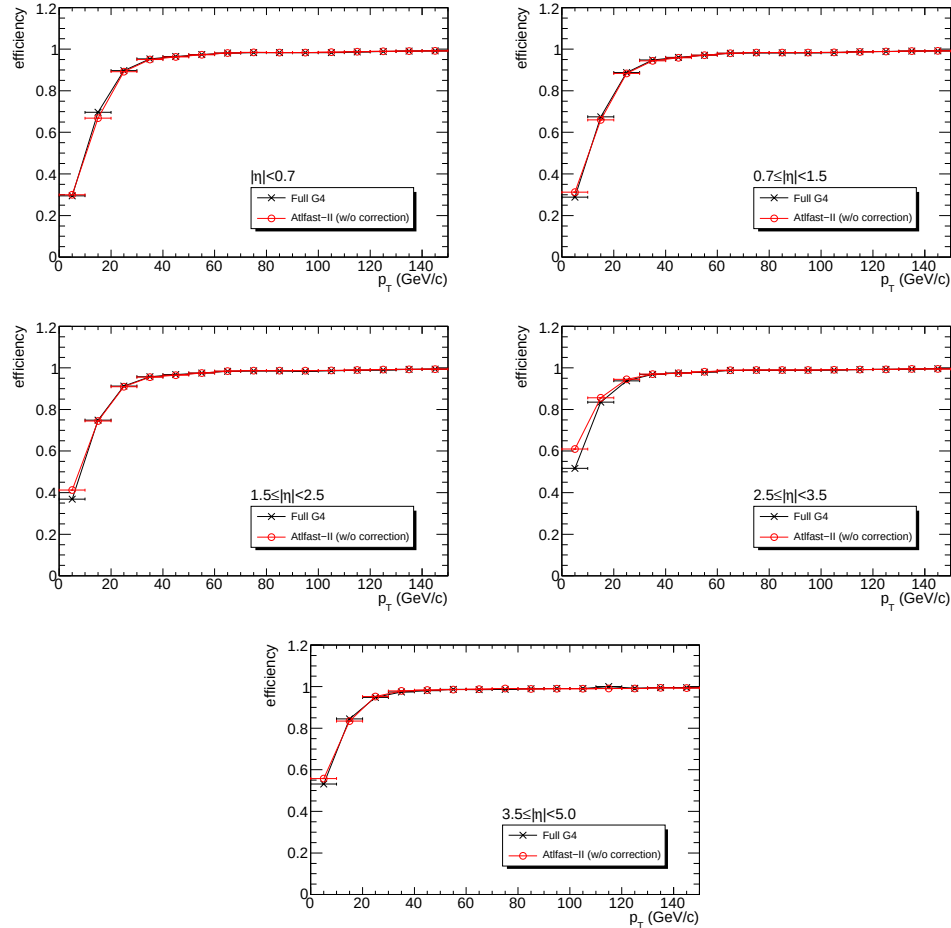


Figure 7.19: Efficiencies for the jet reconstruction as a function of p_T for different $|\eta|$ regions [31]: the barrel ($|\eta| < 0.7$), the extended barrel ($|\eta| = 0.7$ – 1.5), the end cap ($|\eta| = 1.5$ – 2.5), the transition region ($|\eta| = 2.5$ – 3.5) and the forward region ($|\eta| > 3.5$). The results are shown for full simulation (black) and Atfast-II (red).

Because of the independent jet calibration, the jet energy scale of both simulations, as shown in Figure 7.20, agrees well and is within 1% of the true jet energy for jets with a transverse momentum $p_T \gtrsim 100$ GeV. For jets below 100 GeV

the H1 jet calibration method gives deviations from the true jet energy of up to 5% for both simulations. However, the differences between full and fast simulation, that appear in this region, are smaller than the deviations from the true jet energy and can be seen as uncertainties from the jet calibration method.

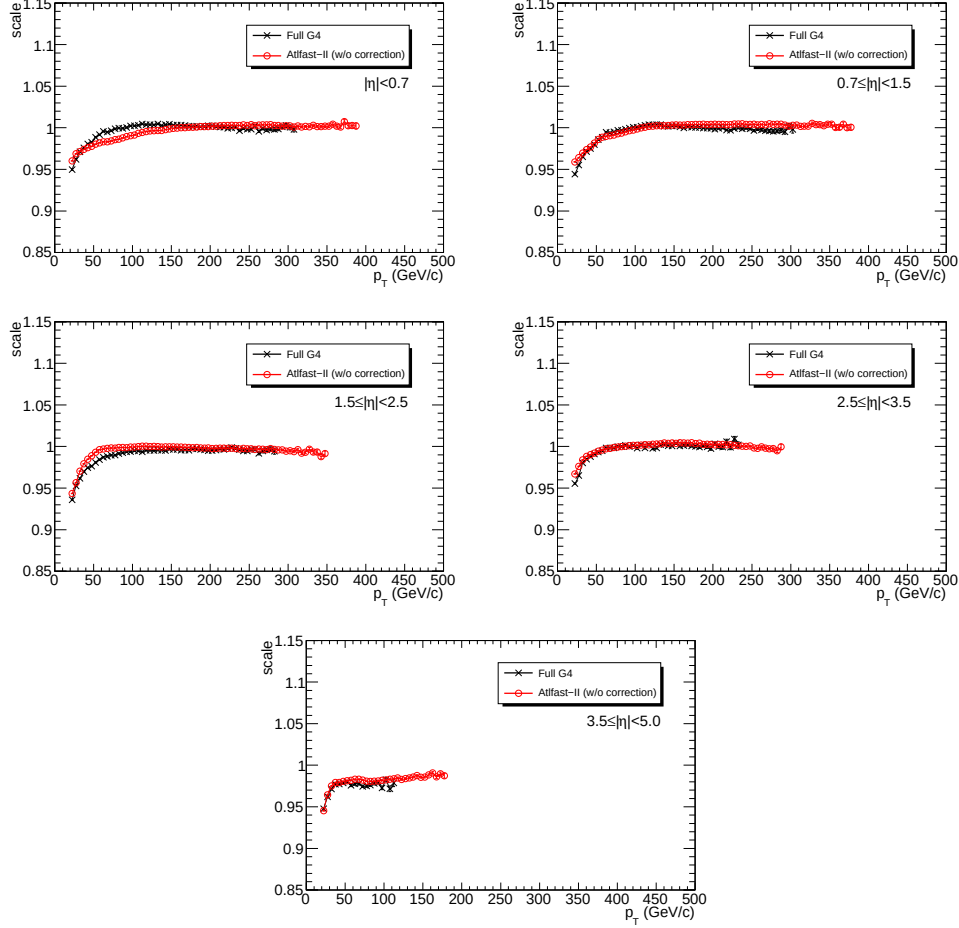


Figure 7.20: The mean values of the jet energy scale $p_T^{\text{rec}}/p_T^{\text{true}}$ as a function of p_T for light flavour jets from $t\bar{t}$ events for different $|\eta|$ regions [31]: the barrel ($|\eta| < 0.7$), the extended barrel ($|\eta| = 0.7\text{--}1.5$), the end cap ($|\eta| = 1.5\text{--}2.5$), the transition region ($|\eta| = 2.5\text{--}3.5$) and the forward region ($|\eta| > 3.5$). The results are shown for full simulation (black) and Atlfast-II (red).

The jet energy resolution is shown in Figure 7.21 as function of the jet p_T for the same pseudorapidity as above. For $|\eta| < 2.5$ the agreement is very good. Above $|\eta| = 2.5$ the jet resolution is underestimated by $\approx 10\%$. However, jets from this region are usually not used for precision jet observables and this resolution difference should hence not cause deviations in analysis.

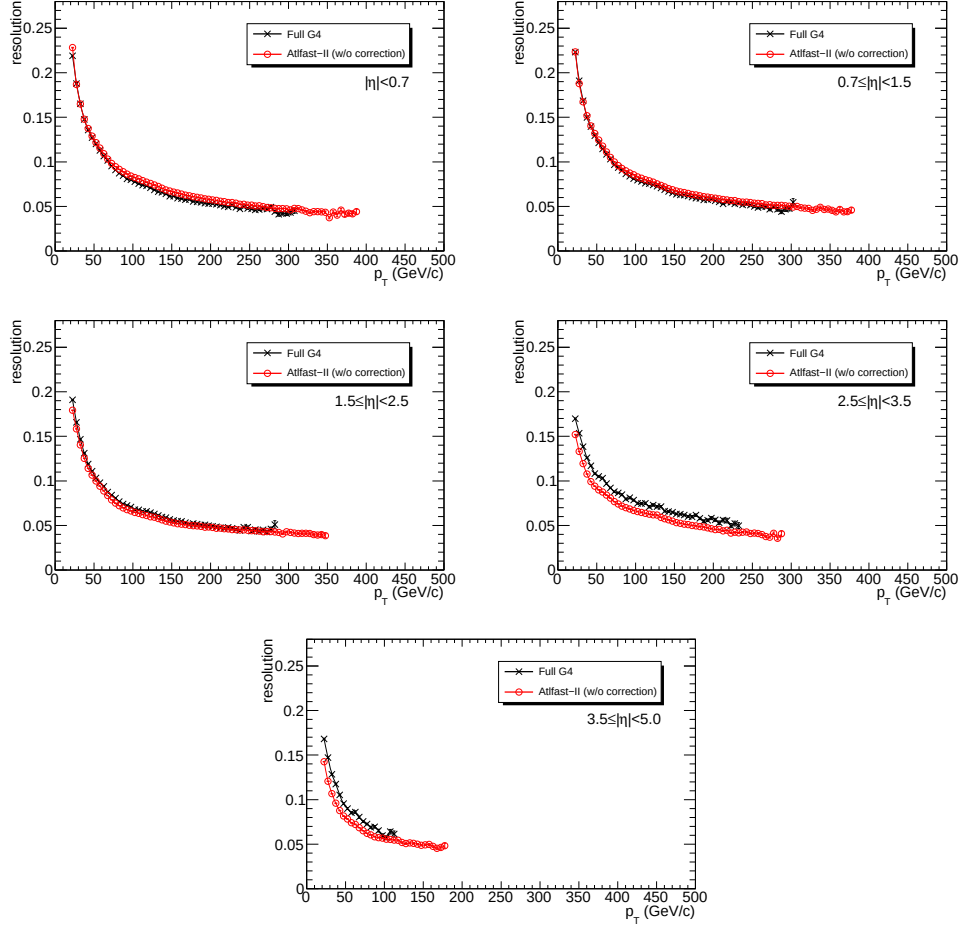


Figure 7.21: The jet energy resolution as a function of p_T for the light flavor jets in $t\bar{t}$ events in different $|\eta|$ regions [31]: the barrel ($|\eta| < 0.7$), the extended barrel ($|\eta| = 0.7\text{--}1.5$), the end cap ($|\eta| = 1.5\text{--}2.5$), the transition region ($|\eta| = 2.5\text{--}3.5$) and the forward region ($|\eta| > 3.5$). The results are shown for full simulation (black) and Atlfast-II (red).

7.2.4 Performance of the missing transverse momentum

The missing transverse momentum $\cancel{E}_T$ is reconstructed from the vectorial sum of the calibrated transverse momenta of all visible objects in the detector [15]. In addition, clustered energy not associated to a reconstructed electron, photon, hadronic τ or jet is included in the calculation of the missing transverse momentum and calibrated using the H1 calibration method. The missing transverse momentum is very sensitive to the overall quality of the detector simulation, because all detector components and reconstructed objects are included in the calculation.

Figure 7.22 shows the resolution of the reconstruction of the missing transverse momentum for QCD jet events, where no intrinsic $\cancel{E}_T$ is expected, and $W \rightarrow e\nu$ events, where intrinsic $\cancel{E}_T$ is expected due to the neutrinos in the final state. The agreement is very good for the core component of the distribution. However, differences are found in the extreme tail of the distribution for $|\Delta\cancel{E}_T| \gtrsim 100$ GeV.

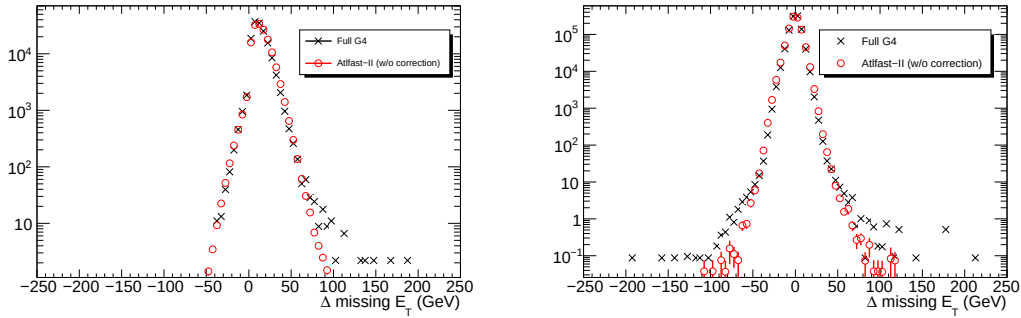


Figure 7.22: The distributions of $\Delta\cancel{E}_T = \cancel{E}_T(\text{rec}) - \cancel{E}_T(\text{true})$ for events from QCD jet production (left) and $W \rightarrow e\nu$ events (right) [31]. The results from full simulation (black) and Atfast-II (red) are superimposed.

7.3 Summary

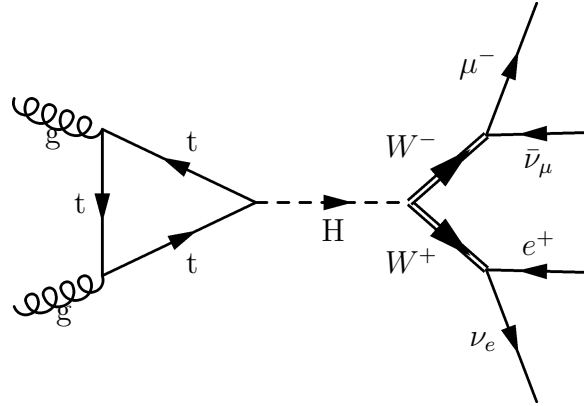
The fast calorimeter simulation FastCaloSim included in the Atfast-II detector simulation gives a good description of object energy response and resolution functions and of reconstruction and identification efficiencies. If efficiency corrections for reconstructed objects are included, the remaining differences to the full simulation can be reduced in most cases to the 1–2% level, with the exception of misidentified electrons and muons, where larger differences remain due to intrinsic limitations in the fast simulation.

Nevertheless, Atfast-II is well suited for the simulation of large Monte Carlo samples and allows to analyse processes that cannot be studied in detail with full simulation due to CPU time limitations.

8 Study of the inclusive Higgs boson production in the $H \rightarrow WW$ final state

Current experimental data and theoretical constraints favour a relatively light Higgs boson with a mass below 200 GeV, as it was shown in the introduction to the properties of the Standard Model and the Higgs boson at the LHC in Chapters 1 and 3. In the mass range of $114 \text{ GeV} < m_H \lesssim 200 \text{ GeV}$ the dominant Higgs boson production mode is gluon fusion and the dominant decay mode is into a pair of W bosons¹. The LHC estimates of the Higgs boson discovery potential show that the combination of these two dominant modes in the $gg \rightarrow H \rightarrow W^-W^+ \rightarrow l^-\bar{\nu} l^+\nu$ final state is the most promising discovery channel for a Higgs boson in the mass range between 130 GeV and 190 GeV (see Figure 3.5).

The signature of the channel consists of two isolated leptons² and large missing transverse energy from the decays of the two W bosons, as can be seen in the basic leading order Feynman diagram:



No additional physics objects are visible from the production mode. However, if higher order QCD corrections and soft QCD activity are taken into account,

¹For Higgs boson masses below $\approx 135 \text{ GeV}$ the W pair final state is only the second largest decay mode after the $b\bar{b}$ pair decay. Unfortunately, the $b\bar{b}$ final state is very hard to separate from QCD backgrounds at the LHC and most likely impossible to observe in gluon fusion production.

²In the following chapters leptons will always refer to electrons or muons. τ leptons and their hadronic decay products are considered separately.

a picture as sketched in Figure 2.5 is more appropriate and additional detector activity in form of (soft) jets will be visible in some fraction of the signal events.

In the following sections the properties of the signal and the dominant background processes are described in the context of a simple cut based analysis as it would be applied to first data. The analysis is neither optimised to give a maximal significance nor are all backgrounds and all possible systematic uncertainties included. The intention of this chapter is to illustrate the main challenges in the analysis and to highlight the problem of reliably estimating the background, especially if Monte Carlo simulations and data turn out to be not in perfect agreement.

In Chapter 9 two data driven and thus Monte Carlo simulation independent methods to determine the shape and rate of the dominant dilepton background processes are developed.

Finally, Chapter 10 uses one of the data driven background determination methods to estimate the discovery significance for the $gg \rightarrow H \rightarrow W^-W^+ \rightarrow l^-\bar{\nu} l^+\nu$ final state in the early data taking phase of the LHC running at 10 TeV collision energy³.

8.1 Simple cut based analysis

For the simple cut based analysis only the signal process and the most important Standard Model background processes with two leptons in the final state are considered. For the signal process two Higgs boson masses of $m_H = 130$ GeV and $m_H = 170$ GeV are chosen to illustrate the mass dependency. Table 8.1 shows a summary of all processes taken into account in this analysis and their cross section at a centre-of-mass energy of $\sqrt{s} = 10$ TeV. A discussion of the individual processes and the corresponding choice of a suited Monte Carlo generator is given in Appendix B together with additional background processes that can contribute to the observed final state.

The emphasis of this simple cut based analysis is to illustrate the kinematic properties of the signal and of the dominant background processes. Detector specific aspects of the final state selection are briefly mentioned here and, if relevant, further elaborated in Appendix B on the properties of the various processes. Therefore the analysis presented here combines the e^+e^- , $e^\pm\mu^\mp$ and $\mu^+\mu^-$ final states instead of treating them separately to reach the maximal significance.

³The LHC schedule foresees a start with a lower collision energy and an increase to 14 TeV later. The collision energy of $\sqrt{s} = 10$ TeV was the nominal value assumed in the Monte Carlo simulations of the last two years.

Generator	Process	$O(\alpha_s)$	Parton level selection	$\sigma \cdot \text{BR}$
MC@NLO [13] + Herwig [34]	$gg \rightarrow H \rightarrow WW$ ($m_H = 130$ GeV)	NLO	2 leptons, $l = e, \mu, \tau$	273 fb
MC@NLO	$gg \rightarrow H \rightarrow WW$ ($m_H = 170$ GeV)	NLO	2 leptons, $l = e, \mu, \tau$	468 fb
MC@NLO	$q\bar{q} \rightarrow WW$	NLO	2 leptons, $l = e, \mu, \tau$	7,440 fb
gg2WW [35] + Herwig	$gg \rightarrow WW$	LO	2 leptons, $l = e, \mu, \tau$	284 fb
MC@NLO	$q\bar{q}/gg \rightarrow t\bar{t}$	NLO	≥ 1 lepton, $l = e, \mu, \tau$	207,000 fb
AcerMC [36] + Pythia [37]	$gb \rightarrow Wt$	LO	2 leptons, $l = e, \mu, \tau$	2,790 fb
Pythia	$\gamma^{(*)}/Z \rightarrow ll$	LO	2 leptons, $l = e, \mu$ $m_{ll} > 10$ GeV, $1 \cdot e, \mu: p_T > 10$ GeV $ \eta < 2.7$	3,540,000 fb

Table 8.1: Summary of the considered processes and their cross section at $\sqrt{s} = 10$ TeV. To increase the generator efficiency, parton level lepton selection cuts are applied to some background processes. τ final states are not directly analysed, but the leptonic τ decay modes are included. The cross sections for $t\bar{t}$ and $\gamma^{(*)}/Z$ include the parton level selection efficiency.

8.1.1 Selection of the leptonic final state

The first selection step requires the presence of exactly two isolated leptons of opposite sign in the electric charge. The selection cuts for the individual lepton flavours are summarised in Table 8.2. Figure 8.1 shows the transverse momentum (p_T) and the pseudorapidity (η) distribution of the two leptons after applying these selection cuts for the signal and the dominant dilepton backgrounds. The pseudorapidity of the leptons nicely shows the difference between mainly $q\bar{q}$ and mainly gg initiated processes.

The tight lepton identification and isolation requirements are motivated by the need to reject the large W +jet background, where one jet is misidentified as a second lepton. The lepton identification and isolation requirements are discussed in the context of the W +jet background rejection in Appendix B.4.

A summary of the accepted Monte Carlo cross section for all processes after each selection step is shown at the end of this chapter in Table 8.3. The selection of the leptonic final state is listed as cut 1 in this table.

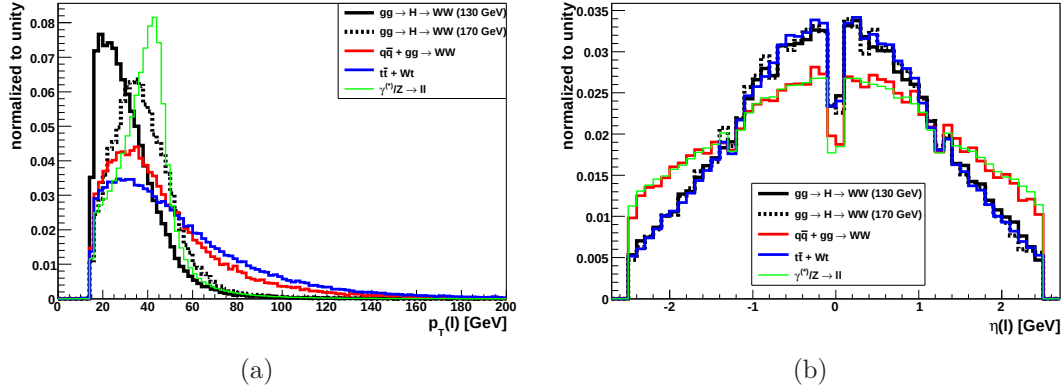


Figure 8.1: The lepton p_T (a) and η (b) distributions after the lepton selection cuts for Higgs boson signal events ($m_H = 130$ GeV and $m_H = 170$ GeV) and the dominant dilepton background processes.

8.1.2 Event selection to suppress reducible backgrounds

After the preselection of events with two leptons, the background is dominated by $\gamma^{(*)}/Z$ events. For the signal process, the presence of two neutrinos in the final state leads to a significant amount of missing transverse energy E_T^{miss} . Figure 8.2(a) shows the normalised distribution of E_T^{miss} .

While all backgrounds with two W bosons in the final state have a similar amount of E_T^{miss} as the signal, a cut on this variable is effective to suppress backgrounds without intrinsic E_T^{miss} due to neutrinos, notably $\gamma^{(*)}/Z \rightarrow ll$ (see Appendix B.3).

Selection of events with sizeable transverse momentum of the dilepton system

In processes with two intermediate W bosons, the missing transverse momentum is balanced by the transverse momentum

$$p_T(l_1 + l_2) = |\vec{p}_T(l_1) + \vec{p}_T(l_2)|$$

of the vectorial sum of the two leptons as long as there is no significant additional event activity. Figure 8.2(b) shows the distribution of $p_T(l_1 + l_2)$, which overall is similarly distributed as E_T^{miss} . The differences come mainly from the p_T and η selection criteria applied on the individual leptons, which is obviously not the case for the individual neutrinos.

The correlation of these two variables is shown in Figure 8.3 after the lepton selection cuts (a) and after an additional cut (see the following subsection) to veto jets with $p_T > 20$ GeV (b) in the event. This cut reduces the extra event activity considerably.

Lepton flavour	Electron	Muon
Identification	tight electron selection	combined muon ID track match $\chi^2 < 100$
Kinematic cuts	$p_T > 15 \text{ GeV} ; \eta < 2.5$	
Calorimeter isolation	$\Sigma E_T^{\Delta R < 0.3} < 10 \text{ GeV} ; \Sigma E_T^{\Delta R < 0.3} / p_T(l) < 0.2$ Sum over all calorimeter without the lepton energy	
Track isolation	$\Sigma p_T^{\Delta R < 0.3} < 10 \text{ GeV} ; \Sigma p_T^{\Delta R < 0.3} / p_T(l) < 0.1$ Sum over all tracks with $p_T(\text{track}) > 1 \text{ GeV}$, except the lepton track	
Track origin	$ \text{IP}/\sigma(\text{IP}) < 10$, IP=Impact parameter	

Table 8.2: Summary of the electron and muon selection cuts. Details are discussed in Appendix B.4.

Especially after this jet veto the clear correlation of E_T^{miss} and $p_T(l_1 + l_2)$ is evident. While the measurement of E_T^{miss} has large uncertainties and depends on a very good performance of the whole detector, the momentum of the leptons is measured with high precision. Therefore, a cut on $p_T(l_1 + l_2) > 30 \text{ GeV}$ is applied and listed as cut 2 in Table 8.3.

Veto of QCD jet activity

A veto on additional jet activity serves two purposes. On the one hand, it allows to replace to some extent the cut on the uncertain E_T^{miss} by a cut on the well measured leptons. On the other hand, there are several backgrounds, notably $t\bar{t}$ and Wt , that have large intrinsic jet activity. Figure 8.4 shows the p_T distribution of the leading transverse momentum jet for the signal and background processes before (a) and after (b) the cut on $p_T(l_1 + l_2) > 30 \text{ GeV}$.

The leading jet p_T shows a clear distinction between the $t\bar{t}$ and Wt background on the one hand and the signal processes on the other hand. A veto on any jet with $p_T(\text{jet}) > 20 \text{ GeV}$ rejects $> 99\%$ of the $t\bar{t}$ and $> 95\%$ of the Wt background, while most of the signal survives (listed as cut 3 in Table 8.3). Figure 8.4 also shows that the kinematic distribution of $\gamma^{(*)}/Z$ events is changed considerably by the cut on the p_T of the two leptons⁴ so that this background is further reduced by the jet veto.

⁴For the $\gamma^{(*)}/Z$ background, a cut on $p_T(l_1 + l_2)$ is a selection of $\gamma^{(*)}/Z$ bosons with sizeable transverse momentum. As the boson needs to be kinematically balanced by other event activity, recoiling jets are usually found in the opposite detector hemisphere.

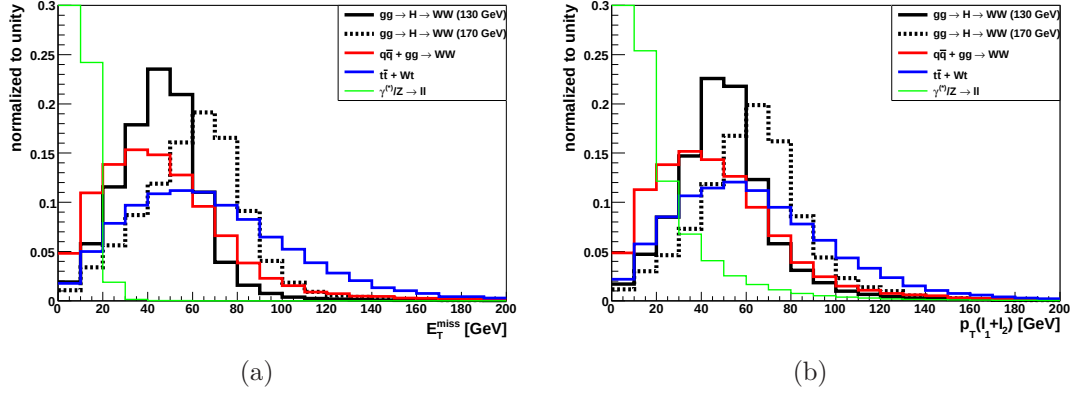


Figure 8.2: (a) Distribution of the missing transverse energy E_T^{miss} (b) of the transverse momentum $p_T(l_1 + l_2)$ of the vectorial sum of the two leptons after the lepton selection cuts for Higgs boson signal events ($m_H = 130$ GeV and $m_H = 170$ GeV) and the dominant dilepton background processes. Please note that the lowest bin $x = [0, 10]$ GeV on the x-axis of the $\gamma^{(*)}/Z \rightarrow ll$ distribution is at 0.74 (a) respectively 0.44 (b) and hence not fully shown.

Selection of events with missing transverse energy

After the jet veto, the $\gamma^{(*)}/Z$ background is still the dominant dilepton background source and E_T^{miss} , shown in Figure 8.5(a) after the jet veto, is the only remaining variable with sufficient suppression power. Hence it is unavoidable to apply a selection cut on $E_T^{\text{miss}} > 30$ GeV to reduce the $\gamma^{(*)}/Z$ background (listed as cut 4 in Table 8.3).

However, due to the previous cut on $p_T(l_1 + l_2) > 30$ GeV the E_T^{miss} cut has a high efficiency for the signal process and all other backgrounds but $\gamma^{(*)}/Z$. Thus systematic uncertainties in the measurement of E_T^{miss} will only have a minor influence on the overall selection efficiency for these processes. On the other hand, the cut reduces the $\gamma^{(*)}/Z$ background by a large, but uncertain (see Section B.3.1), fraction.

Finally, Figure 8.5(b) shows that after the E_T^{miss} cut the $\gamma^{(*)}/Z$ background retains the characteristic Z boson resonance mass peak. If a sufficiently large number of $\gamma^{(*)}/Z$ events remains after the E_T^{miss} cut, their amount can be estimated from the invariant dilepton mass distribution. The dominant part of the remaining $\gamma^{(*)}/Z$ events could be removed with a veto cut on the resonant Z peak. However, in view of the background normalisation strategy of the Chapters 9 and 10, such a cut is not applied.

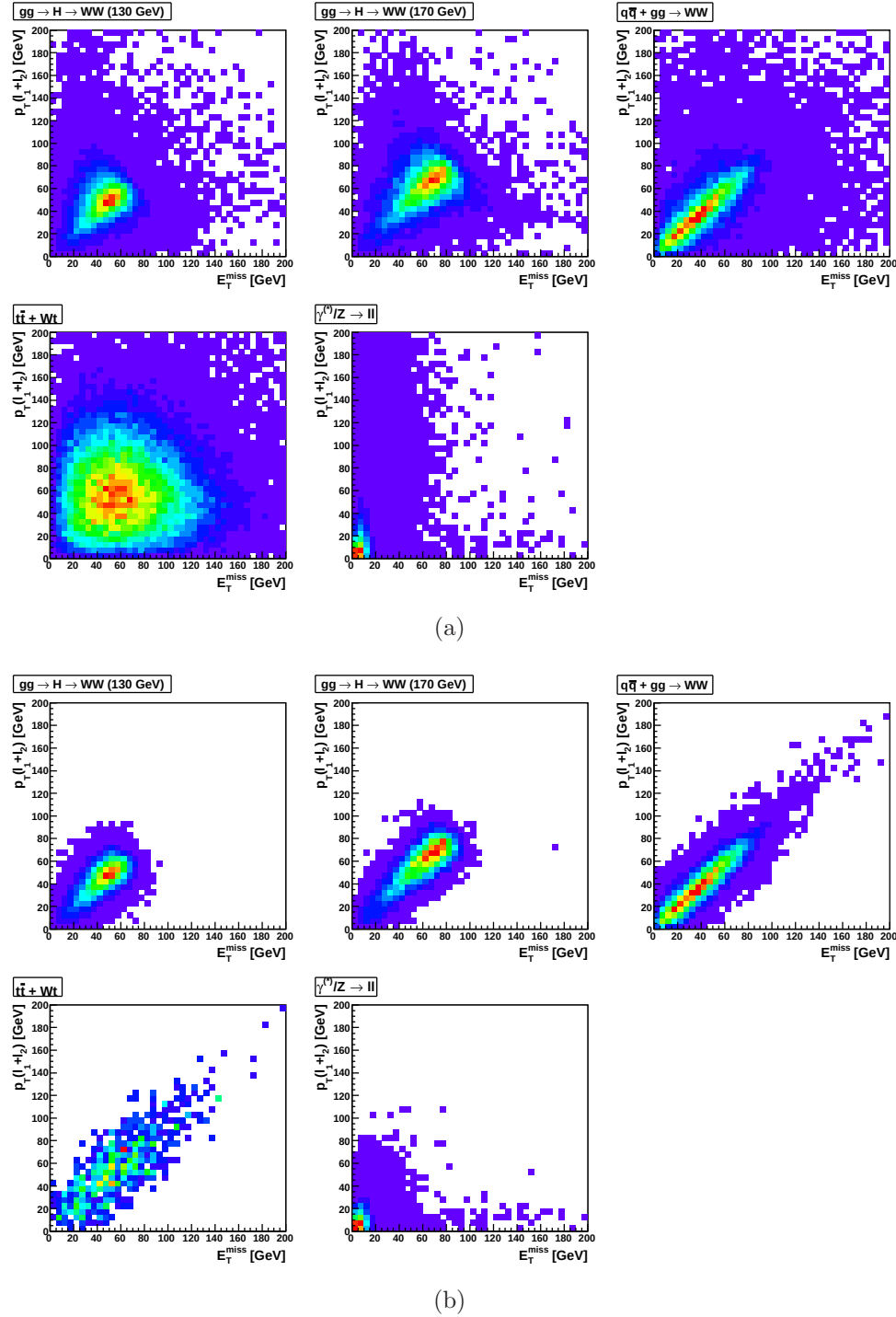


Figure 8.3: Correlation between the missing transverse energy E_T^{miss} and the transverse momentum $p_T(l_1 + l_2)$ of the vectorial sum of the two leptons after the lepton selection cuts (a) and after an additional jet veto cut (b) for Higgs boson signal events ($m_H = 130$ GeV and $m_H = 170$ GeV) and the dominant dilepton background processes.

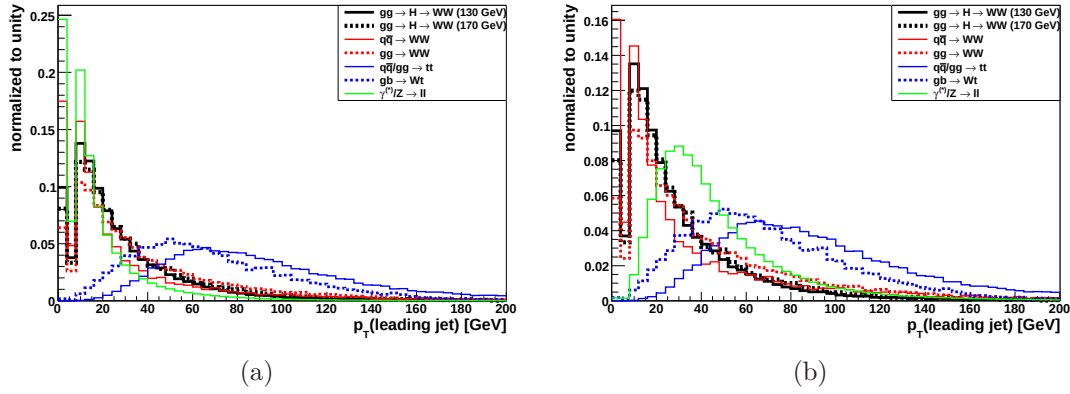


Figure 8.4: Transverse momentum p_T distribution of the leading p_T jet before (a) and after (b) the cut on $p_T(l_1 + l_2) > 30$ GeV is applied for Higgs boson signal events ($m_H = 130$ GeV and $m_H = 170$ GeV) and the dominant dilepton background processes. Events without a jet are filled into the lowest p_T bin.

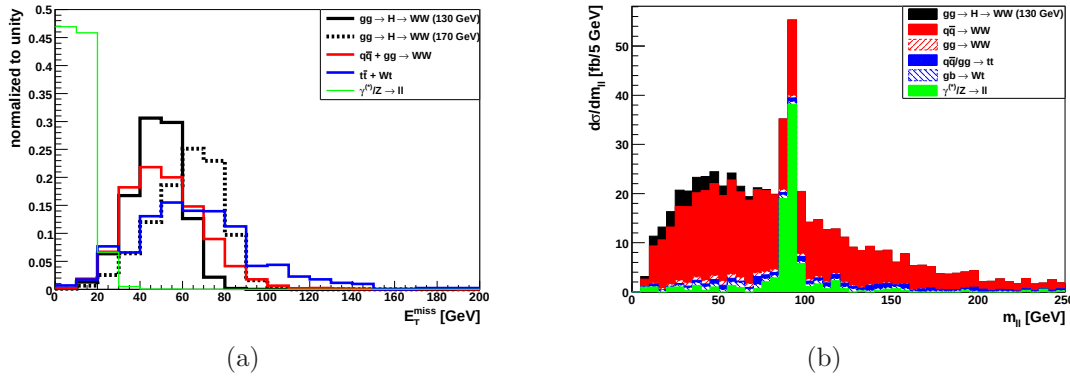


Figure 8.5: (a) Distribution of the missing transverse energy E_T^{miss} after the jet veto cut for Higgs boson signal events ($m_H = 130$ GeV and $m_H = 170$ GeV) and the dominant dilepton background processes. (b) Distribution of the total sum of all contributions to m_{ll} after the cut on $E_T^{\text{miss}} > 30$ GeV for Higgs boson signal events ($m_H = 130$ GeV) and the dominant dilepton background processes.

Additional safety cuts

Three more cuts are applied to the dilepton and E_T^{miss} system. A cut on the transverse opening angle of $\Delta\phi_{ll} < 3$ vetoes back-to-back topologies (cut 5 in Table 8.3). This is the case for some of the remaining $\gamma^{(*)}/Z \rightarrow ll$ events and especially for $\gamma^{(*)}/Z \rightarrow \tau\tau$ events (see Section B.3.2).

A cut on the transverse vectorial sum of the two leptons and p_T^{miss} , the transverse momentum vector of E_T^{miss} , at $p_T(l_1 + l_2 + p_T^{\text{miss}}) < 60$ GeV is used to suppress any large off-diagonal elements in the correlation of E_T^{miss} and $p_T(l_1 + l_2)$ (cut 6 in Table 8.3). As there are no such off-diagonal elements in the Monte Carlo simulation, the cut has almost no effect.

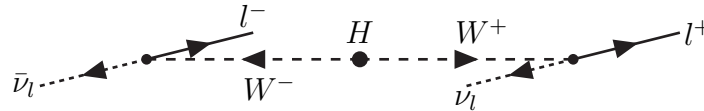
Finally, a cut on the pseudorapidity difference of the two leptons at $\Delta\eta_l < 3$ is intended to possibly suppress low mass QCD backgrounds ($c\bar{c}$, $b\bar{b}$) and backgrounds due to false lepton identification (W +jets), as this region is not populated by resonant or non-resonant high p_T processes⁵, but might get contributions from these processes (cut 7 in Table 8.3).

All three cuts have only a small influence on the current Monte Carlo event selection, but remain as safety measures against unforeseen detector effects or backgrounds.

8.1.3 Kinematic selection to suppress the irreducible WW background

At this stage of the event selection the reducible backgrounds from $t\bar{t}$, Wt and $\gamma^{(*)}/Z \rightarrow ll$ are relatively small and suppressed approximately by a factor of 5-10 with respect to the irreducible $q\bar{q} \rightarrow WW$ process. Any suppression of the $q\bar{q} \rightarrow WW$ background itself against the signal has to come from the different production characteristics of the two W bosons.

Figure 8.6 shows the Feynman diagrams for the signal and the $q\bar{q} \rightarrow WW$ process. While the WW system is produced mostly non-resonant and with an undefined total angular momentum state for the background, it is produced from a narrow spin 0 resonance for the Higgs boson signal. The spin 0 decay of the Higgs boson occurs into two oppositely charged W bosons, travelling in opposite direction and with opposite relative spin orientation in the Higgs boson rest frame.



Finally, the weak decay of the oppositely charged W bosons with opposite spin orientation results in two leptons that tend to travel in the same direction and two neutrinos that tend to go into the opposite direction adding up to a sizeable E_T^{miss} .

⁵extremely high mass systems could populate this region, but are not relevant in this analysis

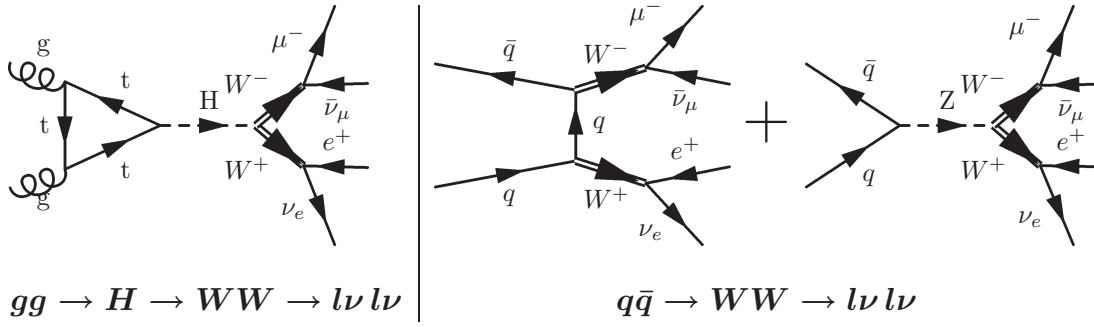


Figure 8.6: Feynman diagrams for the signal process (left) and the dominant irreducible WW background (right).

The transverse opening angle $\Delta\phi_u$

This spin correlation in the $H \rightarrow W^-W^+ \rightarrow l^-\bar{\nu} l^+\nu$ decay is best observable in the transverse opening angle $\Delta\phi_u$ of the two leptons, as shown in the comparison of the normalised shapes in Figure 8.7(a). The relative importance of the individual background components is visible in the stacked distributions shown in Figure 8.7(b). If not indicated otherwise, only the signal distribution for a Higgs boson mass of $m_H = 130$ GeV is shown in the following figures.

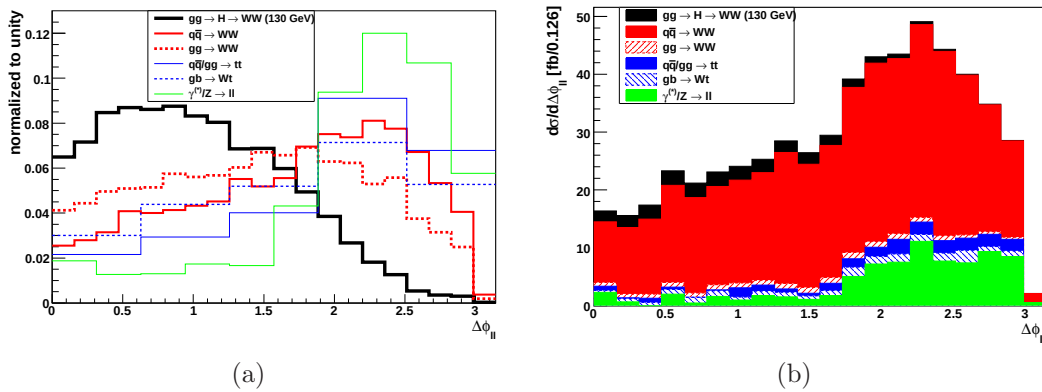


Figure 8.7: Distribution of the transverse opening angle $\Delta\phi_u$ for Higgs boson signal events ($m_H = 130$ GeV) and the dominant dilepton background processes after all cuts against reducible backgrounds. The normalised distribution is shown in (a), the total sum of all contributions to the $\Delta\phi_u$ distribution is shown in (b). Due to low Monte Carlo statistics, a coarse binning is chosen for the top and $\gamma^{(*)}/Z$ processes.

As discussed above, the signal exhibits leptons that are close together in phase space and hence have a small transverse opening angle while most backgrounds produce leptons at larger angles. However, even after a tight cut on $\Delta\phi_u$ the

$q\bar{q} \rightarrow WW$ background would remain approximately a factor 2 larger compared to the signal even for a Higgs boson mass of $m_H = 170$ GeV, where the signal event yield is largest (see cut 8a in Table 8.3).

The transverse mass m_T

The second characteristic difference in the production of the WW system in the signal process is the narrow invariant mass resonance of the Higgs boson, while all background processes produce a wide continuum of invariant WW boson masses starting at the WW threshold energy of ≈ 160 GeV. If it were possible to reconstruct the invariant WW mass, this would be a striking signature. But, due to the presence of the two neutrinos in the final state, it is unfortunately not possible at the LHC to reconstruct the invariant WW system mass⁶.

However, instead of the invariant mass it is possible to define several types of transverse mass for the system consisting of the two leptons and the missing transverse momentum. The definition of a transverse mass is not unique, but all definitions of the transverse mass result in a mass peak of the signal process close to the true Higgs boson mass with a narrow edge towards high masses and a long tail towards low masses. An overview of some different choices for the definition of the transverse mass with their advantages and disadvantages is given in Appendix C.

In the following, the exclusive transverse mass $m_T \equiv M_T^{\text{true,T}}$ is used and labelled just as m_T . The easiest way of calculating the transverse mass $M_T^{\text{true,T}}$ is to calculate the invariant mass of purely transverse Lorentz vectors:

$$\begin{aligned}\hat{p}^{l_1} &= (p_T^{l_1}, p_x^{l_1}, p_y^{l_1}, 0)^T \\ \hat{p}^{l_2} &= (p_T^{l_2}, p_x^{l_2}, p_y^{l_2}, 0)^T \\ \hat{p}^{\text{miss}} &= (E_T^{\text{miss}}, p_x^{\text{miss}}, p_y^{\text{miss}}, 0)^T \\ \left(M_T^{\text{true,T}}\right)^2 &= (\hat{p}^{l_1} + \hat{p}^{l_2} + \hat{p}^{\text{miss}})_\mu \cdot (\hat{p}^{l_1} + \hat{p}^{l_2} + \hat{p}^{\text{miss}})^\mu \quad . \quad (8.1)\end{aligned}$$

The distribution of the transverse mass after the cuts against reducible backgrounds is shown in Figure 8.8, once as a comparison of the normalised shapes in Figure 8.8(a) and once stacked to show the relative importance of the individual components in Figure 8.8(b).

Overall, the distribution of m_T is qualitatively narrower for the signal than for the backgrounds, but the shapes are not sufficiently different to allow a reduction of the backgrounds by a large factor.

⁶The final state consists of 4 approximately massless fermions. To reconstruct the invariant mass one needs $4 \cdot 3 = 12$ independent kinematic measurements. Measurable are the two lepton 3-momenta and the missing transverse momentum 2-vector: $2 \cdot 3 + 2 = 8$ degrees of freedom. Hence 4 measurements are missing for the Higgs boson mass reconstruction.

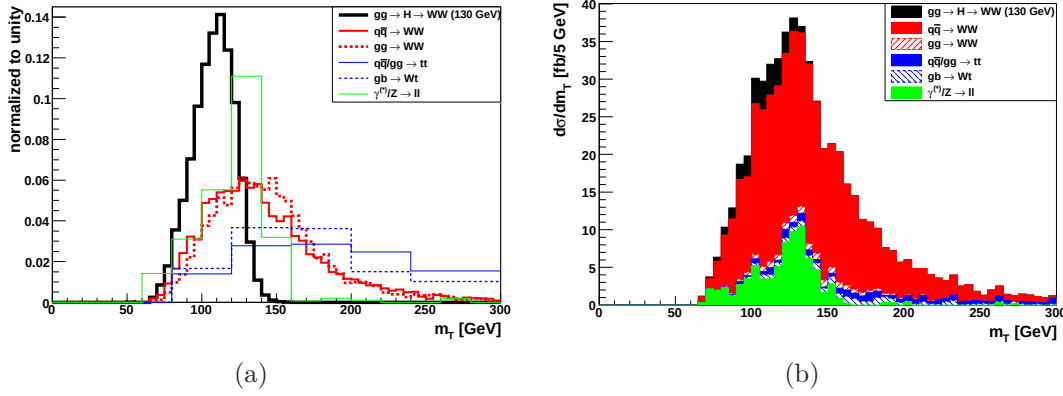


Figure 8.8: Distribution of the transverse mass m_T for Higgs boson signal events ($m_H = 130$ GeV) and the dominant dilepton background processes after all cuts against reducible backgrounds. The normalised distribution is shown in (a), the total sum of all contributions to the m_T distribution in (b). Due to low Monte Carlo statistics, a coarse binning is chosen for the top and $\gamma^{(*)}/Z$ processes.

Event selection

When looking at the dependency of the transverse mass m_T on the transverse opening angle $\Delta\phi_{ll}$, as depicted in Figure 8.9, it appears that these two variables are almost uncorrelated for the signal and the $q\bar{q} \rightarrow WW$ background for $\Delta\phi_{ll} \lesssim 2$. Some correlation is present in the background for $\Delta\phi_{ll} \gtrsim 2$, but this region can be rejected as its contribution to the total signal is very small.

Thus, as a selection cut to reduce the irreducible WW background, a cut on $\Delta\phi_{ll} < 1$ is applied (listed as cut 8a in Table 8.3). The distribution of the transverse mass m_T after this cut is shown in Figure 8.10 for a Higgs boson mass of $m_H = 130$ GeV and in Figure 8.11 for $m_H = 170$ GeV. As final signal region in the transverse mass m_T the range $m_H - 40 \text{ GeV} < m_T < m_H - 10 \text{ GeV}$ is chosen (listed in Table 8.3 as cut 9a for an analysis optimised for $m_H = 130$ GeV and cut 9b for $m_H = 170$ GeV). This last cut selects the Jacobian peak of the transverse mass m_T for the signal and can reduce the dominant WW background by a factor of 2 - 3 depending on the tested Higgs boson mass.

It should be noted, that due to this last cut the event selection depends on the Higgs boson mass. In the following, all statements about numbers of signal or background events after this cut will include the tested Higgs boson mass to indicate the used mass range in m_T .

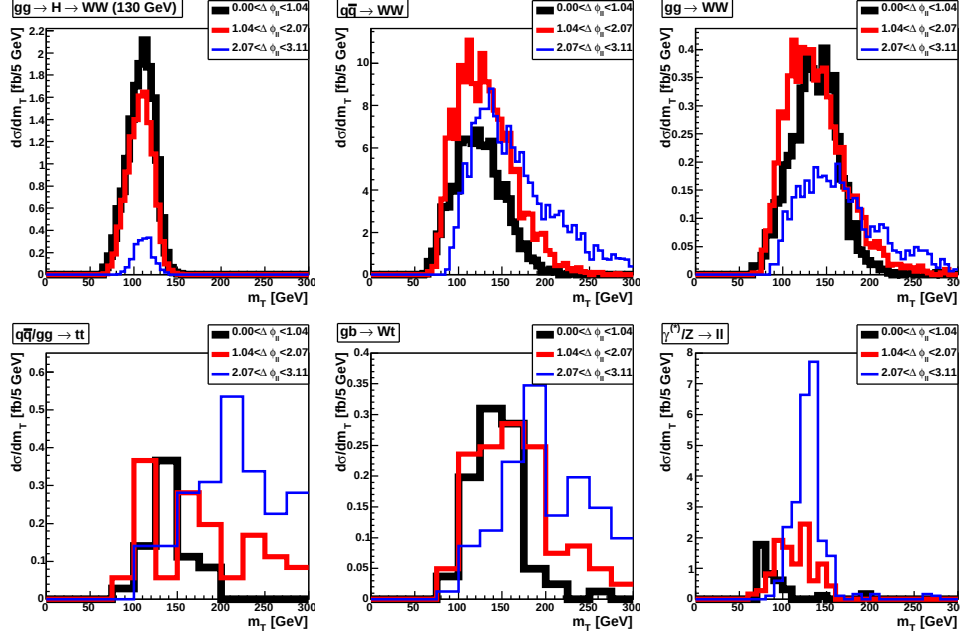


Figure 8.9: Dependency of the transverse mass m_T on the transverse opening angle $\Delta\phi_U$ for Higgs boson signal events ($m_H = 130$ GeV) and the dominant dilepton background processes after all cuts against reducible backgrounds. For each signal and background process the distribution of m_T is shown for three regions of $\Delta\phi_U$.

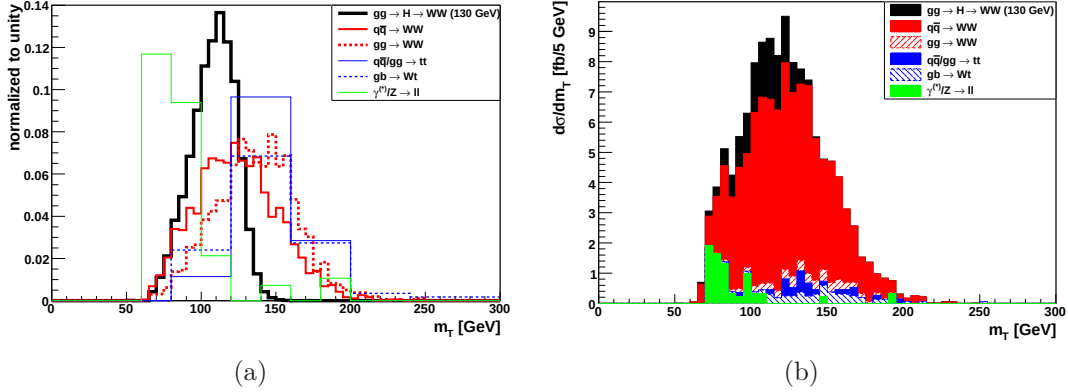


Figure 8.10: Distribution of the transverse mass m_T for Higgs boson signal events ($m_H = 130$ GeV) and the dominant dilepton background processes after the cut on $\Delta\phi_U < 1$. The normalised distribution is shown in (a), the total sum of all contributions to the m_T distribution in (b).

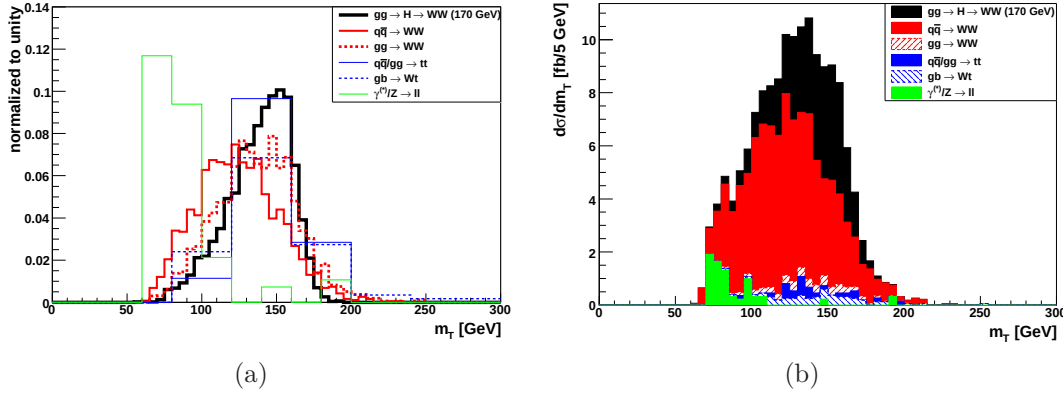


Figure 8.11: Distribution of the transverse mass m_T for Higgs boson signal events ($m_H = 170$ GeV) and the dominant dilepton background processes after the cut on $\Delta\phi_{ll} < 1$. The normalised distribution is shown in (a), the total sum of all contributions to the m_T distribution in (b).

Possible background control samples

The event selection cuts on $\Delta\phi_{ll}$ and m_T are effective in suppressing the backgrounds using the kinematical differences to the signal process. A modification or reversal of these cuts can be used for an almost signal free selection of some of the backgrounds:

- After the selection of $\Delta\phi_{ll} < 1$ a mass window cut on the invariant dilepton mass m_{ll} of $75 \text{ GeV} < m_{ll} < 105 \text{ GeV}$ can be used to control the size of the $\gamma^{(*)}/Z$ background. The current Monte Carlo prediction for this process gives a small contribution. However, a large remaining $\gamma^{(*)}/Z$ background due to an inefficient suppression from the $E_T^{\text{miss}} > 30 \text{ GeV}$ selection cut, would show up as clear peak in m_{ll} . This is listed as cut 8b in Table 8.4.
- The dominant WW background can be estimated by selecting the region $\Delta\phi_{ll} > 2$ (listed as cut 8c in Table 8.4). Due to the correlation of $\Delta\phi_{ll}$ and m_T in this region (see Figure 8.9), an additional cut on m_T is not applied. The size of the $\gamma^{(*)}/Z$ background in this control region can again be determined from the m_{ll} mass distribution. As this is intended as a simple cut based analysis to illustrate the main properties of the event selection, the details of how different backgrounds contribute to the individual control samples are not elaborated here.

8.1.4 Summary of the cut based analysis

A summary of all cuts and the accepted cross sections after the cuts for each signal and background is given in Table 8.3. For a Higgs boson mass of $m_H = 130$ GeV a signal-to-background ratio of $S/B \approx 1/4$ can be achieved. For a mass of $m_H = 170$ GeV the signal-to-background ratio is much better, mainly due to the larger $H \rightarrow WW$ branching ratio and the smaller background contribution at high transverse masses. Hence, a ratio of $S/B \approx 1$ is reachable.

8.2 Signal significance and systematic uncertainties

The aim of any Higgs boson analysis is to be as sensitive as possible to the possible presence of a Higgs boson signal. The sensitivity or the discovery significance is normally given by the improbability of the background processes to fluctuate in such a way, that it appears as the additional contribution of the expected signal (see Section 3.3).

In the event counting cut based analysis, a number of n_{S+B} events in the signal region, selected by the cuts given in Section 8.1, is measured. The number of possible signal events is extracted by subtracting the expected number of background events $\hat{n}_B$. Usually, $\hat{n}_B$ is determined from a separate measurement in some control region, as a direct prediction from Monte Carlo, in most cases, is not sufficiently precise.

In order to calculate the signal significance, the following three main sources of fluctuations to the expected number of background events $\hat{n}_B$ need to be considered:

- (i) The statistical fluctuation of the number of background events in the signal region. The fluctuation follows a Poissonian distribution with the number of expected background events $\hat{n}_B$ as central value. For the purpose of simple illustrations the approximation $\Delta(\hat{n}_B) = \sqrt{\hat{n}_B}$ is used in the following as statistical uncertainty on $\hat{n}_B$.
- (ii) The statistical fluctuation of the measured number of events n_C^i in each control region i , that is used to determine the expected number of background events $\hat{n}_B$. For the purpose of a simple illustration only the $\Delta\phi_U > 2$ control region for the WW background is considered here, again using the approximation $\Delta(n_C) = \sqrt{n_C}$.
- (iii) The systematic uncertainty $\Delta(f_{n_C \rightarrow \hat{n}_B})$ on the factor $f_{n_C \rightarrow \hat{n}_B} = \hat{n}_B/n_C$ that relates the measured number of events n_C in the control region to the expected number of background events $\hat{n}_B$ in the signal region.

While the effect of fluctuations of type (i) and (ii) on the significance is straightforward to evaluate, the uncertainty on the factor $f_{n_C \rightarrow \hat{n}_B}$ in item (iii) can be

Selection Cut (all cross sections in fb)	$H \rightarrow WW$ $m_H = 130$	$H \rightarrow WW$ $m_H = 170$	$q\bar{q} \rightarrow WW$	$gg \rightarrow WW$	$t\bar{t}$	Wt	$\gamma^{(*)}/Z$
1: 2 isolated leptons	76.4±0.4	200.7±1.4	1334±10	63.4±0.4	8993±41	672±6	751223±451
2: $p_T(l_1 + l_2) > 30$ GeV	65.1±0.3	183.1±1.4	931±8	50.4±0.4	7504±38	582±6	141561±196
3: $p_T(\text{jet}) < 20$ GeV	31.46±0.25	80.2±1.0	481±6	17.61±0.23	28.6±2.2	21.8±1.2	15754±66
4: $E_T^{\text{miss}} > 30$ GeV	28.99±0.24	77.6±0.9	439±6	16.09±0.22	24.9±2.1	20.3±1.1	91±5
5: $\Delta\phi_{ll} < 3$	28.95±0.24	77.5±0.9	426±5	15.82±0.22	23.5±2.0	19.7±1.1	78±5
6: $p_T(l_1 + l_2 + p_T^{\text{miss}}) < 60$ GeV	28.94±0.24	77.5±0.9	426±5	15.81±0.22	23.4±2.0	19.6±1.1	77±5
7: $\Delta\eta_{ll} < 3$	28.94±0.24	77.5±0.9	410±5	15.54±0.22	22.8±2.0	19.1±1.1	76±5
8a: $\Delta\phi_{ll} < 1$	14.78±0.17	43.2±0.7	91.4±2.5	4.91±0.12	3.1±0.8	4.5±0.5	7.2±2.1
9a: 90 GeV < m_T < 120 GeV	9.64±0.14	7.29±0.29	31.7±1.5	1.12±0.06	0.3±0.3	0.74±0.21	2.1±1.1
9b: 130 GeV < m_T < 160 GeV	0.78±0.04	23.6±0.5	27.8±1.4	2.02±0.08	1.5±0.5	2.0±0.4	0.22±0.22

Table 8.3: Accepted cross section in fb for the simple $H \rightarrow WW \rightarrow l\nu l\nu$ analysis after the application of the successive selection cuts as discussed in the text. The signal process is shown for Higgs boson masses of $m_H = 130$ GeV and 170 GeV. Only the dominant dilepton backgrounds are listed. The errors give the Monte Carlo statistical uncertainty.

Selection Cut (all cross sections in fb)	$H \rightarrow WW$ $m_H = 130$	$H \rightarrow WW$ $m_H = 170$	$q\bar{q} \rightarrow WW$	$gg \rightarrow WW$	$t\bar{t}$	Wt	$\gamma^{(*)}/Z$
8b: $\Delta\phi_{ll} < 1$ & 75 GeV < m_{ll} < 105 GeV	0.00±0.04	0.4±0.6	6.1±2.5	0.4±0.7	0.3±0.8	0.4±0.6	2.1±1.4
8c: $\Delta\phi_{ll} > 2$	2.22±0.07	4.76±0.23	169.9±3.4	4.37±0.12	13.2±2.7	8.4±0.7	53.3±4.4

Table 8.4: The same as Table 8.3 but for the background control sample selection cuts as discussed in the text.

difficult to assess, as it is not of a statistical nature. There is often no definite procedure to assign an exact number to uncertainties of this kind and quantitative statements will depend to some degree on choices made in the evaluation of the size of the primary fluctuations. Especially theory uncertainties on Monte Carlo predictions contributing to the evaluation of $f_{n_C \rightarrow \hat{n}_B}$ are in many cases non-Gaussian and depend on values assigned to unphysical parameters. To avoid that a possible discovery depends strongly on uncertainties of type (iii), they should be small compared to the statistical uncertainties (i) and (ii).

8.2.1 The influence of systematic uncertainties

For the simple $H \rightarrow W^-W^+ \rightarrow l^-\bar{\nu} l^+\nu$ analysis the signal region consists of events with small transverse lepton opening angle $\Delta\phi_{ll}$ and transverse masses m_T close to the Higgs boson mass (see Section 8.1.3). The number of signal events n_S in this region is given by the selection cuts 9a for the $m_H = 130$ GeV analysis and 9b for the $m_H = 170$ GeV analysis in the summary Table 8.3. As control region, in which the dominant WW background can be measured, events with a transverse lepton opening angle $\Delta\phi_{ll} > 2$ are selected (listed as cut 8c in Table 8.4). The control region for the $\gamma^{(*)}/Z$ background is ignored in the following. Hence the background extrapolation factor is given by:

$$f_{n_C \rightarrow \hat{n}_B} = \frac{\int d\sigma (\Delta\phi_{ll} < 1, m_H - 40 < m_T < m_H - 10 \text{ GeV})}{\int d\sigma (\Delta\phi_{ll} > 2)} . \quad (8.2)$$

Using the Gaussian approximation for all Poisson uncertainties, the significance S can be calculated using

$$S = \frac{n_S}{\sqrt{[\Delta(\hat{n}_B)]^2 + \left[\frac{\Delta(n_C)}{n_C} \cdot \hat{n}_B\right]^2 + \left[\frac{\Delta(f_{n_C \rightarrow \hat{n}_B})}{f_{n_C \rightarrow \hat{n}_B}} \cdot \hat{n}_B\right]^2}} . \quad (8.3)$$

The dependency of the significance S on the relative uncertainty on the extrapolation factor $\Delta(f_{n_C \rightarrow \hat{n}_B})/f_{n_C \rightarrow \hat{n}_B}$ is shown in Figure 8.12 for Higgs boson masses of $m_H = 130$ GeV and $m_H = 170$ GeV and assumed integrated luminosities of 1 and 10 fb^{-1} . Within the scope of the simple cut based analysis only the dominant dilepton backgrounds are included in the total background expectation $\hat{n}_B$ and the significance S is therefore overestimated.

For an integrated luminosity of 1 fb^{-1} the statistical uncertainties dominate and hence the dependency of the significance on the size of the uncertainty of $f_{n_C \rightarrow \hat{n}_B}$ is relatively small. On the other hand, the statistical uncertainty is too large to reach a 5σ discovery. For an integrated luminosity of 10 fb^{-1} the dependency on $f_{n_C \rightarrow \hat{n}_B}$ is more pronounced and already a relative uncertainty of 5% gives a sizeable reduction of the significance by 25%. As stated above, the uncertainty on $f_{n_C \rightarrow \hat{n}_B}$ is not always well defined and similarly the magnitude of a sizeable

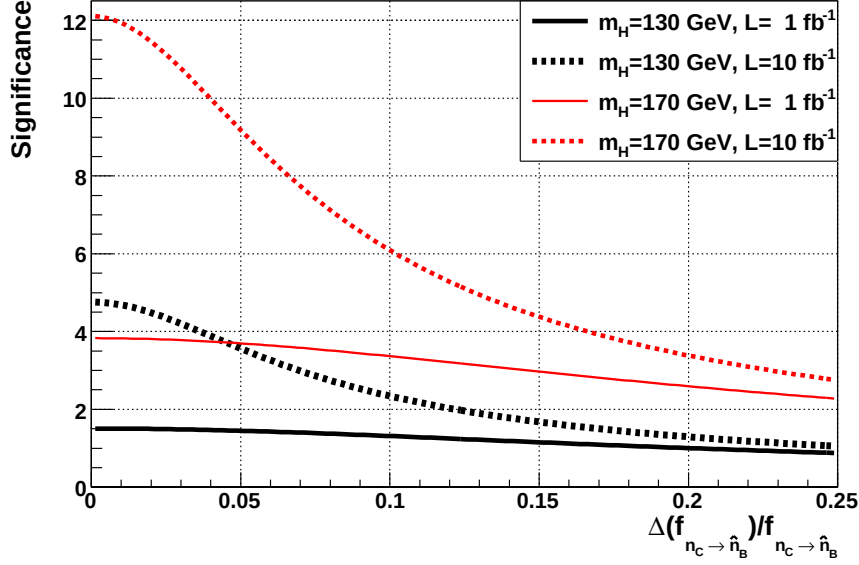


Figure 8.12: Dependency of the significance S on the systematic uncertainty of the background extrapolation factor $f_{n_C \rightarrow \hat{n}_B}$ for Higgs boson masses of $m_H = 130$ GeV and $m_H = 170$ GeV and for assumed integrated luminosities of 1 and 10 fb^{-1} . This significance is optimistic, as only the dominant dilepton backgrounds are included in the calculation.

reduction of the significance is not well defined. However, a relative uncertainty of $\approx 5\%$ on $f_{n_C \rightarrow \hat{n}_B}$ and hence 25% on S seems to be a sensible upper limit which the total systematic background uncertainty should not exceed.

In the following sections different sources of systematic uncertainties on $f_{n_C \rightarrow \hat{n}_B}$ are briefly illustrated and the corresponding precisions are discussed which are needed to stay within the $\approx 5\%$ limit. The WW background is clearly the dominant background source, so the following systematic uncertainties on the total background measurement are considered:

- The systematic uncertainties on the shape of the transverse opening angle $\Delta\phi_u$ and the transverse mass m_T for the $q\bar{q} \rightarrow WW$ background.
- The systematic uncertainties on the relative sizes of the top backgrounds $t\bar{t}$ and Wt compared to the $q\bar{q} \rightarrow WW$ background.
- The systematic uncertainties on the relative size of the $gg \rightarrow WW$ background compared to the $q\bar{q} \rightarrow WW$ background.

In order to study uncertainties on a level of 5% systematic variation, large Monte Carlo samples are needed to reduce the Monte Carlo statistical uncertainty well below the 5% systematic limit. In many cases these large samples are not

available with the full ATLAS simulation and fast simulation methods need to be used. Wherever fast simulations are used within this chapter, it is mentioned. The comparisons between fast and full simulation that show the validity of this approach are shown, where needed, in Appendix B.

8.3 Systematic uncertainties on the shape of the $q\bar{q} \rightarrow WW$ background

The shapes of the transverse opening angle $\Delta\phi_u$ and of the transverse mass m_T for the dominant WW background need to be determined from Monte Carlo simulations. Two main factors have to be considered:

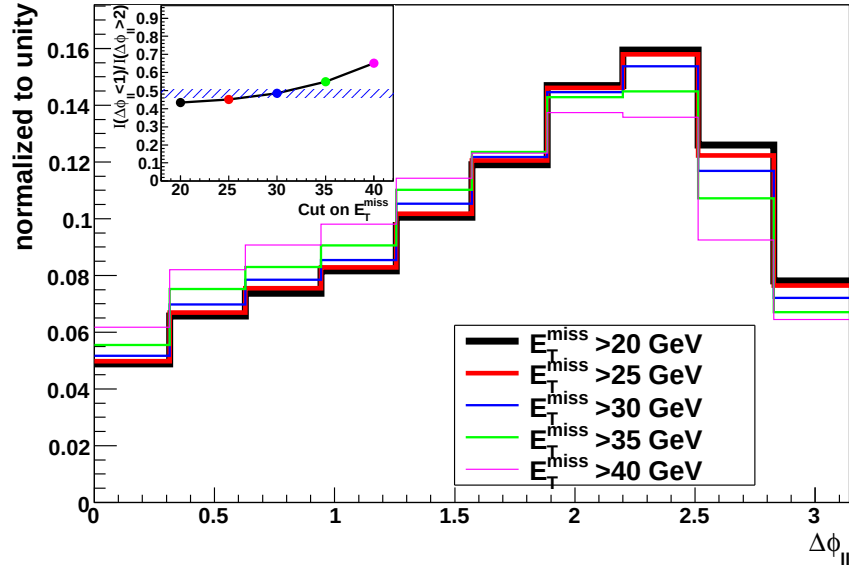
- The influence of the cuts on $p_T(l_1 + l_2)$ and E_T^{miss} on the shape of $\Delta\phi_u$.
- The intrinsic uncertainty on the shape of $\Delta\phi_u$ and m_T due to the primary event generator.

8.3.1 Uncertainty on $\Delta\phi_u$ from the missing transverse momentum

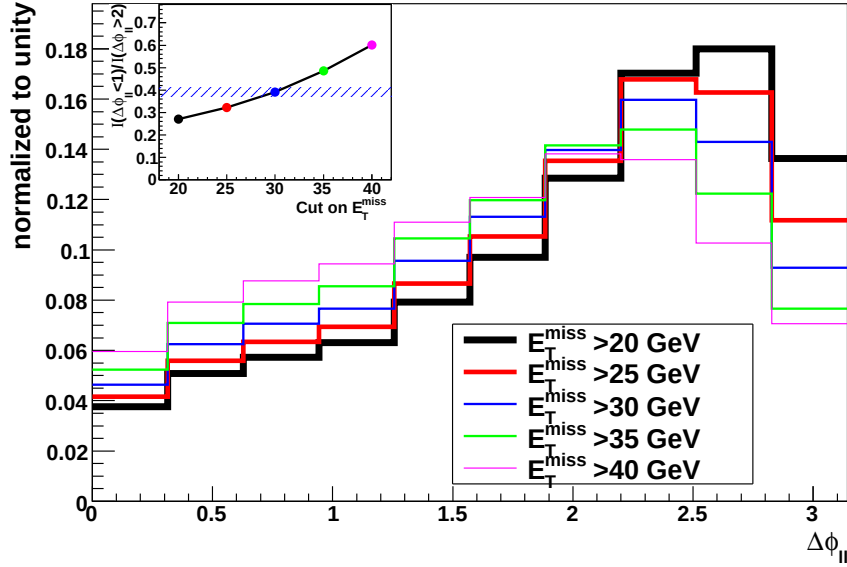
The selection cuts on $p_T(l_1 + l_2)$ and E_T^{miss} strongly influence the shape of $\Delta\phi_u$, because events passing these cuts are biased towards small opening angles where the leptons and neutrino momenta add up in the transverse direction. For large $\Delta\phi_u$ the momenta often counterbalance each other and the events fail the selection cuts.

After the application of the jet veto the uncertainty on the determination of the missing transverse momentum is dominated by low p_T physics processes (see Section 2.5.5) and the calibration of soft energy depositions in the calorimeter. The presence of pile-up (see Section 2.5.6) can lead to further complications.

The missing transverse momentum is one of the most difficult event properties to measure and calibrate, and hence predictions for the performance of E_T^{miss} in the early data taking phase are almost impossible to make before data taking actually starts. Uncertainties can affect the resolution and scale of E_T^{miss} , but for the prediction of $\Delta\phi_u$ from Monte Carlo, uncertainties on the scale are most important. Figure 8.13(a) shows the variation of the $\Delta\phi_u$ distribution under variation of the E_T^{miss} cut value by ± 5 GeV and ± 10 GeV. These variations are most probably unrealistically large, but as the inset plot shows, already variations of 2 - 3 GeV can change the ratio of events with $\Delta\phi_u < 1$ to $\Delta\phi_u > 2$ by $\approx 5\%$. Figure 8.13(b) shows that for the same distribution without the cut on $p_T(l_1 + l_2) > 30$ GeV the effect of E_T^{miss} scale uncertainties would be strongly enhanced.



(a)



(b)

Figure 8.13: Normalised distribution of $\Delta\phi_{ll}$ after the lepton selection and a jet veto cut as function of the E_T^{miss} cut value for the dominant $q\bar{q} \rightarrow WW$ background with (a) and without (b) the cut on $p_T(l_1 + l_2) > 30$ GeV. The inset plot shows how the ratio of the number of events with $\Delta\phi_{ll} < 1$ to the number of events with $\Delta\phi_{ll} > 2$ depends on the E_T^{miss} cut value.

Figure 8.14 shows the dependency of $f_{n_C \rightarrow \hat{n}_B}$ only for the $q\bar{q} \rightarrow WW$ background after all cuts for Higgs boson masses of $m_H = 130$ GeV and $m_H = 170$ GeV⁷. A variation of $\approx 5\%$ on $f_{n_C \rightarrow \hat{n}_B}$ is reached for uncertainties on the E_T^{miss} scale of $\approx 2 - 3$ GeV. For low true E_T^{miss} values, as they are relevant for the efficiency of this selection cut, this is a realistic uncertainty even in the case of a well performing detector, as can be seen in Reference [15].

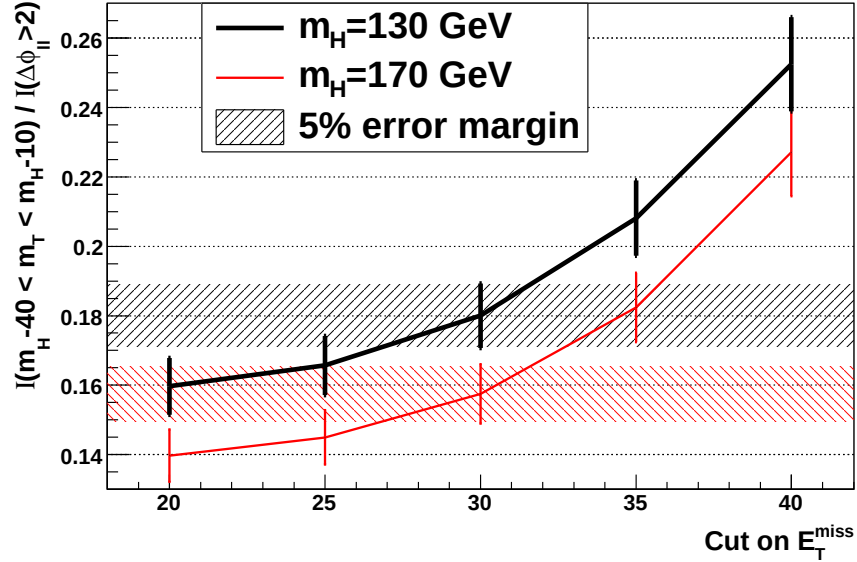


Figure 8.14: Systematic variation of the ratio $f_{n_C \rightarrow \hat{n}_B}$ for the dominant $q\bar{q} \rightarrow WW$ background after all event selection cuts as function of the variation of the E_T^{miss} cut value. The analysis for $m_H = 130$ GeV is shown in black and for $m_H = 170$ GeV in red. The light hashed area indicates the size of a 5% variation of the central $f_{n_C \rightarrow \hat{n}_B}$ value. The error bars give the Monte Carlo statistical uncertainty.

The performance of the E_T^{miss} reconstruction can be best tested with measurements on $Z \rightarrow ll$ events, where no intrinsic E_T^{miss} is expected. If these measurements indicate large uncertainties on the E_T^{miss} scale, asymmetric cut values for $p_T(l_1 + l_2)$ and E_T^{miss} (for example $p_T(l_1 + l_2) > 30$ GeV, $E_T^{\text{miss}} > 20$ GeV) could be used to remove some of the dependency of the analysis cuts on the E_T^{miss} scale. This will of course reduce the significance, either due to an increase of the number of background events or a reduction of signal events.

⁷It should be noted that the variation of E_T^{miss} was not propagated into the transverse mass m_T , as this would only cause a shift of the reconstructed mass. While such a shift changes the number of reconstructed events inside the fixed mass window, it can easily be corrected by changing the mass window position and would thus suggest a different Higgs boson mass.

8.3.2 Generator uncertainties on the shape of the $\Delta\phi_{ll}$ and m_T distributions

The best available predictions for the $q\bar{q} \rightarrow WW$ process are obtained with MC@NLO using a NLO matrix element matched to a parton shower (see Appendix B.2.1). Remaining uncertainties on the prediction of $\Delta\phi_{ll}$ can be analysed by varying the PDFs and the choices for the renormalization (μ_R) and factorization (μ_F) scale (see Chapter 2). To draw significant conclusions from such variations, very large Monte Carlo datasets are needed (altogether $\approx 500\text{M}$ events). As these variations concern only the primary event generator, datasets simulated with Atlast-I (see Section 6.1.2) with a collision energy of $\sqrt{s} = 14$ TeV were used. However, this choice does not bias the conclusions drawn in the following sections, as the distributions of the Atlast-I events are very similar (see Figure B.2) and only the total cross section, which is irrelevant for this study, shows differences.

For the variation of the PDFs, the CTEQ group has published a set of 40 error PDFs that correspond to $\pm 1\sigma$ (68%) variations along the 20 eigenvectors of the error matrix of the PDF fit. Figure 8.15 shows that the ratio of the number of events with $\Delta\phi_{ll} < 1$ to the number of events with $\Delta\phi_{ll} > 2$ is within Monte Carlo statistical uncertainties not affected by the choice of PDFs.

Figure 8.16 shows the same ratio as function of the renormalization scale μ_R and factorization scale μ_F . μ_R and μ_F are varied symmetrically ($\mu_F = \mu_R$) and asymmetrical ($\mu_F = 1/\mu_R$) by factors of up to 8. For factors of 8 the ratio varies by about 3%. However, it is arguable what uncertainties should be assigned to such a variation⁸ of $\mu_{R,F}$, but the factor of 8 is considerably larger than the canonical factor 2, which is usually chosen.

The systematic uncertainty on $f_{n_C \rightarrow \hat{n}_B}$ after all cuts is shown in Figures 8.17 and 8.18. Due to the increased Monte Carlo statistical uncertainty the size of the variation of $f_{n_C \rightarrow \hat{n}_B}$ cannot be extracted exactly. However, the variations remain all within 5% of the central value and thus PDF and QCD scale uncertainties seem to be uncritical.

8.4 Systematic uncertainties from top backgrounds

Figures 8.7(a) and 8.8(a) show that the top backgrounds $t\bar{t}$ and Wt have a distinctly different shape in the transverse opening angle $\Delta\phi_{ll}$ and the transverse mass m_T compared to the dominant $q\bar{q} \rightarrow WW$ background. A change of the relative contribution of the top backgrounds to the total background will thus

⁸Unfortunately, there seems to be no way to assign a quantitative number, as theory uncertainties are definitely non-Gaussian. The only reliable way to calculate the NLO uncertainty would be to calculate the process at least at NNLO, where the whole argument is continued at the next level.

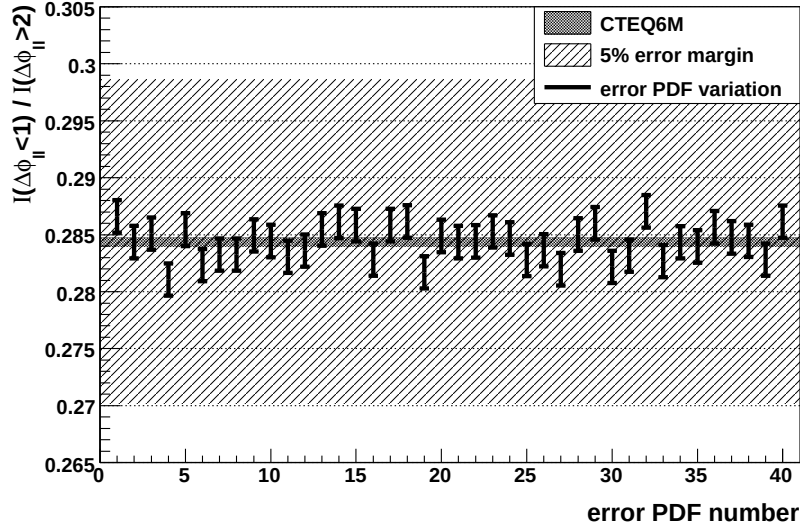


Figure 8.15: Systematic uncertainty on the ratio of the number of events with $\Delta\phi_u < 1$ to the number of events with $\Delta\phi_u > 2$ after the lepton selection cuts (cut 1) due to PDF variations of the $q\bar{q} \rightarrow WW$ background using the 40 CTEQ6M 1σ error PDFs. The central dark line shows the central prediction of the ratio, the width of the line is the Monte Carlo statistical uncertainty. The error bars give the prediction after variation, the size of the error bar is the Monte Carlo statistical uncertainty. The light hashed area indicates the size of a 5% variation of the central ratio.

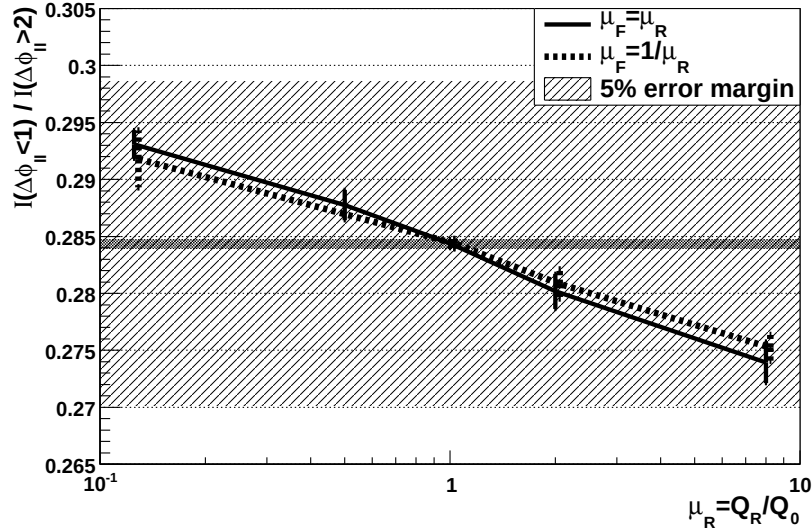


Figure 8.16: The same as Figure 8.15 but for a variation of the QCD renormalization (μ_R) and factorization (μ_F) scale. Considered are the cases $\mu_F = \mu_R$ and $\mu_F = 1/\mu_R$.

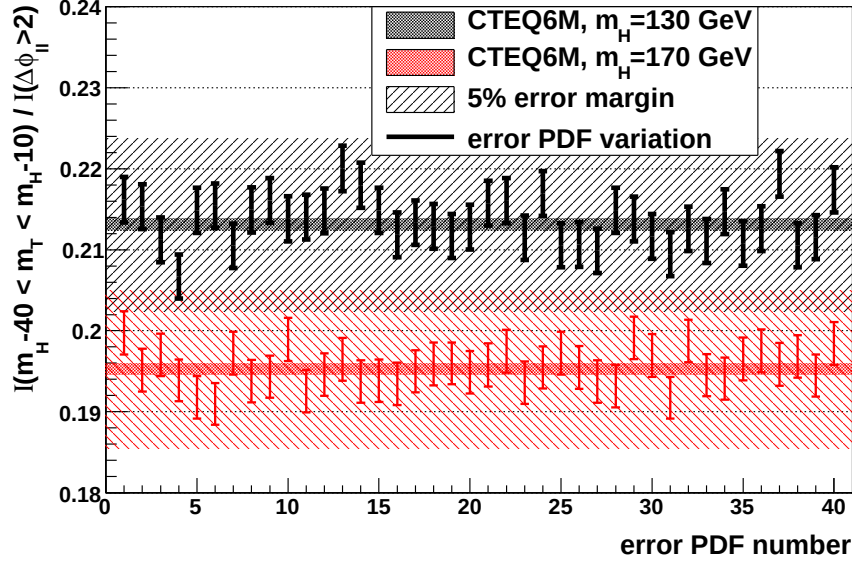


Figure 8.17: Systematic uncertainty on the ratio $f_{n_C \rightarrow \hat{n}_B}$ due to PDF variations of the $q\bar{q} \rightarrow WW$ background after all event selection cuts. The analysis for $m_H = 130$ GeV is shown in black and for $m_H = 170$ GeV in red. The same variations as in Figure 8.15 are applied. The PDF variations are extracted from high statistics $\sqrt{s} = 14$ TeV Atlfast-I samples. Therefore the central value is different to Figure 8.14.

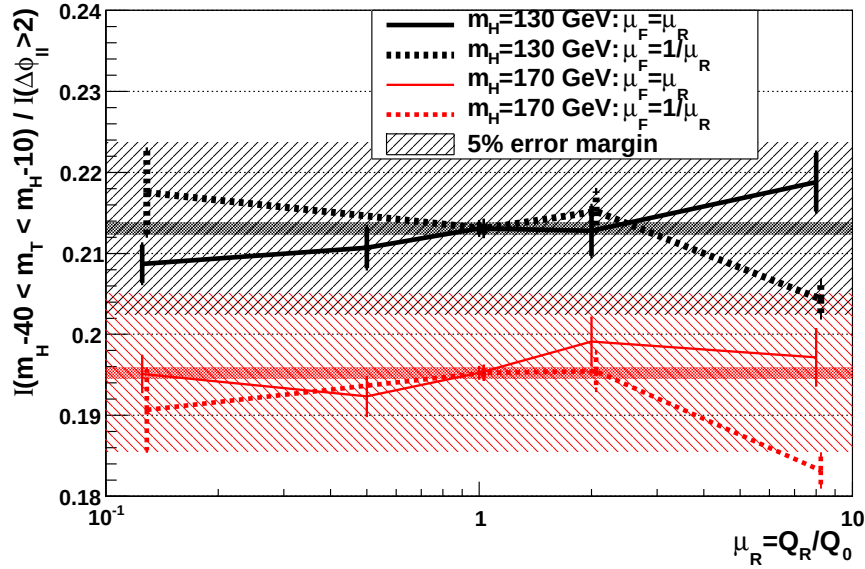


Figure 8.18: The same as Figure 8.17 but for QCD scale variations. The same variations as in Figure 8.16 are applied.

also affect the expected normalisation factor $f_{n_C \rightarrow \hat{n}_B}$, that is relying on precise Monte Carlo knowledge of the shape of the sum of all backgrounds. Fortunately, the $t\bar{t}$ and Wt backgrounds have a sufficiently similar shape in $\Delta\phi_{ll}$ and m_T among each other (see Appendix B.2.4), such that a combined treatment of both backgrounds is possible, as far as their influence on the main control region is concerned.

The top backgrounds are mainly affected by the jet veto cut at $p_T > 20$ GeV that removes the largest fraction of the events, such that only $\approx 0.4\%$ ($t\bar{t}$) respectively $\approx 4\%$ (Wt) of the events remain after this selection cut (see Section 8.1.2). Due to this large rejection, already small changes in the jet activity can have a large effect on the accepted cross sections:

- Especially low p_T jets have large systematic uncertainties on the jet energy scale and jet efficiency [38]. While ATLAS aims at an overall final jet energy scale precision of 1%, in the following a systematic uncertainty of the order of 10% on the jet energy scale for $p_T \approx 20$ GeV jets is assumed, as it is very difficult to calibrate these low p_T jets and without data realistic estimates of the jet energy scale uncertainty are not yet available.

Figure 8.19 shows that the accepted cross section of the top backgrounds varies by $\approx 50\%$ when the jet energy scale at 20 GeV is changed by $\approx 10\%$, while the dominant $q\bar{q} \rightarrow WW$ background is almost constant.

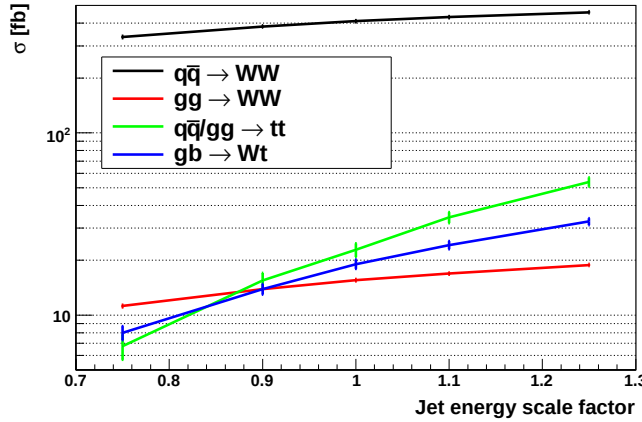


Figure 8.19: Dependency of the accepted cross section after all cuts against reducible backgrounds on the jet energy scale for the dominant dilepton background contributions.

- There are no firm predictions from the event generators for the low p_T jet activity and no clear strategy of assigning a systematic uncertainty to the current predictions. It is planned that LHC measurements will be used to

improve the description of jet activity by the generators, but one has to expect that these improvements will take a considerable amount of time.

- Non-resonant Wb contributions to both processes are not included in the Monte Carlo simulations but could be enhanced by the jet veto and show a higher jet veto survival probability than the top quark resonant processes (see Appendix B.2.4). In the parton level study in Reference [39] a relative increase by a factor ≈ 2 of the accepted LO $t\bar{t}$ cross section in the full calculation is found after a jet veto cut of $p_T(\text{jet}) > 15$ GeV and $|\eta(\text{jet})| < 3.2$.
- The relative contribution of $t\bar{t}$ to Wt is also subject to the jet uncertainties above and thus independent estimates or measurements for both backgrounds are needed, as far as it is possible to treat these backgrounds independently (see Appendix B.2.4).

While it is possible to define control samples for the top backgrounds, especially when using b -tagging tools to identify b -jets, these control samples can have a different composition of $t\bar{t}$ and Wt than the events in the signal region and hence the extrapolation to the phase space region with a complete jet veto needs to use Monte Carlo simulations again. Furthermore it can be expected that top control samples have different non-resonant contributions than the top events in the signal region.

In the following it is not tried to assign quantitative uncertainties to the effects above, as there is no firm basis to do this. Figure 8.20 shows the systematic uncertainty on $f_{n_C \rightarrow \hat{n}_B}$ from variations of the accepted top background cross section by one order of magnitude. While a reduced cross sections is uncritical, an increase by an approximate factor of 2-3 is sufficient to change $f_{n_C \rightarrow \hat{n}_B}$ by 5%.

In this discussion it should be noted that jets with a p_T of ≈ 20 GeV could also be produced from underlying event activity and pile-up. If the jet multiplicity due to these effects causes a too large rejection of the signal process, the jet veto threshold will need to be raised. In this case top backgrounds become more important rapidly and the systematic uncertainty gets enhanced by the factor by which the accepted top background cross section increases due to the raised cut value.

8.5 Systematic uncertainties from the $gg \rightarrow WW$ background

The loop induced $gg \rightarrow WW$ background, strictly seen, is part of the NNLO correction (see Section 2.3) to the inclusive $q\bar{q} \rightarrow WW$ processes. The Feynman diagrams for both processes are shown in Figure 8.21.

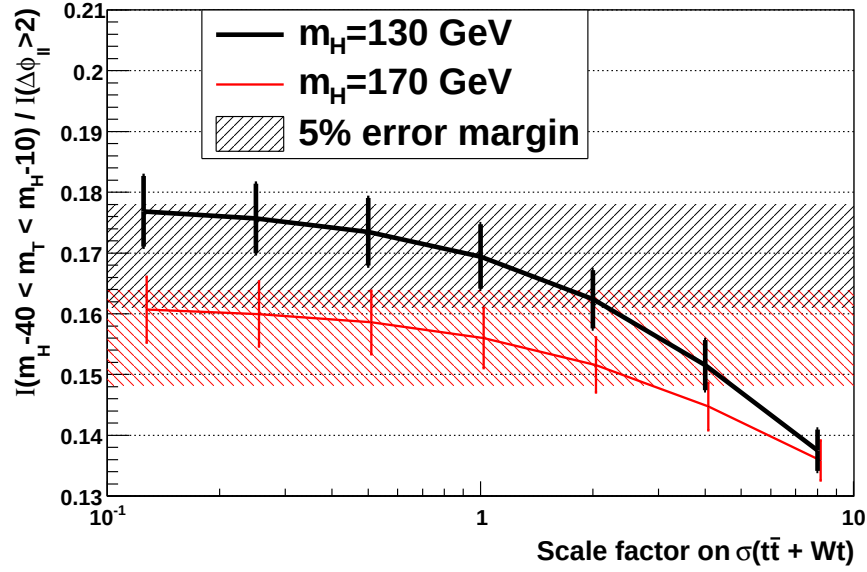


Figure 8.20: Systematic variation of the ratio $f_{n_C \rightarrow \hat{n}_B}$ for the sum of all backgrounds due to a variation of the size of the top backgrounds $t\bar{t}$ and Wt . Otherwise the same as Figure 8.14.

As part of the NNLO corrections the $gg \rightarrow WW$ process normally would not be considered as a separate background, but just added as one contribution to the total correction. However, so far the $q\bar{q} \rightarrow WW$ process is calculated only up to NLO and a complete NNLO calculation is not available yet. Due to the large gluon luminosity at the LHC, the cross section of the $gg \rightarrow WW$ contribution is enhanced by a large factor. Furthermore, this is the lowest order for the gluon induced production of WW and thus the process gives a finite cross section, such that an individual consideration of this contribution is possible. This is also necessary, because any uncertainties derived from the calculation of $q\bar{q} \rightarrow WW$ does not include a possible contribution from $gg \rightarrow WW$.

The $gg \rightarrow WW$ process was first calculated and discussed in Reference [35] and in Reference [40] as part of this thesis and its relevance at the LHC assessed in the context of the $gg \rightarrow H \rightarrow WW$ analysis. Before selection cuts, the $gg \rightarrow WW$ process contributes only $\approx 4\%$ to the total $q\bar{q} \rightarrow WW$ cross section (see Table 8.1). However, the kinematic distribution of the $gg \rightarrow WW$ events is more signal like than the $q\bar{q} \rightarrow WW$ background and thus the event selection cuts in Section 8.1.3 increase the relative contribution of $gg \rightarrow WW$ to almost 10%.

For the evaluation of the systematic uncertainty due to the $gg \rightarrow WW$ contribution, the following two factors need to be considered:

- The calculation of the $gg \rightarrow WW$ cross section is at LO and thus not very precise. A NLO calculation is not available and will most likely not be

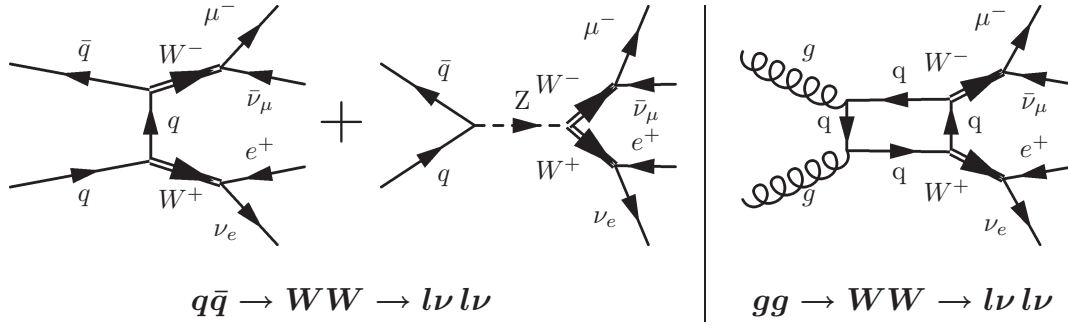


Figure 8.21: Feynman diagrams for the dominant $q\bar{q} \rightarrow WW$ background (left) and the $gg \rightarrow WW$ contribution (right).

available for the next years. A systematic uncertainty on the cross section can be derived from the usual scale variations, but one should keep in mind that a variation of the renormalization scale μ_R only affects α_s and a variation of the factorization scale μ_F only enters the evaluation of the gluon luminosity. The matrix element itself is not affected. Numerically, the variation of both QCD scales by a factor of two results in a change of the cross section by $\pm 25\%$ [35]. However, it is arguable how predictive this uncertainty is. Comparisons to related processes, that are calculated at NLO, indicate that an uncertainty of $\approx 50\%$ on the leading order cross section is justified.

- For all processes but $gg \rightarrow WW$ the leading p_T jet is generated either with a matrix element ($gg \rightarrow H \rightarrow WW$, $q\bar{q} \rightarrow WW$, $t\bar{t}$, Wt) or with a matrix element correction to the parton shower ($\gamma^{(*)}/Z$) (see Appendix B).

For $gg \rightarrow WW$ no matrix element correction is available and hence the leading jet p_T depends only on the parton shower parameters. The Herwig parton shower used for the $gg \rightarrow WW$ Monte Carlo simulation predicts that $\approx 35\%$ of the events are accepted by the jet veto cut of $p_T < 20$ GeV, while $\approx 50\%$ are accepted for $q\bar{q} \rightarrow WW$ and $gg \rightarrow H \rightarrow WW$ (see Figure 8.4(b)). In addition, when using independently generated events using the Pythia parton shower for $gg \rightarrow WW$, the jet veto efficiency for $gg \rightarrow WW$ and $q\bar{q} \rightarrow WW$ is approximately identical. As long as no sufficiently precise generator predictions for the low p_T jet activity are available, this causes an uncertainty of $\approx 30\%$, if the difference between the generators is taken as a scale of the uncertainty.

Both sources of uncertainty on the accepted cross section would be mitigated, if an isolated measurement of the $gg \rightarrow WW$ process were possible in some control region. Due to the large cross section of the other backgrounds, a measurement of $gg \rightarrow WW$ is only possible, if at all, after all cuts against reducible backgrounds.

In order to locate a phase space region where an enhanced fraction of $gg \rightarrow WW$ events is observable, the correlation between $\Delta\phi_{ll}$ and m_T is shown in Figure 8.22 for Higgs boson signal events ($m_H = 170$ GeV) and the $gg \rightarrow WW$ and $q\bar{q} \rightarrow WW$ background processes. For the illustration of this correlation high Monte Carlo statistics $\sqrt{s} = 14$ TeV Atfast-I samples are used.

For small transverse opening angles $\Delta\phi_{ll}$, the shape of the $gg \rightarrow WW$ process is almost identical to the shape of the signal process for a Higgs boson mass of $m_H \approx 170$ GeV, as it is also visible in the m_T projection in Figure 8.11. Hence it is very difficult to distinguish between both processes in this region and essential to know the $gg \rightarrow WW$ cross section exactly to avoid a false identification of this background contribution as signal.

In the region $1.5 \lesssim \Delta\phi_{ll} \lesssim 2.5$, where a significant difference between the Higgs boson signal process and the $gg \rightarrow WW$ background is visible, the $q\bar{q} \rightarrow WW$ process has its dominant contribution and is a factor of ≈ 50 larger than the $gg \rightarrow WW$ process. Thus a measurement of the $gg \rightarrow WW$ cross section in this region would require an absolute knowledge, either from Monte Carlo predictions or independent measurements, of the $q\bar{q} \rightarrow WW$ cross section at the 2% level or better, which is unrealistic for the first years of data taking.

Therefore it is necessary to rely on the Monte Carlo prediction of the $gg \rightarrow WW$ cross section to estimate the size of this background. Figure 8.23 shows the systematic uncertainty on the background extrapolation factor $f_{n_C \rightarrow \hat{n}_B}$ from a variation of the $gg \rightarrow WW$ cross section by one order of magnitude. Especially the analysis looking for a Higgs boson with a mass of $m_H = 170$ GeV is very sensitive to the $gg \rightarrow WW$ process and the cross section needs to be known from Monte Carlo simulations within a factor of 2 to stay within the 5% uncertainty limit.

8.6 The challenge of the background normalisation

In the previous Sections 8.3 - 8.5 the dominant systematic uncertainties on the $q\bar{q} \rightarrow WW$, $t\bar{t}$, Wt and $gg \rightarrow WW$ backgrounds were shown and compared to the acceptable limit of $\approx 5\%$ overall systematic uncertainty on the background normalisation.

Several sources of systematic uncertainty are close to this limit and will need very detailed control measurements to make sure that the backgrounds are sufficiently well known. It is likely that with increasing experience during the LHC data taking phase, such control measurements will be possible for most uncertainties. However, the question will remain if all uncertainties are correctly, and sufficiently precisely, estimated. Furthermore, all systematic uncertainties add up and the total uncertainty should stay below $\approx 5\%$. This puts even tighter bounds on the individual sources of systematic uncertainties.

In the case of the $gg \rightarrow WW$ background, no clear strategy of measuring the

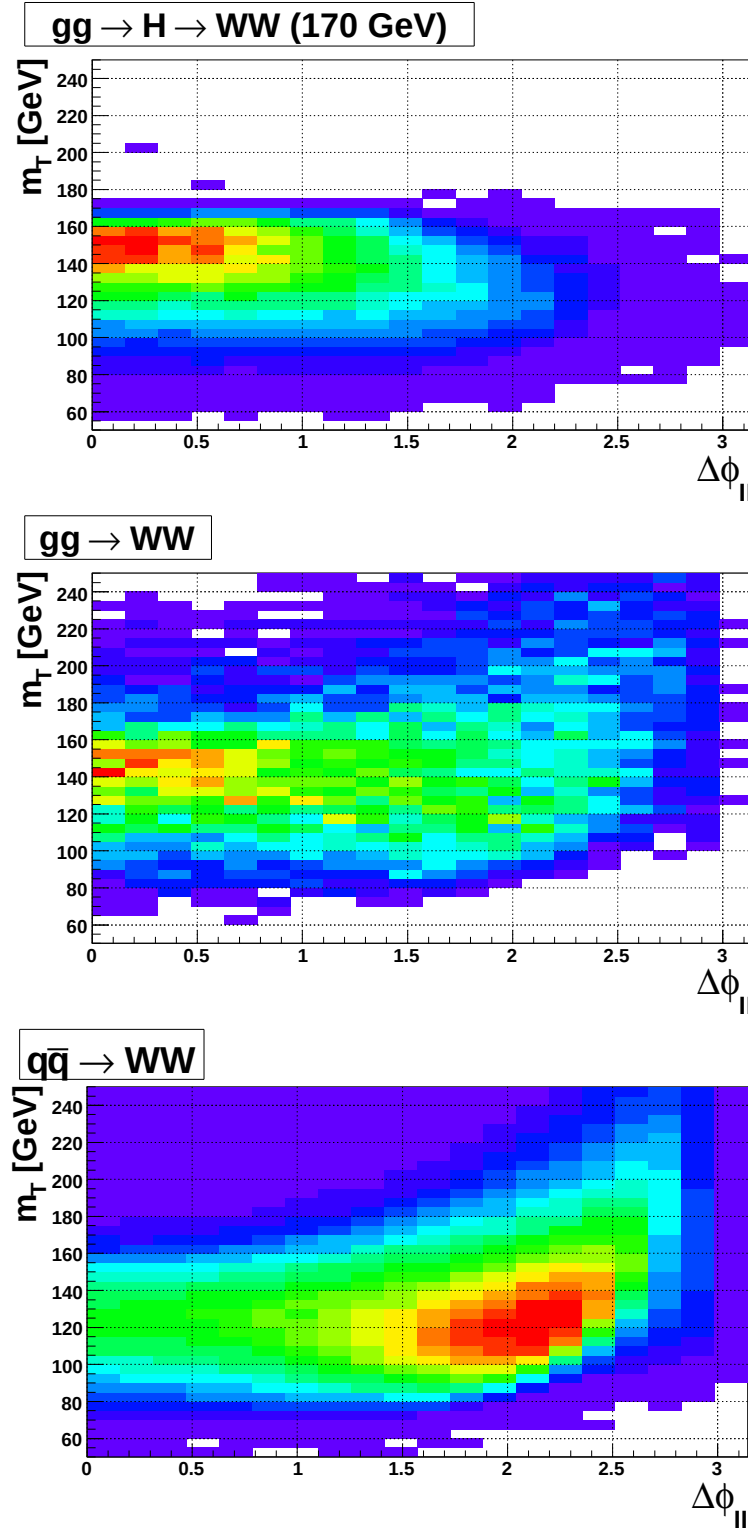


Figure 8.22: Correlation between $\Delta\phi_{II}$ and m_T after all cuts against reducible backgrounds for Higgs boson signal events ($m_H = 170$ GeV) and the $gg \rightarrow WW$ and $q\bar{q} \rightarrow WW$ background processes. High Monte Carlo statistics $\sqrt{s} = 14$ TeV Atfast-I samples are used.

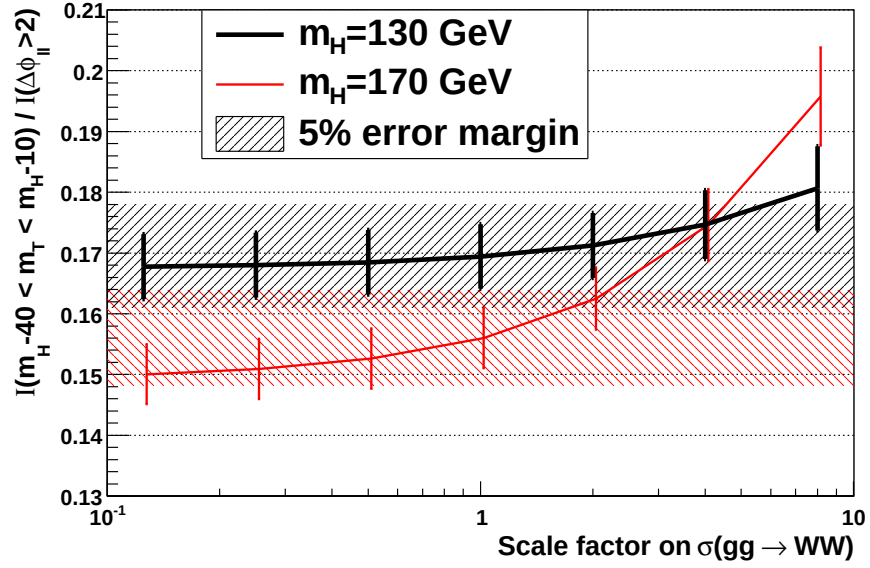


Figure 8.23: Systematic variation of the ratio $f_{n_C \rightarrow \hat{n}_B}$ for the sum of all backgrounds due to a variation of the cross section of the $gg \rightarrow WW$ background. Otherwise the same as Figure 8.14.

cross section is established and the used value and uncertainty relies exclusively on Monte Carlo predictions. Especially for a background source that is very similar to the signal this is an unsatisfactory situation.

9 Measurement of the shape and rate of the dominant dilepton backgrounds

The previous chapter illustrated that the normalisation of the background in the $gg \rightarrow H \rightarrow WW$ channel is challenging and can significantly limit the discovery potential. To overcome this problem, a background normalisation strategy is needed that does not rely on Monte Carlo predictions for the background extrapolation into the signal region, but takes all factors, in an as unbiased way as possible, from the measured distributions themselves. In the optimal case, this is equivalent to a measurement of the shape and absolute rate of the backgrounds.

In order to illustrate the background normalisation method, very large Monte Carlo samples are needed. Therefore the numbers and distributions in this chapter were obtained with Monte Carlo samples corresponding to an integrated luminosity of $\gtrsim 100 \text{ fb}^{-1}$ and are shown at the generator level without any detector simulation. Because of the high integrated luminosity, the samples are produced with the foreseen final LHC centre-of-mass energy of $\sqrt{s} = 14 \text{ TeV}$.

9.1 Separation into exclusively longitudinal and exclusively transverse variables

After the selection cuts against the reducible backgrounds, the main discrimination power between the signal and the background processes comes from the transverse $\Delta\phi_{ll}$ and m_T distributions. If one aims at measuring the shape of these variables for the backgrounds, one needs a variable or set of variables, that

- is mostly uncorrelated to $\Delta\phi_{ll}$ and m_T , to allow an unbiased extraction of the shape
- behaves differently for the signal and background processes. Otherwise the extracted shape in the normalisation region contains the same fraction of signal events as in the signal region and possible signal events are included in the measured background shape.

Due to the correlations between all transverse variables it is unlikely to find another transverse variable that fulfils the above mentioned criteria. However, so

far longitudinal variables are not used in the final event selection¹ and thus offer a possibility for uncorrelated event properties. However, one needs to keep in mind that in the longitudinal direction no components of the neutrinos from the W decay is measurable and the longitudinal W momentum cannot be reconstructed.

9.1.1 The rapidity of the WW system

The principal longitudinal event property is the rapidity

$$Y = \frac{1}{2} \ln \frac{E + p_z}{E - p_z} \quad (9.1)$$

of the WW system, as it is shown in Figure 9.1(a) for the Higgs boson signal at $m_H = 130$ GeV and for the $q\bar{q} \rightarrow WW$, $gg \rightarrow WW$ and $t\bar{t}$ backgrounds before any cuts are applied.

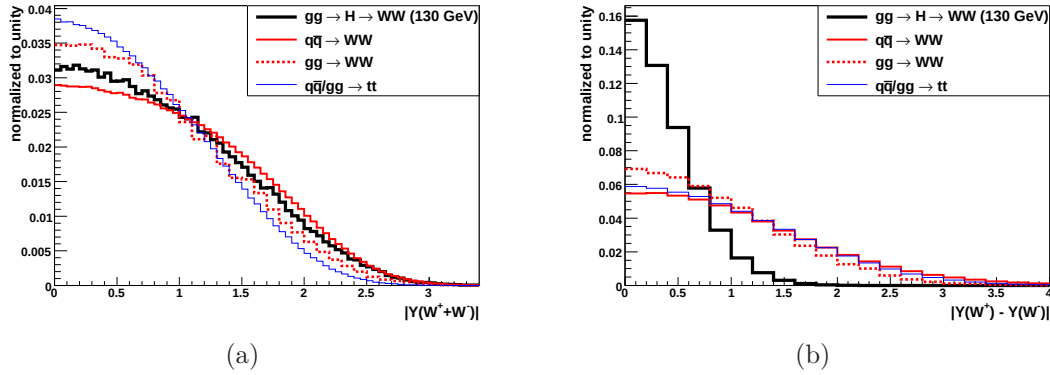


Figure 9.1: Normalised distributions of the rapidity $|Y(W^+ + W^-)|$ (a) and of the rapidity difference $|Y(W^+) - Y(W^-)|$ (b) of the WW system for the Higgs boson signal process at $m_H = 130$ GeV and for the dominant non-resonant backgrounds.

The rapidity $Y(W^+ + W^-)$ of the produced WW system is mainly determined by the proton PDFs (see Section 2.2). In order to produce a WW system with a mass of ≈ 160 GeV at a $\sqrt{s} \approx 10$ TeV collider, parton momentum fractions of $x \approx 10^{-2}$ are needed. The heavier the system gets, the larger are the average x values. For non zero rapidities $Y(W^+ + W^-)$ the x values of the partons need to be asymmetric. As can be seen in Figure 2.1, the PDFs in general are steeply falling functions with increasing x , so large x values are improbable and thus large rapidities of the WW system are disfavoured. Hence, for the gluon initiated processes, the lightest WW system from the Higgs boson decay has the widest rapidity distribution and the heaviest system from $t\bar{t}$ the narrowest.

¹The definition of the transverse mass m_T is deliberately chosen to be exclusively transverse at the cost of a slightly reduced separation power of the signal from the backgrounds (see Appendix C.3).

The production of $q\bar{q} \rightarrow WW$ involves one valence quark with a maximum in the PDFs for $x \approx 10^{-1}$. This leads to a wider rapidity distribution than for the signal process, although the WW system mass for $q\bar{q} \rightarrow WW$ is in general considerably higher than for $gg \rightarrow H \rightarrow WW$.

The dependency of the shapes of the transverse variables $\Delta\phi_U$ and m_T on $|Y(W^+ + W^-)|$ is shown in Figures 9.2 and 9.3 for Higgs boson signal events and the dominant non-resonant backgrounds. The transverse variables are mostly uncorrelated to $Y(W^+ + W^-)$, because a change in Y corresponds to a longitudinal Lorentz boost of the system of the hard interaction and transverse variables should not be affected by such a boost.

A small correlation can be seen for large values of $|Y(W^+ + W^-)|$, especially in the transverse mass m_T , for all processes. This is caused by a suppression of large WW system masses for high rapidities due to the steeply falling PDFs.

9.1.2 The rapidity difference of the W bosons

Apart from the $\gamma^{(*)}/Z$ process, all other backgrounds have two non-resonant W bosons. Therefore, one can expect that the rapidity difference of the two W bosons can be exploited to distinguish the resonant signal from the background processes.

The rapidity difference $|Y(W^+) - Y(W^-)|$ of the two W bosons is shown in Figure 9.1(b) for the same signal and background processes as in Figure 9.1(a). The distribution exhibits a clear difference between the resonant Higgs boson production and the non-resonant background processes.

In the same way as $Y(W^+ + W^-)$, the rapidity difference $|Y(W^+) - Y(W^-)|$ is mainly determined by the proton PDFs in the case of the non-resonant background processes. For a given rapidity Y of the WW system, the relative shapes of the gluon and sea quark PDFs are sufficiently similar (see Figure 2.1) that the increase in mass of the hard interacting system due to an increase in $|Y(W^+) - Y(W^-)|$ results in a very similar shape of the cross section as function of $|Y(W^+) - Y(W^-)|$ for all backgrounds. By coincidence, the shape of the $q\bar{q} \rightarrow WW$ and the $t\bar{t}$ background process are almost identical, while the $gg \rightarrow WW$ loop process is a bit more signal like.

For the resonant production of the WW system in a scalar Higgs boson decay the rapidity difference of the W bosons is a pure phase space function and determined only by the mass of the Higgs boson and therefore independent of the proton PDFs.

The dependency of the shapes of the two transverse variables on $|Y(W^+) - Y(W^-)|$ is shown in Figures 9.4 and 9.5. For the background processes there is again almost no correlation between $\Delta\phi_U$ and m_T on the one hand and $|Y(W^+) - Y(W^-)|$ on the other hand. The spin correlations of the W bosons are mediated by a t-channel quark exchange. The transverse component of this spin correlation, which is important for the transverse variables, is independent of the longitudi-

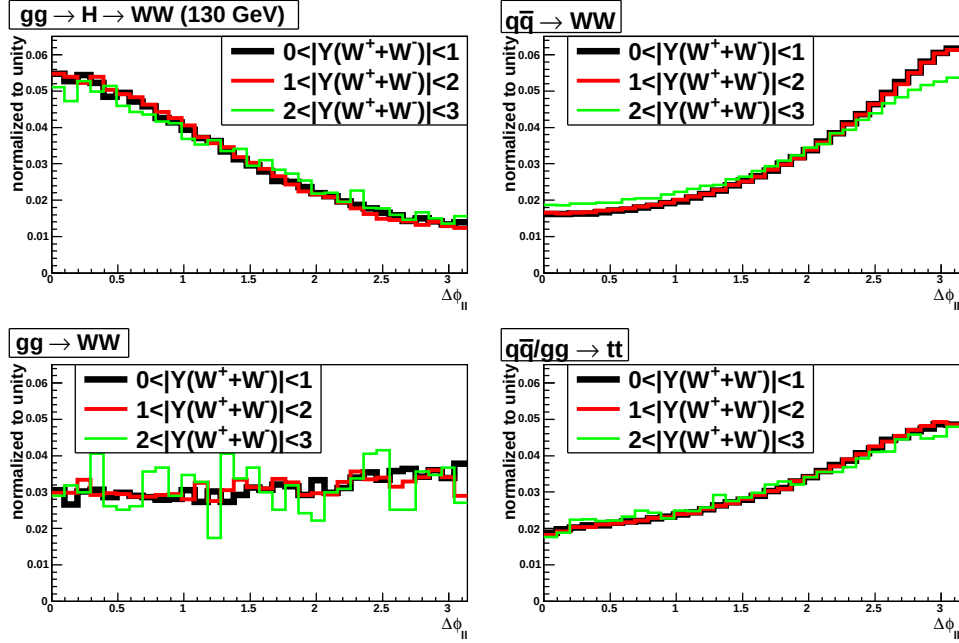


Figure 9.2: Distribution of the transverse opening angle $\Delta\phi_{ll}$ in three bins of the longitudinal variable $|Y(W^+ + W^-)|$ for Higgs boson signal events ($m_H = 130$ GeV) and the dominant non-resonant backgrounds.

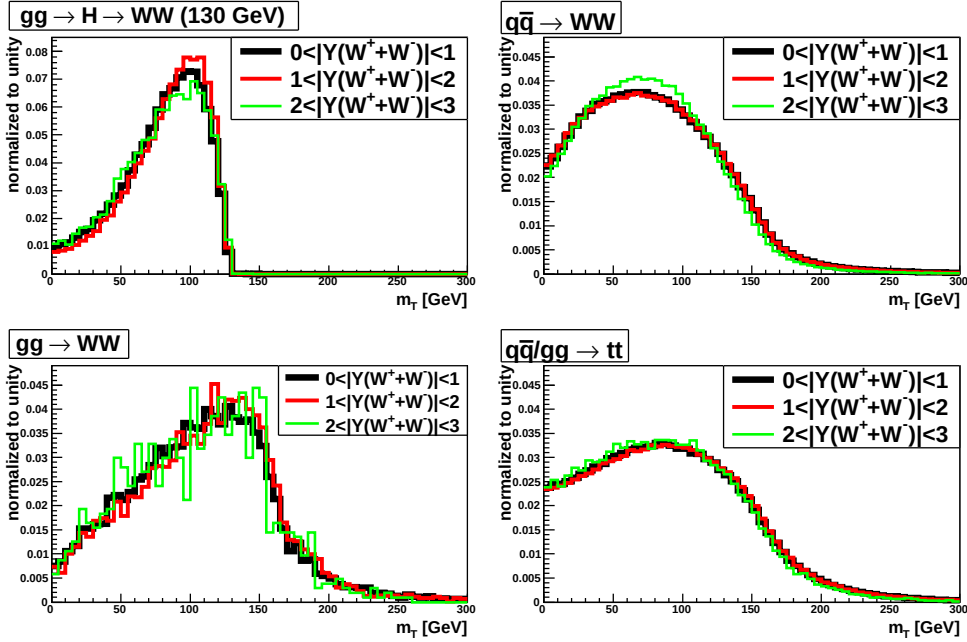


Figure 9.3: The same as Figure 9.2 for the transverse mass m_T .

nal momentum of this t-channel exchange and thus the transverse variables are (mostly) independent of $|Y(W^+) - Y(W^-)|$.

Small observable correlations are introduced again by the PDFs for high masses of the hard interacting system and in the case of the $q\bar{q} \rightarrow WW$ background from the fact that this is not an exclusively non-resonant process. As can be seen in Figure 8.21, the $q\bar{q} \rightarrow \gamma^{(*)}/Z \rightarrow WW$ amplitude contributes to the WW production and introduces some correlations. The effect of this amplitude will be further elaborated in Section 9.4.2.

For the resonant Higgs boson process a large difference $|Y(W^+) - Y(W^-)|$ causes a small transverse momentum of the W bosons in the 4π symmetric decay and thus leads to a smaller transverse mass.

9.1.3 Observable longitudinal variables

Unfortunately, the longitudinal variables $Y(W^+ + W^-)$ and $|Y(W^+) - Y(W^-)|$ are not observable because of the unmeasurable longitudinal neutrino momentum. All observable longitudinal variables have to be derived from the lepton momenta.

Similarly as for the definition of the transverse mass m_T , several choices of defining longitudinal variables from the lepton momenta exist. Two definitions that are independent of the relative transverse orientation of the leptons are:

- The rapidity $Y(l^+ + l^-)$ of the vectorial sum of the two leptons as estimator of $Y(W^+ + W^-)$.
- The difference in pseudorapidity $\eta(l^+) - \eta(l^-)$ of the two leptons as estimator of $Y(W^+) - Y(W^-)$.

Figure 9.6 shows the correlation between $Y(W^+ + W^-)$ and $Y(l^+ + l^-)$ at the generator level for the signal and background processes. Detector acceptance cuts of $|\eta(l)| < 2.5$ are applied to the leptons. In all cases the correlation is very strong and $Y(l^+ + l^-)$ is a good substitute. Due to the limited lepton acceptance, the region $|Y(W^+ + W^-)| > 2.5$ is not accessible, but also hardly populated and should be removed anyway to avoid the small correlations mentioned in Section 9.1.1.

The correlation between $\eta(l^+) - \eta(l^-)$ and $Y(W^+) - Y(W^-)$ is shown in Figure 9.7. It is not as strong as for $Y(l^+ + l^-)$, but sufficient to use the variable for the reconstruction. The bias in the correlation is mainly due to a detector acceptance bias on $\eta(l^+) - \eta(l^-)$. For large $Y(W^+) - Y(W^-)$, both leptons are only observable within the detector acceptance, if the WW system is either central ($Y(W^+ + W^-) \approx 0$) or if the two $W \rightarrow l\nu$ decays are oriented to give a small $|\eta(l^+) - \eta(l^-)|$. Otherwise, one lepton is outside $|\eta(l)| < 2.5$ and the event is not recorded.

The distributions of $|Y(l^+ + l^-)|$ and $|\eta(l^+) - \eta(l^-)|$ after all lepton acceptance cuts are shown in Figures 9.8(a) and 9.8(b) for the Higgs boson signal process

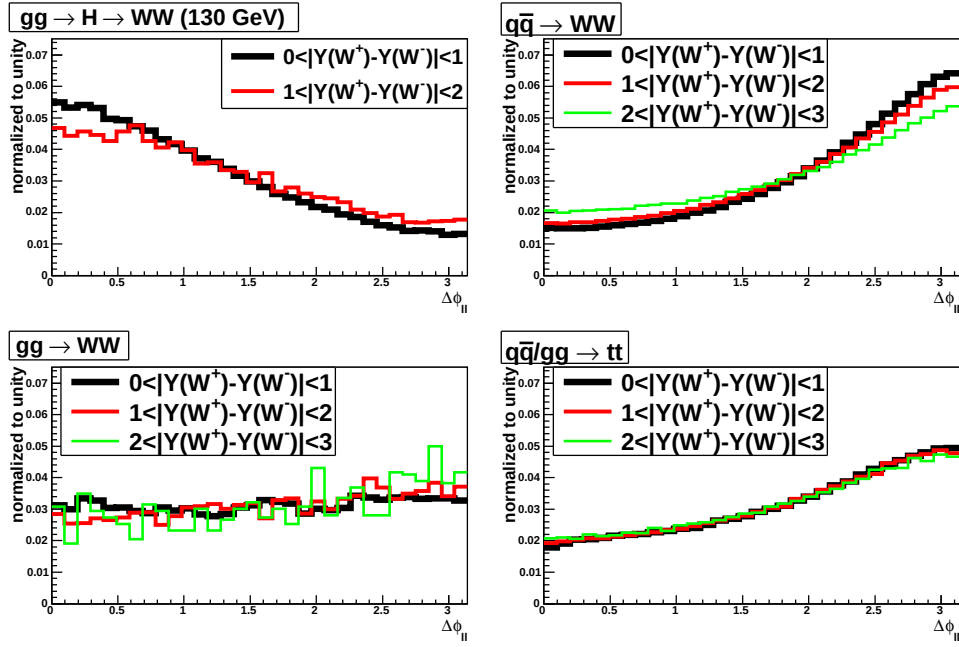


Figure 9.4: Distribution of $\Delta\phi_{ll}$ in three bins of the longitudinal variable $|Y(W^+) - Y(W^-)|$ for Higgs boson signal events ($m_H = 130$ GeV) and the dominant non-resonant backgrounds. For the Higgs boson signal process the contribution to the $|Y(W^+) - Y(W^-)| > 2$ region is negligible small and not shown.

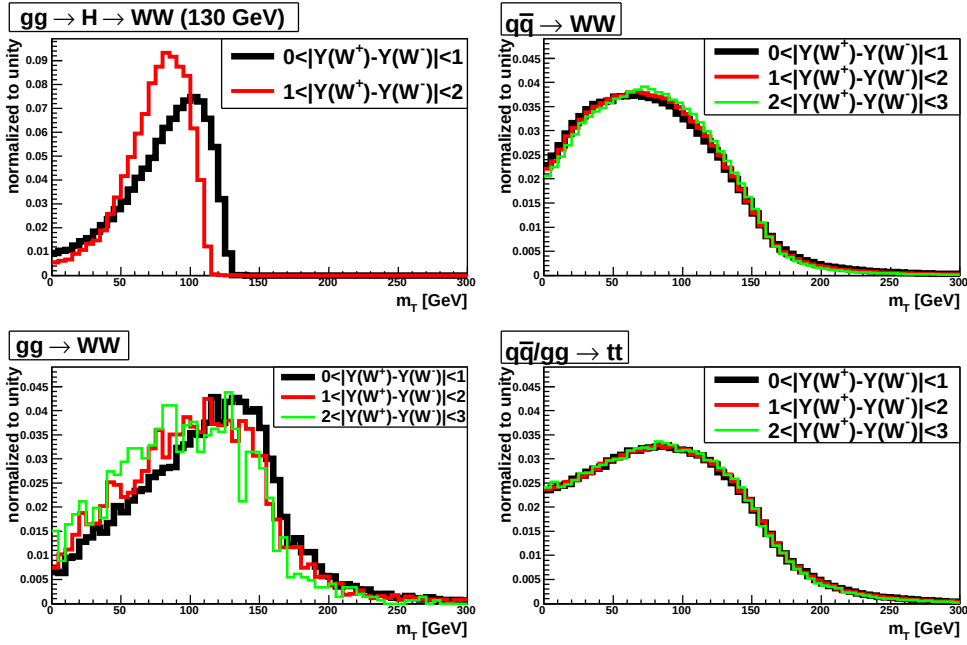


Figure 9.5: The same as Figure 9.4 for the transverse mass m_T .

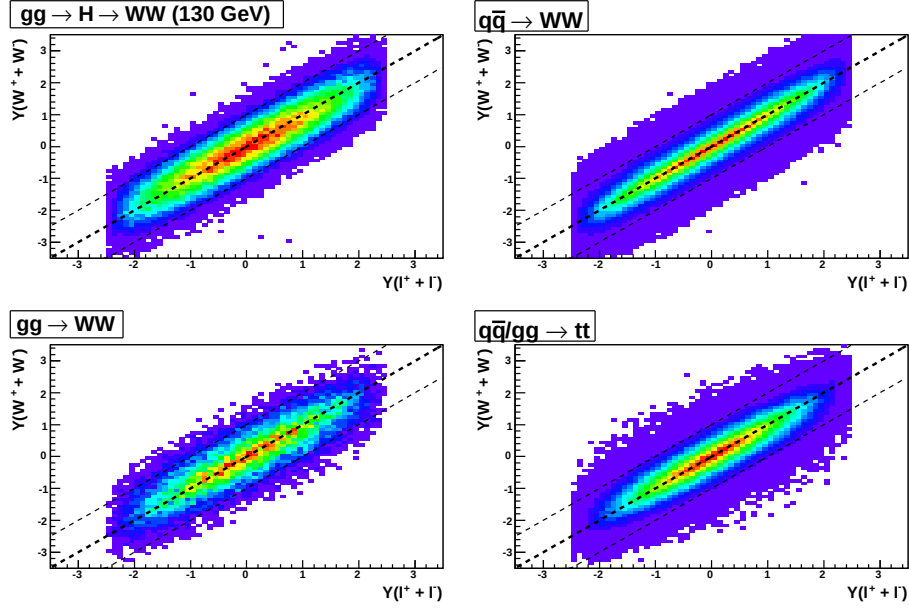


Figure 9.6: Correlation between $Y(W^+ + W^-)$ and $Y(l^+ + l^-)$ for Higgs boson signal events ($m_H = 130$ GeV) and the dominant non-resonant backgrounds. The hashed lines are intended to guide the eye.

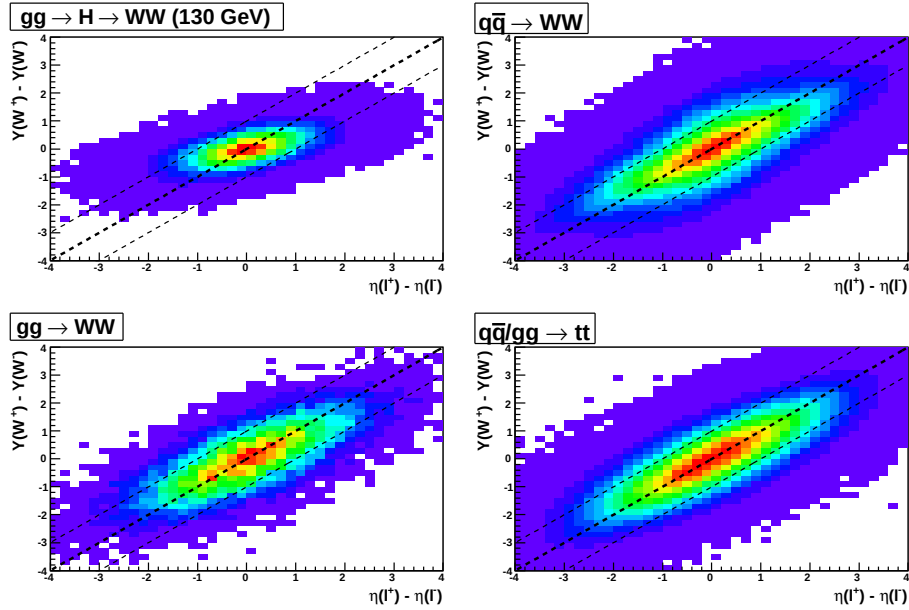


Figure 9.7: Correlation between $Y(W^+) - Y(W^-)$ and $\eta(l^+) - \eta(l^-)$ for Higgs boson signal events ($m_H = 130$ GeV) and the dominant non-resonant backgrounds. The hashed lines are intended to guide the eye.

at $m_H = 130$ GeV and for the dominant non-resonant backgrounds at generator level. The differences between the various processes that were visible for $Y(W^+ + W^-)$ and $|Y(W^+) - Y(W^-)|$ are still present, but weaker.

One reason for the increased similarity between signal and background is the bias from the lepton η acceptance. This bias can be removed to some extent by considering regions where no bias is expected. For $|Y(l^+ + l^-)|$ one can expect no bias, if the separation $|\eta(l^+) - \eta(l^-)|$ of the leptons is small. Vice versa one expects no bias on $|\eta(l^+) - \eta(l^-)|$ if the sum of both leptons is central in the detector and hence $|Y(l^+ + l^-)|$ is small.

The distribution of $|Y(l^+ + l^-)|$ for events with $|\eta(l^+) - \eta(l^-)| < 1$ and the distribution of $|\eta(l^+) - \eta(l^-)|$ for events with $|Y(l^+ + l^-)| < 1$ is shown in Figures 9.8(c) and 9.8(d). After this additional selection, the $|Y(l^+ + l^-)|$ distribution shows a clear distinction between the gluon and the quark initiated processes. Also for the $|\eta(l^+) - \eta(l^-)|$ distribution the shape difference between the signal and the background processes gets slightly enhanced.

Finally, Figures 9.8(e) and 9.8(f) show both distributions after all selection cuts against reducible backgrounds (cut 7 in Table 8.3). The cuts do not alter the distributions substantially.

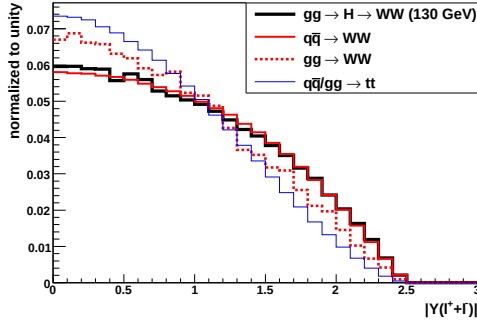
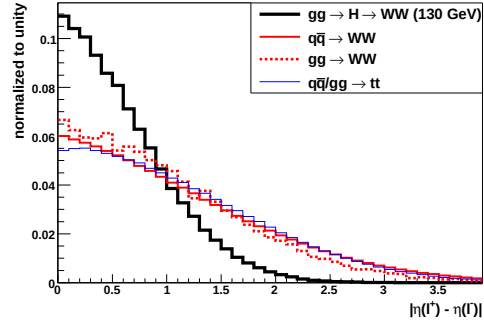
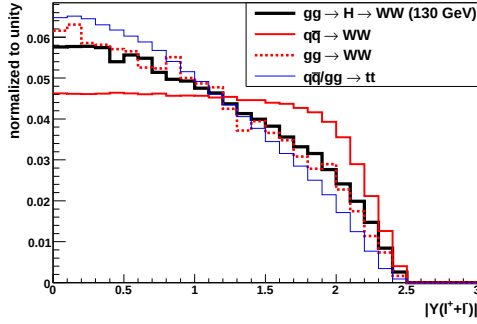
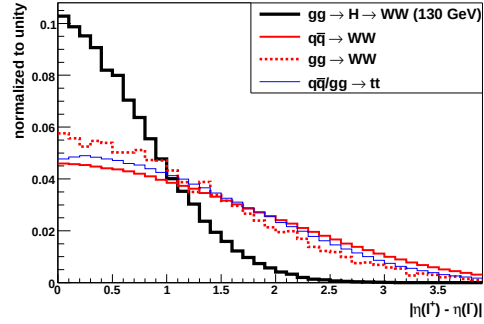
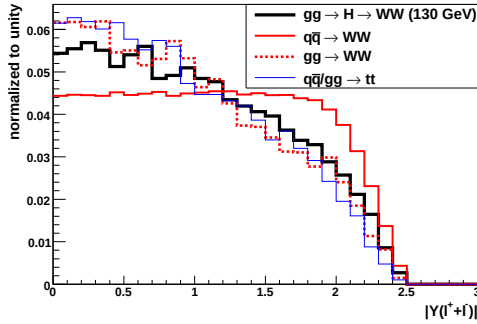
9.1.4 Correlation between the observable longitudinal and transverse variables

As a final step, before a measurement of background shapes can be considered, it needs to be shown to what extent the two observable longitudinal variables introduced in the previous section are uncorrelated to the transverse variables.

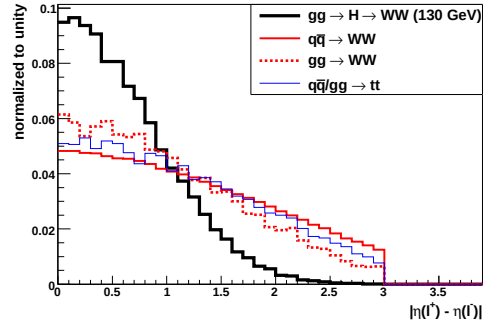
Therefore, Figures 9.9 and 9.10 show the normalised distributions of $\Delta\phi_{ll}$ and m_T for three equidistant bins in the $|Y(l^+ + l^-)|$ distribution after all selection cuts against reducible backgrounds and only for events with $|\eta(l^+) - \eta(l^-)| < 1$. Both transverse variables are completely uncorrelated to $|Y(l^+ + l^-)|$ for all processes.

The same transverse distributions are shown in Figures 9.11 and 9.12 in three equidistant bins of $|\eta(l^+) - \eta(l^-)|$ for events with $|Y(l^+ + l^-)| < 1$. As expected, the resonant signal shows a strong correlation with the transverse mass. Also the background processes have a small correlation both for $\Delta\phi_{ll}$ and m_T . While the $gg \rightarrow WW$ and the $t\bar{t}$ process become slightly more signal like with increasing $|\eta(l^+) - \eta(l^-)|$, the $q\bar{q} \rightarrow WW$ background shows the opposite behaviour.

Because of this reversed behaviour of the different background processes, it is not possible to apply a global correction to remove the remaining correlation with $|\eta(l^+) - \eta(l^-)|$. For the small $gg \rightarrow WW$ and $t\bar{t}$ backgrounds this is of no further importance, but especially for the dominant $q\bar{q} \rightarrow WW$ background the difference in the shape between the distributions for small and large $|\eta(l^+) - \eta(l^-)|$ can be up to $\approx 10\%$.


 (a) Rapidity $|Y(l^+ + l^-)|$ of the sum of the two leptons

 (b) Pseudorapidity difference $|\eta(l^+) - \eta(l^-)|$ of the two leptons

 (c) Same as (a) but for $|\eta(l^+) - \eta(l^-)| < 1$

 (d) Same as (b) but for $|Y(l^+ + l^-)| < 1$


(e) Same as (c) but after all cuts against reducible background



(f) Same as (d) but after all cuts against reducible background

 Figure 9.8: Normalised distributions of $|Y(l^+ + l^-)|$ and $|\eta(l^+) - \eta(l^-)|$ for the Higgs boson signal process at $m_H = 130$ GeV and for the dominant non-resonant backgrounds

(a) and (b): after lepton acceptance cuts.

 (c) and (d): after the additional selection of $|\eta(l^+) - \eta(l^-)| < 1$, respectively $|Y(l^+ + l^-)| < 1$, to avoid the bias from the detector acceptance.

(e) and (f): in addition after all cuts against reducible backgrounds (see Chapter 8).

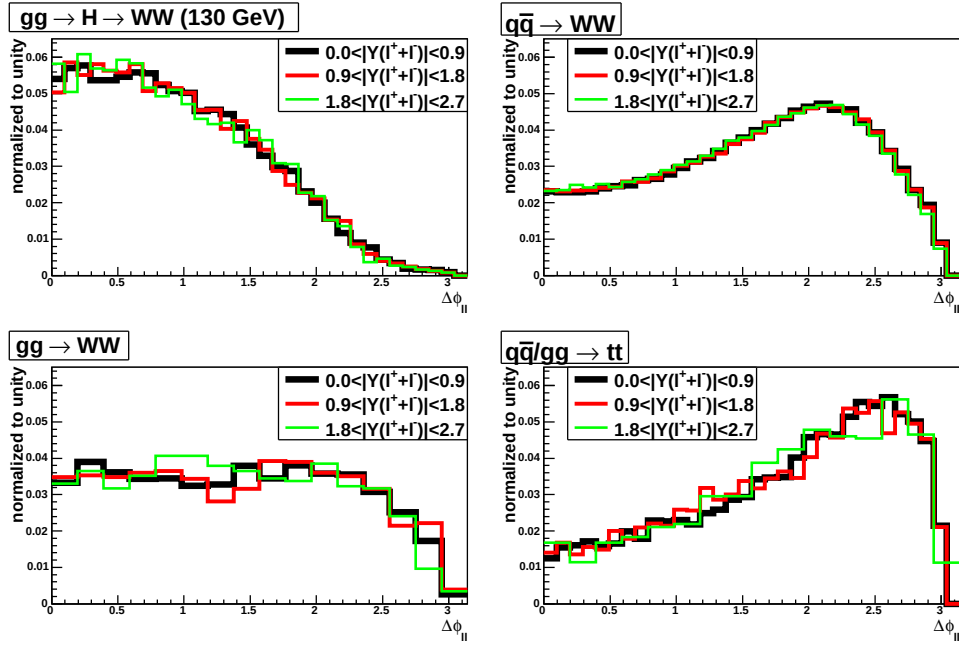


Figure 9.9: Distribution of $\Delta\phi_{ll}$ in three bins of the longitudinal variable $|Y(l^+ + l^-)|$ for Higgs boson signal events ($m_H = 130$ GeV) and the dominant non-resonant backgrounds after all cuts against reducible backgrounds (see Chapter 8) and an additional selection of events with $|\eta(l^+) - \eta(l^-)| < 1$.

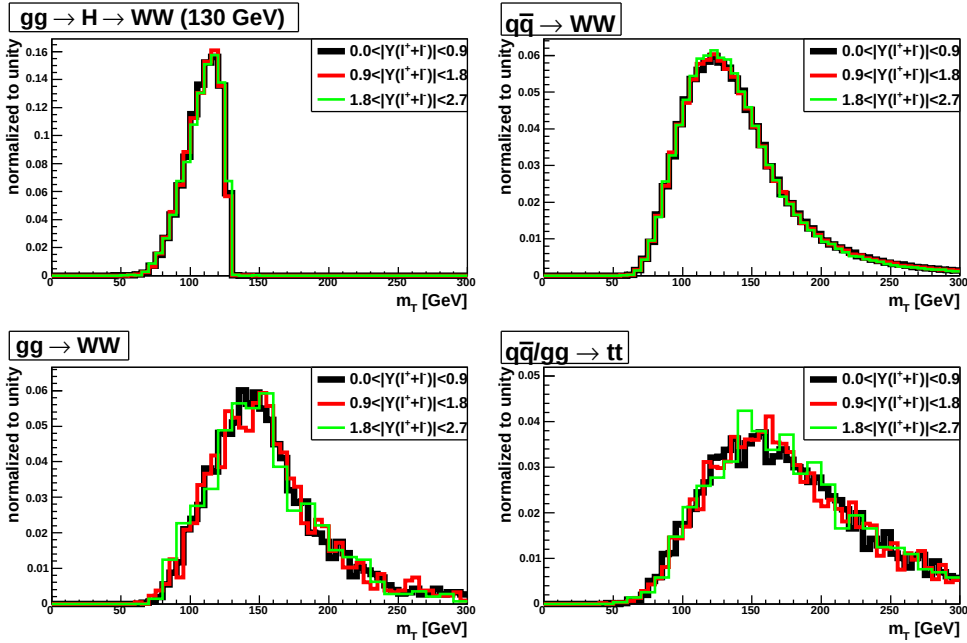


Figure 9.10: The same as Figure 9.9 but for the transverse mass m_T .

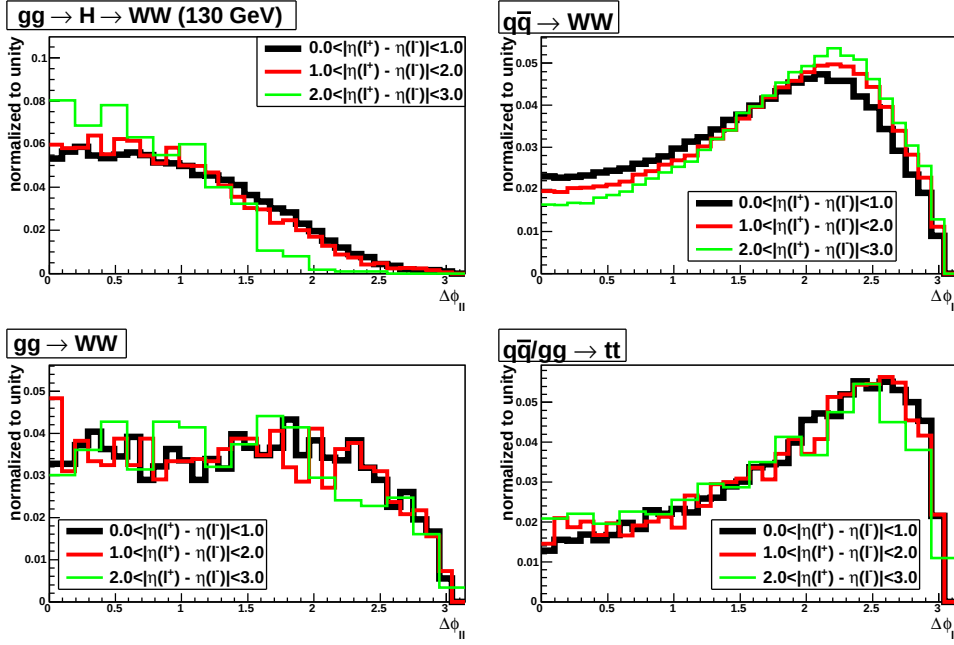


Figure 9.11: Distribution of $\Delta\phi_{ll}$ in three bins of the longitudinal variable $|\eta(l^+) - \eta(l^-)|$ for Higgs boson signal events ($m_H = 130$ GeV) and the dominant non-resonant backgrounds after all cuts against reducible backgrounds (see Chapter 8) and an additional selection of events with $|Y(l^+ + l^-)| < 1$.

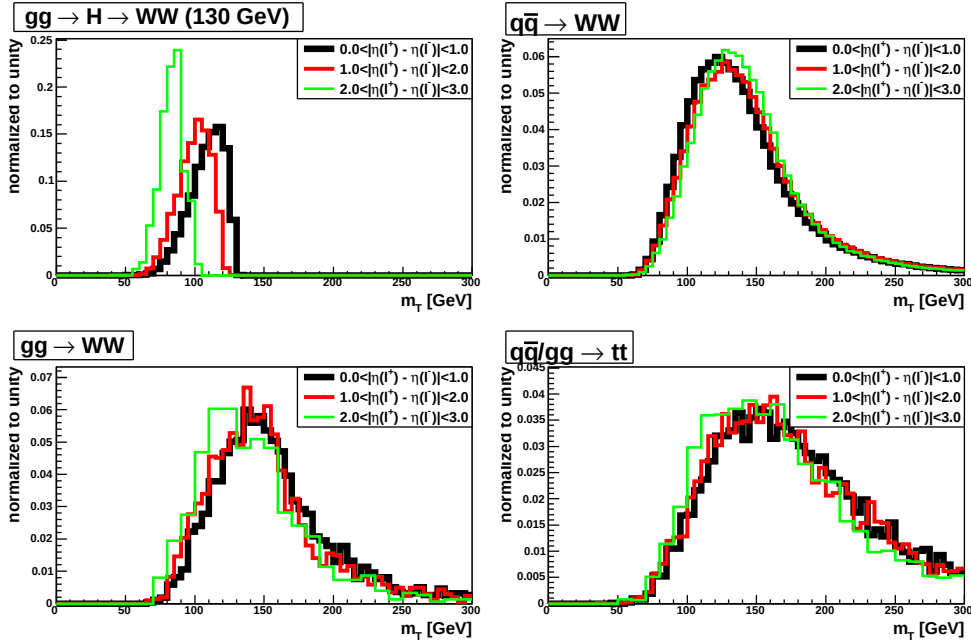


Figure 9.12: The same as Figure 9.11 but for the transverse mass m_T .

9.2 Using uncorrelated longitudinal variables to measure transverse distributions

The shape of a variable y for a class of signal processes S can be extracted from a measurement of y that includes an unknown contribution from an independent class of background processes B, if the following requirements are met:

1. A second variable x is uncorrelated to y for the background processes B.
2. The distribution of x is sufficiently different for the two process classes S and B.
3. The relative shape of the distribution of x is known, or measurable, for process class B.

9.2.1 Measuring the distribution of a signal process

If the requirements above are met, it is possible to extract a shape of y for process class S using the following steps. For convenience, the variable x is divided into two² intervals X_B and X_S . In the interval X_B a large fraction of events from class B and a small fraction of event from class S is expected. In the interval X_S only a small fraction of events from class B and hence a relatively larger fraction of events from class S is expected.

4. The total distribution of y is measured for events in the intervals X_B and X_S . The relative composition of events from class B and S is not known.
5. The shape x_B as function of x is known for class B (see requirement 3.), thus the ratio r of the integral of events in the intervals X_S and X_B can be calculated for the background B:

$$r = x_B(X_S) / x_B(X_B)$$

$$\text{with } x_B(X_B) = \int_{X_B} x_B(x) dx \text{ and } x_B(X_S) = \int_{X_S} x_B(x) dx.$$

6. The whole measured distribution $y(X_B)$ is scaled by the ratio r .
7. This scaled distribution is subtracted from the measured distribution $y(X_S)$:

$$y_S = y(X_S) - r \cdot y(X_B) \quad .$$

In the difference y_S all contributions from class B should have been subtracted and the remaining events are exclusively signal events from class S. Because x and y were not required to be uncorrelated for the signal S and the relative size of the signal components in X_S and X_B are not known, the final distribution y_S is biased and represents neither the y shape of S in X_S nor in X_B .

²In principle the method can be extended to an arbitrary number of intervals.

9.2.2 Measuring the unbiased distribution for two process classes

If in addition the variables x and y are uncorrelated also for the signal process class S, y_S is an unbiased distribution of y for the signal processes. In this case a global fit can be used to extract also the unbiased distributions y_B of the class B and the relative normalisation of S and B in the intervals X_S and X_B .

In the following, $N(X_S)$ and $N(X_B)$ are the numbers of the total observed events in the intervals X_S , respectively X_B . $N_S(X_{S/B})$ and $N_B(X_{S/B})$ are the number of events from process class S, respectively B, in the intervals $X_{S/B}$. The distributions $\hat{y}_S$, $\hat{y}_B$, $\hat{y}(X_S)$ and $\hat{y}(X_B)$ are defined as normalised distributions with integral 1, e.g. $\hat{y}_S = y_S / \int y_S(y) dy$. For the global fit, all distributions $\hat{y}$ are represented by histograms with n bins.

8. The unknown normalisation variables are $N_B(X_B)$, $N_S(X_B)$ and $N_S(X_S)$. The normalisation factor $N_B(X_S)$ can be calculated from r and $N_B(X_B)$:

$$N_B(X_S) = r \cdot N_B(X_B) \quad .$$
9. The unknown distribution $\hat{y}_B$ and the normalisation variables can then be obtained from a fit of the two conditions:

$$\begin{aligned} N(X_B) \cdot \hat{y}(X_B) &= N_B(X_B) \cdot \hat{y}_B + N_S(X_B) \cdot \hat{y}_S \\ N(X_S) \cdot \hat{y}(X_S) &= N_B(X_S) \cdot \hat{y}_B + N_S(X_S) \cdot \hat{y}_S \quad . \end{aligned}$$

10. For n bins in the normalised $\hat{y}$ distributions $2n$ equations are obtained. The unknown distribution $\hat{y}_B$ has $n - 1$ degrees of freedom, resulting in a total number of $3 + (n - 1)$ free variables in the fit. Hence the $\hat{y}$ distributions need to be represented by at least two bins in the histograms for the fit to converge.

9.2.3 Measurement of all distributions and normalisation factors

For the fit procedure described above the ratio $r = N_B(X_S) / N_B(X_B)$ is still needed as external input. When two uncorrelated distributions y are available (e.g. $y^I = \Delta\phi_{ll}$ and $y^{II} = m_T$), one could assume that this external information of r is no longer needed. Two normalised distributions $\hat{y}$ with n bins each give $4n$ equations:

$$\begin{aligned} N(X_B) \cdot \hat{y}^I(X_B) &= N_B(X_B) \cdot \hat{y}_B^I + N_S(X_B) \cdot \hat{y}_S^I \\ N(X_S) \cdot \hat{y}^I(X_S) &= N_B(X_S) \cdot \hat{y}_B^I + N_S(X_S) \cdot \hat{y}_S^I \\ N(X_B) \cdot \hat{y}^{II}(X_B) &= N_B(X_B) \cdot \hat{y}_B^{II} + N_S(X_B) \cdot \hat{y}_S^{II} \\ N(X_S) \cdot \hat{y}^{II}(X_S) &= N_B(X_S) \cdot \hat{y}_B^{II} + N_S(X_S) \cdot \hat{y}_S^{II} \end{aligned}$$

Without r there are four normalisation variables and altogether $4(n - 1)$ bins in the histograms to be determined, resulting in a total of $4n$ free variables. Hence, the system should be solvable.

However, when adding up all equations that belong to one distribution y , it turns out that not all equations are independent and the following system of equations is obtained:

$$\begin{aligned} N(X_B) &= N_B(X_B) + N_S(X_B) \\ N(X_S) &= N_B(X_S) + N_S(X_S) \\ N(X_B) &= N_B(X_B) + N_S(X_B) \\ N(X_S) &= N_B(X_S) + N_S(X_S) \quad . \end{aligned}$$

Two equations each are identical, because y^I and y^{II} are distributions of different variables for the same events and thus have the same integral. Hence the system unfortunately cannot be solved, and additional external information, either from independent measurements or Monte Carlo simulations, about r or an equivalent variable is needed.

9.3 Measuring the $q\bar{q} \rightarrow WW$ background

The first case considered is the measurement of the shape of the $q\bar{q} \rightarrow WW$ background using the rapidity $|Y(l^+ + l^-)|$ of the sum of the two leptons. To reduce the bias from detector acceptance, only events with a small pseudorapidity difference $|\eta(l^+) - \eta(l^-)| < 1$ of the leptons are selected.

Using the definitions from the previous section, the process class B are all gluon induced processes ($gg \rightarrow H \rightarrow WW$, $gg \rightarrow WW$ and $t\bar{t}$) and the process class S is the $q\bar{q} \rightarrow WW$ process. The uncorrelated variable is $x = |Y(l^+ + l^-)|$. The distribution y can either be $\Delta\phi_{ll}$ or m_T . For the cut based analysis shown in Chapter 8, $y = \Delta\phi_{ll}$ is more relevant and hence considered here. The intervals X_B and X_S are chosen as $X_B : |Y(l^+ + l^-)| < 1$ and $X_S : |Y(l^+ + l^-)| > 1$ (see Figure 9.8(e)).

In addition the ratio r for the gluon initiated processes B is needed in order to follow the procedure introduced in Section 9.2.2. Unfortunately, this ratio cannot easily be measured, because neither in $\Delta\phi_{ll}$ nor in m_T there is a region where no events from the $q\bar{q} \rightarrow WW$ process are expected to allow an isolated measurement of the gluon processes.

Hence, the ratio r can only be determined from a Monte Carlo simulation. Figure 9.13 shows the shape of the $\Delta\phi_{ll}$ and m_T distribution for the $q\bar{q} \rightarrow WW$ process and the gluon induced backgrounds (including the Higgs boson signal) after following the steps 1 – 7 of Section 9.2 with $r = 1.258 \pm 0.012$ determined from the Monte Carlo sum of all gluon processes.

Neither the scaled shapes of the $\Delta\phi_{ll}$ nor of the m_T distribution agree very well between the regions $|Y(l^+ + l^-)| < 1$ and $|Y(l^+ + l^-)| > 1$ for the gluon

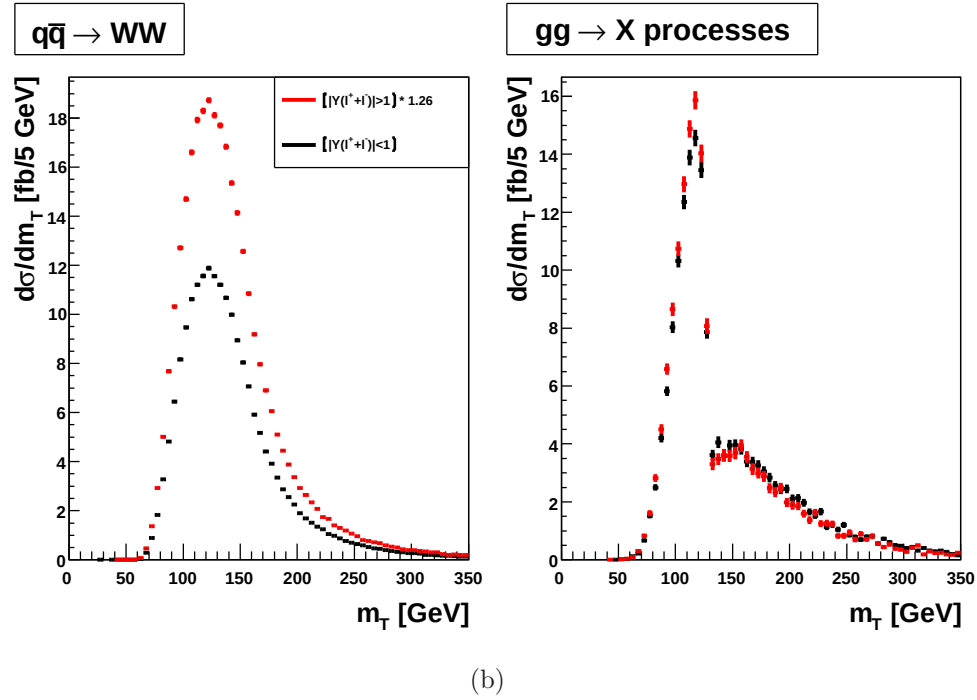
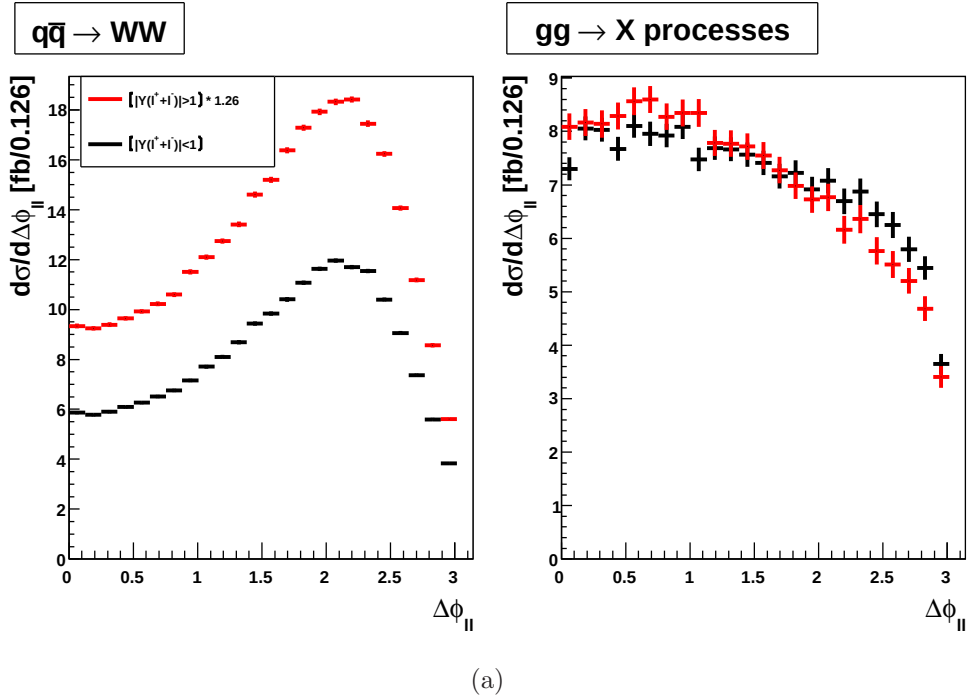


Figure 9.13: Comparison of the $\Delta\phi_{||}$ (a) and m_T (b) shape for the $q\bar{q} \rightarrow WW$ process and the gluon induced background contribution after the normalisation procedure (steps 1 – 7 of Section 9.2). The peak at $m_T \approx 130$ GeV in (b) is from the Higgs boson process. All Distributions for $|Y(l^+ + l^-)| > 1$ are scaled by the factor $r = 1.258 \pm 0.012$ determined from the Monte Carlo sum of all gluon processes.

induced processes. While there is no correlation between $|Y(l^+ + l^-)|$ and the transverse variables for each individual gluon induced process (see Figures 9.9 and 9.10), this is not the case for the sum of all gluon induced processes if their relative contributions in the two regions $|Y(l^+ + l^-)| < 1$ and $|Y(l^+ + l^-)| > 1$ are not identical. As can be seen in Figure 9.8(e), the shape of the $|Y(l^+ + l^-)|$ distribution is different for the gluon induced processes and thus the necessary scaling factor r is slightly different for each individual gg process. The scaling factor r used in Figure 9.13 gives the correct normalisation for the integral of the sum of all processes and thus has a small bias for the individual processes, especially for the relatively large $gg \rightarrow H \rightarrow WW$ contribution.

The final shape of $\Delta\phi_{ll}$ and m_T extracted from the subtraction of the measured shapes in $|Y(l^+ + l^-)| < 1$ and $|Y(l^+ + l^-)| > 1$ (step 7 of Section 9.2) is shown in Figure 9.14. The effect of the disagreement of the $|Y(l^+ + l^-)|$ distributions for the gluon induced processes can be seen in a residual contribution of the $gg \rightarrow H \rightarrow WW$ process to the measured $q\bar{q} \rightarrow WW$ shapes.

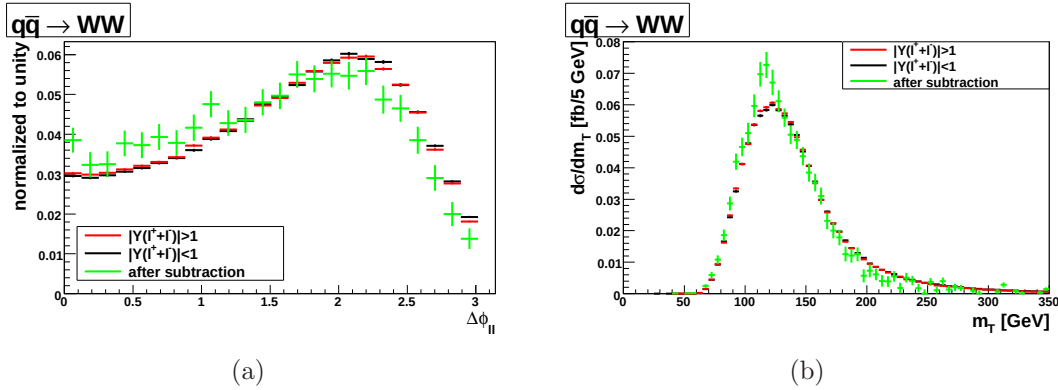


Figure 9.14: Comparison of the $\Delta\phi_{ll}$ (a) and m_T (b) shape for the $q\bar{q} \rightarrow WW$ process for the regions $|Y(l^+ + l^-)| < 1$ and $|Y(l^+ + l^-)| > 1$ as well as the measured shape after subtraction of the gluon induced processes, including the Higgs boson signal ($m_T = 130$ GeV).

If no Higgs boson exists, this disagreement is largely removed and the $q\bar{q} \rightarrow WW$ shape can be extracted unbiased, as it is shown in Figure 9.15.

At this stage the steps 8 – 10 of Section 9.2 could be followed to extract the shape of the gluon induced processes and the normalisation factors. However, the generator level results with a perfect Monte Carlo knowledge of the factor r show already sufficient disagreement to question such an approach, if a Higgs boson signal is present.

The differences of the $|Y(l^+ + l^-)|$ distributions for the gluon induced processes is mainly caused by the different masses of the hard interacting systems and thus by the proton PDFs. It might be possible to use the per event measurement of

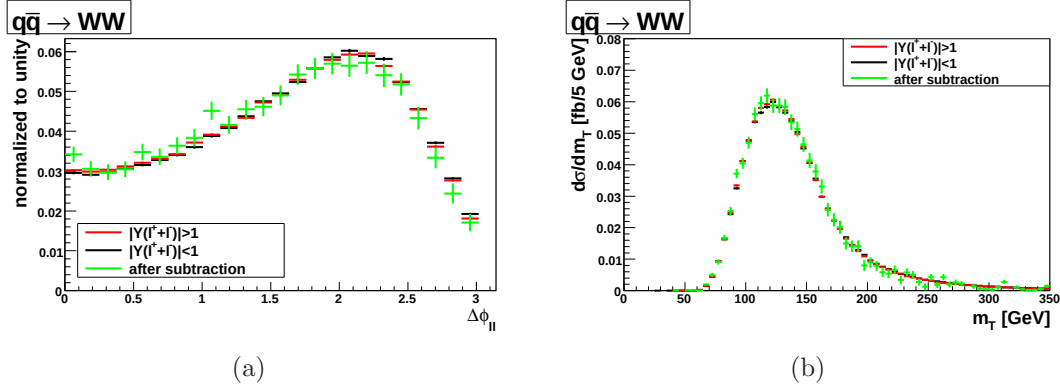


Figure 9.15: The same as Figure 9.14, but without the Higgs boson signal contribution to the gluon induced processes.

m_T and $|\eta(l^+) - \eta(l^-)|$ to obtain an estimate of this mass in order to apply a Monte Carlo correction to the $|Y(l^+ + l^-)|$ distribution.

Furthermore, part of the difference for the $t\bar{t}$ process is caused by the $\approx 10\%$ contribution of the $q\bar{q}$ initial state to the cross section. On the other hand, especially the $t\bar{t}$ process would allow a measurement of the $|Y(l^+ + l^-)|$ distribution for a gluon induced process, if the jet veto is relaxed and possibly an additional b -jet tag requirement applied.

Despite the challenges mentioned above, it is reckoned that an iterative approach to the extraction of the shape and normalisation of the $q\bar{q} \rightarrow WW$ process on the one hand and the gluon induced processes on the other hand, involving Monte Carlo corrections for known differences in the intermediate steps, is possible. However, this can neither be further elaborated in the context of this thesis, nor be tested for feasibility within the required uncertainty limits.

Finally, the measurement of the $q\bar{q} \rightarrow WW$ cross section and the shape of the transverse variables after the event selection of Chapter 8 can be used to estimate the dominant background source to the Higgs boson production, but this still is not sufficient to discover the Higgs boson. Especially the problem of the $gg \rightarrow WW$ background is not addressed at all.

9.4 Measuring the backgrounds to the $gg \rightarrow H \rightarrow WW$ signal process

Instead of pursuing in more detail the measurement of the $q\bar{q} \rightarrow WW$ background, a second background normalisation approach involving the longitudinal $|\eta(l^+) - \eta(l^-)|$ distribution is followed, because of the promising large difference in this variable between the resonant signal process and all backgrounds.

Using the definitions of Section 9.2, the process class S is the $gg \rightarrow H \rightarrow WW$ Higgs boson signal process. The process class B consists of the $q\bar{q} \rightarrow WW$, the $gg \rightarrow WW$ and the $t\bar{t}$ backgrounds. The uncorrelated variable is $x = |\eta(l^+) - \eta(l^-)|$ after an additional selection cut on $|Y(l^+ + l^-)| < 1$. As signal and background enhanced regions the intervals $X_S : |\eta(l^+) - \eta(l^-)| < 1$ and $X_B : |\eta(l^+) - \eta(l^-)| > 1$ are chosen.

The longitudinal variable $|\eta(l^+) - \eta(l^-)|$ does not show an as complete lack of correlations to the transverse variables as it is the case for $|Y(l^+ + l^-)|$. Spin correlations in the leptonic decay of the two W bosons link the longitudinal separation of the leptons to their relative transverse orientation. While this effect is sufficiently small for the $gg \rightarrow WW$ and $t\bar{t}$ backgrounds, it is sizeable for $q\bar{q} \rightarrow WW$ and especially visible in the $\Delta\phi_{ll}$ distribution (see Figure 9.11).

This correlation also appears in the dependency of the background extrapolation ratio r of the number of events between the regions X_S and X_B on the Higgs boson selection cuts. In Figure 9.16 the Monte Carlo prediction for r is shown for the individual processes and the combined background for each selection cut introduced in Chapter 8. While the ratio is basically unchanged by the cuts suppressing the reducible backgrounds, the lepton angular cuts have a strong effect on r .

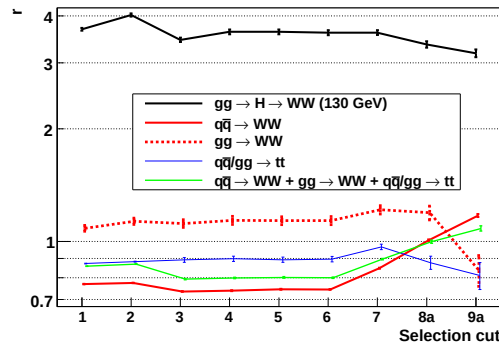


Figure 9.16: Dependency of the background extrapolation ratio r between the number of events in X_S and X_B on the Higgs boson selection cuts introduced in Chapter 8. A Higgs boson mass of 130 GeV is assumed. The error bars show the Monte Carlo statistical error.

One way of evaluating this dependency is to use extensive Monte Carlo simulations. However, the problem of determining r is mitigated, if the ratio can be directly measured. To do this, r needs to be measured after all event selection cuts and hence from the m_T distribution itself after the selection cut on $\Delta\phi_{ll} < 1$.

As an illustration of the shape and absolute size of the signal and background contributions in X_S and X_B , a comparison of the distributions of $\Delta\phi_{ll}$ and m_T is shown in Figures 9.17 and 9.18. The $\Delta\phi_{ll}$ distributions are shown after all cuts against reducible backgrounds, the m_T distributions after the additional cut on

$\Delta\phi_u < 1$. Following the normalisation procedure in steps 1 – 7 of Section 9.2, a Monte Carlo background extrapolation ratio of $r = 1.00 \pm 0.01$, as it is found after the $\Delta\phi_u < 1$ selection (see cut 8a in Figure 9.16), is used for the X_B region in both sets of distributions.

Several points should be noted:

- The value of $r = 1.00$ is pure coincidence and depends on the chosen cuts. While a similar size of the background in the two regions X_S and X_B is advantageous in order to optimize the final statistical error, it is not mandatory or needed for the method.
- For the background only case the subtraction results in an almost complete compensation of the X_S distribution by the X_B distribution. The remaining structure is caused by the relative shift in the transverse mass m_T between the distributions in X_S and X_B . This shift is a bias due to the $|\eta(l^+) - \eta(l^-)|$ selection and not consistent for the different backgrounds and therefore cannot be corrected globally. It could be corrected with Monte Carlo information, if the size of the $q\bar{q} \rightarrow WW$ process were to be known from an independent measurement, using for example the method from Section 9.3.
- For the signal plus background case the optimal value of r should be extracted from all background events in a narrow transverse mass interval enclosing the expected Higgs boson signal peak (e.g. as given by cut 9a/b).
- For the signal plus background case the small signal contribution in X_B is subtracted from the large signal peak in X_S . Due to the correlation of the resonant signal process with $|\eta(l^+) - \eta(l^-)|$, this subtraction gives a biased m_T distribution. The bias is such, that after the subtraction only signal events in the range $m_H - 40 \text{ GeV} \lesssim m_T < m_H$ remain, whatever the width of the m_T distribution for the signal was before the subtraction.

9.4.1 Measuring the background extrapolation ratio

After the subtraction of the measured X_B distribution, the Higgs boson signal shows a narrow peak with a sharp edge below the true Higgs boson mass. Outside this narrow peak, there is almost no contribution from the signal. Hence, the region outside the peak can be used as a sideband to measure the total background size in X_S and X_B and to derive r from this measurement.

The largest part of the Higgs boson signal events can be found in the interval $M_T^{\text{signal}} : m_H - 40 \text{ GeV} < m_T < m_H$ of the distribution of the transverse mass m_T . Hence, the background extrapolation ratio r can be measured in the regions $m_T > m_H$ and $m_T < m_H - 40 \text{ GeV}$. Most of the events from the dominant dilepton backgrounds are located in the interval $70 \text{ GeV} < m_T < 200 \text{ GeV}$. In

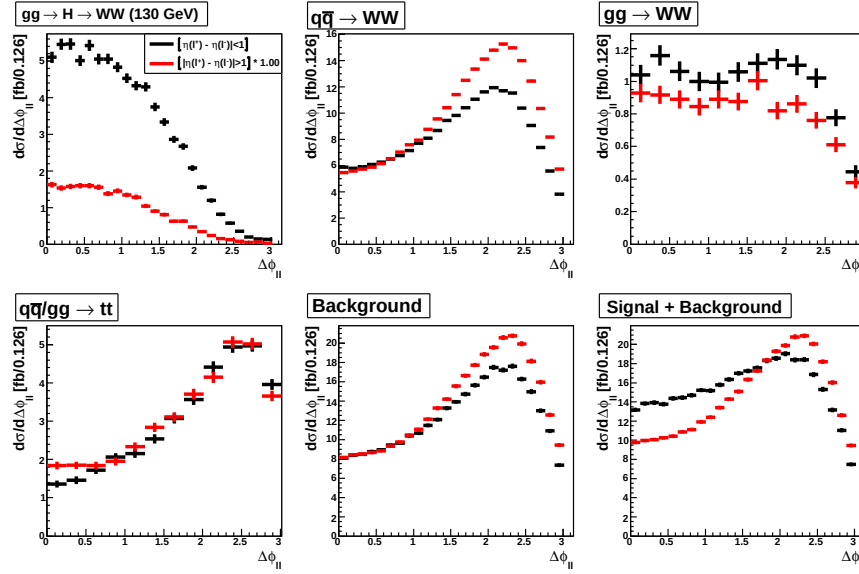


Figure 9.17: Comparison of the $\Delta\phi_{ll}$ shape for the $gg \rightarrow H \rightarrow WW$ signal process ($m_H = 130$ GeV) and the dominant dilepton backgrounds for the regions $|\eta(l^+) - \eta(l^-)| < 1$ and $|\eta(l^+) - \eta(l^-)| > 1$ after all cuts against reducible backgrounds. For the relative normalisation of the $|\eta(l^+) - \eta(l^-)| < 1$ and $|\eta(l^+) - \eta(l^-)| > 1$ distributions see the text.

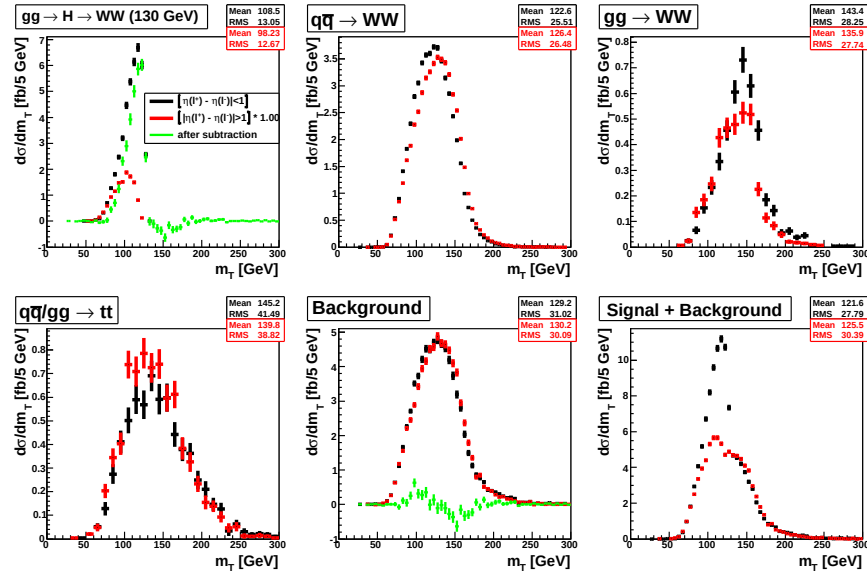


Figure 9.18: The same as Figure 9.17 for the m_T shape after the additional selection cut on $\Delta\phi_{ll} < 1$. The subtraction of the m_T distribution for the sum of the signal and background is drawn in green on top of the Higgs boson signal. The subtraction of the background only case is drawn on top of the summed background distribution.

order to reduce the influence of other background sources, the measurement of r is further restricted to this interval. Thus, the interval in which r is measured is $M_T^{\text{norm.}} : 70 \text{ GeV} < m_T < m_H - 40 \text{ GeV} \cup m_H < m_T < 200 \text{ GeV}$.

The extraction of the Higgs boson signal is done in the following steps:

- $I_{X_S}^N$ is the integral of the measured events in X_S in the interval $M_T^{\text{norm.}}$.
 $I_{X_B}^N$ is the same integral in X_B in the the interval $M_T^{\text{norm.}}$.
- The measured background extrapolation ratio is $r = I_{X_S}^N / I_{X_B}^N$.
- The total measured m_T distribution in X_B is scaled by r and subtracted from the measured m_T distribution in X_S .
- The remaining events in the interval M_T^{signal} are considered as Higgs boson candidate events.

If no Higgs boson signal is present, the events in X_S and X_B should cancel each other after scaling with r and no excess of events should remain. Therefore the observable ratio $r^{\text{obs.}}$ in the interval M_T^{signal} after subtraction should be identical to the true value r^{true} in the same interval.

Figure 9.19 shows the variation of r , r^{true} and $r^{\text{obs.}}$ as function of the Higgs boson mass assumption only for the Standard Model backgrounds. Due to the correlations of the transverse variables with $|\eta(l^+) - \eta(l^-)|$, the cancellation is not perfect. However, the difference between r^{true} and $r^{\text{obs.}}$ is less than 5% in the whole mass range. Thus, any systematic uncertainty on the transverse distributions of the backgrounds that influence X_S and X_B in the same way will be suppressed by a factor of at least 20. It should be noted especially that also the $gg \rightarrow WW$ background cancels to a large degree. Thus a precise knowledge of the cross section is no longer needed.

If a Higgs boson signal is present, the measured background extrapolation ratio r is affected as well. A comparison of the resulting distributions of m_T is shown in Figure 9.20 for a Higgs boson signal of $m_H = 130 \text{ GeV}$ and $m_H = 170 \text{ GeV}$, together with the background only distributions in X_S before and after the background subtraction.

The Higgs boson signal shows up as a clear mass peak after subtraction, while the background only case results in an almost zero distribution. It should be noted that for a Higgs boson mass of $m_H = 170 \text{ GeV}$ the m_T mass peak in X_S and X_B is considerably broader than the 40 GeV signal interval. However, the low mass tail of the signal has almost the same ratio r between X_S and X_B as the background and hence this tail does not bias the measurement of r and is removed by the subtraction. The obvious drawback is that a considerable part of the signal events are lost.

The corresponding comparison of the m_T distributions for Higgs boson masses in the full range from 120–190 GeV is shown in Figures D.1 - D.4. With increasing

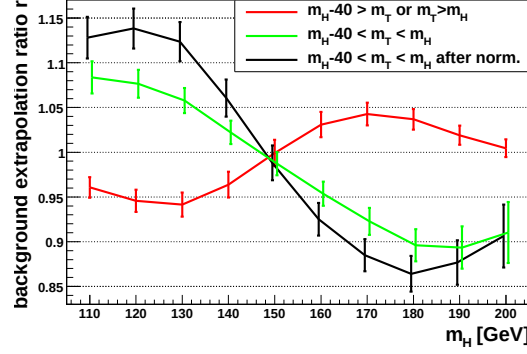


Figure 9.19: Measured background extrapolation ratio r (red) as function of the Higgs boson mass used in the analysis for a background only experiment without a Higgs boson signal contribution. The true background extrapolation ratio r^{true} in the interval M_T^{signal} is shown in green. The observable ratio $r^{\text{obs.}}$ in the interval M_T^{signal} after rescaling of the X_B distribution with r is shown in black. The error bars give the Monte Carlo statistical uncertainty.

$$r = r(m_T < m_H - 40 \text{ GeV or } m_T > m_H)$$

$$r^{\text{true}} = r(m_H - 40 \text{ GeV} < m_T < m_H)$$

$$r^{\text{obs.}} = r(m_H - 40 \text{ GeV} < m_T < m_H)/r$$

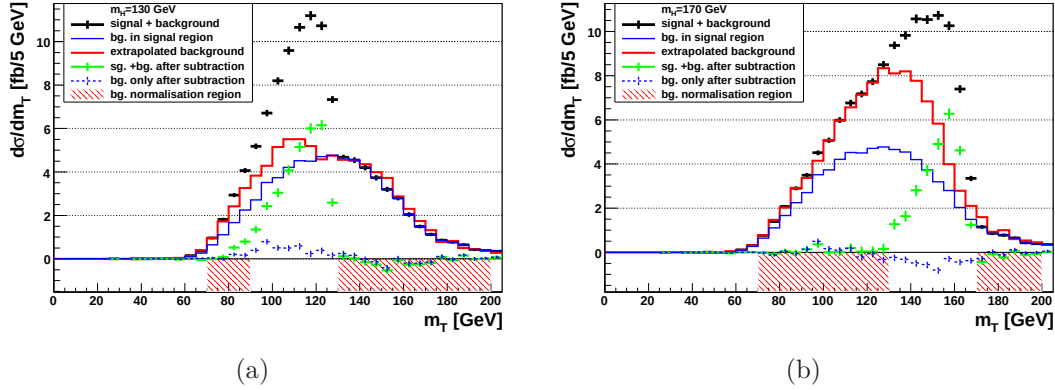


Figure 9.20: Comparison of the true, measured and subtracted m_T distributions for $m_H = 130$ GeV (a) and $m_H = 170$ GeV (b). Black: the measured distribution of m_T in X_S . Red: the measured distribution in X_B scaled with the measured r . Green: the m_T distribution after the subtraction of the first two. Solid blue: the true background distribution in X_S . Dashed blue: The subtracted m_T distribution if no signal is present. The hashed area indicates the m_T interval $M_T^{\text{norm.}}$, in which r is measured. The error bars show the Monte Carlo statistical uncertainty.

Higgs boson mass, the differences between the measured distributions in X_S and X_B get smaller and hence the remaining signal decreases rapidly. The main reason is that above the kinematic threshold for the production of two real W bosons ($m_H > 2m_W$) the $H \rightarrow WW$ decay can produce two W bosons with a large rapidity difference. This decreases the visible difference in the $|\eta(l^+) - \eta(l^-)|$ distribution between the signal process and the backgrounds considerably and thus causes one of the necessary conditions of the background normalisation method not to be fulfilled any longer.

9.4.2 Impact of the $q\bar{q} \rightarrow Z^{(*)} \rightarrow WW$ amplitude on the $q\bar{q} \rightarrow WW$ background

The largest bias in the background subtraction is caused by the correlation between $\Delta\phi_{ll}$ and m_T on the one hand and $|\eta(l^+) - \eta(l^-)|$ on the other hand. In order to understand this correlation, the dominant $q\bar{q} \rightarrow WW$ production amplitudes are analysed individually. The two leading order Feynman diagrams for the $q\bar{q} \rightarrow WW$ process, the non-resonant t-channel production of two W bosons radiated from a quark line and the resonant s-channel production of $\gamma^{(*)}/Z$ with a subsequent decay into two W bosons, are shown in Figure 8.21.

Both diagrams are divergent at large energies if considered individually. However, both diagrams converge if the production is regularized by requiring $p_T(W) < 250 \text{ GeV}^3$ and the selection cut on $\Delta\phi_{ll} < 1$ is applied. Using a MadGraph [41] matrix element and the MadEvent [41] event generator, it is possible to consistently produce $q\bar{q} \rightarrow WW \rightarrow l\nu l\nu$ events for the full $q\bar{q} \rightarrow WW$ process and the s- and t-channel components individually.

Figure 9.21 shows that after these cuts the distributions of the W bosons for the components agree sufficiently well with the full process to assume that the remaining differences have a physical meaning and are not a residual effect of the divergence. MadGraph/MadEvent uses a leading order matrix element without initial state radiation, therefore the WW system has no transverse momentum.

The accepted cross section after the regularization cuts is 55 fb for the full WW process, 43 fb for the s-channel and 108 fb for the t-channel. Hence, the t-channel amplitude is the bigger component and the s-channel the smaller one. There is a non-trivial relative phase between the two components to produce the full result by interference. It must be assumed that the relative size of the two components and the relative phase varies depending on the phase space point.

Figure 9.22 shows the correlation of the transverse mass m_T with $|\eta(l^+) - \eta(l^-)|$ for the full WW process as well as for the s- and t-channel components and their scalar sum. The comparison indicates that there is no correlation present between m_T and $|\eta(l^+) - \eta(l^-)|$ for the non-resonant t-channel and the visible correlation for the full WW process is caused by the s-channel component. Unfortunately

³The contribution of the full $q\bar{q} \rightarrow WW$ process to the $p_T(W) > 250 \text{ GeV}$ region is negligible.

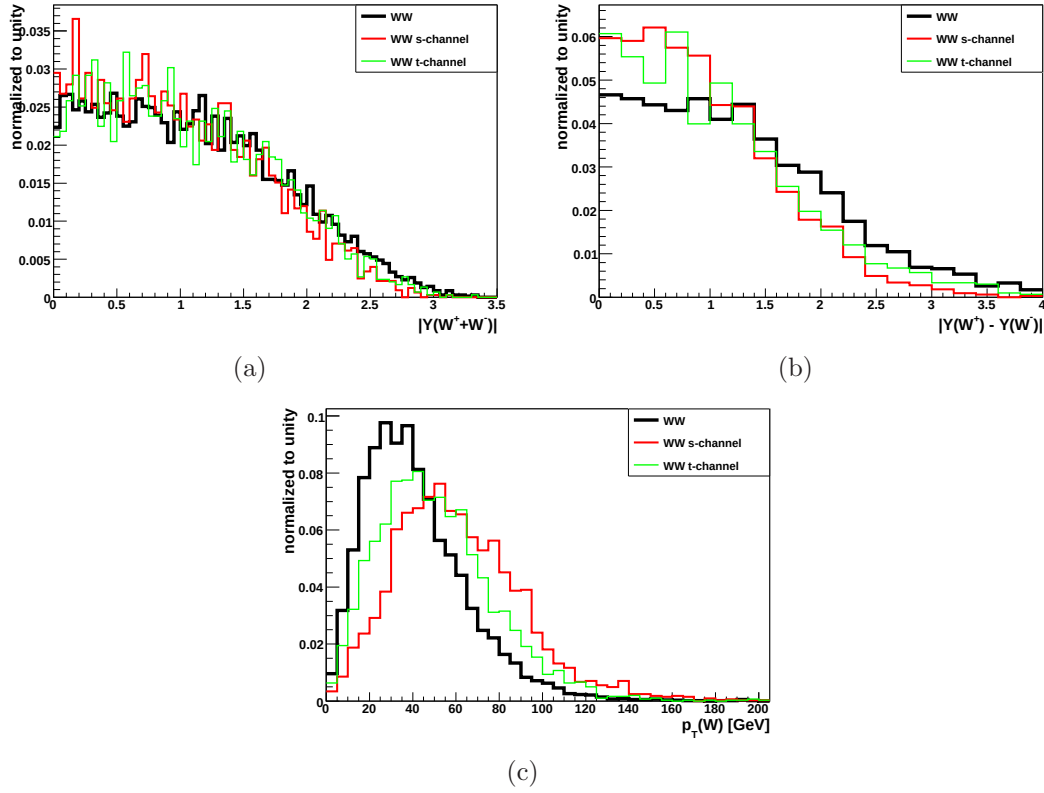


Figure 9.21: Control distributions for the MadGraph Monte Carlo simulation of the $q\bar{q} \rightarrow WW$ process and the s- and t-channel components after all selection cuts against reducible backgrounds and additional cuts on $p_T(W) < 250$ GeV and $\Delta\phi_{ll} < 1$.

- (a): Distribution of $|Y(W^+ + W^-)|$;
- (b): Distribution of $|Y(W^+) - Y(W^-)|$;
- (c): Distribution of $p_T(W)$.

it is not possible to test the correlations of $\Delta\phi_{ll}$, as shown in Figure 9.17, in the same way, due to the divergent behaviour of the two components at large $\Delta\phi_{ll}$. But one can assume that the $\Delta\phi_{ll}$ correlations are also caused by the resonant s-channel component.

In principle it should be possible to measure the production of $q\bar{q} \rightarrow \gamma^{(*)}/Z^{(*)} \rightarrow l^+l^-$ events at large invariant masses ($m_{ll} > 2m_W$) and combine this with the well calculable $\gamma^{(*)}/Z^{(*)} \rightarrow W^+W^-$ decay process in order to improve the Monte Carlo description of the full $q\bar{q} \rightarrow \gamma^{(*)}/Z^{(*)} \rightarrow W^+W^-$ amplitude. However, whether this information can finally be used to correct for the visible correlation in the total $q\bar{q} \rightarrow WW$ process cannot be judged without a more detailed theoretical analysis.

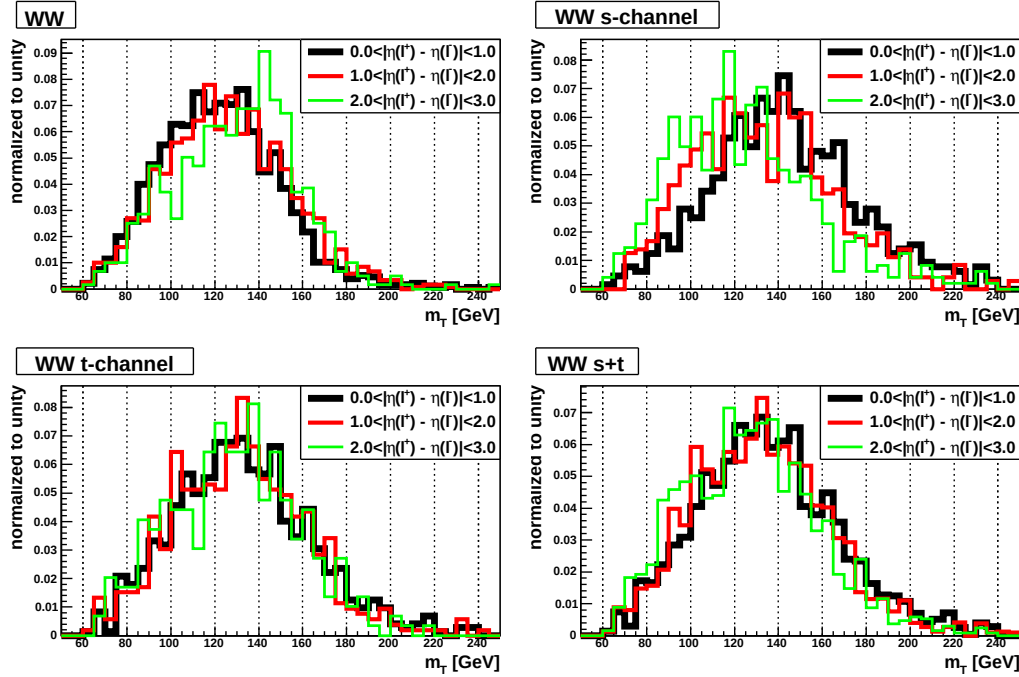


Figure 9.22: Normalised distribution of the transverse mass m_T in three bins of $|\eta(l^+) - \eta(l^-)|$ for the full WW process, the s- and t-channel components and the scalar sum of the two components.

9.5 Combining both background normalisation methods

The background normalisation method introduced in Section 9.4 is suited to discover a Higgs boson, but not the optimal technique to make a precision measurement of the cross section. As it was shown in Section 4.3, the measurement of the $gg \rightarrow H \rightarrow WW$ cross section is essential for precision measurements of the Higgs boson couplings and this will be one of the major goals of Higgs physics at the LHC after an initial discovery.

For a precise cross section measurement it is necessary to know the efficiencies of all acceptance and selection cuts to calculate the total cross section from the final number of observable events. An accurate knowledge of all efficiencies is only achievable with extensive comparisons between observed distributions and Monte Carlo simulations. With the background normalisation method of Section 9.4 it is possible to obtain an almost clean sample of signal events that can be used to adjust the signal Monte Carlo. This signal Monte Carlo can then be directly used to determine the Higgs boson component of all gluon induced processes to the $q\bar{q} \rightarrow WW$ measurement procedure in Section 9.3, because the region

$|Y(l^+ + l^-)| < 1$ and $|\eta(l^+) - \eta(l^-)| < 1$ is identical in both normalisation methods.

Once the bias from the large Higgs boson contribution is removed (at least to some degree), the method of Section 9.3 can be used to obtain the unbiased shape and size of the dominant $q\bar{q} \rightarrow WW$ background and of all gluon induced processes.

The knowledge of the unbiased size and shape of the $q\bar{q} \rightarrow WW$ background can then be used to adjust the $q\bar{q} \rightarrow WW$ Monte Carlo simulation to feed the result back into the method of Section 9.4 in order to remove the largest drawback of this method: the effect of the correlation between $|\eta(l^+) - \eta(l^-)|$ and m_T for the $q\bar{q} \rightarrow WW$ background.

At this point the whole procedure can be reiterated to increase the overall precision. The final result should be very well adjusted Monte Carlo samples for the two dominant processes, $gg \rightarrow H \rightarrow WW$ and $q\bar{q} \rightarrow WW$, that allow a calculation of the total cross section with small systematic uncertainties.

The whole procedure will be more difficult as independent measurements and control samples for other backgrounds, notably $\gamma^{(*)}/Z$, $t\bar{t}$, Wt and W +jet, are needed as well. However, these are expected to be small backgrounds compared to $q\bar{q} \rightarrow WW$ and hence less important to know very precisely. Most of these backgrounds also contain additional or different objects in the final state that allow to select an independent control sample.

9.5.1 Optimising statistical versus systematic uncertainties

Obviously the combination discussed above could also be used to discover the Higgs boson. However, one needs to keep in mind that any reduction in the systematic uncertainty offered by the methods in Section 9.3 and 9.4 comes at the price of the statistical uncertainty from additional measurements of distributions and ratios.

For the background normalisation methods using longitudinal variables the final signal region of the event selection described in Chapter 8 was further divided into subsamples in $|Y(l^+ + l^-)|$ and $|\eta(l^+) - \eta(l^-)|$. Already this division into subsamples increases the relative statistical uncertainty. The statistical uncertainty on the measured value of r and the small size of the background samples due to the additional cut on $\Delta\phi_{ll} < 1$ leads to a further increase of the total statistical uncertainty.

If both background normalisation methods are combined the statistical uncertainty for low integrated luminosities, where one hopes to discover the Higgs boson, is even further increased and hence the method is no longer attractive for an early discovery.

Therefore, in the next chapter only the background measurement strategy from Section 9.4 is applied to the $H \rightarrow WW$ analysis as introduced in Chapter 8.

10 Analysis results based on $\sqrt{s}=10$ TeV simulations

The previous chapter was focused on the development of methods to measure the dominant dilepton backgrounds to the $gg \rightarrow H \rightarrow WW$ analysis. In this chapter the background normalisation method discussed in Section 9.4 is applied to the simple cut based analysis presented in Chapter 8. In addition to the dominant dilepton backgrounds, also the backgrounds from the processes $q\bar{q} \rightarrow WZ$, $q\bar{q} \rightarrow Z \rightarrow \tau\tau$, $q\bar{q} \rightarrow W + \text{jet} \rightarrow l\nu + \text{jet}$ and $q\bar{q}/gg \rightarrow b\bar{b}$ are considered.

For the analysis Monte Carlo samples produced with the ATLAS full simulation are used. In order to avoid a bias from limited Monte Carlo statistics, normally a factor of ≈ 10 more Monte Carlo events than expected events in data are used. For some backgrounds the available fully simulated Monte Carlo statistics is too low¹ to be used for an analysis involving a background shape measurement. In these cases fast simulation samples produced with either Atfast-I or Atfast-II are used and rescaled to agree with the full simulation as far as possible.

10.1 Event selection cuts

The analysis follows the cut based analysis from Chapter 8 up to the cut on $\Delta\phi_{ll} < 1$ (cut 8a). In order to use the method of Section 9.4, an additional cut on the longitudinal separation of the leptons, $|Y(l^+ + l^-)| < 1$, is used that reduces the bias of the lepton acceptance cuts (see Section 9.1.3). This selection cut is listed as cut 9 in the event selection summary tables of this chapter.

So far lepton pairs produced in the decay of heavy hadronic resonances were neglected as background source. In inelastic proton-proton collisions a large number of J/Ψ and Y hadrons are produced that decay into two leptons. The inclusive cross section for both processes is large but not very well known and the rate of two isolated leptons from the $J/\Psi \rightarrow ll$ and $Y \rightarrow ll$ decays is even more uncertain².

¹The Monte Carlo statistics of the W +jets sample corresponds to an integrated luminosity of $\approx 500 \text{ pb}^{-1}$, the Drell-Yan $\gamma^{(*)}/Z^{(*)}$ process to $\approx 1 \text{ fb}^{-1}$. The analysis is aimed at an integrated luminosity of $1 - 10 \text{ fb}^{-1}$.

²Both particles are also produced inside jets in the hadronization phase. How many of these pairs remain as background after isolation requirements is very difficult to predict.

If J/Ψ and Y processes contribute to the total background after all selection cuts, they should be observable as narrow peaks in the invariant dilepton mass and hence it should be possible to isolate and measure these background components. As a precaution, usually a cut of $m_{ll} \gtrsim 10$ GeV is applied. However, this cut would break the strict separation into longitudinal and transverse variables needed for the background measurement, so instead a cut on the transverse dilepton mass $m_T^{\ell\ell} > 10$ GeV could be applied (listed as cut 10 in the tables in this chapter), if this source of background proves to be non negligible.

Applying the cut on $m_T^{\ell\ell} > 10$ GeV to suppress remaining J/Ψ and Y events reduces all other signal and background processes by $\approx 30\%$ (see Tables 10.1–10.3) but leaves all important distributions essentially unchanged. Hence the cut is equivalent to a $\approx 30\%$ increase of the needed integrated luminosity.

However, in the following sections it is assumed that this safety cut is not needed and hence it is not applied. The selection cuts on $p_T(l_1 + l_2)$, E_T^{miss} and $p_T(\text{jet})$ and the final requirement of a transverse mass m_T close to a Higgs boson candidate mass should be sufficient to suppress low mass hadronic resonances which are produced dominantly at low p_T . However, without either data or reliable Monte Carlo predictions, this is impossible to demonstrate.

10.2 Expected signal and background rate

The expected event rate after the additional cuts of Section 10.1 are listed in Table 10.1 for the Higgs bosons signal process with $m_H = 130$ GeV and $m_H = 170$ GeV and for the dominant dilepton backgrounds.

Because of the limited Monte Carlo statistics in full simulation, an Atfast-II sample is used for the $q\bar{q} \rightarrow WW$ background in this chapter and scaled down by 3% to agree³ with the full simulation of $q\bar{q} \rightarrow WW$ after cut 8a.

Also for the $t\bar{t}$ and Wt backgrounds only a few Monte Carlo events remain after all selection cuts. Unfortunately no sufficiently large Atfast-II samples exists to be used. Instead a $\sqrt{s} = 14$ TeV Atfast-I sample is used and scaled to agree to the sum of the full simulation of $t\bar{t}$ and Wt after cut 7 (see Appendix B.2.4 for details).

10.2.1 Additional dilepton backgrounds

So far the processes $q\bar{q} \rightarrow WZ$, $q\bar{q} \rightarrow Z \rightarrow \tau\tau$ and $q\bar{q}/gg \rightarrow b\bar{b}$ were not considered. The summary of the event selection for all three processes is given in Table 10.2.

³The comparison of the full and fast simulation Monte Carlo shapes of the $\Delta\phi_{ll}$ and m_T distributions in Figure B.2 demonstrates that also the distributions of the transverse variables agree sufficiently well.

Selection Cut (all cross sections in fb)	$H \rightarrow WW$		$q\bar{q} \rightarrow WW$		$gg \rightarrow WW$
	$m_H = 130$	$m_H = 170$	full simulation	Atlfast-II	
1-7: reducible bg. cuts	28.94±0.24	77.5±0.9	410±5	427.8±1.3	15.54±0.22
8a: $\Delta\phi_{ll} < 1$	14.78±0.17	43.2±0.7	91.4±2.5	91.4±0.6	4.91±0.12
9: $ Y(l^+ + l^-) < 1$	9.00±0.13	26.6±0.5	48.1±1.9	50.4±0.5	3.15±0.10
10: $m_T^{\mu} > 10$ GeV	5.97±0.11	18.8±0.5	34.4±1.6	35.3±0.4	2.29±0.08

Selection Cut (all cross sections in fb)	$t\bar{t}$	Wt	$t\bar{t}$	$\gamma^{(*)}/Z$
	full simulation		Atlfast-I	
1-7: reducible bg. cuts	22.8±2.0	19.1±1.1	41.9±0.3	76±5
8a: $\Delta\phi_{ll} < 1$	3.1±0.8	4.5±0.5	7.94±0.13	7.2±2.1
9: $ Y(l^+ + l^-) < 1$	2.3±0.7	2.6±0.4	5.06±0.10	5.8±2.1
10: $m_T^{\mu} > 10$ GeV	1.7±0.6	2.0±0.4	3.66±0.09	4.1±1.8

Table 10.1: Accepted cross section in fb for the Higgs boson signal process ($m_H = 130$ GeV and $m_H = 170$ GeV) and the dominant dilepton backgrounds after the application of the successive selection cuts of Chapter 8 and 10 (extension of Table 8.3). For the $q\bar{q} \rightarrow WW$ process also the accepted cross section for the Atlfast-II sample (scaled to agree to the full simulation sample after cut 8a) is shown. For the $t\bar{t}$ and Wt background also the Atlfast-I $t\bar{t}$ sample scaled to agree to the sum of $t\bar{t}$ and Wt after cut 7 is shown.

Selection Cut (all cross sections in fb)	$q\bar{q} \rightarrow WZ$	$q\bar{q} \rightarrow Z \rightarrow \tau\tau$	$q\bar{q}/gg \rightarrow b\bar{b}$
1: 2 isolated leptons	211.6±1.9	5887±25	410±28
2: $p_T(l_1 + l_2) > 30$ GeV	164.0±1.7	566.9±7.8	28.9±7.5
3: $p_T(\text{jet}) < 20$ GeV	36.9±0.8	43.1±2.1	3.8±2.7
4: $E_T^{\text{miss}} > 30$ GeV	32.3±0.8	13.5±1.2	–
5: $\Delta\phi_{ll} < 3$	31.0±0.7	3.0±0.6	–
6: $p_T(l_1 + l_2 + p_T^{\text{miss}}) < 60$ GeV	30.0±0.7	3.0±0.6	–
7: $\Delta\eta_{ll} < 3$	29.5±0.7	3.0±0.6	–
8a: $\Delta\phi_{ll} < 1$	4.13±0.28	–	–
9: $ Y(l^+ + l^-) < 1$	2.45±0.20	–	–
10: $m_T^{\mu} > 10$ GeV	1.91±0.17	–	–

Table 10.2: Accepted cross sections in fb for the additional dilepton background contributions to the $H \rightarrow WW \rightarrow l\nu l\nu$ analysis.

The $q\bar{q} \rightarrow WZ \rightarrow l\nu ll$ process appears as a background, if the third lepton is either outside the detector acceptance or not identified. On the other hand, if events with a third lepton are selected it is expected that a clean sample of WZ events can be obtained and therefore no difficulties due to this background are expected.

The $q\bar{q} \rightarrow Z \rightarrow \tau\tau$ process appears as a background, if both τ leptons decay leptonically into electrons or muons. In this case the neutrinos from the τ decays lead to a sizeable E_T^{miss} . However, both τ leptons are produced preferentially in a back-to-back topology and the visible decay leptons are still back-to-back due to the small proper mass of the τ . Hence the $\Delta\phi_{ll} < 1$ cut reduces this background to a negligible small level (see Appendix B.3.2).

The cross section of the $q\bar{q}/gg \rightarrow b\bar{b}$ background is large and leptons can be produced in leptonic b hadron decays. The b to lepton fragmentation function has a non negligible tail towards high lepton momentum fractions and thus the rate of two isolated leptons from $b\bar{b}$ cannot be ignored. Due to the large cross section, only a sufficiently large Atfast-II Monte Carlo sample is available⁴. As long as the initial lepton p_T selection cuts remain at $p_T > 15$ GeV the Monte Carlo prediction of the $b\bar{b}$ background yields a negligible contribution after all selection cuts⁵.

10.2.2 The W +jet background

The inclusive production of W bosons is a background to the $gg \rightarrow H \rightarrow WW$ analysis, if a jet from QCD initial state radiation is misidentified as a lepton. While the lepton identification efficiencies can be expected to be predicted with high accuracy by the current Monte Carlo simulations, the uncertainties on the misidentification of leptons as isolated electrons or muons are large [15] and hence differences between data and the current Monte Carlo simulations are expected.

An additional complication is that the cross section for W production is large and on the order of 10^8 Monte Carlo events are needed for this analysis. With the ATLAS full simulation, only $5 \cdot 10^6$ events were produced for both, $W \rightarrow e\nu$ and $W \rightarrow \mu\nu$. This is sufficient for a rough estimate of the accepted Monte Carlo cross section after all cuts, but not sufficient to estimate the results of an analysis measuring event distributions. The simulation of a high Monte Carlo statistics sample with Atfast-II was not attempted because of the differences for misidentified leptons shown in Section 7.2.

The accepted cross section for the $W \rightarrow e\nu$ and $W \rightarrow \mu\nu$ processes after each successive cut is shown in Table 10.3 for the standard cut selection and

⁴The sample was produced with a filter of $p_T > 9$ GeV on both leptons at the parton level to reduce the needed amount of CPU time for the detector simulation.

⁵Both leptons are real leptons and the uncertainties from the use of Atfast-II are small in this case.

with weak lepton isolation and event selection cuts⁶, in order to increase the accepted cross section after each cut and thus reduce the relative Monte Carlo statistical uncertainty. If only weak lepton isolation requirements are used, the initial number of dilepton events is increased by a factor ≈ 4 . However, this increase is lost again after the tight jet veto on $p_T(\text{jet}) < 20$ GeV. This also shows that the lepton isolation in the Monte Carlo is strongly correlated to the overall event activity and hence a measurement of the W +jet background using only relaxed isolation cuts is difficult.

If also the jet veto cut is increased to 40 GeV the accepted W +jet cross section increases by a factor of 3–4. The results obtained with these relaxed cuts are scaled to agree with the W +jet background with the standard selection after cut 7 (last column in Table 10.3) and used for this chapter. Within the Monte Carlo statistical uncertainty the standard cut selection and the scaled relaxed W +jet selection agree.

Even with relaxed selection requirements the Monte Carlo statistics is still not sufficient for the shape analysis, but relaxing the selection cuts further does not lead to a significant reduction of the uncertainties.

10.2.3 The influence of detector effects on the measurement of the transverse variables

The background normalisation methods described in Chapter 9 were developed using generator level objects without taking detector efficiency and resolution effects into account.

The ATLAS detector is built almost completely symmetric in ϕ , hence detector effects should affect mainly the longitudinal variables and hence the transverse variables only via correlations between transverse distributions and η dependent efficiency and resolution functions.

The background normalisation method in Section 9.4 directly measures the normalisation ratio from the longitudinal $|\eta(l^+) - \eta(l^-)|$ distribution and is thus insensitive to global detector efficiency and resolution effects that influence the selection cut efficiencies. Furthermore, the selection of events in the region $|Y(l^+ + l^-)| < 1$ ensures a central WW system in both $|\eta(l^+) - \eta(l^-)|$ regions. The effect of QCD radiation is, to first order, the same if a particle or jet is emitted from an incoming parton or a t-channel parton in the case of non-resonant processes. Hence it can be expected that the jet structure of events (and also the structure of soft particles) in the two $|\eta(l^+) - \eta(l^-)|$ regions is very similar.

Figure 10.1 compares the shapes of the p_T and η distributions of the leading p_T jet after selection cut 2 for the Higgs boson signal process and the dominant dilepton backgrounds. Apart from small differences, the distributions are identical for all processes and thus all jet and E_T^{miss} related detector effects should affect

⁶For a discussion of the lepton selection cuts, see Appendix B.4

Selection Cut (all accepted cross sections in fb)	standard lepton selection	$q\bar{q} \rightarrow W + \text{jet} \rightarrow e\nu + \text{jet}$		
		medium e ; weak lepton isolation	medium e ; weak iso.; $p_T(\text{jet}) < 40$ GeV	scaled by 0.23 $\pm$ 0.08
1: 2 isolated leptons	91 $\pm$ 14	403 $\pm$ 30	403 $\pm$ 30	93.8 $\pm$ 6.9
2: $p_T(l_1 + l_2) > 30$ GeV	50 $\pm$ 10	255 $\pm$ 23	255 $\pm$ 23	59.2 $\pm$ 5.5
3: $p_T(\text{jet}) < 20(40)$ GeV	38.8 $\pm$ 9.2	45.3 $\pm$ 9.9	134 $\pm$ 17	31.1 $\pm$ 4.0
4: $E_T^{\text{miss}} > 30$ GeV	28.0 $\pm$ 7.8	34.5 $\pm$ 8.6	106 $\pm$ 15	24.6 $\pm$ 3.5
5: $\Delta\phi_{ll} < 3$	28.0 $\pm$ 7.8	34.5 $\pm$ 8.6	101 $\pm$ 15	23.6 $\pm$ 3.4
6: $p_T(l_1 + l_2 + p_T^{\text{miss}}) < 60$ GeV	25.9 $\pm$ 7.5	32.4 $\pm$ 8.4	97 $\pm$ 14	22.6 $\pm$ 3.4
7: $\Delta\eta_{ll} < 3$	21.6$\pm$6.8	28.0 $\pm$ 7.8	93 $\pm$ 14	21.6$\pm$3.3
8a: $\Delta\phi_{ll} < 1$	12.9 $\pm$ 5.3	12.9 $\pm$ 5.3	43.2 $\pm$ 9.6	10.0 $\pm$ 2.2
9: $ Y(l^+ + l^-) < 1$	8.6 $\pm$ 4.3	8.6 $\pm$ 4.3	23.7 $\pm$ 7.2	5.5 $\pm$ 1.7
10: $m_T^{\mu} > 10$ GeV	8.6 $\pm$ 4.3	8.6 $\pm$ 4.3	19.4 $\pm$ 6.5	4.5 $\pm$ 1.5

Selection Cut (all accepted cross sections in fb)	standard lepton selection	$q\bar{q} \rightarrow W + \text{jet} \rightarrow \mu\nu + \text{jet}$		
		medium e ; weak lepton isolation	medium e ; weak iso.; $p_T(\text{jet}) < 40$ GeV	scaled by 0.31 $\pm$ 0.07
1: 2 isolated leptons	169 $\pm$ 19	609 $\pm$ 36	609 $\pm$ 36	188 $\pm$ 11
2: $p_T(l_1 + l_2) > 30$ GeV	98 $\pm$ 14	392 $\pm$ 29	392 $\pm$ 29	120.7 $\pm$ 8.8
3: $p_T(\text{jet}) < 20(40)$ GeV	67 $\pm$ 12	79 $\pm$ 13	240 $\pm$ 22	73.9 $\pm$ 6.9
4: $E_T^{\text{miss}} > 30$ GeV	54 $\pm$ 11	65 $\pm$ 12	182 $\pm$ 19	55.9 $\pm$ 6.0
5: $\Delta\phi_{ll} < 3$	54 $\pm$ 11	65 $\pm$ 12	177 $\pm$ 19	54.6 $\pm$ 5.9
6: $p_T(l_1 + l_2 + p_T^{\text{miss}}) < 60$ GeV	54 $\pm$ 11	65 $\pm$ 12	177 $\pm$ 19	54.6 $\pm$ 5.9
7: $\Delta\eta_{ll} < 3$	50.1$\pm$10	61 $\pm$ 11	163 $\pm$ 18	50.1$\pm$5.7
8a: $\Delta\phi_{ll} < 1$	18.8 $\pm$ 6.3	18.8 $\pm$ 6.3	52 $\pm$ 10	16.1 $\pm$ 3.2
9: $ Y(l^+ + l^-) < 1$	8.3 $\pm$ 4.2	8.3 $\pm$ 4.2	23.0 $\pm$ 6.9	7.1 $\pm$ 2.1
10: $m_T^{\mu} > 10$ GeV	4.2 $\pm$ 3.0	4.2 $\pm$ 3.0	14.6 $\pm$ 5.5	4.5 $\pm$ 1.7

Table 10.3: Accepted cross sections in fb for the W +jet background in the $H \rightarrow WW \rightarrow l\nu l\nu$ analysis with standard lepton identification cuts (1st column) and with relaxed lepton (2nd column) and jet selection requirements (3rd column, jet veto at $p_T(\text{jet}) < 40$ GeV). The 4th column shows the accepted cross sections of the 3rd column rescaled to agree with the standard selection after cut 7.

both $|\eta(l^+) - \eta(l^-)|$ regions in the same way, regardless of potential correlations between η dependent detector effects and transverse variables.

For the leptons the selection of the different $|\eta(l^+) - \eta(l^-)|$ intervals for the signal and background regions results in different lepton distributions in both samples. The p_T and η distributions of the selected leptons are shown in Figure 10.2 for the total sum of the signal and background processes after selection cut 9. While mainly central leptons in the barrel detectors are selected in the signal region $|\eta(l^+) - \eta(l^-)| < 1$, the background region $|\eta(l^+) - \eta(l^-)| > 1$ also contains an enhanced fraction of leptons from the endcap detectors.

As can be seen from the electron and muon efficiency distributions in Figure B.9, the lepton efficiencies, especially for electrons, are different for the barrel and endcap regions of the detector and thus the transverse distributions of $\Delta\phi_{ll}$ and m_T can potentially be affected. However, the p_T and η dependency of the lepton efficiencies can be measured with $Z \rightarrow ll$ events with high precision and can thus be corrected, if needed. Furthermore, a correction for the lepton reconstruction and identification efficiency⁷ based on the prediction of the detector Monte Carlo simulation does not lead to a significant change in the $\Delta\phi_{ll}$ and m_T distributions in the signal and background region, apart from their absolute normalisation. Hence, although this is not an inherent feature of the background normalisation method, the shapes of the $\Delta\phi_{ll}$ and m_T distributions in both $|\eta(l^+) - \eta(l^-)|$ intervals are independent of the lepton reconstruction and identification efficiencies.

10.2.4 The transverse mass distribution for the signal and background processes

Before the systematic uncertainties of the background normalisation method and the discovery potential can be discussed in Sections 10.3 and 10.4, it is necessary to show to which degree the shapes of the transverse mass in the signal ($|\eta(l^+) - \eta(l^-)| < 1$) and background ($|\eta(l^+) - \eta(l^-)| > 1$) regions agree after detector simulation and how they behave for the background processes not considered so far.

Therefore, the comparison of m_T is shown in Figures 10.3–10.7 for the Higgs boson signal processes ($m_H = 130$ GeV and $m_H = 170$ GeV) and for all background processes. Where it is necessary to use fast simulation samples or samples with relaxed selection cuts due to the low Monte Carlo statistics, the comparison for these samples is also shown.

For the comparison of the shapes the background region ($|\eta(l^+) - \eta(l^-)| > 1$) of all samples is weighted by a factor $r = 0.91 \pm 0.08$. r is calculated from the sum of all background processes and chosen as the Monte Carlo ratio of the integral of

⁷The correction was applied by weighting each event with the product of the inverse p_T and η dependent lepton efficiencies as given by Figure B.9.

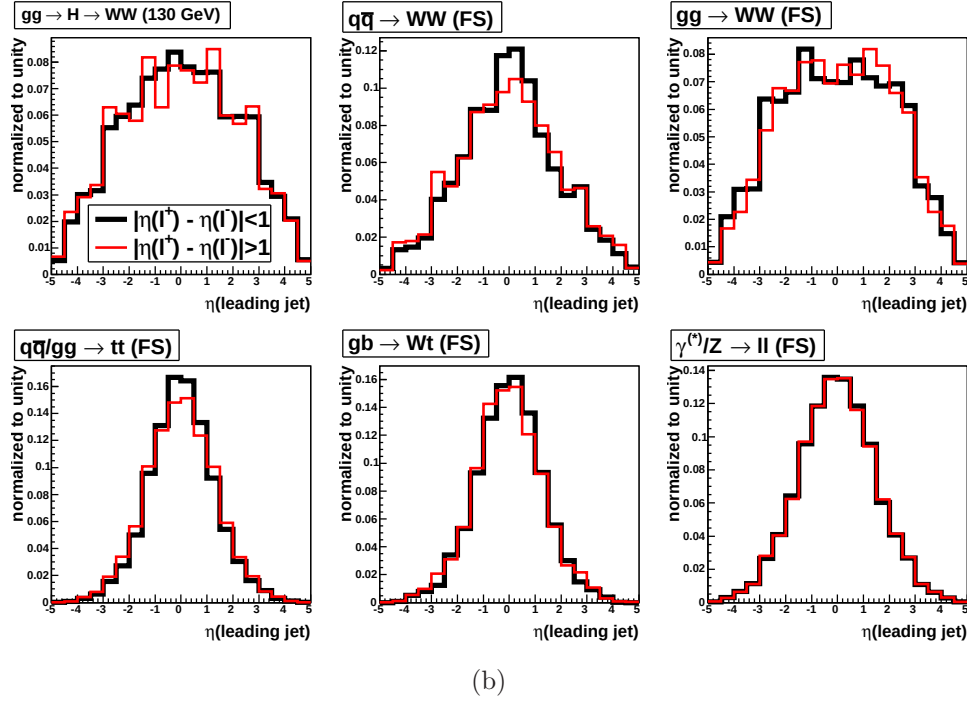
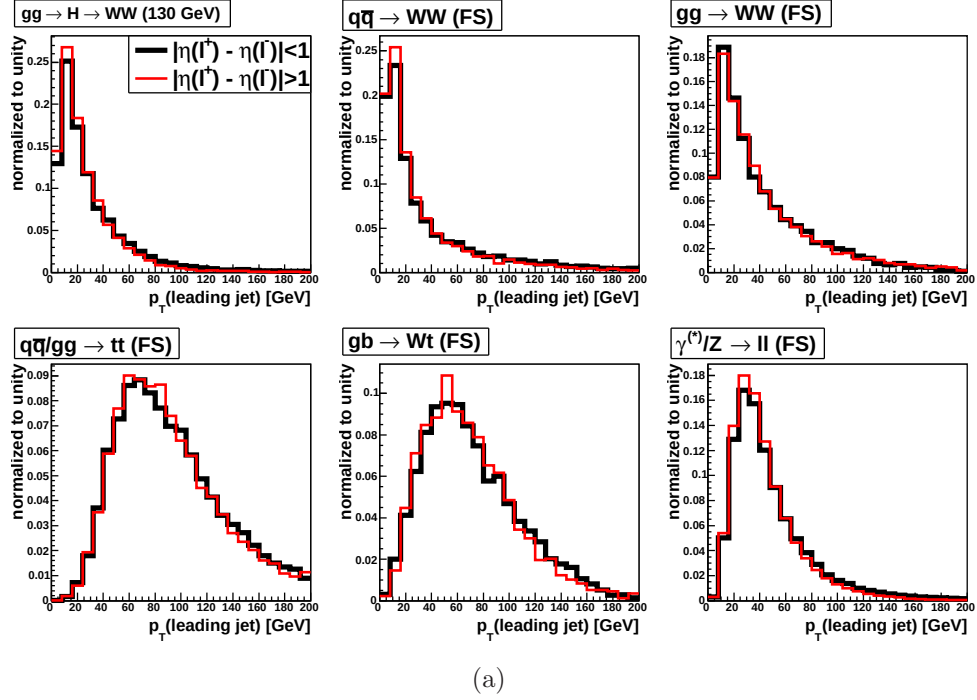


Figure 10.1: Comparison of the leading jet p_T (a) and η (b) distribution between the regions $|\eta(l^+) - \eta(l^-)| < 1$ and $|\eta(l^+) - \eta(l^-)| > 1$ after selection cut 2 for the Higgs boson signal process ($m_H = 130$ GeV) and the dominant dilepton backgrounds.

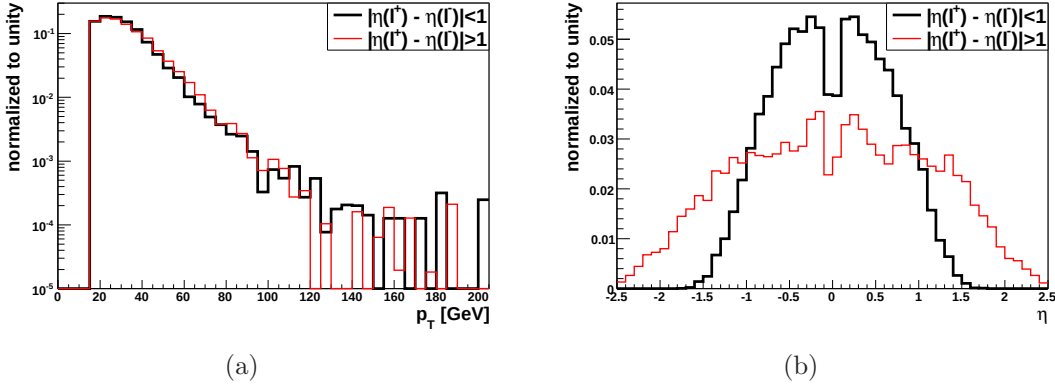


Figure 10.2: Comparison of the lepton p_T (a) and η (b) distribution between the regions $|\eta(l^+) - \eta(l^-)| < 1$ and $|\eta(l^+) - \eta(l^-)| > 1$ for the total sum of the signal and background processes.

events in the interval $m_T = [70, 200]$ GeV between the signal and the background region.

In summary, the comparisons show that the m_T distributions in the signal and background region are mostly uncorrelated for all background processes and that the ratio r calculated from the sum of all background processes is also valid for the individual backgrounds. However, a few differences should be pointed out and discussed:

- For the signal process the observable difference between the number of events in the signal and background region is decreased with respect to the generator level study in Section 9.4. The main reason is that two isolated leptons are required in selection cut 1. As can be seen from the comparison of the $\Delta\phi_{ll}$ distribution in Figure 10.8, the signal region has a large efficiency drop for $\Delta\phi_{ll} < 0.3$ where the two leptons cause each other not to be isolated⁸. For the $|\eta(l^+) - \eta(l^-)| > 1$ background region this is obviously not the case.

Unfortunately it is not trivial to remove the energy deposits of one lepton from the calculation of the isolation properties of the other lepton. Fortunately, this isolation inefficiency has no impact on the m_T distribution or the background normalisation method because of the very weak correlation between $\Delta\phi_{ll}$ and m_T for small values of $\Delta\phi_{ll}$ for the signal and background processes and hence causes no further problems, apart from a reduced signal rate.

⁸The lepton isolation requirements are based on observables in a cone $\Delta R < 0.3$ around the lepton (see Appendix B.4).

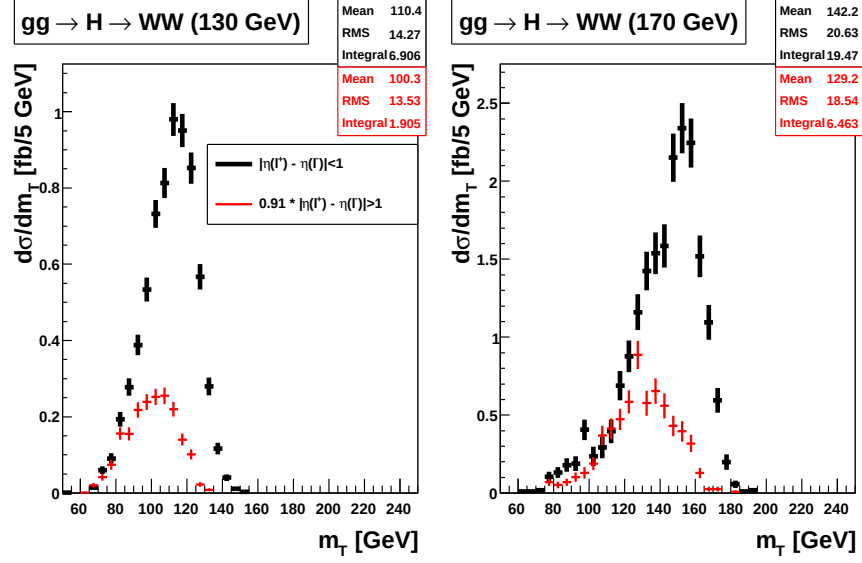


Figure 10.3: Comparison of the m_T distribution in the $|\eta(l^+) - \eta(l^-)| < 1$ signal region to the $|\eta(l^+) - \eta(l^-)| > 1$ background region for the Higgs boson signal processes with $m_H = 130$ GeV and $m_H = 170$ GeV after the selection cut 8a ($\Delta\phi_{ll} < 1$). The background region is scaled by the factor $r = 0.91$ derived from the Monte Carlo ratio between all background events between the two $|\eta(l^+) - \eta(l^-)|$ regions.

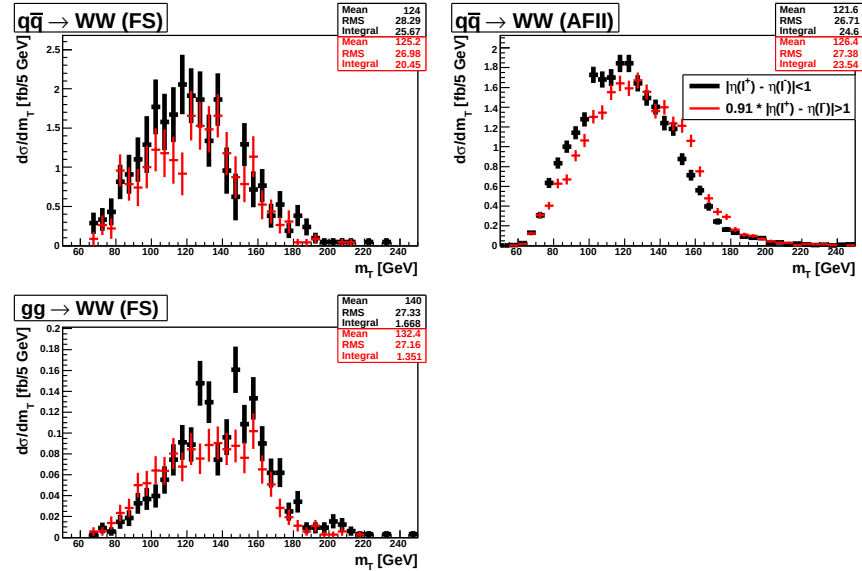


Figure 10.4: The same as Figure 10.3 but for the $q\bar{q} \rightarrow WW$ process in full simulation (FS) and Atfast-II (AFII) and for the $gg \rightarrow WW$ process.

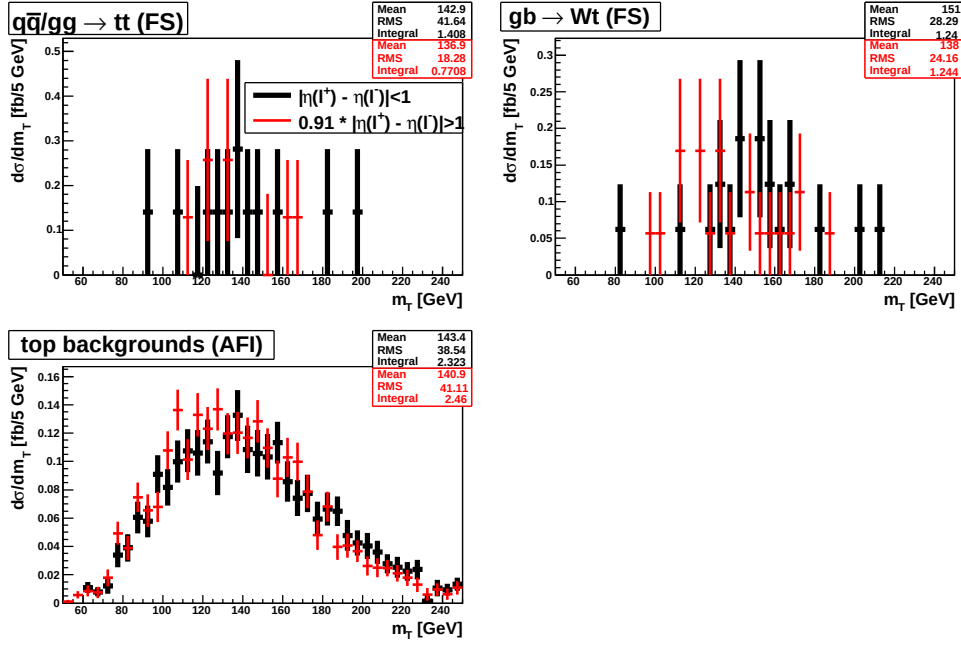


Figure 10.5: The same as Figure 10.3 but for the $t\bar{t}$ and Wt backgrounds in full simulation (FS) and the $t\bar{t}$ background with AtIfast-I (AFI).

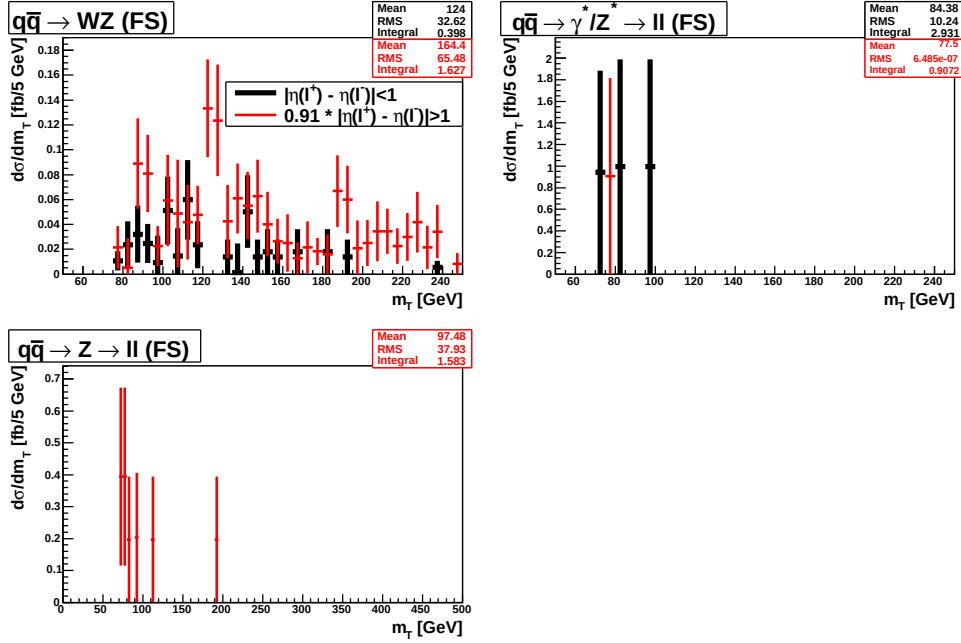


Figure 10.6: The same as Figure 10.3 but for the $q\bar{q} \rightarrow WZ$ process and separately for the Drell-Yan $q\bar{q} \rightarrow \gamma^{(*)}/Z^{(*)}$ ($10 < m_{ll} < 60$ GeV) process and the Z boson production $q\bar{q} \rightarrow \gamma^{(*)}/Z$ ($m_{ll} > 60$ GeV) process. For the Drell-Yan and Z boson process only samples with limited Monte Carlo statistics are available.

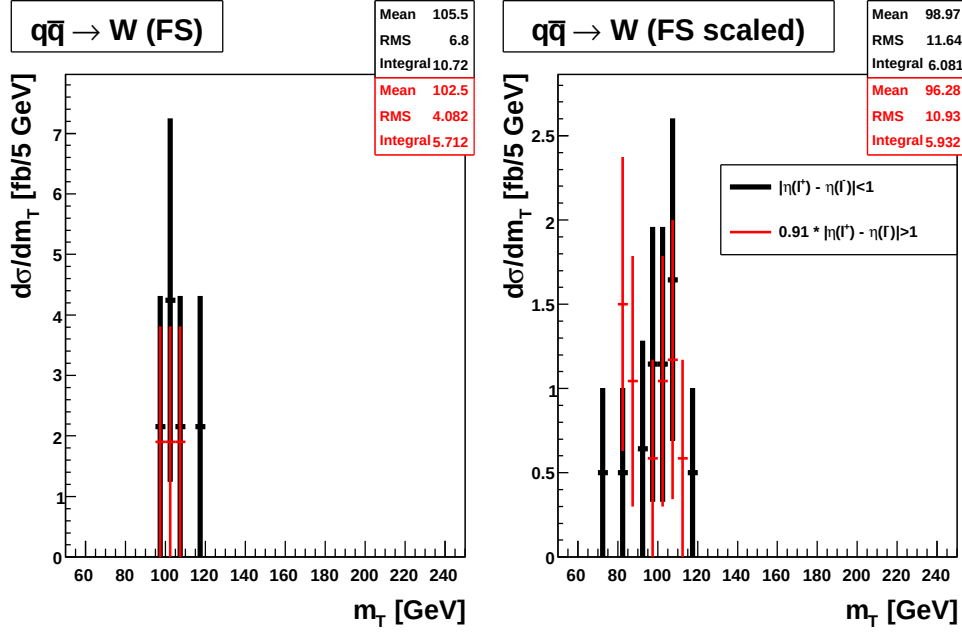


Figure 10.7: The same as Figure 10.3 but for the W +jet background process with standard selection cuts and with relaxed cuts scaled to agree with standard cuts after cut 7.

- The background normalisation does not work very well for the $q\bar{q} \rightarrow WZ$ and $q\bar{q} \rightarrow \gamma^{(*)}/Z$ processes shown in Figure 10.6.

For the $q\bar{q} \rightarrow WZ$ process the ratio r is wrong by a factor of ≈ 4 and also the shape of the m_T distribution does not seem to agree well for the signal and background region.

The Drell-Yan $q\bar{q} \rightarrow \gamma^{(*)}/Z^{(*)}$ process was simulated with a generator level filter for a dilepton invariant mass in the range $10 \text{ GeV} < m_{ll} < 60 \text{ GeV}$. The Monte Carlo statistics of the sample is very limited and no final conclusion can be drawn. However, the transverse mass m_T is very low (also if selection cuts are relaxed) and hence it can be assumed that the contribution of this background to the m_T region where Higgs boson signal events are expected is very small.

The combination of the $\Delta\phi_{ll} < 1$ selection cut with the selection of the signal and background region causes all events from the resonant $q\bar{q} \rightarrow Z \rightarrow ll$ process⁹ (produced with a generator level filter of $m_{ll} > 60 \text{ GeV}$ for the $q\bar{q} \rightarrow \gamma^{(*)}/Z$ process) to be found in the background region and hence the background subtraction does not work.

⁹This was checked with a large Atfast-II Monte Carlo sample

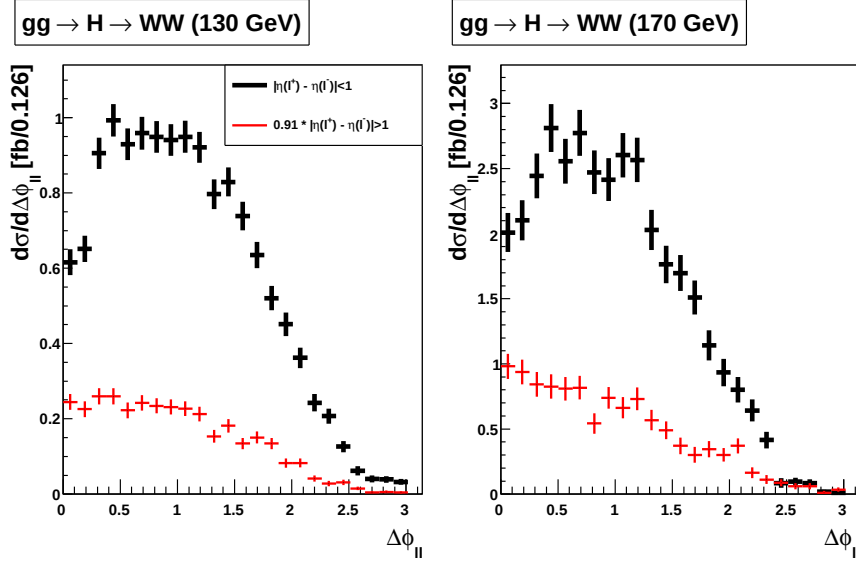


Figure 10.8: Comparison of the $\Delta\phi_{\ell\ell}$ distributions in the $|\eta(l^+) - \eta(l^-)| < 1$ signal region to the $|\eta(l^+) - \eta(l^-)| > 1$ background region for the Higgs boson signal processes with $m_H = 130$ GeV and $m_H = 170$ GeV before the selection cut 8a ($\Delta\phi_{\ell\ell} < 1$). The background region is scaled by the factor $r = 0.91$ derived from the Monte Carlo ratio of all background events between the two $|\eta(l^+) - \eta(l^-)|$ regions after cut 8a.

The invariant dilepton mass $m_{\ell\ell}$ for all three processes above and for the sum of the signal and background processes is shown in Figure 10.9. For the WZ process two cases need to be considered. If the lepton from the W decay is not reconstructed, events from this process behave similarly to $Z \rightarrow \ell\ell$ events, with the exception of having intrinsic E_T^{miss} and thus a high efficiency to pass the event selection. If on the other hand one of the leptons from the Z decay is not reconstructed, these events do not show a peak at the Z boson mass. Fortunately, the background normalisation method works in this case and a similar amount of events with a compatible m_T shape is found in the signal and background region.

Hence, with sufficient luminosity, a fit of the resonant $m_{\ell\ell}$ contribution in the signal and background region can be used to subtract the problematic component of the $(W)Z$ backgrounds with the help of Monte Carlo simulations and the remaining events can then be treated in a similar way as all other backgrounds by the background normalisation method.

For the analysis in this chapter, the total Monte Carlo prediction of the resonant Z contribution of these backgrounds is so small that no special

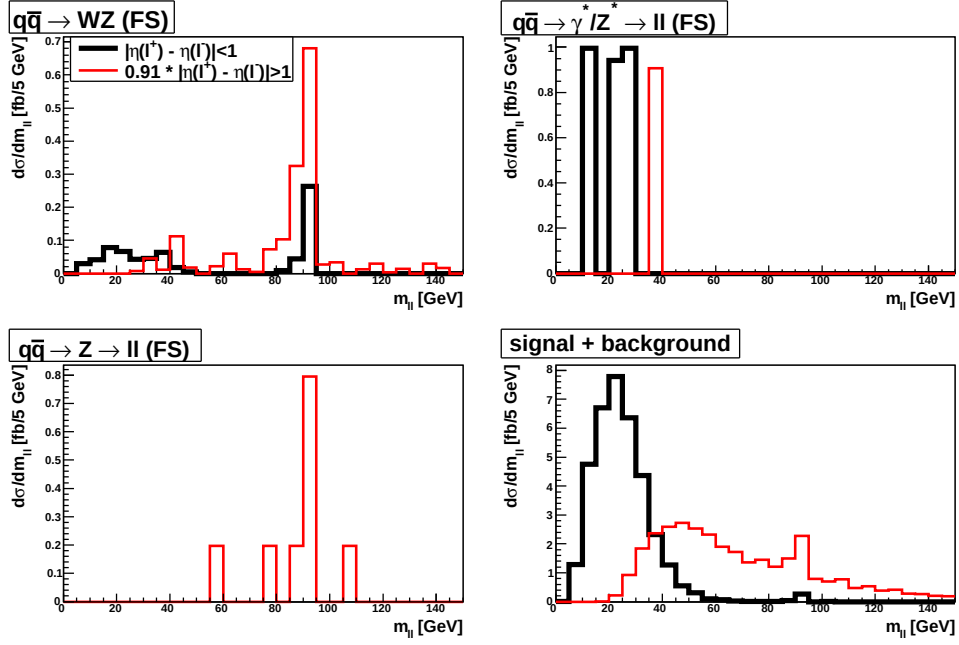


Figure 10.9: Comparison of the dilepton invariant mass m_{ll} distribution in the $|\eta(l^+) - \eta(l^-)| < 1$ signal region to the $|\eta(l^+) - \eta(l^-)| > 1$ background region for the $q\bar{q} \rightarrow WZ$, the Drell-Yan $q\bar{q} \rightarrow \gamma^{(*)}/Z^{(*)}$ ($10 < m_{ll} < 60$ GeV) and the Z boson production $q\bar{q} \rightarrow \gamma^{(*)}/Z$ ($m_{ll} > 60$ GeV) processes as well as for the sum of the Higgs boson signal ($m_H = 130$ GeV) and all background processes.

treatment is necessary.

- The Monte Carlo statistical uncertainty for the W +jet background is large and also for this background no final conclusion can be drawn. However, it appears that the background normalisation does work although this is not an intrinsic feature of the normalisation method. While one can assume that the m_T distributions in the signal and background regions are at least similar, the ratio r for this background is strongly influenced by the differences between the η distributions of jets misidentified as isolated leptons in the two $|\eta(l^+) - \eta(l^-)|$ regions. The uncertainty on the η dependency of the misidentification efficiency is large and it cannot be assumed that the background normalisation of the W +jet background will work on data in the same way as for the current Monte Carlo simulation.

However, like for the Drell-Yan background, the transverse mass m_T of the W +jet background is very low and has only a small overlap with the transverse mass region where Higgs boson signal events are expected. Hence the presence of the W +jet background might make it necessary to adjust

the interval, in which the ratio r is measured when applying the method from Section 9.4, but should not affect the signal extraction otherwise. Nevertheless, an independent method of measuring the W +jet background will be needed to verify this conclusion.

The final comparison of the m_T distributions between the signal and the background region is shown in Figure 10.10 for the sum of all background processes and for the total sum of the Higgs boson signal ($m_H = 130$ GeV) and all background processes. Because of the large statistical fluctuations from the $q\bar{q} \rightarrow \gamma^{(*)}/Z$ and $q\bar{q} \rightarrow W$ +jet processes, both comparisons are also shown without the contribution of these backgrounds. The ratio r is almost unaffected by this change and is $r = 0.89 \pm 0.015$ without and $r = 0.91 \pm 0.08$ with these background contributions, respectively.

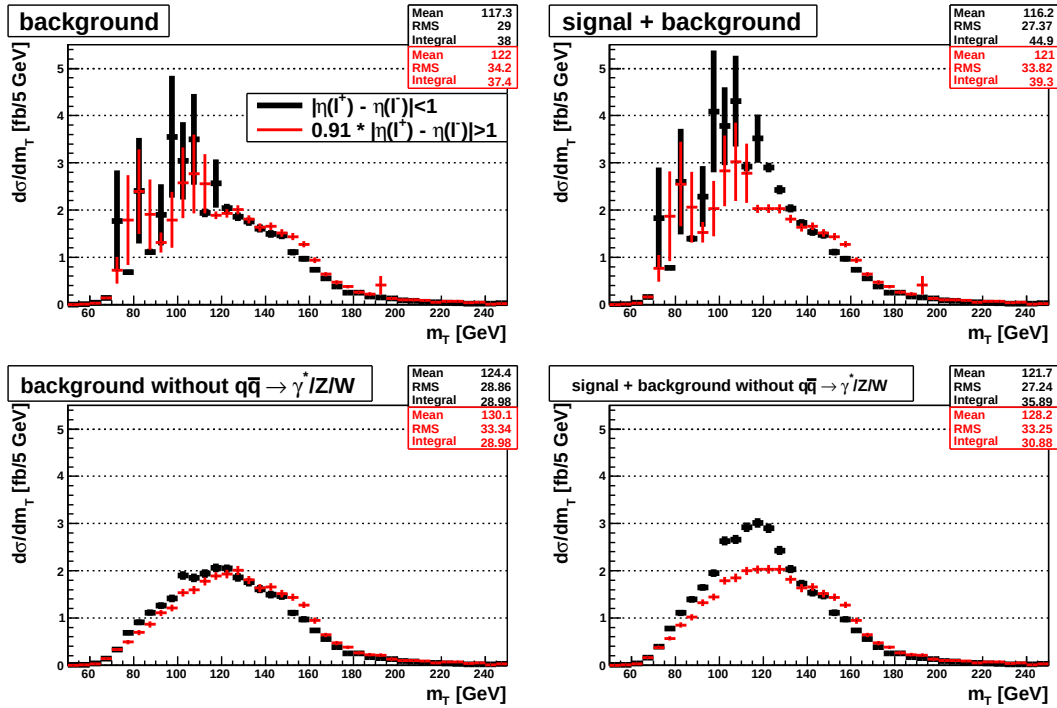


Figure 10.10: The same as Figure 10.3 but for the sum of all background processes and the sum of the Higgs boson signal ($m_H = 130$ GeV) and all background processes. In addition, the distributions where the $q\bar{q} \rightarrow \gamma^{(*)}/Z$ and W +jet backgrounds are omitted because of their large Monte Carlo statistical uncertainty are shown.

If a Higgs boson signal is present, the ratio r is measured outside the m_T interval where signal events are expected. The measured background extrapolation ratio r as well as the needed background ratio r^{true} in the signal region and the observed

ratio $r^{\text{obs.}}$ after the background extrapolation are shown in the same way as for the generator level analyses (see Figure 9.19) in Figure 10.11 as function of the assumed Higgs boson mass for all background components and for all backgrounds without the $q\bar{q} \rightarrow \gamma^{(*)}/Z$ and $W+\text{jet}$ background because of their large Monte Carlo statistical uncertainty.

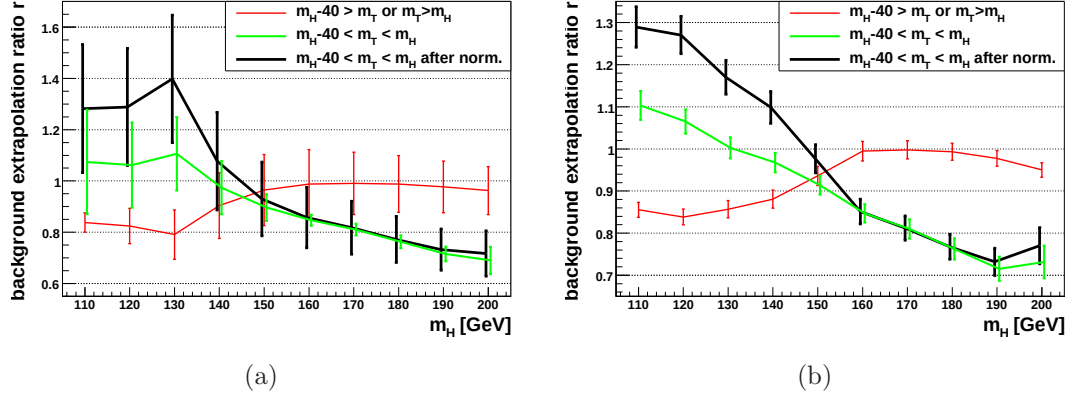


Figure 10.11: Measured background extrapolation ratio r (red) as function of the Higgs boson mass used in the analysis for a background only experiment without a Higgs boson signal contribution. The true background extrapolation ratio r^{true} in the interval M_T^{signal} is shown in green. The observable ratio $r^{\text{obs.}}$ in the interval M_T^{signal} after rescaling of the X_B distribution with r is shown in black. In Figure (a) all background components are included, in Figure (b) the $q\bar{q} \rightarrow \gamma^{(*)}/Z$ and $q\bar{q} \rightarrow W+\text{jet}$ processes with large statistical fluctuations are excluded. The error bars give the Monte Carlo statistical uncertainty.

$$\begin{aligned}
 r &= r(m_T < m_H - 40 \text{ GeV or } m_T > m_H) \\
 r^{\text{true}} &= r(m_H - 40 \text{ GeV} < m_T < m_H) \\
 r^{\text{obs.}} &= r(m_H - 40 \text{ GeV} < m_T < m_H)/r
 \end{aligned}$$

If a Higgs boson signal is present, the measured background extrapolation ratio r is affected as well. A comparison of the resulting distributions of m_T is shown in Figure 10.12 for a Higgs boson signal of $m_H = 130$ GeV and $m_H = 170$ GeV, together with the background only distributions in the signal region before and after the background subtraction. Due to the large statistical fluctuations the $q\bar{q} \rightarrow \gamma^{(*)}/Z$ and $q\bar{q} \rightarrow W+\text{jet}$ processes are not included.

The m_T distributions for all Higgs boson masses in the range from 110 GeV to 200 GeV are shown in Appendix D in Figures D.5–D.9 without the contribution of the $q\bar{q} \rightarrow \gamma^{(*)}/Z$ and $q\bar{q} \rightarrow W+\text{jet}$ backgrounds and in Figures D.10–D.14 if these backgrounds are included.

The comparison with the generator level distributions in Section 9.4 shows no

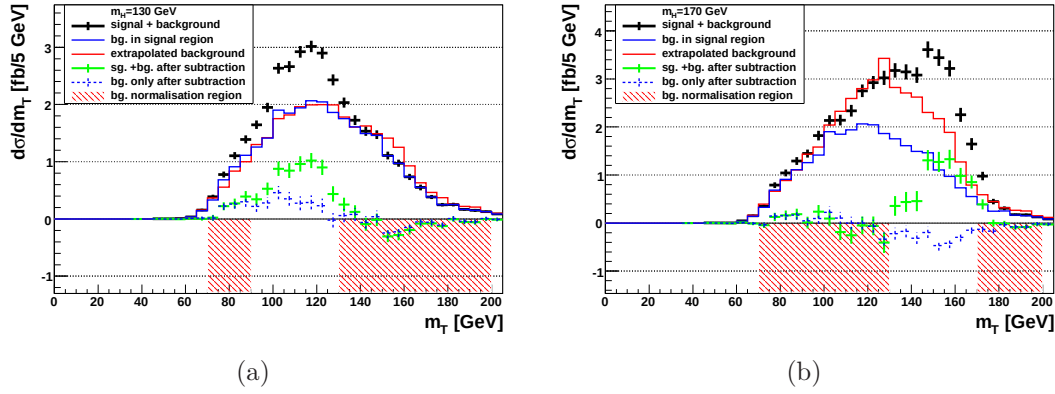


Figure 10.12: Comparison of the true, measured and subtracted m_T distributions for $m_H = 130$ GeV (a) and $m_H = 170$ GeV (b). Black: the measured distribution of m_T in X_S . Red: the measured distribution in X_B scaled with the measured r . Green: the m_T distribution after the subtraction of the first two. Solid blue: the true background distribution in X_S . Dashed blue: The subtracted m_T distribution if no signal is present. The hashed area indicates the m_T interval $M_T^{\text{norm.}}$, in which r is measured. The error bars show the Monte Carlo statistical uncertainty.

qualitative difference apart from the smaller signal yield due to the lepton isolation and hence it can be concluded that the background normalisation method is almost unaffected by reconstruction and identification efficiencies and energy resolution effects.

As discussed before, the influence of the W +jet and $q\bar{q} \rightarrow \gamma^{(*)}/Z$ backgrounds could affect the analysis for $m_H \lesssim 130$ GeV severely. Due to the large statistical fluctuations, no conclusion can be reached on the background normalisation in this range, but because of the small extracted signal for small Higgs boson masses (see Figure D.5), this region is not promising.

For Higgs boson masses $m_H \gtrsim 130$ GeV it might be necessary to adjust the interval in which r is measured, if large contributions from the W +jet background are expected from independent control measurements in the relevant m_T regions.

10.2.5 The accumulation of statistical uncertainties

Due to the additional selection cuts compared to the analysis in Chapter 8 and the separation of events into a signal and a background region, the statistical uncertainties are increased. Figure 10.13 shows the dependency of the significance on the systematic uncertainty of the factor r for Higgs boson masses of $m_H = 130$ GeV and $m_H = 170$ GeV and for assumed integrated luminosities of 10 and 30 fb^{-1} (compared to 1 and 10 fb^{-1} in Chapter 8).

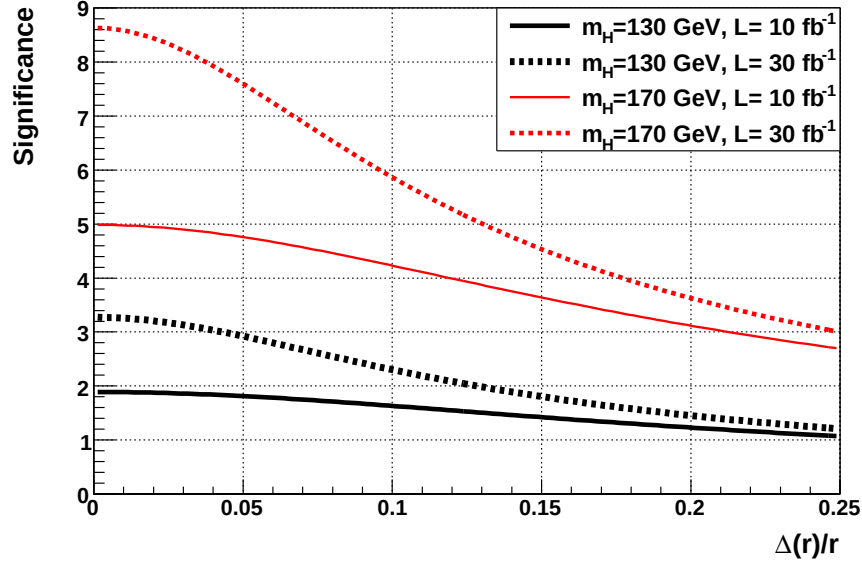


Figure 10.13: Dependency of the significance S on the systematic uncertainty of the background extrapolation factor r for Higgs boson masses of $m_H = 130$ GeV and $m_H = 170$ GeV and assumed integrated luminosities of 10 and 30 fb^{-1} .

The significance is calculated as

$$S = \frac{n_{S+B}^{\text{sub.}} - n_B^{\text{sub.}}}{\sqrt{[\Delta(n_B^{\text{sub.}})]^2 + \left[\frac{\Delta(r)}{r} \cdot r \cdot n_B^{\text{bg.}}\right]^2}}, \quad (10.1)$$

where $n_{S+B}^{\text{sub.}}$ is the number of events after subtraction when a Higgs boson signal is present, $n_B^{\text{sub.}}$ is the number of events after subtraction when no signal is present and $n_B^{\text{bg.}}$ is the number of events in the background region when no signal is present. $\Delta(n_B^{\text{sub.}})$ is the total statistical uncertainty on $n_B^{\text{sub.}}$ calculated, for simplicity, in the Gaussian approximation and contains the statistical uncertainty from the measurement of r .

The comparison to the significance obtained in the simple cut based analysis in Figure 8.12 shows that the overall significance for both Higgs boson masses is reduced by an approximate factor of 2.5 due to the accumulation of statistical uncertainties. A reasonable upper limit on the systematic uncertainty on the background extrapolation factor r , such that the systematic uncertainties are much smaller than the statistical uncertainties, would again be 5% to 10%.

If it were possible to recover the signal efficiency loss from the failed selection of isolated leptons at low $\Delta\phi_{ll}$, the signal and hence also the significance would increase by 10% to 15%. Furthermore, the significance is calculated only from the number of events after subtraction without consideration of the signal peak

structure. Using the full information from the expected signal and the measured background shape should give a considerable improvement of the significance. However, it is unlikely that these possible improvements could completely recover the loss in significance due to the statistical uncertainties of the additional measurements.

10.3 Systematic background uncertainties

The same systematic uncertainties as for the cut based analysis in Chapter 8 are considered. The systematic uncertainties are separated in factors influencing the shape of the dominant $q\bar{q} \rightarrow WW$ background and factors influencing the size of the small background contributions from top production ($t\bar{t}$ and Wt), as well as from the $gg \rightarrow WW$ background. A study of the systematic uncertainty from the W +jet background is unfortunately not possible due to the limited Monte Carlo statistics.

In the Sections 8.3–8.5 the dependency of the background normalisation factor $f_{n_C \rightarrow \hat{n}_B}$ on the variation of one systematic background parameter was studied for the analysis of Higgs boson masses of $m_H = 130$ GeV and $m_H = 170$ GeV. The relevant factor in the context of the background normalisation method of Section 9.4 is the observable ratio $r^{\text{obs.}}$ between the number of background events in the signal region and the extrapolated number of background events in the signal region for the background only case where no Higgs boson signal is present.

10.3.1 Systematic uncertainties on the shape of the $q\bar{q} \rightarrow WW$ background

In Section 8.3 the dependency of the background normalisation on the E_T^{miss} energy scale and generator variations of the PDFs and QCD scales were analysed.

Figure 10.14 shows the systematic variation of $r^{\text{obs.}}$ when the E_T^{miss} cut value is varied by up to ± 10 GeV. $r^{\text{obs.}}$ is almost constant and one can conclude that the uncertainty due the E_T^{miss} energy scale is well below 5%.

The variation of the ratio $r^{\text{obs.}}$ when the 40 CTEQ6M 1σ error PDFs are used is shown in Figure 10.15(a) and the variation of $r^{\text{obs.}}$ when the renormalization (μ_R) and factorization (μ_F) scale is varied by factors of up to 8 is shown in Figure 10.15(b). All samples are statistically independent and the observable variation of $\approx 5\%$ is consistent with the expected statistical fluctuations. Hence it can be concluded that the systematic uncertainty of $r^{\text{obs.}}$ due to PDF and QCD scale variations of the dominant $q\bar{q} \rightarrow WW$ background is at most of the order of 5%.

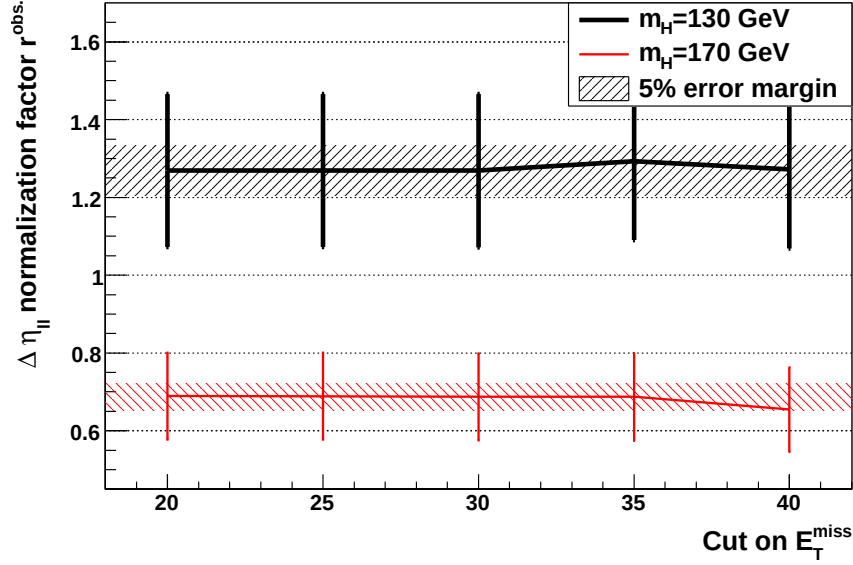


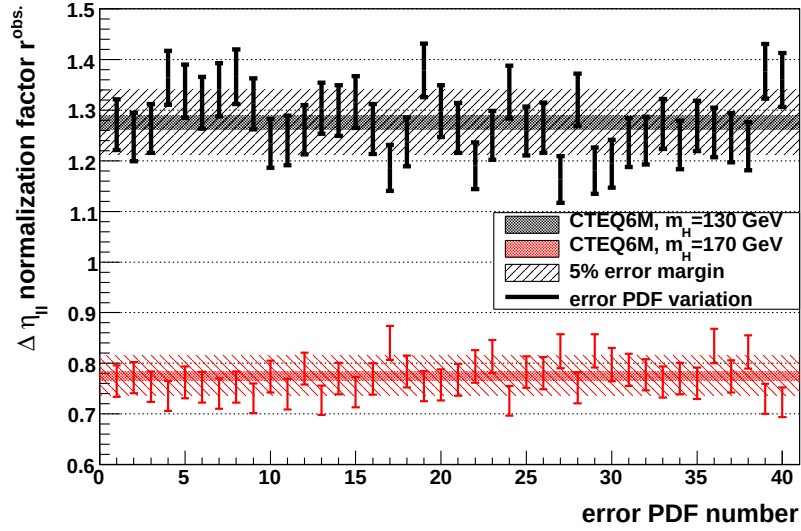
Figure 10.14: Systematic variation of the ratio $r^{\text{obs.}}$ after all event selection cuts as a function of the E_T^{miss} cut value. The variation of the normalisation factor for $m_H = 130$ GeV is shown in black and for $m_H = 170$ GeV in red. The light hashed area indicates the size of a 5% variation of the central $f_{n_C \rightarrow \hat{n}_B}$ value.

10.3.2 Systematic uncertainties due to variations of the size of the top and $gg \rightarrow WW$ backgrounds

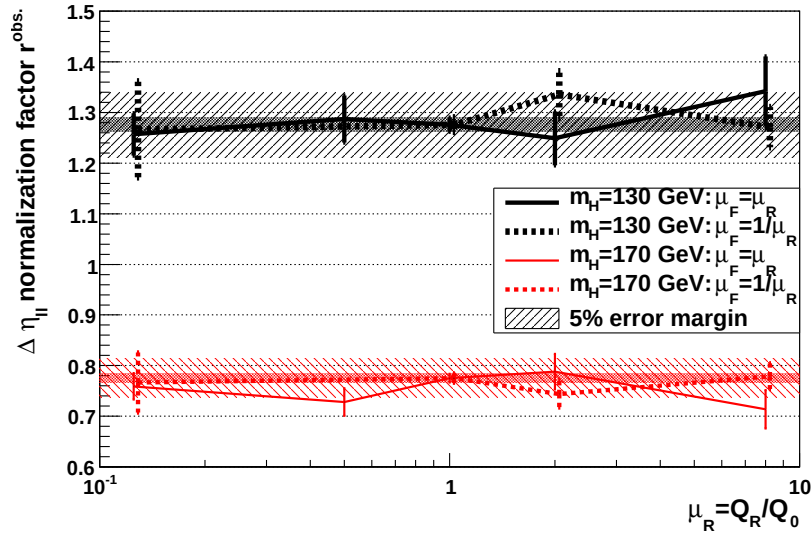
In Sections 8.4 and 8.5 the systematic uncertainty due to a variation of the accepted cross section for the top backgrounds ($t\bar{t}$ and Wt) and the $gg \rightarrow WW$ background were considered. The variation of $r^{\text{obs.}}$ when the accepted top and $gg \rightarrow WW$ background cross sections are varied by factors of up to 8 are shown in Figure 10.16.

The dependency of the background normalisation on the accepted cross section of these backgrounds is considerably reduced by the normalisation method of Section 9.4 with respect to the analysis from Chapter 8. In order to stay within a systematic uncertainty of $\approx 5\%$, the accepted cross section of the top backgrounds needs to be known within a factor of ≈ 8 and the $gg \rightarrow WW$ background within a factor of ≈ 4 .

Especially in the case of the $m_H = 170$ GeV analysis, where the m_T distribution of the $gg \rightarrow WW$ background is expected to be very similar to the signal process, this reduced dependency of the background extrapolation on the knowledge of the $gg \rightarrow WW$ cross section is essential.



(a)



(b)

Figure 10.15: Systematic variation of the ratio $r^{\text{obs.}}$ after all event selection cuts due to PDF (a) and QCD scale variations (b). The central dark bar shows the central prediction of the ratio, the width of the bar is the Monte Carlo statistical uncertainty. The error bars give the prediction after variation, the size of the error bar is the Monte Carlo statistical uncertainty. The light hashed area indicates the size of a 5% variation of the central ratio. In (a) the 40 CTEQ6M 1σ error PDFs are used. In (b) the renormalization (μ_R) and factorization (μ_F) scale is varied. Considered are the cases $\mu_F = \mu_R$ and $\mu_F = 1/\mu_R$.

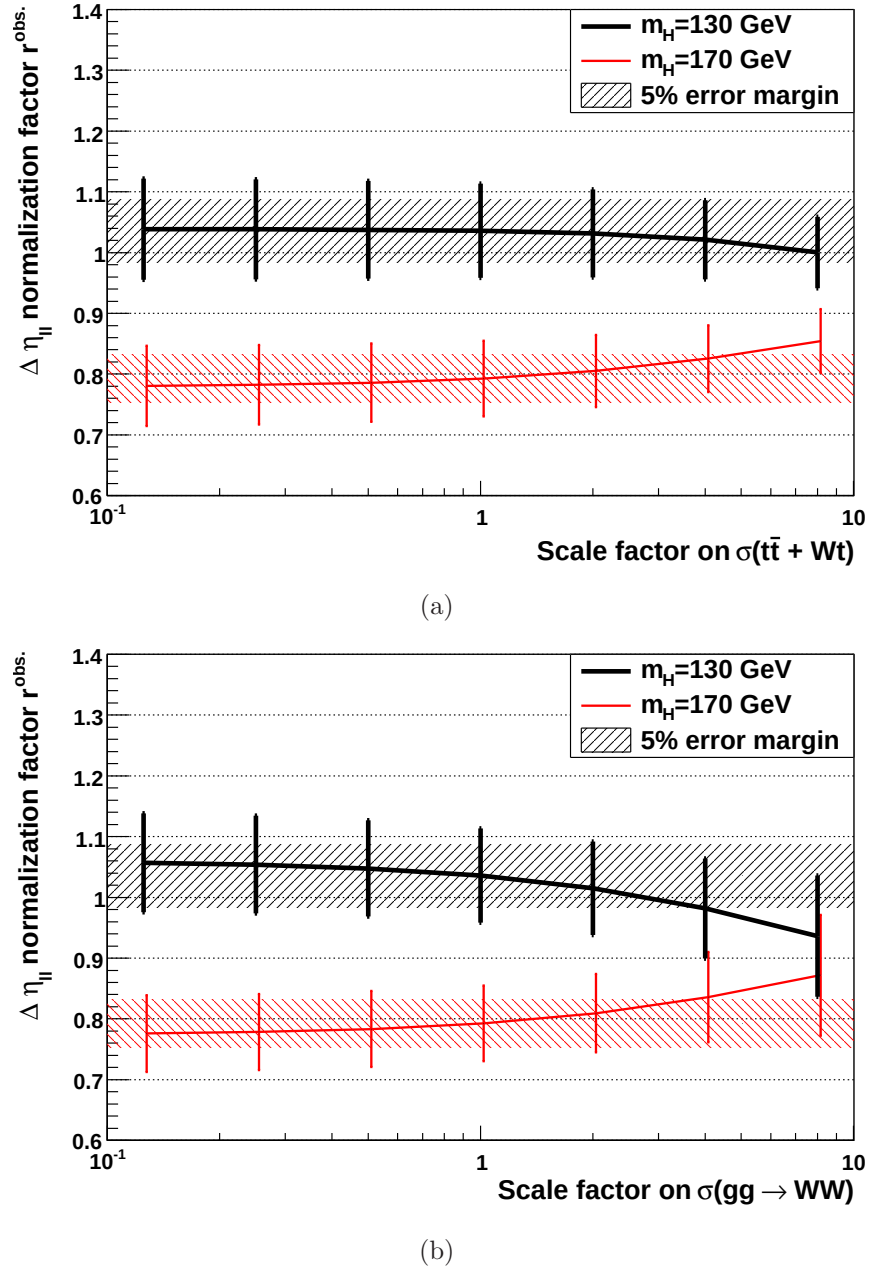


Figure 10.16: Systematic variation of the ratio $r^{\text{obs.}}$ after all event selection cuts due to variations of the accepted top (a) and $gg \rightarrow WW$ (b) background cross section by factors of up to 8.

10.3.3 Total systematic uncertainty

The variation of the various sources of systematic uncertainties by large factors show that r can be well predicted. The variation due to the E_T^{miss} energy scale and the relative size of the top and $gg \rightarrow WW$ background should give a contribution to r well below the 5% level. Due to Monte Carlo statistical limitations in the evaluation of the PDF and scale dependency, it is not possible to set a limit on the systematic uncertainty of r below the 5% level, hence a total systematic uncertainty of 5% on r is used in the following sections.

10.3.4 Further reduction of systematic uncertainties

The dominant part of the residual systematic variations of $r^{\text{obs.}}$ is caused by the relative shift in the m_T distribution of the $q\bar{q} \rightarrow WW$ background between the signal and background regions. Without this shift, the dependency of r on m_H in Figure 10.11 would be almost flat and changes in the relative size and shape of individual background components would be less pronounced.

If the $q\bar{q} \rightarrow WW$ background can be measured, as it was discussed in Section 9.5, this shift can be corrected and it can be expected that the remaining systematic uncertainties are considerably reduced. However, the independent measurement of the $q\bar{q} \rightarrow WW$ background also comes with its own statistical uncertainties, which makes this approach not feasible for integrated luminosities of $\approx 10 \text{ fb}^{-1}$.

10.4 Discovery potential

The discovery potential of a Higgs boson channel is evaluated in terms of the probability that a fluctuation of the background is seen as signal at a specific mass (see Section 3.3). Because of the large number of statistically independent measurements in the signal and background region, a calculation of the discovery significance using the Gaussian approximation for the Poisson distributions and Gaussian error propagation is not sufficiently accurate.

Instead, a large number of Monte Carlo experiments is used, in which the measurements in the signal and background region are performed using histograms generated with a Poisson distribution from the Monte Carlo predictions. The systematic uncertainty on the background extrapolation factor r is taken into account by using a Gaussian nuisance parameter with $\Delta(r)/r = 0.05$ in the generation of the Monte Carlo experiments. The figure of merit for the distinction between signal and background is the number of events $n^{\text{sub.}}$ in the signal region $m_H - 40 \text{ GeV} < m_T < m_H$ after subtraction of the scaled m_T distribution from the background region. The significance is extracted from the probability for a background only Monte Carlo experiment to yield more events $n^{\text{sub.}}$ than the median expectation μ from signal plus background Monte Carlo experiments.

The distribution of $n^{\text{sub.}}$ produced from a large number of Monte Carlo experiments for the background only and the signal plus background hypothesis is shown in Figures 10.17 and 10.18 for Higgs boson masses between 120 GeV and 190 GeV, assuming an integrated luminosity of 10 fb^{-1} . The fraction of background only Monte Carlo experiments that show more events $n^{\text{sub.}}$ than the median expectation for the signal plus background case is indicated by a red area. The probability of this fraction and the corresponding Gaussian standard deviation for this probability are given in Table 10.4.

Higgs Boson mass m_H	120 GeV	130 GeV	140 GeV	150 GeV
# exp. signal events μ	46.3	62.3	66.6	84.6
# exp. background events	26.4	20.8	13.7	-2.7
$P(n^{\text{sub.}} > \mu)$	0.157	0.028	0.011	$4.6 \cdot 10^{-5}$
σ	1.4	2.2	2.6	4.1

Higgs Boson mass m_H	160 GeV	170 GeV	180 GeV	190 GeV
# exp. signal events μ	84.4	73.5	29.0	2.3
# exp. background events	-20.3	-21.5	-20.5	-15.9
$P(n^{\text{sub.}} > \mu)$	$1.2 \cdot 10^{-7}$	$1.1 \cdot 10^{-9}$	$3.8 \cdot 10^{-4}$	0.064
σ	5.3	6.1	3.5	1.9

Table 10.4: Expected number of events after the subtraction of the background region for the signal plus background hypothesis and for the background only hypothesis for Higgs boson masses between 120 GeV and 190 GeV and an integrated luminosity of 10 fb^{-1} . $P(n^{\text{sub.}} > \mu)$ gives the probability of the fraction of background only Monte Carlo experiments that show more events $n^{\text{sub.}}$ than the median expectation μ for the signal plus background hypothesis. σ gives the Gaussian significance for this fraction and hence the discovery potential.

Due to CPU time constraints, a maximum of $5 \cdot 10^8$ Monte Carlo experiments were performed, which is sufficient to reach a $\approx 6\sigma$ discovery significance. For Higgs boson masses of 160 GeV and 170 GeV only a few events in the Monte Carlo experiments for the background only case exceed the median expectation. Hence the significance is also calculated from a fit of a Gaussian distribution to the tail of the background only Monte Carlo experiments.

It should be noted that the 1σ fluctuations of the background only experiments, indicated by a black hashed area in Figures 10.17 and 10.18, are compatible with $n^{\text{sub.}} = 0$ for all Higgs boson masses. Hence the systematic structure in the shape of m_T after the subtraction of the background region is not observable within the statistical uncertainties.

In Figure 10.19 the discovery significance is shown as a function of the Higgs

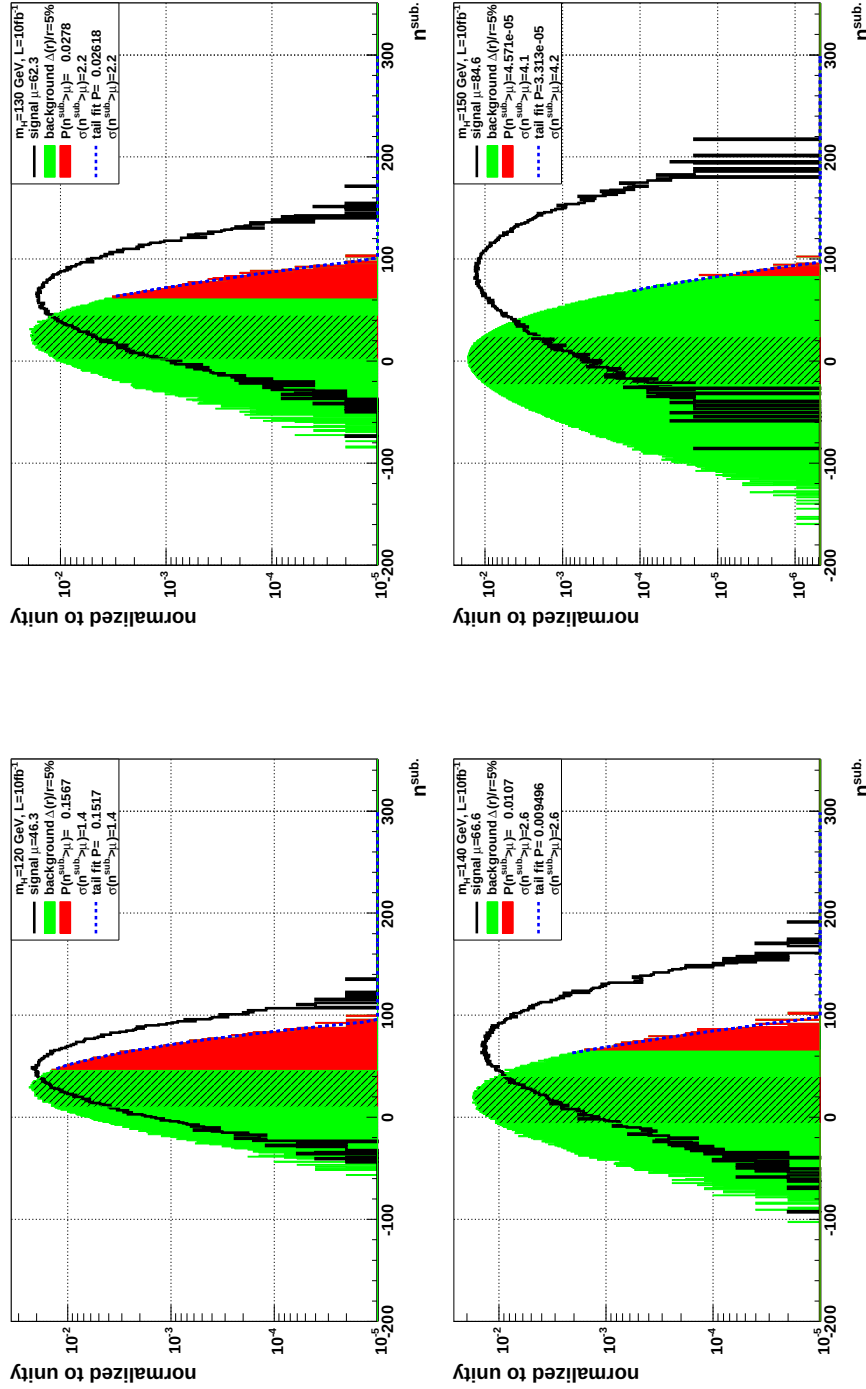


Figure 10.17: Distribution of $n^{\text{sub.}}$ over a large number of Monte Carlo experiments for the background only and for the signal plus background hypothesis for Higgs boson masses of 120, 130, 140 and 150 GeV. The signal+background distribution is shown in black. The bulk of the background only distribution is shown in green and the tail, that extends beyond the median signal+background expectation μ , is shown in red. The dashed blue curve is a Gaussian fit of the positive tail of the background only distribution. The black on green hashed area indicates the statistical 1σ fluctuation of the background.

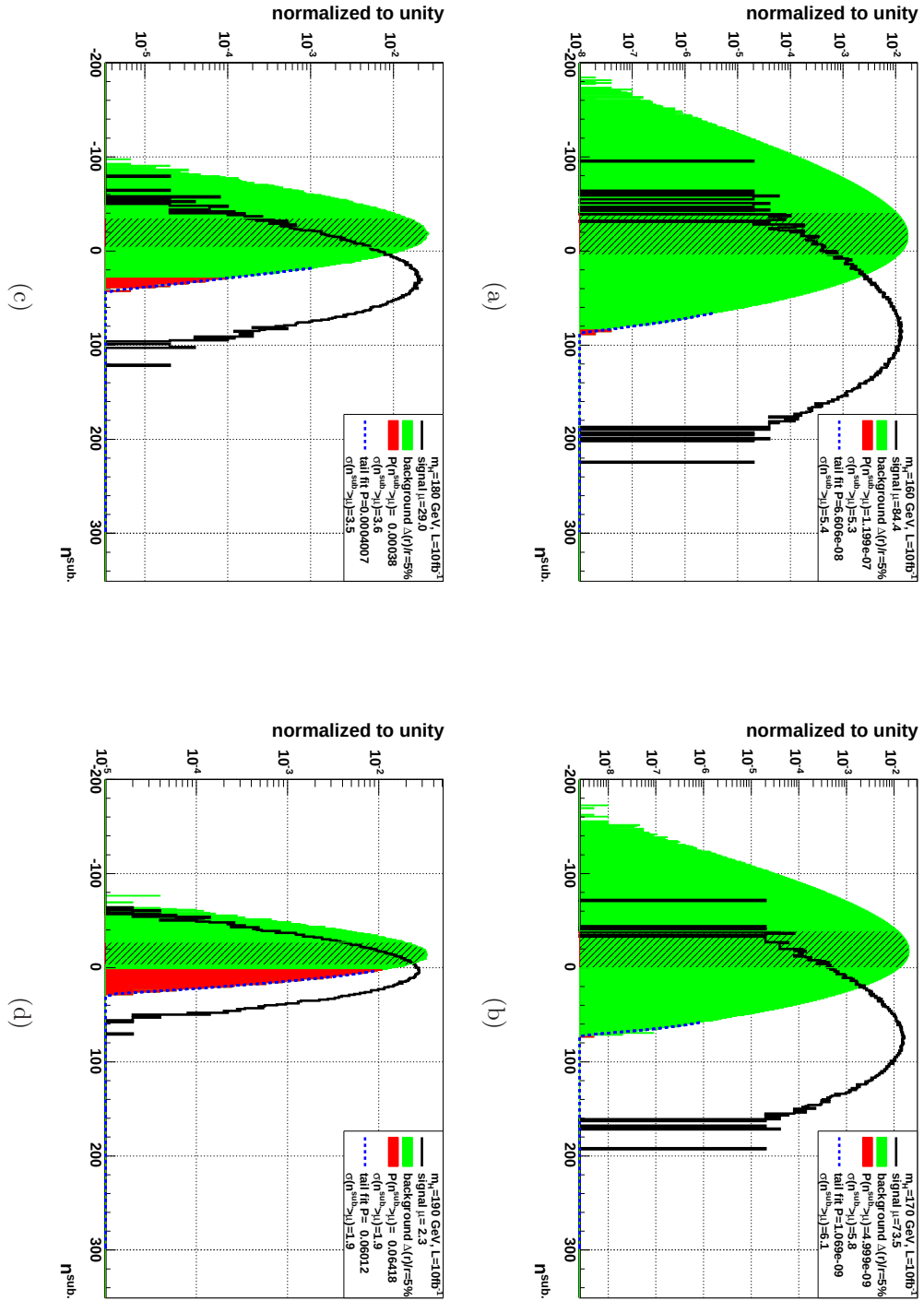


Figure 10.18: The same as Figure 10.17 for Higgs boson masses of 160, 170, 180 and 190 GeV.

boson mass for an integrated luminosity of 10 fb^{-1} . A significance of 5σ is reached for Higgs boson masses of 160 GeV and 170 GeV, where the Higgs boson branching ratio is maximal. Towards lower Higgs boson masses the significance decreases quickly due to the decreasing signal and increasing background rates. Towards higher Higgs boson masses, the significance decreases because the signal cross section decreases and the differences in the Higgs boson signal distributions between the signal and background regions disappear and only a small number of signal events remain after the background subtraction.

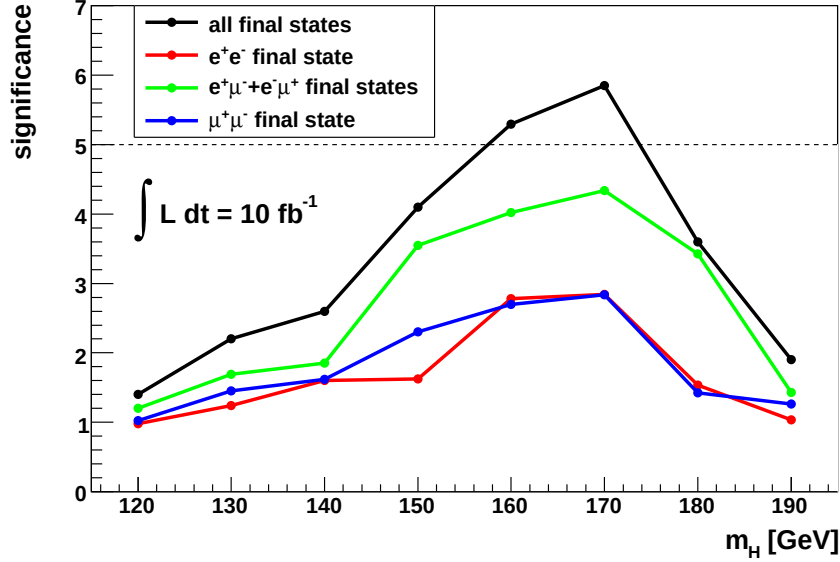


Figure 10.19: Higgs boson discovery significance in the $gg \rightarrow H \rightarrow W^-W^+ \rightarrow l^-\bar{\nu} \ l^+\nu$ final state as function of the Higgs boson mass for an integrated luminosity of 10 fb^{-1} , shown for all final states and separately for the e^+e^- , $e^\pm\mu^\mp$ and $\mu^+\mu^-$ final states.

The separation of the discovery potential into the three observable e^+e^- , $e^\pm\mu^\mp$ and $\mu^+\mu^-$ final states is also shown in Figure 10.19. The dominant part of the discovery potential is gained by the mixed $e^\pm\mu^\mp$ final state.

Although the quality of the measurement and subtraction of the background was shown for Monte Carlo simulations, no distributions or scale factors from the Monte Carlo simulation are included in the analysis method nor is the Monte Carlo information on the relative size of the different background contributions needed. Only in the final evaluation of the discovery significance Monte Carlo information is used, when the statistical and systematical fluctuations of the background only experiments are evaluated with the central value from the Monte Carlo prediction.

However, for Higgs boson masses $m_H \gtrsim 150 \text{ GeV}$, where a significant sensitivity

to the presence if a Higgs boson is found, the background only prediction of $n^{\text{sub.}}$ is negative due to the relative shift in the m_T distribution for the $q\bar{q} \rightarrow WW$ background. This shift is unaffected by the systematic variations of the $q\bar{q} \rightarrow WW$ background itself. If the relative size of either the top or the $gg \rightarrow WW$ background is increased, the effect of the shift is reduced and the expectation for $n^{\text{sub.}}$ approaches zero.

Furthermore, it can be expected that any additional non-resonant background source, that is not included in this analysis, shows a behaviour that is very similar to all other non-resonant backgrounds and would thus not alter r or the background subtraction method in a significant way. As the number of resonances in the standard model is limited, no additional relevant dilepton resonances are expected.

Hence, as a save baseline analysis for the sensitive mass range, a conservative background expectation of $n^{\text{sub.}} = 0$ can be used and no Monte Carlo information is needed at all¹⁰. For $m_H \approx 150$ GeV, where the m_T shift is not relevant, the discovery significance is almost unaffected by fixing $n^{\text{sub.}} = 0$ for the background only case. Above 150 GeV the significance is reduced by $\approx 1\sigma$.

10.5 Increasing the discovery potential

While an analysis without the need of Monte Carlo information is very appealing, a comparison shows that the discovery potential of the analysis presented in the previous section is approximately equivalent to the simple cut based analysis in Chapter 7 with a systematic background uncertainty of 10% to 15%, which is reachable with high probability before the needed integrated luminosity for a discovery with the analysis method described in this chapter is accumulated.

A dominant part of the additional statistical uncertainties of the analysis methods in this and the last chapter is caused by the measurement of the background distributions. While the Monte Carlo knowledge of these distributions might not be at the 1% level, Monte Carlo information can be used to considerably reduce the total statistical uncertainty.

An appropriate method for an analysis that includes Monte Carlo information for the signal and background processes is a combined fit of the transverse $\Delta\phi_{ll}$ and m_T and the longitudinal $|Y(l^+ + l^-)|$ and $|\eta(l^+) - \eta(l^-)|$ distributions for each process individually using the Monte Carlo predictions of the shapes. Because of uncertainties in the Monte Carlo knowledge of the various distributions, the challenge in such a fit is the definition of parametrised background functions that are:

- on the one hand sufficiently flexible to model systematic differences between

¹⁰This is of course only true if the contributions of the W +jet and $\gamma^{(*)}/Z^{(*)}$ backgrounds to high transverse masses is really negligible.

the Monte Carlo expectation and the real shape.

- on the other hand sufficiently constrained to significantly reduce the overall statistical uncertainty.
- unbiased by the Monte Carlo expectation and model correlations in distributions, as it is for example the case between $|\eta(l^+) - \eta(l^-)|$ and m_T for the dominant $q\bar{q} \rightarrow WW$ background, without enforcing a correlation that is not present in data or restricting an existing correlation, that is not present in the Monte Carlo simulation.

If these requirements can be satisfied, the combined fit is statistically more powerful than the simple cut based analysis, controls systematic uncertainties in a similar way as described in Chapter 9 and uses the full shape information of all distributions to maximise the discovery potential. However, the definition of such a combined fit and especially of shape functions for the backgrounds, including an evaluation of the systematic uncertainties and possible biases from the shape definition, is outside the scope of this thesis.

Conclusion

Monte Carlo simulations of physics events, including the detailed simulation of the detector, are an indispensable tool in the development and implementation of analysis strategies for high energy physics experiments. Despite massive world-wide computing resources, it is expected that only the equivalent of a few percent of the recorded ATLAS data can be simulated with the full ATLAS detector simulation, while most analyses require significantly more Monte Carlo events than data events for signal and important background processes.

In the course of this thesis the fast calorimeter simulation FastCaloSim was developed that is combined in the Atlfast-II detector simulation package with full simulations of the other detector components. Atlfast-II reduces the event simulation time by a factor of 10–20 without severely compromising the simulation quality of critical variables. During the last two years, Atlfast-II was validated by the ATLAS collaboration and used to simulate several hundred million events, many of which could never have been produced with the full detector simulation. It is expected that Atlfast-II will also play an important role in the ATLAS simulation strategy for the upcoming data taking era.

The inclusive Higgs boson $H \rightarrow W^-W^+ \rightarrow l^-\bar{\nu} l^+\nu$ decay is one of the most promising channels for an early discovery of a Higgs boson. However, due to the neutrinos from the W boson decays, a complete reconstruction of the final state is not possible. In the standard analyses for this channel, the presence of a signal can only be established with the help of Monte Carlo simulations for all background processes and hence the influence of systematic uncertainties from the use of Monte Carlo simulations limits the discovery potential severely. Especially the loop induced WW background process $gg \rightarrow WW$, that was first analysed within this thesis, is very similar to the Higgs boson signal and therefore needs to be known well to avoid a misinterpretation of background events as signal.

In order to avoid a strong dependency of a Higgs boson discovery on Monte Carlo predictions, two background normalisation methods were developed in this thesis. They allow to isolate and measure the two dominant processes in the analysis, the Higgs boson signal and the $q\bar{q} \rightarrow WW$ background, exclusively from recorded data without using Monte Carlo information. The method of the isolation of the Higgs boson signal was applied to the current Monte Carlo predictions of signal and background processes. It was shown that a discovery of a Standard Model Higgs boson is possible in the $H \rightarrow WW$ channel in the mass range $\approx 160 - 170$ GeV with data corresponding to an integrated luminosity of

10 fb^{-1} even without the use of any Monte Carlo information.

In the early phase of this PhD thesis it was analysed how well the nature of an observed particle as Higgs boson can be established after an initial discovery. This requires a measurement of the coupling strengths of the observed particle to other Standard Model particles, because it is this coupling strength that generates the particle masses in the Higgs theory. Due to limited observables at the LHC, only relative measurements of coupling parameters are possible without assumptions. However, using moderate theoretical assumptions, absolute measurements of the Higgs boson coupling strengths to the most important heavy Standard Model particles, W , Z and the top quark, are possible with accuracies of $\approx 50\%$ or better using data corresponding to an integrated luminosity of 30 fb^{-1} .

A Bijective transformation of histograms into Gaussian distributions

For the calculation of correlation matrices in Section 6.3.3 a transformation is needed, that transforms an arbitrary histogram into a normalised Gaussian distribution. The transformation should work in such a way that probabilities are conserved. For the mapping $x \rightarrow x'$ the integral of the normalised histogram in the interval $[-\infty, x]$ must be equal to the normalised Gaussian integral in $[-\infty, x']$.

In order to illustrate the transformation, the distribution of the first Tile calorimeter layer (TileBar0) from Figure 6.10 as shown in Figure A.1(a) is used as example. Furthermore it is assumed that a shower depth within the 4th bin of this histogram is chosen in the first simulation step (black box in Figure A.1(a)). The example value of an energy fraction $x = 0.29$, for which the Gauss transformation is explained, is marked by the small red box.

The distribution of the energy fraction in the 4th bin of the shower depth is shown in Figure A.1(b), the red line corresponds again to the example value $x = 0.29$. The mapping takes the following steps:

1. Figure A.1(b): The value to be mapped is $x = 0.29$.
2. Figure A.1(c): The normalised cumulative distribution of the energy fraction is built. The value $x = 0.29$ is mapped to $y = 0.425$.
3. Figure A.1(d): The y -axis of this cumulative distribution is - by construction - a uniform distribution between 0 and 1 with the mapped value $y = 0.425$.
4. Figure A.1(e): This uniform distribution can be projected to the y -axis of a cumulative Gaussian distribution. The value $y = 0.425$ is mapped to $x' = -0.188$.
5. Figure A.1(f): The Gauss projected value is $x' = -0.188$.

This is a reversible transformation scheme that assigns any value in an arbitrary source histogram a value in a Gaussian distribution and vice versa.

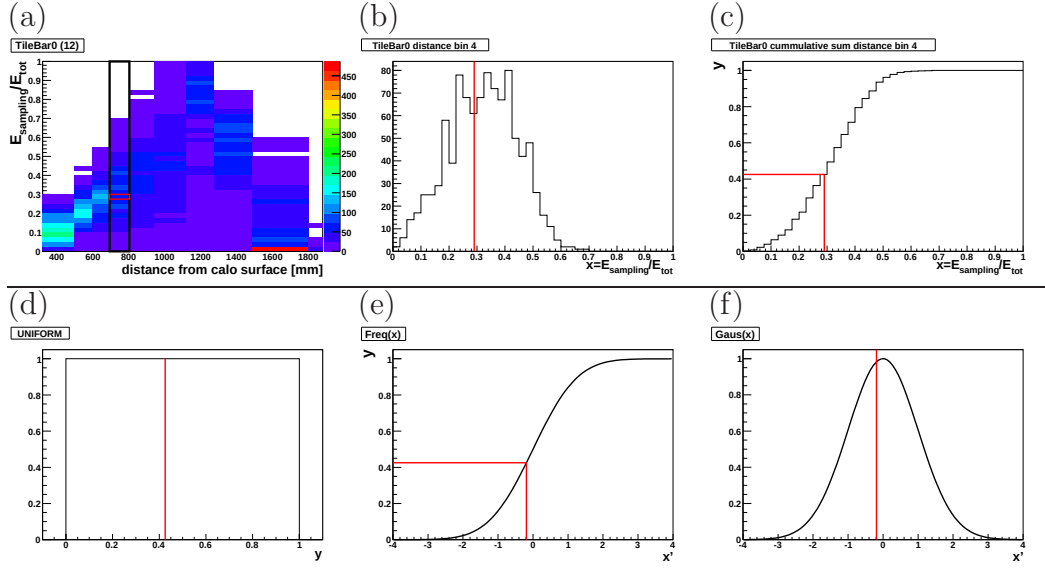


Figure A.1: Illustration of the Gauss transformation for each shower depth bin of each histogram in Figures 6.9 and 6.10. The transformation is demonstrated for a value of $x = 0.29$ (indicated by the red box in (a) and the red line in (b)-(f)) located in the 4th bin (indicated by the black box in (a)) of the shower depth distribution.

B Simulation and properties of the signal and background processes

This appendix gives a short overview of the Monte Carlo generators used for the production of the signal and background samples and gives details on important properties of these processes.

The choice of a suited Monte Carlo generator for a process depends on the selection cuts of the analysis and in particular on the phase space region that is selected. The most important selection cut for the $gg \rightarrow H \rightarrow WW$ analysis discussed in Chapters 8–10, as far as it concerns the choice of Monte Carlo generators, is the jet veto on $p_T(\text{jet}) < 20$ GeV. Due to this cut, Monte Carlo generators are chosen that give an accurate prediction of the inclusive process at NLO, while the correct treatment of multiple (hard) QCD radiation is less important. Currently only the MC@NLO event generator is able to produce hadron level events, which are needed for a detector simulation, with a NLO matrix element. Hence this generator is chosen wherever possible. MC@NLO uses the Herwig parton shower matched to the NLO matrix element. The leading p_T jet is generated with the matrix element, subleading jets are produced with the parton shower.

A second important consideration is the availability of sufficiently large Monte Carlo samples produced with the full ATLAS simulation.

B.1 The Higgs boson signal process

$$gg \rightarrow H \rightarrow WW$$

The inclusive Higgs boson production process is available in MC@NLO and a comparison of MC@NLO to calculations with NNLO matrix elements and NNLL resummations in Reference [42] shows a very good agreement.

The decay of the Higgs boson in the MC@NLO signal process is performed by Herwig, which has an incomplete description of spin correlations in the $H \rightarrow W^-W^+ \rightarrow l^-\bar{\nu} l^+\nu$ decay in the officially available version. For the full simulation signal samples, an updated version of Herwig was used that includes the full spin correlations.

B.1.1 Treatment of spin correlations in the MC@NLO Higgs signal process

The Higgs boson signal samples in fast simulation were produced with the officially available version of Herwig. Hence, a correction of the spin correlation is necessary for these samples. For the following discussion three angles are needed:

- In the rest frame of the W^+ boson the angle θ_{l^+,W^+} between the positively charged lepton and the original direction of flight of the W^+ in the H rest frame (W^+ boost direction). For convenience we use $\cos(\theta)_{l^+,W^+}$.
- The same angle θ_{l^-,W^-} for the negatively charged lepton in the W^- rest frame.
- The opening angle $d\phi$ between the decay planes of the two W bosons in the rest frame of the Higgs boson H .

Figure B.1(a) shows the correlation between $\cos(\theta)_{l^+,W^+}$ and $\cos(\theta)_{l^-,W^-}$. This correlation is identical for Herwig, Pythia and a MadGraph $H \rightarrow WW$ matrix element. Figure B.1(b) shows the distribution of the opening angle $d\phi$ between the W boson decay planes. This angle is modelled with a flat phase space distribution in Herwig while the MadGraph matrix element (and identically Pythia) shows a strong tendency towards large plane opening angles. Finally, Figure B.1(c) demonstrates that there is a non trivial correlation between $\cos(\theta)_{l^+,W^+}$ and $d\phi$ in the MadGraph matrix element (and identically in Pythia), so a simple reweighting of the Herwig events with the correct distribution of $d\phi$ is not sufficient.

For the analysis each MC@NLO $H \rightarrow WW$ signal event is reweighted with the ratio of the MadGraph matrix element and an event weight calculated from the implementation of the partial Herwig spin correlations. After reweighting, the parton level MC@NLO $H \rightarrow WW$ decay distributions agree in all variables with Pythia and the MadGraph matrix element.

B.2 Non-resonant dilepton backgrounds

B.2.1 $q\bar{q} \rightarrow WW$

A MC@NLO matrix element is also available for the dominant $q\bar{q} \rightarrow WW$ background and again MC@NLO is the best suited generator. The full spin correlations between the W boson decay products are implemented in the version of MC@NLO used to produce the WW samples in full and fast simulation.

For the analysis of systematic uncertainties very large $q\bar{q} \rightarrow WW$ samples are needed ($\approx 0.5 \cdot 10^9$ events in total) in order to reduce the statistical Monte Carlo uncertainties below the level of the systematic uncertainties under study. Due to CPU time constraints, this is only possible with fast detector simulations.

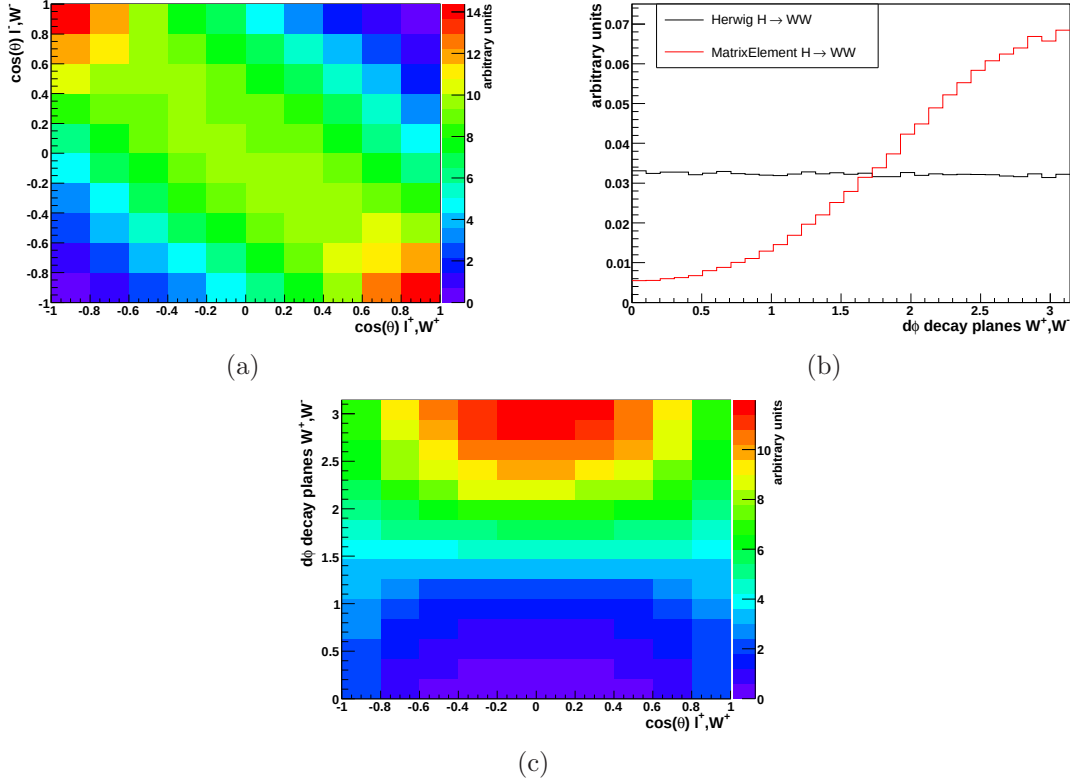


Figure B.1: Spin correlations in the $H \rightarrow WW$ decay.

- (a): Correlation between $\cos(\theta)_{l^+, W^+}$ and $\cos(\theta)_{l^-, W^-}$.
- (b): Comparison of the opening angle $d\phi$ between the W decay planes in the Higgs boson rest frame.
- (c): Correlation between the decay plane opening angle $d\phi$ and $\cos(\theta)_{l^+, W^+}$ in MadGraph/Pythia.

Figure B.2 shows a comparison of the critical shapes for the transverse $\Delta\phi_u$ and m_T distributions for the full and the fast Atfast-I and Atfast-II simulations. Within the statistical uncertainties the agreement to the full simulation is sufficiently good that the fast simulations can be used for the analysis of systematic uncertainties. The (large) differences in global efficiencies between the full and fast simulations are corrected by global scale factors.

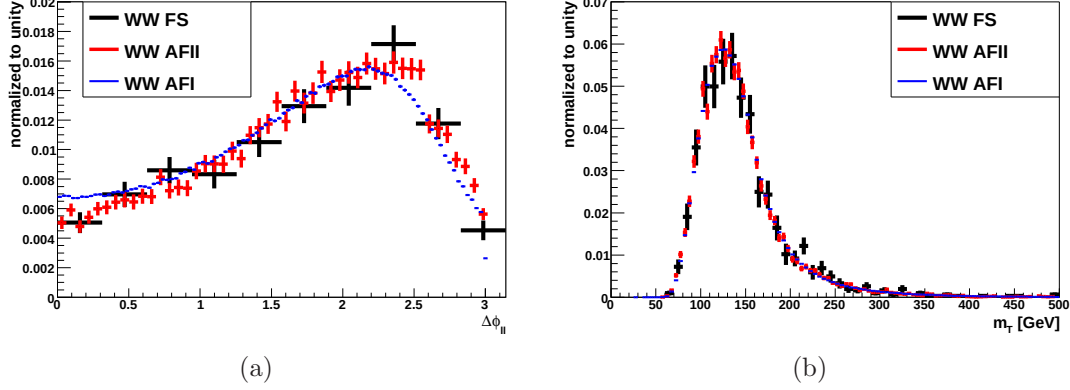


Figure B.2: Comparison of the Monte Carlo prediction for the shapes of $\Delta\phi_u$ (a) and m_T (b) for the $q\bar{q} \rightarrow WW$ process after selection cut 7 in Table 8.3. Compared are full simulation (FS), Atfast-II (AFII) and Atfast-I (AFI).

B.2.2 $q\bar{q} \rightarrow WZ$

The $q\bar{q} \rightarrow WZ$ background is small and no large uncertainties on this background are expected. A sufficiently large MC@NLO sample in full simulation is available and suitable for the analysis.

B.2.3 $gg \rightarrow WW$

For the $gg \rightarrow WW$ background only leading order matrix elements are available. This background and its influence and properties are discussed in Section 8.5.

B.2.4 $t\bar{t}$ and Wt

The simulation of the $t\bar{t}$ background to the $gg \rightarrow H \rightarrow WW$ analysis is particularly difficult due to the jet veto on 20 GeV jets. Without this jet veto, the resonant production of two top quarks is clearly the dominant process at the LHC for WW final states. However, after the jet veto, processes like Wt with less intrinsic jet activity and non-resonant contributions with softer jets are enhanced.

Some of the Feynman diagrams for the $t\bar{t}$ and Wt related processes are shown in Figure B.3. The diagram (a) shows the dominant $t\bar{t}$ process. The diagrams (b) and (c) are non-resonant contributions where successively one of the top quarks is replaced by the non-resonant production of Wb . The b quarks from the non-resonant contributions are softer and hence the relative importance of these effects is enhanced by the jet veto [39].

For very soft b quarks in the final state it is more appropriate to include the $g \rightarrow b\bar{b}$ splitting into the definition of a b quark proton PDF and consider the resonant and non-resonant Wt processes in the diagrams (d) and (e).

Finally, the extreme case in diagram (f), where both b quarks are not visible as external partons, is part of the inclusive $q\bar{q} \rightarrow WW$ production process.

Currently no Monte Carlo generator is available that is able to produce all top induced processes consistently, including non-resonant and NLO effects. Due to the jet veto, the analysis is sensitive to the difficult phase space region where it is not clear whether the simulation of a b quark as an external particle or the simulation of a process with a b quark in the initial state is more appropriate. The non-resonant contributions make this distinction even more difficult.

For the analysis two Monte Carlo generators were used: MC@NLO for the simulation of the top quark resonant $t\bar{t}$ process and AcerMC for the simulation of the top quark resonant Wt process. Because of the missing treatment of non-resonant contributions, a large systematic uncertainty should be assumed on the accepted cross section after the jet veto.

Full and fast simulation samples

The Monte Carlo statistics provided in full simulation for both samples is not sufficient for the purpose of a shape analysis. Figure B.4 shows the comparison of the transverse $\Delta\phi_{ll}$ and m_T distributions for both processes in full simulation with the fast simulation of $t\bar{t}$. For the Wt process no fast simulation sample with considerably larger statistics is available.

Due to the very limited statistics in the full simulation, it is hard to draw a conclusion from this comparison, apart from an agreement within huge statistical uncertainties. Therefore, Figure B.5 shows the same comparison where the jet veto threshold is increased to 30 GeV. The transverse distributions should not be affected much by this increase. The agreement of the $\Delta\phi_{ll}$ and m_T distributions of the fast $t\bar{t}$ sample with both top backgrounds in full simulation, $t\bar{t}$ and Wt , is sufficiently good to justify the use of the Atfast-I sample for the shape analysis.

Forward jet veto

The jet veto cut for 20 GeV jets is applied to all jets, even those detected in the forward region of the detector. As the forward calorimeters are in a very exposed environment at the LHC, it is currently not clear how well forward jets with a

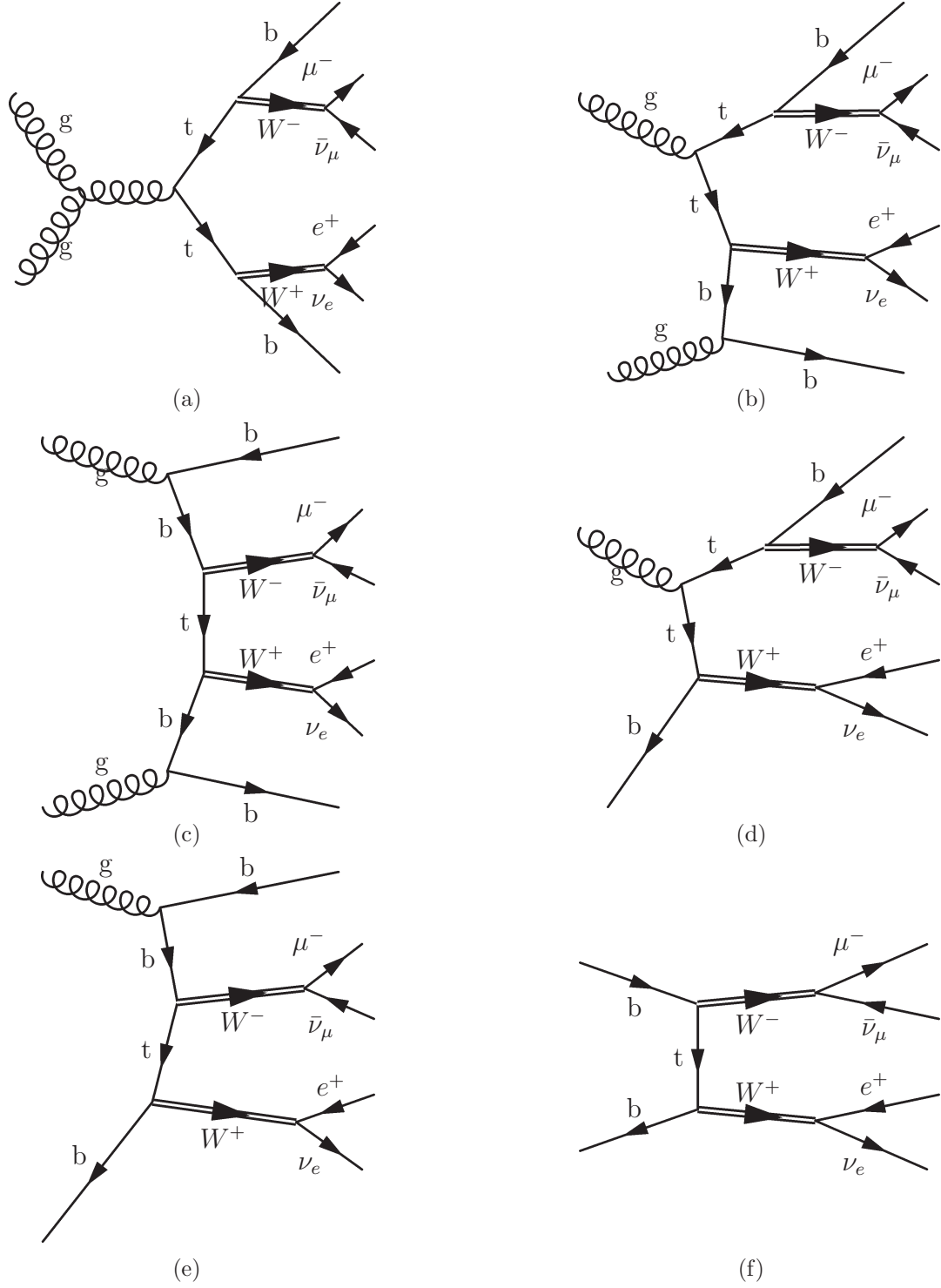


Figure B.3: Evolution of the resonant and non-resonant $t\bar{t}$ and Wt Feynman diagrams when external b -jets are vetoed down to the $b\bar{b} \rightarrow WW$ process.

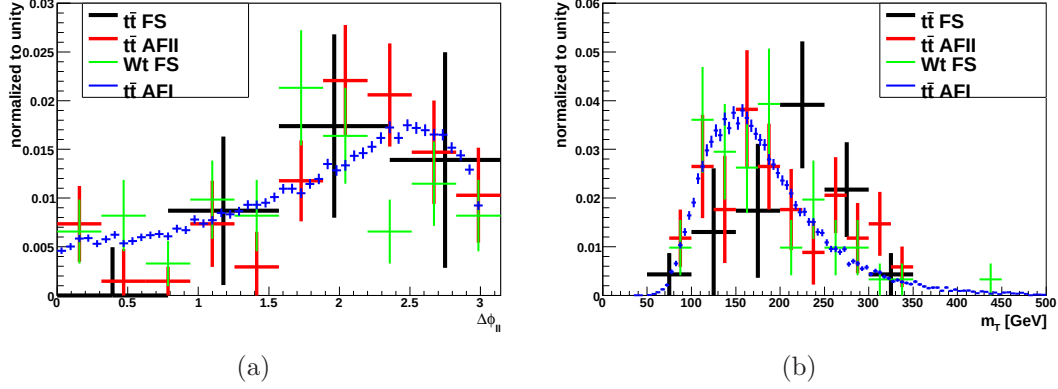


Figure B.4: Comparison of the Monte Carlo prediction for the shapes of $\Delta\phi_{ll}$ (a) and m_T (b) for the top processes after selection cut 7 in Table 8.3. Compared are full simulation (FS), Atlfast-II (AFII) and Atlfast-I (AFI).

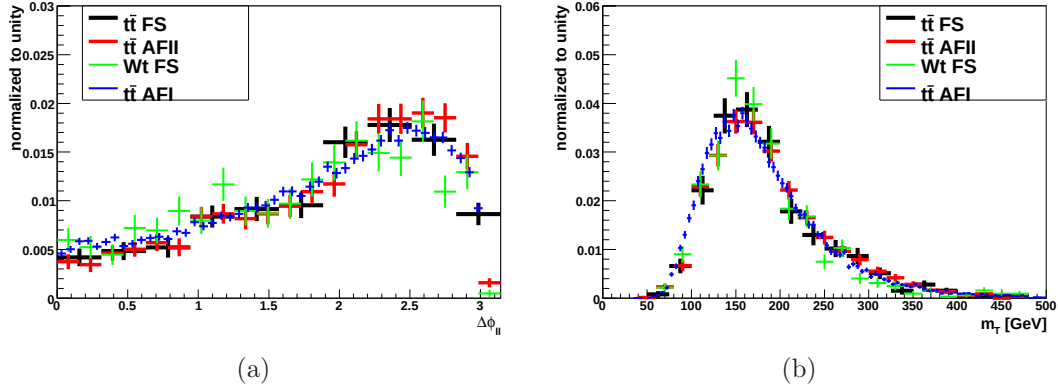


Figure B.5: The same comparison as in Figure B.4 but with a relaxed jet veto threshold of $p_T > 30$ GeV instead of the standard cut at $p_T > 20$ GeV.

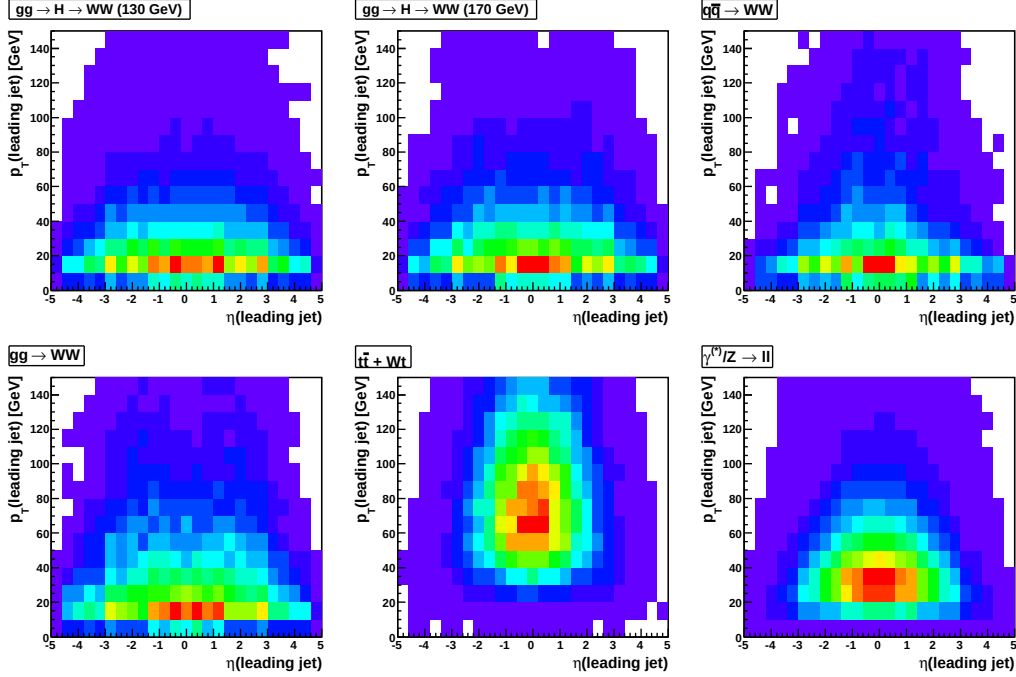


Figure B.6: Correlation of the leading jet transverse momentum and pseudorapidity after the cut on $p_T(l_1 + l_2) > 30$ GeV for the Higgs boson signal process ($m_H = 130$ GeV and $m_H = 170$ GeV) and the dominant dilepton backgrounds.

p_T as low as 20 GeV can be detected and how often underlying event and pile-up effects lead to falsely detected jets.

Figure B.6 shows the correlation of p_T and η for the leading jet. Especially for the $t\bar{t}$ and Wt background, the forward regions of the detector ($|\eta| > 3.2$) are hardly populated. While a restricted veto only on the leading jet within $|\eta| < 3.2$ will increase the top quark backgrounds, it will not change the overall picture in a significant way, if it turns out to be necessary to restrict the η range in which the jet veto is applied.

B.3 Resonant dilepton backgrounds

B.3.1 Drell Yan $\gamma^*/Z \rightarrow ee/\mu\mu$ and $Z \rightarrow ee/\mu\mu$

For the $q\bar{q} \rightarrow \gamma^*/Z$ background a NNLO prediction for the inclusive cross section exists and a NLO event generation is possible with MC@NLO. However, the biggest uncertainty for the $q\bar{q} \rightarrow \gamma^*/Z$ background is not from the cross section but comes from the efficiency of the E_T^{miss} selection cut. A Monte Carlo evalu-

ation of this efficiency is only possible with full simulation (and with Atfast-II with some additional uncertainty). Due to the large cross section and the low event selection efficiency, very large Monte Carlo samples are needed and only a Pythia sample is available in full simulation with sufficient statistics. The Pythia implementation of the process has a matrix element correction for the leading p_T jet and hence the choice of Pythia should not be a drawback for this analysis and the application efficiency of the jet veto.

The production of the full $\gamma^*/Z \rightarrow ee/\mu\mu$ process is split (arbitrarily) into two subsamples. The Drell-Yan sample is filtered at the generator level for invariant dilepton masses in the range $10 \text{ GeV} < m_{ll} < 60 \text{ GeV}$. The produced statistics is limited and one Monte Carlo event corresponds to a cross section of $\approx 1 \text{ fb}$. The Z sample is filtered at the generator level for invariant dilepton masses $m_{ll} > 60 \text{ GeV}$. One Monte Carlo event corresponds to a cross section of $\approx 0.2 \text{ fb}$.

One of the most effective selection cuts against the $q\bar{q} \rightarrow \gamma^*/Z$ background is the $E_T^{\text{miss}} > 30 \text{ GeV}$ cut. Figure B.7 shows the reconstruction resolution of the x component of E_T^{miss} before and after the cut on $E_T^{\text{miss}} > 30 \text{ GeV}$. For all processes with intrinsic E_T^{miss} due to neutrinos in the final state, the resolution is almost not affected. However, for the $q\bar{q} \rightarrow \gamma^*/Z$ background events from the tail of the E_T^{miss} resolution are selected. These resolution tails are caused by (non Gaussian) detector effects and it cannot be expected that even the full detector simulation gives an adequate description of the tails. Hence predictions for the accepted $q\bar{q} \rightarrow \gamma^*/Z$ cross section after the $E_T^{\text{miss}} > 30 \text{ GeV}$ cut have a very large systematic uncertainty and it will be necessary to measure the size of this background after all selection cuts from data.

B.3.2 $Z \rightarrow \tau\tau$

The most important difference between the $Z \rightarrow \tau\tau$ and the $Z \rightarrow ee/\mu\mu$ process is the presence of intrinsic E_T^{miss} from the τ decay neutrinos. However, due to the small $\tau\tau \rightarrow l\nu\nu l\nu\nu$ branching ratio and the low transverse momentum of the leptons, the accepted cross section after the lepton selection is only $\approx 1\%$ of the $\gamma^*/Z \rightarrow ee/\mu\mu$ cross section. The two leptons from the τ decays are preferentially back-to-back and the contribution of $Z \rightarrow \tau\tau$ can be neglected after the $\Delta\phi_{ll} < 1$ selection cut, as can be seen from the comparison of the $\Delta\phi_{ll}$ distribution for the $|\eta(l^+) - \eta(l^-)| < 1$ and $|\eta(l^+) - \eta(l^-)| > 1$ region in Figure B.8.

B.4 Lepton reconstruction and the rejection of the W +jets background

For the suppression of the large W +jets background tight lepton identification and isolation requirements are essential. The W +jets background appears as

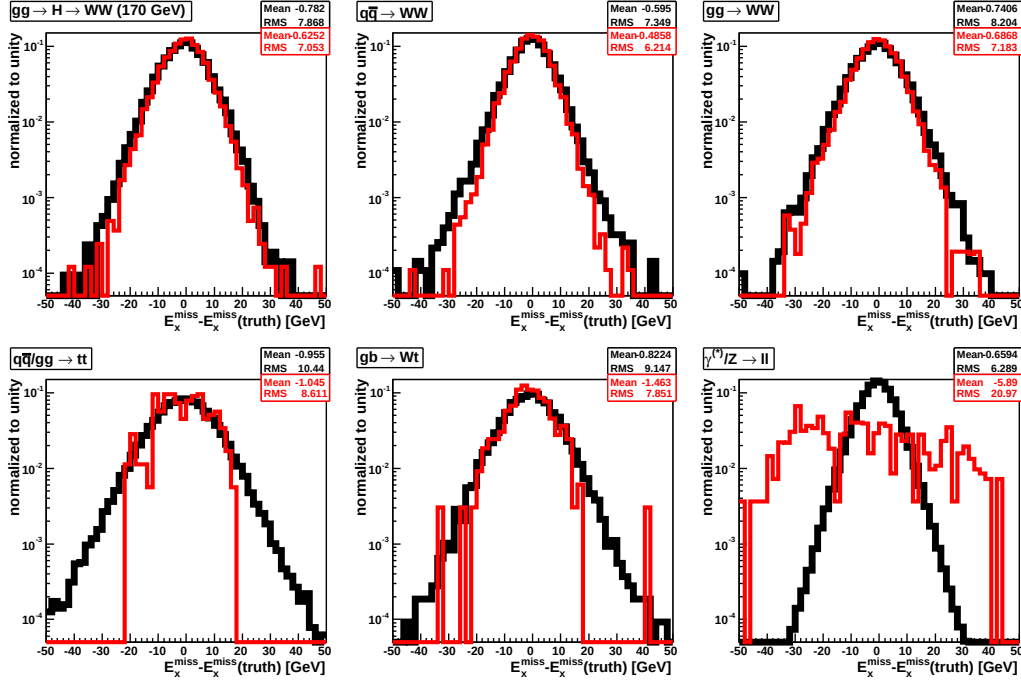


Figure B.7: Resolution of the x -component of the missing transverse energy vector. For each signal and background, the black distribution is after the lepton selection cuts, the red distribution after the cut on $p_T(l_1 + l_2) > 30$ GeV, jet veto and $E_T^{\text{miss}} > 30$ GeV.

background due to a jet from QCD initial state radiation being misidentified as a lepton.

B.4.1 Electron reconstruction

Electrons are reconstructed in the ATLAS detector from a combination of calorimeter and inner detector track measurements [15]. The best rejection against jets being misidentified as electrons is provided by the so called “tight” electron identification cuts. These cuts combine requirements on the measured electron shower shape in all electromagnetic calorimeter layers with requirements on the track reconstruction quality. The electron reconstruction and identification efficiency as function of the true electron p_T and η is shown in Figure B.9(a).

For the analysis of the shape of the W +jets background a relaxed electron identification is used. The “medium” electron identification cuts without calorimeter isolation use the same requirements on the calorimeter shower shape as the tight cuts except for the isolation requirements on the reconstructed calorimeter cluster that are applied as part of the tight electron identification procedure. In addition the track quality requirements are relaxed and especially no TRT re-

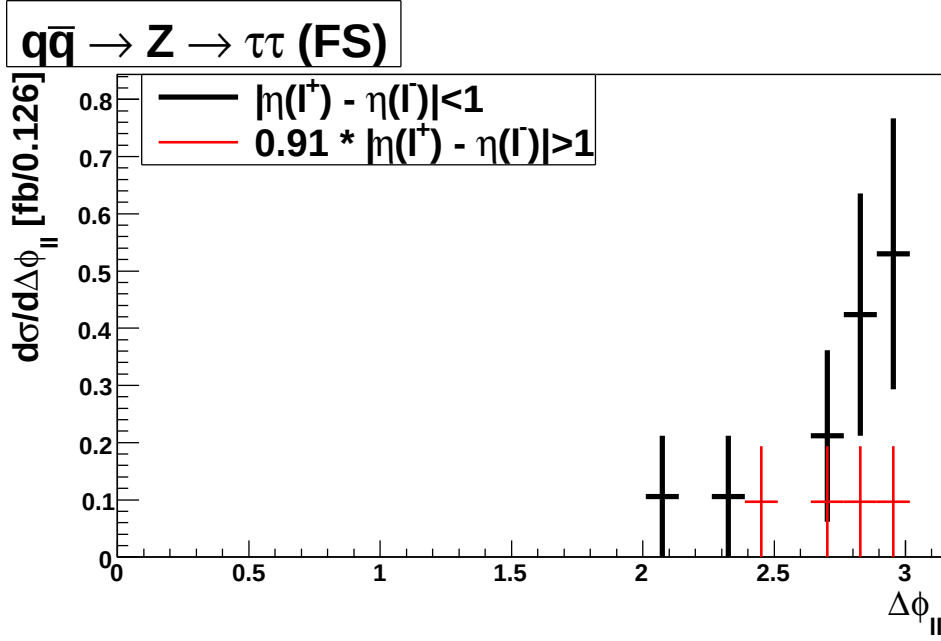


Figure B.8: The same as Figure 10.8 but for the $q\bar{q} \rightarrow Z \rightarrow \tau\tau$ process.

quirements are applied to distinguish between tracks from electrons and hadrons. This increases the jet to electron misidentification rate considerably.

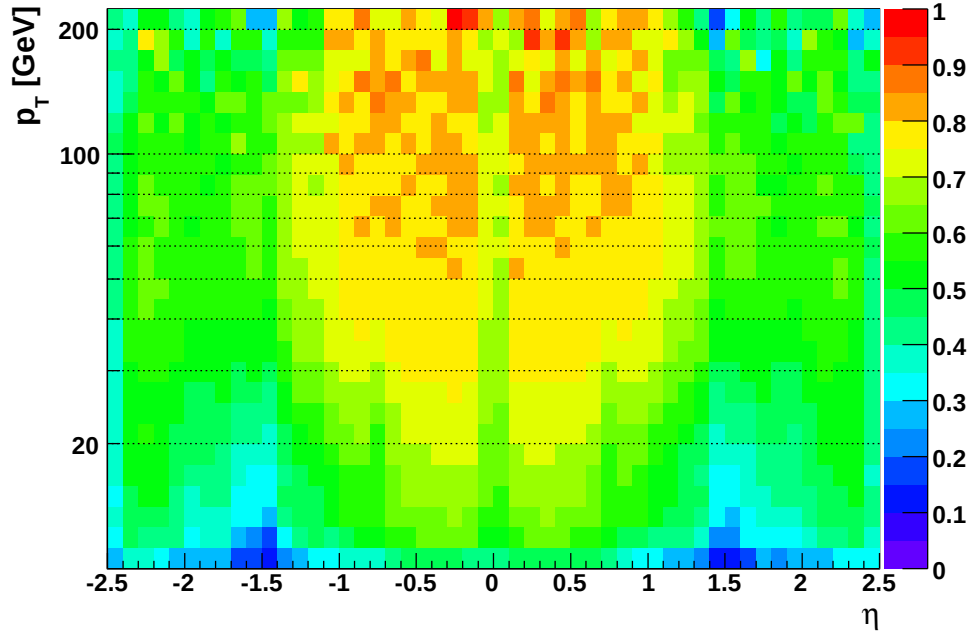
B.4.2 Muon reconstruction

Muons are reconstructed in the ATLAS detector from a combination of a reconstructed good quality track in the muon system with a good quality track in the inner detector [15]. If the χ^2 of the statistical matching of these two tracks is less than 100, the muon is accepted. The muon reconstruction and identification efficiency as function of the true muon p_T and η is shown in Figure B.9(b).

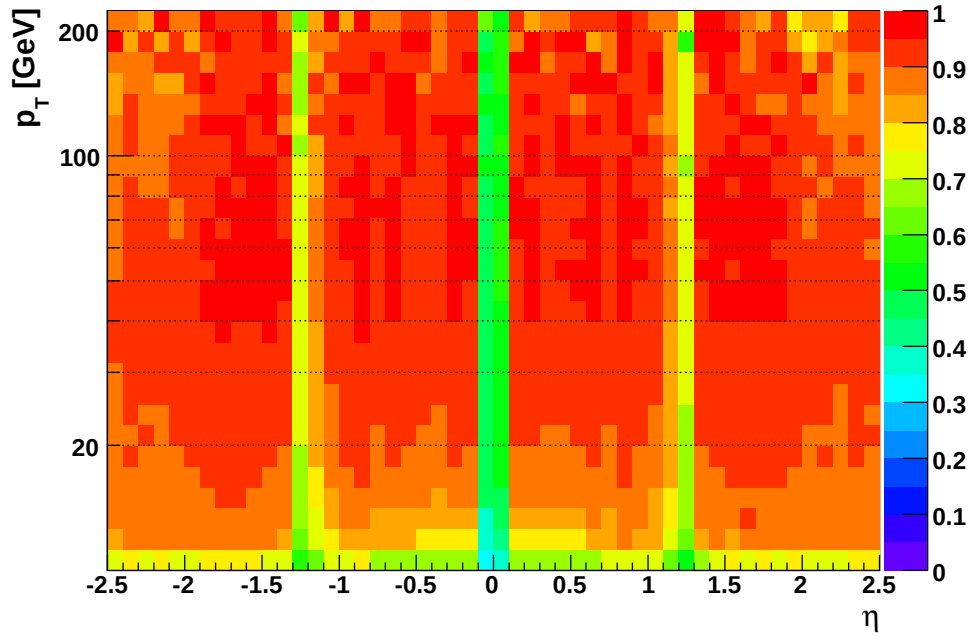
B.4.3 Lepton isolation

Most misidentified leptons are caused by rare developments inside a jet. Therefore, additional event activity in the form of either additional energy in the calorimeter or additional tracks can be expected close to the misidentified lepton. This additional activity is measured using two primary isolation variables:

- The transverse energy sum $\Sigma E_T^{\Delta R < 0.3}$ in the calorimeter in a cone of $\Delta R < 0.3$ around the lepton. The calorimeter energy of the lepton itself is subtracted from the sum.
- The scalar sum $\Sigma p_T^{\Delta R < 0.3}$ of the transverse track momenta of all tracks with



(a)



(b)

Figure B.9: Lepton identification and reconstruction efficiency as function of the true lepton p_T and η for electrons (a) and muons (b) in $q\bar{q} \rightarrow Z \rightarrow l\bar{l}$ events.

$p_T > 1$ GeV in a cone of $\Delta R < 0.3$ around the lepton. The track momentum of the lepton itself is subtracted from the sum.

Both isolation variables are used directly and additionally normalised to the transverse momentum of the lepton.

In the following leptons from W decays are compared to misidentified leptons. In the case of electrons, the electrons from a $W \rightarrow e\nu$ sample are compared to all electrons from a $W \rightarrow \mu\nu$ sample. For the purpose of the analysis, all electrons in the $W \rightarrow \mu\nu$ sample are defined as misidentified electrons regardless of their origin. Also shown are the misidentified electrons with relaxed identification cuts. For muons the role of the samples is reversed and muons from a $W \rightarrow \mu\nu$ sample are compared to muons from a $W \rightarrow e\nu$ sample.

The distribution of all four resulting isolation variables for electrons and muons is shown in Figures B.10 and B.11 for identified leptons as defined above.

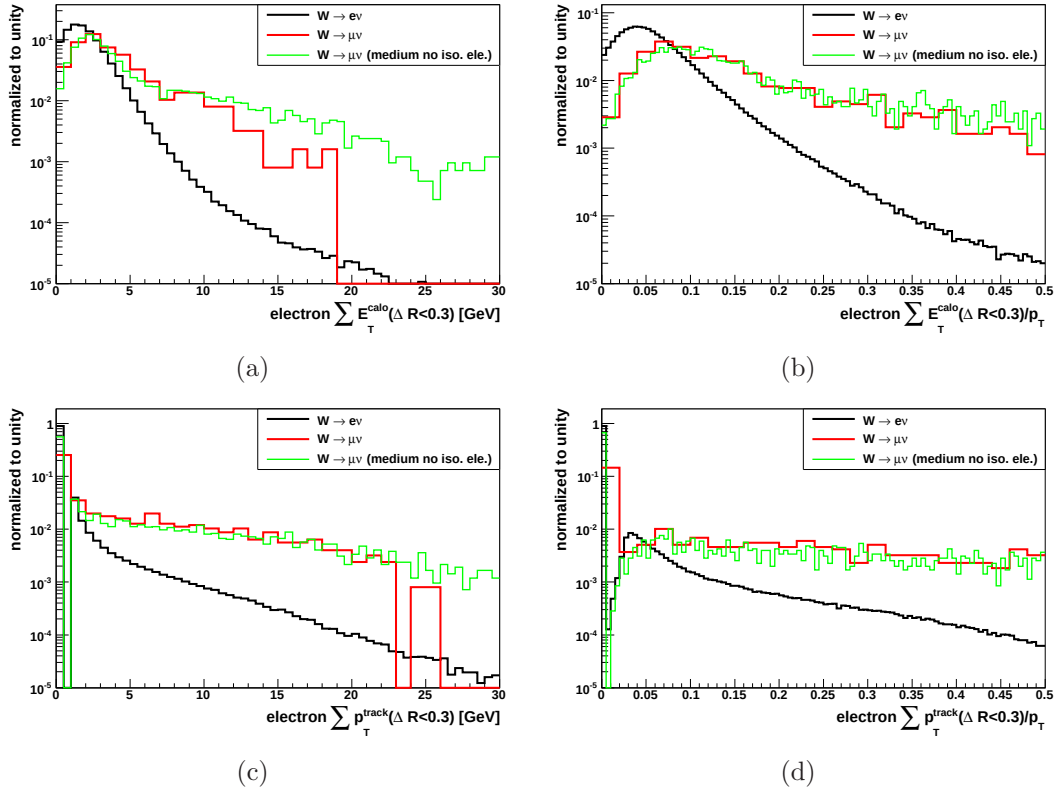


Figure B.10: Absolute and normalised electron calorimeter isolation energy sum and track isolation momentum sum in a cone of $\Delta R < 0.3$. Electrons from $W \rightarrow e\nu$ are real electrons, electrons from $W \rightarrow \mu\nu$ are caused by jets misidentified as electrons.

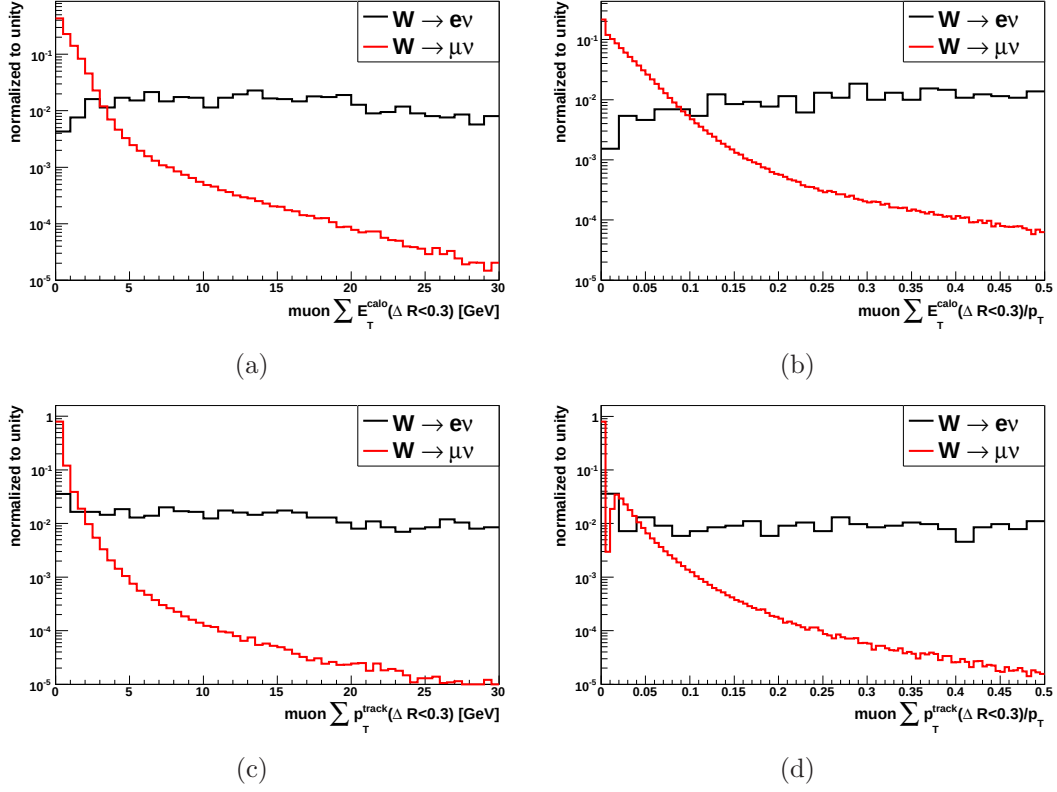


Figure B.11: Same as Figure B.10 for muons. Muons from $W \rightarrow \mu\nu$ are real muons, muons from $W \rightarrow e\nu$ are caused by jets misidentified as muons.

All isolation variables show a clear difference between leptons from W decays and misidentified leptons and can hence be used to considerably reduce the accepted W +jets cross section.

B.4.4 Displaced leptons

If (misidentified) leptons are produced in interactions in the inner detector or in the decay of particles with long lifetime, the measurement of the transverse impact parameter significance can be used to reject these leptons. The distribution of the transverse impact parameter significance is shown in Figure B.12 for electrons and muons. For misidentified leptons from jets originating from light quark flavours, as it is the case in the W +jets background, no strong difference to leptons from W decays is expected.

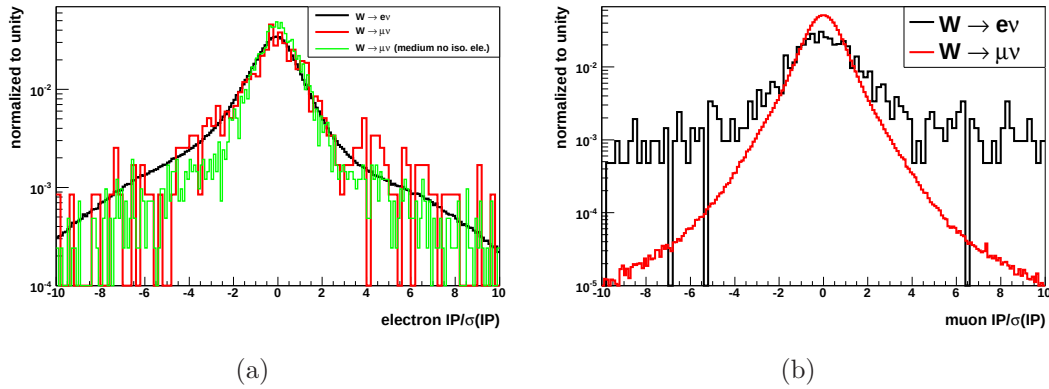


Figure B.12: Electron and muon impact parameter significance. Left: Electrons from $W \rightarrow e\nu$ are real electrons, electrons from $W \rightarrow \mu\nu$ are caused by jets misidentified as electrons. Right: Muons from $W \rightarrow \mu\nu$ are real muons, muons from $W \rightarrow e\nu$ are caused by jets misidentified as muons.

B.4.5 Lepton selection requirements

The lepton selection and isolation cuts listed in Table 8.2 are chosen to give a high rejection against misidentified leptons while keeping the lepton identification efficiencies at an acceptable level. Tighter isolation requirements would reduce the efficiency of selecting leptons from W decays severely, without reducing the misidentification rate significantly. The cut on the transverse impact parameter significance of 10 is chosen to accept almost all leptons from W decays and is intended as safety measure against leptons produced in decays or interactions in the inner detector.

Relaxed W +jets selection

The number of Monte Carlo events for the W +jets background is not sufficient for the analysis described in Chapter 10 and a fast simulation of this background is not possible, since the correct simulation of detector effects is crucial. Therefore, the lepton identification and isolation cuts are relaxed to allow an increased number of events to pass the event selection. The relaxed selection cuts are summarized in Table B.1.

Lepton flavour	electron	muon
Identification	medium(tight) electron selection without isolation (with iso.)	combined muon ID track match $\chi^2 < 100$
Kinematic cuts	$p_T > 15 \text{ GeV} ; \eta < 2.5$	
Calorimeter isolation	$\Sigma E_T^{\Delta R < 0.3} < 20(10) \text{ GeV} ; \Sigma E_T^{\Delta R < 0.3} / p_T(l) < 0.4(0.2)$ Sum over all calorimeter cells outside the core lepton	
Track isolation	$\Sigma p_T^{\Delta R < 0.3} < 20(10) \text{ GeV} ; \Sigma p_T^{\Delta R < 0.3} / p_T(l) < 0.2(0.1)$ Sum over all tracks with $p_T(\text{track}) > 1 \text{ GeV}$	
Track origin	$ \text{IP}/\sigma(\text{IP}) < 10$, IP=Impact parameter	

Table B.1: Summary of the relaxed electron and muon selection cuts. The standard selection requirements from Table 8.2 are given in brackets.

B.4.6 Kinematic distribution

The p_T and η distributions of the selected leptons, using the standard and the relaxed selection requirements, are shown in Figure B.13 for electrons and in Figure B.14 for muons. Misidentified leptons in general have a much softer p_T distribution than the leptons from W decays. For electrons, an accumulation in the crack region of the detector at $|\eta| \approx 1.3$ is observable.

However, it can be assumed that the Monte Carlo prediction for the jet to lepton misidentification efficiency is not accurate and a large systematic uncertainty should be assigned to the rate and distribution of misidentified leptons. Once data is available the misidentification efficiency needs to be measured and it can be expected that the lepton selection cuts must be modified to adjust for an underestimated misidentification rate, either globally or locally in some p_T or η regions.

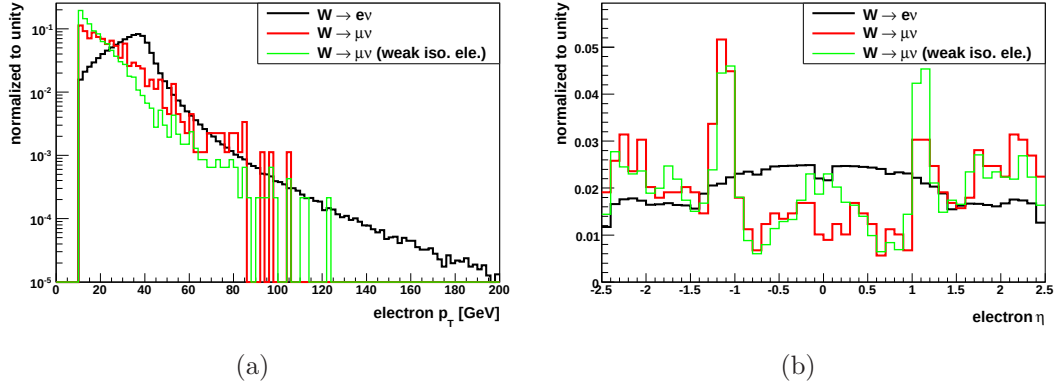


Figure B.13: Electron p_T and η distribution. Electrons from $W \rightarrow e\nu$ are real electrons, electrons from $W \rightarrow \mu\nu$ are caused by jets misidentified as electrons.

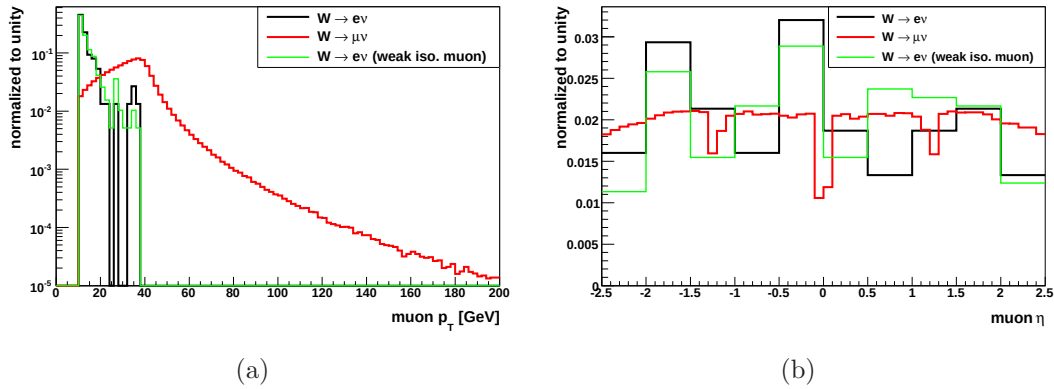


Figure B.14: Muon p_T and η distribution. Muons from $W \rightarrow \mu\nu$ are real muons, muons from $W \rightarrow e\nu$ are caused by jets misidentified as muons. Due to the low Monte Carlo statistics only a coarse binning for misidentified muons is possible.

C The definition of the transverse mass

The presence of two neutrinos in the final state allows for several definitions of a transverse mass, depending on the assumptions used to include information from the missing individual neutrino components. In the following, four transverse mass definitions are introduced and compared with regard to their properties for the signal and background processes studied in Chapters 8–10.

C.1 Variants of the transverse mass

C.1.1 The classical transverse mass

The classical transverse mass M_T^{classic} is defined in analogy to the transverse mass of the W boson [43, 44]:

$$(M_T^{\text{classic}})^2 = 2 \cdot |p_T^{\text{ll}}| \cdot E_T^{\text{miss}} \cdot [1 - \cos \angle(\vec{p}_T^{\text{ll}}, \vec{p}_T^{\text{miss}})] \quad (\text{C.1})$$

where $\vec{p}_T^{\text{ll}}$ is the vectorial sum of the transverse momentum vectors of the two leptons.

C.1.2 The transverse mass assuming a WW system at rest

Under the assumption that the W bosons decay at rest in the Higgs boson rest frame, the mass of the dineutrino system is identical to the mass of the dilepton system due to momentum conservation. This leads to the alternative definition of M_T^{WW} as introduced in Reference [45]:

$$(M_T^{WW})^2 = (E_T^{\text{ll}} + E_T^{\nu\nu})^2 - (\vec{p}_T^{\text{ll}} + \vec{p}_T^{\text{miss}})^2 \quad (\text{C.2})$$

$$E_T^{\text{ll}} = \sqrt{(p_T^{\text{ll}})^2 + m_{\text{ll}}^2} \quad (\text{C.3})$$

$$E_T^{\nu\nu} = \sqrt{(E_T^{\text{miss}})^2 + m_{\text{ll}}^2} \quad (\text{C.4})$$

C.1.3 The true transverse mass

The definition of M_T^{WW} has the disadvantage of no sharp mass edge close to the true mass of the WW system. Thus M_T^{true} is defined in Reference [46] to fulfil

the condition $M_T^{\text{true}} < m_{WW}$:

$$(M_T^{\text{true}})^2 = (E_T^{\text{ll}} + E_T^{\text{miss}})^2 - (\vec{p}_T^{\text{ll}} + \vec{p}_T^{\text{miss}})^2 \quad (\text{C.5})$$

C.1.4 The exclusive transverse mass

For this thesis a slightly modified version of the previously defined M_T^{true} is used, that is better suited to the background normalisation strategy used in Chapter 9, as it is using only transverse event properties. It can be calculated by replacing consistently m_{ll} in Equation (C.5) by the transverse dilepton mass m_T^{ll} . However, the simplest way of calculating this mass $M_T^{\text{true,T}}$ is to calculate the invariant mass of purely transverse Lorentz vectors:

$$\hat{p}^{l_1} = (p_T^{l_1}, p_x^{l_1}, p_y^{l_1}, 0)^T \quad (\text{C.6})$$

$$\hat{p}^{l_2} = (p_T^{l_2}, p_x^{l_2}, p_y^{l_2}, 0)^T \quad (\text{C.7})$$

$$\hat{p}^{\text{miss}} = (E_T^{\text{miss}}, p_x^{\text{miss}}, p_y^{\text{miss}}, 0)^T \quad (\text{C.8})$$

$$(M_T^{\text{true,T}})^2 = (\hat{p}^{l_1} + \hat{p}^{l_2} + \hat{p}^{\text{miss}})_\mu \cdot (\hat{p}^{l_1} + \hat{p}^{l_2} + \hat{p}^{\text{miss}})^\mu \quad (\text{C.9})$$

C.2 Properties of the transverse mass definitions

Figure C.1 shows a comparison of the four transverse mass definitions for the signal process at $m_H = 130$ GeV and $m_H = 170$ GeV as well as for the most relevant background processes. The comparison is performed after the selection cut on $\Delta\phi_{\text{ll}} < 1$, to show the distributions on which the transverse mass selection would be performed (see Section 8.1.3).

Figures C.2 - C.5 show the correlations between $\Delta\phi_{\text{ll}}$ and the transverse mass definitions for the Higgs boson signal process at $m_H = 130$ GeV and for the dominant $q\bar{q} \rightarrow WW$ background before the selection cut on $\Delta\phi_{\text{ll}} < 1$.

Finally, Figures C.6 - C.9 show comparisons of the shape of the transverse mass definitions for the two longitudinal regions $\Delta\eta_{\text{ll}} < 1$ and $\Delta\eta_{\text{ll}} > 1$ after the selection cut on $\Delta\phi_{\text{ll}} < 1$.

C.3 Advantages and disadvantages of the different transverse mass definitions

The following list of advantages and disadvantages for the four transverse mass definitions can be compiled from the figures of the previous section:

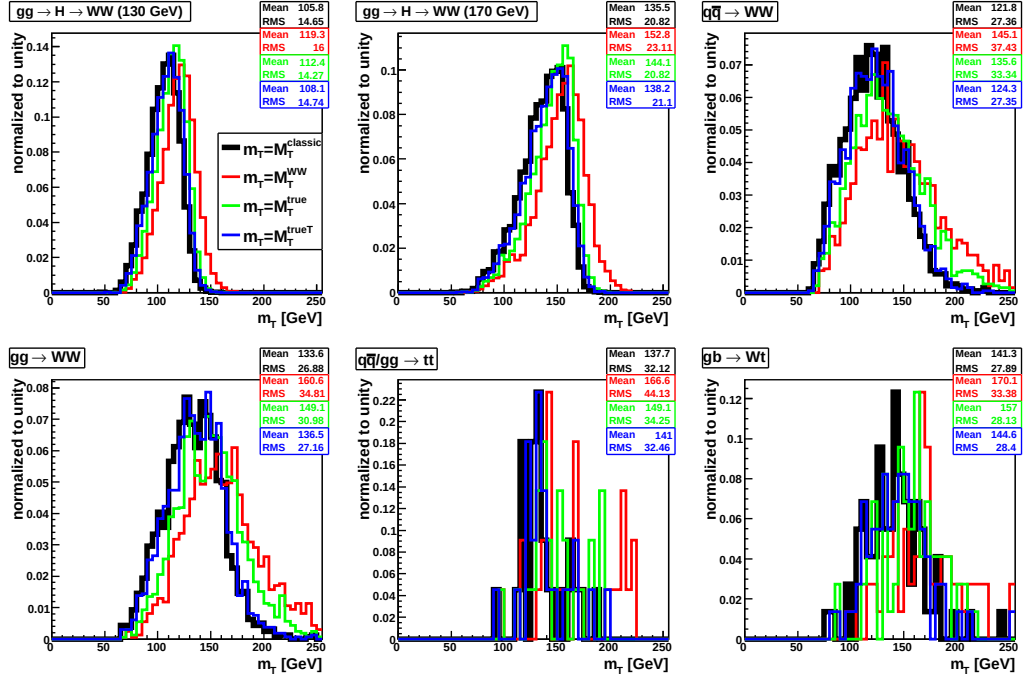


Figure C.1: Comparison of the transverse mass definitions for signal and background processes after the selection cut on $\Delta\phi_{\ell\ell} < 1$.

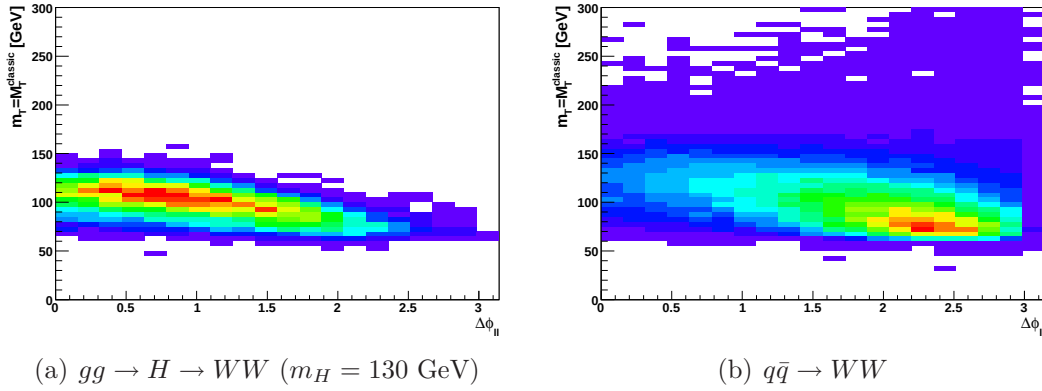


Figure C.2: Correlation of $\Delta\phi_{\ell\ell}$ and M_T^{classic} .

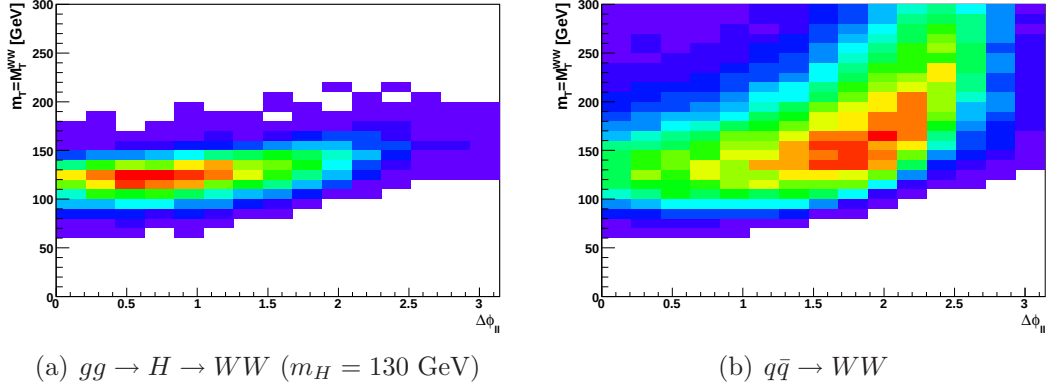


Figure C.3: Correlation of $\Delta\phi_{ll}$ and M_T^{WW} .

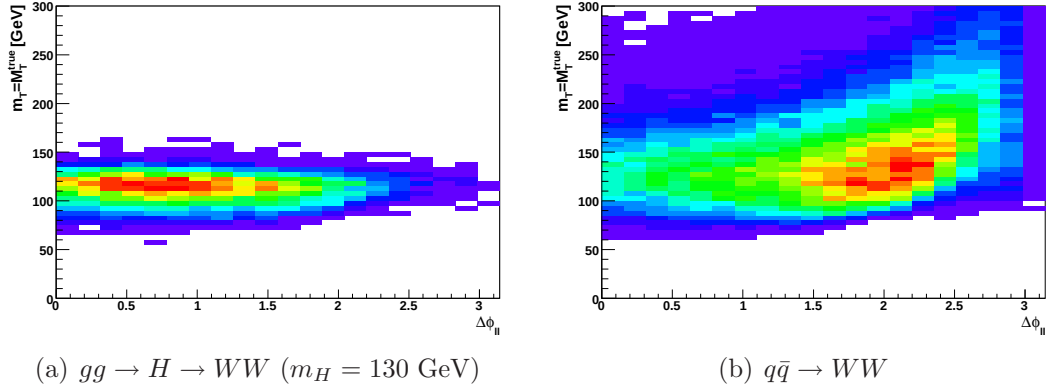


Figure C.4: Correlation of $\Delta\phi_{ll}$ and M_T^{true} .

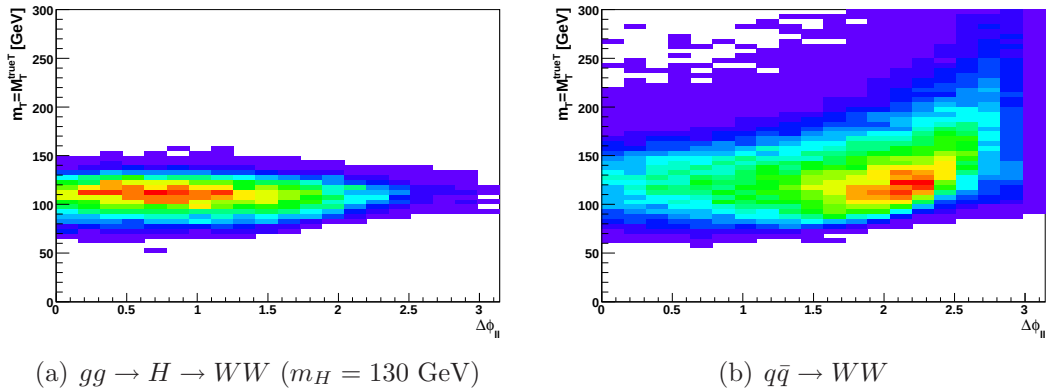


Figure C.5: Correlation of $\Delta\phi_{ll}$ and $M_T^{\text{true,T}}$.

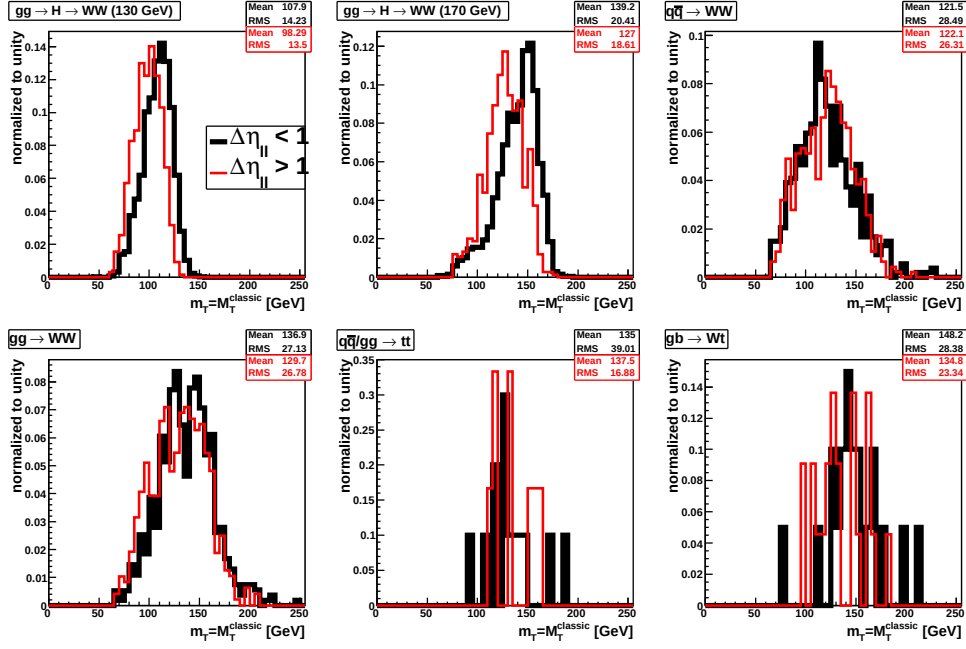


Figure C.6: Normalised distribution of M_T^{classic} for $\Delta\eta_u < 1$ and $\Delta\eta_u > 1$ after the selection cut on $\Delta\phi_u < 1$.

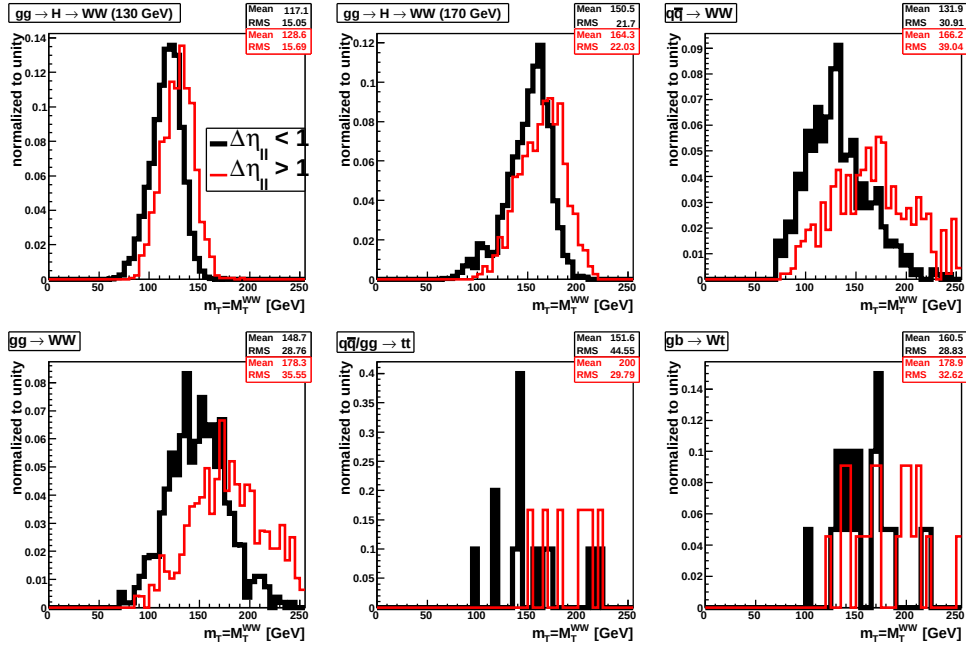


Figure C.7: Normalised distribution of M_T^{WW} for $\Delta\eta_u < 1$ and $\Delta\eta_u > 1$ after the selection cut on $\Delta\phi_u < 1$.

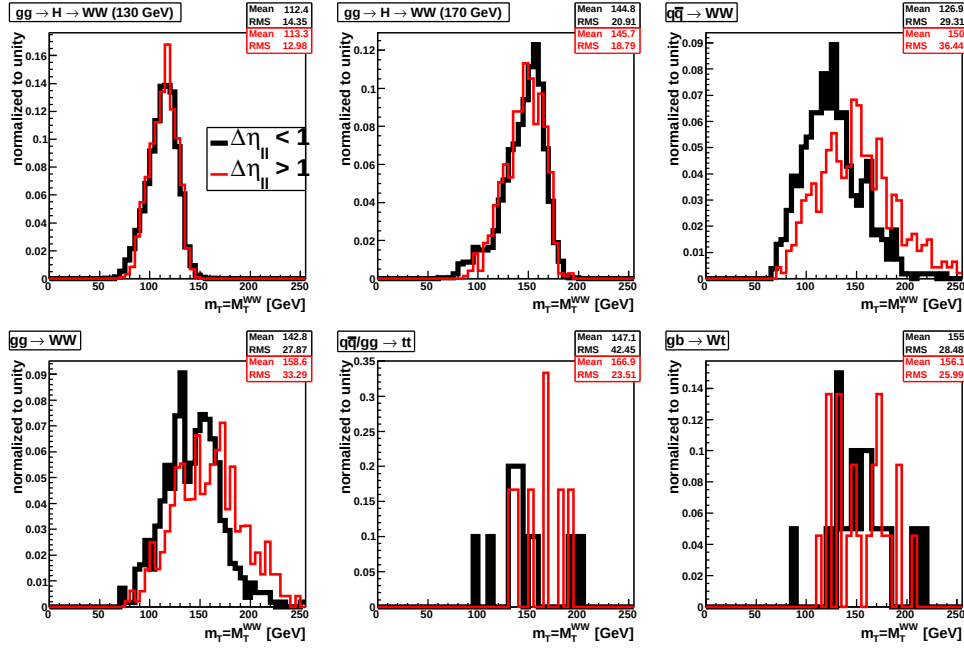


Figure C.8: Normalised distribution of M_T^{true} for $\Delta\eta_{||} < 1$ and $\Delta\eta_{||} > 1$ after the selection cut on $\Delta\phi_{||} < 1$.

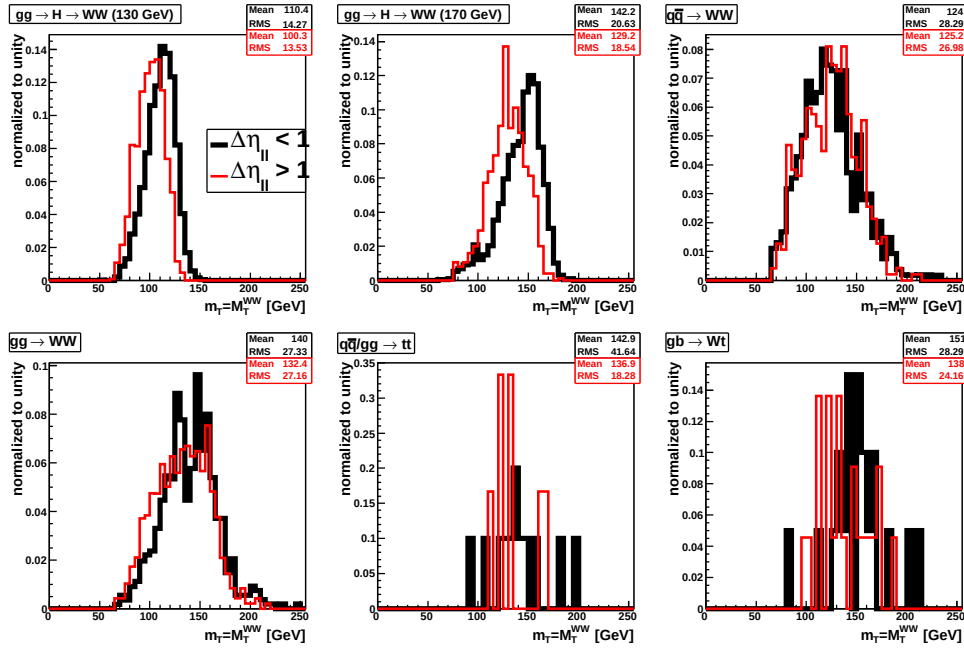


Figure C.9: Normalised distribution of $M_T^{\text{true},T}$ for $\Delta\eta_{||} < 1$ and $\Delta\eta_{||} > 1$ after the selection cut on $\Delta\phi_{||} < 1$.

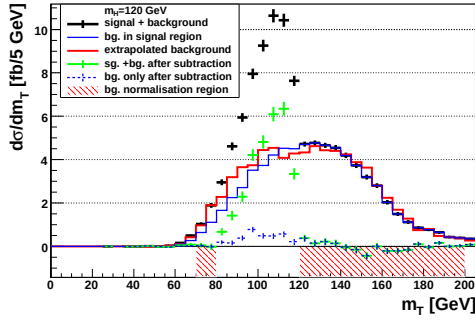
M_T^{classic}	+ Sharp mass edge below the Higgs boson mass m_H for the signal process – Strong correlation with $\Delta\phi_{ll}$ for the Higgs boson signal process and the $q\bar{q} \rightarrow WW$ background – Correlation with $\Delta\eta_{ll}$ for the Higgs boson signal process + No correlation with $\Delta\eta_{ll}$ for the background processes
M_T^{WW}	– No mass edge close to the Higgs boson mass m_H for the signal process – Strong correlation with $\Delta\phi_{ll}$ for the Higgs boson signal process and the $q\bar{q} \rightarrow WW$ background – Correlation with $\Delta\eta_{ll}$ for the Higgs boson signal process and for the background processes
M_T^{true}	+ Sharp mass edge below the Higgs boson mass m_H for the signal process + No correlation with $\Delta\phi_{ll}$ for the Higgs boson signal process + No correlation with $\Delta\phi_{ll}$ for the $q\bar{q} \rightarrow WW$ background for $\Delta\phi_{ll} \lesssim 1.5 - 2$ + No correlation with $\Delta\eta_{ll}$ for the Higgs boson signal process – Correlation with $\Delta\eta_{ll}$ for the background processes
$M_T^{\text{true,T}}$	+ Sharp mass edge below the Higgs boson mass m_H for the signal process + No correlation with $\Delta\phi_{ll}$ for the Higgs boson signal process + No correlation with $\Delta\phi_{ll}$ for the $q\bar{q} \rightarrow WW$ background for $\Delta\phi_{ll} \lesssim 1.5 - 2$ – Correlation with $\Delta\eta_{ll}$ for the Higgs boson signal process + No correlation with $\Delta\eta_{ll}$ for the background processes

For a good separation of the signal and the background, an asymmetric peak with a sharp mass edge is needed to get a clear distinction from the broad symmetric peak structure of the backgrounds. This requirement rules out M_T^{WW} , because of the more symmetric and wider peak structure.

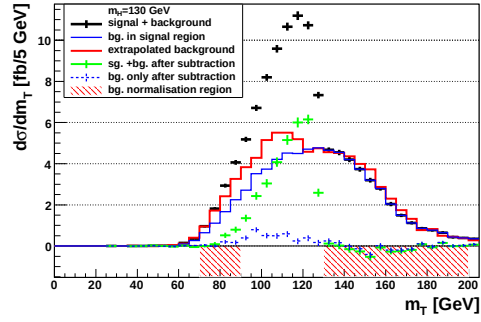
While there is no strict demand for uncorrelated $\Delta\phi_{ll}$ and transverse mass distributions, having these two variables uncorrelated allows to change the cut on $\Delta\phi_{ll}$ without influencing the shape of the transverse mass for the individual processes. Only the relative size of the signal and background contributions would be affected. In this case M_T^{classic} is ruled out as well.

The comparison of the remaining definitions, M_T^{true} and $M_T^{\text{true,T}}$, shows that M_T^{true} gives a narrower signal peak and a wider $q\bar{q} \rightarrow WW$ background distribution, as the signal peak is also uncorrelated to $\Delta\eta_{ll}$. Thus this definition gives the best signal to background ratio. However, the background normalisation technique introduced in Chapter 9 requires that there is no correlation between $\Delta\eta_{ll}$ and the transverse mass for the background processes. This is only the case for $M_T^{\text{true,T}}$ and hence the definition $m_T \equiv M_T^{\text{true,T}}$ is used.

D Event distributions

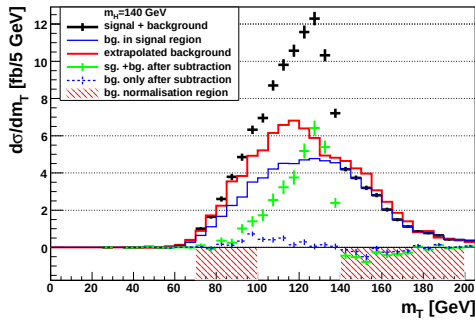


(a)

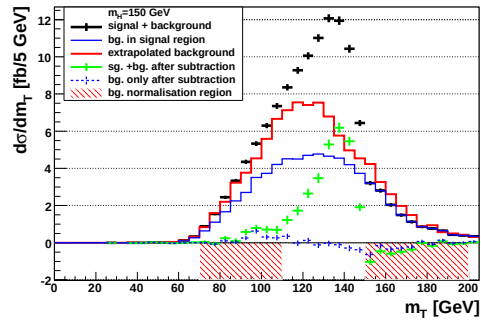


(b)

Figure D.1: Same as Figure 9.20 for Higgs boson masses of 120 GeV and 130 GeV.



(a)



(b)

Figure D.2: Same as Figure 9.20 for Higgs boson masses of 140 GeV and 150 GeV.

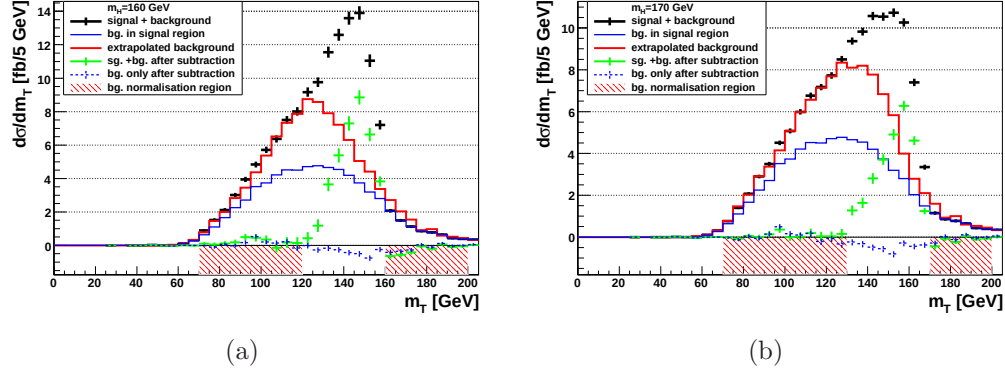


Figure D.3: Same as Figure 9.20 for Higgs boson masses of 160 GeV and 170 GeV.

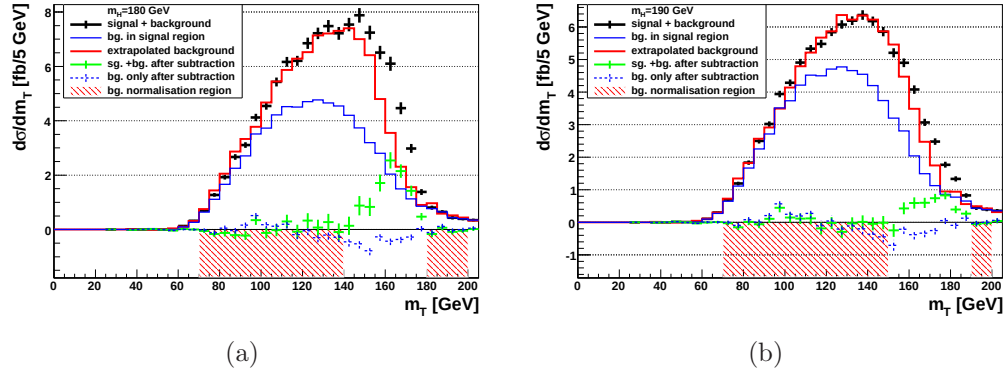


Figure D.4: Same as Figure 9.20 for Higgs boson masses of 180 GeV and 190 GeV.

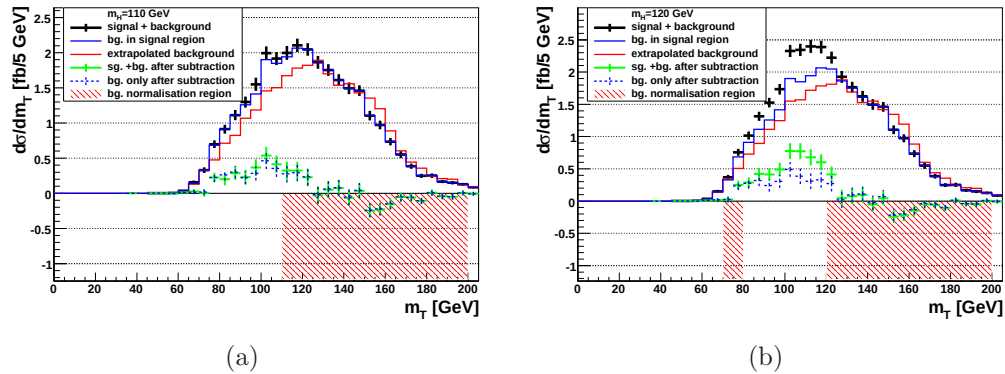
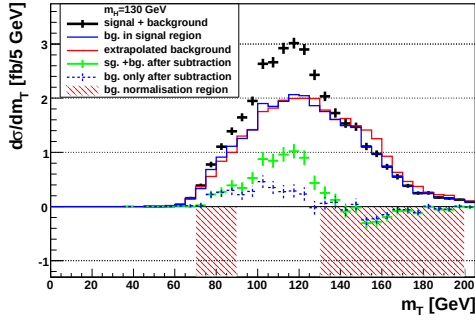
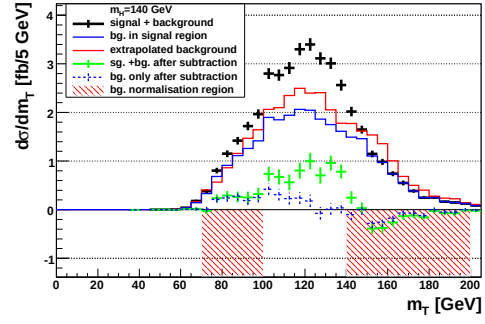


Figure D.5: Same as Figure 10.12 for Higgs boson masses of 110 GeV and 120 GeV.

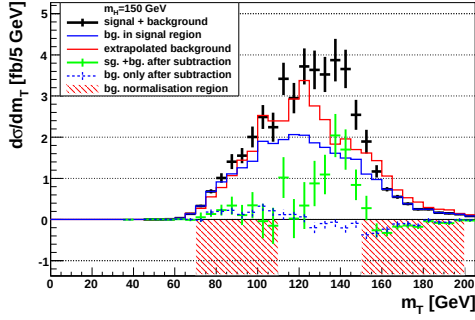


(a)

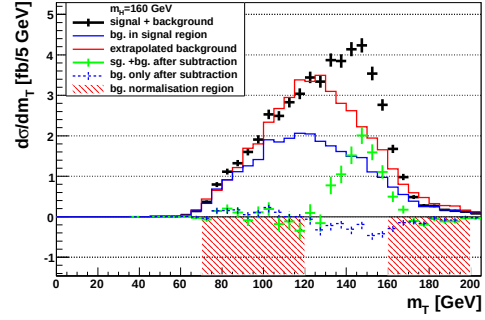


(b)

Figure D.6: Same as Figure 10.12 for Higgs boson masses of 130 GeV and 140 GeV.

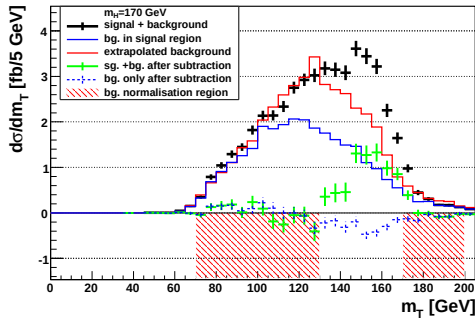


(a)

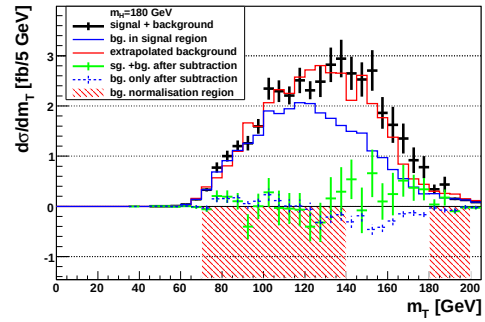


(b)

Figure D.7: Same as Figure 10.12 for Higgs boson masses of 150 GeV and 160 GeV.



(a)



(b)

Figure D.8: Same as Figure 10.12 for Higgs boson masses of 170 GeV and 180 GeV.

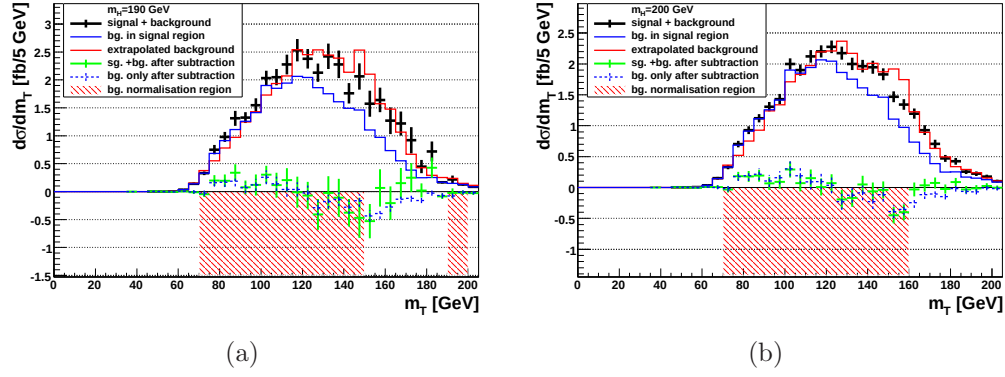


Figure D.9: Same as Figure 10.12 for Higgs boson masses of 190 GeV and 200 GeV.

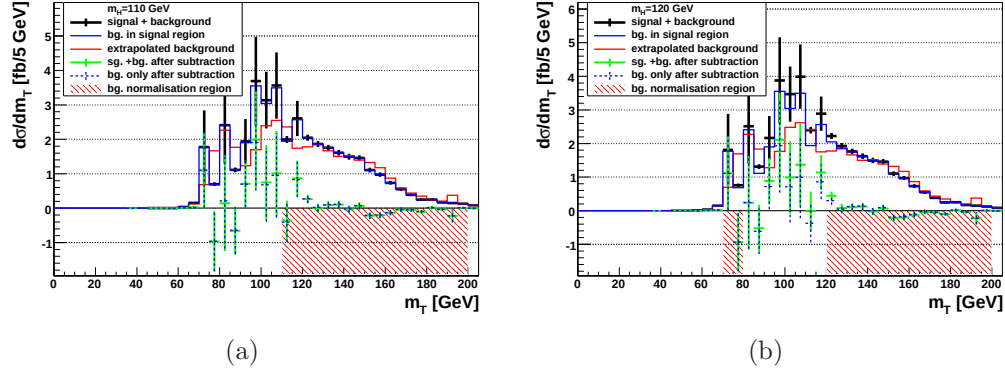


Figure D.10: Same as Figure 10.12 for Higgs boson masses of 110 GeV and 120 GeV, but all backgrounds are included.

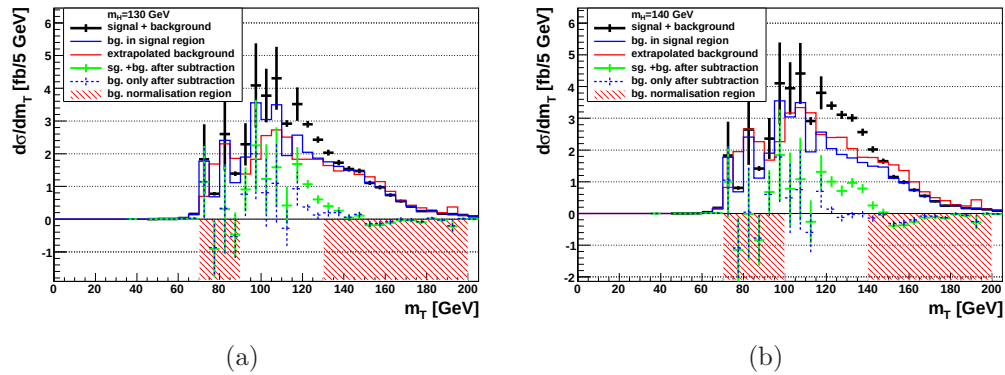
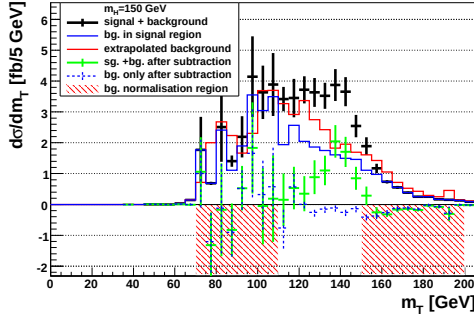
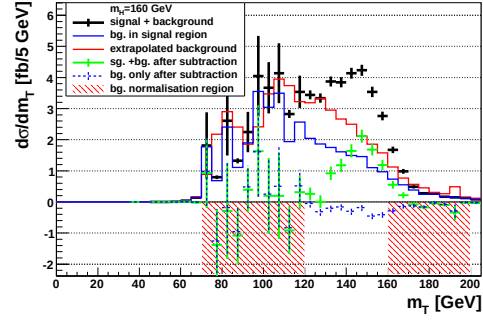


Figure D.11: Same as Figure 10.12 for Higgs boson masses of 130 GeV and 140 GeV, but all backgrounds are included.

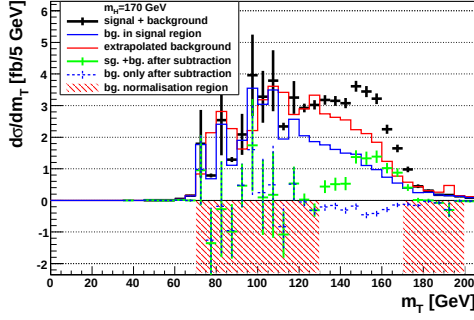


(a)

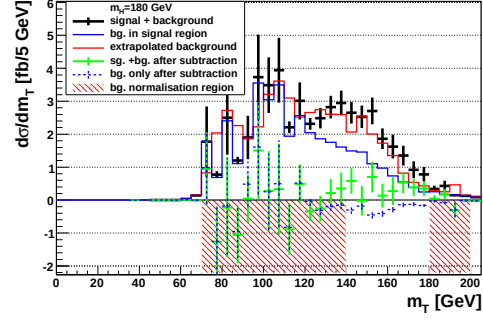


(b)

Figure D.12: Same as Figure 10.12 for Higgs boson masses of 150 GeV and 160 GeV, but all backgrounds are included.

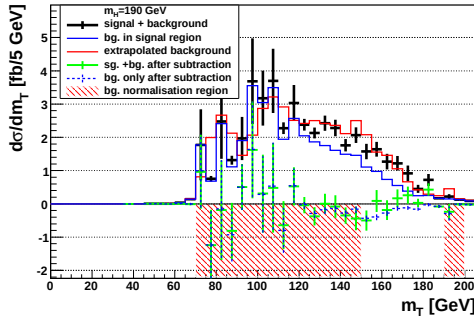


(a)

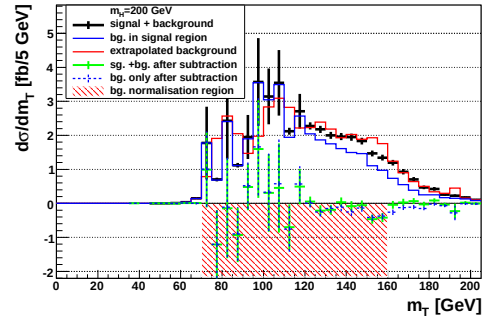


(b)

Figure D.13: Same as Figure 10.12 for Higgs boson masses of 170 GeV and 180 GeV, but all backgrounds are included.



(a)



(b)

Figure D.14: Same as Figure 10.12 for Higgs boson masses of 190 GeV and 200 GeV, but all backgrounds are included.

List of Tables

1.1	Table of the fundamental particles of the Standard Model	6
4.1	List of all ATLAS studies used as input for the combined fit of the Higgs boson coupling parameters	33
4.2	Overview of the systematic uncertainties assumed for the fit of relative Higgs boson couplings [18]. The uncertainties on W - and Z -fusion (α_{WF} and α_{ZF}) are treated as 100% correlated.	36
5.1	Design parameters of the LHC.	44
5.2	Position and coverage of inner detector silicon sub systems	49
5.3	Main parameters of the ATLAS calorimeter system	52
6.1	Generated particle energies and pseudorapidities used as input for the FastCaloSim parametrisation	67
8.1	Summary of the considered processes and their cross section at 10 TeV	132
8.2	Summary of the electron and muon selection cuts. Details are discussed in Appendix B.4.	134
8.3	Accepted cross section in fb for the simple $H \rightarrow WW \rightarrow l\nu l\nu$ analysis after the application of the successive selection cuts	145
8.4	The same as Table 8.3 but for the background control sample selection cuts as discussed in the text.	145
10.1	Extension of the accepted cross sections of Table 8.3 with the additional selection cuts of Chapter 10	189
10.2	Accepted cross sections in fb for the additional dilepton background contributions to the $H \rightarrow WW \rightarrow l\nu l\nu$ analysis.	189
10.3	Accepted cross section in fb for the W +jet background in the $H \rightarrow WW \rightarrow l\nu l\nu$ analysis	192
10.4	Expected number of events after the subtraction of the background region	210
B.1	Summary of the relaxed electron and muon selection cuts	XVIII

List of Figures

1.1	SM Higgs Potential	9
1.2	Higgs boson mass limit from electroweak precision measurements .	12
1.3	The triviality (upper) bound and the vacuum stability (lower) bound on the Higgs boson mass as a function of the new physics or cut-off scale Λ	13
2.1	Parton Distribution Functions	15
2.2	Factorisation in high p_T processes	16
2.3	Higher order corrections to Higgs production	17
2.4	Gluon fusion cross sections at the LHC	18
2.5	Event generation	20
3.1	Dominant Standard Model Higgs boson production modes at the LHC. The VBF production process is sometimes also called Vector Boson Fusion (VBF).	26
3.2	Higgs boson production cross sections at the LHC ($\sqrt{s} = 14\text{TeV}$) .	27
3.3	Experimentally accessible Standard Model Higgs boson decay modes at the LHC.	27
3.4	Calculated Standard Model Higgs boson branching ratios	28
3.5	Standard Model Higgs boson discovery potential of the ATLAS experiment with an integrated luminosity of 10 fb^{-1}	29
3.6	Standard Model Higgs boson discovery potential of the ATLAS experiment as function of the integrated luminosity	29
3.7	Observed and expected 95% CL upper limits on the ratios to the Standard Model cross section, as functions of the Higgs boson mass for the combined CDF and D0 analyses	30
4.1	Reconstructed Higgs boson mass resonance and mass resolution .	31
4.2	Expected accuracy of the ATLAS experiment for the determination of ratios of Higgs boson coupling parameters	35
4.3	Lower 1σ limit on the total Higgs boson width for the ATLAS experiment	38
4.4	Expected accuracy of two LHC experiments in the determination of absolute Higgs boson coupling parameters using an integrated luminosity of $2 \cdot 30 \text{ fb}^{-1}$	40

4.5	The same as Figure 4.4, but loop contributions from unknown particles to the gluon fusion production process and the $H \rightarrow \gamma\gamma$ decay mode are allowed.	41
5.1	Schematic view of the position of the four LHC experiments and the layout of the CERN accelerator complex	43
5.2	Cross sections and production rates for important processes at the LHC	45
5.3	Cut-away view of the ATLAS detector	47
5.4	Cut-away view of the ATLAS inner detector	48
5.5	Schematic view of a quarter-section of the ATLAS inner detector .	49
5.6	Material distribution (X_0, λ) at the exit of the inner detector envelope	51
5.7	Cut-away view of the ATLAS calorimeter system	51
5.8	Cumulative amounts of material in front of and in the electromagnetic calorimeters	53
5.9	Sketch of a electromagnetic barrel module	54
5.10	Schematic drawing of the segmentation in depth and η of the tile-calorimeter modules	55
5.11	Schematic R- ϕ (left) and R-z (right) views of the hadronic end-cap calorimeter	56
5.12	Cumulative amount of material, in units of hadronic interaction length λ , as a function of $ \eta $	57
5.13	Cut-away view of the ATLAS muon system	58
5.14	Block diagram of the ATLAS trigger and data acquisition systems	59
5.15	Schematic diagram of the ATLAS trigger system	60
6.1	Visualisation of the reconstruction geometry of the calorimeter system in the r-z plane	66
6.2	Total energy response of the calorimeter to photons and charged pions	69
6.3	Fitted calorimeter calibration constants as function of the particle energy for photons in the range $0.20 < \eta < 0.25$	70
6.4	The same as Figure 6.3 for pions.	71
6.5	Fitted calorimeter calibration constants as function of the pseudo-rapidity for photons of 200 GeV energy.	72
6.6	The same as Figure 6.5 for pions.	73
6.7	Reconstructed shower depth and correlation to the total calibrated energy response for photons	75
6.8	The same as Figure 6.7 for pions.	75
6.9	Energy response of the individual calorimeter layers to photons . .	76
6.10	The same as Figure 6.9 for pions.	77
6.11	Correlation matrix of the calorimeter energy fractions	78
6.12	Consistency test for the shower depth	80

6.13	Consistency test of the energy parametrisation for photons	81
6.14	The same as Figure 6.13 for pions.	81
6.15	Pseudorapidity intervals within each calorimeter layer for which a parametrisation of the radial shower profile was prepared for photons	84
6.16	The same as Figure 6.15 for pions.	84
6.17	Displacement in η and ϕ between the energy weighted centre of the shower in the 2 nd EM barrel layer and the extrapolated position of the true particle trajectory for photons	85
6.18	The same as Figure 6.17 for pions.	86
6.19	Energy distribution in the 2 nd EM barrel calorimeter around the expected point of impact for photons	87
6.20	Energy distribution in the 2 nd EM barrel calorimeter around the expected point of impact for pions	88
6.21	Illustration of the η asymmetry seen in the shower shape of a particle inside one calorimeter layer	89
6.22	Illustration of the EM calorimeter and energy deposition of muons in the EM calorimeter	91
6.23	Average ratio between the cell energies and the fitted energy shape for photons	93
6.24	The same as Figure 6.23 for pions.	94
6.25	Ratios between the cell energies and the fitted energy shape for photons	95
6.26	The same as Figure 6.25 for pions.	96
6.27	Schematic view of the implementation of the fast calorimeter sim- ulation FastCaloSim in the ATLAS reconstruction software.	98
7.1	Comparison between full and fast simulation for photons with an energy of 20 GeV and $0.2 < \eta < 0.25$	106
7.2	The same as Figure 7.1 for photons with an energy of 100 GeV and $2.0 < \eta < 2.05$ [30].	107
7.3	Comparison between full and fast simulation for charged pions with an energy of 20 GeV and $0.2 < \eta < 0.25$	110
7.4	The same as Figure 7.3 for charged pions with an energy of 100 GeV and $2.0 < \eta < 2.05$	110
7.5	The same as Figure 7.1 for charged pions with an energy of 20 GeV and $0.2 < \eta < 0.25$	111
7.6	The same as Figure 7.1 for charged pions with an energy of 100 GeV and $2.0 < \eta < 2.05$	112
7.7	Comparison of the jet EM energy scale between full simulation and FastCaloSim in four different pseudorapidity regions [30].	115
7.8	Comparison of the jet EM transverse energy resolution between full simulation and FastCaloSim in four different pseudorapidity regions [30].	116

7.9	Comparison of the jet reconstruction efficiency between full simulation and FastCaloSim in four different pseudorapidity regions [30].	117
7.10	Jet linearity (left) and jet resolution (right) as a function of the jet transverse energy in three pseudorapidity bins [30].	118
7.11	Jet linearity (left) and jet resolution (right) as a function of the jet pseudorapidity for three transverse energy bins [30].	118
7.12	The electron identification efficiencies as a function of p_T (left) and $ \eta $ (right)	120
7.13	The linearity $(E_T^{\text{rec}} - E_T^{\text{true}})/E_T^{\text{true}}$ for electrons from $Z \rightarrow ee$ events as function of p_T in the barrel region ($ \eta < 1.3$) (upper left) and endcap region (upper right) and as a function of $ \eta $ (lower plot) .	121
7.14	The energy resolution as function of p_T for electrons from $Z \rightarrow ee$ events for the barrel (left) and endcap (right)	122
7.15	Probability for a jet to fake an electron as a function of p_T and $ \eta $	123
7.16	The identification efficiencies for muons as a function of p_T for $ \eta < 2.5$ (left), and as a function of η for $p_T > 10$ GeV (right) . .	124
7.17	The relative muon momentum resolution $\Delta p_T/p_T$ for full simulation (black) and Atlfast-II (red) for the standard STACO combined muon selection for muons with $p_T > 10$ GeV from $Z \rightarrow \mu\mu$ decays.	124
7.18	The p_T distributions of misidentified muons in $t\bar{t}$ events for full simulation (black) and Atlfast-II (red) [31].	125
7.19	Efficiencies for the jet reconstruction as a function of p_T for different $ \eta $ regions	126
7.20	The mean values of the jet energy scale $p_T^{\text{rec}}/p_T^{\text{true}}$ as a function of p_T for light flavour jets from $t\bar{t}$ events for different $ \eta $ regions . .	127
7.21	The jet energy resolution as a function of p_T for the light flavor jets in $t\bar{t}$ events in different $ \eta $ regions	128
7.22	The distributions of $\Delta E_T = E_T(\text{rec}) - E_T(\text{true})$ for events from QCD jet production (left) and $W \rightarrow e\nu$ events (right)	129
8.1	The lepton p_T (a) and η (b) distributions after the lepton selection cuts for Higgs boson signal events ($m_H = 130$ GeV and $m_H = 170$ GeV) and the dominant dilepton background processes. . . .	133
8.2	Distribution of the missing transverse energy and the transverse momentum of the vectorial sum of the two leptons after the lepton selection cuts	135
8.3	Correlation between the missing transverse energy E_T^{miss} and the transverse momentum $p_T(l_1 + l_2)$ of the vectorial sum of the two leptons	136
8.4	Transverse momentum p_T distribution of the leading p_T jet before (a) and after (b) the cut on $p_T(l_1 + l_2) > 30$ GeV is applied for Higgs boson signal events ($m_H = 130$ GeV and $m_H = 170$ GeV) and the dominant dilepton background processes	137

8.5	(a) Distribution of the missing transverse energy E_T^{miss} after the jet veto cut for Higgs boson signal events ($m_H = 130$ GeV and $m_H = 170$ GeV) and the dominant dilepton background processes. (b) Distribution of the total sum of all contributions to m_{ll} after the cut on $E_T^{\text{miss}} > 30$ GeV for Higgs boson signal events ($m_H = 130$ GeV) and the dominant dilepton background processes. . . .	137
8.6	Feynman diagrams for the signal process (left) and the dominant irreducible WW background (right).	139
8.7	Distribution of the transverse opening angle $\Delta\phi_{ll}$ for Higgs boson signal events ($m_H = 130$ GeV) and the dominant dilepton background processes after all cuts against reducible backgrounds . . .	139
8.8	Distribution of the transverse mass m_T for Higgs boson signal events ($m_H = 130$ GeV) and the dominant dilepton background processes after all cuts against reducible backgrounds	141
8.9	Dependency of the transverse mass m_T on the transverse opening angle $\Delta\phi_{ll}$ for Higgs boson signal events ($m_H = 130$ GeV) and the dominant dilepton background processes after all cuts against reducible backgrounds	142
8.10	Distribution of the transverse mass m_T for Higgs boson signal events ($m_H = 130$ GeV) and the dominant dilepton background processes after the cut on $\Delta\phi_{ll} < 1$	142
8.11	Distribution of the transverse mass m_T for Higgs boson signal events ($m_H = 170$ GeV) and the dominant dilepton background processes after the cut on $\Delta\phi_{ll} < 1$	143
8.12	Dependency of the significance S on the systematic uncertainty of the background extrapolation factor $f_{n_C \rightarrow \hat{n}_B}$ for Higgs boson masses of $m_H = 130$ GeV and $m_H = 170$ GeV and for assumed integrated luminosities of 1 and 10 fb $^{-1}$	147
8.13	Normalised distribution of $\Delta\phi_{ll}$ after the lepton selection and a jet veto cut as function of the E_T^{miss} cut value for the dominant $q\bar{q} \rightarrow WW$ background	149
8.14	Systematic variation of the ratio $f_{n_C \rightarrow \hat{n}_B}$ as function of the E_T^{miss} cut value	150
8.15	Systematic uncertainty on the ratio of the number of events with $\Delta\phi_{ll} < 1$ to the number of events with $\Delta\phi_{ll} > 2$ after the lepton selection cuts (cut 1) due to PDF variations of the $q\bar{q} \rightarrow WW$ background using the 40 CTEQ6M 1σ error PDFs	152
8.16		152
8.17	Systematic uncertainty on the ratio $f_{n_C \rightarrow \hat{n}_B}$ due to PDF variations of the $q\bar{q} \rightarrow WW$ background after all event selection cuts	153
8.18	The same as Figure 8.17 but for QCD scale variations	153

8.19	Dependency of the accepted cross section after all cuts against reducible backgrounds on the jet energy scale for the dominant dilepton background contributions.	154
8.20	Systematic variation of the ratio $f_{n_C \rightarrow \hat{n}_B}$ for the sum of all backgrounds due to a variation of the size of the top backgrounds $t\bar{t}$ and Wt	156
8.21	Feynman diagrams for the dominant $q\bar{q} \rightarrow WW$ background and the $gg \rightarrow WW$ contribution	157
8.22	Correlation between $\Delta\phi_{ll}$ and m_T after all cuts against reducible backgrounds for Higgs boson signal events ($m_H = 170$ GeV) and the $gg \rightarrow WW$ and $q\bar{q} \rightarrow WW$ background processes	159
8.23	Systematic variation of the ratio $f_{n_C \rightarrow \hat{n}_B}$ for the sum of all backgrounds due to a variation of the cross section of the $gg \rightarrow WW$ background	160
9.1	Normalised distributions of the rapidity and of the rapidity difference of the WW system	162
9.2	Distribution of the transverse opening angle $\Delta\phi_{ll}$ in three bins of the longitudinal variable $ Y(W^+ + W^-) $ for Higgs boson signal events ($m_H = 130$ GeV) and the dominant non-resonant backgrounds.	164
9.3	The same as Figure 9.2 for the transverse mass m_T	164
9.4	Distribution of $\Delta\phi_{ll}$ in three bins of the longitudinal variable $ Y(W^+) - Y(W^-) $ for Higgs boson signal events ($m_H = 130$ GeV) and the dominant non-resonant backgrounds	166
9.5	The same as Figure 9.4 for the transverse mass m_T	166
9.6	Correlation between $Y(W^+ + W^-)$ and $Y(l^+ + l^-)$ for Higgs boson signal events ($m_H = 130$ GeV) and the dominant non-resonant backgrounds	167
9.7	Correlation between $Y(W^+) - Y(W^-)$ and $\eta(l^+) - \eta(l^-)$ for Higgs boson signal events ($m_H = 130$ GeV) and the dominant non-resonant backgrounds	167
9.8	Normalised distributions of $ Y(l^+ + l^-) $ and $ \eta(l^+) - \eta(l^-) $	169
9.9	Distribution of $\Delta\phi_{ll}$ in three bins of the longitudinal variable $ Y(l^+ + l^-) $ for Higgs boson signal events ($m_H = 130$ GeV) and the dominant non-resonant backgrounds after all cuts against reducible backgrounds (see Chapter 8) and an additional selection of events with $ \eta(l^+) - \eta(l^-) < 1$	170
9.10	The same as Figure 9.9 but for the transverse mass m_T	170

9.11	Distribution of $\Delta\phi_u$ in three bins of the longitudinal variable $ \eta(l^+) - \eta(l^-) $ for Higgs boson signal events ($m_H = 130$ GeV) and the dominant non-resonant backgrounds after all cuts against reducible backgrounds (see Chapter 8) and an additional selection of events with $ Y(l^+ + l^-) < 1$	171
9.12	The same as Figure 9.11 but for the transverse mass m_T	171
9.13	Comparison of the $\Delta\phi_u$ and m_T shape for the $q\bar{q} \rightarrow WW$ signal process and the gluon induced backgrounds	175
9.14	Comparison of the $\Delta\phi_u$ (a) and m_T (b) shape for the $q\bar{q} \rightarrow WW$ process for the regions $ Y(l^+ + l^-) < 1$ and $ Y(l^+ + l^-) > 1$ as well as the measured shape after subtraction of the gluon induced processes, including the Higgs boson signal ($m_T = 130$ GeV). . . .	176
9.15	The same as Figure 9.14, but without the Higgs boson signal contribution to the gluon induced processes.	177
9.16	Dependency of the background extrapolation ratio r between the number of events in X_S and X_B on the Higgs boson selection cuts	178
9.17	Comparison of the $\Delta\phi_u$ shape for the $gg \rightarrow H \rightarrow WW$ signal process and the dominant dilepton backgrounds	180
9.18	The same as Figure 9.17 for the m_T shape	180
9.19	Measured background extrapolation factor as function of the Higgs boson mass	182
9.20	Comparison of the true, measured and subtracted m_T distributions	182
9.21	Control distributions for the MadGraph Monte Carlo simulation of the $q\bar{q} \rightarrow WW$ process and the s- and t-channel components . .	184
9.22	Normalised distribution of the transverse mass m_T in three bins of $ \eta(l^+) - \eta(l^-) $ for the full WW process, the s- and t-channel components and the scalar sum of the two components.	185
10.1	Comparison of the leading jet p_T (a) and η (b) distribution between the regions $ \eta(l^+) - \eta(l^-) < 1$ and $ \eta(l^+) - \eta(l^-) > 1$	194
10.2	Comparison of the lepton p_T (a) and η (b) distribution between the regions $ \eta(l^+) - \eta(l^-) < 1$ and $ \eta(l^+) - \eta(l^-) > 1$ for the total sum of the signal and background processes.	195
10.3	Comparison of the m_T distribution in the $ \eta(l^+) - \eta(l^-) < 1$ signal region to the $ \eta(l^+) - \eta(l^-) > 1$ background region for the Higgs boson signal processes	196
10.4	The same as Figure 10.3 but for the $q\bar{q} \rightarrow WW$ process in full simulation (FS) and Atfast-II (AFII) and for the $gg \rightarrow WW$ process.	196
10.5	The same as Figure 10.3 but for the $t\bar{t}$ and Wt backgrounds in full simulation (FS) and the $t\bar{t}$ background with Atfast-I (AFI). . . .	197

10.6	The same as Figure 10.3 but for the $q\bar{q} \rightarrow WZ$ process and separately for the Drell-Yan $q\bar{q} \rightarrow \gamma^{(*)}/Z^{(*)}$ ($10 < m_{ll} < 60$ GeV) process and the Z boson production $q\bar{q} \rightarrow \gamma^{(*)}/Z$ ($m_{ll} > 60$ GeV) process	197
10.7	The same as Figure 10.3 but for the W +jet background process with standard selection cuts and with relaxed cuts scaled to agree with standard cuts after cut 7.	198
10.8	Comparison of the $\Delta\phi_{ll}$ distribution between the $ \eta(l^+) - \eta(l^-) < 1$ signal region and the $ \eta(l^+) - \eta(l^-) > 1$ background region for the Higgs boson signal processes with $m_H = 130$ GeV and $m_H = 170$ GeV before the selection cut 8a ($\Delta\phi_{ll} < 1$)	199
10.9	Comparison of the dilepton invariant mass m_{ll} distribution between the $ \eta(l^+) - \eta(l^-) < 1$ signal region and the $ \eta(l^+) - \eta(l^-) > 1$ background region	200
10.10	The same as Figure 10.3 but for the sum of all background processes and the sum of the Higgs boson signal ($m_H = 130$ GeV) and all background processes	201
10.11	Measured background extrapolation factor as function of the Higgs boson mass	202
10.12	Comparison of the true, measured and subtracted m_T distributions	203
10.13	Dependency of the significance S on the systematic uncertainty of the background extrapolation factor r for Higgs boson masses of $m_H = 130$ GeV and $m_H = 170$ GeV and assumed integrated luminosities of 10 and 30 fb ⁻¹	204
10.14	Systematic variation of the ratio $r^{\text{obs.}}$ after all event selection cuts as a function of the E_T^{miss} cut value	206
10.15	Systematic variation of the ratio $r^{\text{obs.}}$ after all event selection cuts due to PDF (a) and QCD scale variations (b)	207
10.16	Systematic variation of the ratio $r^{\text{obs.}}$ after all event selection cuts due to variations of the accepted top (a) and $gg \rightarrow WW$ (b) background cross section by factors of up to 8.	208
10.17	Distribution of $n^{\text{sub.}}$ over a large number of Monte Carlo experiments for the background only and for the signal plus background hypothesis for Higgs boson masses of 120, 130, 140 and 150 GeV	211
10.18	The same as Figure 10.17 for Higgs boson masses of 160, 170, 180 and 190 GeV.	212
10.19	Higgs boson discovery significance in the $gg \rightarrow H \rightarrow W^-W^+ \rightarrow l^-\bar{\nu} l^+\nu$ final state as function of the Higgs boson mass	213
A.1	Illustration of the Gauss transformation	II
B.1	Spin correlations in the $H \rightarrow WW$ decay	V

B.2	Comparison of the Monte Carlo predictions for the shapes of $\Delta\phi_u$ and m_T for the $q\bar{q} \rightarrow WW$ process	VI
B.3	Evolution of the resonant and non-resonant $t\bar{t}$ and Wt Feynman diagrams when external b -jets are vetoed down to the $b\bar{b} \rightarrow WW$ process.	VIII
B.4	Comparison of the Monte Carlo prediction for the shapes of $\Delta\phi_u$ and m_T for the top processes	IX
B.5	The same comparison as in Figure B.4 but with a relaxed jet veto threshold of $p_T > 30$ GeV	IX
B.6	Correlation of the leading jet transverse momentum and pseudo-rapidity after the cut on $p_T(l_1 + l_2) > 30$ GeV for the Higgs boson signal process ($m_H = 130$ GeV and $m_H = 170$ GeV) and the dominant dilepton backgrounds.	X
B.7	Resolution of the x -component of the missing transverse energy vector	XII
B.8	The same as Figure 10.8 but for the $q\bar{q} \rightarrow Z \rightarrow \tau\tau$ process. . . .	XIII
B.9	Lepton identification and reconstruction efficiency as function of the true lepton p_T and η for electrons (a) and muons (b) in $q\bar{q} \rightarrow Z \rightarrow ll$ events.	XIV
B.10	Absolute and normalised electron calorimeter isolation energy sum and track isolation momentum sum in a cone of $\Delta R < 0.3$	XV
B.11	Same as Figure B.10 for muons	XVI
B.12	Electron and muon impact parameter significance	XVII
B.13	Electron p_T and η distribution	XIX
B.14	Muon p_T and η distribution	XIX
C.1	Comparison of the transverse mass definitions for signal and background processes after the selection cut on $\Delta\phi_u < 1$	XXII
C.2	Correlation of $\Delta\phi_u$ and M_T^{classic}	XXII
C.3	Correlation of $\Delta\phi_u$ and M_T^{WW}	XXIII
C.4	Correlation of $\Delta\phi_u$ and M_T^{true}	XXIII
C.5	Correlation of $\Delta\phi_u$ and $M_T^{\text{true,T}}$	XXIII
C.6	Normalised distribution of M_T^{classic} for $\Delta\eta_u < 1$ and $\Delta\eta_u > 1$ after the selection cut on $\Delta\phi_u < 1$	XXIV
C.7	Normalised distribution of M_T^{WW} for $\Delta\eta_u < 1$ and $\Delta\eta_u > 1$ after the selection cut on $\Delta\phi_u < 1$	XXIV
C.8	Normalised distribution of M_T^{true} for $\Delta\eta_u < 1$ and $\Delta\eta_u > 1$ after the selection cut on $\Delta\phi_u < 1$	XXV
C.9	Normalised distribution of $M_T^{\text{true,T}}$ for $\Delta\eta_u < 1$ and $\Delta\eta_u > 1$ after the selection cut on $\Delta\phi_u < 1$	XXV
D.1	Same as Figure 9.20 for Higgs boson masses of 120 GeV and 130 GeV.	XXVII
D.2	Same as Figure 9.20 for Higgs boson masses of 140 GeV and 150 GeV.	XXVII

D.3	Same as Figure 9.20 for Higgs boson masses of 160 GeV and 170 GeV.	XXVIII
D.4	Same as Figure 9.20 for Higgs boson masses of 180 GeV and 190 GeV.	XXVIII
D.5	Same as Figure 10.12 for Higgs boson masses of 110 GeV and 120 GeV.	XXVIII
D.6	Same as Figure 10.12 for Higgs boson masses of 130 GeV and 140 GeV.	XXIX
D.7	Same as Figure 10.12 for Higgs boson masses of 150 GeV and 160 GeV.	XXIX
D.8	Same as Figure 10.12 for Higgs boson masses of 170 GeV and 180 GeV.	XXIX
D.9	Same as Figure 10.12 for Higgs boson masses of 190 GeV and 200 GeV.	XXX
D.10	Same as Figure 10.12 for Higgs boson masses of 110 GeV and 120 GeV, but all backgrounds are included.	XXX
D.11	Same as Figure 10.12 for Higgs boson masses of 130 GeV and 140 GeV, but all backgrounds are included.	XXX
D.12	Same as Figure 10.12 for Higgs boson masses of 150 GeV and 160 GeV, but all backgrounds are included.	XXXI
D.13	Same as Figure 10.12 for Higgs boson masses of 170 GeV and 180 GeV, but all backgrounds are included.	XXXI
D.14	Same as Figure 10.12 for Higgs boson masses of 190 GeV and 200 GeV, but all backgrounds are included.	XXXI

Bibliography

- [1] S.L. Glashow, Nucl. Phys. **22** (1961) 579, S. Weinberg, Phys. Rev. Lett. **19** (1967) 1264, A. Salam, in Elementary Particle Theory, ed. N. Svartholm, Stockholm, “Almqvist and Wiksell” (1968), 367.
- [2] F. Halzen, A. D. Martin, Quarks & Leptons: An Introductory Course in Modern Particle Physics, (John Wiley & Sons, 1984).
- [3] A. Djouadi, The Anatomy of electro-weak symmetry breaking. I: The Higgs boson in the standard model, Phys. Rept. **457** (2008) 1–216.
- [4] R. Barate et al., Search for the standard model Higgs boson at LEP, Phys. Lett. **B565** (2003) 61–75.
- [5] The LEP Electroweak Working Group, <http://lepewwg.web.cern.ch/LEPEWWG/>, 2009.
- [6] T. Hambye, K. Riesselmann, Matching conditions and Higgs mass upper bounds revisited, Phys. Rev. **D55** (1997) 7255–7262.
- [7] V. Barger, R.J.N. Philips, Collider Physics, (Addison-Wesley, 1987).
- [8] J. Pumplin et al., New generation of parton distributions with uncertainties from global QCD analysis, JHEP **07** (2002) 012.
- [9] D. de Florian, M. Grazzini, Higgs production through gluon fusion: updated cross sections at the Tevatron and the LHC, Phys. Lett. **B674** (2009) 291–294.
- [10] A. D. Martin, W. J. Stirling, R. S. Thorne, G. Watt, Parton distributions for the LHC, arXiv:0901.0002v3 [**hep-ph**] (2009).
- [11] F. Ambrogini et al., Proceedings of the Workshop on Monte Carlo’s, Physics and Simulations at the LHC PART I, arXiv:0902.0293 [**hep-ph**] (2009).
- [12] F. Ambrogini et al., Proceedings of the Workshop on Monte Carlo’s, Physics and Simulations at the LHC PART II, arXiv:0902.0180 [**hep-ph**] (2009).
- [13] S. Frixione, B. R. Webber, Matching NLO QCD computations and parton shower simulations, JHEP **06** (2002) 029.

- [14] U. Aglietti et al., Tevatron for LHC report: Higgs, hep-ph/0612172 [**hep-ph**] (2006).
- [15] G. Aad et al., Expected Performance of the ATLAS Experiment: Detector, Trigger and Physics, (CERN, Geneva, 2008), CERN-OPEN-2008-020.
- [16] Combined CDF and DZero Upper Limits on Standard Model Higgs-Boson Production with up to 4.2 fb⁻¹ of Data, arXiv:0903.4001 [**hep-ex**] (2009).
- [17] E. Accomando et al., Workshop on CP Studies and Non-Standard Higgs Physics, hep-ph/0608079 [**hep-ph**] (2006).
- [18] M. Dührssen, Prospects for the measurement of Higgs boson coupling parameters in the mass range from 110 - 190 GeV, Internal Report ATL-PHYS-2003-030, CERN, Geneva, Jul 2003.
- [19] M. Dührssen et al., Extracting Higgs boson couplings from LHC data, Phys. Rev. **D70** (2004) 113009.
- [20] C. Lefevre, The CERN accelerator complex, June 2006.
- [21] W. W. Armstrong et al., ATLAS: Technical proposal for a general-purpose p p experiment at the Large Hadron Collider at CERN, (1994), CERN-LHCC-94-43.
- [22] CMS, the Compact Muon Solenoid: Technical proposal, (1994), CERN-LHCC-94-38.
- [23] H. P. Beck, Triggering at high luminosity colliders, New Journal of Physics **9** (2007) 334.
- [24] G. Aad et al., The ATLAS Experiment at the CERN Large Hadron Collider, JINST **3** (2008) S08003.
- [25] J. Allison et al., Geant4 developments and applications, IEEE Transactions on Nuclear Science **53 No. 1** (2006) 270–278.
- [26] S. Agostinelli et al., G4 - a simulation toolkit, Nuclear Instruments and Methods **A 506** (2003) 250–303.
- [27] K. Assamagan et al., The ATLAS Monte Carlo Project, Internal Report ATL-COM-SOFT-2008-024, CERN, Geneva, Dec 2008, To be submitted to JINST.
- [28] D. Cavalli et al., Performance of the ATLAS fast simulation ATL-FAST, Internal Report ATL-PHYS-INT-2007-005; ATL-COM-PHYS-2007-012, CERN, Geneva, Jan 2007.

-
- [29] A. Salzburger et al., The Fast ATLAS Track Simulation (FATRAS), Internal Report ATL-SOFT-PUB-2008-001; ATL-COM-SOFT-2008-002, CERN, Geneva, Mar 2008.
 - [30] D. Cavalli et al., Validation of the fast calorimeter simulation FastCaloSim against the Athena release 10 and 13.0.40 full simulation, Internal Report ATL-PHYS-INT-2009-073. ATL-COM-PHYS-2009-164, CERN, Geneva, Aug 2009.
 - [31] D. Adams et al., The ATLFAST-II performance in release 14 -particle signatures and selected benchmark processes-, Internal Report ATL-COM-PHYS-2009-554, CERN, Geneva, Oct 2009.
 - [32] G. Grindhammer, S. Peters, The parameterized simulation of electromagnetic showers in homogeneous and sampling calorimeters, 1993, hep-ex/0001020.
 - [33] A. Camard, F. Hubaut, B. Laforge, P. Schwemling, Study of the EM Barrel Module 0 with muons, Internal Report ATL-LARG-2001-017, CERN, Geneva, May 2001.
 - [34] G. Corcella et al., HERWIG 6.5: an event generator for Hadron Emission Reactions With Interfering Gluons (including supersymmetric processes), JHEP **01** (2001) 010.
 - [35] T. Binoth, M. Ciccolini, N. Kauer, M. Kramer, Gluon-induced WW background to Higgs boson searches at the LHC, JHEP **03** (2005) 065.
 - [36] B. P. Kersevan, E. Richter-Was, The Monte Carlo event generator AcerMC version 2.0 with interfaces to PYTHIA 6.2 and HERWIG 6.5, (2004).
 - [37] T. Sjostrand et al., High-energy physics event generation with PYTHIA 6.1, Comput. Phys. Commun. **135** (2001) 238–259.
 - [38] ATLAS Hadronic Calibration Workshop, <http://www.lip.pt/atlas-workshop/>, 2009.
 - [39] N. Kauer, Top pair production beyond double pole approximation: pp , $pp \rightarrow p\bar{p} \rightarrow$ six fermions and zero, one or two additional partons, Phys. Rev. **D67** (2003) 054013.
 - [40] M. Dührssen, K. Jakobs, J. J. van der Bij, P. Marquard, The process $gg \rightarrow WW$ as a background to the Higgs signal at the LHC, JHEP **05** (2005) 064.
 - [41] J. Alwall et al., MadGraph/MadEvent v4: The New Web Generation, JHEP **09** (2007) 028.

- [42] C. Anastasiou, G. Dissertori, F. Stockli, B. R. Webber, QCD radiation effects on the $H \rightarrow WW \rightarrow l \nu l \nu$ signal at the LHC, *JHEP* **03** (2008) 017.
- [43] G. Arnison et al., Experimental observation of isolated large transverse energy electrons with associated missing energy at $s^{*}(1/2) = 540\text{-GeV}$, *Phys. Lett.* **B122** (1983) 103–116.
- [44] M. Banner et al., Observation of single isolated electrons of high transverse momentum in events with missing transverse energy at the CERN anti-p p collider, *Phys. Lett.* **B122** (1983) 476–485.
- [45] D. L. Rainwater, D. Zeppenfeld, Observing $H \rightarrow W^{(*)}W^{(*)} \rightarrow e^{\pm}\mu^{\mp}\cancel{p}_T$ in weak boson fusion with dual forward jet tagging at the CERN LHC, *Phys. Rev.* **D60** (1999) 113004.
- [46] A. J. Barr, B. Gripaios, C. G. Lester, Measuring the Higgs boson mass in dileptonic W-boson decays at hadron colliders, *JHEP* **07** (2009) 072.