

Search for additional Higgs bosons decaying to tau
leptons and measurement of the $\mathcal{CP}$ properties of
the Higgs Yukawa coupling to top quarks using the
CMS detector

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Abstract

In 2012 the CMS and ATLAS experiments announced the discovery of a boson with mass around 125 GeV and properties consistent with that of the Higgs boson predicted by the standard model (SM). Despite this success, many observations about the universe appear to be inconsistent with the SM phenomenology, implying the need for beyond the standard model (BSM) physics. Many BSM scenarios require extended Higgs sectors, which usually include additional Higgs bosons and predict different properties of the 125 GeV state. Two approaches may be used to test such scenarios: search for the additional bosons, or measure the properties of the 125 GeV state to test the comparability with the SM. This thesis explores both of these options. A search for additional Higgs bosons in the mass region 90 GeV–3.2 TeV decaying to pairs of tau leptons is presented. This analysis uses the dataset collected by CMS in 2016, which corresponds to an integrated luminosity of 35.9 fb^{-1} . No statistically significant excesses are observed. Exclusion limits are set on the cross sections times branching fractions. The results are also interpreted in the context of the minimal supersymmetric standard model (MSSM). A measurement of the $\mathcal{CP}$ -properties of the Higgs Yukawa coupling to top quarks is also performed using Higgs boson production via gluon fusion in association with two jets. The analysis of the azimuthal angular correlations between the jets is sensitive to the $\mathcal{CP}$ mixing angle α_{hgg} . This analysis uses the datasets collected by CMS in 2016 and 2017 which correspond to an integrated luminosity of 71.8 fb^{-1} . The best-fit value of α_{hgg} is measured to be $-59^{\circ}{}^{+47^{\circ}}_{-50^{\circ}}$. This constitutes the first direct measurement of the $\mathcal{CP}$ -properties of a fermionic Higgs coupling.

Declaration

I declare that the work presented in this thesis is my own. Figures and results taken from other sources are indicated by a reference in the text or figure caption. Figures labelled "CMS" are sourced from CMS publications. Figures labelled "CMS Supplementary" have been made public as supplementary material accompanying a publication. All Figures with these labels, including those made by myself, also include the relevant reference in the caption. The analyses presented in this document were developed in collaboration with other members of the CMS experiment. Chapters 1 and 2 and the first three sections in Chapter 3 do not contain original work by myself, but are intended to give an overview of the theoretical background, the CMS detector, and the standard object reconstruction techniques that are essential to the work presented in subsequent sections. I contributed to the measurements of corrections described in Chapter 3, where the corrections measured by myself are indicated in the text. The background methods used in the analyses, which are also described in this Chapter, were developed and tested in collaboration with other CMS members. For the embedded method, I was responsible for measuring several of the corrections, and for measuring the muon selection efficiencies needed to estimate the background yields. The fake-factor method was developed by other members of the $H \rightarrow \tau\tau$ group, and my only contribution to this work was in validation. I was solely responsible for the QCD method in the $\tau_e\tau_\mu$ channel, but it should be noted that this work builds on the methodology used for previous CMS analyses. For the MSSM analysis described in Chapter 4, I contributed to the optimisation of event selections and categorisation, studies of background methods, evaluating systematic uncertainties, and the statistical interpretation of the result. These results were published in Ref. [1]. For the $\mathcal{CP}$ analysis described in Chapter 5, I was responsible for all stages of the analysis including: optimising event selections and categorisations, studying background methods, deriving corrections, evaluating systematics, and the statistical interpretation of the results. This analysis is currently not published, but a publication is planned in the near future. This will include additional data collected by CMS in 2018, and a combination with other Higgs final states. For the trigger developments described in Chapter 6, I was responsible for: studies of variables, optimisations of cut thresholds, rate measurements, and performance studies. I also had a small contribution to the publication that established the observation of the Higgs boson decaying to a pair of tau leptons [2]. My contribution to this analysis was in the evaluation of theoretical uncertainties on the signal, but this work will not be further discussed in this thesis.

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Chapter 1

Theory and motivation

In order to adequately describe the Higgs searches and measurements detailed in subsequent chapters, it is first necessary to introduce the theory and motivations behind them. This chapter begins by giving an overview of the standard model (SM) of particle physics with emphasis on the Higgs mechanism. This is followed by a discussion of several issues faced by the SM, which highlights the need for beyond the standard model (BSM) physics. The phenomenology of the Higgs sectors in a few examples of such BSM scenarios will then be described. This includes the Higgs sectors in the minimal supersymmetric standard model (MSSM) and in $\mathcal{CP}$ -violating Higgs models.

1.1 The standard model of particle physics

The SM of particle physics is a renormalisable quantum field theory (QFT), that encapsulates our best understanding of all known elementary particles and their interactions. Fundamental to the standard model is the principle of local gauge invariance, where demanding the Lagrangian to be invariant under gauge transformations gives rise to conserved currents and interactions. The gauge symmetry group of the SM is $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$.

In the SM there are 12 fundamental fermions, which are summarised in Table 1.1. This includes six quarks, three charged leptons, and three neutrinos. Each of these particles also has an antimatter partner with the same mass but opposite quantum numbers. The fermions are further organised into three generations, where across generations particles of the same type have identical properties apart from their mass. The fermions interact by four fundamental forces: the strong and weak nuclear forces, the electromagnetic force, and gravity. All forces except gravity are described by the SM, and are mediated by the gluons (g), the $W^\pm$ and Z bosons, and the photon (γ), as summarised in Table 1.2. All fermions interact via the weak force, and those that are charged interact electromagnetically. Only the quarks carry colour charge, the quantum number associated with the strong force, and hence also have strong interactions.

The theory that describes the strong interaction is quantum chromodynamics (QCD), governed by a $SU(3)_C$ gauge symmetry. All particles with colour charge, C , interact by the strong force. The fact that the gluons that mediate the strong force are themselves coloured leads to two characteristic properties of QCD: asymptotic freedom and confinement. Due to asymptotic freedom the strength of the strong interaction decreases with energy. Confinement requires that observed states have neutral colour charges, and as a result quarks are always found in bound states known as hadrons.

The unification of the electromagnetic and weak forces was first proposed by Glashow [4],

Table 1.1: The fundamental fermions organised by generation.

Type	Charge	Generation		
		1	2	3
Quarks				
up-type	$+\frac{2}{3}$	up (u)	charm (c)	top (t)
down-type	$-\frac{1}{3}$	down (d)	strange (s)	bottom (b)
Leptons				
neutral	0	electron neutrino (ν_e)	muon neutrino (ν_μ)	tau neutrino (ν_τ)
charged	-1	electron (e)	muon (μ)	tau (τ)

Table 1.2: The fundamental forces described in the SM and the corresponding vector bosons. The charges and masses of the bosons are also summarised [3].

Force	Boson	Charge	Mass (GeV)
Electromagnetic	photon (γ)	0	0
Weak	Z	0	91.2
	$W^\pm$	± 1	80.4
Strong	gluon (g)	0	0

Weinberg [5], and Salam [6] in the 1960s. The unified theory, known as electroweak theory, has symmetry group $SU(2)_L \otimes U(1)_Y$, where the subscript "L" indicates that the weak force couples only to left handed fermions. The weak hypercharge, Y, corresponding to the $U(1)_Y$ symmetry, is related to the electric charge, Q, and the third component of the weak isospin, t_3 , by

$$Y = 2(Q - t_3). \quad (1.1)$$

The corresponding gauge fields are the three weak isospin fields from $SU(2)_L$, W_μ^i ($i=1,2,3$), and the weak hypercharge field from $U(1)_Y$, B_μ . These fields are related to the physical $W^\pm$ ($W_\mu^\pm$), Z (Z_μ), and γ (A_μ) states by

$$\begin{aligned} W_\mu^\pm &= \frac{1}{\sqrt{2}} (W_\mu^1 \mp iW_\mu^2), \\ Z_\mu &= W_\mu^3 \cos \theta_w - B_\mu \sin \theta_w, \text{ and} \\ A_\mu &= W_\mu^3 \sin \theta_w + B_\mu \cos \theta_w, \end{aligned} \quad (1.2)$$

where θ_w is the weak mixing angle, which is related to the weak (g) and electromagnetic (g') coupling constants by

$$\sin \theta_w = \frac{g'}{\sqrt{g^2 + g'^2}} \text{ and } \cos \theta_w = \frac{g}{\sqrt{g^2 + g'^2}}. \quad (1.3)$$

The issue with the electroweak theory discussed thus far is that it is not possible to simply write down mass terms for the $W^\pm$ and Z bosons of the form $\frac{1}{2}m_Z^2 Z_\mu Z^\mu$ and $m_W^2 W_\mu^- W^{+\mu}$, without breaking the gauge invariance of the Lagrangian. A similar problem also occurs if we try to add mass terms for the fermions. This problem implied the need for electroweak symmetry breaking in the theory, which is achieved using the Higgs mechanism as described in the next section. The Higgs mechanism was subsequently incorporated in the electroweak theory prior to the publications by Weinberg and Salam.

1.2 The Higgs mechanism

The Higgs mechanism was first proposed in the 1960s by Englert and Brout [7]; Higgs [8–10]; Guralnik, Hagen and Kibble [11]; and Kibble [12]. By this mechanism the electroweak symmetry is spontaneously broken. This is achieved through the introduction of a field that is symmetric under gauge transformations, but has a non-zero vacuum expectation value (VEV). According to Goldstone's theorem [13], a consequence of spontaneous symmetry breaking is the emergence of a massless scalar boson for each broken generator. These so called Goldstone bosons become the longitudinal degrees of freedom of the $W^\pm$ and Z bosons.

The simplest field that can be introduced to break the electroweak symmetry is a complex scalar $SU(2)_L$ doublet

$$\phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix}. \quad (1.4)$$

The corresponding Lagrangian for this field is

$$\mathcal{L}_{\text{Higgs}} = (D_\mu \phi)^\dagger (D^\mu \phi) - V(\phi^\dagger \phi), \quad (1.5)$$

where the covariant derivative is given by

$$D_\mu = \partial_\mu - \frac{i}{2} g \vec{W}_\mu \cdot \vec{\sigma} - \frac{i}{2} g' B_\mu. \quad (1.6)$$

The potential term is chosen as

$$V(\phi^\dagger \phi) = \lambda(\phi^\dagger \phi)^2 + \mu^2 \phi^\dagger \phi, \quad (1.7)$$

where λ and μ^2 are both real, and $\lambda > 0$ is required for the vacuum to be stable. For $\mu^2 < 0$ the field acquires a non-zero VEV and we get spontaneous symmetry breaking. The potential has a circle of minima, of which we arbitrarily choose the minimum

$$\langle 0 | \phi | 0 \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ \sqrt{\frac{-\mu^2}{\lambda}} \end{pmatrix} = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v \end{pmatrix}, \quad (1.8)$$

where $v \sim 246$ GeV is the VEV [3]. By choosing the appropriate gauge in which the Goldstone bosons are eliminated, we can write the field ϕ as an expansion about this minimum,

$$\phi = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + H \end{pmatrix}. \quad (1.9)$$

Substituting this expression into Equation 1.5 and ignoring all but the terms quadratic in the fields gives

$$\mathcal{L}_{\text{Higgs}} = \frac{1}{2} \partial_\mu H \partial^\mu H + \frac{1}{4} g^2 v^2 W_\mu^+ W^{-\mu} + \frac{1}{8} \frac{g^2 v^2}{\cos^2 \theta_w} Z_\mu Z^\mu - \mu^2 H^2, \quad (1.10)$$

where we have used Equation 1.2 to relate the B_μ and W_μ^i fields to the physical Z_μ and $W_\mu^\pm$ states. This Lagrangian contains mass terms for the $W^\pm$ and Z bosons, with $m_W = \frac{gv}{2}$ and $m_Z = \frac{1}{2} \frac{gv}{\cos \theta_w}$, and importantly no mass term for the photon field. There is also a massive scalar Higgs field H with $m_H = \sqrt{-2\mu^2}$. It is also worth noting that the interaction terms omitted from the above equation give couplings of the Higgs to gauge bosons $\propto m_V^2$.

The quarks and charged leptons also acquire mass via the Higgs mechanism through Yukawa interactions. Taking the first generation as an example, the Yukawa part of the Lagrangian is

$$\mathcal{L}_{\text{Yukawa}} = -y_e \bar{l}_L \phi e_R - y_d \bar{q}_L \phi d_R - y_u \bar{q}_L \epsilon_{ij} \phi_j^* u_R + \text{h.c.}, \quad (1.11)$$

where y_f is the Yukawa coupling strength for fermion f , $\epsilon_{ij} = \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}$, $\bar{l}_L = \begin{pmatrix} \bar{\nu}_L \\ \bar{e}_L \end{pmatrix}$, $\bar{q}_L = \begin{pmatrix} \bar{u}_L \\ \bar{d}_L \end{pmatrix}$, and $e_{L/R}$, ν_L , $u_{L/R}$, and $d_{L/R}$ denote the left and right handed fermion fields. After expanding the above expression about the minimum we get terms of the form $-y_f \frac{v}{\sqrt{2}} \bar{f}_L f_R + \text{h.c}$ corresponding to mass terms with $m_f = y_f \frac{v}{\sqrt{2}}$, and interaction terms $\propto m_f$. Note the absence of a ν_R field means that the neutrinos do not acquire their mass through this mechanism. Generalising to three generations, we get Yukawa terms that mix quark generations of the form $-(Y_d)_{ij}(q_{Li}\phi)d_{Rj}$, where i and j are generation indices. If we change basis to diagonalise Y_{ij} we get generation mixing in charged current weak interactions, governed by the CKM matrix.

1.3 Higgs physics at the LHC

In 2012 both ATLAS and CMS experiments announced the discovery of a boson with a mass around 125 GeV [14, 15]. Measurements of the properties of this boson, such as its spin and parity [16, 17], and its couplings to the other particles confirmed its consistency with the SM Higgs boson. Since the discovery observations of decays to $\gamma\gamma$ [18, 19], ZZ [20, 21], WW [22, 23], $\tau\tau$ [2, 24], and $b\bar{b}$ [25, 26] have been reported. Both CMS and ATLAS have also recently announced observations of Higgs boson production in association with $t\bar{t}$ pairs [27, 28], providing the first direct measurement of the Higgs coupling to top quarks. All current measurements of the Higgs boson couplings are found to be consistent with the SM predictions, as illustrated in Figure 1.1.

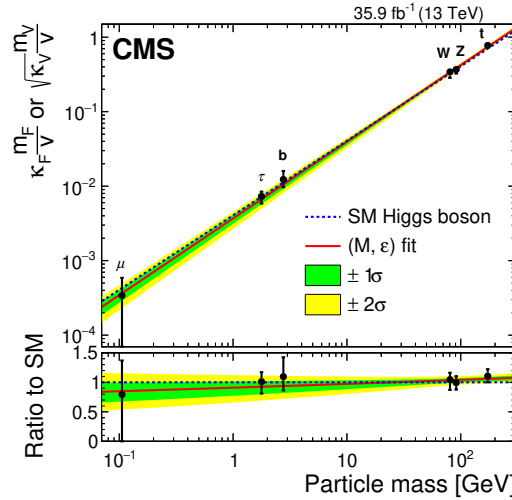


Figure 1.1: The measured coupling strengths of the Higgs boson to fermions (t , b , τ) as a function of particle mass, where $\kappa_f = y_f/y_f^{\text{SM}}$. The current upper limit on the Higgs coupling to muons is also shown by the leftmost data point. The "reduced" coupling strengths of the Higgs to vector bosons are also shown, which are defined as $\sqrt{\kappa_V} m_V/v$, where κ_V are the bosonic couplings relative to the SM expectations [29].

The Feynman diagrams for the dominant Higgs production processes at the Large Hadron Collider (LHC) are illustrated in Figure 1.2. These production modes include: gluon fusion (GF), vector boson fusion (VBF), W/Z-associated production (VH), and $t\bar{t}$ -associated production ($t\bar{t}H$). The cross sections for each process at $\sqrt{s} = 13$ TeV are shown in Table 1.3. The dominant production mode is GF, which has a cross section one to two orders of magnitude larger than the other processes. Although smaller in cross section, the other production modes have more distinct topologies that can be exploited to reduce backgrounds. For example, the VBF production process is characterised by two energetic jets, at least one of which is typically found in forward detector

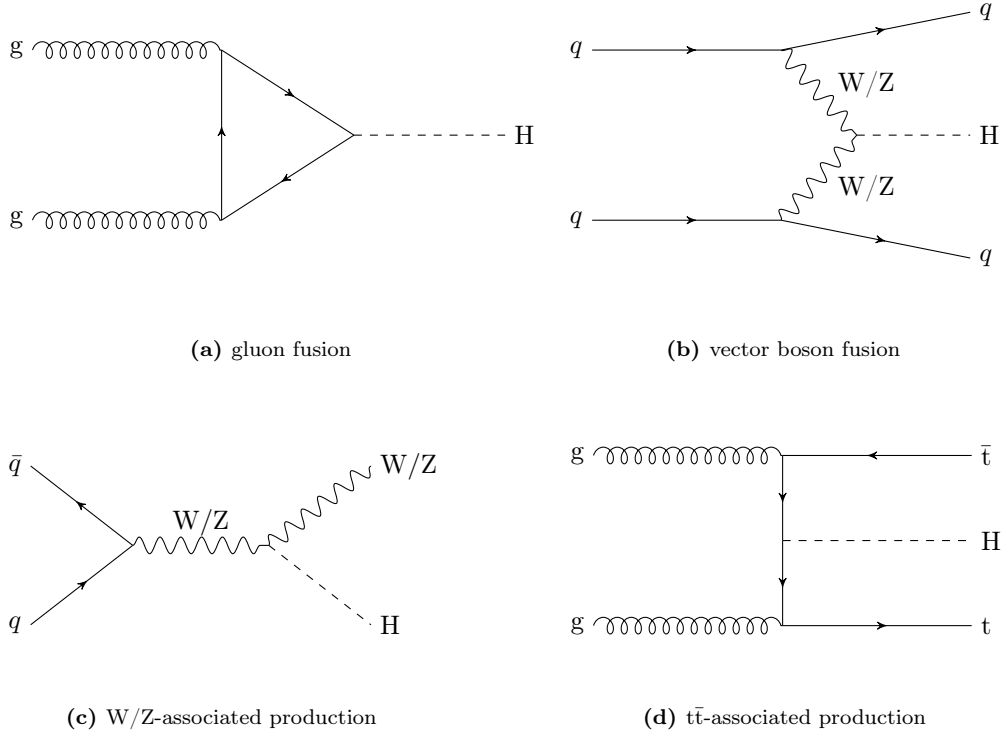


Figure 1.2: Feynman diagrams showing the dominant Higgs boson production modes at the LHC.

regions. These dijets tend to have a large separation in pseudorapidity, $|\Delta\eta_{jj}|$, and high invariant mass, m_{jj} .

Table 1.3: The Higgs boson production cross sections at $\sqrt{s} = 13$ TeV for $m_H = 125$ GeV [30]. The theoretical uncertainties on the cross sections are also shown.

Process	Cross section (pb)	Uncertainty (%)
GF	48.58	5.0
VBF	3.78	2.1
W-associated	1.37	2.0
Z-associated	0.88	4.1
$t\bar{t}$ -associated	0.51	6.8

The SM Higgs branching ratios are shown in Table 1.4. The dominant decay mode is $H \rightarrow b\bar{b}$, with a branching ratio of almost 60%. Despite this sizeable branching ratio, the $b\bar{b}$ final state is plagued by large backgrounds from QCD multi-jet events, which limits its sensitivity. The $\gamma\gamma$ and $ZZ \rightarrow 4l$ decays on the other hand have much smaller branching ratios, but were instrumental in the Higgs discovery thanks to their clean experimental signatures. The $\tau\tau$ channel, of interest in this thesis, has a relatively large branching ratio of about 6%. Although not as clean as the ZZ or $\gamma\gamma$ final states, it is less impaired by backgrounds than $b\bar{b}$. The $\tau\tau$ channel also has good sensitivity to the VBF phase space, which is advantageous for the $\mathcal{CP}$ analysis described in Chapter 5.

1.4 Beyond the standard model

With the discovery of the Higgs boson all hypothesised SM particles have now been observed. Although tested to a high degree of precision, the SM does not provide an explanation for several

observations, which include:

- The discovery of neutrino oscillations [31] which implies that neutrinos must have mass. The SM currently does not include a mechanism for generating these masses. Mechanisms have been proposed [32] but are yet to be experimentally verified and incorporated.
- Several astronomical observations provide strong evidence for the existence of dark matter, which is expected to make up about 85% of the matter in the universe [33]. The SM does not include a candidate particle for this dark matter.
- The running couplings of the strong and electroweak interactions do not intersect at a high energy scale, as we would expect for a grand unified theory (GUT) [34].
- The amount of $\mathcal{CP}$ violation predicted by the SM does not account for the large matter-anti-matter asymmetry observed in the universe.
- The force of gravity does not appear in the SM.

Aside from these observations, the issue most relevant to the Higgs sector is the hierarchy problem [35]. The SM is accepted to be an effective theory that is valid up to an energy scale Λ , where new physics effects become important. If we take the value of Λ to be the Plank scale, where quantum gravitational effects become relevant, then $\Lambda \sim 10^{19}$. The Higgs mass receives corrections from the virtual effects of every particle that couples directly or indirectly to the Higgs field. The physical mass is related to the bare mass m_H^0 and the correction Δm_H^2 by

$$m_H^2 = m_H^{0\,2} + \Delta m_H^2. \quad (1.12)$$

The size of the Δm_H^2 correction depends on Λ . For example, the Feynman diagram for the fermionic (top quark) loop correction shown in Figure 1.3 (a) yields a correction that diverges quadratically with Λ :

$$\Delta m_H^2 = -\frac{y_f^2}{8\pi^2} \Lambda^2 + \dots \quad (1.13)$$

[35]. Obtaining a value of $m_H = 125$ GeV would therefore require extreme fine tuning of m_H^0 to cancel the $\Delta m_H^2 \sim \mathcal{O}(10^{38})$ correction.

Table 1.4: The Higgs boson branching ratio for $m_H = 125$ GeV [30].

Decay mode	Branching ratio (%)
$b\bar{b}$	58.24
$\tau\tau$	6.27
$\mu\mu$	0.02
$c\bar{c}$	2.89
WW	21.37
ZZ	2.61
gg	8.19
$\gamma\gamma$	0.23
$Z\gamma$	0.15

One popular solution to the hierarchy problem is supersymmetry (SUSY) [36, 37], which postulates a symmetry that relates fermions and bosons. According to this theory, every SM particle has at least one supersymmetric partner ¹ known as a superpartner. The superpartners of

¹The left- and right-handed components of the SM fermions each have superpartners denoted $\tilde{f}_L$ and $\tilde{f}_R$ which mix to form two mass eigenstates, $\tilde{f}_1$ and $\tilde{f}_2$.

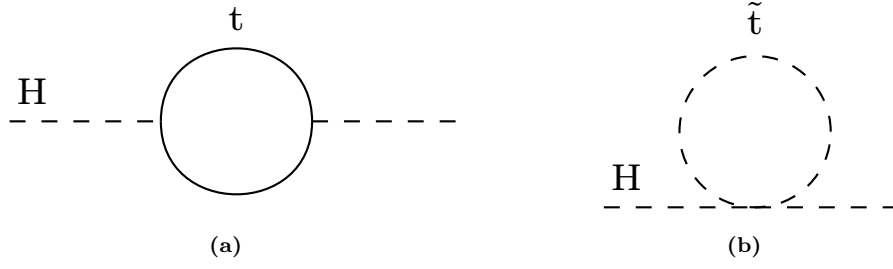


Figure 1.3: Feynman diagrams for the one loop corrections to the Higgs boson mass for a fermionic top quark loop (a) and a scalar stop squark loop (b).

the fermions are scalar bosons, which are named by adding a prefix "s" to the fermion name (e.g. the superpartner of the top is the stop). The superpartners of the bosons are fermions, which are named by appending "ino" to the root of the boson name (e.g. the superpartner of the gluon is the gluino). The symbols for the superpartners are the same as for the corresponding SM particles, except with a tilde placed above the symbol (e.g. the stop is denoted $\tilde{t}$).

The superpartner loop corrections to the Higgs mass, such as the sfermion loop shown in Figure 1.3 (b), give terms that contribute to Δm_H^2 with opposite signs relative to the equivalent SM terms. For example, adding the corresponding sfermion term to Equation 1.13 gives

$$\Delta m_H^2 = -\frac{y_f^2}{8\pi^2}\Lambda^2 + \frac{\lambda_S}{8\pi^2}\Lambda^2 + \dots, \quad (1.14)$$

and the terms will exactly cancel as long as the coupling strength λ_S is related to the fermion Yukawa coupling by $\lambda_S = y_f^2$ [35]. The cancellation is also achieved for all terms contributing to Δm_H^2 . If SUSY were unbroken, then the masses of the SUSY particles would be the same as their SM partners, and the relation $\lambda_S = y_f^2$ would hold. However, the lack of any superpartner detections implies that, if SUSY is realised in nature, it must be broken. In this case, the cancellation in the Δm_H^2 terms is not exact. Hence, the SUSY breaking scale, and consequently the masses of the superpartners, cannot be much larger than $\mathcal{O}(1 \text{ TeV})$, without requiring further fine tuning of m_H^0 [35].

SUSY also provides solutions to two of the other problems listed at the start of the section. Most notably, SUSY modifies the running couplings of the strong and electroweak interactions, which allows them to intersect at a high energy [34]. For R -parity [35] conserving SUSY models, the lightest SUSY particle is also a candidate for dark matter. The simplest supersymmetric extension of the SM is the MSSM [38, 39], which includes only the smallest number of new particles and interactions consistent with phenomenology [40].

1.5 The MSSM Higgs sector

The MSSM Higgs sector requires (at least) two complex scalar doublets to generate mass terms for both up-type quarks and down-type fermions [41],

$$\phi_d = \begin{pmatrix} \phi_d^0 \\ \phi_d^- \end{pmatrix}, \quad \phi_u = \begin{pmatrix} \phi_u^+ \\ \phi_u^0 \end{pmatrix}. \quad (1.15)$$

As for the SM, the potential associated with these fields induces electroweak symmetry breaking and the neutral components of the fields acquire vacuum expectation values:

$$\langle 0|\phi_d|0\rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} v_d \\ 0 \end{pmatrix}, \quad \langle 0|\phi_u|0\rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_u \end{pmatrix}. \quad (1.16)$$

The values of v_d and v_u are related to the SM VEV by $v^2 = v_u^2 + v_d^2$. The parameter $\tan\beta$ is defined as the ratio of the two vacuum expectation values, $\tan\beta = \frac{v_u}{v_d}$.

After spontaneous symmetry breaking, three of the degrees of freedom result in Goldstone bosons, which become the longitudinal components of the $W^\pm$ and Z bosons. The remaining five degrees of freedom correspond to five physical Higgs bosons. In the $\mathcal{CP}$ -conserving MSSM, this includes two neutral scalars, h and H , one neutral pseudoscalar, A , and two charged Higgs bosons, $H^\pm$. These states arise from the mixing of the fields:

$$\begin{aligned} H^\pm &= \phi_d^\pm \sin\beta + \phi_u^\pm \cos\beta \\ A &= \sqrt{2} (\text{Im } \phi_d^0 \sin\beta + \text{Im } \phi_u^0 \cos\beta) \\ h &= -(\sqrt{2}\text{Re } \phi_d^0 - v_d) \sin\alpha + (\sqrt{2}\text{Re } \phi_u^0 - v_u) \cos\alpha \\ H &= (\sqrt{2}\text{Re } \phi_d^0 - v_d) \cos\alpha + (\sqrt{2}\text{Re } \phi_u^0 - v_u) \sin\alpha. \end{aligned} \quad (1.17)$$

These relations depend on v_u , v_d , and a mixing angle α which is found by diagonalising the mass matrix of which the h and H bosons are eigenstates,

$$\mathcal{M}_0^2 = \begin{pmatrix} m_A^2 \sin^2\beta + m_Z^2 \cos^2\beta & -(m_A^2 + m_Z^2) \sin\beta \cos\beta \\ -(m_A^2 + m_Z^2) \sin\beta \cos\beta & m_A^2 \cos^2\beta + m_Z^2 \sin^2\beta \end{pmatrix}. \quad (1.18)$$

At tree-level the MSSM Higgs sector can be described in terms of two free parameters which are conventionally chosen to be the mass of the A boson, m_A , and $\tan\beta$. The squared-masses of the h and H bosons are the eigenvalues of $\mathcal{M}_0^2$,

$$m_{H,h}^2 = \frac{1}{2} \left(m_A^2 + m_Z^2 \pm \sqrt{(m_A^2 + m_Z^2)^2 - 4m_Z^2 m_A^2 \cos^2 2\beta} \right), \quad (1.19)$$

and the masses of the charged Higgs bosons are given by

$$m_{H^\pm}^2 = m_A^2 + m_W^2. \quad (1.20)$$

From Equation 1.19 we also obtain an expression relating α to the free parameters,

$$\cos^2(\beta - \alpha) = \frac{m_h^2(m_Z^2 - m_h^2)}{m_A^2(m_H^2 - m_h^2)}, \quad (1.21)$$

and an upper bound on the tree-level mass of the lightest Higgs boson,

$$m_h \leq m_Z |\cos 2\beta| \leq m_Z. \quad (1.22)$$

Even in cases where h is interpreted as the observed 125 GeV Higgs boson, the fact that this value is at most m_Z is not a problem, as the radiative corrections to the tree-level mass allow for values of $m_h \sim 125$ GeV. The properties of the MSSM Higgs sector are affected by these radiative corrections, which depend on several of the SUSY-breaking parameters. These parameters/corrections are usually fixed in representative benchmark scenarios. The analysis described in this thesis considers two of these scenarios: hMSSM [42] and $m_h^{\text{mod}+}$ [43], which will be described in detail in Section 1.5.2.

Table 1.5 shows the tree-level couplings of the three neutral Higgs bosons to the $W^\pm/Z$ and the up- and down-type fermions, relative to the equivalent couplings to the Higgs boson in the SM. In the so-called decoupling limit, where $m_A \gg m_Z$, the masses of the A , H , and $H^\pm$ bosons are almost degenerate ($m_A \approx m_H \approx m_{H^\pm}$) and $\alpha \approx \beta - \frac{\pi}{2}$. The magnitudes of the tree-level coupling in this limit are shown in Table 1.6. On inspection of these couplings, we can see that the h boson

has very similar properties to the SM Higgs boson in the decoupling limit. The couplings of the A and the H bosons to the down-type particles also scale with $\tan \beta$, and conversely the couplings to up-type quarks scale with $\frac{1}{\tan \beta}$ in this limit. This results in an enhancement in the branching ratios to $b\bar{b}$ and $\tau\tau$ for large values of $\tan \beta$. Searches in the ditau final state are therefore well motivated for the high $\tan \beta$ regions of the parameter space. The branching ratios to taus in the hMSSM is shown as a function of m_A - $\tan \beta$ in Figure 1.4. Even for the largest $\tan \beta$ regions, the branching ratios to taus are at most $\mathcal{O}(10\%)$. While this is typically ≈ 10 times smaller than the $b\bar{b}$ branching ratio, the ditau final state is experimentally much cleaner and thus more sensitive to MSSM Higgs searches. For low $\tan \beta$ values the dominant decays are to bosons or $t\bar{t}$ pairs. However, the branching ratio of the A boson to taus can also be sizeable in the region $m_A \lesssim 300$ GeV where the decays to $t\bar{t}$ pairs are kinematically forbidden, and larger still in the region $m_A \lesssim 200$ GeV where the $A \rightarrow hZ$ decay is also restricted.

Table 1.5: The tree-level couplings of the neutral MSSM bosons relative to the couplings to the Higgs boson in the SM [44].

Boson	Couplings		
	g_{VV}	g_{uu}	g_{dd}/g_{ll}
A	0	$\cot \beta$	$\tan \beta$
H	$\cos(\beta - \alpha)$	$\frac{\sin \alpha}{\sin \beta}$	$\frac{\cos \alpha}{\cos \beta}$
h	$\sin(\beta - \alpha)$	$\frac{\cos \alpha}{\sin \beta}$	$-\frac{\sin \alpha}{\cos \beta}$

Table 1.6: The tree-level couplings of the neutral MSSM bosons in the decoupling limit ($m_A \gg m_Z$) relative to the couplings to the Higgs boson in the SM.

Boson	Couplings		
	g_{VV}	g_{uu}	g_{dd}/g_{ll}
A	0	$\cot \beta$	$\tan \beta$
H	0	$\cot \beta$	$\tan \beta$
h	1	1	1

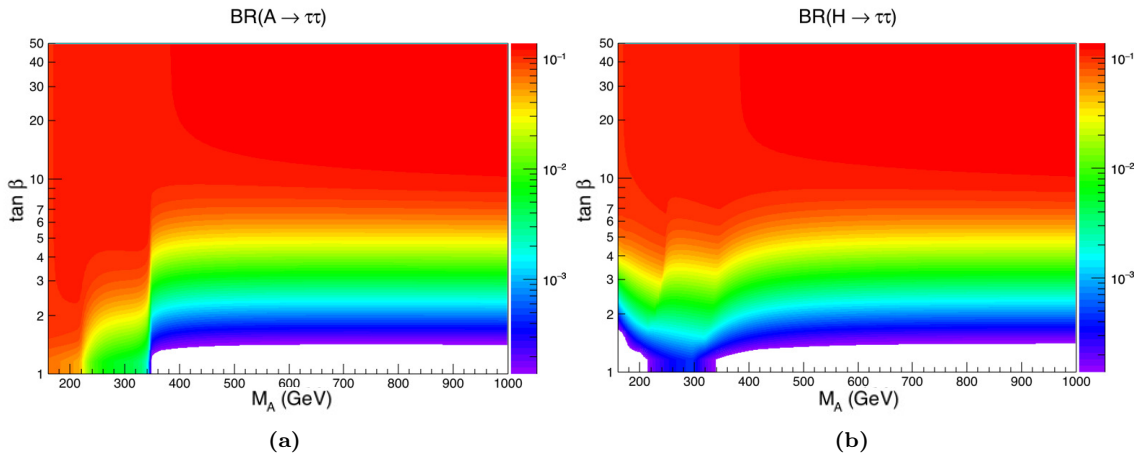


Figure 1.4: The $\tau\tau$ branching ratios of the A (left) and H (right) bosons as a function of m_A - $\tan \beta$ in the hMSSM benchmark scenario [45].

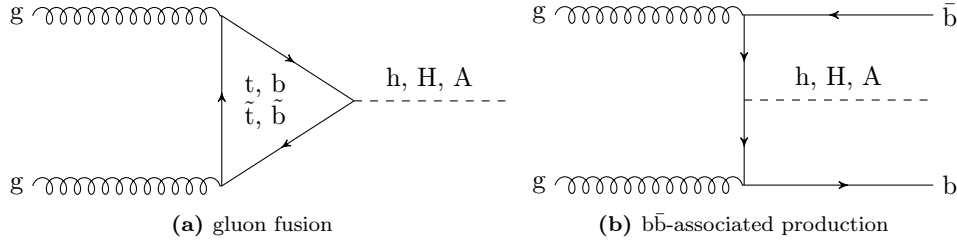


Figure 1.5: Feynman diagrams showing the dominant A and H boson production modes at the LHC.

1.5.1 MSSM Higgs boson production at the LHC

The dominant processes producing A and H bosons at the LHC are GF and $b\bar{b}$ -associated production. The Feynman diagrams for these processes are shown in Figure 1.5. The VBF and W/Z-associated production modes, which are relevant for SM Higgs production, are suppressed over much of the parameter space. This is due to the lack of tree-level couplings of the A boson to WW and ZZ, and the fact that the equivalent H boson couplings tend to zero in the decoupling limit.

The cross sections for the GF and $b\bar{b}$ -associated production processes in the hMSSM are shown as a function of m_A - $\tan\beta$ in Figure 1.6. At low $\tan\beta$ the GF production process is dominant. The GF loop has contributions from both top and bottom quarks, in addition to SUSY particles such as the stop and sbottom. For low values of $\tan\beta$, the loop is dominated by the t quark contribution, whereas at high $\tan\beta$ the b quark contribution is larger as a result of the enhanced coupling to down-type fermions. The enhanced b quark coupling also increases the cross section for $b\bar{b}$ -associated production with $\tan\beta$. This process is therefore dominant in the high $\tan\beta$ regions of the parameter space.

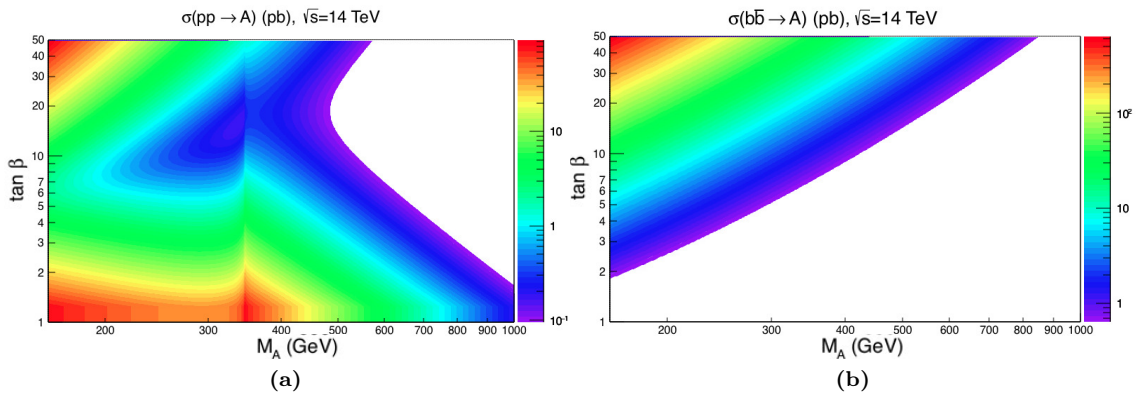


Figure 1.6: Cross section for GF (left) and the $b\bar{b}$ -associated (right) production of the A boson at $\sqrt{s} = 14$ TeV as a function of m_A - $\tan\beta$ in the hMSSM benchmark scenario [45].

1.5.2 MSSM benchmark scenarios

Although at tree-level the MSSM Higgs sector can be described in terms of the two free parameters, the higher order radiative corrections are sizeable and depend on several of the SUSY-breaking

parameters. Benchmark scenarios are therefore defined that depend only on m_A - $\tan\beta$. In these scenarios the radiative corrections must be estimated appropriately, to exhibit interesting phenomenology and be consistent with the observation of the 125 GeV Higgs boson. The analysis described in this thesis considers the $m_h^{\text{mod}+}$ and hMSSM scenarios. For the former, the 125 GeV Higgs is interpreted as the h boson, and the radiative corrections are estimated by setting the SUSY parameters to sensible choices based on experimental and theoretical arguments to give $m_h \sim 125$ GeV. The hMSSM uses a different method to estimate the corrections, whereby the mass of the lightest scalar is fixed by construction to the mass of the observed Higgs boson, and then the measured value of m_h is used to estimate the dominant radiative corrections. Both of the considered scenarios will be described in more detail in the next sections.

1.5.2.1 The $m_h^{\text{mod}+}$ scenario

The $m_h^{\text{mod}+}$ scenario [43] is a modification of the m_h^{max} [46] scenario. For the m_h^{max} scenario, the SUSY parameters were chosen so as to maximise the value of m_h as a function of $\tan\beta$. This was chosen to give conservative exclusion limits on $\tan\beta$ from the Large Electron-Positron (LEP) searches. Values of m_h up to ~ 135 GeV are allowed in this case. However, in the advent of the 125 GeV Higgs observation, this scenario is found to be inconsistent with $m_h \sim 125$ GeV over most of the parameter space. For the $m_h^{\text{mod}+}$ scenario, several of the parameter values were therefore modified to maximise the m_A - $\tan\beta$ parameter space consistent with $m_h = 125 \pm 3$ GeV, where the ± 3 GeV reflects the uncertainty on the m_h prediction. Only small regions of phase space at low values of $\tan\beta$ and/or low m_A are inconsistent with $m_h = 125 \pm 3$ GeV in this scenario.

The parameters that have the largest effects on the MSSM Higgs sector predictions are: the top quark mass (which is taken to be 173.2 GeV), the left- and right-handed mass parameters of the third generation squarks which are given as $M_{\tilde{t}_L} = M_{\tilde{b}_L} = M_{\tilde{t}_R} = M_{\tilde{b}_R} =: M_{\text{SUSY}}$, and the stop and sbottom mixing parameters X_t and X_b . There are also significant dependencies on: the stau mixing parameter X_τ , the higgsino mass parameter μ , the SUSY-breaking scale in the stau/sneutrino sector $M_{\tilde{l}_3}$, the gluino mass $m_{\tilde{g}}$, and the U(1) and SU(2) gaugino mass parameters M_1 and M_2 .

The value of M_{SUSY} is set to 1 TeV. The mixing parameter X_t is set to $1.5M_{\text{SUSY}}$, which can be related to X_b and X_τ using the following relations for the trilinear Higgs-stop, Higgs-sbottom, and Higgs-stau couplings, A_t , A_b , and A_τ :

$$\begin{aligned} X_t &= A_t - \mu^* \cot\beta \\ X_b &= A_b - \mu^* \tan\beta \\ X_\tau &= A_\tau - \mu^* \tan\beta, \end{aligned} \tag{1.23}$$

where $A_t = A_b = A_\tau$ is assumed. The values of M_1 and M_2 are assumed to be related by the GUT relation,

$$M_1 = \frac{5 \sin^2 \theta_w}{3 \cos^2 \theta_w} M_2, \tag{1.24}$$

with M_2 set to 200 GeV. The values of μ , $m_{\tilde{g}}$, and $M_{\tilde{l}_3}$ are also set to 200 GeV, 1.5 TeV, and 1 TeV respectively.

Lastly, there are also several parameters that only have a minor impact on the MSSM Higgs sector predictions. These include the SUSY-breaking masses in the first and second generation squark (slepton) sectors, which are set to 1.5 TeV (500 GeV), and the trilinear couplings to the first and second generation squarks and sleptons, which are set to zero.

1.5.2.2 The hMSSM scenario

The hMSSM scenario [42, 45] is formulated based on several assumptions. The first assumption is that the observed 125 GeV state corresponds to the h boson. The neutral $\mathcal{CP}$ -even Higgs mass matrix can be written as

$$\mathcal{M}^2 = \mathcal{M}_0^2 + \begin{pmatrix} \Delta\mathcal{M}_{11}^2 & \Delta\mathcal{M}_{12}^2 \\ \Delta\mathcal{M}_{12}^2 & \Delta\mathcal{M}_{22}^2 \end{pmatrix}, \quad (1.25)$$

where $\mathcal{M}_0^2$ is the tree-level mass matrix defined in Equation 1.18, and the radiative corrections are given by the 2×2 matrix with elements $\Delta\mathcal{M}_{ij}^2$. The second assumption of the hMSSM is that only the $\Delta\mathcal{M}_{22}^2$ term, that involves the dominant stop-top sector correction, needs to be taken into account, $\Delta\mathcal{M}_{22}^2 \gg \Delta\mathcal{M}_{11}^2, \Delta\mathcal{M}_{12}^2$. Under this assumption, the mass eigenstates of the matrix are given by

$$m_{H,h}^2 = \frac{1}{2} \left(m_A^2 + m_Z^2 + \Delta\mathcal{M}_{22}^2 \pm \sqrt{m_A^4 + m_Z^4 - 2m_A^2 m_Z^2 \cos 4\beta + \Delta\mathcal{M}_{22}^4 + 2\Delta\mathcal{M}_{22}^2(m_A^2 - m_Z^2) \cos 2\beta} \right). \quad (1.26)$$

By substituting the known value of m_h the $\Delta\mathcal{M}_{22}^2$ correction can be determined,

$$\Delta\mathcal{M}_{22}^2 = \frac{m_h^2(m_A^2 + m_Z^2 - m_h^2) - m_A^2 m_Z^2 \cos^2 2\beta}{m_Z^2 \cos^2 \beta + m_A^2 \sin^2 \beta - m_h^2}. \quad (1.27)$$

This can then be used to estimate the value of m_H in terms of only $\tan \beta$, m_A , and m_h ,

$$m_H^2 = \frac{(m_A^2 + m_Z^2 - m_h^2)(m_Z^2 \cos^2 \beta + m_A^2 \sin^2 \beta) - m_A^2 m_Z^2 \cos^2 2\beta}{m_Z^2 \cos^2 \beta + m_A^2 \sin^2 \beta - m_h^2}. \quad (1.28)$$

The mixing angle α is also given by these three parameters,

$$\tan \alpha = -\frac{(m_Z^2 + m_A^2) \cos \beta \sin \beta}{m_Z^2 \cos^2 \beta + m_A^2 \sin^2 \beta - m_h^2}. \quad (1.29)$$

The third assumption of the hMSSM is that the couplings of all Higgs bosons to fermions and gauge bosons are given in terms of $\tan \beta$ and α only, as listed in Table 1.5. The couplings of two Higgs bosons to one gauge boson, as well as the trilinear self-couplings among Higgs bosons are also given in terms of only $\tan \beta$ and α . The only relevant trilinear Higgs couplings for LHC phenomenology are the Hhh and hhh self-couplings, which are given by

$$\begin{aligned} \lambda_{hhh} &= \lambda_{hhh}^{\text{tree}} + 3 \frac{\Delta\mathcal{M}_{22}^2 \cos^3 \alpha}{m_Z^2 \sin \beta} \\ \lambda_{Hhh} &= \lambda_{Hhh}^{\text{tree}} + 3 \frac{\Delta\mathcal{M}_{22}^2 \sin \alpha}{m_Z^2 \sin \beta} \cos^2 \alpha, \end{aligned} \quad (1.30)$$

where $\lambda_{hhh,Hhh}^{\text{tree}}$ are the corresponding tree-level couplings [45].

Finally, the superparticles are assumed to have masses too heavy to be detected at the LHC, and apart from their effects on $\mathcal{M}^2$ do not impact the Higgs sector. This means that only decays of the Higgs bosons into SM particles are allowed in the hMSSM. Altogether, this gives enough information to determine the masses of the five Higgs bosons and the cross sections and branching ratios as a function of only $\tan \beta$ and m_A .

The assumptions made in formulating the hMSSM are reasonable for low and moderate $\tan \beta$ values, where the bottom Yukawa coupling is not significantly enhanced and consequently the $\Delta\mathcal{M}_{22}^2 \gg \Delta\mathcal{M}_{11}^2, \Delta\mathcal{M}_{12}^2$ approximation holds. The hMSSM is also not well defined in cases where the denominators of Equations 1.27–1.29 tend to zero. In such regions of the parameter space $\mathcal{M}_{22}^2$ diverges and we get $m_H \rightarrow \infty$.

1.6 $\mathcal{CP}$ violation in the Higgs sector

The SM assumes the simplest possible scalar sector which includes just one Higgs doublet. This does not allow for the possibility of $\mathcal{CP}$ -violation in the Higgs sector, as there exists only one state with $\mathcal{CP}$ -even couplings to the other particles. However, there are numerous examples of BSM physics scenarios with more complex Higgs sectors, and many of these models include $\mathcal{CP}$ -violating Higgs couplings [47]. Such models could have important implications for explaining the matter anti-matter asymmetry observed in the universe, which is one of the deficiencies of the SM highlighted in Section 1.4.

The simplest extension of the SM Higgs sector that can allow $\mathcal{CP}$ -violation is a two Higgs doublet model (2HDM) [48]. In these models the SM field content is extended by including an additional complex scalar doublet. The MSSM Higgs sector described previously is an example of such a model. As for the MSSM, a generic 2HDM includes five physical Higgs bosons, three of which are neutral and two of which are charged. For $\mathcal{CP}$ -conserving models these bosons correspond to the familiar h , H , A , and $H^\pm$. Conversely, for $\mathcal{CP}$ -violating models the three neutral bosons are not $\mathcal{CP}$ eigenstates, but rather admixtures of the $\mathcal{CP}$ -even and $\mathcal{CP}$ -odd components. These bosons are denoted h_1 , h_2 , and h_3 , where $m_{h_1} \leq m_{h_2} \leq m_{h_3}$. One of these states may be interpreted as the observed 125 GeV Higgs boson; a hypothesis that can be tested by measuring the $\mathcal{CP}$ -properties of the Higgs couplings.

Since we have discussed the MSSM extensively in the previous section, it is worth noting that at tree-level the Higgs couplings in MSSM are constrained by SUSY and are consequently $\mathcal{CP}$ -conserving. However, loop effects mediated predominantly by the third generation squarks can lead to sizeable $\mathcal{CP}$ -violation [47]. In terms of the Higgs particle content though the end result is the same, namely three neutral bosons with indefinite $\mathcal{CP}$ parities.

In addition to the 2HDM extensions, there are many other models with rich phenomenologies that predict $\mathcal{CP}$ -violation in the Higgs sector. A review of several such scenarios is given in Ref. [47].

1.6.1 Higgs $\mathcal{CP}$ measurements at the LHC

The $\mathcal{CP}$ nature of the Higgs can be probed using the experimentally clean $H \rightarrow ZZ$ final state, where variables such as the azimuthal angle between the two Z decay planes are sensitive to the Higgs $\mathcal{CP}$ -state [49]. Such measurements provided the first tests of the $\mathcal{CP}$ properties of the Higgs [16, 17], by directly comparing the pure $\mathcal{CP}$ -even (0^+) and pure $\mathcal{CP}$ -odd (0^-) hypotheses as illustrated in Figure 1.7. These initial measurements were able to exclude the 0^- hypothesis at confidence levels in excess of 95%. The analysis strategy has since been extended to include also information from the production vertices in the VBF and VH production modes, where the associated jet topologies are also sensitive to the $\mathcal{CP}$ nature of the HVV vertex [50]. These techniques have been used to test $\mathcal{CP}$ -violating scenarios, in which there is an admixture of $\mathcal{CP}$ -even and $\mathcal{CP}$ -odd components. For example, a recent CMS publication [51] set bounds on the fractional $\mathcal{CP}$ -odd $H \rightarrow ZZ$ cross section f_{a_3} . Values of f_{a_3} larger than 0.092 were excluded at the 95% confidence level (CL). Although these bounds appear relatively stringent, the sensitivity to mixed $\mathcal{CP}$ scenarios in most models is actually quite weak. This is because the $\mathcal{CP}$ -odd component is usually heavily suppressed due to the lack of tree-level AVV couplings. In such cases the HVV vertex effectively projects out the $\mathcal{CP}$ -even component, which results in observable distributions that are practically identical to the pure $\mathcal{CP}$ -even case even if the state has a sizeable $\mathcal{CP}$ -odd contribution [52].

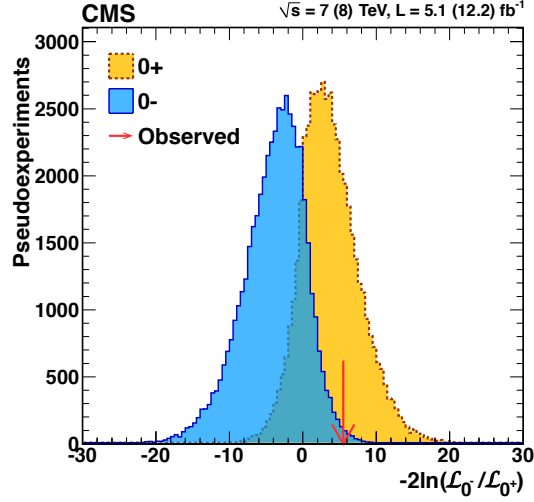


Figure 1.7: Comparison of the pure $\mathcal{CP}$ -even (0^-) and pure $\mathcal{CP}$ -odd (0^+) hypotheses in the $H \rightarrow ZZ$ final state [16].

For fermionic Higgs couplings the situation is different. In such cases there exists both Hff and Aff tree-level couplings, meaning the magnitudes of the $\mathcal{CP}$ -even and $\mathcal{CP}$ -odd components can be comparable and $\mathcal{CP}$ -violating effects can be large. Several methods have been proposed for probing these fermionic couplings, most of which consider the $H\tau\tau$ or $Ht\bar{t}$ vertices. The $H\tau\tau$ vertex is probed using $H \rightarrow \tau\tau$ decays, where angular correlations between tau decay products are sensitive to the $\mathcal{CP}$ state. These measurement are however made more difficult by the presence of neutrinos in the final state. Several procedures have been proposed to extract the $\mathcal{CP}$ information, which usually centre around the hadronic τ decay modes [53–55]. Such decays involve fewer neutrinos and often the taus decay to multiple visible particles which allows more angular observables to be defined.

The $\mathcal{CP}$ properties of the $Ht\bar{t}$ vertex can be probed directly using correlations of the t quark decay products in $t\bar{t}$ -associated Higgs production [56]. It is also possible to probe the $\mathcal{CP}$ -nature of the $Ht\bar{t}$ Yukawa coupling using the GF process, under the assumption that the quark loop is dominated by the top quark. The current experimental measurements of the GF cross section, coupled with the direct measurements of the top Yukawa coupling using $t\bar{t}H$ events, indicate that the loop is indeed dominated by the top quark contribution. The $\mathcal{CP}$ information is extracted using Higgs bosons produced in association with two jets, where the azimuthal angular correlations of the jets is sensitive to the $\mathcal{CP}$ properties. As these properties are measured in the production vertex events can be targeted using any Higgs final state. The analysis described in Chapter 5 focuses on the $H \rightarrow \tau\tau$ channel, which is expected to have good sensitivity to the dijet phase space compared to other final states. The methodology and parametrisation is based on Ref. [57] and will be detailed in the next section.

1.6.2 Measuring the Higgs $\mathcal{CP}$ properties using gluon fusion

The $\mathcal{CP}$ properties of the Yukawa couplings to fermions f ($f = t, b, \tau, \dots$) are given by the Lagrangian

$$\mathcal{L}_{\text{Yukawa}} = y_f H \bar{\psi}_f \psi_f + \tilde{y}_f A \bar{\psi}_f \gamma_5 \psi_f, \quad (1.31)$$

where H and A demote the scalar and pseudoscalar Higgs fields. The coupling strengths are given by the y_f and $\tilde{y}_f$ terms, which are assumed to be of SM strength $y_f \approx \tilde{y}_f \approx y_f^{SM} = m_f/v$, where

m_f is the mass of the fermion and v is the VEV of the SM Higgs field. Under this assumption GF production is dominated by the top quark loop and can be described by the effective Lagrangian

$$\mathcal{L}_{\text{eff}} = \kappa_t \cdot \frac{\alpha_s}{12\pi v} \cdot H G_{\mu\nu}^a G^{a\mu\nu} + \tilde{\kappa}_t \cdot \frac{\alpha_s}{16\pi v} \cdot A G_{\mu\nu}^a G_{\rho\sigma}^a \epsilon^{\mu\nu\rho\sigma}, \quad (1.32)$$

where $G_{\mu\nu}^a$ is the gluon field strength, α_s is the strong coupling constant, $\kappa_t = \frac{y_t}{y_t^{\text{SM}}}$, and $\tilde{\kappa}_t = \frac{\tilde{y}_t}{y_t^{\text{SM}}}$ [58, 59].

From this effective Lagrangian emerge Hgg, Hggg, and Hgggg vertices, which are illustrated in Figure 1.8. The tensor structure of the Hgg vertex is

$$T^{\mu\nu} = a_2(q_1 \cdot q_2 g^{\mu\nu} - q_1^\nu q_2^\mu) + a_3 \epsilon^{\mu\nu\rho\sigma} q_{1\rho} q_{2\sigma}, \quad (1.33)$$

where q_1 and q_2 are the 4-momenta of the incoming gluons, and a_2 and a_3 are scalar form factors related to the Yukawa interactions [57],

$$a_2 = \kappa_t \frac{\alpha_s}{3\pi v}, \quad a_3 = \tilde{\kappa}_t \frac{\alpha_s}{2\pi v}. \quad (1.34)$$

The ratio of these form factors can be used to define a mixing angle,

$$\tan \alpha_{\text{hgg}} = \frac{a_3}{a_2}, \quad (1.35)$$

which can take any value between $-90^\circ < \alpha_{\text{hgg}} \leq 90^\circ$. Values of $\alpha_{\text{hgg}} = 0^\circ$ and $\alpha_{\text{hgg}} = 90^\circ$ correspond to the pure $\mathcal{CP}$ -even and pure $\mathcal{CP}$ -odd scenarios respectively. All values in the range $0^\circ < |\alpha_{\text{hgg}}| < 90^\circ$ imply $\mathcal{CP}$ -violation, where $|\alpha_{\text{hgg}}| = 45^\circ$ corresponds to the maximally $\mathcal{CP}$ -violating case when $a_2 = a_3$.

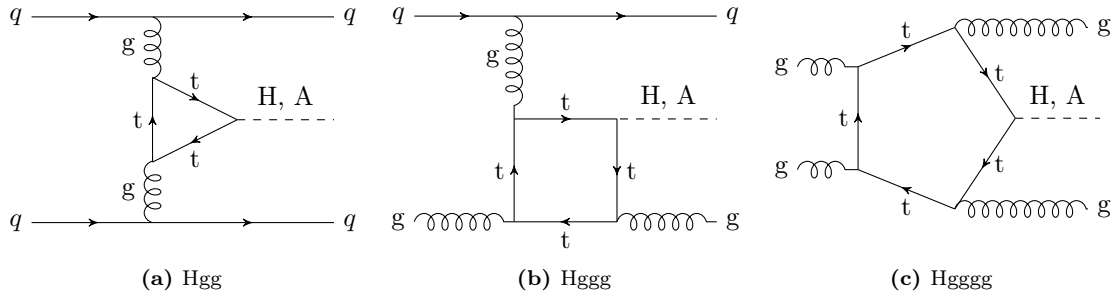


Figure 1.8: Feynman diagrams contributing to GF Higgs production in association with two jets. These diagrams are associated with the Hgg (a), the Hggg (b), and the Hgggg (c) vertices.

The structure of the Feynman diagram associated with the Hgg vertex in Figure 1.8 is very similar to the VBF production process shown in Figure 1.2 (b). The azimuthal angle between the two tagging jets $\Delta\phi_{jj}$ in this case is sensitive to the tensor structure of the Hgg vertex, which is related to the $\mathcal{CP}$ -structure of the $Ht\bar{t}$ Yukawa coupling. The absolute value of $\Delta\phi_{jj}$ is a powerful tool for discriminating between the two tensor structures of Equation 1.33. However, as $|\Delta\phi_{jj}|$ is a $\mathcal{CP}$ -even observable it is not sensitive to the interference between the $\mathcal{CP}$ -even and $\mathcal{CP}$ -odd contributions. This means that the tensor structure cannot be unambiguously determined in cases where there are comparable $\mathcal{CP}$ -even and $\mathcal{CP}$ -odd contributions. The missing information is encoded in the sign of $\Delta\phi_{jj}$, which is defined by ordering the jets according to their pseudorapidity [57]. The signed $\Delta\phi_{jj}$ distributions are shown in Figure 1.9 for the pure $\mathcal{CP}$ -even, pure $\mathcal{CP}$ -odd, and $\mathcal{CP}$ -mixed scenarios. The signed $\Delta\phi_{jj}$ is a $\mathcal{CP}$ -odd observable and therefore sensitive to the interference. Observing an asymmetry in $\Delta\phi_{jj}$ would directly demonstrate that $\mathcal{CP}$ is violated.

The contributions from the Hggg and Hgggg couplings illustrated in Figures 1.8 (b) and (c) dilute the sensitivity to the structure of the Hgg vertex. There are also diagrams with Hgg vertices contributing to the H+2j events that are not sensitive to the Higgs $\mathcal{CP}$ -state. An example is shown in Figure 1.10, where a gluon splits to form two outgoing quarks. The analysing power can be improved by applying VBF-like topological selections such as large m_{jj} and/or $|\Delta\eta_{jj}|$ requirements. This enhances the contribution of the inherently VBF-like Hgg diagrams such as in Figure 1.8 (a). Comparisons of the $\Delta\phi_{jj}$ variable after applying different m_{jj} selections will be shown in Chapter 5.

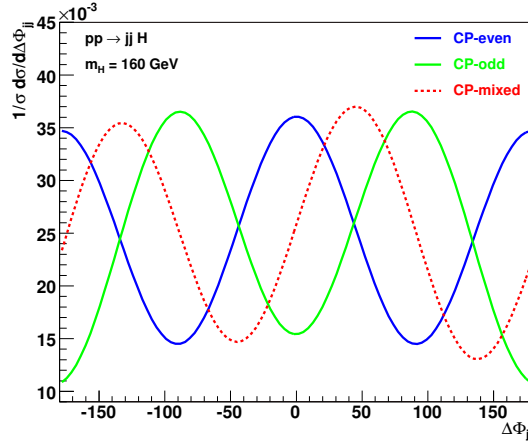


Figure 1.9: Normalised distributions of the signed jet azimuthal angle distribution. The distributions are shown for the pure $\mathcal{CP}$ -even case (blue), the pure $\mathcal{CP}$ -odd (green), and a $\mathcal{CP}$ -mixed case (red). A $|\Delta\eta_{jj}| > 3$ cut is applied for the two tagging jets [57].

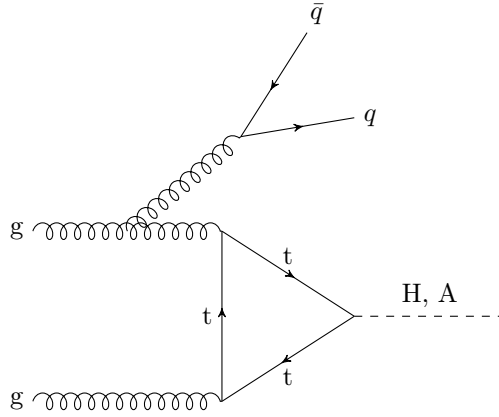


Figure 1.10: An example Feynman diagram for a process contributing to the H+2j events that is not sensitive to the Higgs $\mathcal{CP}$ properties. In this example a gluon splits to form two outgoing quarks.

Chapter 2

The LHC and CMS Detector

2.1 The LHC

The LHC [60] is a 27 km long synchrotron located at CERN just outside Geneva. It is designed to collide proton beams with centre of mass energies ($\sqrt{s}$) up to 14 TeV and instantaneous luminosity of about $10^{34} \text{ cm}^{-2}\text{s}^{-1}$. The LHC also has the capability to collide heavy ions with energies up to 2.8 TeV per nucleon and a peak luminosity of about $10^{27} \text{ cm}^{-2}\text{s}^{-1}$.

A schematic showing the layout of the LHC and CERN accelerator complex is shown in Figure 2.1. The LHC is situated in the tunnel originally constructed to house the LEP collider [61]. The tunnel contains two adjacent parallel beam pipes that intersect at four points. The LHC is supplied with protons from the injector chain Linac2 – Proton Synchrotron Booster (PSB) – Proton Synchrotron – Super Proton Synchrotron (SPS), raising the proton energies to 50 MeV, 1.4 GeV, 25 GeV, and 450 GeV respectively [62]. Protons from the SPS are injected into the LHC beam pipes as two counter-rotating beams. Protons in the beams are bunched together, with each bunch containing $\mathcal{O}(10^{11})$ protons. At design operation each beam consists of 2808 bunches with a bunch spacing of 25 ns. The protons are accelerated to collision energy using eight 400 MHz radio frequency (RF) cavities. To keep the beams following a circular trajectory 1232 niobium-titanium superconducting dipole magnets are used. These magnets must be kept at their operating temperature of 1.9 K, requiring 96 tonnes of superfluid helium. At the four intersection points bunches are collided at a rate of 40 MHz. Collision events are then recorded by the ALICE [63], ATLAS [64], CMS [65], and LHCb [66] detectors.

The event rate for a process produced in LHC collisions is given by

$$N_{\text{event}} = L\sigma_{\text{event}}, \quad (2.1)$$

where σ_{event} is the cross section of the process and L is the machine luminosity. The cross sections for many of the processes of interest are very small and therefore require a high luminosity in order to accumulate a sufficient number of events. Taking Higgs production by GF as an example, with a production cross section of 55 pb at $\sqrt{s} = 14 \text{ TeV}$ [30] the LHC operating at its design luminosity ($10^{34} \text{ cm}^{-2}\text{s}^{-1}$) will produce about one Higgs boson every two seconds.

The LHC began its first physics collisions in 2010, initially colliding beams at $\sqrt{s} = 7 \text{ TeV}$ which was increased to 8 TeV throughout 2012. At the end of this data-taking period, collectively known as Run 1, the LHC began the two year Long Shutdown 1 (LS1). During this time the LHC and its detector were upgraded to prepare for an increase in $\sqrt{s}$. The LHC was re-started in 2015 to begin its Run 2 era of data-taking at $\sqrt{s} = 13 \text{ TeV}$, lasting until the end of 2018. During

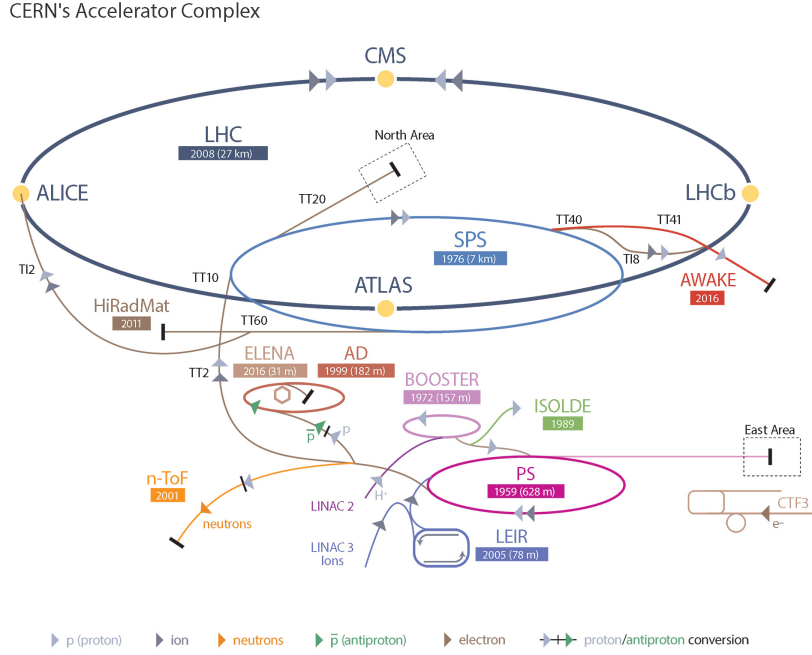


Figure 2.1: Schematic showing the layout of the CERN accelerator complex [67].

Run 2 the LHC reached and surpassed its initial design luminosity of $10^{34} \text{ cm}^{-2}\text{s}^{-1}$, achieving a record peak luminosity of $2.1 \times 10^{34} \text{ cm}^{-2}\text{s}^{-1}$ in 2018 [68]. Figure 2.2 shows the total integrated luminosity of proton-proton collisions delivered to CMS by the LHC at the time of writing [68]. This includes 6.1 fb^{-1} of data delivered at 7 TeV, 23.3 fb^{-1} at 8 TeV, and 163.2 fb^{-1} at 13 TeV. The datasets used for the analyses described in this thesis are those collected during the 2016 and 2017 data-taking periods at 13 TeV. Only data recorded by CMS in which all sub-detectors were functioning correctly is certified as good for use in physics analyses. This corresponds to 35.9 fb^{-1} and 41.9 fb^{-1} of data collected in 2016 and 2017 respectively.

2.2 The CMS detector

The CMS detector was designed to achieve high sensitivity in the searches for the Higgs boson and new physics at the TeV scale. It consists of several sub-detectors built round the central beam axis, as illustrated in Figure 2.3. The detector is 21.6 m long, 14.6 m in diameter, and weighs 12500 tonnes [65]. The central feature of the detector is the superconducting solenoid operated at 3.8 T. Located within the bore of the magnet are the inner tracker and electromagnetic and hadronic calorimeters. The muon system, comprised of three types of gaseous muon detectors, is located outside of the superconducting coil embedded in the iron return yoke of the solenoid. These sub-detectors will be described in more detail in the following sections.

The coordinate system employed by CMS has its origin centred at the nominal collision point within the detector. The y -axis points vertically upwards and the x -axis points radially inwards towards the centre of the LHC. The z -axis points along the beam axis. The transverse energy and momentum, E_T and p_T , are defined in the x - y plane. The azimuthal angle ϕ is also measured in the x - y plane with respect to the x -axis. The polar angle θ is measured from the z -axis. It is common practise to define the polar coordinate in terms of the pseudorapidity defined as $\eta = -\ln[\tan(\frac{\theta}{2})]$. Relative distances in the ϕ - η plane are given by $\Delta R = \sqrt{\Delta\phi^2 + \Delta\eta^2}$.

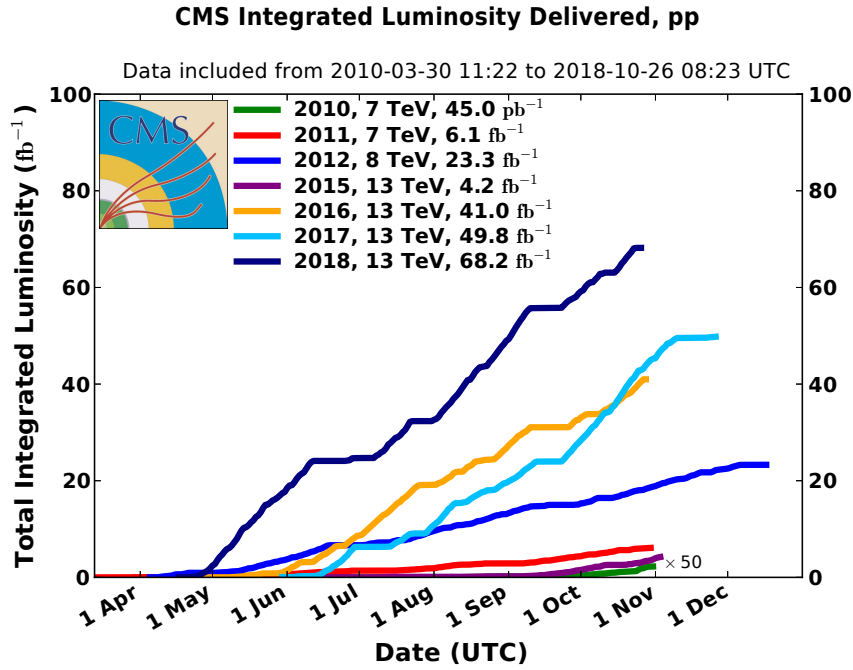


Figure 2.2: The total integrated luminosity for proton-proton collisions delivered to CMS by the LHC between 2010–2018 [68].

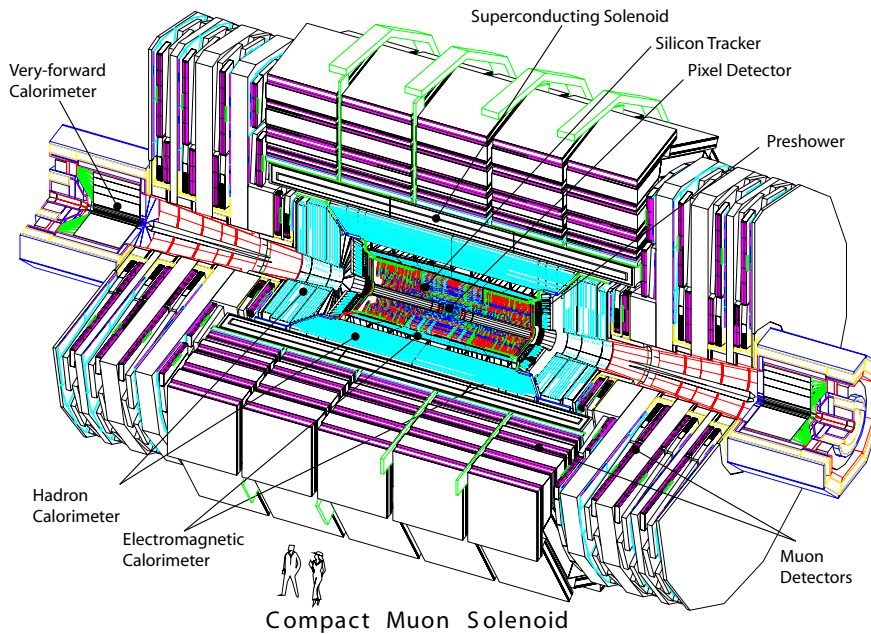


Figure 2.3: A perspective view of the CMS detector showing the layout of main sub-detector components [65].

2.3 Superconducting magnet

A high magnetic field strength is key to providing a large bending power and consequently precise momentum measurements of charged particles. In order to achieve this CMS uses a superconducting solenoidal magnet [65], operated at 3.8 T. The solenoid is 12.5 m long with a 6.3 m free bore. The coil is made up of over 2000 turns of stabilised reinforced niobium-titanium wire. Liquid helium is used to cool the magnet to its operating temperature of about 4.5 K. The flux is returned through a

10000 tonne iron yoke consisting of five barrel wheels and two endcaps. Each wheel is made up of three concentric iron rings, and each endcap consists of three iron disks.

2.4 Tracker

The sub-detector located closest to the interaction point is the tracker [65, 69, 70]. It is used to measure the trajectories of charged particles and precisely determine the position of primary and secondary vertices. At design luminosity $\mathcal{O}(1000)$ particles traverse the tracker for each LHC bunch crossing requiring a tracking detector which has high granularity and fast response. Silicon technology was chosen in order to meet these requirements, while at the same time being radiation-hard in order to withstand damage caused by the large flux of incoming particles. The CMS tracker is comprised of silicon pixel and strip detectors, a schematic of the tracking detector is shown in Figure 2.4.

The pixel detector spans the pseudorapidity range $|\eta| < 2.5$. The original pixel detector consisted of three 53-cm-long cylindrical pixel detector modules surrounding the interaction point at radii 4.4, 7.3, and 10.2 cm with two disk pixel modules at each side with longitudinal positions ± 34.5 cm and ± 46.5 cm extending in radius from ≈ 6 cm to ≈ 15 cm. In total 66 million pixels covered an area of approximately 1 m^2 . Each pixel had dimension $100 \times 150 \mu\text{m}^2$ and provided a spacial resolution of $15\text{--}20 \mu\text{m}$ in both the $r\text{--}\phi$ plane and the z direction, allowing for 3D vertex reconstruction. The original detector was designed to handle luminosities up to $1 \times 10^{34} \text{ cm}^2\text{s}^{-1}$ and an average of 25 inelastic interactions per bunch crossing, referred to as event pile-up (PU).

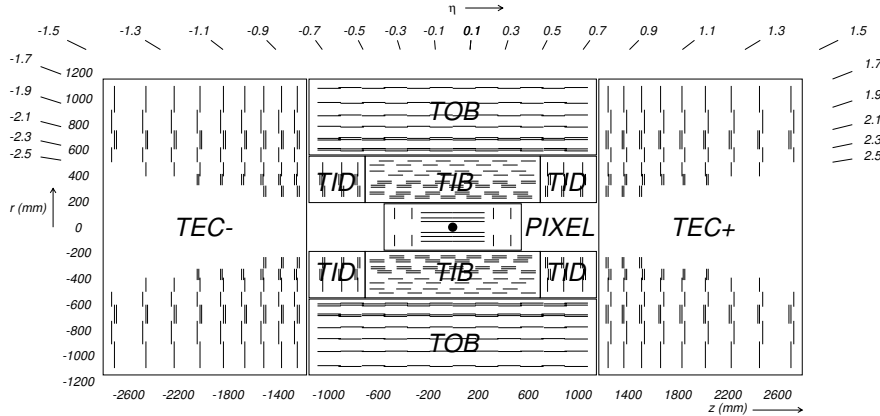


Figure 2.4: Schematic of the CMS tracker (pre- pixel upgrade). Each line represents a detector module and double lines indicate back-to-back modules delivering stereo hits [65].

In the 2016/2017 end of year technical stop the pixel detector was upgraded to ensure high tracking performance in luminosities up to $2 \times 10^{34} \text{ cm}^2\text{s}^{-1}$ and PU up to and exceeding 50 [70]. The diameter of the CMS beam pipe was reduced to allow the placement of the inner pixel layer closer to the interaction point. The three barrel modules were replaced with four module layers at radii of 3.0, 6.8, 10.2 and 16.0 cm. The two endcap disks were also replaced with three disks with longitudinal position ± 29.1 cm, ± 39.6 cm, and ± 51.6 cm extending in radius from 4.5 cm to 16.1 cm. The upgraded pixel detector consists of 124 million pixels with dimensions $100 \times 150 \mu\text{m}^2$. Figure 2.5 shows a comparison of the layouts of the original and upgraded pixel detectors.

Surrounding the pixel detector is the silicon strip detector, comprised of four subsystems [65]. The tracker inner barrel (TIB) and tracker inner disks (TID) consist of four barrel layers and three endcap disks which extends in radius between 20 cm to 55 cm. The TIB/TID use $320 \mu\text{m}$ thick

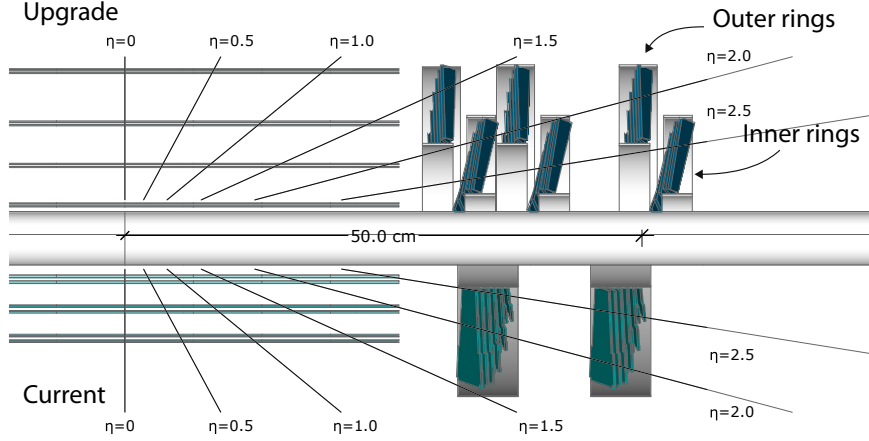


Figure 2.5: Comparison of the layout of the original and upgraded pixel detectors. The lower portion of the diagram labelled "current" shows the pixel detector used prior to the upgrade in 2016, while the upper portion labelled "upgrade" shows the pixel detector used after the 2016 upgrade [70].

silicon micro-strip sensors arranged parallel to the beam axis in the barrel and radial to the beam axis on the disks. The strip pitch for the TIB is 80 and 120 μm on the inner and outer two layers respectively, leading to a single point resolution between 23–35 μm . For the TID the mean strip pitch varies from 100 to 141 μm . The TIB/TID delivers up to four r - ϕ measurements.

Surrounding the TIB/TID is the tracker outer barrel (TOB) subsystem. It extends in the z direction to ± 118 cm and has an outer radius of 116 cm. The TOB consists of six barrel layers using 500 μm thick strips with pitches of 183 μm and 122 μm on the inner four and outer two layers respectively. The strips are arranged parallel to the beam axis providing up to six r - ϕ measurements with a single point resolution between 35–53 μm .

Located at each end of the TID/TOB are the tracker endcaps (TEC) which range between 124 cm $< |z| < 282$ cm and extend in radius between 22.5 to 113.5 cm. Each TEC is composed of nine disks each carrying up to seven concentric rings of micro-strip sensors. The inner four and outer three disks use 320 and 500 μm thick strips respectively, with an average pitch of 97–184 μm . Each TEC provides nine ϕ measurements per trajectory.

In the first two layers of the TIB and TOB, the first two rings of the TID, and rings 1, 2, and 5 of the TEC a second micro-strip detector module is mounted back-to-back with stereo angle 100 mrad to provide additional measurements of the z coordinate in the barrel and the r coordinate on the disks. The single point resolution on these measurements is 230 μm and 530 μm in the TIB and TOB respectively. The resolution on r coordinate measurement in the TID and TEC varies with the pitch. The tracker layout ensures at least three hits in the pixel detector (at least four hits for the upgraded detector) and nine hits in the silicon-strip tracker, with at least four of the latter being 2D measurements.

2.5 Electromagnetic calorimeter

The CMS electromagnetic calorimeter (ECAL) [65, 71] is a hermetic homogeneous calorimeter used to detect high energy electrons and photons. It consists of 75848 lead tungstate (PbWO_4) crystals arranged in a barrel and endcap configuration. In addition to its radiation-hardness,

PbWO_4 was chosen as the crystal material for its high density, short radiation length, and small Molière radius. These properties allow the construction of an ECAL which is compact while maintaining a fine granularity. Scintillation light emitted in the crystals by showering electrons and photons is converted into an electrical signal by photodetectors. Avalanche photodiodes and vacuum phototriodes are used as photodetectors in the barrel and endcaps respectively. The scintillation decay time of the crystals is of the same order as the 25 ns LHC bunch crossing time ensuring that $\approx 80\%$ of the light is emitted between bunch crossings.

The CMS ECAL is positioned outside of the tracker, inside the bore of the magnet, and spans the pseudorapidity range $|\eta| < 3.0$. Figure 2.6 shows a schematic of the detector layout. The ECAL barrel (EB) covers the region $|\eta| < 1.479$ and has 360-fold granularity in ϕ and (2×85) -fold granularity in η , giving total of 61200 crystals. The crystals have a tapered shape with a front face area of 0.0174×0.0174 in η - ϕ ($22 \times 22 \text{ mm}^2$) and a length of 230 mm, corresponding to 25.8 radiation lengths. The front faces of the crystals are positioned at a radius of 1.29 m from the beam axis. The crystals are aligned with their axes tilted by 3° with respect to the vector from the nominal collision point, in both ϕ and η projections, to prevent the gaps between crystals being aligned with particle trajectories.

The ECAL endcaps (EE) cover the region $1.479 < |\eta| < 3.0$. Each endcap contains 7324 crystals arranged in a rectangular x - y grid. The longitudinal positions of the endcap envelopes from the nominal interaction point is $\pm 315.4 \text{ cm}$. The crystals have a front face area of $28.62 \times 28.62 \text{ mm}^2$ and a length of 220 mm, corresponding to 24.7 radiation lengths. Sampling calorimeters known as Preshower detectors are placed at the front of each endcap. These consist of two layers of lead radiators, to initiate showering, with silicon-strip sensors placed after each radiator. The total thickness of the Preshower detectors is 20 cm, corresponding to about three radiation lengths. The main purpose of the preshower detectors is to identify photons originating from neutral pion decays in the region $1.653 < |\eta| < 2.6$ but they also aid in electron identification and position resolution.

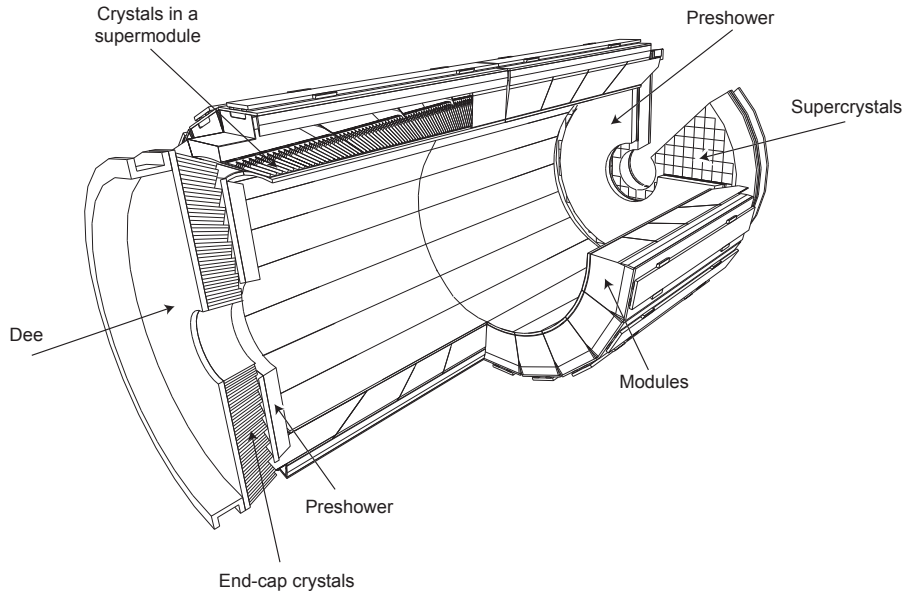


Figure 2.6: The layout of the CMS ECAL showing the positions of the crystal modules, endcaps, and Preshower system [65].

The ECAL energy resolution can be parametrized as a function of the energy of the incident

particle, E , as

$$\frac{\sigma}{E} = \frac{S}{\sqrt{E}} \oplus \frac{N}{E} \oplus C, \quad (2.2)$$

where S is the stochastic term, N is the noise term, and C is the constant term. The stochastic term encapsulates fluctuations in the lateral shower containment and the photon yield; the noise term is due to noise from electronics, digitisation, and PU; and the constant term originates from non-uniformity of the longitudinal response, leakage of energy from the back of the crystal, and inter-calibration errors. These terms have been measured in an electron beam test as $S = 0.028 \text{ GeV}^{1/2}$, $N = 0.12 \text{ GeV}$, and $C = 0.003$ using electrons with energies between 20 to 250 GeV [65].

2.6 Hadronic calorimeter

Located outside of the ECAL and spanning the pseudorapidity region $|\eta| < 5.2$ is the hadronic calorimeter (HCAL) [65, 72]. The HCAL is a sampling calorimeter used to detect strongly interacting particles which is important for measurements of hadronic jets and missing energy signals. The HCAL consists of four sub-detectors, the layout of the sub-detectors is shown in Figure 2.7.

The hadron barrel (HB) and hadron endcap (HE) calorimeters are located within the bore of the magnet covering the pseudorapidity regions $|\eta| < 1.3$ and $1.3 < |\eta| < 3.0$ respectively. These detectors consist of alternating layers of brass absorber plates inter-spaced with layers of plastic scintillator. The choice of brass for the absorber material is motivated by its relatively short nuclear interaction length and it being non-magnetic, the latter is necessary as the HB/HE are contained within the solenoid of the magnet. The HB provides a minimum of 5.82 interaction lengths of absorber material at $\eta = 0$ which increases as $1/\sin\theta$ to a maximum of 10.6 interaction lengths at $|\eta| = 1.3$. The HE provides ≈ 10 interaction lengths of absorber material. The plastic scintillator is divided into tiles with η - ϕ areas of 0.087×0.087 for $|\eta| < 1.6$ and $\approx 0.17 \times 0.17$ for $|\eta| > 1.6$. Charged particles from the hadronic showers emit scintillation light in the tiles, which is then collected by wavelength shifting fibres and read out using hybrid photodiodes.

In the central detector region ($|\eta| < 1.3$) the combined stopping power of the EB and HB systems does not provide sufficient containment of hadronic showers. To remedy this problem the hadron outer (HO) calorimeter is used to extend the HCAL outside of the solenoid coil. The HO uses tiles of plastic scintillator, with η - ϕ areas of 0.087×0.087 , placed in between the solenoid coil and iron return yoke. The barrel section of the return yoke consist of five 2.536 m wide rings. A layer of tiles is positioned inside of each of these rings. For the central ring, an additional layer of tiles is used with a 19.5 cm thick layer of iron sandwiched between the two tile layers. The magnet coil as well as the additional iron layer for the central ring serve as the absorbers for the calorimeter. The depth of the HCAL is increased to a minimum of 11.8 interaction lengths by including this sub-system. The HO provides additional sampling of hadronic showers outside of the HB, which is important for measuring late starting or highly-penetrating showers.

The hadron forward (HF) detector consist of two calorimeters positioned $\pm 11.2 \text{ m}$ from the interaction point. This extends the pseudorapidity range of the HCAL to $|\eta| < 5.2$. For these forward regions of the detector the flux of incoming particles is extremely large, with $\approx 8 \times$ as much energy deposited as for the rest of the detector. Consequently, the HF needs to be built from materials which are extremely radiation-hard and therefore quartz fibres are used as the active material, which are embedded in steel absorbers. The absorbers consist of 165 cm long and 5 mm thick grooved plates, corresponding to ≈ 10 interaction lengths of absorber material. The fibres are situated along the grooves, half of which run along the full length of the absorber and half of which

start at a depth of 22 cm from the front of the detectors. This arrangement allows for differentiation between electrons/photons and hadrons, as the former deposit most of their energies in the first 22 cm of the absorbers whereas the latter produce roughly equal deposits in the first 22 cm and the remaining 143 cm of the absorber. A signal is generated when charged particles contained in the showers emit Cherenkov light in the quartz fibres, which is then read out by photomultiplier tubes.

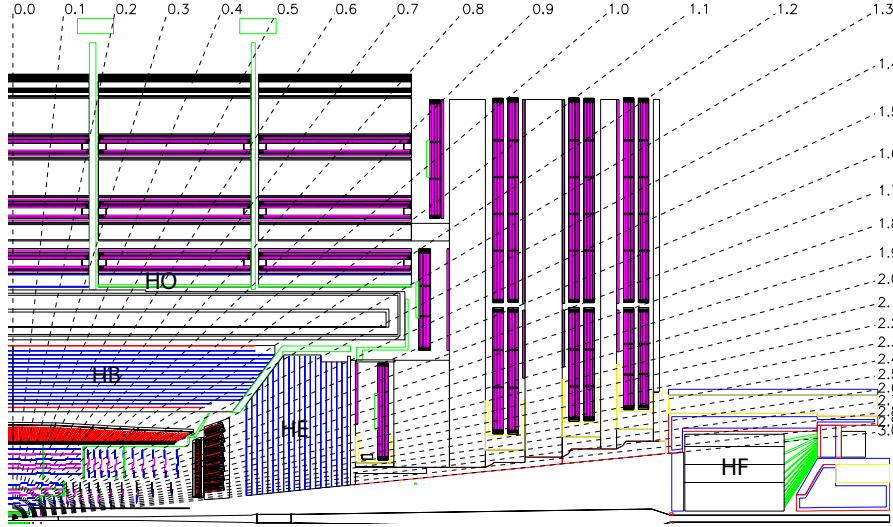


Figure 2.7: Longitudinal view of the CMS HCAL showing the positions of the HB, HE, HO and HF calorimeters [65].

Several upgrades to the HCAL photo-detectors and electronics are planned for the future, as described in Ref. [73]. As part of this program the photomultiplier tubes used by the HF were upgraded during LS1, replacing the previous single-channel phototubes with multi-anode phototubes.

The HCAL energy resolution can be parametrized as a function of the energy of the incident particle, E , as

$$\frac{\sigma}{E} = \frac{S}{\sqrt{E}} \oplus C, \quad (2.3)$$

where the stochastic and constant terms have been measured in a test beam as $S = 0.943 \text{ GeV}^{\frac{1}{2}}$ and $C = 0.084$ [74].

2.7 Muon system

The CMS muon system [65, 75] was designed to have three functions: efficient muon identification, precise momentum measurement, and triggering capability. These functions are enabled by the high-field solenoidal magnet and flux-return yoke [65]. The muon system is located outside of the solenoid coil and HO calorimeter, inter-spaced between the iron plates forming the return yoke of the magnet. The muon system spans the pseudorapidity range $|\eta| < 2.4$ and is comprised of three types of gaseous detectors: drift tube (DT) chambers, cathode strip chambers (CSC), and resistive plate chambers (RPC). The layout of one quadrant of the CMS detector, showing the position of the muon detectors, is shown in Figure 2.8

The DT chambers provide coverage for the $|\eta| < 1.2$ region and are organised into four stations. These include inner and outer stations which are positioned inside and outside of the magnet return yoke and two stations inter-spaced between the three iron layers of the yoke. Each

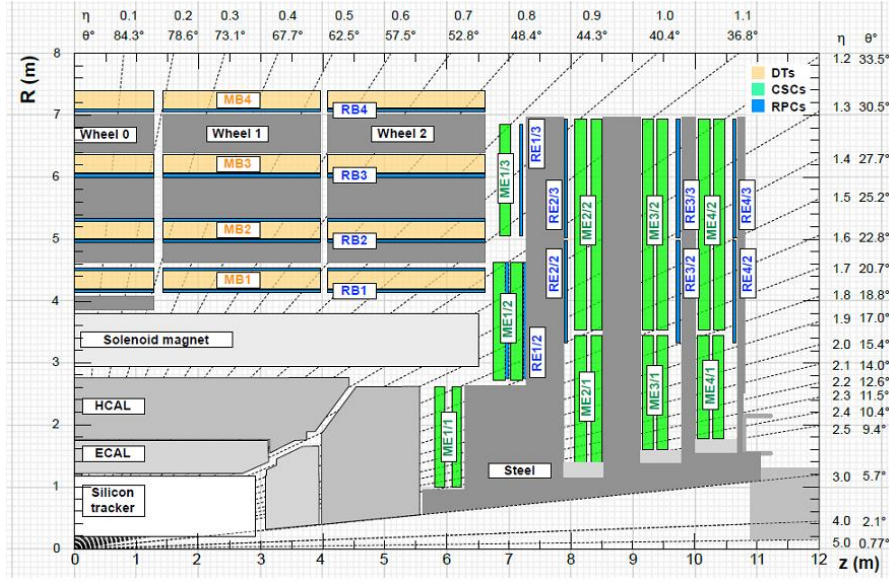


Figure 2.8: The layout of one quadrant of the CMS detector showing the DT, CSC, and RPC components of the muon system [76]. The DT chambers shown in yellow span the region $|\eta| < 1.2$. The CSC chambers are shown in green and cover the region $0.9 < |\eta| < 2.4$. The RPCs shown in blue augment the DTs and CSCs in the region $|\eta| < 1.6$.

DT chamber is made up of a series of rectangular drift cells. Each cell has a cross section of 13×42 mm² and is about 2.4 m long. The cells contain an anode wire running centrally along the length of the cell, and are filled with a mixture of argon and carbon dioxide gas. The cells are arranged into superlayers, each containing four layers of cells. The DT chambers in the three innermost muon stations contain three superlayers. The two outer superlayers are arranged such that the anode wires are parallel to the beam axes, providing a measurement of the ϕ coordinate, and the inner superlayer is positioned with the wires perpendicular to the beam axes to measure the z coordinate. For the outermost station the DT chambers contains only two superlayers with wires arranged parallel to the beam axis and therefore do not measure the z coordinate. The spacial resolution per chamber of the DTs for the ϕ and z measurements are 77–123 μ m and 133–393 μ m respectively [77].

The CSCs cover the region $0.9 < |\eta| < 2.4$. The larger flux of muons coupled with the higher radiation levels in this region motives the usage of CSCs which are radiation resistant, have a fast response, and have a fine segmentation. The CSC are trapezoidal-shaped multiwire proportional chambers containing panels dividing the chamber into six layers of gas. The gas mixture used includes argon, carbon dioxide and tetrafluoromethane. Layers of cathode strips on the panel surfaces run radially outwards providing measurements of the muon ϕ coordinate. Anode wires in between the panels run perpendicular to provide a measurement of the r coordinate. Each chamber contains six layers of anode wires interleaved among seven layers of cathode strips. The CSCs are organised into four stations inter-spaced between the three endcap layers of the return yoke of the magnet. The spacial resolution per chamber of the CSCs is 45–143 μ m [78].

The DTs and CSCs are augmented with RPCs in the $|\eta| < 1.6$ region. The RPCs are gaseous parallel-plate detectors. Each RPC is a double-gap module consisting of two layer of gas each separated by anode and cathode plates. The anodes of each stack are positioned back-to-back with an array of copper strips running parallel to the beam axis situated in between. A total of six layers of RPC chambers are embedded the barrel iron yoke. This includes two layers in each of the

two innermost muon stations and one in each of the two outer stations. In the original detector design three layers of RPCs were embedded in each endcap iron yoke, which was increased to four layers during LS1 [76]. The layers are positioned at either side of the three iron disks that make up the yoke endcap. The spacial resolution per chamber of the RPC is 0.78–1.38 cm [78]. While this resolution is poor compared to the resolution of DTs and CSCs the response time is very fast. This allows them to be used as a dedicated muon trigger with the capability to correctly identify the bunch crossing from which a muon originates.

The best muon momentum resolution is obtained by combining information from the tracker with the muon systems. The momentum resolution of the standalone muon systems is about 9% for muons with p_T up to 200 GeV, and varies with η between 15–40% for 1 TeV muons.

2.8 Triggering and computing

At design operation the LHC collides protons at a rate of 40 MHz. This rate is too high for every event to be read out by CMS, and an infeasible amount of storage capacity would be required to write it all out. Therefore, a dedicated trigger system is used to preferentially select the most interesting events, reducing the recorded rate to $\mathcal{O}(1 \text{ kHz})$.

CMS uses a two stage trigger system [65, 79, 80]. The first stage is the Level-1 (L1) trigger, which is a hardware-based trigger consisting of custom-built programmable electronics. The event rate is reduced to about 100 kHz by this trigger. The L1 system was upgraded during the end of year technical stop in 2015/2016 to help cope with the increase in instantaneous luminosity and PU [80]. The electronics for the calorimeter trigger, muon trigger, and global trigger were upgraded, and additional interconnections between the three systems was added. This has led to the following improvements in the L1 trigger capabilities:

- Improved muon momentum resolution and inclusion of muon isolation.
- Improved electron/photon isolation.
- Improved hadronic tau identification and isolation.
- Improved jet finding algorithms with PU subtraction.
- Improved position resolution for electrons/photons, taus, and jets.
- Improved global L1 trigger menu with a greater number of triggers including more complicated multi-object relations such as invariant mass cuts.

The upgraded L1 trigger uses a time-multiplexed architecture [80]. An overview of the data flow is shown in Figure 2.9. The L1 trigger uses information from the calorimeters and muon systems. Energy deposits recorded in the ECAL and HCAL which are passed to the first layer of the calorimeter trigger (Calo Trigger Layer 1) consisting of several FPGA cards. Each card maps onto a column of the detector, but receives data for several bunch crossings. The data is then passed to the second layer of the calorimeter trigger (Calo Trigger Layer 2) such that each FPGA in this layer receives data for the entire calorimeter, but for only one bunch crossing. Objects are then identified and a sorted list of the best candidates (those with the highest p_T) are passed to the global trigger. At the same time recorded hits from the three muon detectors are passed to the first layer of the muon trigger (Muon Tracking-Finder Layer), where muon tracks are reconstructed after combining hits from the DTs, CSCs, and RPCs. The track finding is performed regionally in sectors of ϕ and three regions of η (barrel, endcap and overlap). The second layer of the muon trigger (Sorting/Merging Layer) then combines the ϕ sectors. The output is passed to the global

muon trigger along with calorimeter information from Calo Trigger Layer 2 needed to compute isolation. The η regions are then combined and a list of the top muon candidates for the entire detector is passed to the global trigger. The global trigger finally combines the information from the calorimeter and muon triggers and makes a decision on whether the event is passed to the next stage of the CMS trigger.

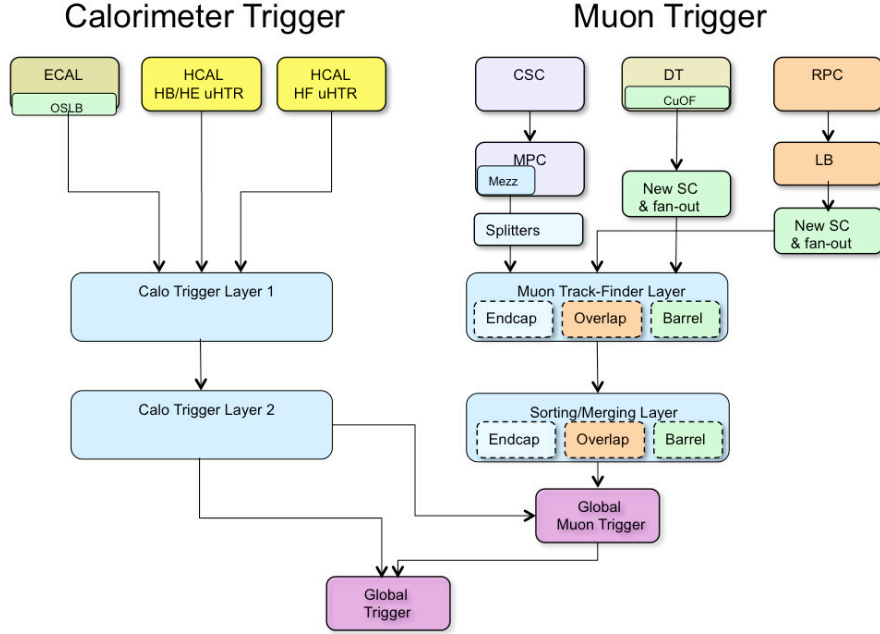


Figure 2.9: Schematic of the L1 trigger system showing the data flow for the reconstruction of objects [80].

The second stage is the high-level trigger (HLT) [79, 81], which is a software based trigger operated at a multi-processor farm including over 13000 CPU cores. The HLT uses information from the full detector, including the tracker, calorimeters, and muon systems. Having access to this information allows the algorithms and selections used in the HLT to closely follow the ones used offline. This means that objects at HLT level typically have good momentum resolution and high identification efficiencies. Of the $\mathcal{O}(100 \text{ kHz})$ of events received from the L1 trigger the HLT reduces the output rate to about 1 kHz.

The events selected by the HLT produce $\mathcal{O}(10 \text{ PB})$ of data every year which needs to be processed, stored, and made easily retrievable by analysts worldwide. To achieve this CMS, in collaboration with the other LHC experiments, operates a global computing infrastructure known as the Worldwide LHC Computing Grid (WLCG) [82]. The WLCG combines the computing resources of participating research institutions from around the world into a computing infrastructure organised into several tiers. Tier-0 includes two computing centres, the CERN Data Center and the Wigner Research Centre in Budapest. The data is copied at these sites before being processed and distributed to Tier-1 sites. Tier-1 consists of 13 large computing centres and is responsible for the storage of raw, reconstructed, and simulated data, as well the distribution of data to Tier-2 sites. Tier-2 sites are typically universities and other scientific institutions, totalling around 160 sites globally. These sites handle the production and reconstruction of simulated data and provide computing resources for analysts.

Chapter 3

Object reconstruction and analysis tools

This chapter reviews the tools that are common to the two analyses described in this thesis. This begins with a description of the various algorithms used to reconstruct physics objects in Section 3.1. Section 3.2 then outlines the datasets used in the analyses and the triggers used to record the events. Sections 3.3 and 3.4 describe the simulation of proton-proton collision events and the procedures used to correct the simulated events to improve the description of the data. Finally, the methods used to model backgrounds are detailed in Section 3.5

3.1 Object reconstruction

The reconstruction of taus is particularly challenging due to the diverse assortment of particles which they decay into. This includes charged and neutral hadrons, electrons, muons, and neutrinos. Excellent performance is required for all algorithms used to reconstruct the visible decay products, as well as the missing transverse energy which is used to infer the presence of the undetectable neutrinos. In addition to these objects, jets are powerful tools for differentiating between different Higgs production mechanisms for both SM and BSM scenarios. Information about the Higgs $\mathcal{CP}$ properties is also extracted from the GF jet topology and thus the accurate and efficient reconstruction of jets is essential. This section describes the algorithms used to reconstruct physics objects in the CMS detector. Emphasis is given to the objects used for the analyses described in this thesis.

3.1.1 Tracks and vertices

CMS reconstructs particle tracks using the Combinatorial Track Finder (CTF) algorithm [83]. This algorithm performs four steps:

1. Track seeds are formed using hits in the first few layers of the tracker (two or three hits). A seed defines an initial estimate of a trajectory's parameters and their uncertainties.
2. The track seeds are extrapolated using a Kalman filter (KF) [84]. This extrapolates outwards along the expected trajectory and searches for additional hits along the flight path in successive detector layers. At each layer associated hits are added to the track candidate and the trajectory parameters are updated. This process continues until the final detector layer is reached.
3. The final trajectory is fitted using a KF and smoother. The track fitting is performed iteratively. After each iteration spurious hits are searched for and removed. This continues

until no more spurious hits are found. This fit provides the best possible estimates of the trajectory parameters.

4. Tracks are finally required to pass a set of quality criteria and any failing tracks are discarded.

Six iterations of the CTF algorithm are performed, with the hits associated with a reconstructed track removed after each iteration. The first iterations start with easy to identify tracks, such as those with high p_T and/or close to the interaction region. The latter iterations are then used to reconstruct more difficult tracks, such as those originating far from the interaction region, and recover tracks missed by previous iterations. By removing associated hits after each iteration the combinatorial complexity is reduced making it easier to search for the more complicated tracks in the latter iterations. The tracking efficiencies for muons and pions with $p_T > 500$ MeV have been measured using $\sqrt{s} = 7$ TeV collision data and found to be $> 99.3\%$ and $> 98.5\%$ respectively [85].

After all tracks have been reconstructed it is possible to determine the positions of all interaction vertices, including the primary (signal) vertex as well as additional PU vertices. The reconstruction of vertices starts by selecting tracks that are consistent with being promptly produced close to the centre of the luminous region, which is known as the beamspot. Tracks are then clustered into groups that appear to originate from the same vertex, based on the z coordinate of the point of closest approach to the beamspot. The clustering is performed using a deterministic annealing algorithm [86], which identifies candidate vertices and assigns tracks to them. Candidate vertices are retained if at least two of their associated tracks are incompatible with other vertices. Candidate vertices are then fitted using an adaptive vertex fitter [87] to determine the best estimate of the 3D vertex positions. The primary interaction vertex is taken as the vertex with the largest value of the summed physics-object p_T^2 . The physics objects in this case are defined as jets clustered from tracks associated with the candidate vertex using the anti- k_T algorithm [88] (which will be described in Section 3.1.5), and the associated missing transverse energy given as the negative vector sum of the p_T of those jets [89]. The vertex reconstruction efficiencies were measured using $\sqrt{s} = 7$ TeV collision data and found to be $> 98.7\%$ for all vertices containing at least two tracks, and $> 99.9\%$ for vertices with four or more tracks [83]. The resolution of the primary vertex x and z coordinates has been measured using $\sqrt{s} = 13$ TeV data data, as shown in Figure 3.1 as a function of the summed track p_T [90]. The position resolution improves rapidly as the summed p_T of the associated tracks increases.

3.1.2 Particle flow

The reconstruction of all stable particles in an event, which includes electrons, photons, muons, and hadrons, is performed using the particle flow (PF) algorithm [91–93]. This algorithm combines the information from all of the CMS sub-detectors in order to achieve optimal measurements of particle energies, directions, and types. The output of the algorithm is a list of particles that may then be used to build higher-level objects such as jets, hadronically decaying taus, and missing transverse energy E_T^{miss} .

The PF algorithm starts with a set of elements which includes tracks and calorimeter clusters. Calorimeter clusters are needed to measure the energies and directions of neutral particles, to separate the energy deposits from neutral and charged hadrons, and to identify bremsstrahlung photons originating from electrons. Clusters are built using a dedicated clustering algorithm, driven by the need to reconstruct low energy particles and close-by energy deposits, which are typical of jet constituents. No clustering is performed for the HF sub-detector, instead each cell gives rise to one cluster. For all other sub-detectors, the clustering is performed separately in each sub-detector

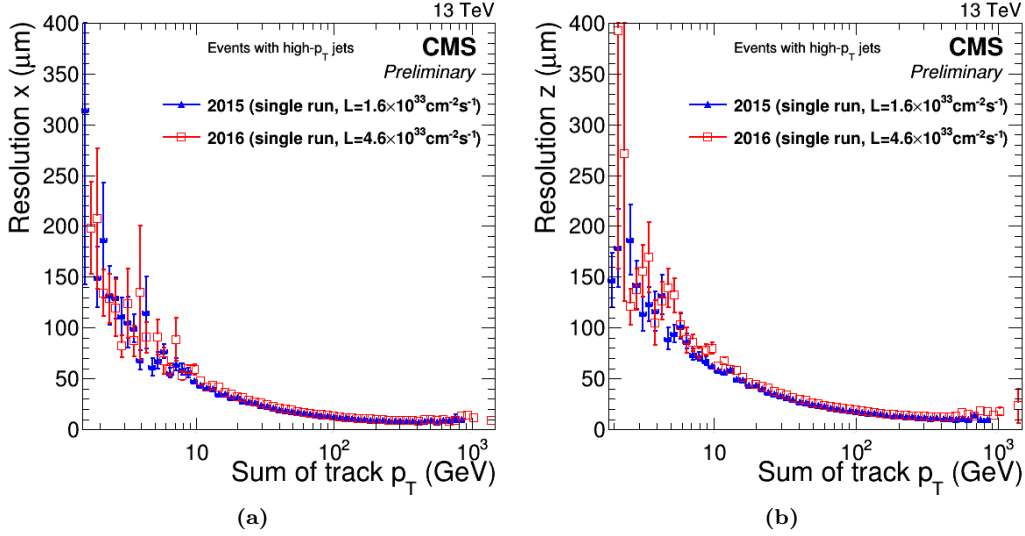


Figure 3.1: The primary vertex position resolution in x (left) and z (right) as a function of the summed track p_T of the associated tracks. The measurements are shown separately for collision data collected in 2015 (blue) and 2016 (red) data-taking periods [90].

component. Clustering begins with "cluster seeds" consisting of cells of local energy maxima above a given threshold. "Topological clusters" are then grown by aggregating neighbouring cells into the seeds provided their energies are at least two standard deviations above the expected electronics noise levels. Each topological cluster gives rise to as many "particle flow clusters" as there are cluster seeds. The energy of each cell may therefore be shared between several particle flow clusters, where the share of the cell's energy received by each cluster depends on the cell-cluster distance.

A linking algorithm is performed for each pair of elements, which defines a distance parameter between them to quantify the quality of the link. To find links between tracks and clusters the tracks are first extrapolated: to the two layers of the Preshower detector, to the expected maximum of a typical longitudinal electron shower profile in the ECAL, and to a depth of one interaction length in the HCAL. A track is linked to a cluster if the extrapolated track position is within the boundaries of the cluster. At each intersection point between the tracker layers and the track, tangents to the tracks are also extrapolated to the ECAL. If an extrapolated tangent point is within an ECAL cluster it is linked as a potential bremsstrahlung photon. The cluster boundaries can be enlarged by up to the size of a cell in each direction to account for gaps between cells/modules and uncertainties on the positions of the shower maxima. The link distance parameter is defined as the (η, ϕ) distance between the extrapolated track/tangent point and the centre of the cluster. A link is made between two calorimeter clusters if the position of the more granular cluster is within the boundaries of the less granular calorimeter. The link distance parameter in this case is defined as the (η, ϕ) distance between the two cluster positions. The cluster boundaries can again be enlarged as described previously. A link is made between a track in the inner tracker and a track in the muon system if a global fit between the two tracks returns an acceptable χ^2 . The result of such a link is known as a global muon. If multiple links are possible between a given muon track and several tracker tracks, only the link that gives the smallest χ^2 is retained. The linking distance in this case is given as the χ^2 value.

Next the PF algorithm combines elements that are linked directly or indirectly into blocks, each typically containing 1–3 elements. A set of particles is then identified from each block using

the following procedure:

- If the block contains a global muon this gives rise to a PF muon provided its momentum is consistent with the momentum of the tracker-only track within three standard deviations. The corresponding track is subsequently removed from the block and an estimate of the energy deposited in the calorimeters is subtracted from the clusters.
- The remaining tracks are submitted to a pre-identification stage that exploits the tendency of electrons to give rise to short tracks and lose energy by bremsstrahlung in the tracker to pre-identify candidate electron tracks. These candidate tracks are refit with a Gaussian sum filter (GSF) [94]. A multivariate discriminant including both tracking and calorimetric variables is then used to identify electrons. If the candidate passes the discriminant it is identified as a PF electron and its track and ECAL clusters are removed from the block.
- The remaining tracks give rise to PF charged hadrons, with momentum given by the track momentum and mass set to that of a charged pion. If the calorimetric energy is consistent with the track momentum within uncertainties the momentum is redefined by fitting the tracker and the calorimeter measurements. The matched cluster is subsequently removed from the block.
- If in the previous step the calorimetric energy is not consistent with the track momentum, the excess energy is attributed to overlapping photons and/or neutral hadrons. If the excess is larger than the ECAL energy a PF photon is created with energy equal to the ECAL energy. A PF neutral hadron is also created with the remaining excess energy. Alternatively if the excess is smaller than the ECAL energy only a PF photon is created with energy equal to the energy excess.
- Finally, the remaining ECAL and HCAL deposits give rise to PF photons and PF neutral hadrons respectively.

3.1.3 Muons

Muon reconstruction is performed by matching tracks in the inner tracker to tracks in the muon system [95, 96]. The tracks are first reconstructed independently in the inner tracker (as described in Section 3.1.1) and in the muon system. The reconstruction of tracks in the muon system, known as standalone muon tracks, starts from seeds consisting of groups of CSC or DT segments. These seeds are used to define an initial estimate of the trajectory, which is then extrapolated using a KF to gather hits from the DT, CSC, and RPC sub-detectors.

The matching of tracker and muon tracks follows two approaches: tracker muon reconstruction, and global muon reconstruction. The tracker muon reconstruction is an inside-out approach, which starts by selecting tracks in the inner tracker with $p_T > 0.5$ GeV and total momentum $p > 5$ GeV and extrapolating them to the muon system. A candidate qualifies as a tracker muon if the tracker track is matched to at least one muon segment. The matching is performed in a local x - y coordinate frame, where x is defined as the better measured coordinate in the r - ϕ plane and y is orthogonal to it. Tracks are matched if the difference in their x positions is less than 3 cm, or if the ratio of this distance to its uncertainty is less than 4. The global muon reconstruction is an outside-in approach which starts with standalone muon tracks and matches them to tracks in the inner detector. The matching is performed by comparing parameters of the two tracks propagated to a common surface. If the tracks are matched, a combined fit with the KF is performed using the hits from both the tracker and standalone muon tracks.

Global muons typically require muon segments in at least two muon stations, whereas tracker muons require matching to only one muon segment. The tracker muon reconstruction is therefore more efficient than the global reconstruction for muons with $p < 5$ GeV, as many of these low p muons will only penetrate through to the first muon station. The global muon reconstruction improves the momentum resolution for muons with $p_T > 200$ GeV, as shown in Figure 3.2. Below 200 GeV the momentum resolution is driven by the inner tracker.

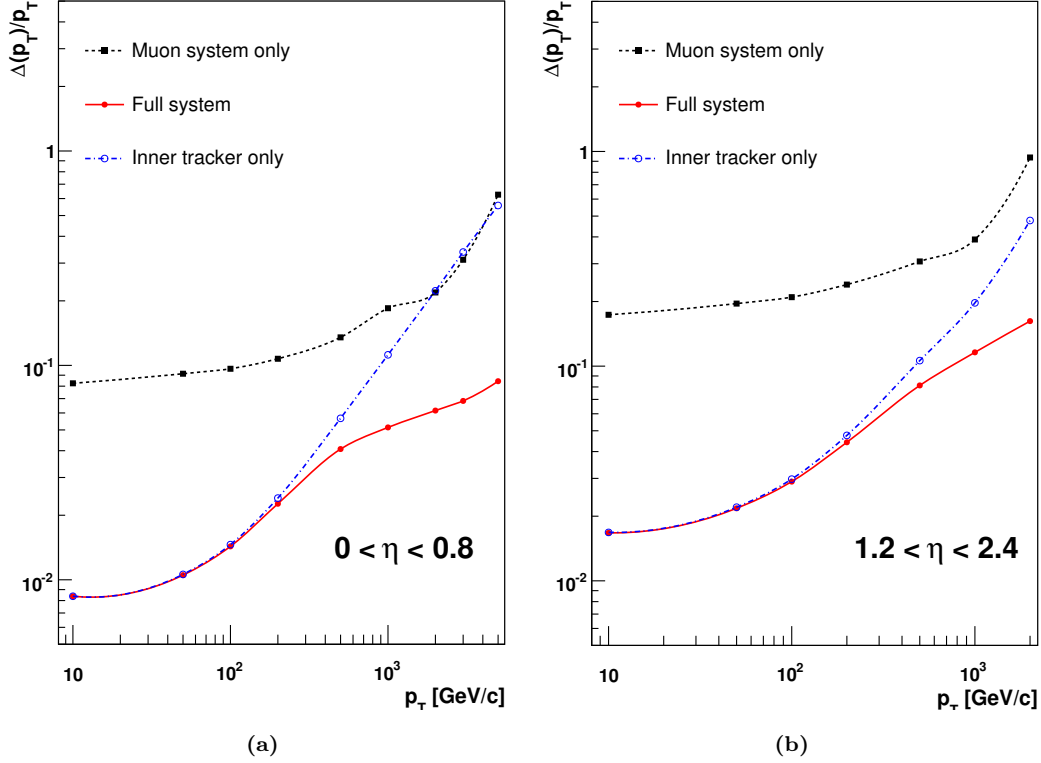


Figure 3.2: The muon p_T resolution as a function of p_T is shown for the region $|\eta| < 0.8$ (left) and $1.2 < |\eta| < 2.4$ (right), using only the muon system (black), only the inner tracker (blue), and both sub-detectors (red). [65].

Fake muons originate from hadron shower remnants that reach the muon system (hadron punch-through). This background is usually small for global muons but is larger for tracker muons, since they may contain hits only in the innermost muon station. A larger background contribution comes from real muons produced by, for example, in-flight hadron decays, and heavy-flavour decays. These backgrounds can be reduced by applying additional identification and isolation requirements. The cut based muon identification discriminant used for the analyses described in this thesis is the *medium* ID described in Ref. [96]. This discriminant aims to reduce the backgrounds from both fake muons and real muons from in-flight hadron decays. The selections used by this discriminant are as follows:

- The muon must be either a tracker or a global muon and also be identified by the PF algorithm.
- The track must include hits in more than 80% of the inner tracker layers it traverses.
- If the muon is only a tracker muon and not a global muon then the muon segment compatibility (defined in Ref. [97]) must be greater than 0.451.
- If the muon is both a tracker and a global muon then:

- the muon segment compatibility must be greater than 0.303,
- the χ^2 per degree of freedom of the global track fit must be less than 3,
- the position match between the tracker and standalone muon [97] must have $\chi^2 < 12$, and
- the maximum χ^2 computed by the kink-finding algorithm [96] must be less than 20.

The efficiency for muons with $p_T > 20$ GeV to pass the *medium* ID requirement is $> 98\%$, while the probability for hadrons to be misidentified as muons is $\mathcal{O}(0.1\%)$ [96].

The backgrounds (including real muons from heavy-flavour decays) can be further reduced by applying isolation requirements on the muons. The muon isolation is computed by summing together the energy deposited within a cone with size $\Delta R = 0.4$ centred on the muon direction, and is defined relative to the muon p_T (p_T^μ). Different isolation definitions are used in CMS. The isolation variable used in the analyses described in this thesis is defined as:

$$I_{\text{rel}} = \frac{1}{p_T^\mu} \left(\sum p_T^{h^\pm} + \max \left(0, \sum p_T^{h^0} + \sum p_T^\gamma - \Delta\beta \sum p_T^{h^\pm, \text{PU}} \right) \right) \quad (3.1)$$

[96, 98], where the sums run over all charged hadrons associated with the primary vertex ($h^\pm$), neutral hadrons (h^0), photons (γ), and charged hadrons not associated with the primary vertex ($h^\pm, \text{PU}$). The last is used to estimate the contribution of PU to the neutral energy sums. This is scaled by a factor $\Delta\beta = 0.5$ to account for the fraction of neutral to charged energy in PU interactions.

3.1.4 Electrons

Electrons are reconstructed by matching tracks with energy clusters in the ECAL. In addition to the PF algorithm described in Section 3.1.2, a standalone electron reconstruction algorithm is used to reconstruct electrons [99]. Before electrons reach the ECAL they must traverse a significant amount of material that makes up the tracker, which consists of between ≈ 0.4 to 2 radiation lengths of material depending on η . Electrons deposit a significant fraction of their energy through the emission of bremsstrahlung photons in this material. At η regions where the thickness of the material is maximal electrons lose an average of 86% of their energy. This energy is spread out predominantly in the ϕ direction due to the bending of the electron by the magnetic field.

In order to accurately measure the initial electron energy it is necessary to collect the energy of all the radiated photons. This is performed by supercluster (SC) algorithms. Two SC algorithms known as the "hybrid" and the "multi-5 $\times$ 5" algorithms perform the clustering in the barrel and endcaps respectively. Both algorithms begin with seed crystals, which are crystals of local energy maxima above a minimum threshold (1 GeV in the barrel and 0.18 GeV in the endcaps). In the hybrid method, arrays of 5×1 crystals in $\eta \times \phi$ are added around the seed, provided their energy is > 0.1 GeV. This starts with the first array centred on the seed then steps repeatedly in both ϕ directions up to $\approx \pm 0.3$ rad. Contiguous arrays are then grouped into clusters and merged into the SC if their energy exceeds 0.35 GeV. In the multi-5 $\times$ 5 method, the energy is collected around the seeds in sub-clusters of 5×5 crystal, starting with the seed with the largest energy. The sub-clusters with energy > 1 GeV, within a η range of ± 0.07 and a ϕ range of ± 0.3 rad are merged into the SC. The energy weighted positions of all sub-clusters are extrapolated to the Preshower, with the most energetic cluster used as a reference point. The maximum ϕ distance between the sub-clusters and the reference point, extended by ± 0.15 rad, defines the Preshower cluster ϕ range, with the η range set to ± 0.15 rad. The energy deposited in the Preshower within this boundary is added to the SC.

For the reconstruction of electron tracks, the standard CTF algorithm is compromised by the large radiative losses in the tracker material. These losses lead to a reduced hit-collection efficiency and a poor estimate of the trajectory parameters. This motivates the use of a dedicated tracking procedure for electrons. This procedure starts by finding seeds of two or three hits in the tracker. Two complementary methods are used to select the seeds. The first method, known as ECAL-based seeding, uses the SC energy and position to estimate the trajectory in the first layers of the tracker. Seeds are selected if they contain hits within windows around the extrapolated trajectory. The second method is a tracker-based seeding, which starts with tracks reconstructed using the CTF algorithm. If a CTF track's direction is compatible with the position of the closest PF cluster, and the ratio of the cluster energy to the track momentum is within a specified range, then the seeds from that track are selected. The seeds are then extrapolated and smoothed using the same procedure described in Section 3.1.1 but with the GSF used instead of the KF. This filter models the energy loss by bremsstrahlung using a Gaussian mixture rather than a single Gaussian (as is used by the KF), giving a better estimate of the trajectory parameters.

Electron identification discriminants are used to distinguish genuine electrons from backgrounds such as misidentified jets, electrons from photon conversions, and electrons produced in heavy-flavour decays. These include both MVA-based and cut based discriminants. The variables used in these discriminants include:

- Shower shape variables such as the width of the shower in the η direction.
- The fraction of energy deposited in the HCAL.
- The fraction of energy lost through bremsstrahlung defined as $f_{\text{brem}} = (p_{\text{in}} - p_{\text{out}}/p_{\text{in}})$, where p_{in} is the momentum at the point of closest approach to the beam spot, and p_{out} is the momentum extrapolated to the ECAL surface.
- Track quality variables.
- The variable $1/E_{\text{SC}} - 1/p$, where E_{SC} is the SC energy and p is the momentum of the track at the point of closest approach to the vertex. Distributions of this variable for genuine electrons and background electrons are shown in Figure 3.3 (a).
- Variables comparing the geometrical matching between the track and the SC, such as the $\Delta\eta$ separation between the SC and the trajectory extrapolated to its position of closest approach to the SC. Distributions of this variable for genuine electrons and background electrons are shown in Figure 3.3 (b).

Several working points on the MVA discriminant are defined using cuts on the MVA score. The analyses described in this thesis use two of these working points, with efficiencies of $\approx 90\%$ and 80% for genuine electrons for misidentification probabilities $\sim 2\text{--}5\%$ [99]. These will be referred to as the *loose* and *tight* working points, respectively. In addition to these MVA IDs a very loose cut based ID is used on occasion to veto events with leptons, which will be referred to as the *veto* electron ID.

Isolation discriminants are used to further reduce the backgrounds. The same isolation variable defined in Equation 3.1, except using a cone size of $\Delta R = 0.3$, is used for the dataset collected during 2016. For the 2017 dataset, an alternate method is used to subtract the PU contribution to the neutral p_{T} sums. The relative isolation in this case is defined as:

$$I_{\text{rel}} = \frac{1}{p_{\text{T}}^{\text{e}}} \left(\sum p_{\text{T}}^{\text{h}^{\pm}} + \max \left(0, \sum p_{\text{T}}^{\text{h}^0} + \sum p_{\text{T}}^{\gamma} - \rho A_{\text{eff}} \right) \right), \quad (3.2)$$

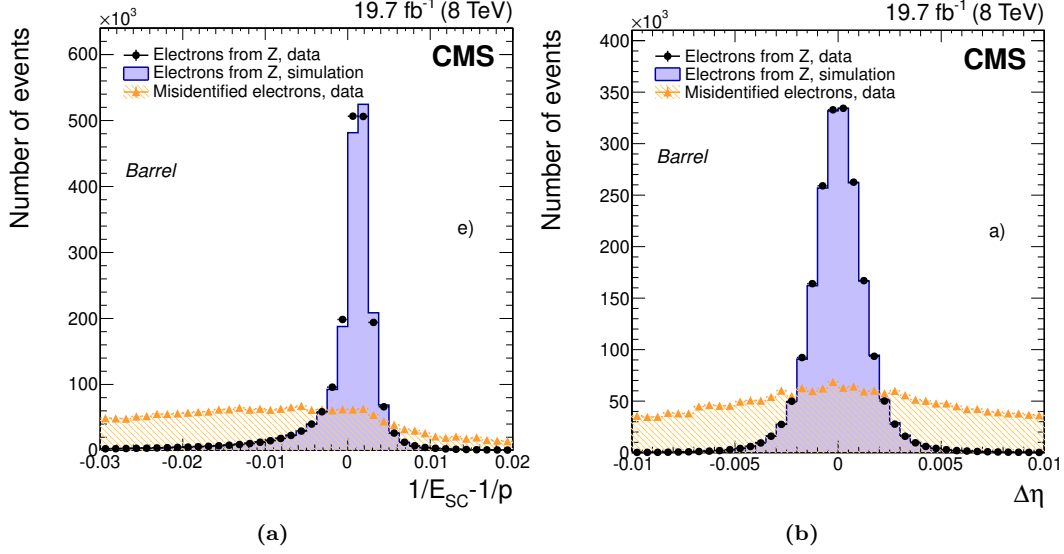


Figure 3.3: Two of the variables used to define the electron identification discriminants, $\Delta\eta$ (left) and $1/E_{SC} - 1/p$ (right), which are defined in the text. The distributions are shown for genuine electrons in simulation (solid blue) and data (black dots), and for background electrons in data (orange triangles) [99].

where the sums run over all charged hadrons associated with the primary vertex ($h^\pm$), neutral hadrons (h^0), and photons (γ). The PU is estimated by the ρA_{eff} term, where ρ is the average p_T density for particles within the area of any jet with $p_T > 3$ GeV and $|\eta| < 2.5$, clustered using the k_T algorithm [100] (which will be described in Section 3.1.5) with a distance parameter of 0.6, and A_{eff} is the effective area, as defined in Ref. [99].

3.1.5 Jets

The production of quarks and gluons is followed almost immediately by their fragmentation and hadronisation. This results in a collimated sprays of hadrons known as a jet [101]. In order to accurately measure the kinematics of the initial quark or gluon, it is crucial to be able to detect and combine a jet's constituents.

The combination of particles into jets is performed by jet clustering algorithms [101]. These algorithms begin by defining two distances parameters: d_{ij} , which is the distance between objects i and j , and d_{iB} , which is the distance between object i and the beam. The objects referred to here can be either individual particles or clusters of particles. The distance parameters are defined as:

$$\begin{aligned} d_{ij} &= \min(p_{Ti}^{2p}, p_{Tj}^{2p}) \frac{\Delta R_{ij}^2}{R^2} \\ d_{iB} &= p_{Ti}^{2p}, \end{aligned} \quad (3.3)$$

where the constant $p = -1, 0, 1$ defines the algorithm's behaviour, ΔR_{ij} is the ΔR separation of objects i and j , and R is the radius parameter that controls the typical size of the jets. The clustering is then performed as follows:

1. d_{ij} and d_{iB} are computed for all objects and the minimum distance is determined.
2. If the minimum is a d_{ij} then the objects i and j are merged into a single object.
3. If the minimum is a d_{iB} then object i is declared a jet and removed from list of objects.
4. The procedure is repeated until no objects remain.

CMS uses the anti- k_T jet clustering algorithm [88] implemented in **Fastjet** [102]. This corresponds to setting $p = -1$ in Equation 3.3, which has the tendency to cluster jets around the hardest particles. Conversely, the k_T algorithm introduced in the definition of the electron isolation in Section 3.1.4 uses $p = +1$. The anti- k_T algorithm is infrared- and collinear-safe, which means that the set of hard jets is unchanged by the emission of soft gluons or by collinear parton splittings. The analyses described in this thesis use jets clustered from PF candidates with a radius parameter of $R = 0.4$. Simulated jets are also constructed by clustering all stable generator-level particles, with the exception of neutrinos.

The effect of PU on the jet energy and substructure is partially mitigated using the charged hadron subtraction (CHS) procedure [103], whereby charged hadrons not associated with the primary vertex are not included in the jet clustering. A loose identification discriminant is also used to reject misidentified jets arising from calorimeter and/or electronics noise (noise jets). This discriminant uses variables such as the energy fractions carried by different types of PF candidates, and the charged and neutral particle multiplicities. For data collected in 2017, the $2.65 < |\eta| < 3.14$ region is affected by large amounts of ECAL noise, resulting in a significant increase in the number of noise jets. These jets have a similar shape-structure to PU jets, with both being more diffuse compared to genuine jets. Therefore, the jets in this η region are required to pass a PU jet identification discriminant [104]. This discriminant uses a combination of vertex information and jet shape information. However, as no tracking is available beyond $|\eta| > 2.5$, the jet shape variables provide the full discriminating power against PU jets, and likewise against noise jets, in the problematic η region.

3.1.5.1 Jet Energy Corrections

The energies of the jets are corrected such that they equal, on average, the energies of the corresponding true particle jets at hadron level. The correction [105] is applied to each component of the uncorrected jet four-momentum, p_μ^{raw} . The corrected jet four-momentum is calculated as:

$$p_\mu^{\text{corr}} = C_{\text{offset}}(p_\mu^{\text{raw}}) \cdot C_{\text{MC}}(p'_T, \eta) \cdot C_{\text{rel}}(\eta) \cdot C_{\text{abs}}(p''_T) \cdot p_\mu^{\text{raw}}, \quad (3.4)$$

where C_i are the correction factors described below, p_T^{raw} is the p_T of the uncorrected jet, p'_T is the p_T of the jet after applying the C_{offset} correction, and p''_T is the p_T of the jet after applying all previous corrections.

The offset correction, C_{offset} , removes the contribution of electronics noise and PU. This is estimated using the average p_T density per unit area ρ (as defined in Section 3.1.4 except including k_T jets in the full detector range $|\eta| < 5$ in this case), and the jet area A as defined in Ref. [106]. In addition to PU and noise, ρ contains also a contribution from the underlying event (UE) of the hard interaction. The latter should not be subtracted from the jet energy and so the average UE contribution $\langle \rho_{\text{UE}} \rangle$, is subtracted from ρ . The energy removed from the jet is therefore $(\rho - \langle \rho_{\text{UE}} \rangle) \cdot A$.

The MC calibration factor, C_{MC} , is based on simulation, and corrects the reconstructed jet energy to equal, on average, the energy of the generator-level jets. To determine this correction, reconstruction-level jets are matched to generator-level jets with $\Delta R < 0.25$, and the variable $\mathcal{R} = p_T^{\text{reco}}/p_T^{\text{gen}}$ is measured in bins of p_T^{gen} . The correction is then defined as the inverse of the average of $\mathcal{R}$, $C_{\text{MC}} = \frac{1}{\langle \mathcal{R} \rangle}$, expressed as a function of the average p_T^{reco} .

The relative residual correction factor, C_{rel} , corrects the response to be flat as a function of η . It is measured using a dijet p_T balancing method. A tag jet with $|\eta| < 1.3$ and probe jet with any

η , separated by $|\Delta\phi| > 2.7$ rad, are selected. The $|\eta| < 1.3$ region is chosen for its uniformity in jet response. The p_T balancing method is based on transverse momentum conservation, as the two back-to-back jets should have approximately equal p_T . The balance is defined as $\mathcal{B} = \frac{p_T^{\text{probe}} - p_T^{\text{tag}}}{p_T^{\text{ave}}}$, where p_T^{ave} is the average p_T of the tag and probe jets. The correction is determined in bins of η and p_T^{ave} by averaging the $\mathcal{B}$ distribution, $C_{\text{rel}} = \frac{2 + \langle \mathcal{B} \rangle}{2 - \langle \mathcal{B} \rangle}$.

The absolute jet energy response correction, C_{abs} , is designed to make the response flat as a function of p_T . This correction is measured using the missing transverse energy projection fraction (MPF) method, using samples of γ +jet and Z+jet events. The MPF method is based on the fact that at parton-level the p_T of the γ/Z is perfectly balanced by the p_T of the recoiling hadrons, meaning there is no intrinsic E_T^{miss} . At reconstruction-level, an imbalance due to the detector response to the γ/Z or the jets gives rise to E_T^{miss} . The measured E_T^{miss} and $p_T^{\gamma/Z}$ can therefore be used to calibrate the p_T response of the jets. The E_T^{miss} reconstruction will be described in Section 3.1.7.

3.1.6 b-jets

Jets originating from the hadronisation of b quarks (b-jets) are important signatures for many SM and BSM topologies, including the analyses described in this thesis. For the $\phi \rightarrow \tau\tau$ search, the presence of b-jets is used to tag the $b\bar{b}$ -associated production mode. A veto on events with b-jets is also used to reduce the $t\bar{t}$ background in the $\mathcal{CP}$ analysis, exploiting the near 100% branching ratio for the $t \rightarrow bW$ decay.

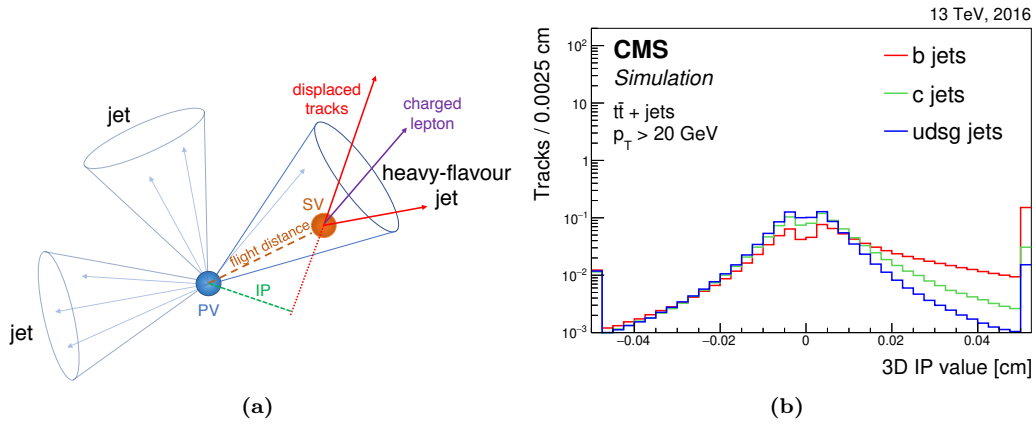


Figure 3.4: Left: Illustration of a heavy flavour decay giving rise to displaced tracks and a secondary vertex [89]. Right: The distributions of the 3D impact parameters for tracks associated with different types of jets [89]. The sign of the impact parameter is defined as the sign of the scalar product of the vector pointing from the primary vertex to the point of closest approach and the vector that defines the jet direction [107].

The b hadrons have relatively long lifetimes (≈ 1.5 ps), which leads to flight distances of between 1–10 mm from their production points [89]. This means that the tracks of the decay products are usually displaced from the primary vertex and often secondary vertices are reconstructed, as illustrated in Figure 3.4 (a). The flight distance significance of the secondary vertex, defined as the displacement from the primary vertex divided by the uncertainty, is a powerful variable for differentiating b-jets from light-flavour jets. A track’s impact parameter, defined as the distance of closest approach to the primary vertex, is also typically large, as is its significance (the impact parameter divided by its uncertainty). Distributions of impact parameters for tracks associated with different jet types are shown in Figure 3.4 (b). Algorithms used to identify b-jets often rely on

the presence of these displaced tracks and/or secondary vertices, but also may include variables sensitive to other properties of the b hadrons, such as their large masses and hard fragmentations.

The b-tagging algorithm used in the analyses described in this thesis is an upgraded version of the combined secondary vertex (CSV) algorithm [107] used in Run 1, known as CSVv2 [89]. Jets with $p_T > 20$ GeV and $|\eta| < 2.4$ are taken as inputs. This algorithm reconstructs secondary vertices using the inclusive vertex finding (IVF) algorithm, which includes all tracks with $p_T > 0.8$ GeV, impact parameters $> 50 \mu\text{m}$, and impact parameter significances > 1.2 in the vertex reconstruction. The CSVv2 algorithm then combines information relating to the displaced tracks and associated secondary vertices using a multivariate technique. Several working points are defined for the resulting discriminant. The analyses described here use the *medium* working point, which has an efficiency of $\approx 60\%$ for a misidentification probability of $\approx 1\%$ for light-flavour jets as measured in simulated $t\bar{t}$ events [89].

3.1.7 Missing transverse energy

Stable or long-lived weakly interacting particles, such as neutrinos as well as hypothetical BSM particles, pass through the CMS detector without any interactions and consequently cannot be directly observed. Their presence is instead inferred by a momentum imbalance in the transverse plane. The missing transverse momentum is defined as

$$\vec{E}_T^{\text{miss}} = - \sum \vec{p}_T^i. \quad (3.5)$$

where the sum runs over all the reconstructed PF particles in the event. The missing transverse energy, E_T^{miss} , is defined as the magnitude of $\vec{E}_T^{\text{miss}}$. The accuracy of the E_T^{miss} measurement is affected by minimum energy thresholds in the calorimeter, nonlinearity of the calorimeter response for hadronic particles, and tracker inefficiencies [108]. The impact of these effects on the E_T^{miss} measurement is reduced by propagating the jet energy corrections to the $\vec{E}_T^{\text{miss}}$ according to

$$\vec{E}_T^{\text{miss,corr}} = \vec{E}_T^{\text{miss,raw}} - \sum_{\text{jets}} (\vec{p}_{T,\text{jet}}^{\text{corr}} - \vec{p}_{T,\text{jet}}^{\text{raw}}), \quad (3.6)$$

where "corr" and "raw" refer to the corrected and uncorrected values respectively, and the sum runs over all jets with $p_T > 15$ GeV that have less than 90% of their energies deposited in the ECAL [108].

Recoil corrections are also applied to some of the simulated samples in order to improve the agreement with data. The hadronic recoil is defined as

$$\vec{U}_T = \vec{E}_T^{\text{miss}} - \sum \vec{p}_T^\nu, \quad (3.7)$$

where the sum runs over all neutrinos in the event. This quantity is determined using a sample of $Z \rightarrow \mu\mu$ events by measuring the $\vec{p}_T$ of the Z boson and the total transverse momentum of the recoiling jets, $\vec{H}_T$:

$$\vec{U}_T = -\vec{H}_T - \vec{p}_T^Z. \quad (3.8)$$

The $\vec{U}_T$ is projected onto the plane perpendicular, $\vec{U}_\perp$, and parallel, $\vec{U}_\parallel$, to the Z boson $\vec{p}_T$. The $\vec{U}_\perp$ and $\vec{U}_\parallel$ distributions are fitted using a superposition of Gaussians in data and simulation. Corrections are applied to the simulation as

$$\begin{aligned} U'_\parallel &= \langle U_\parallel \rangle_{\text{data}} + (U_\parallel - \langle U_\parallel \rangle_{\text{MC}}) \frac{\sigma_{\text{data}}(U_\parallel)}{\sigma_{\text{MC}}(U_\parallel)} \\ U'_\perp &= U_\perp \frac{\sigma_{\text{data}}(U_\perp)}{\sigma_{\text{MC}}(U_\perp)}, \end{aligned} \quad (3.9)$$

where $\langle U_i \rangle$ and σ_i are the mean values and resolutions of the fitted functions respectively. The $\vec{E}_T^{\text{miss}}$ is then recalculated using Equation 3.7.

3.1.8 Hadronic taus

Tau leptons are unstable and with a lifetime of 0.3 ps [3] decay close to the primary vertex before reaching the detector. This means that they cannot be directly detected and must instead be reconstructed from their decay products. With a mass of 1.8 GeV, taus are massive enough to decay both leptonically and hadronically. In the leptonic case, the tau decays to either an electron or muon accompanied by two neutrinos. The branching fractions for these decays are 17.8% and 17.4%, respectively [3]. The remaining 64.8% of decays are hadronic, where the tau decays to a mixture of charged and neutral mesons accompanied by a single neutrino. The hadronic decay modes are summarised in Table 3.1. For the most common decay modes, a tau will decay to either one or three charged mesons ($\pi^\pm$, $K^\pm$), and up to two neutral pions. In this document the hadronically decaying taus will be denoted as τ_h . Likewise, the leptonically decaying taus will be denoted as τ_e and τ_μ for decays into electrons and muons respectively, or τ_l when referring to the leptonic decay modes collectively.

Table 3.1: Summary of the most common hadronic tau decay modes and the corresponding branching fractions. The generic symbol $h^\pm$ represents either a $\pi^\pm$ or a $K^\pm$. Intermediate resonances are listed where relevant [3].

Decay mode	Meson resonance	Branching fraction (%)
$\tau^- \rightarrow h^- \nu_\tau$		11.5
$\tau^- \rightarrow h^- \pi^0 \nu_\tau$	$\rho(770)$	26.0
$\tau^- \rightarrow h^- \pi^0 \pi^0 \nu_\tau$	$a_1(1260)$	9.5
$\tau^- \rightarrow h^- h^+ h^- \nu_\tau$	$a_1(1260)$	9.8
$\tau^- \rightarrow h^- h^+ h^- \pi^0 \nu_\tau$		4.8
Other hadronic decay modes		3.2
Total		64.8

Hadronic tau leptons are reconstructed using the hadrons plus strips (HPS) algorithm [109, 110]. This algorithm is seeded by jets with $p_T > 14$ GeV and $|\eta| < 2.5$. For a given jet, the HPS algorithm generates all possible combinations of hadrons for the decay modes: $h^\pm$, $h^\pm \pi^0$, $h^\pm \pi^0 \pi^0$, and $h^\pm h^\mp h^\pm$. These modes cover all the largest branching fractions in Table 3.1. The jets are required to contain at least one charged hadron with $p_T > 0.5$ GeV that is compatible with originating from the primary vertex. The π^0 mesons decay almost exclusively to a pair of photons, with a branching fraction of $\approx 99\%$. These photons have a high probability of converting to e^+e^- pairs in the tracker. The resulting electrons are bent in the magnetic field dispersing them predominantly in the ϕ direction. The e/γ constituents of the jets are therefore clustered into rectangular $\Delta\eta \times \Delta\phi$ strips. To account for the e/γ candidates being more collimated for higher momentum π^0 's the size of the strips are dynamically varied as a function of p_T .

The procedure used to cluster e/γ candidates into strips is as follows:

1. The e/γ of highest p_T not clustered into a strip is used to seed a new strip. The initial η and ϕ values of the strip are set to that of the seed e/γ .
2. The e/γ with the next highest p_T within a $\eta \times \phi$ window centred on the strip is merged into

the strip. The size of the strip window is given by

$$\begin{aligned}\Delta\eta &= f(p_T^{e/\gamma}) + f(p_T^{\text{strip}}) \text{ and} \\ \Delta\phi &= g(p_T^{e/\gamma}) + g(p_T^{\text{strip}}),\end{aligned}\tag{3.10}$$

where $p_T^{e/\gamma}$ is the p_T of the candidate e/γ being merged into the strip, p_T^{strip} is the current p_T of the strip, and

$$\begin{aligned}f(p_T) &= 0.20 \cdot p_T^{-0.66} \text{ and} \\ g(p_T) &= 0.35 \cdot p_T^{-0.71},\end{aligned}\tag{3.11}$$

where the p_T is in GeV. The upper and lower limits on the strip size is set to $\eta \times \phi = 0.15 \times 0.30$ and $\eta \times \phi = 0.05 \times 0.05$ respectively. The new strip η and ϕ positions are computed as the p_T weighted average of all e/γ constituents in the strip.

3. The above stage is repeated until there are no more e/γ candidates within the $\Delta\eta \times \Delta\phi$ window. The process then continues by selecting the next seed and the above procedure is repeated until all e/γ candidates have been merged into strips.

All strips with $p_T > 2.5$ GeV are considered as π^0 candidates. The algorithm then proceeds to form all combinations of: $h^\pm$, $h^\pm\pi^0$, $h^\pm\pi^0\pi^0$, and $h^\pm h^\mp h^\pm$. Further constraints are placed on the visible mass of the decay products to ensure that $h^\pm\pi^0$ decays are compatible with the intermediate $\rho(770)$ resonance, and the $h^\pm\pi^0\pi^0$ and $h^\pm h^\mp h^\pm$ decays are compatible with the $a_1(1260)$ resonance. The mass constraints are defined as:

- $0.3 - \Delta m_{\tau_h} < m_{\tau_h} < 1.3\sqrt{p_T/(100 \text{ GeV})} + \Delta m_{\tau_h}$ GeV for the $h^\pm\pi^0$ mode, with the upper limit constrained to lie between 1.3 and 4.2 GeV,
- $0.4 - \Delta m_{\tau_h} < m_{\tau_h} < 1.2\sqrt{p_T/(100 \text{ GeV})} + \Delta m_{\tau_h}$ GeV for the $h^\pm\pi^0\pi^0$ mode, with the upper limit constrained to lie between 1.2 and 4.0 GeV, and
- $0.8 < m_{\tau_h} < 1.5$ GeV for the $h^\pm h^\mp h^\pm$ mode,

where Δm_{τ_h} is the change in the mass of the τ_h candidate brought about by the addition of the e/γ candidates to its strip as defined in Ref. [110]. For the $h^\pm\pi^\mp\pi^\pm$ decay, the three charged hadrons must also originate within $\Delta z < 4$ cm from the same vertex and have a total charge of ± 1 .

Once all possible τ_h candidates have been formed, the $h^\pm$ and π^0 candidates are required to be contained within the signal cone, with a size defined as $R_{\text{sig}} = (3 \text{ GeV})/p_T$ bounded between $0.05 < R_{\text{sig}} < 0.10$. Candidates with additional charged particles or π^0 candidates outside of this signal cone are discarded. If there are multiple τ_h candidates for the same jet remaining after these selections then the candidate with the largest p_T is retained. The $h^\pm\pi^0$ and $h^\pm\pi^0\pi^0$ decay modes are finally consolidated into the $h^\pm\pi^0$ mode to be analysed together. In this document the three tau decay modes will be referred to as the 1-prong ($h^\pm$), 1-prong+1- π^0 ($h^\pm\pi^0$ and $h^\pm\pi^0\pi^0$), and 3-prong ($h^\pm\pi^\mp\pi^\pm$) decay modes.

Even after the HPS reconstruction the jet $\rightarrow \tau_h$ fake rate is significant. Requiring the τ_h candidates to be isolated helps to reduce this background. The isolation of the τ_h candidates, I_{τ_h} , is computed by summing the scalar p_T of the charged particles and photons within a cone with size $\Delta R = 0.5$ centred on the direction of the τ_h candidate. The charged hadron and e/γ constituents of the τ_h candidates are excluded from the sum. The formula for I_{τ_h} is:

$$I_{\tau_h} = \underbrace{\sum p_T^{\text{charged}}(d_z < 0.2\text{cm})}_{(1)} - \max\left(0, \underbrace{\sum p_T^\gamma}_{(2)} - \underbrace{\Delta\beta \sum p_T^{\text{charged}}(d_z > 0.2\text{cm})}_{(3)}\right).\tag{3.12}$$

Term (1) in the above formula is the charged isolation, the contribution of PU to this term is minimised by requiring the charged particles in the sum to originate within a distance $d_z < 0.2$ cm from the τ_h production vertex. Term (2) is the neutral isolation term. The contribution of the PU to the neutral isolation is subtracted using term (3). This is estimated by summing together the p_T of charged particles not originating from the τ_h vertex, within a cone of size $\Delta R = 0.8$, and scaling this by the constant term, $\Delta\beta = 0.2$. This constant accounts for the fraction of neutral to charged energy and the difference in the two cone sizes (0.5 vs 0.8).

An MVA-isolation discriminant combines isolation variables with variables sensitive to the finite lifetime of the tau, to give the best possible separation between genuine τ_h decays and jets. The variables related to the tau isolation include:

- The charged and neutral isolation sums defined by terms (1) and (2) in Equation 3.12, as well as the PU correction defined by term (3), included as separate variables.
- The number of e/γ candidates with $p_T > 0.5$ GeV in the signal and isolation cones.
- The variable $p_T^{\text{strip, outer}}$, defined as:

$$p_T^{\text{strip, outer}} = \sum p_T^{e/\gamma} (\Delta R > R_{\text{sig}}), \quad (3.13)$$

which is the scalar p_T sum of the e/γ candidates clustered into strips but located outside of the signal cone.

- The p_T weighted ΔR , $\Delta\eta$, and $\Delta\phi$ of the e/γ candidates clustered in strips relative to the direction of the τ_h .

The variables sensitive to the τ_h lifetime include:

- The signed 3D impact parameter of the leading track of the τ_h candidate relative to the primary vertex and its significance. The primary vertex is refitted after excluding the tracks from the τ_h candidate.
- The distance between the τ_h production (primary) and decay (secondary) vertices and its significance. A flag indicating whether the secondary vertex was reconstructed or not is also included. The primary vertex is refitted as described above.

These are combined also with the tau p_T , and η , and decay mode.

The working points on the MVA-isolation discriminant are defined using p_T dependent cuts on the MVA score in order to ensure the signal efficiencies are uniform with p_T . In total six working points are defined with efficiencies between 40–90%, in steps of 10%, which are named: *vloose*, *loose*, *medium*, *tight*, *vtight*, and *vvtight*. The efficiency of genuine τ_h identification and misidentification probability for $\text{jet} \rightarrow \tau_h$ fakes is shown in Figure 3.5, for each of the six working points.

In addition to the MVA-isolation variable, further discriminants are used to reject $e \rightarrow \tau_h$ and $\mu \rightarrow \tau_h$ fakes. The anti-electron discriminant is a MVA-based discriminant using variables such as:

- Energy/momentum fractions: the electromagnetic energy fraction, the ECAL and HCAL energies relative to the track momentum, and the fraction of the τ_h energy carried by the photons.
- Photon related variables: The p_T weighted root-mean-square ϕ and η distances between each of the photons and the τ_h , the number of photons in the strips, the ratio of energies of bremsstrahlung photons measured in the ECAL and the tracker.

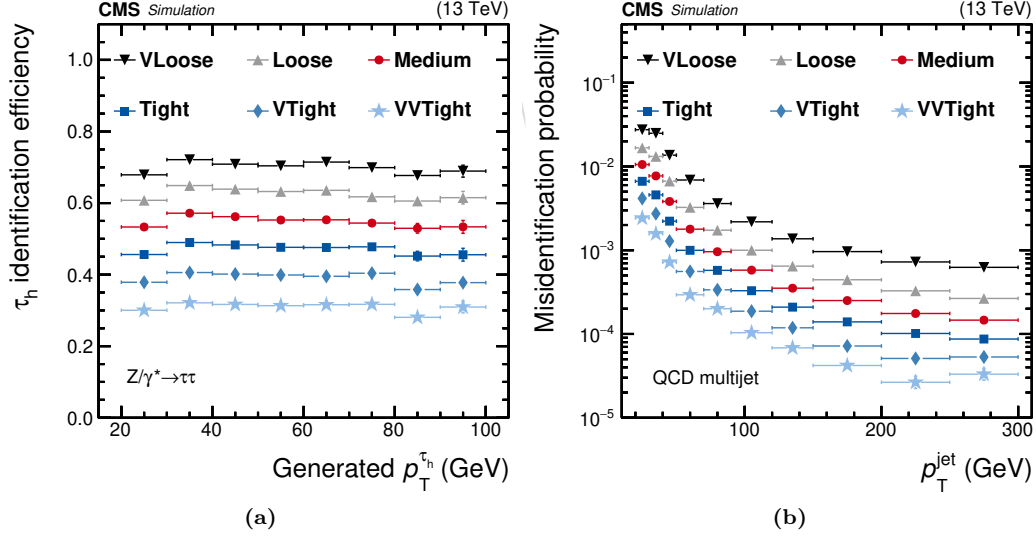


Figure 3.5: The efficiency of genuine τ_h identification (left) and the misidentification probability for $\text{jet} \rightarrow \tau_h$ fakes (right) is shown for each of the six MVA-isolation working points [110].

- Variables relating to the quality of track such as the χ^2 of the track fit.
- Kinematic variables of the τ_h candidate including p_T , η , and mass.

Five working points for the anti-electron discriminator are defined (*loose–vtight*) with efficiencies that range between $\approx 70\text{--}85\%$ for misidentification probabilities between $\mathcal{O}(10^{-4})\text{--}\mathcal{O}(10^{-1})$.

The anti-muon discriminator is a cut-based variable with two working points, *loose* and *tight*. This discriminant is based on vetoing candidates when track segments or hits in the muon detector are found near the τ_h direction, and comparing the calorimeter energy to the track momentum. The efficiency for the discriminator is $> 99\%$ for both working points with a misidentification probability $\mathcal{O}(10^{-3})$.

3.2 Datasets and triggers

The datasets collected during 2016 and 2017 are used for the analyses described in this thesis. These datasets correspond to integrated luminosities of 35.9 fb^{-1} and 41.9 fb^{-1} , respectively. A set of triggers requiring combinations of electrons, muons, and hadronically decaying taus were used to collect interesting events. The chosen triggers are designed to select events with tau pairs decaying to the $\tau_h\tau_h$, $\tau_\mu\tau_h$, $\tau_e\tau_h$, and $\tau_e\tau_\mu$ final states, which will be referred to as "channels".

A double-tau trigger was used to select $\tau_h\tau_h$ events. This trigger requires two isolated hadronically decaying taus with $p_T^{\tau_h} > 35 \text{ GeV}$.

The $\tau_\mu\tau_h$ final state was selected using single-muon and muon+tau triggers. The single-muon trigger requires an isolated muon candidate with a p_T threshold between 22 to 27 GeV, depending on the data taking period. In 2016 the p_T threshold was 22 GeV throughout the year, whereas for 2017 the threshold was either 24 or 27 GeV depending on the run period. Muon+tau triggers were also available for both 2016 and 2017 data taking periods. These triggers require both an isolated muon and an isolated hadronic tau. In 2016 the muon was required to have $p_T^\mu > 19 \text{ GeV}$ and the tau was required to have $p_T^{\tau_h} > 20 \text{ GeV}$. In 2017 the p_T thresholds on the muon and tau candidates were increased to 20 GeV and 27 GeV respectively.

Single-electron and electron+tau triggers target the $\tau_e\tau_h$ final state; although the latter was only used in 2017. The single-electron trigger requires an isolated electron candidate with a p_T threshold between 24 to 32 GeV, depending on the data taking period. In 2016 the p_T threshold was 24 GeV throughout the year, whereas in 2017 the threshold was either 27 or 32 GeV depending on the run period. The electron+tau cross trigger utilised for 2017 data taking requires the presence of an isolated electron with $p_T^e > 24$ GeV, and an isolated hadronic tau with $p_T^{\tau_h} > 30$ GeV. Finally, electron+muon triggers, requiring one isolated muon and one isolated electron, were used to select $\tau_e\tau_\mu$ events. Two such triggers were used, each placing different p_T requirements on the leptons - the first requires $p_T^\mu > 8$ GeV and $p_T^e > 23$ GeV, and the second requires $p_T^\mu > 23$ GeV and $p_T^e > 12$ GeV.

Although all the triggers described above may be used to select ditau events, the final combination of triggers is analysis specific. The exact combination of triggers used for each analysis will be indicated in the relevant chapters.

3.3 Event simulation

The following text gives a brief overview of tools used to simulate proton-proton collision events based on the more detailed review given in Ref. [111]. The generation of a simulated physics event starts with the process of interest known as the hard subprocess. This process involves the scattering of two incoming partons producing a few outgoing particles. The fraction of the proton momenta carried by an incoming parton is stochastic and is therefore sampled from parton distribution functions (PDFs). The PDFs are functions $f_i(x, \mu_F)$ which describe the probability of a parton species i to carry a momentum fraction x probed at a chosen factorisation scale μ_F . The hard subprocess involves either large momentum transfer or heavy particle production, resulting in energetic outgoing partons that are asymptotically free. These processes are therefore generated according to matrix-elements using perturbation theory. The matrix-element computations involve the cancellation of ultraviolet and infrared divergences, which introduces a dependence on μ_F and the renormalisation scale μ_R . Current generators are usually either leading order (LO) or next-to-leading order (NLO) accurate, but progress on next-to-next-to-leading order (NNLO) generation has been made in recent years such as in Ref. [112].

Parton-shower generators evolve the events from the hard scale to the low scale $\mathcal{O}(1 \text{ GeV})$. This involves the emission of many additional soft or collinear partons. At the lowest scale hadronisation occurs, where outgoing partons interaction non-perturbatively to form hadrons. The non-perturbative nature of this problem calls for the use of models with tunable parameters to describe the data. These tools also provide a description of secondary interactions of proton constituents not involved in the hard subprocess, known as the UE. The input parameters of these models are tuned using collision data. Unstable particles such as taus and heavy-flavour mesons must then also be decayed, producing a collection of stable final-state particles. The response of the detector to these final-state particles is finally simulated, to obtain a simulated event that is analogous to an event recorded during data-taking.

3.3.1 Monte Carlo samples

The Monte Carlo (MC) samples used to model background processes are common to both analyses described in this thesis and as such will be described in full here. The analysis-specific signal samples will be described in more detail in the following chapters. Separate MC samples were produced to model 2016 and 2017 conditions. As both 2016 and 2017 datasets were collected at the

same centre of mass energy, there should be no change to the true event topologies. The samples should therefore only vary due to different detector conditions including, for example, different detector responses and/or different numbers of PU interactions. However, updates to the PDFs and UE tunes were made in 2017 and subsequently included in the 2017 MC production. The generator-level event topologies can therefore also differ slightly between eras.

The QCD-induced W +jets and Drell-Yan Z/γ^* +jets processes are modelled at LO accuracy using **MadGraph5** [113] with MLM jet matching [114], including 0–4 outgoing partons in the matrix element calculations. In addition to the inclusive samples, exclusive samples were generated each with exactly 1, 2, 3, or 4 outgoing partons to increase the event statistics. The exclusive and inclusive samples are "stitched" together by applying event weights, such that the weighted number of events for each parton multiplicity matches that of the unweighted inclusive sample. For the $\mathcal{CP}$ analysis, the electroweak production of W and Z bosons is also considered, which is modelled using **MadGraph5** at LO accuracy. The single-top and $t\bar{t}$ processes are simulated at NLO accuracy using **POWHEG** [115–118]. Diboson processes including ZZ , WZ , and WW are modelled at NLO accuracy using **Madgraph5_aMC@NLO** [119] with FxFx jet merging scheme [120]. The $W + \gamma$ and $W + \gamma^*$ processes are also considered in some instances, which are modelled using **Madgraph5_aMC@NLO** with FxFx jet merging at NLO and **MadGraph5** with MLM jet matching at LO respectively.

The NNPDF3.0 [121] and NNPDF3.1 [122] PDFs were used in the generation of the 2016 and 2017 samples respectively. The parton showering, hadronisation, and tau decays are modelled using **PYTHIA** 8.2 [123]. The UE simulation is also performed using **PYTHIA** with the parametrisation set to the CUETP8M1 [124] tune for the 2016 samples, and to the CP5 [125] tune for the 2017 production. The generated events are then passed through a **GEANT4**-based [126] simulation of the CMS detector, and are reconstructed using the same algorithms used to reconstruct the data. Additional inelastic proton-proton interactions (PU) are generated using **PYTHIA** and mixed with the simulation to match the expected PU profile of the data.

3.3.2 Embedded samples

The description of the $Z/\gamma^* \rightarrow \tau\tau$ process can be improved using an embedding technique [127]. This method involves selecting $Z/\gamma^* \rightarrow \mu\mu$ events in data and replacing the muons with simulated tau leptons, thus exploiting lepton universality in $Z/\gamma^* \rightarrow ll$ decays. This literal embedding of the taus into the event gives the procedure its name. As illustrated in Figure 3.6, the production of the embedded samples can be factorised into four stages:

1. Selection: Dimuon events are selected in data. The selection requirements are deliberately kept as loose as possible to minimise any biases. Both muons are required to be global muons and to have been identified by the PF algorithm. The muons must have also passed a dimuon trigger with online p_T thresholds of 17 and 8 GeV for the leading and sub-leading muon legs respectively.
2. Cleaning: The hits and energy deposits associated with the selected muons are removed from the event.
3. Simulation: A $Z/\gamma^* \rightarrow \tau\tau$ event is simulated with the kinematics of the taus set to that of the removed muons. The taus are then decayed using **PYTHIA** and the final state particles are passed through the detector simulation as described in the previous section.
4. Merging: All hits and energy deposits from the simulated tau pair are combined with the cleaned event.

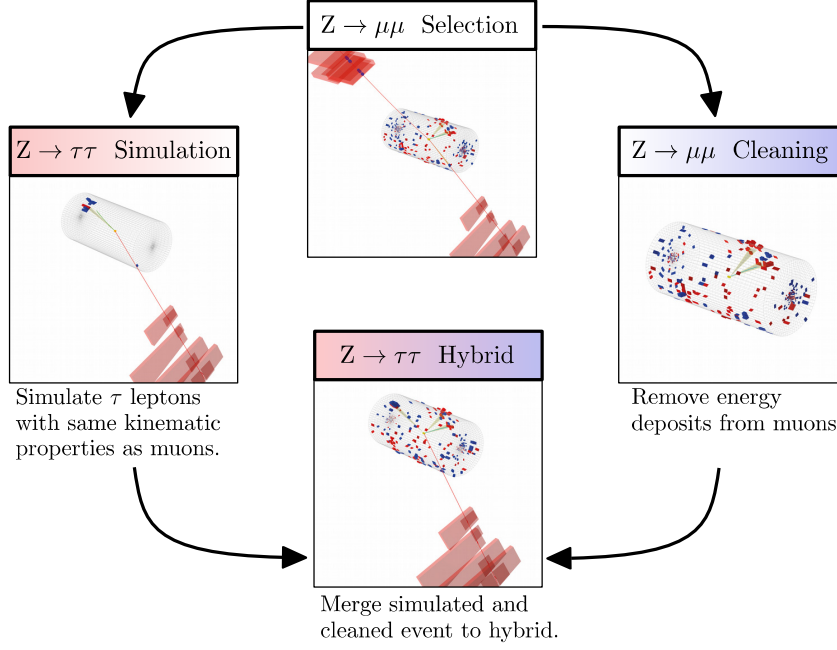


Figure 3.6: Illustration of the four main stages in the production of an embedded event: selection, cleaning, simulation, and merging [127].

Approximately 99% of the selected dimuons come from Z/γ^* decays, with the remaining 1% predominantly from $t\bar{t}$ and diboson events. The muons produced in these processes typically originate from decays of W and Z bosons, which also respect lepton universality. Therefore, the embedding procedure not only models the $Z/\gamma^* \rightarrow \tau\tau$ process, but provides an estimate for all backgrounds producing a pair of genuine tau leptons.

The main advantage of using the embedded samples is that the bulk of the event content comes directly from data, with the exception of the well described tau leptons. This includes the often difficult to model jet topology, PU, and UE structure. By construction these properties are extremely well modelled by the embedded samples. The Z boson kinematics can also be poorly described by MC generators (as will be discussed in Section 3.4.5). This problem is again remedied by the embedded technique, as the $Z/\gamma^* \rightarrow \tau\tau$ and $Z/\gamma^* \rightarrow \mu\mu$ kinematics are near identical.

The statistics of the embedded samples increase linearly with the amount of collected collision data, which removes the need for the computationally demanding simulation of large MC samples. The sample statistics are further increased by requiring the visible decay products of the taus to pass loose p_T requirements. The neutrinos from the tau decay carry off a share of the energy, and often this results in visible decay products that are too soft to pass the online/offline p_T cuts. To avoid unnecessary discarding of such events, every tau pair decay is repeated 1000 times and only if the taus fail to meet the p_T requirements on every repetition is the event discarded. The last tau pair to pass the selections is retained for the subsequent detector simulation. An event weight defined as the number of times the tau pair successfully passes the p_T requirements divided by 1000 and multiplied by the branching ratio of the corresponding $\tau\tau$ decay is also stored. Application of these weights removes any resulting bias to the variable distributions, and allows the yields to be estimated correctly as will be described in Section 3.5.2.

3.4 Simulation corrections

The objects derived from simulation are prone to mismodelling which must be corrected for. For the MC this of course includes all the objects in the event. The embedded samples only model the tau decay products (electrons, muons, and hadronic taus) using simulation, and therefore corrections are derived only for these objects. The applied corrections can be grouped into three categories:

- MC generator corrections: Due to the mismodelling of the true event topology by MC generators. An example is the poor description of the Z boson p_T distribution by LO generators.
- Efficiency corrections: Due to differences in the rates at which objects pass selection requirements such as the identification, isolation, and triggers. These corrections are applied using scale factors defined as the ratio of the efficiencies measured in data and simulation, $SF = \frac{\epsilon_{\text{data}}}{\epsilon_{\text{sim}}}$.
- Energy scale corrections: Due to differences in the detector response to objects in data and simulation. These corrections are applied by shifting an objects energy to match the data.

In CMS it is common practice to distribute corrections centrally where possible, so that they may benefit several analyses. However, centrally distributed scale factors do not always match the specific needs of an analysis (e.g if a different combination of triggers is required), and in these instances it is necessary for analysts to measure corrections themselves. This section describes the corrections used in the analyses described in this thesis. More details will be provided in cases where the corrections were derived by the author, which will also serve as examples of how the centrally distributed scale factors are measured.

In general separate corrections are required for both MC and embedded objects, as often the size of the corrections differ significantly. Embedded corrections are derived either by comparing the embedded samples directly with data, or by comparing the embedded samples with MC. For the latter, the corresponding MC→data correction is applied in addition to the embedding→MC correction. This procedure is slightly more convoluted, however, the embedding→MC corrections are typically much easier to measure than embedding/MC→data corrections, and in this way centrally distributed MC→data corrections may be utilised.

3.4.1 E_T^{miss} recoil corrections

The E_T^{miss} recoil corrections described in Section 3.1.7 are applied to the W+jets, Z/ γ^* +jets, and signal events. Figure 3.7 shows the improvement in the data/MC agreement by the application of the recoil corrections for $Z \rightarrow \mu\mu$ events.

3.4.2 Pile-up reweighting

The PU profile for the MC samples does not exactly match the data. The MC events are reweighted so that the number of PU interactions match the data. The number of interactions is estimated from the measured instantaneous luminosity per bunch crossing.

3.4.3 B-tag efficiency corrections

Scale factors are used to correct for differences in the rates at which jets pass the CSVv2 algorithm in data and simulation. The scale factors are measured as a function of the jet p_T and η and separately for different jet-types (b-, c-, or light-flavour jets), as described in Ref. [89]. The

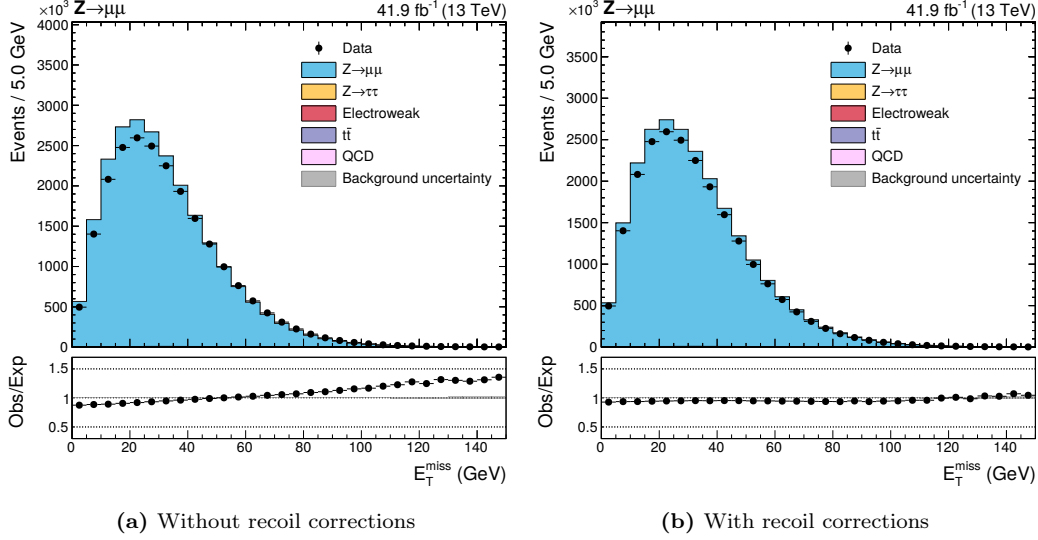


Figure 3.7: Comparisons of the data/MC agreement of the E_T^{miss} distributions for $Z/\gamma \rightarrow \mu\mu$ events before (left) and after (right) applying recoil corrections for data collected in 2017. The uncertainty band shown in shaded grey includes only statistical uncertainties.

measured scale factors are applied using a "promote-demote" method, whereby a simulated jet not passing the b-tagger can be "promoted" to a tagged jet, and similarly a jet passing the b-tagger can be "demoted" to an untagged jet depending on the following:

- For a given jet the promote/demote probabilities are computed as

$$P(\text{promote}) = \frac{\epsilon_{\text{sim}}(\text{SF} - 1)}{1 - \epsilon_{\text{sim}}} \quad \text{and} \quad (3.14)$$

$$P(\text{demote}) = 1 - \text{SF},$$

where $\epsilon_{\text{sim}} = \epsilon_{\text{sim}}(p_T, \eta, \text{jet-type})$ is the b-tagging efficiency measured in simulation.

- A random number between 0–1 is generated.
- If the jet is untagged, $\text{SF} > 1$, and the random number is greater than $P(\text{promote})$ then the jet is promoted.
- If the jet is b-tagged, $\text{SF} < 1$, and the random number is greater than $P(\text{demote})$ then the jet is demoted.

3.4.4 Top p_T reweighting

The agreement of the $t\bar{t}$ MC with the data is improved by reweighting the top p_T distribution. The weights are given as

$$\text{weight} = \sqrt{\text{SF}(p_T^t) \cdot \text{SF}(p_T^{\bar{t}})} \quad (3.15)$$

$$\text{SF}(p_T) = e^{a - b/(\text{GeV}) \cdot p_T},$$

where $p_T^{t/\bar{t}}$ are the p_T of the top quarks, and a and b are constants. The weights applied for the $\phi \rightarrow \tau\tau$ analysis were derived based on data/MC comparisons made in Ref. [128], using data collected at 8 TeV. This corresponds to setting $a = 0.156$ and $b = 0.00137 \text{ GeV}^{-1}$ in Equation 3.15. The weights were updated prior to the $\mathcal{CP}$ analysis based on the comparisons in Ref. [129, 130] using 13 TeV data, corresponding to setting $a = 0.0615$ and $b = 0.0005 \text{ GeV}^{-1}$ in Equation 3.15.

Figure 3.8 shows the improvement in the data/MC agreement of the $\tau_e\tau_\mu$ invariant mass distribution in a $t\bar{t}$ control region.

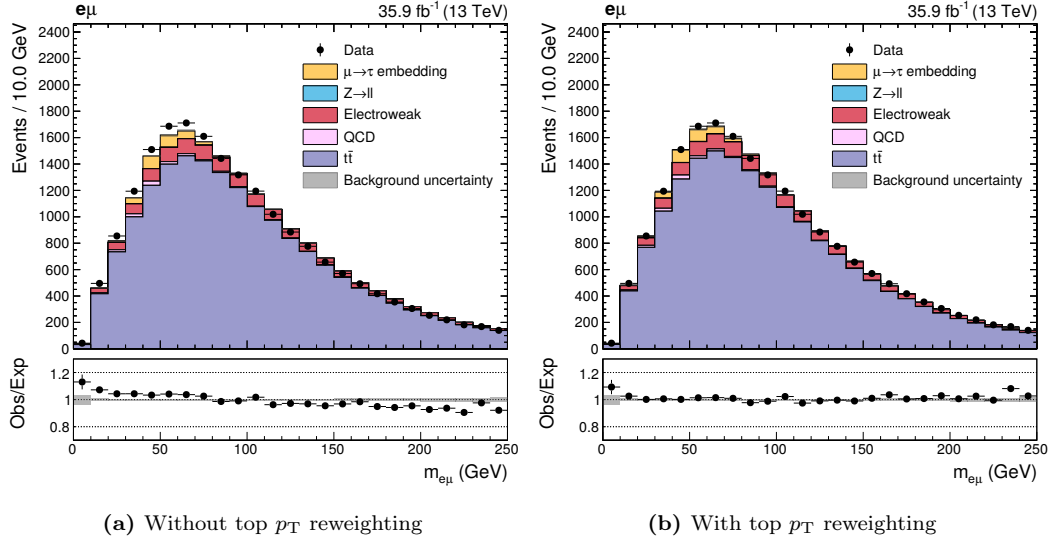


Figure 3.8: Comparison of the data/MC agreement of the $\tau_e\tau_\mu$ invariant mass distribution in a $t\bar{t}$ control region before (left) and after (right) applying the top p_T reweighting. The uncertainty band shown in shaded grey includes only the statistical uncertainties.

3.4.5 Z/γ^* p_T -mass reweighting

The LO generator used to simulate the Drell-Yan Z/γ^* process does not describe well events where the Z/γ^* boson has high p_T or large mass. These regions of phase space are important for $\tau\tau$ analyses, as the tight p_T cuts on the tau decay products imposed by the trigger preferentially select more boosted Z/γ^* bosons. The high mass regions are of course also relevant for high mass $\phi \rightarrow \tau\tau$ searches. The modelling of these variables can be significantly improved using NLO generators, however, the simulation of NLO samples with large event statistics is impractical.

The data/MC agreement can alternatively be improved by deriving weights from a sample of $Z/\gamma^* \rightarrow \mu\mu$ events and applying them to the LO samples. Muon pairs with opposite-sign charges separated by $\Delta R > 0.5$ are selected. Each muon is required to have $p_T > 10$ GeV, $I_{\text{rel}} < 0.15$, and pass the *medium* ID. The leading muon must also pass the single-muon trigger with an online p_T threshold between 22–24 GeV depending on the run period, and pass an additional offline p_T cut 1 GeV higher than the online requirement.

The weights are measured in bins of the dimuon p_T ($p_T^{\mu\mu}$) and invariant mass ($m_{\mu\mu}$). For each bin the weight is computed as,

$$\text{weight}(p_T^{\mu\mu}, m_{\mu\mu}) = \frac{N_{\text{data}} - N_{\text{MC}}^{\text{non-}Z/\gamma^*}}{N_{\text{MC}}^{Z/\gamma^*}}, \quad (3.16)$$

where N_{data} is the yield measured in data, $N_{\text{MC}}^{\text{non-}Z/\gamma^*}$ is the contribution of processes other than $Z/\gamma^* \rightarrow \mu\mu$ estimated from MC, and $N_{\text{MC}}^{Z/\gamma^*}$ is the $Z/\gamma^* \rightarrow \mu\mu$ yield as predicted by the LO generator. The weights are derived imposing

$$\sum_{\text{bins}} (N_{\text{data}} - N_{\text{MC}}^{\text{non-}Z/\gamma^*}) = \sum_{\text{bins}} N_{\text{MC}}^{Z/\gamma^*} = 1, \quad (3.17)$$

where the sums run over all p_T -mass bins, such that the overall normalisation of the selected $Z \rightarrow \mu\mu$ events is not changed by the application of the weights. The weights are applied depending on the generator-level boson p_T and mass. Therefore, the procedure relies upon the $p_T^{\mu\mu}$ and $m_{\mu\mu}$ being similar to the equivalent generator-level quantities, which is a good approximation due to the excellent muon resolution of the CMS detector.

The reweighting procedure is tested using samples of $Z/\gamma^* \rightarrow \mu\mu$ and $Z/\gamma^* \rightarrow ee$ events. As the former consists of the same events used to derive the weights good agreement with the data is somewhat expected, however, this allows the assumption that the generator- and reconstruction-level variables are similar to be directly tested. The $Z/\gamma^* \rightarrow ee$ sample provides an additional check of the reweighting procedure for an orthogonal event selection. A comparison of the $p_T^{\mu\mu}$ and p_T^{ee} distributions before and after the reweighting is shown in Figure 3.9. The reweighted events are observed to describe the shapes of the distributions well. The data/MC agreement for all variables important for event selection and categorisation (e.g. N_{jets}) is also improved after application of the weights, which verifies that the chosen parametrisation in bins of p_T and mass is sufficient.

3.4.6 Electron and muon efficiency corrections

Corrections are applied to electrons and muons to account for data/simulation differences in the tracking, identification, isolation, and trigger efficiencies. The corrections are derived with a tag-and-probe method using samples of $Z \rightarrow \mu\mu$ and $Z \rightarrow ee$ events. The tracking corrections derived for MC leptons are also applied to the embedded events. The electron tracking efficiency is found to be 2% (1%) higher for the 2016 (2017) embedded samples compared to the MC. This difference is applied as an additional correction to the embedded samples. The muon tracking efficiencies are the same in both MC and embedding and therefore no additional correction is applied. The identification, isolation, and trigger corrections used for the 2017 samples were measured by the author, and as such will be described below.

Dilepton events are selected with one lepton designated as a tag and the other as a probe. The tags are required to pass identification and isolation requirements to reduce the background from fake leptons. The tags are also required to pass a single-lepton trigger, thus ensuring that the probe lepton is completely unbiased by the online selections. The specific selections applied for the electron and muon tags are:

- Electrons: The tag is required to pass the *tight* electron ID and have $I_{\text{rel}} < 0.1$. The tag must also have $p_T > 36$ GeV and pass the single electron trigger with an online p_T threshold of 35 GeV.
- Muons: The tag is required to pass the *medium* muon ID and have $I_{\text{rel}} < 0.15$. The tag must also have $p_T > 25$ GeV and pass the single muon trigger with an online p_T threshold of 24 GeV.

The lepton pairs are additionally required to carry opposite charges and have invariant mass $65 < m_{ll} < 115$ GeV, to further increase the purity of $Z \rightarrow ll$ events. Each lepton is also treated in turn as both the tag and the probe to approximately double the available statistics. The ID, isolation, and trigger scale factors are measured separately but in sequence meaning that the isolation scale factors are measured under the pre-condition that the probe leptons have passed the ID, and similarly the trigger scale-factors are measured under the pre-condition that the probe leptons have passed both the ID and isolation requirements.

Events are split into bins depending on the probe p_T and η . These bins are further divided into

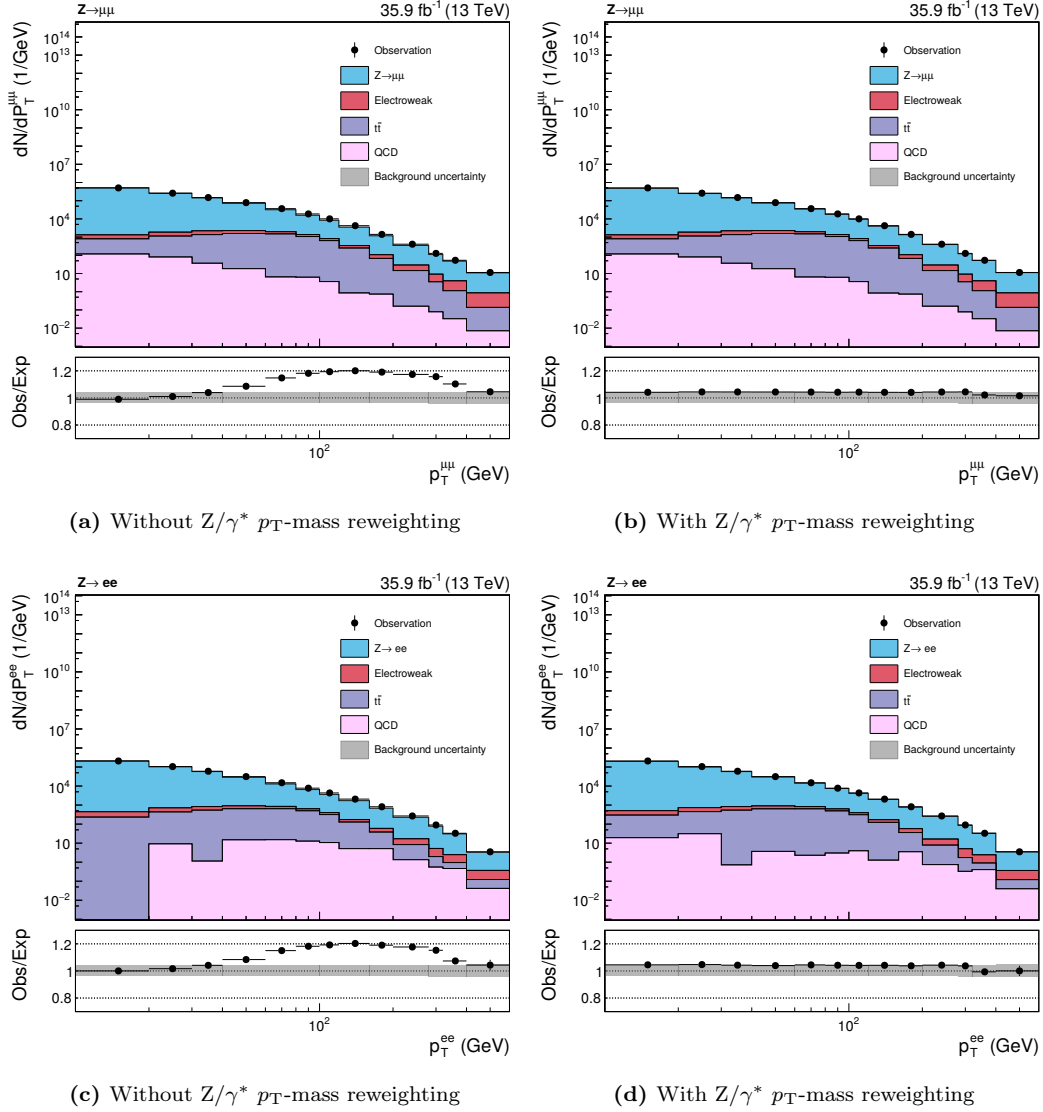


Figure 3.9: Comparison of the data/MC agreement of the Z/γ^* p_T distributions before (left) and after (right) applying the Z/γ^* p_T -mass reweighting. The upper and lower insets show the distributions in the $Z/\gamma^* \rightarrow \mu\mu$ and $Z/\gamma^* \rightarrow ee$ final state. The uncertainty band shown in shaded grey includes the uncertainties on the cross section and luminosity added in quadrature with the statistical uncertainties.

"pass" and "fail" categories, depending on whether the probe lepton passes the ID/isolation/trigger requirement in question. Even after applying the aforementioned selections, some of the events selected originate from background processes. This background contamination is accounted for by fitting the m_{ll} distributions. The signal is modelled using a superposition of two Voigtian functions, with the width of the Breit-Wigner part fixed to the Z natural-width. The background is modelled using a decaying exponential for the isolation and trigger measurements, and using an error function turning into a decaying exponential for the ID measurements. The latter is necessary as the background contamination is significantly larger for the ID measurements, and this model provides a better description for the kinematic turn-on regions of the distributions. In principle this same background model could also be used for the isolation and trigger methods; however, there tends to be too few background events to constrain the model parameters in this case, and often this leads to part of the signal being mistakenly fitted as background. The efficiency is given

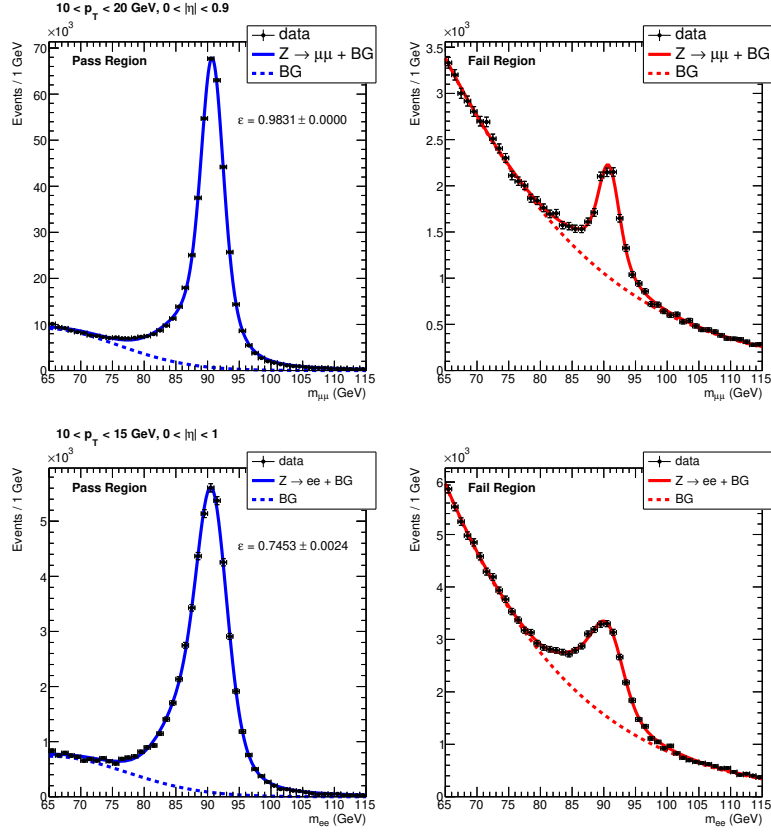


Figure 3.10: Examples of the fits used to measure the muon (upper) and electron (lower) efficiencies for determining scale-factors. In this example the lepton ID efficiencies are being measured.

as $\epsilon = \frac{N_{\text{pass}}}{N_{\text{pass}} + N_{\text{fail}}}$, where $N_{\text{pass/fail}}$ are the best-fit values for the $Z \rightarrow ll$ yields. Figure 3.10 shows an example of the fitted m_{ll} distributions used to measure the electron and muon ID efficiencies.

The efficiencies and scale factors measured for muons in the region $|\eta| < 0.9$ and electrons in the region $|\eta| < 1.0$ are shown in Figure 3.11. The trigger efficiencies are measured for the single-muon trigger with an online p_T threshold that was increased from 24 to 27 GeV during data taking, and for the single-electron trigger with a p_T threshold increased from 27 to 32 GeV. The trigger and electron ID corrections are observed to differ significantly for the embedding and MC samples. The former is due to the trigger being simulated in an empty detector environment, before the merging stage of the embedding procedure. This is necessary as the trigger sequence depends on the precise alignment of the simulated and physical detector geometries, which is difficult to achieve. The empty environment effectively means the electrons and muons are very well isolated and easier to identify, which leads to larger efficiencies. Conversely, the isolation variables are computed after the merging stage and therefore the MC and embedded corrections are much more similar. The electron ID is calculated also after the merging stage, however, this quantity is affected by the slight misalignment of the detector geometries, resulting in a lower efficiency for the embedded events.

3.4.7 Hadronic tau efficiency corrections

Corrections are applied to taus to account for data/simulation differences in the trigger and isolation efficiencies and, in the case of the embedded samples, the tracking efficiency. The latter correction affects the rate at which taus pass the HPS algorithm. This is found to be consistent with a scale

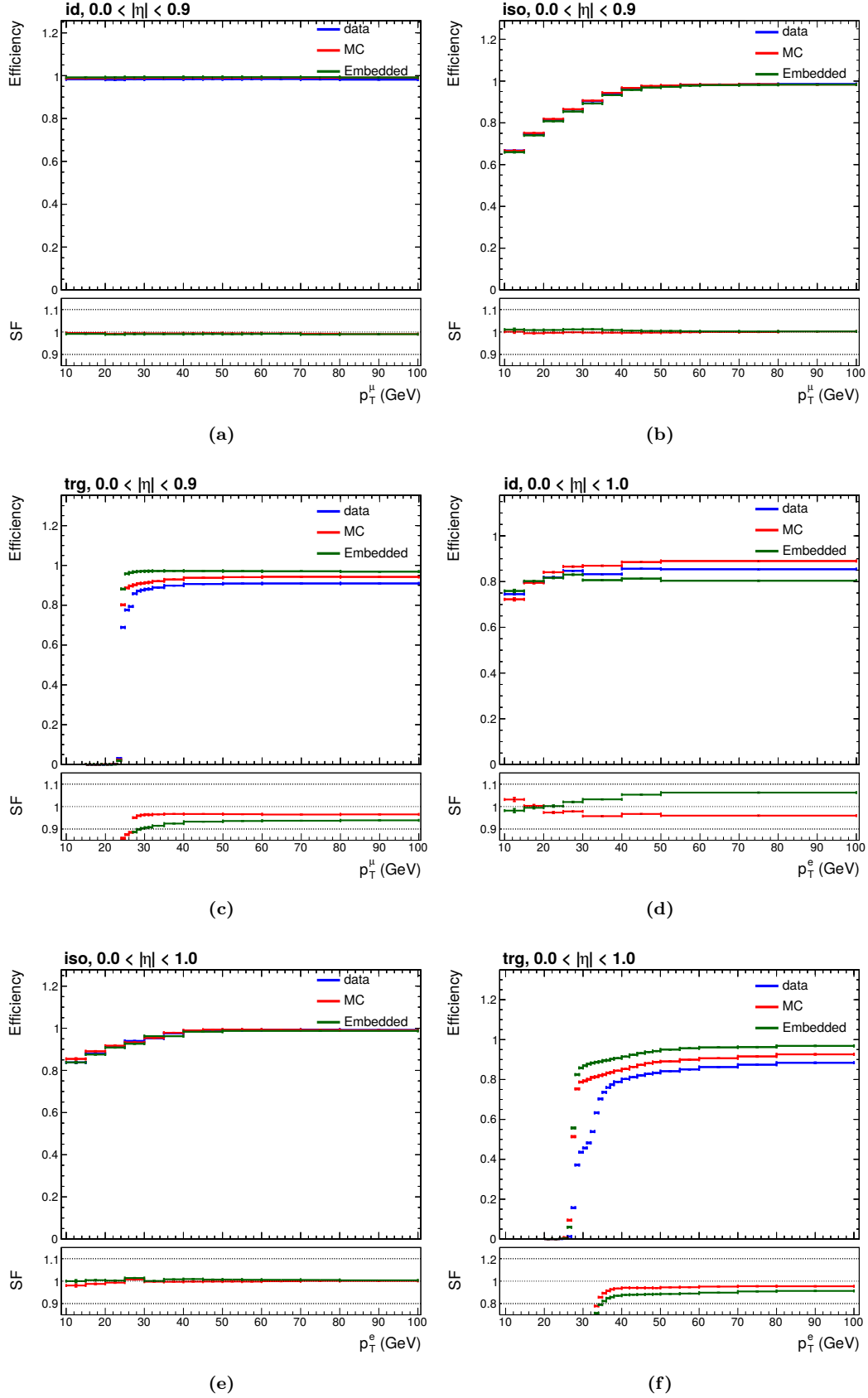


Figure 3.11: The measured muon ID (a), isolation (b), and single-muon trigger (c) efficiencies for the region $|\eta| < 0.9$, and the measured electron ID (d), isolation (e), and single-electron trigger (f) efficiencies for the region $|\eta| < 1.0$. The efficiencies measured in data, MC, and embedded events are shown in blue, red, and green respectively.

factor of 1 for the MC, and so no correction is applied. The tracking efficiency is found to differ for embedded and MC events. The ratio of the MC and embedded efficiencies, which depend on the decay mode of the tau, is applied as a correction. The size of this correction is 0.975 for 1-prong taus, 1.023 for 1-prong+ $1\pi^0$ taus, and 0.927 for 3-prong taus.

The trigger efficiencies are measured with a tag-and-probe method using samples of $Z/\gamma^* \rightarrow \tau_h \tau_\mu$ events, where the τ_μ is used as the tag. Dedicated muon+tau monitoring triggers are used for these measurements, where the online selections applied to the tau leg match exactly those used for the trigger in question. The corrections are measured for the tau legs of lepton+tau triggers and the individual tau legs of the double-tau trigger.

The isolation correction accounts for differences in the rate at which taus pass the MVA-isolation discriminant. A constant scale factor is applied to account for this difference. The scale factors applied to the MC have a size of 0.95 (0.97) for the *tight* (*medium*) ID working point in the 2016 samples, and 0.89 for the *tight* working point in the 2017 samples. The scale factors applied for the embedded samples were measured by the author as described below.

A sample enriched in $Z/\gamma^* \rightarrow \tau_h \tau_\mu$ events is selected. For the 2016 (2017) measurement the following selections are imposed:

- The muon is required to pass the *medium* ID, have $I_{\text{rel}} < 0.15$, have $p_T > 23$ (28) GeV, and pass the single-muon trigger.
- The hadronic tau candidate must have $p_T > 20$ GeV and pass the HPS algorithm, the *tight* working-point of the anti-muon discriminator, and the *very-loose* working-point of the anti-electron discriminator.
- The muon and tau should have opposite charge and be separated by $\Delta R > 0.5$. If more than one pair meet this requirement the most isolated pair is retained.
- The transverse mass $m_T(\mu, E_T^{\text{miss}})$, which is defined as

$$m_T(i, j) = \sqrt{2p_T^i p_T^j (1 - \cos \Delta\phi_{ij})}, \quad (3.18)$$

must be smaller than 40 GeV.

- Events with additional electrons or muons are vetoed.

The events are then split into a "pass" and a "fail" category, depending on whether the tau passes the MVA-isolation discriminant.

The embedded events are treated as the signal process, which will be denoted here as "EmbedZTT". The method used to estimate the yield of this process will be described in Section 3.5.2. The background processes that are considered are QCD, W+jets, $t\bar{t}$, diboson, single-top, and $Z/\gamma^* \rightarrow \mu\mu$. The methods used to estimate the W+jets and QCD backgrounds will be described in Section 3.5.4. All other backgrounds are modelled using MC. The data/simulation corrections described in this section are applied to the MC and embedded events.

A simultaneous maximum-likelihood fit of the pass and fail categories is performed using the tau ID scale factor as the parameter of interest. The invariant mass of the muon and hadronic tau is used as the discriminating variable. The EmbedZTT template is split into two sub-templates depending on whether the τ_h passes or fails the MVA-isolation which are denoted EmbedZTT^{pass} EmbedZTT^{fail} respectively. The signal template in the pass category is constructed as a function of the SF as

$$\text{EmbedZTT}^{\text{total}}(\text{pass}) = \text{SF} \cdot \text{EmbedZTT}^{\text{pass}}, \quad (3.19)$$

and in the fail category the template is constructed as

$$\text{EmbedZTT}^{\text{total}}(\text{fail}) = (1 - \text{SF}) \cdot \text{EmbedZTT}^{\text{pass}} + \text{EmbedZTT}^{\text{fail}}. \quad (3.20)$$

The following systematic uncertainties affecting the normalisations and shapes of the signal and background templates are also included:

Normalisation uncertainties:

- **Luminosity:** The luminosity uncertainty amounts to 2.5% on the data collected during 2016 data-taking and 2.3% on the data collected in 2017. This uncertainty is applied to all processes where the yield is estimated from MC.
- **Muon ID/isolation efficiency:** The uncertainty on the muon ID/isolation efficiency is 2%.
- **Tau tracking/ID efficiency:** The dominant processes producing genuine taus are already included in the embedded templates. However, some smaller processes such as $t\bar{t} \rightarrow \tau\mu$ contribute to the background. A 10% uncertainty on the tau tracking/ID efficiency is applied.
- $j \rightarrow \tau_h$ **fake-rate:** A 20% uncertainty is applied for all processes containing a $j \rightarrow \tau_h$ fake where the yield is estimated from MC.
- $\mu \rightarrow \tau_h$ **fake-rate:** A 20% uncertainty is applied to the $Z \rightarrow \mu\mu$ process.
- **Trigger efficiency:** The uncertainty on the single-muon trigger efficiency amounts to 2%.
- **Background normalisation uncertainties:**
 - A 4% uncertainty on the EmbedZTT yield.
 - A 4% uncertainty on the $Z \rightarrow \mu\mu$ yield to account for the uncertainty on the Drell-Yan cross section.
 - A 5% uncertainty is on the diboson and single-top yields to account for the uncertainties on the production cross sections.
 - A 6% uncertainty is on the $t\bar{t}$ yields to account for the uncertainties on the production cross section.
 - A 10% uncertainty on the W+jets yield which is uncorrelated for the pass and fail categories.
 - A 10% uncertainty on the QCD yield which is uncorrelated for the pass and fail categories.

Shape uncertainties:

- **Hadronic tau energy scale:** Three uncorrelated uncertainties each that affect only one of the three tau decay modes. Each uncertainty varies the hadronic tau energy scale by 1.2 (3)% for the 2016 (2017) MC and embedded samples.
- $\mu \rightarrow \tau_h$ **energy scale:** Two uncorrelated uncertainties each that affect only one of the 1-prong tau decay modes are applied to the $Z \rightarrow \mu\mu$ process. These uncertainties vary the energy scale correction by 1.5% for both the 1-prong and 1-prong+1- π^0 decay modes in the 2016 MC. In 2017 the size of the uncertainties are 0.3 and 1.0% for 1-prong and 1-prong+1- π^0 decay mode respectively.
- **Statistical uncertainties:** Limitations in the statistical sizes of the signal and background templates are accounted for by the Barlow-Beeston method [131, 132].
- **Top p_T reweighting:** Uncertainty between no and twice the correction on the reweighting applied to $t\bar{t}$ events.

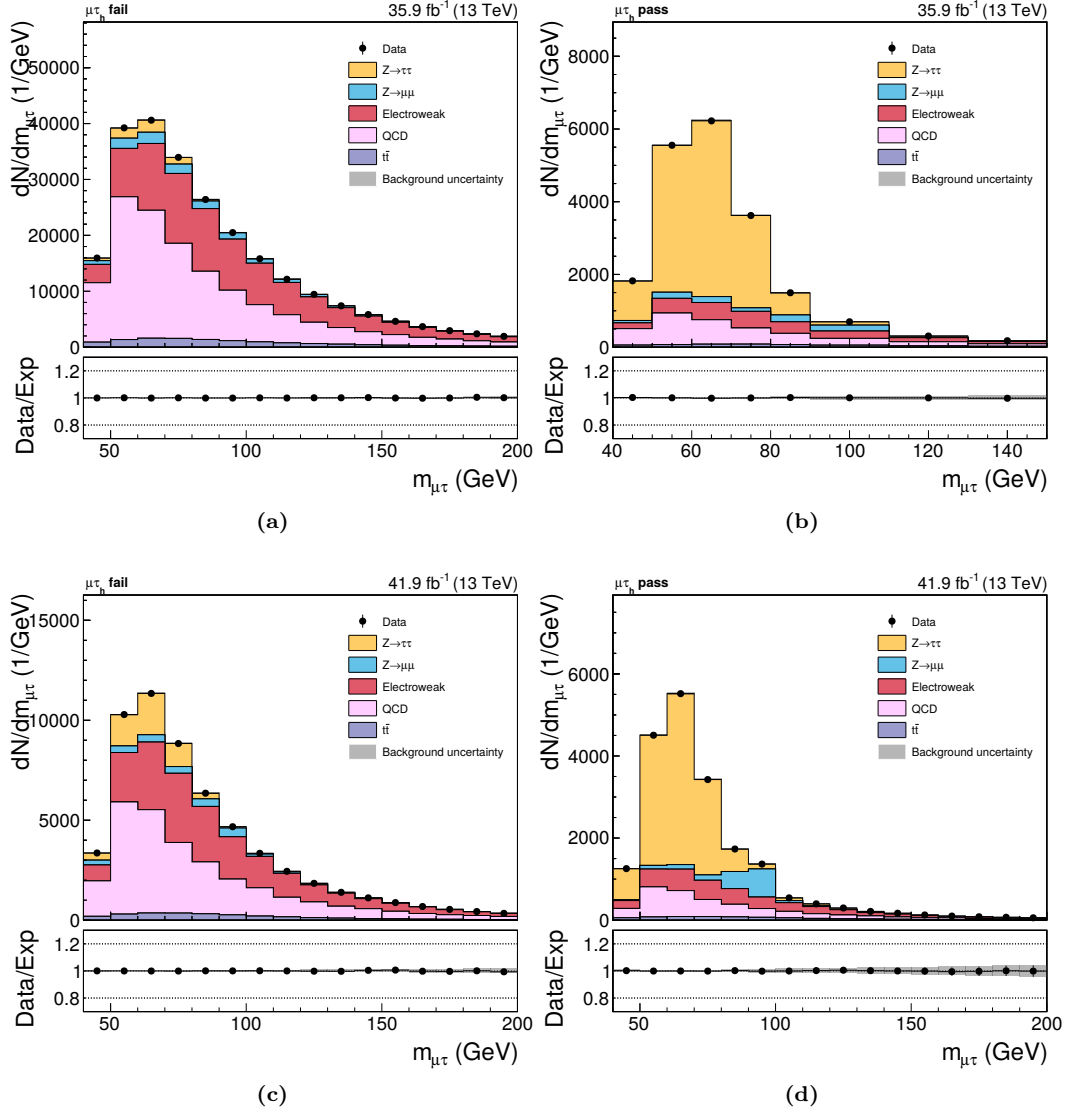


Figure 3.12: The post-fit distributions showing the data/background-prediction agreement for the fail (left) and pass (right) categories for the measurement of the 2016 (upper) and 2017 (lower) embedded tau ID scale-factor. The uncertainty band shown in shaded grey includes both systematic and statistical uncertainties.

- **Drell-Yan Z p_T reweighting:** Uncertainty between no and twice the correction on the reweighting applied to $Z \rightarrow \mu\mu$ events.

The post-fit data/background comparisons of the pass and fail categories are shown in Figure 3.12. Good agreement with the data is observed. The results of the maximum likelihood scans of the SF are shown in Figure 3.13. The measured scale factor is 1.02 ± 0.03 for the 2016 samples, and 0.97 ± 0.02 for the 2017 samples.

3.4.8 $e \rightarrow \tau_h$ and $\mu \rightarrow \tau_h$ fake rate corrections

The rate at which electrons and muons pass the anti-electron and anti-muon discriminants are measured with a tag-and-probe method using samples of $Z/\gamma^* \rightarrow ee$ and $Z/\gamma^* \rightarrow \mu\mu$ events. Scale factors are applied to correct for the observed differences in the fake rates in data and simulation.

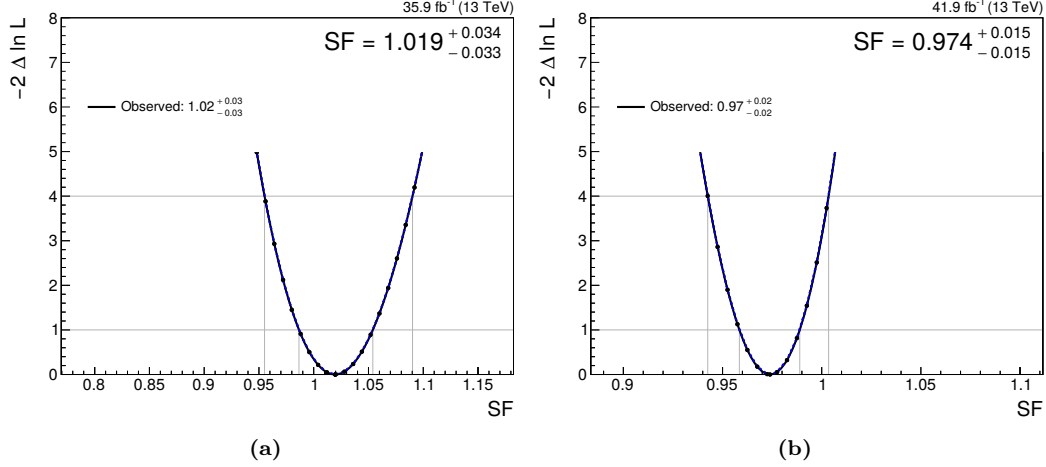


Figure 3.13: The likelihood scans showing the best-fit value of the embedded tau ID SF for 2016 (left) and 2017 (right).

3.4.9 Tau energy scale corrections

The τ_h energy scale is measured separately for each tau decay mode using $Z/\gamma^* \rightarrow \tau_h \tau_\mu$ events. The measurement is performed by fitting mass distributions sensitive to differences in the energy scale, using a similar procedure as described in Section 3.4.10. The $\tau_h \tau_\mu$ mass distribution is used to extract the energy scale for 1-prong taus. Whereas, for 1-prong+ $1-\pi^0$ and 3-prong taus the τ_h mass distribution is utilised. Differences in the energy scales in data and MC are corrected by applying shifts to the τ_h energies. The size of the shifts measured for 2016 (2017) are -0.5% (0.7%), 1.1% (-0.2%), and 0.6% (0.1%) for the 1-prong, 1-prong+ $1-\pi^0$, and 3-prong taus respectively.

The corrections are also applied to the embedded samples along with additional shifts to account for differences in the energy scale in MC and embedded events. These shifts are determined by comparing the tau p_T resolution for MC and embedded events, as illustrated in Figure 3.14. The resolution is defined as the difference between the reconstruction- and generator-level p_T divided by the generator-level p_T . The size of the shifts is then given by

$$\text{shift} = \frac{\langle \text{MC} \rangle + 1}{\langle \text{Embed} \rangle + 1}, \quad (3.21)$$

where $\langle \text{MC} \rangle$ and $\langle \text{Embed} \rangle$ are the mean values of the distributions in Figure 3.14. The 3-prong tau energy scale is observed to be practically identical for MC and embedded events and therefore no correction is applied. A small difference is observed for the 1-prong taus amounting to a correction of -0.4% (-0.2%) for the 2016 (2017) samples. The energy scale of the 1-prong+ $1-\pi^0$ taus shows the largest difference between MC and embedding, where the size of the correction is -0.9% .

3.4.10 $e \rightarrow \tau_h$ and $\mu \rightarrow \tau_h$ energy scale corrections

The energy scale corrections described in the previous section are measured for reconstruction-level taus originating from true hadronic taus and so may not apply to taus originating from lepton fakes. Separate corrections are therefore derived for $e \rightarrow \tau_h$ and $\mu \rightarrow \tau_h$ candidates. The $\tau_l \tau_h$ invariant mass distribution can be fitted to extract the correction, as a shift to the $l \rightarrow \tau_h$ energy scale moves the position of the $Z \rightarrow ll$ mass peak. Separate corrections are measured for 1-prong and 1-prong+ $1-\pi^0$ taus. The 3-prong decay mode is not considered as the misidentification probability is much smaller than for the other decay modes. The size of the $e \rightarrow \tau_h$ corrections applied to the

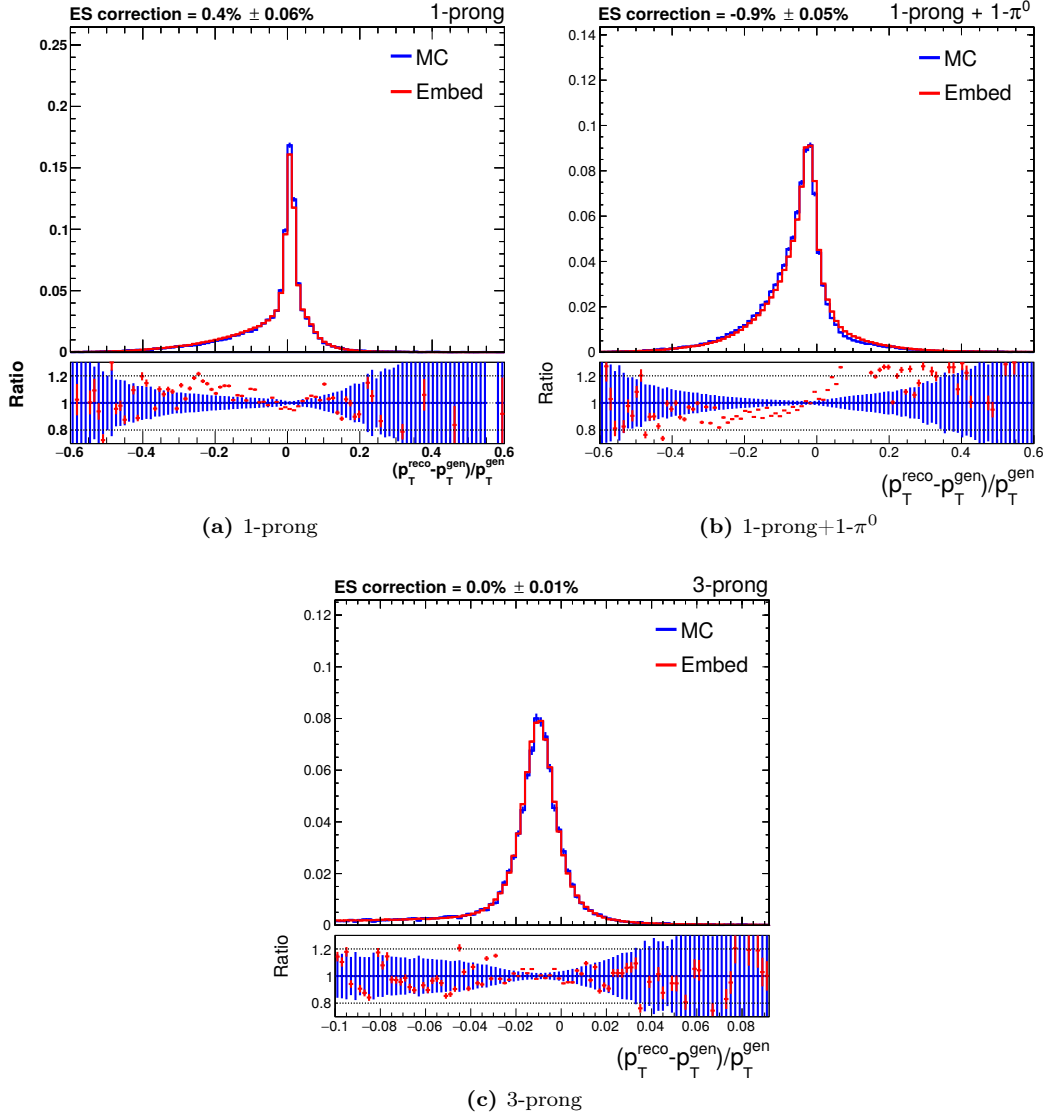


Figure 3.14: Comparison of the τ_h p_T resolutions and energy scales in the 2017 MC and embedded samples is shown for the 1-prong (left), 1-prong+ $1-\pi^0$ (centre), and 3-prong (right) tau decay modes.

2016 MC is 0% and 9.5% for the 1-prong and 1-prong+ $1-\pi^0$ decay modes respectively. The $\mu \rightarrow \tau_h$ corrections are -0.2% and 1.2% for the 1-prong and 1-prong+ $1-\pi^0$ decay modes respectively. The 2017 energy scale corrections were measured by the author as described below.

The corrections are measured using samples of $Z/\gamma^* \rightarrow \tau_e \tau_h$ and $Z/\gamma^* \rightarrow \tau_\mu \tau_h$ events. The $Z/\gamma^* \rightarrow \tau_e \tau_h$ sample is selected by requiring:

- At least one electron with $I_{\text{rel}} < 0.15$ and $p_T > 28$ GeV that passes the *loose* ID and the single-electron trigger.
- At least one hadronic tau candidate with $p_T > 20$ that passes the HPS algorithm, the *tight* working-point of the MVA-isolation, the *loose* working-point of the anti-muon discriminator, and the *tight* working-point of the anti-electron discriminator.

The $Z/\gamma^* \rightarrow \tau_\mu \tau_h$ sample is selected by requiring:

- At least one muon with $I_{\text{rel}} < 0.15$ and $p_T > 25$ GeV, that passes the *medium* ID and the single-muon trigger.
- At least one hadronic tau candidate with $p_T > 20$ that passes the HPS algorithm, the *tight* working-point of the MVA-isolation, the *tight* working-point of the anti-muon discriminator, and the *very-loose* working-point of the anti-electron discriminator.

For both samples the following selections are also imposed:

- The lepton and tau should have opposite charge and be separated by $\Delta R > 0.5$. If more than one pair meet this requirement the most isolated pair is retained.
- $m_T(l, E_T^{\text{miss}})$ must be smaller than 50 GeV.
- Events with additional electrons or muons are vetoed.

Events are split into two categories depending on whether the hadronic tau candidate is a 1-prong or 1-prong+1- π^0 tau.

The $Z/\gamma^* \rightarrow ee$ and $Z/\gamma^* \rightarrow \mu\mu$ processes are treated as signal for the measurements of the $e \rightarrow \tau_h$ and $\mu \rightarrow \tau_h$ energy scales respectively. Many signal templates are generated each with a different shift applied to the $l \rightarrow \tau_h$ energy scale. The shifts range between -10% to $+10\%$ in increments of 1%. For the $\mu \rightarrow \tau_h$ measurement, additional templates are produced in 0.2% increments in the range -2% to $+1\%$ for the 1-prong category and 0% to 3% for the 1-prong+1- π^0 category. These templates are then interpolated between using a horizontal morphing algorithm [133], where the $l \rightarrow \tau_h$ energy scale correction is used as a horizontal morphing parameter.

The background processes that are considered are $Z/\gamma^* \rightarrow \tau\tau$, QCD, W+jets, $t\bar{t}$, diboson, and single-top. The methods used to estimate the QCD and W+jets backgrounds will be described in Section 3.5.4. All other backgrounds are modelled using MC. The data/MC corrections described in this section are also applied.

The measurement of the energy scale correction is performed using a binned maximum likelihood fit using the $\tau_l\tau_h$ mass as the discriminating variable. The following systematic uncertainties affecting the normalisations and shapes of the signal and background templates are also included:

Normalisation uncertainties:

- **Luminosity:** The luminosity uncertainty amounts to 2.3%, which is applied to all processes where the yield is estimated from MC.
- **Muon ID/isolation efficiency:** The uncertainty on the muon ID/isolation efficiency is 2%.
- **Electron ID/isolation efficiency:** The uncertainty on the electron ID/isolation efficiency is 2%.
- **Tau tracking/ID efficiency:** A 5% uncertainty on the tau tracking/ID efficiency is applied.
- **$j \rightarrow \tau_h$ fake-rate:** A 20% uncertainty is applied for all processes containing a $j \rightarrow \tau_h$ fake where the yield is estimated from MC.
- **$\mu \rightarrow \tau_h$ fake-rate:** A 20% uncertainty is applied to the $Z \rightarrow \mu\mu$ process.
- **$e \rightarrow \tau_h$ fake-rate:** A 20% uncertainty is applied to the $Z \rightarrow ee$ process.
- **Trigger efficiency:** A 2% uncertainty on the single-muon and single-electron trigger efficiencies is applied.
- **Background normalisation uncertainties:**

- A 4% uncertainty on the $Z/\gamma^* \rightarrow ll$ and $Z/\gamma^* \rightarrow \tau\tau$ yields to account for the uncertainty on the Drell-Yan cross section.
- A 5% uncertainty is on the diboson and single-top yields to account for the uncertainties on the production cross sections.
- A 6% uncertainty is on the $t\bar{t}$ yields to account for the uncertainties on the production cross section.
- A 10% uncertainty on the W+jets yield.
- A 10% uncertainty on the QCD yield.

Shape uncertainties:

- **Hadronic tau energy scale:** Three uncorrelated uncertainties each that affect only one of the three tau decay modes. Each uncertainty varies the hadronic tau energy scale by 3%.
- **Statistical uncertainties:** Limitations in the statistical sizes of the signal and background templates are accounted for by the Barlow-Beeston method [131, 132].
- **Top p_T reweighting:** Uncertainty between no and twice the correction on the reweighting applied to $t\bar{t}$ events.
- **Drell-Yan Z p_T reweighting:** Uncertainty between no and twice the correction on the reweighting applied to $Z \rightarrow \mu\mu$ events.

The post-fit data/MC comparisons are shown in Figure 3.15. Good agreement between the background-prediction and the data is observed. The results of the maximum likelihood scans of the $l \rightarrow \tau_h$ energy scales are shown in Figures 3.16–3.17. The correction is taken from the minimum of these scans. The uncertainties from the scans are taken as the systematic error on the correction, but for simplicity these are taken as symmetric uncertainties. The measured energy scale shifts and uncertainties for $e \rightarrow \tau_h$ are $1\% \pm 1.4\%$ for the 1-prong decay mode, and $4.1\% \pm 1.8\%$ for 1-prong+1- π^0 decay mode. The measured energy scale shifts and uncertainties for $\mu \rightarrow \tau_h$ are $-0.1\% \pm 0.3\%$ for the 1-prong decay mode, and $1.2\% \pm 1\%$ for 1-prong+1- π^0 decay mode.

3.5 Background methods

3.5.1 Modelling backgrounds using MC

To model background processes using simulation, the generated MC events are first corrected using the procedures described in Section 3.4. The events are then scaled by additional weights given by

$$\text{weight} = \frac{L_{\text{int}} \cdot \sigma \cdot \text{BR} \cdot w_{\text{gen}}}{N_{\text{effective}}}, \quad (3.22)$$

where $\sigma \cdot \text{BR}$ is the cross section times branching ratio for the process, L_{int} is the integrated luminosity, w_{gen} is the generator weight, and $N_{\text{effective}}$ is the effective event number, defined as the sum of the generator weights for all events $N_{\text{effective}} = \sum_i w_{\text{gen}}^i$.

The Drell-Yan Z/γ^* +jets and QCD-induced W+jets processes are scaled to their cross sections computed at NNLO precision using FEWZ [134]. NNLO accurate cross sections are also used for $t\bar{t}$ events computed with Top++ [135], and for single top quark events in the tW-channel [136]. The cross section for single top production in the t-channel is computed using HatHor [137, 138] at NLO precision. The remaining backgrounds, which include diboson and electroweak W/Z/ γ^* +jets events, are scaled to the cross sections computed by their corresponding event generators.

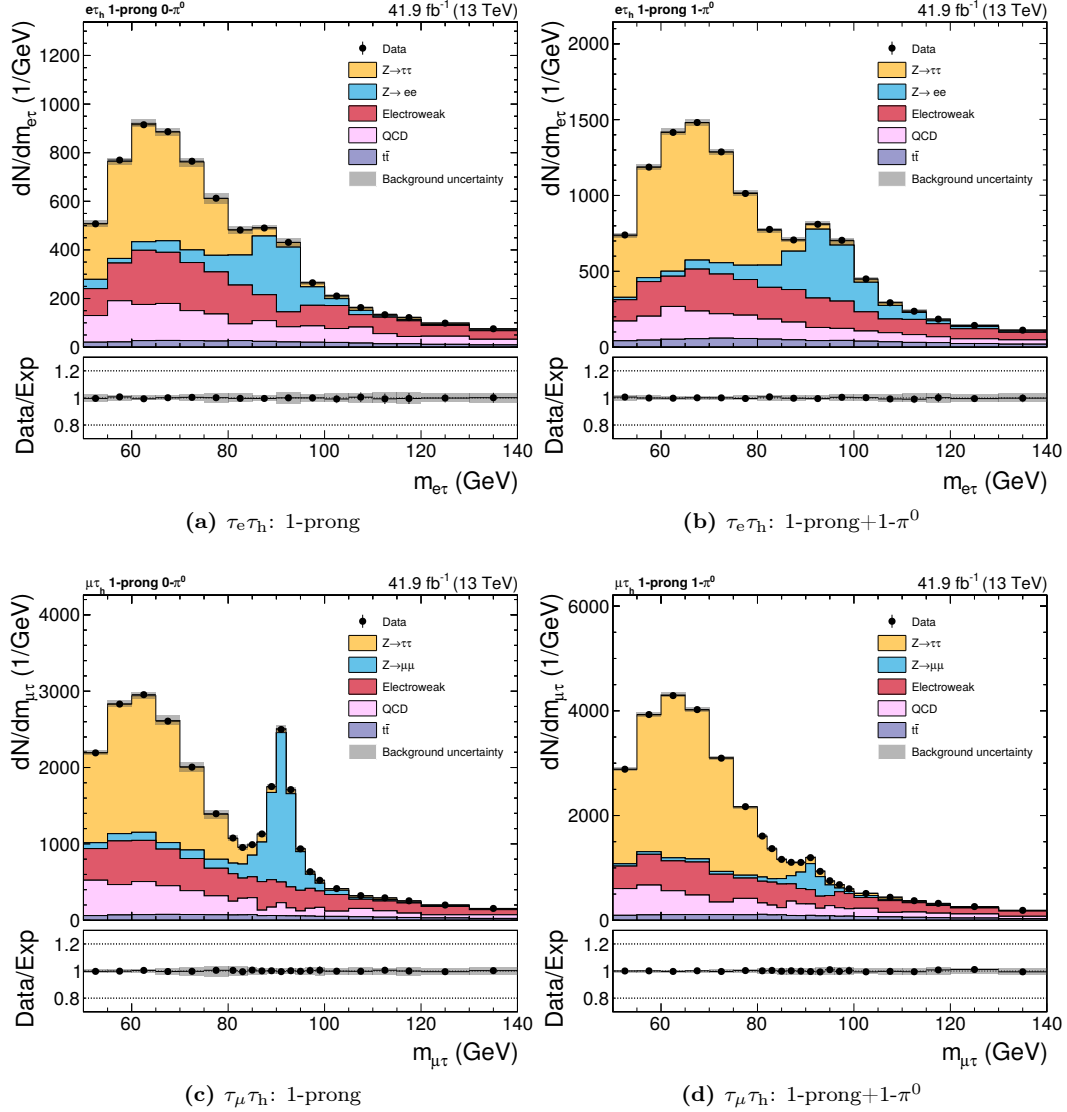


Figure 3.15: The post-fit distributions showing the data/MC agreement for 1-prong (left) and 1-prong+ $1-\pi^0$ (right) categories for the $\tau_e \tau_h$ (upper) and $\tau_\mu \tau_h$ (lower) channels. The uncertainty band shown in shaded grey includes both systematic and statistical uncertainties.

3.5.2 Modelling $Z/\gamma^* \rightarrow \tau\tau$ using embedding

The $Z/\gamma^* \rightarrow \tau\tau$ process is the dominant background for SM Higgs production and therefore a good description of this background is essential. The embedded samples described in Section 3.3.2 are used to improve the modelling of the $Z/\gamma^* \rightarrow \tau\tau$ process. The corrections described in Section 3.4 along with the generator weights described in Section 3.3.2 are applied to the embedded events. Additional weights are also needed to correct for the inefficiencies in the initial dimuon reconstruction. The reconstruction efficiency can be factorised into two components: the efficiency for the muons to pass the ID requirements, ϵ_{ID} , and the double-muon trigger efficiency, ϵ_{trig} . The embedded events are then scaled by

$$\text{weight} = \frac{1}{\epsilon_{\text{ID}} \cdot \epsilon_{\text{trig}}}. \quad (3.23)$$

The ID and trigger efficiencies are measured in data using the same tag-and-probe method described in Section 3.4.6. The efficiencies for the high p_T leg ($p_T > 17$ GeV) of the trigger are

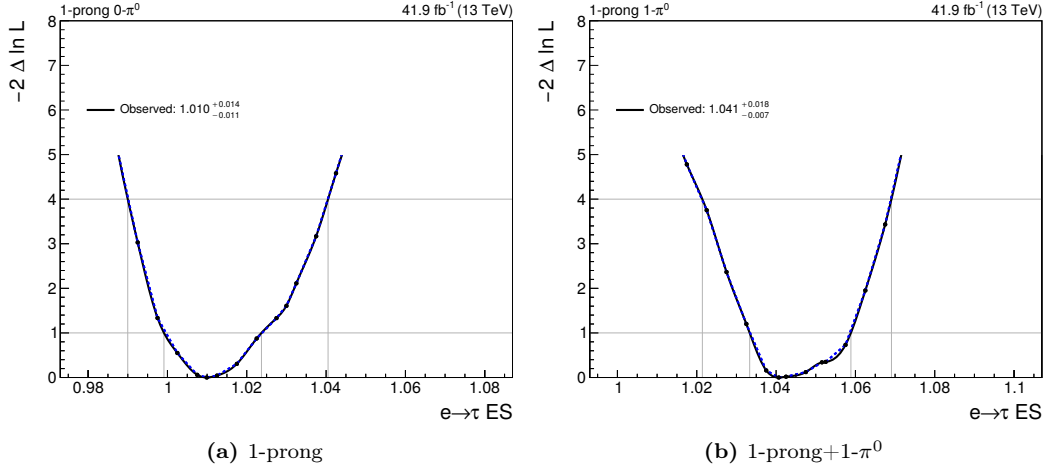


Figure 3.16: The likelihood scans showing the best-fit value of $e \rightarrow \tau_h$ energy scale correction for the 1-prong (left) and 1-prong+1- π^0 (right) measurements.

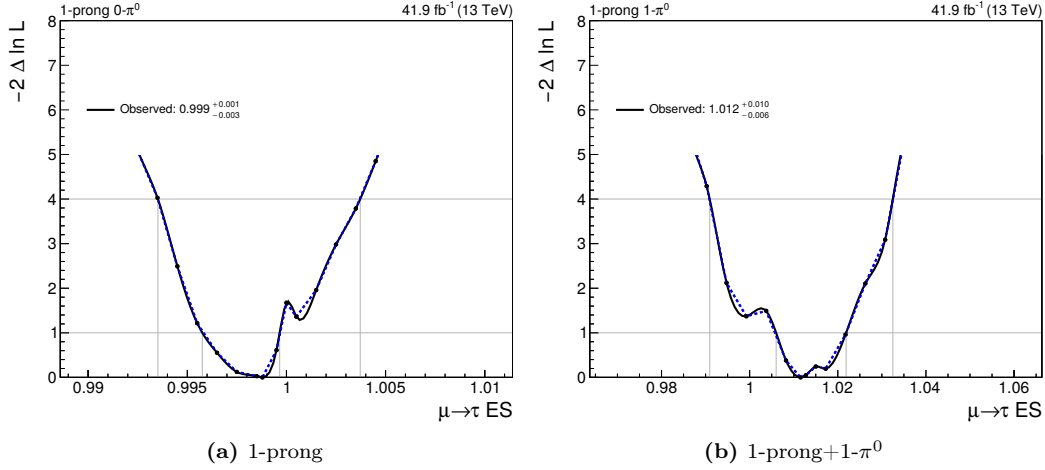


Figure 3.17: The likelihood scans showing the best-fit value of $\mu \rightarrow \tau_h$ energy scale correction for the 1-prong (left) and 1-prong+1- π^0 (right) measurements.

shown as a function of p_T and η in Figures 3.18 (a) and (b) respectively. The efficiency becomes approximately flat with p_T above 20 GeV, but depends quite strongly on η . In addition to the yield estimation, measuring the efficiencies as a function of these variables also improves the description of the η and p_T distributions, as these would otherwise be biased by the differential selection efficiencies.

An important check of the embedded samples is performed by embedding simulated muons rather than taus into the selected $Z/\gamma^* \rightarrow \mu\mu$ events. Since the simulated muons have identical kinematics to the selected muons we expect to see that the embedded distributions match the data closely, and any disagreement would be indicative of biases in the embedding procedure. Figure 3.19 shows the agreement of the muon p_T and η distribution with the data. In these examples the estimated yield is observed to agree with the data to within 2%, and the embedded samples appear to describe the shapes of the p_T and η distributions well.

The agreement of the embedded $Z/\gamma^* \rightarrow \tau\tau$ with the data is checked by selecting events in the $\tau_e\tau_\mu$, $\tau_e\tau_h$, $\tau_\mu\tau_h$, and $\tau_h\tau_h$ channels. These final states have a non-negligible contribution from

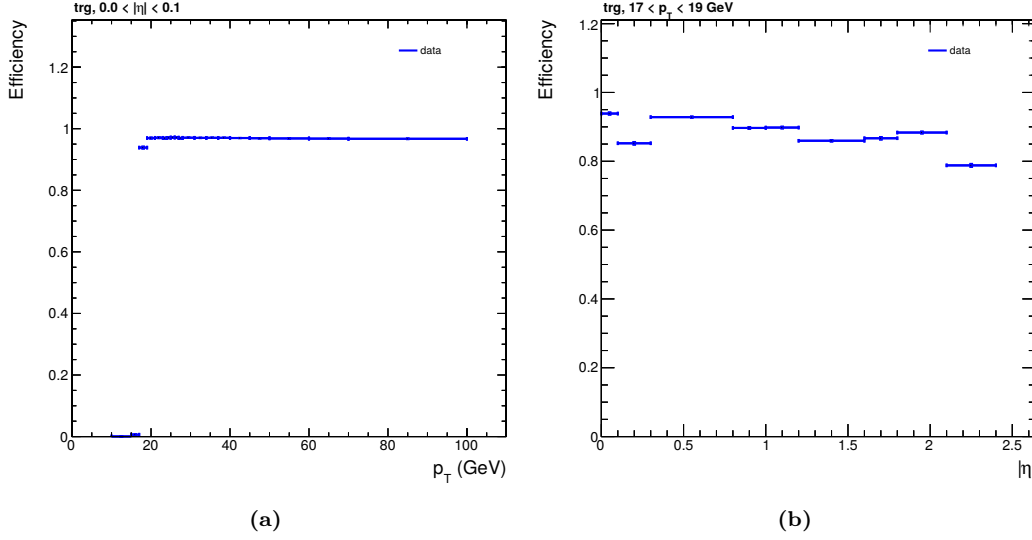


Figure 3.18: The measured efficiency of the high p_T leg of the double-muon trigger used to select $Z/\gamma^* \rightarrow \mu\mu$ events for the embedded samples. The efficiency is shown as a function of p_T (left) and η (right).

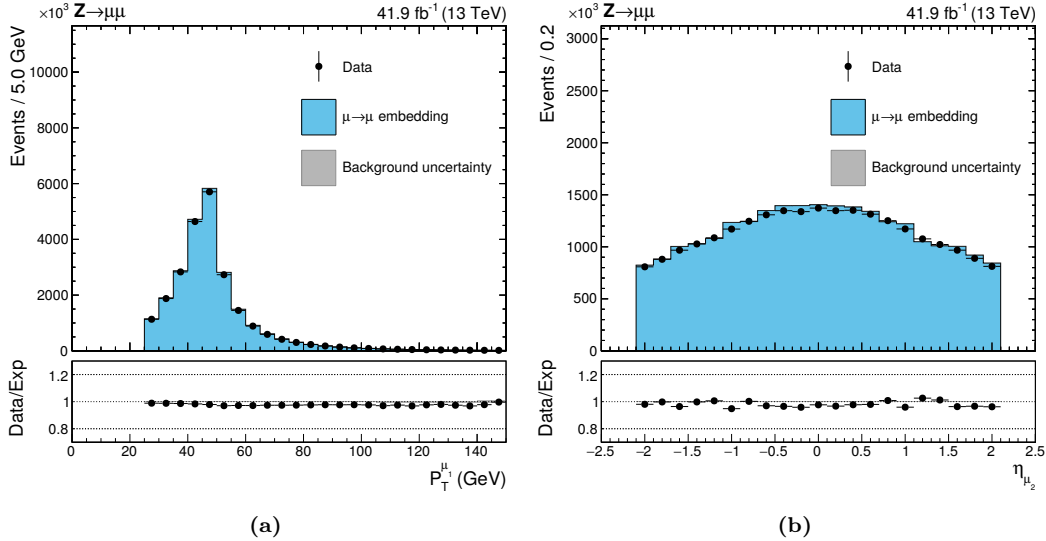


Figure 3.19: Comparisons of the data/embedding agreement for the muon p_T (left) and η (right) distributions are shown for $\mu \rightarrow \mu$ embedded events. The uncertainty band shown in shaded grey includes only the statistical uncertainties.

other background processes but relatively high purity of $Z/\gamma^* \rightarrow \tau\tau$ may be obtained using the $m_{\tau\tau}$ variable (defined in Section 5.3), for $m_{\tau\tau}$ values close to the Z mass. The $m_{\tau\tau}$ distributions in the $\tau_\mu\tau_h$ and $\tau_e\tau_\mu$ final states are shown in Figure 3.20; good agreement with the data in both the yields and the shapes is observed.

3.5.3 The fake-factor method

Backgrounds in which a jet fakes a τ_h can be difficult to model using MC due to the poor description of the jet $\rightarrow \tau_h$ fake rate in simulation. In addition, the small probability of a jet being misidentified as a τ_h necessitates the production of high-statistics MC samples at a significant computational expense. These shortcomings motivate the use of data-driven estimates for these processes. One

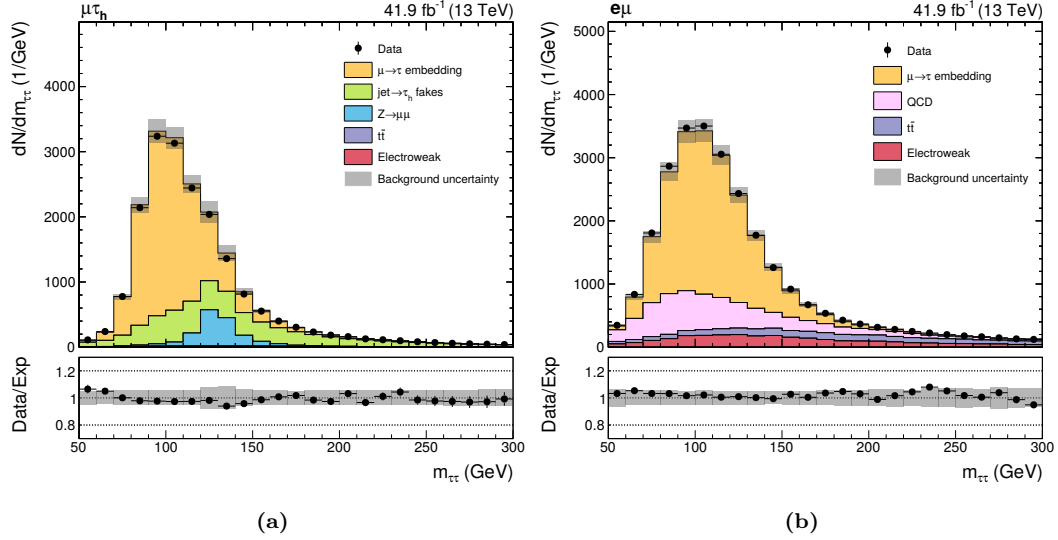


Figure 3.20: Comparison of the data/embedding agreement for the $m_{\tau\tau}$ distributions in the $\tau_\mu\tau_h$ and $\tau_e\tau_\mu$ final states. The uncertainty band shown in shaded grey includes both systematic and statistical uncertainties.

such procedure is the fake-factor (FF) method. The following text is intended to give an overview of the FF method, a more detailed description can be found in Ref. [139].

The FF method models $\text{jet} \rightarrow \tau_h$ backgrounds from side-band regions in data, in which a τ_h candidate fails the nominal *tight/medium* MVA-isolation requirement applied in the signal-regions (SR) but passes the *very-loose* working point. These side-band regions will be referred to here as application-regions (AR). The basic idea of the method is to scale events in the AR by ratios, FF_i , defined as

$$\text{FF}_i = \frac{N(\text{tight}/\text{medium})}{N(\text{very-loose} \&\& \text{!tight}/\text{!medium})}, \quad (3.24)$$

where $N(X)$ is the number of events meeting the identification requirement X . These FF_i ratios are measured in control-regions enriched in $\text{jet} \rightarrow \tau_h$ events referred to here as determination-regions (DR). The FF_i are measured in bins of N_{jets} , and the dependence on p_T is accounted for using fitted functions. For the MSSM analysis described in Chapter 4 the FF_i are also measured in bins of tau decay mode. The FF_i are measured separately for the dominant $\text{jet} \rightarrow \tau_h$ processes, which for the $\tau_l\tau_h$ channels includes QCD, $W+\text{jets}$ and $t\bar{t}$, and for the $\tau_h\tau_h$ channel includes only QCD. For the $\tau_l\tau_h$ channels the FF_i measured for the $W+\text{jets}$ process are also used to model sub-dominant processes such as $Z+\text{jets}$, whereas for the $\tau_h\tau_h$ channel the FF_i measured for QCD events are used to model these processes. The FF_i are also corrected to account for extrapolations from the DR to the SR, and for missing variable dependencies (via non-closure corrections).

The FF_i measured for the different processes are combined into an overall factor, $\overline{\text{FF}}$, using

$$\overline{\text{FF}} = \sum_i f_i \cdot \text{FF}_i, \quad (3.25)$$

where the factor f_i is defined as

$$f_i = \frac{n_i}{\sum_j n_j}, \quad (3.26)$$

which is the fraction of events with a $\text{jet} \rightarrow \tau_h$ originating from process i over the total number of $\text{jet} \rightarrow \tau_h$ events for all processes in the AR. The f_i are also measured as a function of variables such as $p_T^{\tau_h}$ or the ditau mass.

For the $\tau_l\tau_h$ channels the $\text{jet} \rightarrow \tau_h$ background is estimated by applying the measured $\overline{\text{FF}}$ to the events in the AR. The contribution of processes producing genuine τ_h or $l \rightarrow \tau_h$ candidates is subtracted using simulated events. For the $\tau_h\tau_h$ channel the procedure is slightly more complicated due to the presence of two τ_h candidates, where either one or both may originate from a jet fake. The dominant process in this channel is QCD where both τ_h candidates are faked by jets, the other processes overwhelmingly produce only one fake τ_h candidate. The method for this channel is to treat both τ_h candidates sequentially as the fake, performing the same procedure described above in each instance to produce two templates which are then summed together and scaled by a factor of $\frac{1}{2}$. This procedure alone would lead to an underestimation of processes producing only one fake τ_h by a factor of $\frac{1}{2}$. To compensate for this underestimation the FF_i for such processes are scaled up by a factor of 2. Examples of the pre-fit modelling of the leading and sub-leading τ_h p_T in the $\tau_h\tau_h$ final state are shown in Figure 3.21. Good agreement with the data is observed.

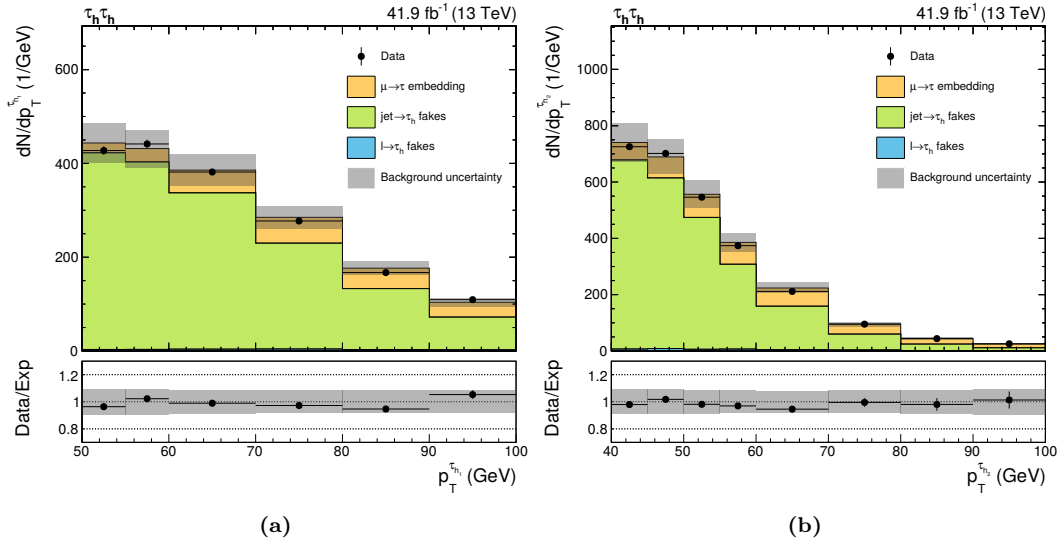


Figure 3.21: Control plots showing the pre-fit modelling of the leading (left) and sub-leading (right) τ_h candidates in the $\tau_h\tau_h$ final state. Events are selected requiring the tau pair to have large mass to enhance the purity of $\text{jet} \rightarrow \tau_h$ events. The uncertainty band shown in shaded grey includes the full systematic uncertainty model.

3.5.4 Alternative $\text{jet} \rightarrow \tau_h$ estimation in the $\tau_l\tau_h$ channels

An alternative procedure to the FF method may be used to estimate the $\text{jet} \rightarrow \tau_h$ background in the $\tau_e\tau_h$ and $\tau_\mu\tau_h$ final states. This will be referred to as the "classical method" in this document. This method is used as a cross check of the FF method for the $\phi \rightarrow \tau\tau$ analysis, and for the measurements of several corrections described in Section 3.4. The two largest processes producing $\text{jet} \rightarrow \tau_h$ fakes in the $\tau_l\tau_h$ channels are QCD and W +jets. In the classical method these backgrounds are estimated using fully-data-driven and semi-data-driven procedures, respectively. The remaining sub-dominant processes producing $\text{jet} \rightarrow \tau_h$ fakes are modelled entirely from simulation.

3.5.4.1 QCD modelling

The QCD background is modelled by inverting the requirement for $\tau_l\tau_h$ pair to be oppositely charged. This exploits the similarity between the processes producing opposite-sign and same-sign events, as both tend to be produced at comparable rates with variable distributions that are very similar. After selecting the same-sign events in data the contribution from other backgrounds

is subtracted using simulation. The resulting template is then scaled by a factor R_{QCD} which accounts for the difference in the rates for the opposite-sign and same-sign events. This factor is measured in a side-band region where the τ_l is anti-isolated. The R_{QCD} factor may vary depending on the event selections but is typically ≈ 1.0 – 1.2 . For the correction measurements described in Sections 3.4.7 and 3.4.10 R_{QCD} is measured as 1.12 ± 0.07 and 1.13 ± 0.13 for the $\tau_\mu\tau_h$ and $\tau_e\tau_h$ channels respectively.

The modelling of the QCD background is illustrated in Figure 3.22 for the $\tau_\mu\tau_h$ final state. The left inset of Figure 3.22 shows the data/background agreement in a side-band region where the τ_μ is anti-isolated. This is the same region used to measure R_{QCD} so good agreement in the yields is expected. However, this region is mainly used to check that the shapes of the variable distributions are modelled well, as is observed. The right inset shows the data/background agreement in a second orthogonal side-band region in which the τ_h is anti-isolated. The W +jets contribution (included in the "Electroweak" template) is larger in this region, making it slightly more difficult to judge the QCD modelling, but nevertheless there is no visible disagreement with the data.

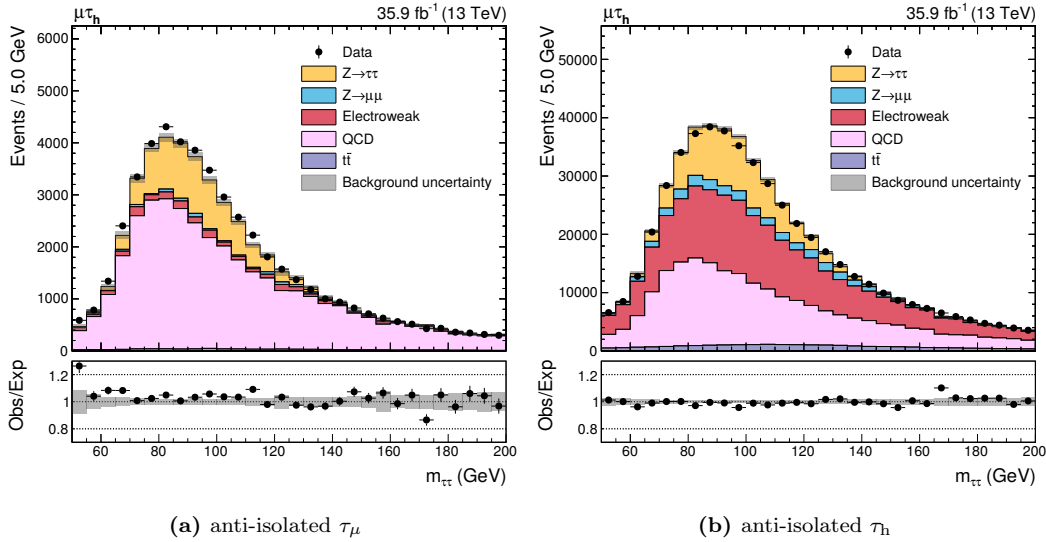


Figure 3.22: Comparisons of the $m_{\tau\tau}$ data/background agreement in the $\tau_\mu\tau_h$ final state are shown for side-band regions in which the τ_μ (left) and τ_h (right) are anti-isolated. The uncertainty band shown in shaded grey includes only statistical uncertainties.

3.5.4.2 W +jets modelling

The shape of the W +jets background is taken from MC and subsequently scaled to a yield estimated using a data-driven method. The procedure is thus referred to as a "semi-data-driven" method. The yield estimation makes use of the $m_T > 70$ GeV region, which has a high purity of W +jets events. The basic idea behind the method is to subtract non- W backgrounds from the data to estimate the W +jets yield in the high- m_T region, which is then extrapolated to the low- m_T signal region. However, the background to be subtracted from the data includes a non-negligible contribution from QCD events, and since the QCD estimate also requires the subtraction of the W +jet events from the same-sign data, the two estimates are tied together.

To take into account the QCD contribution two high- m_T control regions in which the $\tau_l\tau_h$ pair have opposite-sign and same-sign charges are selected. The contribution from backgrounds other than QCD and W +jets are subtracted from the data using MC. Two equations are then

constructed relating the QCD and W+jets yields to the data minus "other" backgrounds in the two regions. The W+jets and QCD yields in the opposite-sign region are additionally related to the yields in the same-sign region using the factors R_W and R_{QCD} , where R_W is the ratio of opposite-sign to same-sign W+jets events estimated from simulation. This results in two equations with two unknowns which can then be solved to estimate the W+jets yield in the same-sign high- m_T region:

$$\begin{aligned}
 N_{\text{data}}^{\text{SS,high-}m_T} - N_{\text{other}}^{\text{SS,high-}m_T} &= N_{\text{QCD}}^{\text{SS,high-}m_T} + N_W^{\text{SS,high-}m_T} \\
 N_{\text{data}}^{\text{OS,high-}m_T} - N_{\text{other}}^{\text{OS,high-}m_T} &= N_{\text{QCD}}^{\text{OS,high-}m_T} + N_W^{\text{OS,high-}m_T} \\
 &= R_{\text{QCD}} \cdot N_{\text{QCD}}^{\text{SS,high-}m_T} + R_W \cdot N_W^{\text{SS,high-}m_T} \\
 \Rightarrow N_W^{\text{SS,high-}m_T} &= \frac{N_{\text{data}}^{\text{OS,high-}m_T} - N_{\text{other}}^{\text{OS,high-}m_T} - R_{\text{QCD}} \cdot (N_{\text{data}}^{\text{SS,high-}m_T} - N_{\text{other}}^{\text{SS,high-}m_T})}{R_W - R_{\text{QCD}}},
 \end{aligned} \tag{3.27}$$

where N_x^y is the yield of process x in region y , and the labels OS and SS indicate the opposite-sign and same-sign regions respectively. This method relies on the values of R_{QCD} and R_W being not too similar, which is usually the case as typically $R_{\text{QCD}} \approx 1$ and $R_W \approx 4$.

Once $N_W^{\text{SS,high-}m_T}$ has been determined the estimate must be extrapolated from the same-sign high- m_T control region to the opposite-sign low- m_T signal region using

$$N_W^{\text{OS,low-}m_T} = N_W^{\text{SS,high-}m_T} \cdot R_W \cdot F_{\text{high} \rightarrow \text{low-}m_T}, \tag{3.28}$$

where $F_{\text{high} \rightarrow \text{low-}m_T}$ is the high- m_T to low- m_T extrapolation factor defined as the ratio $\frac{N_W^{\text{OS,low-}m_T}}{N_W^{\text{OS,high-}m_T}}$ estimated from MC.

The agreement with data is illustrated in Figure 3.23, which shows the m_T distribution after applying a typical signal-region-like selection, and after additionally requiring the τ_h candidate to be anti-isolated to enhance the purity of W+jets events. In both cases the yields and shapes are observed to agree well with the data.

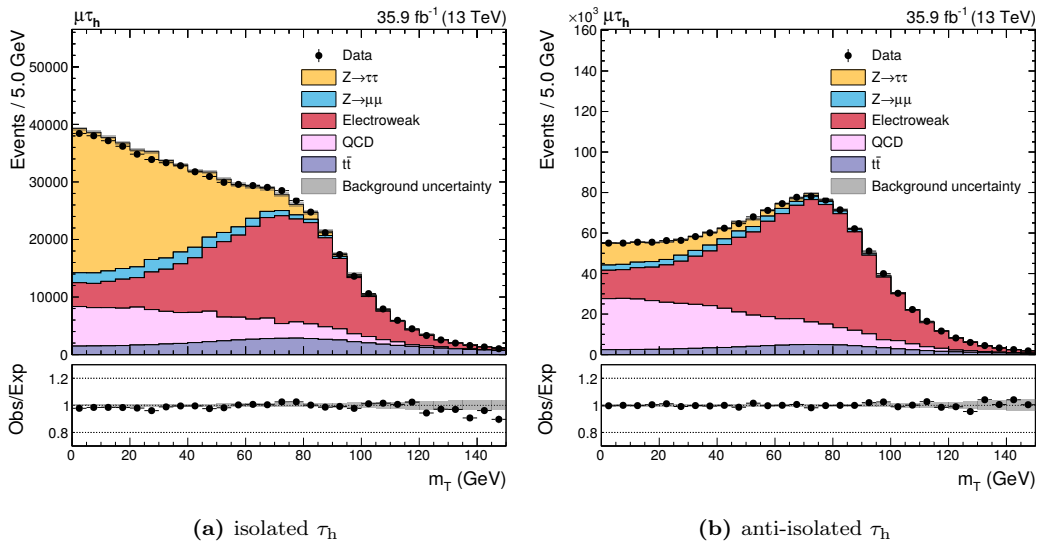


Figure 3.23: Comparisons of the m_T data/background agreement in the $\tau_\mu \tau_h$ final state are shown after applying signal-region-like selections (left) and requiring the τ_h to be anti-isolated to increase the purity of W+jets events (right). The uncertainty band shown in shaded grey includes only statistical uncertainties.

3.5.5 Modelling QCD in the $\tau_e\tau_\mu$ channel

The QCD background in the $\tau_e\tau_\mu$ channel is estimated from same-sign data analogously to the method described in the Section 3.5.4.1. The procedure is more complicated for this channel however, as the ratio of opposite-sign to same-sign events, $\dot{R}_{\text{QCD}}$, can depend strongly on other variables. The size of $\dot{R}_{\text{QCD}}$ is also typically ≈ 2 . The dominant QCD process producing $e\mu$ pairs is $b\bar{b}$, where the leptons originate from semi-leptonic b or c decays. Opposite-sign events are more common as they can originate directly from decays such as $b \rightarrow e^- + X$, $\bar{b} \rightarrow \mu^+ + X$, whereas same-sign events originate from decay chains such as $b \rightarrow e^- + X$, $\bar{b} \rightarrow \bar{c} + X \rightarrow \mu^- + Y$. The lepton arising from the c quark decay tends to be softer in p_T and less isolated, reducing the number of such events passing the p_T and isolation requirements. The softer p_T of the leptons produced in the same-sign configuration also results in the $\dot{R}_{\text{QCD}}$ depending on $p_T^{e/\mu}$.

Although the majority of QCD events are produced via $b\bar{b}$, there is also a contribution from other processes producing fake leptons. These processes tend have $\dot{R}_{\text{QCD}}$ closer to unity. $\dot{R}_{\text{QCD}}$ therefore also depends on any variable that changes the fraction of $b\bar{b}$ events. $\dot{R}_{\text{QCD}}$ is measured as a function of the strongest variable dependencies which are found to be: $p_T^{e/\mu}$, N_{jets} , and the ΔR separation of the $e\mu$ pair, $\Delta R_{e\mu}$. Determining $\dot{R}_{\text{QCD}}$ as a function of these variables is found to be sufficient to describe well all discriminating variables, as well as variables important for event selection and categorisation.

The $\dot{R}_{\text{QCD}}$ are measured in a side-band region in which the muon is anti-isolated ($0.2 < I_{\text{rel}}^\mu < 0.5$), which will be called the DR. This region is chosen over the electron anti-isolated side-band as the dependence of $\dot{R}_{\text{QCD}}$ on I_{rel}^μ is found to be fairly weak compared to the dependence on I_{rel}^e . This is illustrated in Figure 3.24, which shows the dependence of $\dot{R}_{\text{QCD}}$ on $I_{\text{rel}}^{e/\mu}$ as measured in a side-band region in which either the electron or the muon are anti-isolated. The measurement of $\dot{R}_{\text{QCD}}$ is performed by selecting opposite-sign and same-sign events in data before subtracting the contribution from non-QCD processes using MC. $\dot{R}_{\text{QCD}}$ is then given as the ratio $\frac{N_{\text{data}}^{\text{OS}} - N_{\text{MC}}^{\text{OS}}}{N_{\text{data}}^{\text{SS}} - N_{\text{MC}}^{\text{SS}}}$.

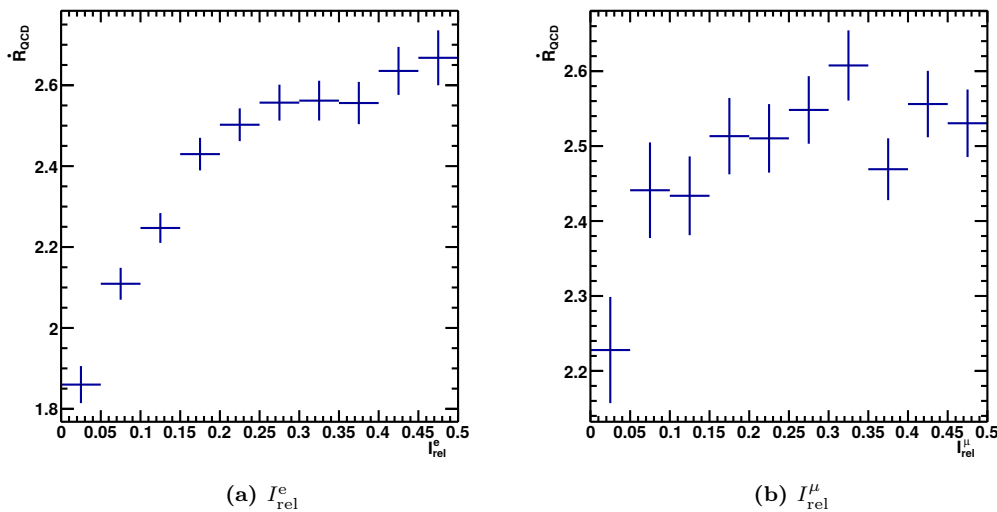


Figure 3.24: The dependence of $\dot{R}_{\text{QCD}}$ on I_{rel}^e (left) and I_{rel}^μ (right). $\dot{R}_{\text{QCD}}$ is measured in a side-band region in which either the electron or the muon is anti-isolated.

There are insufficient event statistics in the DR to simply measure $\dot{R}_{\text{QCD}}$ in many bins of the four parametrising variables. Instead the following procedure is used:

1. The N_{jets} dependence is taken into account using the minimum number of bins. Two bins with $N_{\text{jets}} = 0$ and $N_{\text{jets}} \geq 1$ is found to be sufficient. The dependence on higher jet multiplicities is checked for but found to be consistent within the statistical uncertainties.
2. In each of the two N_{jets} bins the $\Delta R_{e\mu}$ dependence is fitted using linear functions to reduce the influence of statistical fluctuations. The linear fits are illustrated in Figure 3.25.
3. An initial set of $\dot{R}_{\text{QCD}} = \dot{R}_{\text{QCD}}(N_{\text{jets}}, \Delta R_{e\mu})$ is determined as described above. These factors are then applied to the same-sign events in the DR and compared to the opposite-sign events. The resulting non-closure in bins of $p_{\text{T}}^{e/\mu}$ is applied as a residual correction to take into account the $p_{\text{T}}^{e/\mu}$ dependence. Figure 3.26 illustrates the size of these residual corrections.

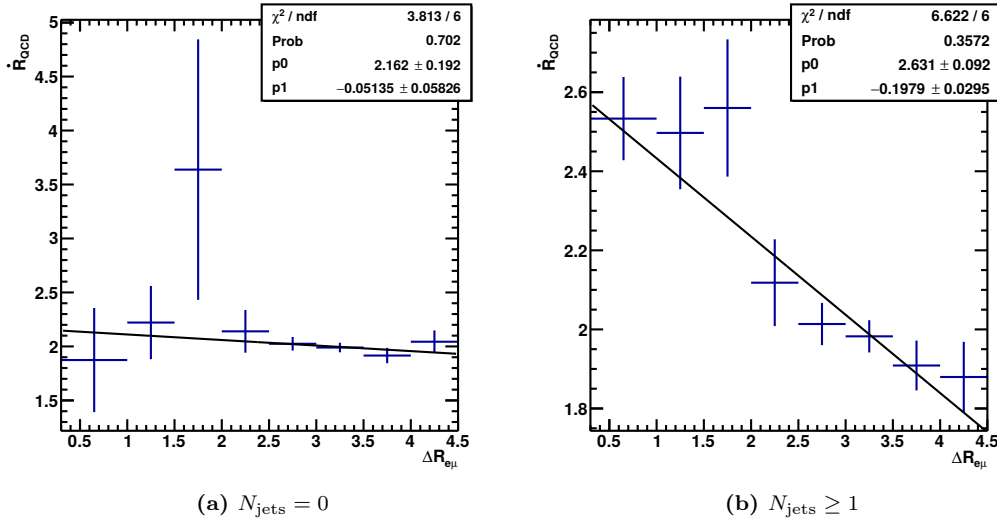


Figure 3.25: The dependence of $\dot{R}_{\text{QCD}}$ on $\Delta R_{e\mu}$ in the two N_{jets} bins. The dependence is fitted using linear functions.

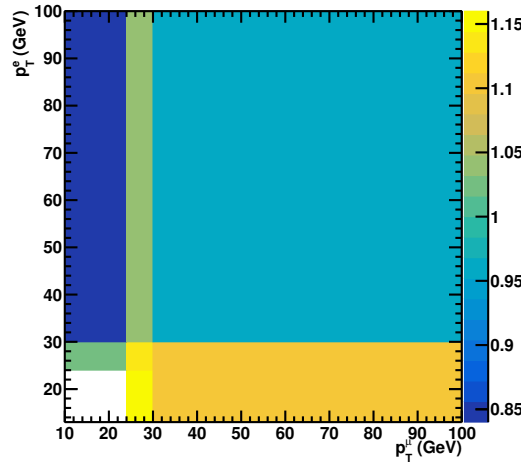


Figure 3.26: The $p_{\text{T}}^{e/\mu}$ dependent residual corrections used to correct for the non-closure of the electron/muon p_{T} distributions after application of $\dot{R}_{\text{QCD}} = \dot{R}_{\text{QCD}}(N_{\text{jets}}, \Delta R_{e\mu})$.

The determined $\dot{R}_{\text{QCD}} = \dot{R}_{\text{QCD}}(N_{\text{jets}}, \Delta R_{e\mu}, p_{\text{T}}^e, p_{\text{T}}^\mu)$ are then applied to the same-sign events in the DR to check that the chosen parametrisations, as well as the above procedure used to measure them, is sufficient to adequately describe all important variables. Examples of the comparisons for

several variables are shown in Figure 3.27. All shown variables with the exception of $\Delta\phi_{jj}$, which is the azimuthal separation of the two leading jets, have been defined previously. The uncertainty band shown in shaded grey is formed by propagating the uncertainties on the linear fit parameters as systematic uncertainties, which are then added in quadrature with the statistical errors. The background templates are observed to match the data well within the uncertainties, and there is no evidence of any residual non-closures.

The QCD method is validated using a second side-band region known as the validation-region (VR), which is orthogonal to both the DR and the signal-region. This region is defined by selecting events where the electron is anti-isolated ($0.15 < I_{\text{rel}}^e < 0.5$) but the muon is isolated ($I_{\text{rel}}^\mu < 0.2$). As discussed above $\dot{R}_{\text{QCD}}$ depends strongly on I_{rel}^e , and as such the application of the $\dot{R}_{\text{QCD}}$ measured previously would not describe the QCD well in this case. Therefore, separate $\dot{R}_{\text{QCD}}$ are measured using exactly the same procedure described previously. These factors are measured using events where both the electron and muon are anti-isolated ($0.15 < I_{\text{rel}}^e < 0.5$ and $0.2 < I_{\text{rel}}^\mu < 0.5$), and then applied to events in the VR. Relatively good agreement with the data is observed, as illustrated in Figure 3.28. There is a slight overestimation $\mathcal{O}(5\%)$ in the yield, but this is fairly small compared to the systematic uncertainties. The small residual non-closure is a consequence of the soft dependence of $\dot{R}_{\text{QCD}}$ on I_{rel}^μ . An additional systematic uncertainty parametrised in bins of $p_{\text{T}}^{e/\mu}$ is used to account for this effect. This is estimated from the observed non-closure in the VR.

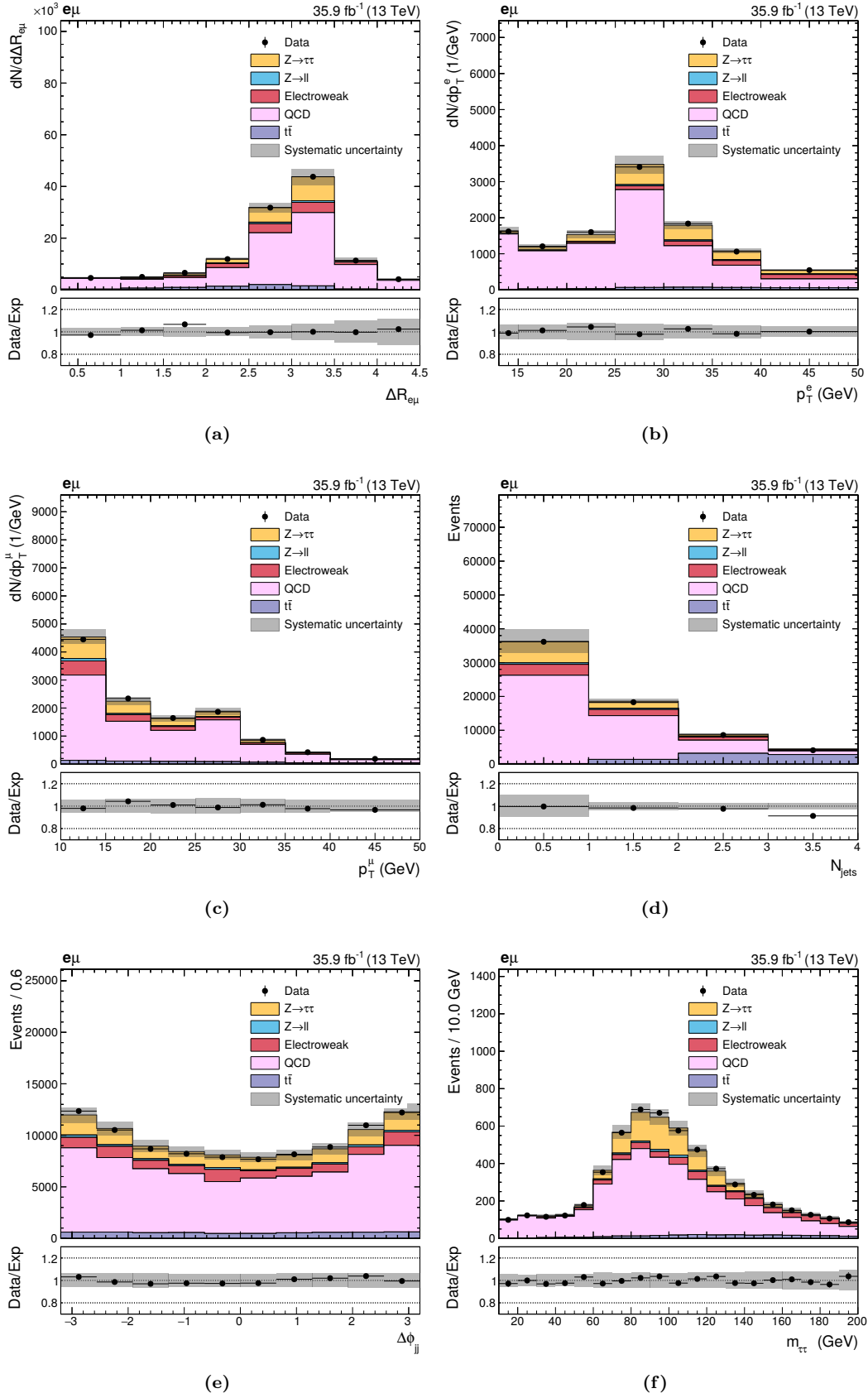


Figure 3.27: Comparisons of the QCD background predictions and data in the muon anti-isolated DR are shown for several key kinematic variables. The uncertainty band shown in shaded grey is formed by propagating the uncertainties on the linear fit parameters as systematic uncertainties, which are then added in quadrature with the statistical errors.

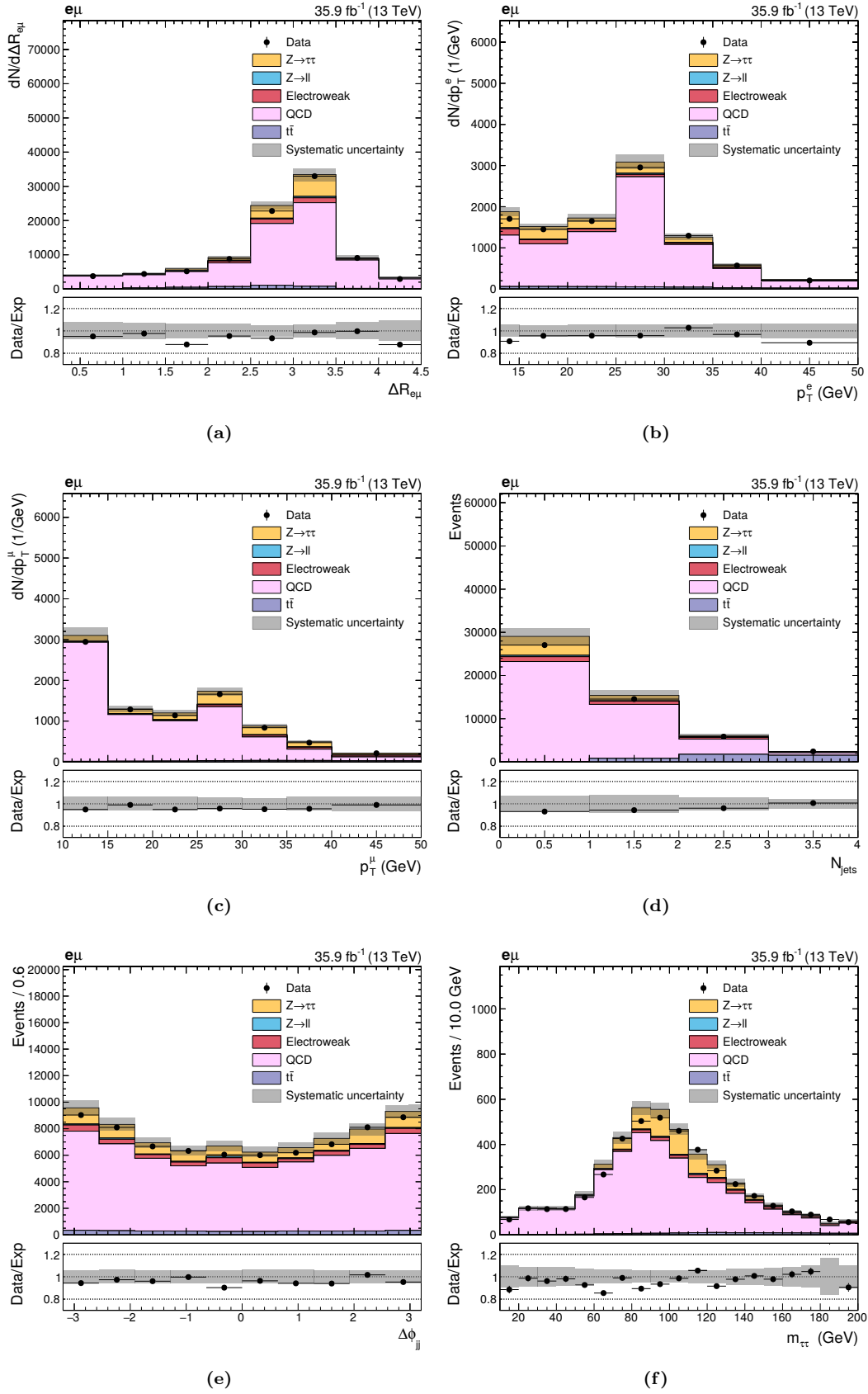


Figure 3.28: Comparisons of the QCD background predictions and data in the electron anti-isolated VR are shown for several key kinematic variables. The uncertainty band shown in shaded grey is formed by propagating the uncertainties on the linear fit parameters as systematic uncertainties, which are then added in quadrature with the statistical errors.

Chapter 4

Search for additional neutral Higgs bosons

This chapter describes the search for additional neutral Higgs bosons in the ditau final state. The search is performed using the dataset collected by CMS in 2016, which corresponds to a total integrated luminosity of 35.9 fb^{-1} . The analysis utilises the four most sensitive ditau final states: $\tau_e\tau_\mu$, $\tau_e\tau_h$, $\tau_\mu\tau_h$, and $\tau_h\tau_h$. The latter three channels, to which the author contributed, will be described in detail here. An overview of the analysis strategy for the $\tau_e\tau_\mu$ channel can be found in Ref. [1]. The final results presented will include all four channels in a combination.

The analysis searches for the h, H, and A bosons, which will be collectively referred to as " ϕ " throughout this chapter. The results are presented in the context of both model independent and model dependent interpretations. For the model independent interpretation, upper limits are set on the $\sigma \times \mathcal{B}^{\tau\tau}$ for neutral Higgs bosons produced by GF ($gg\phi$) and $b\bar{b}$ -associated ($b\bar{b}\phi$) production. The model dependent results are presented in the context of the MSSM, where two benchmark scenarios are considered: $m_h^{\text{mod}+}$ and hMSSM. Exclusion limits are set in the m_A - $\tan\beta$ plane for both scenarios.

4.1 Signal modelling

Signal samples for the $b\bar{b}\phi$ and $gg\phi$ production processes are simulated for a range of mass points between 90 GeV to 3.2 TeV. For the $b\bar{b}\phi$ process, the analysis is not sensitive to differences in pseudoscalar and scalar Higgs boson processes and therefore signal samples simulated for scalar Higgs production are also used to model the pseudoscalar Higgs process. For the $gg\phi$ process the p_T distribution for scalar and pseudoscalar Higgs production can differ. This is accounted for by the reweighting procedure described in the next section. Therefore, signal samples simulated for pseudoscalar Higgs production are also used to model the scalar Higgs processes. The generators used to simulate the two production processes are described in the remainder of this section.

4.1.1 $gg\phi$ signal modelling

The $gg\phi$ signal is simulated at LO accuracy using PYTHIA 8.212. These LO samples are known to poorly describe the ϕ p_T distribution. To improve the description, the LO p_T spectrum is reweighed to the NLO distribution simulated by POWHEG. However, an additional complication is that the p_T distribution depends on the mass of the Higgs boson and $\tan\beta$. This is a result of the m_A - $\tan\beta$ dependence of the coupling strengths to top and bottom quarks, which both contribute to the GF loop. The p_T distribution is harder for the top quark loop, and therefore the Higgs bosons become more boosted for smaller values of $\tan\beta$, as the top contribution becomes more dominant.

The $\tan\beta$ dependence is taken into account by dividing the full p_T spectrum into three components: top-only, bottom-only, and top+bottom interference. In the MSSM the loop may also have contributions from SUSY particles such as stops and sbottoms, but these contributions are typically small (less than a few percent) and are subsequently neglected. In the absence of contributions from SUSY particles the p_T distributions for each of the three components are equivalent to the distributions obtained in the 2HDM, and the p_T spectra are therefore derived in the 2HDM formulation. The top- and bottom-only distributions are generated assuming only top/bottom quarks in the loop. The interference distribution is obtained by generating a top+bottom p_T spectrum and subtracting the top- and bottom-only contributions. These distributions are generated for each considered mass point, and separately for the scalar and pseudoscalar Higgs bosons.

The reweighting procedure essentially modifies the p_T distributions to match the spectra obtained in the 2HDM formulation for top and bottom Yukawa couplings equal to the values predicted in the considered MSSM benchmark scenario. These couplings depend on $\tan\beta$ and the mixing angle α as predicted by the model. The total p_T distribution is formulated as:

$$\begin{aligned} \dot{\sigma} = & \left(\frac{y_t^{\text{MSSM}}}{y_t^{\text{2HDM}}} \right)^2 \dot{\sigma}_{\text{2HDM}}^t(Q_t) + \left(\frac{y_b^{\text{MSSM}}}{y_b^{\text{2HDM}}} \right)^2 \dot{\sigma}_{\text{2HDM}}^b(Q_b) \\ & + \left(\frac{y_t^{\text{MSSM}} y_b^{\text{MSSM}}}{y_t^{\text{2HDM}} y_b^{\text{2HDM}}} \right) (\dot{\sigma}_{\text{2HDM}}^{t+b}(Q_{tb}) - \dot{\sigma}_{\text{2HDM}}^t(Q_{tb}) - \dot{\sigma}_{\text{2HDM}}^b(Q_{tb})), \end{aligned} \quad (4.1)$$

where $y_{t/b}^i$ are the Yukawa coupling strengths in the MSSM and 2HDM, $\dot{\sigma}_{\text{2HDM}}^i$ are the differential distributions ($\dot{\sigma} = \frac{d\sigma}{dp_T}$), and Q_i are resummation scales which depend on the Higgs boson mass as discussed in Ref. [140–142]. The y_i^{2HDM} coupling strengths, given by $(y_t^{\text{2HDM}}, y_b^{\text{2HDM}}) = (\cot\beta, \tan\beta)$ for the pseudoscalar and $(y_t^{\text{2HDM}}, y_b^{\text{2HDM}}) = (\frac{\sin\alpha}{\sin\beta}, \frac{\cos\alpha}{\cos\beta})$ for the scalar, depend on the assumed values of α and $\tan\beta$ used in the generation, where values of $\alpha = \frac{\pi}{4}$ and $\tan\beta = 50(15)$ were chosen for the (pseudo)scalar distributions. The y_i^{MSSM} coupling strengths are computed using the **SusHi** program [143], which scales the three differential distributions as a function of $\tan\beta$ for the model dependent interpretation of the results. For the model independent results, the p_T distribution expected for a $\mathcal{CP}$ -even Higgs boson in the SM for a given mass is assumed.

Figure 4.1 shows examples of the reweighted p_T distribution for three different $\tan\beta$ values: $\tan\beta = 5$, $\tan\beta = 15$, and $\tan\beta = 30$. The unweighted LO distribution is shown also for comparison. There is a significant difference between the LO and NLO distributions, which motivates the use of this procedure. The reweighting parametrised as a function of p_T is found to give good agreement between the reweighted LO and the NLO samples, for all kinematic variables used for event selection and signal extraction.

4.1.2 $bb\phi$ signal modelling

The $bb\phi$ production process is simulated with **Madgraph5_aMC@NLO** at NLO accuracy [144] (using the four-flavour scheme). As illustrated in Figure 4.2, these NLO samples are observed to give a significantly improved description of the N_{jets} and $N_{b\text{-jets}}$ distributions compared to LO samples generated with **PYTHIA 8**. The **PYTHIA 8** samples were used for previous results by CMS such as in [145]. A good description of these variables is important as they are typically used to differentiate between the $bb\phi$ and $gg\phi$ production modes and reduce backgrounds such as $t\bar{t}$.

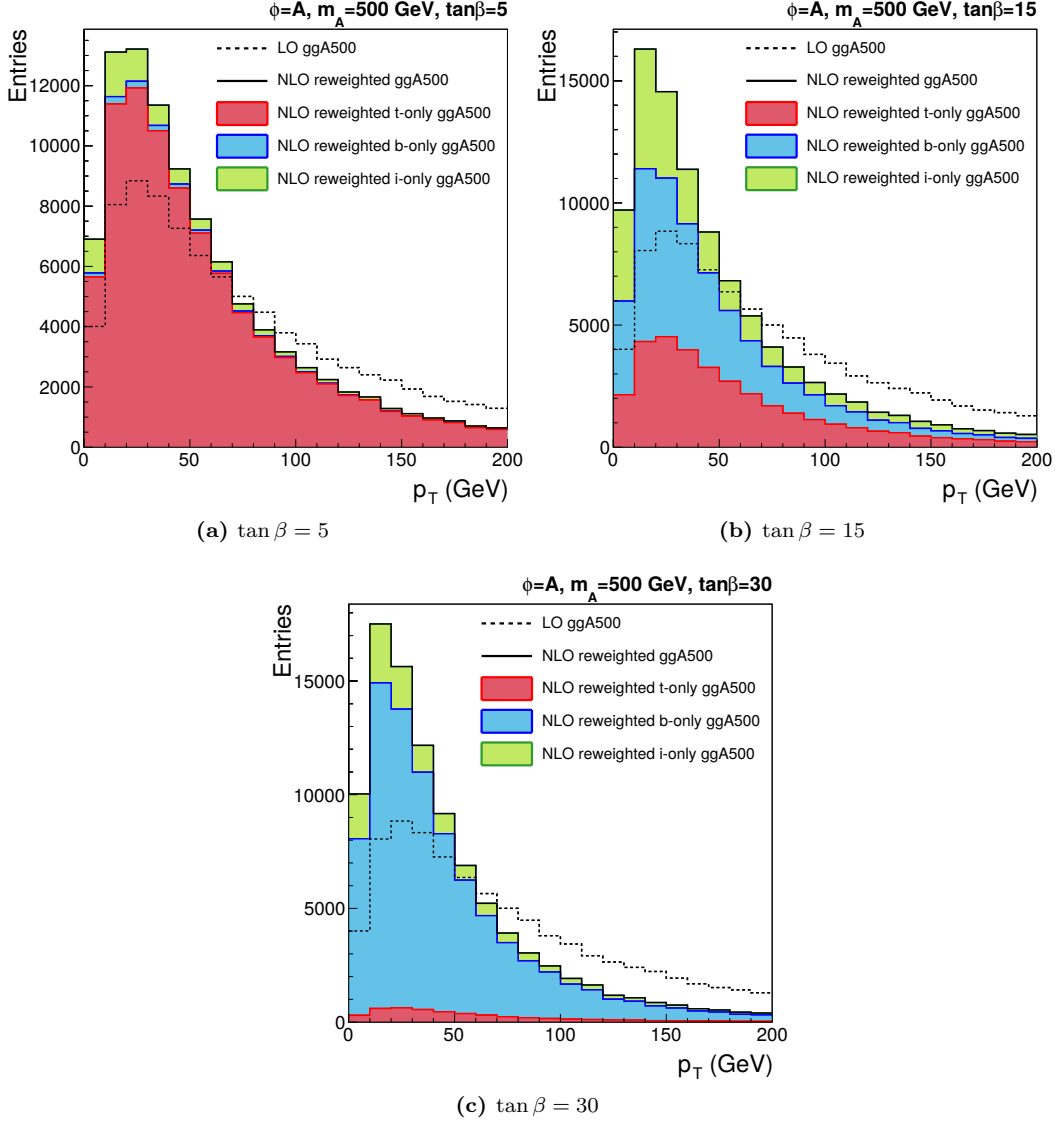


Figure 4.1: Examples of the reweighted p_T distribution for three different $\tan \beta$ values: $\tan \beta = 5$ (a), $\tan \beta = 15$ (b), and $\tan \beta = 30$ (c). The distributions are shown for a 500 GeV pseudoscalar Higgs boson. The reweighted NLO distribution is shown by the solid black line, where the top-only, bottom-only, and interference contributions are shown individually by the red, blue, and green histograms respectively, which are stacked. The unweighted LO distribution is shown by the dashed line for comparison.

4.1.2.1 Theoretical $bb\phi$ acceptance uncertainties

The theoretical uncertainties on the $bb\phi$ signal acceptance recommended in Ref. [30] include contributions from: the μ_R/μ_F scale variations, the shower scale Q_{sh} , and the PDFs. The latter contribution is small compared to the other two and is therefore neglected.

The μ_R/μ_F scale uncertainty is estimated by varying the μ_R and μ_F independently by factors of 2 and $\frac{1}{2}$ under the constraint $\frac{1}{4} < \frac{\mu_R}{\mu_F} < 4$, and taking the envelope of the variations as the uncertainty. For each μ_R/μ_F scale choice the cross section is normalised to the nominal cross section, so that the uncertainty only includes contributions to the signal acceptance and not the inclusive event yield. The size of the scale uncertainty is illustrated in Figure 4.3 for the $N_{b\text{-jets}}$ and N_{jets} distributions for a 700 GeV signal. The scale uncertainty is found to mainly affect these

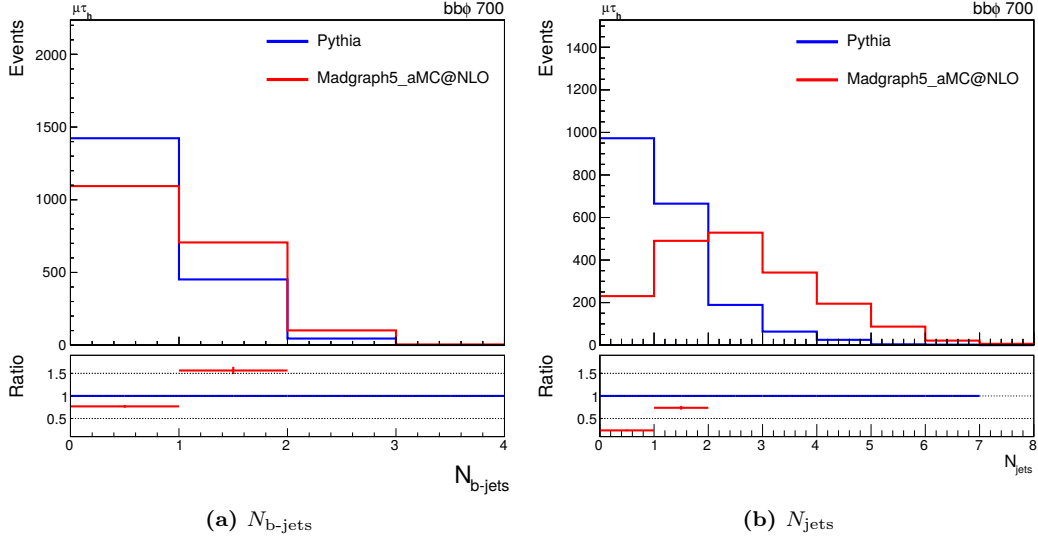


Figure 4.2: Comparison of the LO (blue) and NLO (red) $N_{b\text{-jets}}$ (left) N_{jets} (right) distributions for a 700 GeV Higgs boson.

distributions, and variations in other variables such as the ϕ p_T are typically much smaller.

The Q_{sh} shower scale determines the largest hardness accessible by the parton shower [30]. The uncertainty on this scale choice is determined by varying Q_{sh} by factors of $\sqrt{2}$ and $\frac{1}{\sqrt{2}}$. These variations are found to have the largest effects on the $N_{b\text{-jets}}$ and N_{jets} distributions, as illustrated in Figure 4.4 for a 700 GeV signal. Generation of samples with varied Q_{sh} scales for every considered mass point would be computationally expensive and therefore samples are generated for only a few representative mass points. A fit is then used to interpolate between these mass points as illustrated in Figure 4.5. The uncertainty shown in this case is the uncertainty on the fraction of events passing the $N_{b\text{-jets}} \geq 1$ selection used to define the categories described in Section 4.3.

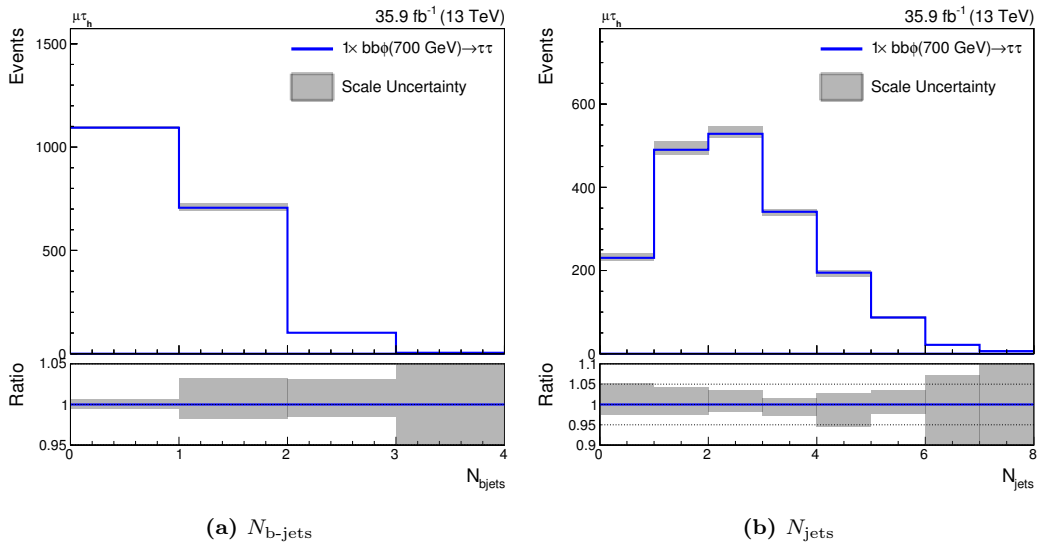


Figure 4.3: The μ_R/μ_F QCD scale uncertainty on the $N_{b\text{-jets}}$ (left) and N_{jets} (right) distributions for a 700 GeV Higgs boson.

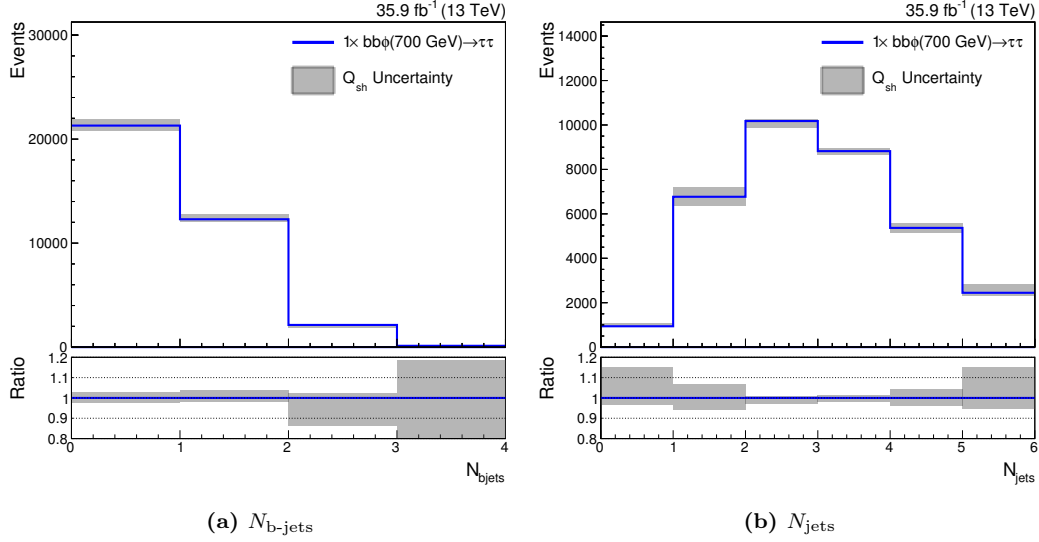


Figure 4.4: The Q_{sh} shower scale uncertainty on the $N_{b\text{-jets}}$ (left) and N_{jets} (right) distributions for a 700 GeV Higgs boson.

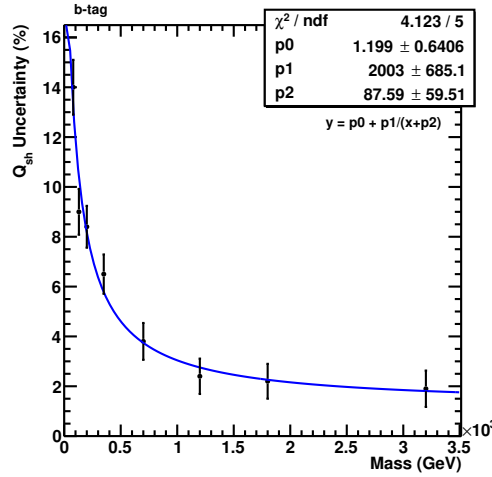


Figure 4.5: The Q_{sh} shower scale uncertainties on the fraction of events passing the $N_{b\text{-jets}} \geq 1$ selection as a function of mass. The fitted function shown by the blue line is used to extrapolate the uncertainty between mass points.

4.1.2.2 MSSM cross sections and branching ratios

The branching ratios of the Higgs bosons to τ leptons are calculated using **FeynHiggs** 2.10.2 [146–150] for the $m_h^{\text{mod}+}$ benchmark scenario, and using **HDECAY** 6.40 [151] for the hMSSM scenario. The GF cross sections are calculated using the **SusHi** v1.4.1 [143] program, which includes NLO QCD corrections in the context of the MSSM [152–157], NNLO QCD corrections for the top quark contribution to the quark loop [158–162], and electroweak corrections due to the light fermions [163, 164]. The $b\bar{b}$ -associated production cross sections are obtained by combining the four-flavour scheme NLO QCD calculations [165, 166], and the five-flavour scheme NNLO QCD calculations implemented in **BBH@NNLO** [167], using the Santander matching scheme [168].

4.2 Event selection

The selection of $\tau_l\tau_h$ and $\tau_h\tau_h$ pairs is described in this section. To reject leptons and hadronic tau candidates originating from PU interactions, the particles are required to be associated with the primary vertex by applying cuts on the transverse and/or longitudinal impact parameters, d_{xy} and d_z . Electrons and muons are required to have $d_{xy} < 0.045$ cm and $d_z < 0.2$ cm, and hadronic taus must have $d_z < 0.2$ cm.

4.2.1 Event selection in the $\tau_\mu\tau_h$ channel

The offline selection requires an oppositely charged muon and hadronic tau pair, which are separated by $\Delta R > 0.5$. The selected muon must pass the *medium* muon ID, have $I_{\text{rel}} < 0.15$, $|\eta| < 2.1$, and $p_T > 23$ GeV. The events are selected online using the single muon trigger, where the online muon is required to be matched to the offline muon with $\Delta R < 0.5$.

The hadronic tau must have $|\eta| < 2.3$ and $p_T > 30$ GeV. The τ_h candidate must also be reconstructed by the HPS algorithm, pass the *tight* MVA-isolation, the *very-loose* working point of anti-electron discriminant, and the *tight* working point of the anti-muon discriminant. The tau MVA-isolation working point is optimised by comparing the expected upper limits on the cross sections times branching ratios for the different working points, as illustrated in Figure 4.6 for the $gg\phi$ process. The very wide range of signal masses considered in the analysis makes the choice of working point non-trivial, as different working points are optimal for different signal mass points. The *tight* working point is chosen as it is observed to give the best sensitivity in both the low mass (≤ 500 GeV) and intermediate mass regions ($500 < m_\phi \leq 1$ TeV), where the latter is the more interesting mass region for MSSM scenarios.

Even after requiring the τ_h candidate to pass the *tight* MVA-isolation, there still remains a significant contribution to the background from W+jets events with a $\text{jet} \rightarrow \tau_h$ fake. As illustrated in Figure 3.23, the $m_T(l, E_T^{\text{miss}})$ variable gives good separation between W+jets events and other processes. The $m_T(\mu, E_T^{\text{miss}})$ is therefore required to be less than 70 GeV to suppress this background.

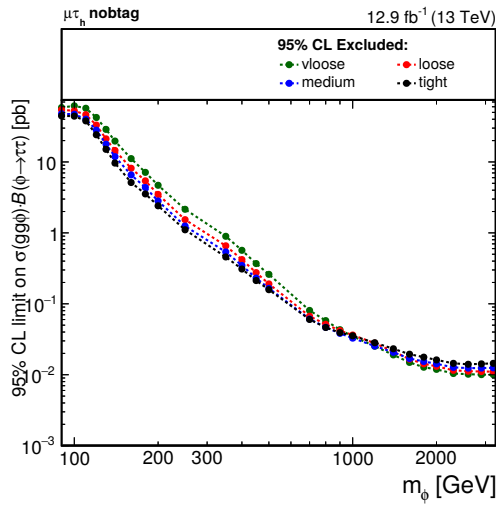


Figure 4.6: The expected limits on the $gg\phi$ cross section times branching ratio in the $\tau_\mu\tau_h$ channel, using different τ_h MVA-isolation working points. The limits are produced using 12.9fb^{-1} of data collected at 13 TeV.

4.2.2 Event selection in the $\tau_e\tau_h$ channel

The offline selection requires an oppositely charged electron and hadronic tau pair, which are separated by $\Delta R > 0.5$. The electron must pass the *tight* working-point of the electron ID, have $I_{\text{rel}} < 0.10$, $|\eta| < 2.1$, and $p_T > 25$ GeV. The events are selected online using the single-electron trigger, where the online electron is required to be matched to the offline electron with $\Delta R < 0.5$.

The hadronic tau must have $|\eta| < 2.3$ and $p_T > 30$ GeV. The τ_h candidate must also be reconstructed by the HPS algorithm, pass the *tight* MVA-isolation, the *tight* working point of anti-electron discriminant, and the *loose* working point of the anti-muon discriminant. As for the $\tau_\mu\tau_h$ channel, the *tight* working point of the MVA-isolation discriminant is chosen as it gives the best sensitivity for masses up to 1 TeV. The W+jets background is also reduced by requiring $m_T(e, E_T^{\text{miss}}) < 70$ GeV.

4.2.3 Event selection in the $\tau_h\tau_h$ channel

The offline selection requires two oppositely charged τ_h candidates with $p_T > 40$ GeV, which are separated by $\Delta R > 0.5$. Both τ_h candidates must pass the HPS algorithm, the *medium* MVA-isolation, the *very-loose* anti-electron discriminant, and the *loose* anti-muon discriminant. The *medium* MVA-isolation working-point is chosen as it is found to give the best expected limits for masses ≥ 400 GeV, which includes the most interesting 500 GeV–1 TeV region. The events are selected online using the double-tau triggers, where both offline τ_h candidates must be matched to the corresponding online τ_h objects with $\Delta R < 0.5$.

4.2.4 Lepton vetoes and pair selection

To ensure the individual channels are completely orthogonal to each other, events with additional electrons or muons are vetoed. This requirement also helps to reduce backgrounds from diboson and $Z/\gamma^* \rightarrow ll$ events which can produce multiple leptons. The p_T and I_{rel} selections applied to veto leptons are looser than for the signal candidates, requiring $p_T > 10$ GeV and $I_{\text{rel}} < 0.3$. For the $\tau_e\tau_h$ and $\tau_\mu\tau_h$ channels additional dilepton vetoes are applied to further reduce the $Z/\gamma^* \rightarrow ll$ background. For the $\tau_e\tau_h$ channel this veto removes any events with two oppositely charged electrons passing the *veto* electron ID with $p_T > 15$ GeV and $I_{\text{rel}} < 0.3$. For the $\tau_\mu\tau_h$ channel the dilepton veto removes any events with two oppositely charged global PF muons with $p_T > 15$ GeV and $I_{\text{rel}} < 0.3$.

The lepton vetoes ensure that there can be only one τ_l candidate for the $\tau_l\tau_h$ channels, as events with multiple candidates would anyway be removed by the vetoes. Multiple $\tau_l\tau_h$ pairs can be formed if more than one τ_h candidate exists. In these cases the pair with the most isolated τ_h is retained - meaning the τ_h candidate with the highest MVA-isolation score. For the $\tau_h\tau_h$ channel the candidate pairs are first sorted according to the leading τ_h MVA-isolation score. It is possible to have multiple pairs with the same leading τ_h candidate, which would rank equally after this initial sorting. In such cases the pairs are further sorted according to the MVA-isolation score of the sub-leading τ_h . The highest ranked pair is then retained.

4.3 Event categorisation

The events are split into two categories based on the number of b-tagged jets to target the $gg\phi$ and $bb\phi$ processes. The two categories are called the "no b-tag" and "b-tag" categories, respectively. The no b-tag category includes all events which contain exactly zero b-tagged jets with $p_T > 20$

GeV. With this selection $\approx 95\%$ of $gg\phi$ events are contained in the no b-tag category. Somewhat unexpectedly, depending on the boson mass approximately 60–80% of $bb\phi$ events are also included in this category. This is partly due to the $\approx 60\%$ b-tagging efficiency for heavy flavour jets. Often the b quarks also hadronise to form jets that are either too soft in p_T to pass the 20 GeV cut, or are emitted outside of the tracker coverage where b-tagging is not possible. This is illustrated in Figure 4.7 which shows the parton-level p_T and η distributions of the leading b quark produced in association with a 120 GeV ϕ boson. The b quarks produced with $p_T \lesssim 20$ GeV or $|\eta| \gtrsim 2.4$ will tend not to be reconstructed as b-jets. The inclusion of the b-tag category is used to target the remaining $bb\phi$ events, and this also allows the $gg\phi$ and $bb\phi$ processes to be differentiated.

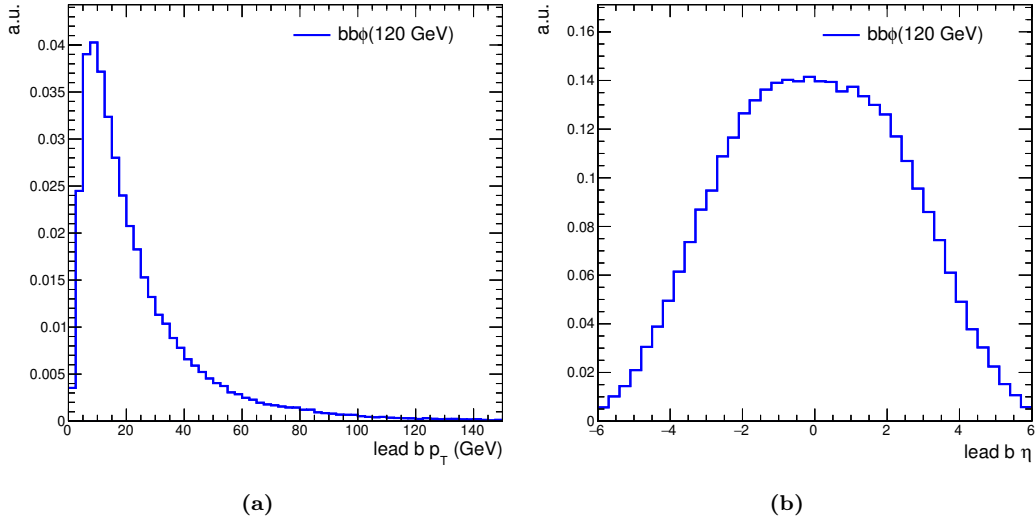


Figure 4.7: Parton-level p_T (a) and η (b) distributions for the leading b quarks produced in association with a 120 GeV ϕ .

The definition of the b-tag category used for previous CMS results such as in Ref. [145] required at least one b-tagged jet and no more than one jet (tagged or untagged) in the event. The latter selection was utilised to reduce the large $t\bar{t}$ background for events with $N_{\text{b-jets}} \geq 1$, which are typically accompanied by several jets. The LO signal samples used at the time suggested that most $bb\phi$ events have ≤ 1 jets, as shown in Figure 4.2. However, the NLO signal samples predict a significantly higher jet multiplicity and therefore this jet-veto may not be as powerful as previously expected.

The b-tag category selections were therefore re-optimised for the NLO signal samples. Three options for this categorisation were considered:

- Option 1: $N_{\text{b-jets}} \geq 1$ and $N_{\text{jets}} \leq 1$.
- Option 2: $N_{\text{b-jets}} \geq 1$.
- Option 3: $N_{\text{b-jets}} = 1$.

Option 1 uses the same selection used for previous CMS results, which is compared to two alternative selections where the jet-veto is removed. These two options require either exactly one or one or more b-jets. The former is an attempt at removing some of the $t\bar{t}$ events, which typically have higher b-jet multiplicities compared to the $bb\phi$ signal. The expected upper limits on the $bb\phi$ cross section times branching ratio as a function of mass are shown in Figure 4.8 using both LO and

NLO $bb\phi$ samples to model the signal. For the LO signal samples, option 1 gives the best expected limits - which is why these selections were used for previous CMS results. However, when using the NLO samples this option performs worse than the other two over most of the mass range. Options 2 and 3 give approximately the same performance, and thus the former is chosen for the analysis as it uses the loosest of the two selections.

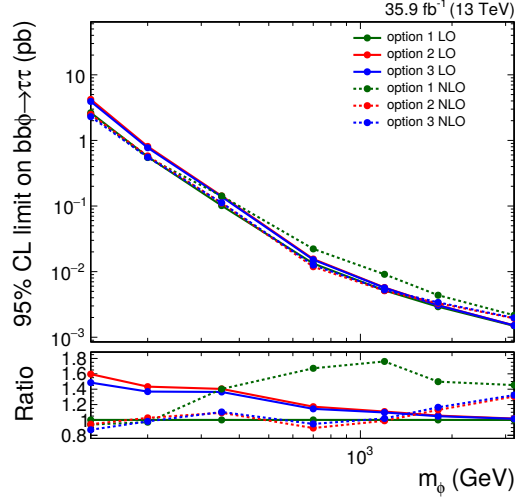


Figure 4.8: Comparison of the expected upper limits on the $bb\phi$ cross section times branching ratios for the three options considered for the b-tag category definition as described in the text. The limits obtained using the NLO signal samples are shown by the broken lines. The equivalent expected limits using the LO signal samples are shown by the solid lines for comparison.

4.3.1 m_T -based sub-categorisation in the $\tau_l\tau_h$ channels

The background in the $\tau_l\tau_h$ channels includes a significant contribution from W +jets events. As discussed in the previous section, the m_T variable can be used to reduce this background. The $m_T < 70$ GeV cuts described previously are quite loose compared to those applied for previous CMS results, which were typically 40–50 GeV. Rather than use a hard m_T cut of this magnitude, events are instead further divided into two sub-categories called the "loose- m_T " and "tight- m_T " categories. The tight- m_T subcategory requires $m_T < 40$ GeV, which selects most of the signal events. The loose- m_T subcategory, which contains events with $40 < m_T < 70$ GeV, is then included to increase the acceptance for signal events with larger m_T . Figure 4.9 shows the m_T distribution in the $\tau_\mu\tau_h$ channel indicating the divisions into the sub-categories.

The improvement in the sensitivity by inclusion of this categorisation is illustrated in Figure 4.10. In this figure the expected upper limits on the $bb\phi$ and $gg\phi$ cross sections times branching ratios are compared for a combination of the loose- m_T and tight- m_T categories, to the tight- m_T category alone. The latter is representative of a likely m_T cut that would be used without the sub-categorisation. The inclusion of the loose- m_T category is observed to improve the $gg\phi$ and $bb\phi$ expected limits by ≈ 2 –8% and 2–14% respectively.

A graphical representation of all the categories used in the analysis is shown in Figure 4.11. This shows the no b-tag and b-tag categories which are included for all channels, and the loose- m_T and tight- m_T subcategories for the $\tau_e\tau_h$ and $\tau_\mu\tau_h$ channels.

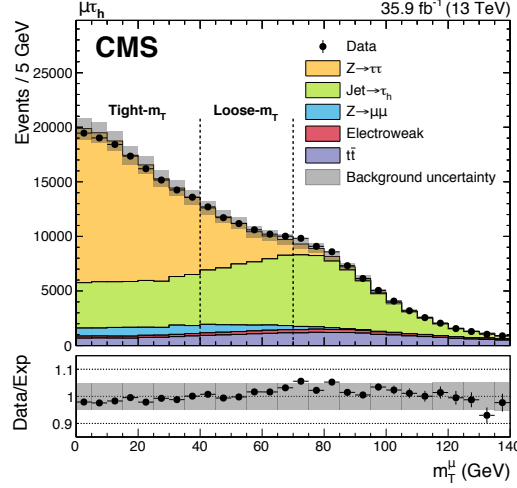


Figure 4.9: Pre-fit observed and expected distributions of the m_T variable in the $\tau_\mu\tau_h$ channel [1]. The distribution is shown prior to any categorisation into no b-tag and b-tag categories. The dashed lines and corresponding labels indicate the definitions of the loose- m_T and tight- m_T sub-categories. The uncertainty band shown in shaded grey includes systematic uncertainties affecting the normalisation of the distribution added in quadrature with the statistical uncertainties.

4.4 Background methods

The backgrounds that are considered in the analysis are: Drell-Yan $Z/\gamma^* + \text{jets}$, $W + \text{jets}$, QCD-multijet, $t\bar{t}$, diboson, and single top. Depending on the interpretation of the results the SM Higgs is treated as either background or part of the signal. The Higgs production processes that are considered are: GF, VBF, and W/Z -associated production. These processes are all simulated using POWHEG at NLO accuracy.

All background processes in which any τ_h candidate originates from a $\text{jet} \rightarrow \tau_h$ fake are modelled by the FF method, as described in Section 3.5.3. For the no b-tag category this includes mainly QCD and $W + \text{jets}$, whereas in the b-tag category there is a larger contribution from $t\bar{t}$. This background method is chosen for its good description of the $\text{jet} \rightarrow \tau_h$ fake rate and the large statistical power of the resulting background templates. The statistical power becomes particularly important for high mass signals where the expected number of background events is small and computationally expensive to model using MC. For such high mass scenarios the expected limits are very sensitive to the coarseness of the discriminating variable binning, and being able to afford finer binning often gives large improvements. All other background processes are modelled using MC.

4.4.1 Generator matching

It is necessary to split the MC processes depending on simulated objects from which the τ_l and τ_h candidates originate from. This ensures that backgrounds that are already covered by the FF method can be removed and are therefore not counted twice. It is also useful to separate Z/γ^* events further depending on whether the τ_h candidates originate from genuine tau lepton decays or from lepton fakes, as these events are subject to different systematic uncertainties. The resulting sub-processes can then be allowed to vary independently in the fit to data.

The reconstructed objects are matched to the generator level objects with $\Delta R < 0.2$. The

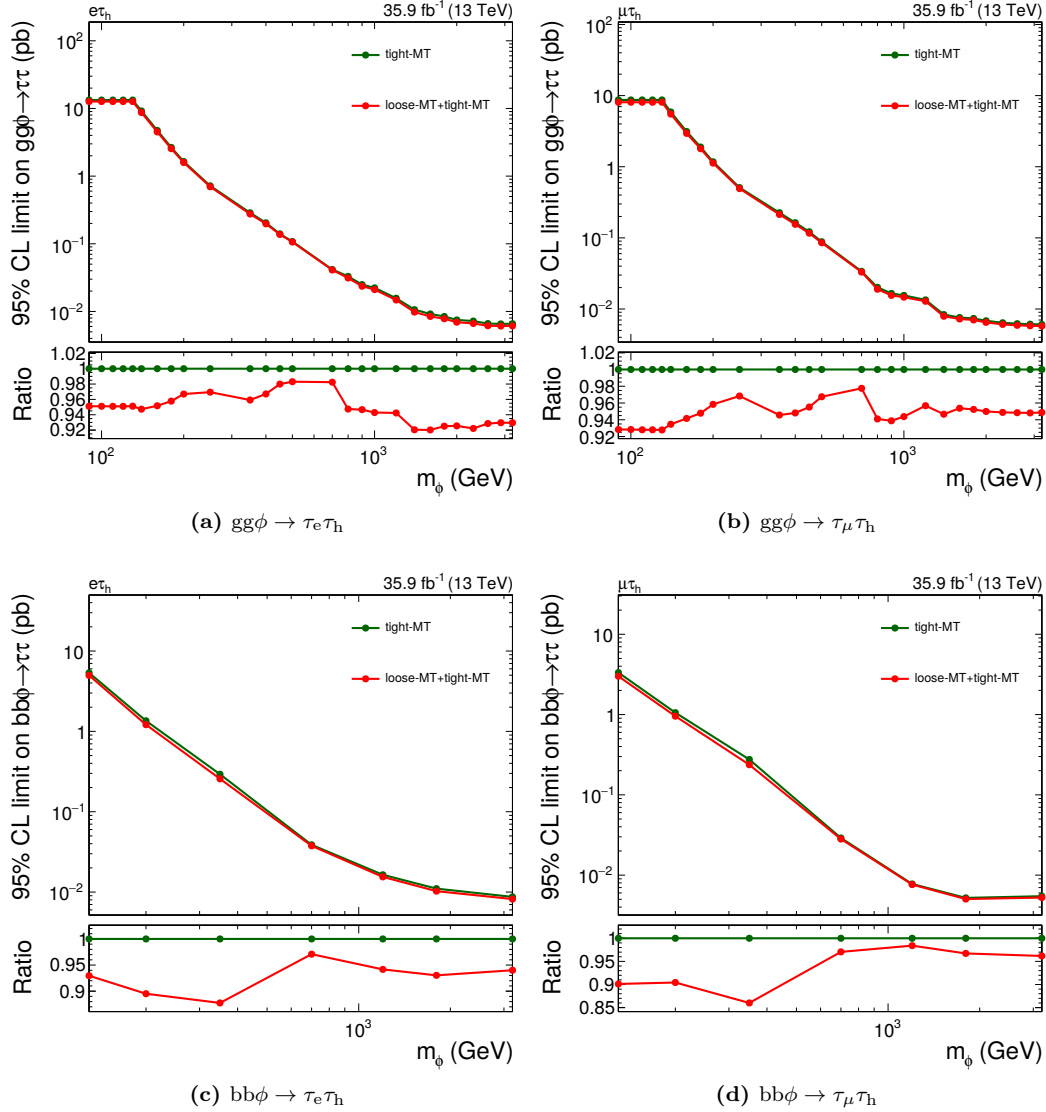


Figure 4.10: Comparisons of the expected upper limits including both the tight- m_T and loose- m_T categories in a combination (red) compared to the tight- m_T category only (green) are shown for the $\tau_e \tau_h$ (left) and $\tau_\mu \tau_h$ (right) channels. The no b-tag and b-tag categories have been combined in all cases. The upper and lower plots show the expected limits for the $gg\phi$ and $bb\phi$ processes respectively.

generator level objects that are considered are: prompt electrons and muons, electron and muons from tau decays, hadronic taus, and jets. A reconstructed object is classified as a jet if it is not matched to any of the other simulated objects. The simulated hadronic taus are formed by summing together the 4-vectors of the visible decay products.

4.4.2 Control regions

The largest backgrounds estimated from MC are $Z/\gamma^* \rightarrow \tau\tau$ and $t\bar{t}$. Additional control regions are included in the fit to data to constrain these backgrounds. For the $Z/\gamma^* \rightarrow \tau\tau$ process, two one-bin control regions are used to constrain the $Z/\gamma^* \rightarrow \tau\tau$ yield in the no b-tag and b-tag categories. These control regions are selected in the dimuon final state to achieve good purity of Z/γ^* events. The selections applied to the muons are chosen to closely mimic those used to select muons in the

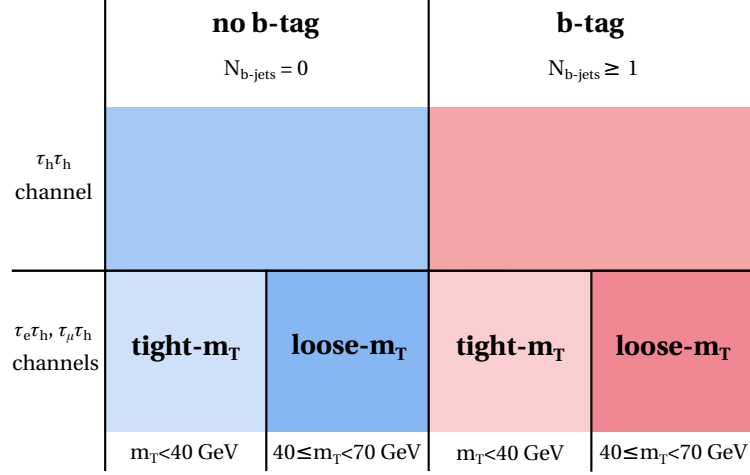


Figure 4.11: Diagram illustrating the categories used in the analysis. This includes no b-tag and b-tag categories for all channels. For the $\tau_l \tau_h$ channels, these categories are further divided into the loose- m_T and tight- m_T subcategories.

$\tau_\mu \tau_h$ channel. This means that the muon ID and isolation requirements are identical, and the events are triggered using the single muon trigger. The (sub)leading muon must also have $p_T > 23(10)$ GeV. The events are additionally required to have $70 < m_{\mu\mu} < 110$ GeV to increase the purity of $Z/\gamma^* \rightarrow \mu\mu$ events. Each of the two control regions apply the same respective requirements on the b-tagged jet multiplicities used to define the no b-tag and b-tag categories. Two unconstrained parameters are used to simultaneously scale the $Z/\gamma^* \rightarrow \tau\tau$ yields in the signal regions and the $Z/\gamma^* \rightarrow \mu\mu$ yields in the control regions - one of these parameters scales the yields for the no b-tag categories, and the other scales the yields in the b-tag categories. Figure 4.12 shows a comparison between the data and the expected background in both control regions.

A $t\bar{t}$ control region is also included to constrain both the shape and the yield of the $t\bar{t}$ process. This control region is selected in the $\tau_e \tau_\mu$ final state, which was not discussed in this document but is described in Ref. [1]. Additional cuts of $E_T^{\text{miss}} > 80$ GeV and $D_\zeta < -50$ GeV are used to increase the purity of $t\bar{t}$ events (the D_ζ variable will be defined in Equation 5.9). A single unconstrained parameter is used to simultaneously scale the yields of the $t\bar{t}$ templates in all control regions and signal regions. Rather than use just a one-bin control region, events are binned as a function of the discriminating variable to allow also the uncertainties affecting the $t\bar{t}$ shape to be constrained. This is done in particular to constrain the uncertainty relating to the top p_T reweighting described in the next section.

4.5 Systematic uncertainties

Two types of systematic uncertainties are considered in the analysis: normalisation uncertainties and shape uncertainties. The normalisation uncertainties affect only the yields of the signal and background templates, whereas shape uncertainties affect both the normalisations and shapes of the discriminating variable distributions. A summary of the systematic uncertainties is given in this section.

4.5.1 Normalisation uncertainties

The uncertainties that affect only the normalisation of the background and signal processes are:

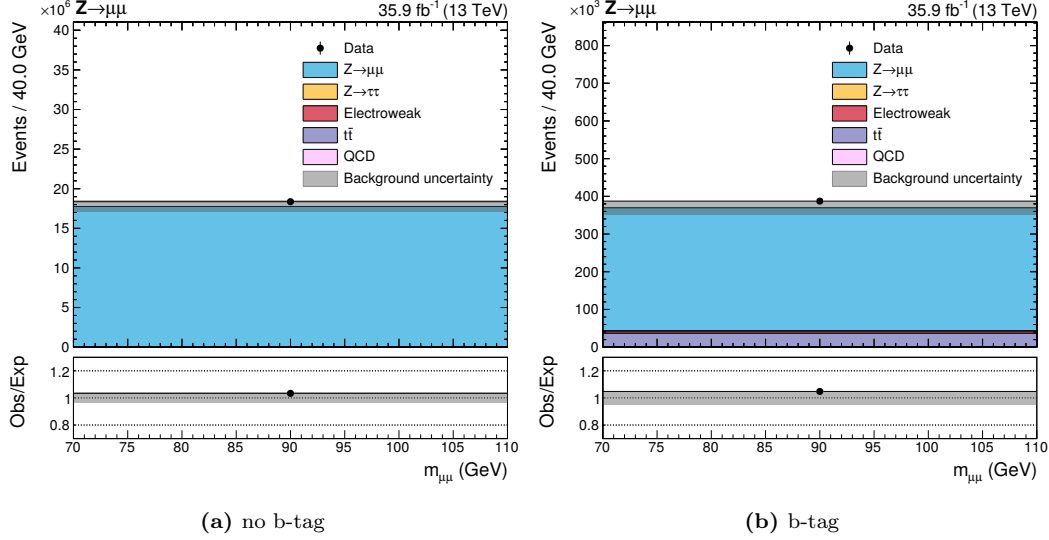


Figure 4.12: Data/background comparisons of the one-bin $Z/\gamma^* \rightarrow \mu\mu$ control regions included in the fit to constrain the $Z/\gamma^* \rightarrow \tau\tau$ yields in the signal regions. Two control regions are included to constrain the yields in the no b-tag (left) and b-tag (right) categories respectively. The uncertainty band shown in shaded grey includes the full systematic uncertainty model added in quadrature with the statistical uncertainties.

- **Luminosity**

The luminosity uncertainty amounts to 2.5% which is applied to all backgrounds and signal processes where the yield is estimated from simulation.

- **Electron, muon and tau ID efficiency**

The combined uncertainties on electron and muon identification, isolation, and trigger efficiencies are 2%. For hadronic taus the uncertainty on the tau identification efficiency is 5%. This is broken down into a part that is correlated between channels with a magnitude of 4%, and a part that is uncorrelated between channels with a magnitude of 3%. The uncorrelated part accounts for the different MVA-isolation and anti-muon/electron discriminator working-points used for each channel. Since two hadronic taus are required in the $\tau_h\tau_h$ channel the magnitudes of the correlated and uncorrelated uncertainties are 8% and 6% respectively. The uncorrelated part in this case is added in quadrature with the uncertainty on tau trigger efficiency which is 7%, giving a total uncorrelated uncertainty of 9.2%.

- **Jet energy scale uncertainty**

Uncertainties on the jet energy corrections are provided as a function of jet p_T and η . There exists 27 independent sources of jet energy scale uncertainties which are grouped together into a single uncertainty by adding the uncertainties for each source in quadrature. The uncertainties on the process yields due to the jet energy scale corrections are then determined by shifting the jet energies up and down by the p_T and η dependent uncertainty on the jet energy corrections. The resulting uncertainties have magnitudes that range between 1–8% depending on the process and the category.

- **B-tagging efficiency**

The uncertainties on the b-tagging efficiency corrections are given as a function of p_T and η . These are propagated by varying the b-tagging scale factors within their uncertainties. The resulting uncertainties on the process yields range between 0–4% depending on the process

and the category.

- **E_T^{miss} response and resolution efficiency**

For samples with recoil corrections applied the uncertainties on the E_T^{miss} scale and resolution are determined by varying the recoil correction parameters within their uncertainties. The uncertainties on the E_T^{miss} response range between 0–5%, and the uncertainties on the resolution range between 0–5%. For all other samples the E_T^{miss} uncertainties due to the clustered (i.e jet energy corrections) and unclustered energy corrections are used to determine the uncertainties. These range between 0–4% and 0–5%, respectively.

- **$l \rightarrow \tau_h$ fake-rate uncertainty**

The uncertainty on the $e \rightarrow \tau_h$ and $\mu \rightarrow \tau_h$ fake-rates is 11% and 12% respectively. These uncertainties are applied only to the $Z/\gamma^* \rightarrow ll$ process.

- **Signal/background normalisation uncertainties**

- A 4% uncertainty is applied to the $Z/\gamma^* \rightarrow ll$ process to account for the uncertainty on the Drell-Yan cross section. The yield for the $Z/\gamma^* \rightarrow \tau\tau$ process is estimated from a control region and therefore the cross section uncertainty is not needed. Instead a theoretical uncertainty on the extrapolation from the control regions to the signal regions is determined by varying the μ_R/μ_F scales. The size of these uncertainties ranges between 1–7% depending on the category.
- Theoretical uncertainty on the $t\bar{t}$ yields due to the extrapolation from the $t\bar{t}$ control region to the signal regions. This is estimated by varying the μ_R/μ_F scales and ranges between 0–2%.
- A 5% uncertainty is applied to the diboson and single-top processes to account for the uncertainties on the production cross sections.
- The uncertainties on the SM GF Higgs production cross section recommended in Ref. [30] are applied. This includes a 3.9% uncertainty from variations of the μ_R/μ_F scales and a 3.2% uncertainty from PDFs+ α_s variations.
- The uncertainties on the VBF Higgs production cross section recommended in Ref. [30] are applied. This includes a 0.4% uncertainty from variations of the μ_R/μ_F scales and a 2.1% uncertainty from PDFs+ α_s variations.
- The uncertainties on the W(Z)-associated Higgs production cross sections recommended in Ref. [30] are applied. This includes a 0.7% (3.8%) uncertainty from the μ_R/μ_F scale variations and a 1.9% (1.6%) uncertainty from PDFs+ α_s variations.
- The uncertainties on the $H \rightarrow \tau\tau$ branching ratio recommended in Ref. [30] are applied for all SM Higgs processes. This includes a 1.7% uncertainty due to missing higher order corrections, a 0.99% uncertainty on the quark masses, and a 0.62% uncertainty on α_s .
- The uncertainties on the MSSM Higgs boson cross sections are included for the model-dependent interpretation. This includes uncertainties due to the μ_R/μ_F scales and the choice of PDFs. The PDFs uncertainties are estimated following the prescription in [169, 170]. The uncertainties are computed for each m_A - $\tan\beta$ point and vary between 15–20%.

- **Theoretical uncertainties on the $bb\phi$ signal acceptance**

The μ_R/μ_F scale uncertainties and Q_{sh} uncertainties are estimated as described in Section 4.1.2.1 and added linearly as recommended in Ref. [30]. The size of the resulting uncertainties range between 3–17% depending on the signal mass and the category.

4.5.2 Shape uncertainties

The uncertainties that affect both the shapes and the yields of the background and signal processes are:

- **Hadronic tau energy scale**

Three uncorrelated uncertainties each that affect only one of the three tau decay modes used in the analysis are applied. Each uncertainty varies the hadronic tau energy scale up and down by 1.2%. Shifts in the energy scales are also propagated to the E_T^{miss} .

- **$e \rightarrow \tau_h$ energy scale**

Two uncorrelated uncertainties each that affect only one of the 1-prong tau decay modes are applied to the $Z/\gamma^* \rightarrow ee$ process in the $\tau_e \tau_h$ channel. These uncertainties vary the energy scale correction by 0.5% for the 1-prong decay mode, and by 1.0% for the 1-prong+1- π^0 decay mode. Shifts in the energy scales are also propagated to the E_T^{miss} .

- **High- p_T tau ID efficiency**

The measurement of the tau ID scale factor is performed using τ_h candidates with p_T up to ≈ 300 GeV and extrapolated to higher p_T values. Although these measurements are consistent with a constant scale factor there is some evidence within the statistical uncertainties for a linearly decreasing scale factor. A p_T dependent asymmetric uncertainty on the tau ID efficiency equal to $^{+5\%}_{-35\%} \times p_T/(1 \text{ TeV})$ is therefore applied to account for the extrapolation to the higher p_T regions.

- **Top p_T reweighting**

An uncertainty between no and twice the correction on the reweighting is applied to $t\bar{t}$ events.

- **Drell-Yan p_T /mass reweighting**

Uncertainties on the Drell-Yan weights include:

- The uncertainty on the muon-energy scale for the dimuon events used to measure the weights.
- An uncertainty due to the subtraction of the non- $Z/\gamma^* \rightarrow \mu\mu$ backgrounds when deriving the weights. The largest background is from $t\bar{t}$. The uncertainty is therefore estimated by varying the $t\bar{t}$ contribution up and down by the uncertainty on the $t\bar{t}$ cross section.
- Statistical uncertainty on the measured weights. All but three of the p_T /mass bins have negligible statistical uncertainties. The three that have large uncertainties affect mainly the high mass regions, and are included as three uncorrelated uncertainties.

- **Fake-factor uncertainties**

Uncertainties on the shapes of the $j \rightarrow \tau_h$ templates come from several sources:

- Statistical uncertainty on the FF measurements, given by the fit uncertainty to the p_T dependence.
- Systematic uncertainties related to the non-closure correction, as a function of the visible mass.
- Systematic uncertainties on the DR→AR extrapolation corrections.
- Systematic uncertainties on the fractions of W +jets and $t\bar{t}$ events with one jet $\rightarrow \tau_h$ in the $\tau_h \tau_h$ channel. This is necessary as the fake-factors for these processes must be scaled by a factor of 2 for this channel as described in Section 3.5.3. For the $\tau_l \tau_h$ channels the uncertainties on the fractions are negligible compared to the other systematic uncertainties.
- An uncertainty due to the subtraction of the genuine τ_h contribution using simulation.

- **Finite statistics in background templates**

Limitations in the statistical sizes of the signal and background templates are accounted for by the Barlow-Beeston method [131, 132], in which the bins of each shape template are allowed to float independently of each other within their respective statistical uncertainty. To reduce the computing time, the statistical uncertainties for a given bin are merged into a subset of the templates. The subset is chosen such that the sum of the squared error of the merged uncertainties is at most 40% of the total sum of the squared errors. During the merging, the templates with the smallest statistical uncertainties are considered first, such that the subset contains as many templates as possible.

4.6 Signal extraction and statistical interpretation

The discriminating variable used for signal extraction is the total transverse mass [171], which is defined as

$$m_T^{\text{tot}} = \sqrt{m_T^2(\vec{E}_T^{\text{miss}}, \tau_1^{\text{vis}}) + m_T^2(\vec{E}_T^{\text{miss}}, \tau_2^{\text{vis}}) + m_T^2(\tau_1^{\text{vis}}, \tau_2^{\text{vis}})}, \quad (4.2)$$

where τ_i^{vis} denote the visible decay products of the tau leptons. As illustrated in Figure 4.13, this variable gives good separation between the signal and the QCD background, which is the dominant background for large signal masses.

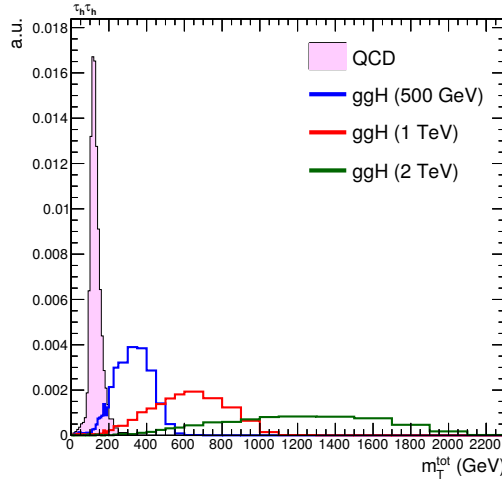


Figure 4.13: Comparison of the m_T^{tot} distributions for the QCD background and three $gg\phi$ signal mass points in the $\tau_h\tau_h$ final state.

Results are extracted by means of a simultaneous binned maximum likelihood fit including all channels, categories, and control regions in the fit. The likelihood function is constructed as

$$\mathcal{L}(\text{data} | \mu, \theta) = \prod_i \text{Poisson}(n_i | \mu \cdot s_i(\theta) + b_i(\theta)) \cdot p(\tilde{\theta} | \theta), \quad (4.3)$$

where the product operator runs over all bins i of the discriminating variables, for all channels and categories. The observed number of data events in each bin is n_i . The signal and background expectations are given by s_i and b_i respectively, which are functions of θ that represent the full set of nuisance parameters corresponding to the systematic uncertainties described in Section 4.5. The signal expectation is scaled by a rate parameter μ , which can represent different quantities depending on how the signal expectation is normalised. In this case the signal is normalised to $\sigma \times \mathcal{B}^{\tau\tau} = 1$ pb, which means μ represents the $\sigma \times \mathcal{B}^{\tau\tau}$ of the signal in units of pb. The Poisson

probabilities are given by

$$\text{Poisson}(n_i | \mu \cdot s_i(\theta) + b_i(\theta)) = \frac{(\mu \cdot s_i(\theta) + b_i(\theta))^{n_i}}{n_i!} e^{-(\mu \cdot s_i(\theta) + b_i(\theta))}. \quad (4.4)$$

The $p(\tilde{\theta}|\theta)$ term represents the full set of probability density functions (pdfs) of the uncertainties on the nominal values of the nuisance parameters $\tilde{\theta}$. Nuisance parameters that affect only the normalisations of the templates are assigned "log-normal" pdfs, which ensures that the template yields can take only positive values. The functional form of the log-normal pdf is

$$f(x) = \frac{1}{x} \frac{1}{\sqrt{2\pi}\sigma^2} \exp\left(-\frac{(\ln x - \mu)^2}{2\sigma^2}\right), \quad (4.5)$$

where μ and σ are the mean and standard deviation of the normally distributed logarithm $\ln x$ of the variable x . Parameters that affect the shapes are assigned Gaussian pdfs. For each of these uncertainties two additional templates are produced corresponding to $\pm 1\sigma$ shifts. A vertical morphing algorithm [132] is used to smoothly interpolate between the $\pm 1\sigma$ shifts and extrapolate beyond them.

4.6.1 Setting of model independent upper limits on $\sigma \times \mathcal{B}^{\tau\tau}$

In the absence of a signal, upper limits on the $\sigma \times \mathcal{B}^{\tau\tau}$ are set using the modified frequentist approach [172, 173]. The test statistic is defined as

$$q_\mu = -2 \ln \left(\frac{\mathcal{L}(\text{data} | \mu, \hat{\theta}_\mu)}{\mathcal{L}(\text{data} | \hat{\mu}, \hat{\theta}_{\hat{\mu}})} \right), \quad 0 \leq \hat{\mu} \leq \mu, \quad (4.6)$$

where $\hat{\mu}$ and $\hat{\theta}_{\hat{\mu}}$ indicate the best-fit values of μ and θ after the fit to data, and $\hat{\theta}_\mu$ is the value of θ that maximises the likelihood for the value of μ being tested. The $0 \leq \hat{\mu}$ boundary prevents unphysical negative signal strengths. The $\hat{\mu} \leq \mu$ constraint ensures that a one-sided confidence interval is constructed, which prevents the exclusion of μ values less than the best-fit $\hat{\mu}$. The probability of finding q_μ at least as large as the observed value q_μ^{obs} is

$$\text{CL}_{\text{s+b}}(\mu) = \int_{q_\mu^{\text{obs}}}^{\infty} f(q_\mu | \mu, \hat{\theta}_\mu) dq_\mu, \quad (4.7)$$

where $f(q_\mu | \mu, \hat{\theta}_\mu)$ is probability density function of q_μ . The exclusion limits are set using the CL_s value defined as

$$\text{CL}_\text{s} = \frac{\text{CL}_{\text{s+b}}}{\text{CL}_\text{b}}, \quad (4.8)$$

where

$$\text{CL}_\text{b} = \text{CL}_{\text{s+b}}(\mu = 0). \quad (4.9)$$

The limits are set at the 95% CL, meaning tested values of μ are said to be excluded if $\text{CL}_\text{s}(\mu) < 0.05$.

The $f(q_\mu | \mu, \hat{\theta}_\mu)$ functions may be determined using toy MC datasets, which are generated according to the expectations with all nuisance parameters fixed to their best-fit values. The effect of systematic uncertainties is accounted for by sampling a set of pseudo-measurements $\tilde{\theta}$ using the chosen pdfs of the nuisance parameters. To remove the need for this computationally expensive generation of toy datasets, the limits are calculated in the asymptotic approximation [174]. In this approximation $f(q_\mu)$ follows a known formula, and the median expected limits and uncertainty contours are estimated using properties of the "Asimov dataset": a dataset constructed by setting the observation in each bin to the expectation.

Exclusion limits are set for both the $gg\phi$ and $bb\phi$ production modes. When setting the limits for one of these processes no assumptions should be made on the presence or absence of the other. Therefore, when setting limits for the $gg\phi$ process the normalisation of the $bb\phi$ processes is allowed to float freely in the fit, and vice versa when setting limits for $bb\phi$.

4.6.2 2D likelihood scans of $gg\phi$ vs $bb\phi$

As the analysis targets two signal processes, the likelihood function can be constructed in terms of both $gg\phi$ and $bb\phi$ signals,

$$\mathcal{L}(\text{data} | \mu_{gg\phi}, \mu_{bb\phi}, \theta) = \prod_i \text{Poisson} \left(n_i | \mu_{gg\phi} \cdot s_i^{gg\phi}(\theta) + \mu_{bb\phi} \cdot s_i^{bb\phi}(\theta) + b_i(\theta) \right) \cdot p(\tilde{\theta} | \theta), \quad (4.10)$$

where $s_i^{gg\phi}$ and $s_i^{bb\phi}$ are the signal expectations for the $gg\phi$ and $bb\phi$ processes, which are scaled by the $\mu_{gg\phi}$ and $\mu_{bb\phi}$ rate parameters that represent the $\sigma \times \mathcal{B}^{\tau\tau}$ of the two processes. A two dimensional likelihood scan is performed using $\mu_{gg\phi}$ and $\mu_{bb\phi}$ as the parameters of interest. The parameters are bounded such that they may take only positive values. The confidence regions are determined from the two dimensional negative log-likelihood defined as

$$-2\Delta \ln \mathcal{L}_{2D} = -2 \cdot (\ln \mathcal{L}(x, y) - \ln \mathcal{L}(\hat{x}, \hat{y})), \quad (4.11)$$

where x and y are the parameters of interest with best fit values $\hat{x}$ and $\hat{y}$. The 68% and 95% confidence regions are given by $-2\Delta \ln \mathcal{L}_{2D} = 2.28$ and 5.99 respectively.

4.6.3 Model dependent limit setting

The results are also interpreted in terms of the two MSSM benchmark scenarios described in Chapter 1. The simulated templates for the single Higgs production processes are combined into a multi-resonance signal model including the h , H , and A bosons. The single Higgs templates are scaled to the model predictions of the production cross sections and branching ratios, that depend on the m_A and $\tan\beta$ point being tested. The masses of the three bosons are also set to those predicted in the model as a function of m_A and $\tan\beta$, where a horizontal morphing algorithm [133] is used to interpolate between the simulated mass points.

The null hypothesis must consider the observed Higgs boson with mass equal to 125 GeV, and the tested models must be able to incorporate a light h with properties similar to this particle. The exclusion limits are therefore set by means of a SM vs MSSM hypothesis test. This ensures a meaningful and well defined problem even when the analysis becomes sensitive to the 125 GeV Higgs boson. The likelihood function for this test is constructed as

$$\mathcal{L}(\text{data} | \mu, \theta) = \mathcal{L}(\text{data} | \mu \cdot s_{\text{MSSM}}(\theta) + (1 - \mu) \cdot s_{\text{SM}}(\theta) + b(\theta)), \quad (4.12)$$

where $\mu = 0, 1$. A value of $\mu = 1$ corresponds to the MSSM hypothesis, whereas $\mu = 0$ corresponds to the SM hypothesis. The test statistic in this case is given by

$$q_\mu = -2 \ln \left(\frac{\mathcal{L}(\text{data} | \mu = 1, \hat{\theta}_{\mu=1})}{\mathcal{L}(\text{data} | \mu = 0, \hat{\theta}_{\mu=0})} \right). \quad (4.13)$$

The asymptotic approximation used to set model independent limits cannot be applied to this test statistic, and therefore the probability density functions are determined using toy MC datasets. These toy datasets are generated for a series of m_A - $\tan\beta$ points. For each point the value of $\text{CL}_s(\mu = 1)$ is computed. The points with $\text{CL}_s < 0.05$ are excluded at the 95% CL.

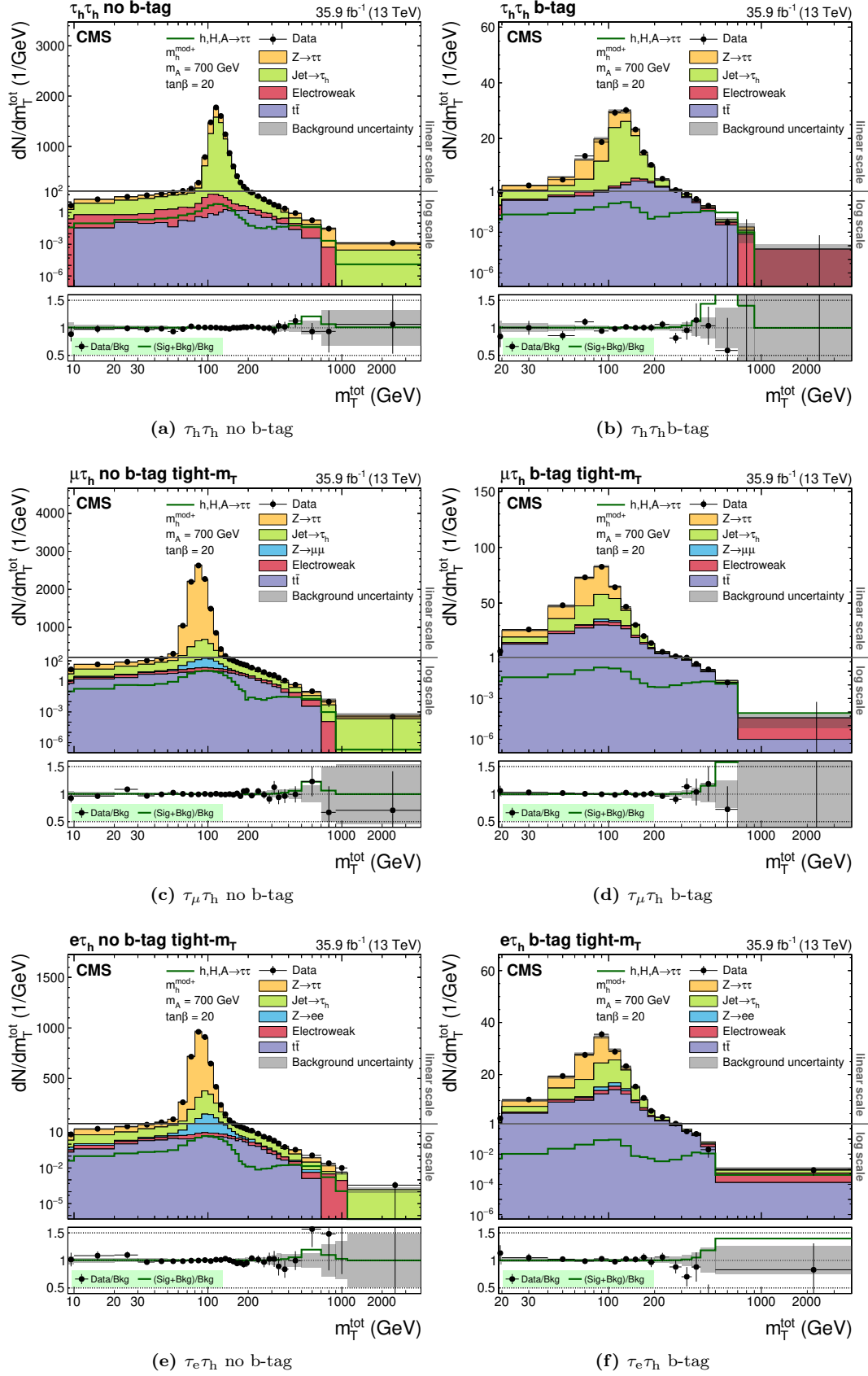


Figure 4.14: Observed and expected distributions of the m_T^{tot} variable in the no b-tag (left) and b-tag (right) categories. The distributions are shown for the $\tau_h \tau_h$ channel in the upper row, the $\tau_\mu \tau_h$ channel on the centre row, and the $\tau_e \tau_h$ channel in the lower row. For the $\tau_l \tau_h$ channels only the most sensitive tight- m_T subcategories are shown. The uncertainty band shown in shaded grey includes both systematic and statistical uncertainties. The expected signal distributions for the three neural Higgs bosons in the $m_h^{\text{mod}+}$ scenario with $m_A = 700$ GeV and $\tan \beta = 20$ are shown by the solid green line [1].

4.7 Results

Before performing the final fit to data, the validity of the model is checked by performing several goodness of fit tests. Tests are chosen that are sensitive to both local deviations from the model in individual bins, and correlated deviations across several bins. All systematic uncertainties and correlations are taken into account when performing these tests. These checks all indicate good compatibility between the data and the model.

The m_T^{tot} distributions after the simultaneous fit to data are shown in Figure 4.14 for the no b-tag and b-tag categories in the $\tau_e\tau_h$, $\tau_\mu\tau_h$, and $\tau_h\tau_h$ channels. For the $\tau_l\tau_h$ channels, only the most sensitive tight- m_T subcategories are shown. The expected background distributions are stacked. For the $\tau_l\tau_h$ channels, the histogram labelled "Electroweak" includes the diboson and single top backgrounds. For the $\tau_h\tau_h$ channel, this histogram also includes the $Z/\gamma^* \rightarrow ll$ process. The signal expected for the $m_h^{\text{mod+}}$ MSSM scenario with $m_A = 700$ GeV and $\tan\beta = 20$ is overlaid by the solid green line. This signal distribution has a twin peak structure, where the peak at $m_T^{\text{tot}} \approx 100$ GeV is due to the 125 GeV Higgs boson, assumed to be the h boson in this scenario. The peak at $m_T^{\text{tot}} \approx 400$ GeV is due to the H and A bosons, which are almost degenerate in mass. For all of the m_T^{tot} distributions the data are observed to agree well with the background-only expectation, and no evidence of a statistically significant excess is observed.

4.7.1 Model independent results

For all model independent results, the SM Higgs boson production by GF, VBF and in association with a vector boson are treated as background processes, with the corresponding cross sections and branching ratios scaled to the SM predictions. The 95% CL upper limits on the $\sigma \times \mathcal{B}^{\tau\tau}$ are shown in Figure 4.15 for both the $gg\phi$ and $bb\phi$ production processes. All channels and categories are combined, including the $\tau_e\tau_\mu$ channel not described in this document.

For the $gg\phi$ process, the expected limits in the low mass region are also shown for two cases where the loop is assumed to be dominated by the bottom quark ("b quark only") and the top quark ("t quark only") contributions. The mass region $\lesssim 110$ GeV is found to be sensitive to the assumptions made on the loop structure. This is due to the p_T distribution for the b quark only case being softer than for the t quark only, which affects the signal acceptance as the softer tau decay products are less likely to pass the online/offline p_T requirements. For higher mass signals, the tau decay products are sufficiently boosted enough to pass these p_T cuts by way of the decaying boson's large mass-energy conversion, and thus are not affected by the assumption made about the loop structure. The t quark only limits are observed to be very similar to the nominal limits. This is because the latter are produced assuming the p_T distribution of a SM Higgs boson, where the loop is dominated by the t quark contribution.

Figure 4.16 shows comparisons of the combined expected limits with the expected limits for the individual channels. For the $gg\phi$ process, the $\tau_\mu\tau_h$ channel is the most sensitive for masses $\lesssim 300$ GeV, where the $\tau_h\tau_h$ channel then becomes more sensitive for masses between 300–700 GeV. For larger masses the $\tau_\mu\tau_h$ channel has either comparable sensitivity or better than the $\tau_h\tau_h$ channel. For the $bb\phi$ process, the $\tau_\mu\tau_h$ channel is the most sensitive for masses $\lesssim 160$ GeV, and $\tau_h\tau_h$ is more sensitive for larger masses. The $\tau_e\tau_\mu$ channel is the least sensitive channel over most of the mass range, with the exception of the regions $\lesssim 140$ (110) GeV and $\gtrsim 2.9$ (2.6) TeV where it is more sensitive than the $\tau_e\tau_h$ channel to the $gg\phi$ ($bb\phi$) production mode.

In Figure 4.17 the expected limits are also compared to the limits set by CMS in a previous

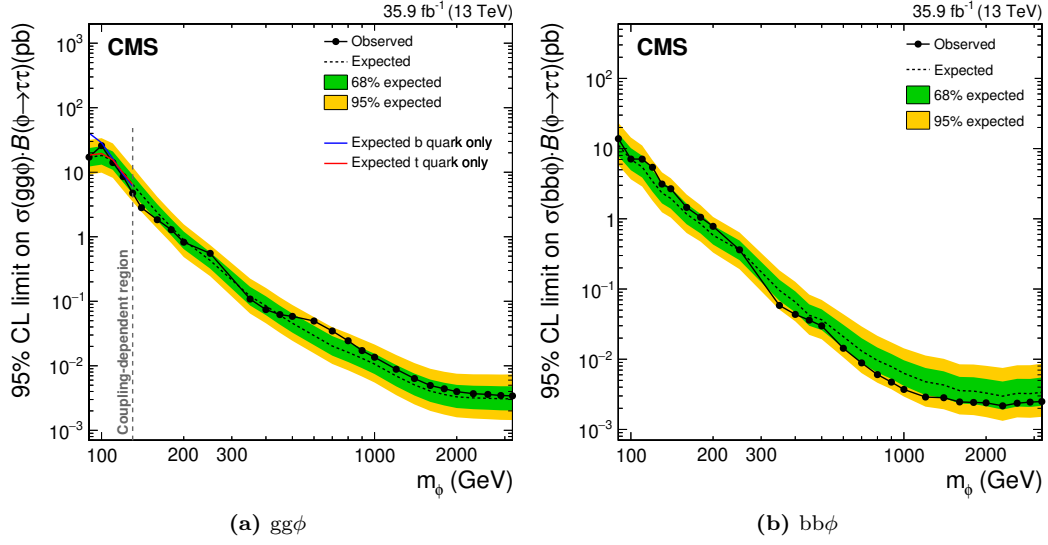


Figure 4.15: Observed and expected 95% CL upper limits on the $gg\phi$ (left) and $bb\phi$ (right) $\sigma \times \mathcal{B}^{\tau\tau}$. The observed limits are shown by the solid black lines and the expected limits are shown by the dashed lines. The 68% and 95% confidence intervals on the expected limits are shown by the green and yellow bands. The left inset also shows the $gg\phi$ expected limits when the quark loop is assumed to be dominated by the top quark (red line) and bottom quark (blue line) contributions [1].

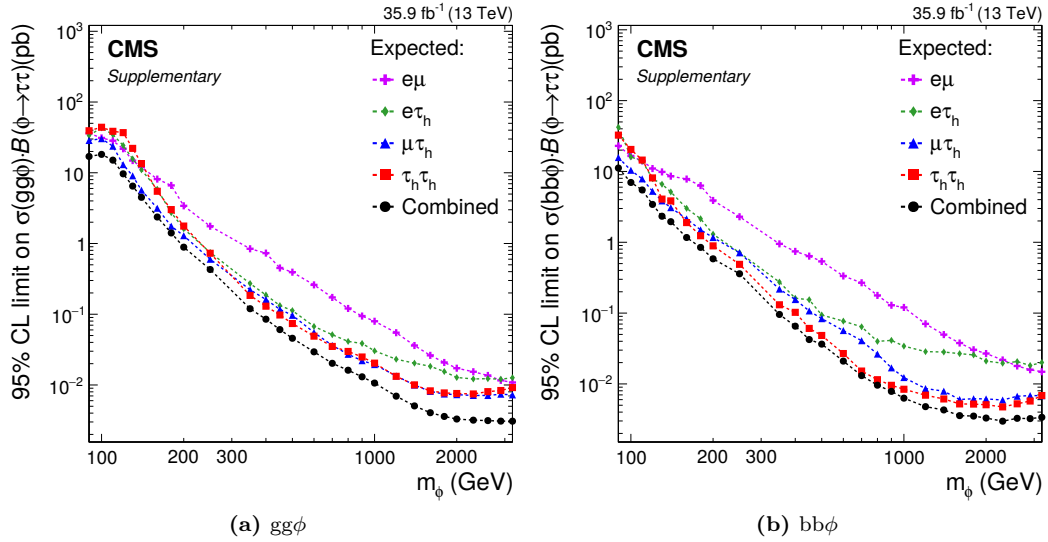


Figure 4.16: The expected 95% CL upper limits on the $gg\phi$ (left) and $bb\phi$ (right) $\sigma \times \mathcal{B}^{\tau\tau}$, comparing the combined limits (black) with the limits for the $\tau_e\tau_\mu$ (magenta), $\tau_e\tau_h$ (green), $\tau_\mu\tau_h$ (blue), and $\tau_h\tau_h$ (red) channels [1].

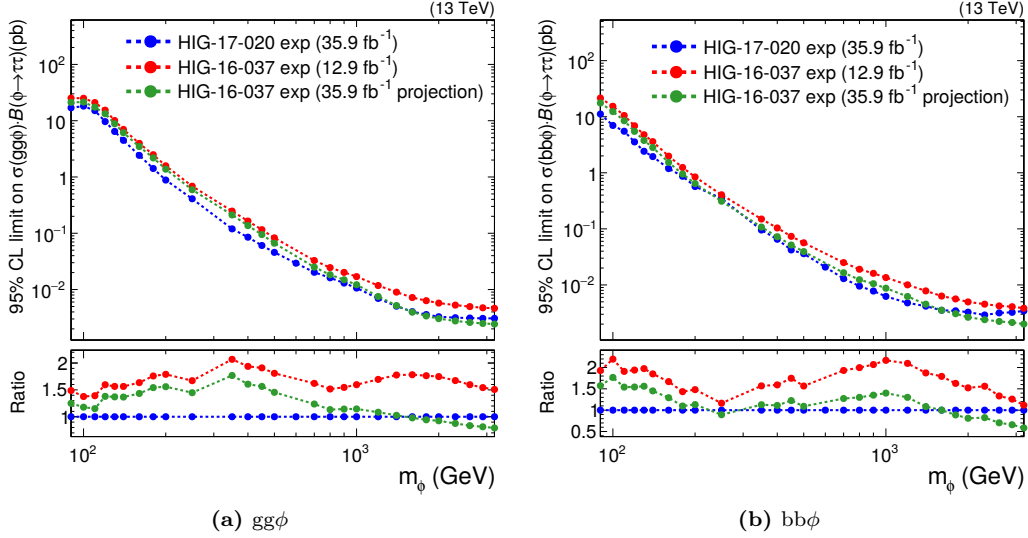


Figure 4.17: Comparison of the expected limits using the full 2016 dataset (blue), the expected limits set in the previous result by CMS using 12.9 fb^{-1} of data collected in 2016 [145] (red), and the projection of the results of the 12.9 fb^{-1} analysis to an integrated luminosity of 35.9 fb^{-1} (green).

analysis using 12.9 fb^{-1} of data collected in the first part of 2016 data-taking [145]. There is a significant improvement in the expected limits by up to a factor of 2 for some mass points. The limits are also compared to the results of the 12.9 fb^{-1} analysis projected to an integrated luminosity 35.9 fb^{-1} . For mass points up to $\sim 1.6 \text{ TeV}$ the current limits are more stringent than the projection. This demonstrates that the re-optimisation of the analysis along with the improvements made to the background modelling (which will be discussed in more detail in Section 4.8) have improved the analysis sensitivity. The expected limits for the mass points above 1.6 TeV are observed to worsen compared to the projection. This is mainly due to the uncertainty on the high- p_T tau ID efficiency described in Section 4.5. This uncertainty was increased prior to the current analysis, when additional measurements of the tau ID efficiency in the high p_T region suggested a more prominent linear trend in the ID scale-factors.

The two dimensional likelihood scans of the $gg\phi$ and $bb\phi$ production rates are shown in Figure 4.18, for a representative subset of the mass points considered. The best-fit values are shown by the black crosses, and the purple bands show the 68% and 95% confidence regions. The measured best-fit values are consistent with the standard model expectations in all cases.

4.7.2 Model dependent results

The results are interpreted in the context of the MSSM for the $m_h^{\text{mod}+}$ and hMSSM benchmark scenarios. Exclusion limits are set at the 95% CL in the m_A - $\tan\beta$ plane, as illustrated in Figure 4.19. The blue shaded areas bounded by the solid black lines indicate the observed exclusion regions. The expected exclusion contour is shown by the dashed black line, where the 68% and 95% confidence intervals are shown by the grey bands. For the $m_h^{\text{mod}+}$ scenario, the shaded red area indicates the region of the parameter space that does not yield a h boson mass compatible with the 125 GeV Higgs within $\pm 3 \text{ GeV}$.

The limits set for the hMSSM scenario are more stringent than those set for the $m_h^{\text{mod}+}$. This

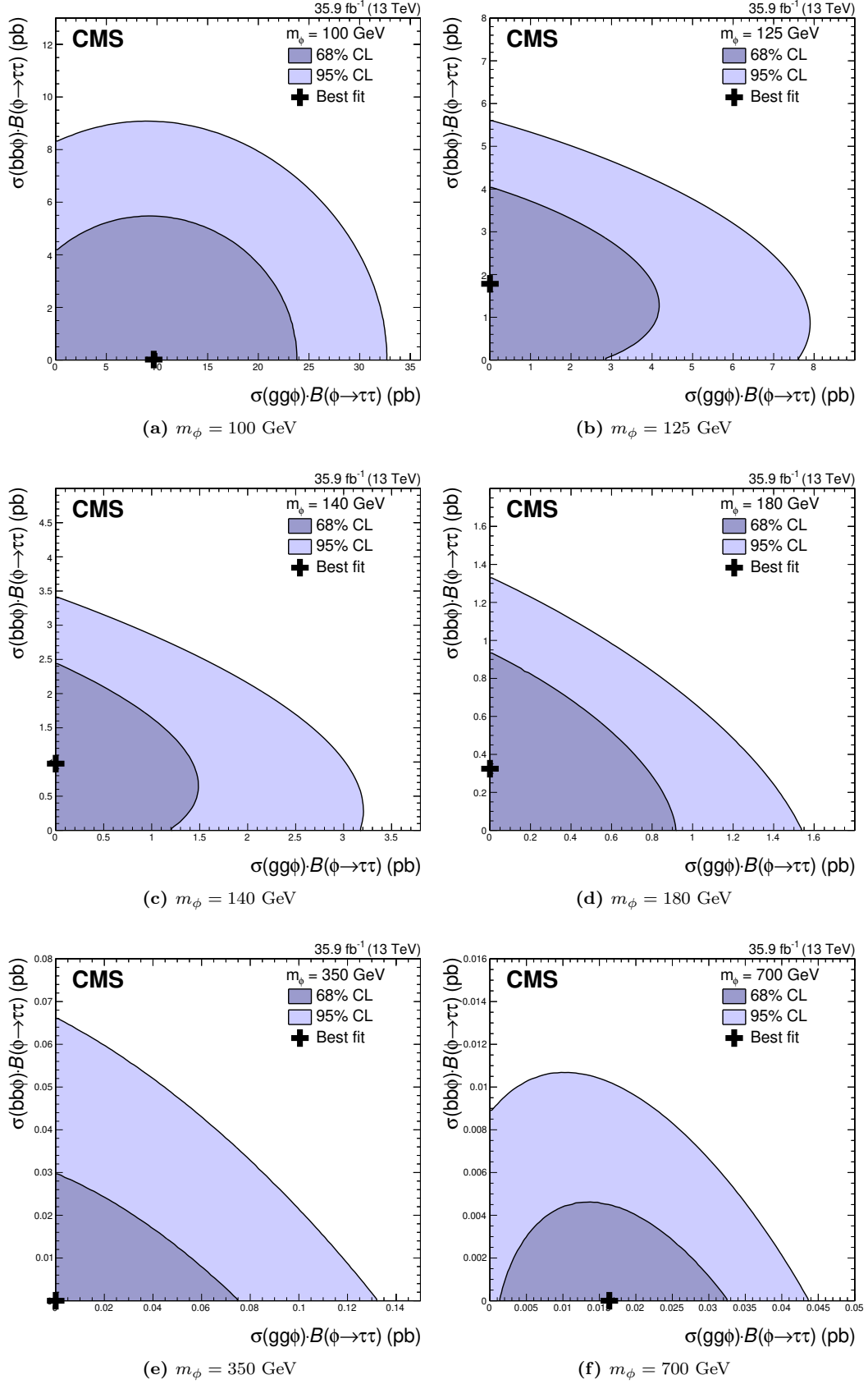


Figure 4.18: Two dimensional likelihood scans of the $\sigma \times \mathcal{B}^{\tau\tau}$ for the $gg\phi$ (x-axis) and $bb\phi$ (y-axis) processes. The black crosses indicate the best fit-values, and the purple bands show the 68% and 95% confidence regions. The scans are shown for a representative subset of the tested mass points [1].

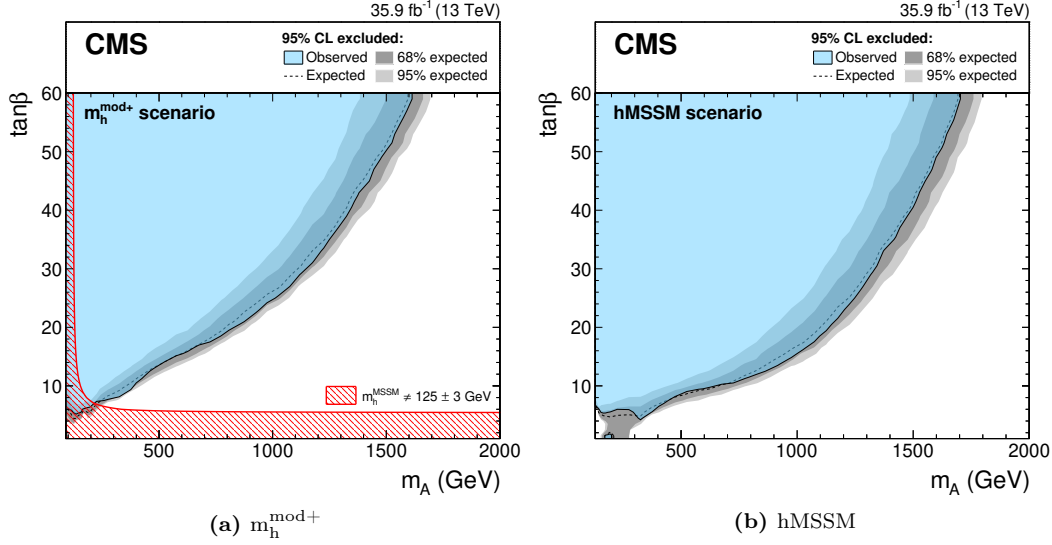


Figure 4.19: The observed and expected 95% CL exclusion contours for the $m_h^{\text{mod}+}$ (left) and hMSSM (right) benchmark scenarios. The observed exclusion regions are shown by the shaded blue areas. The dashed black lines show the expected exclusion region, and the 68% and 95% confidence intervals are shown by the grey bands [1]. For the left inset, the shaded red area indicates the region of the parameter space that does not yield a h boson with mass ± 3 GeV from the mass of the observed Higgs boson at 125 GeV.

is a result of the assumption that the Higgs bosons decay only to SM particles in the hMSSM. The branching ratios to taus are consequently larger in this scenario compared to the $m_h^{\text{mod}+}$, as a share of the total width is taken up by decays into SUSY particles for the latter. The hMSSM limits also show a small region excluded at low $\tan\beta$ close to $m_A = 200$ GeV. This is due to the $gg\phi$ cross section increasing with decreasing $\tan\beta$ for $\tan\beta \lesssim 3$. This is partially offset by the decreasing branching ratio to taus, but overall the $\sigma \times \mathcal{B}^{\tau\tau}$ increases with decreasing $\tan\beta$ in this region.

The model dependent limits can be compared to the results presented by the ATLAS Collaboration in Ref. [175], using a similar integrated luminosity of 36.1 fb⁻¹. The observed limits set by CMS are more stringent in the m_A region $\gtrsim 1.3$ TeV, whereas ATLAS are more sensitive to m_A values between 250–350 GeV and 0.7–1.3 TeV. The observed limits set by both experiments between 350–700 GeV are very comparable. The results presented by ATLAS also do not consider the mass region below 250 GeV, and therefore the CMS analysis solely excludes a region of the parameter space for m_A values between 90–250 GeV and $\tan\beta$ values $\gtrsim 7$.

4.8 Cross check using alternative background methods

This analysis constitutes the first time that the FF method has been used by CMS in searches for $\phi \rightarrow \tau\tau$ resonances. A cross-check of the analysis is performed in the $\tau_e\tau_h$ and $\tau_\mu\tau_h$ channels using alternative background methods to estimate the $\text{jet} \rightarrow \tau_h$ background. These alternative methods are the same as were used for previous CMS results such as in Ref. [145]. The motivations for this cross-check are to give confidence in the background estimation, and to understand potential improvements in the analysis sensitivity.

4.8.1 Alternative background methods

The alternative background methods used are the so called classical methods described in Section 3.5.4, where the QCD is estimated from same-sign data and the W+jets yield is estimated for a high- m_T control region. For the b-tag categories, the W+jet background method is modified slightly with respect to the procedure outlined in Section 3.5.4.2 as described below.

The W+jets method described previously estimates the W normalisation in a high- m_T control region, which is selected by requiring $m_T > 70$ GeV and keeping all other event selections the same. For the b-tag categories, the $m_T > 70$ GeV requirement alone does not give sufficient purity of W+jets events to obtain a reasonable estimate of the W yield. This is due to the large $t\bar{t}$ contribution to the backgrounds in this control region, which must be subtracted using MC. The control region selection is suitably modified by relaxing the requirement for the events to contain at least one b-tagged jet, instead requiring the event to contain at least one jet (tagged or untagged) with $p_T > 20$ GeV. To obtain an estimate for the W yield, Equation 3.28 must then be modified to include an additional extrapolation factor, to extrapolate from the relaxed b-jet selection ($N_{\text{jets}20} \geq 1$) to the b-tag category selection ($N_{\text{b-jets}} \geq 1$):

$$N_W^{\text{OS,low-}m_T} = N_W^{\text{SS,high-}m_T} \cdot R_W \cdot F_{\text{high} \rightarrow \text{low-}m_T} \cdot F_{N_{\text{jets}20 \geq 1} \rightarrow N_{\text{b-jets} \geq 1}}. \quad (4.14)$$

4.8.2 Control regions

The control regions used for the estimation of the W+jet and QCD backgrounds are included in the fit as one-bin control regions to constrain the QCD and W yields. This includes the same-sign regions used to estimate the QCD, and the opposite-sign and same-sign high- m_T control regions used to estimate the W+jets yield. These control regions are added for both the no b-tag and b-tag categories. Figure 4.20 illustrates the control regions added for each of these categories, where the loose- m_T and tight- m_T subcategory signal regions ("SR") are also shown. The loose- m_T and tight- m_T subcategories differ only by their m_T selections, and therefore the same two W+jets control regions are to scale the W+jets yields in both subcategories. These control regions are labelled "W OS CR" and "W SS CR" in Figure 4.20. One QCD control region per subcategory (labelled "QCD CR 1" and "QCD CR 2" in Figure 4.20) are also included. For each category four unconstrained rate parameters are used to scale the QCD and W+jets yields in the signal regions and control regions, which includes:

- One parameter to scale the QCD yields simultaneously in the tight- m_T subcategory and in QCD CR 1.
- One parameter to scale the QCD yields simultaneously in the loose- m_T subcategory and in QCD CR 2.
- One parameter to scale the QCD yields simultaneously in W OS CR and W SS CR.
- One parameter to scale the W+jets yields simultaneously in: the tight- m_T subcategory, the loose- m_T subcategory, QCD CR 1, QCD CR2, W OS CR, and W SS CR.

4.8.3 Systematic uncertainties

The systematic uncertainty template described in Section 4.5 is suitably modified by removing all uncertainties specific to the FF method and applying the following additional uncertainties:

- **QCD normalisation**

The uncertainties on the QCD normalisation are due to the uncertainties on the measured

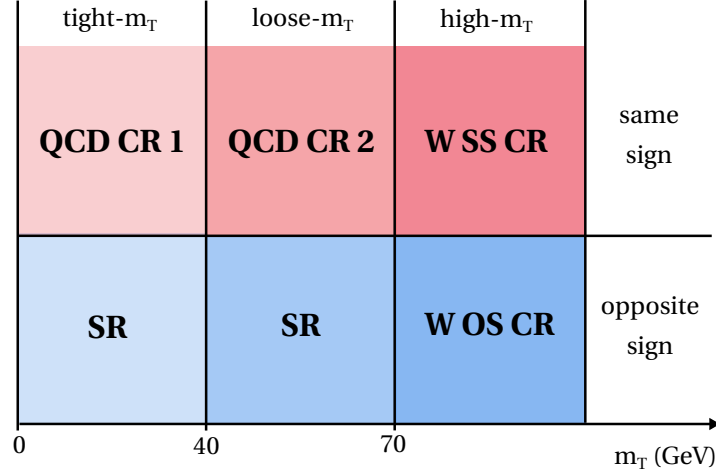


Figure 4.20: Diagram illustrating the one-bin control regions included in the fit to constrain the QCD and W+jet normalisations. The loose- m_T and tight- m_T signal regions are also shown.

R_{QCD} factors used to scale the same-sign data. The R_{QCD} factors are extracted from a fit to the data in the e/μ anti-isolated region - the fit uncertainties are included as systematic uncertainties. Additional uncertainties are also included to account for differences in the R_{QCD} in the e/μ anti-isolated region and the signal region. The size of the total R_{QCD} uncertainties range between 4–32% depending on the channel and the category.

- **W+jet normalisation**

The uncertainty on the W+jets normalisation includes the uncertainty on the R_W and the F_x extrapolation factors described in Section 3.5.4.2 and above. The uncertainties on R_W include both statistical and systematic contributions. The statistical contributions range between 2–3%. The systematic contributions are estimated by comparing the W OS/SS ratios in data and MC in a control region in which the τ_h candidate is anti-isolated. The size of these systematic uncertainties range between 1–2%. The uncertainties on the F_x extrapolation factors also include statistical and systematic contributions. The statistical uncertainties range between 2–11%. The systematic uncertainties are estimated using a sample of $Z \rightarrow \mu\mu$ events, exploiting the similarity between the W+jets and Drell-Yan processes. The dimuon 4-momenta is corrected to account for the differences in the Z and W boson masses. One of the muons is then selected at random and its momentum is added to the E_T^{miss} . The other muon is then combined with a τ_h candidate and required to pass the usual $\tau_\mu\tau_h$ channel selections. The uncertainties are then estimated by comparing the extrapolation factors measured in data and MC for these pseudo-W events. The size of the uncertainties ranges between 2–16%.

- **jet $\rightarrow \tau_h$ fake-rate uncertainty**

The jet $\rightarrow \tau_h$ fake-rate for W+jets events is measured as a function of p_T for both data and MC using $\tau_\mu\tau_h$ events in the high- m_T region. The difference between the two measurements is taken as a p_T dependent shape uncertainty applied to the W+jets process. The size of this uncertainty is $\pm 20\% \times p_T/(100 \text{ GeV})$. For the other processes that are taken entirely from simulation, a constant 20% uncertainty on the jet $\rightarrow \tau_h$ fake-rate is applied.

- **Cross section and luminosity uncertainties**

For jet $\rightarrow \tau_h$ processes that are estimated entirely from simulation the uncertainties on the

luminosity and cross sections described in Section 4.5 are applied.

4.8.4 Results and comparisons

The classical background method suffers more from limited statistics in the background templates in the high m_T^{tot} regions. Before the fit to data, the m_T^{tot} distributions must therefore be re-binned to prevent empty bins in the total background histograms, and to prevent bins having excessively large statistical uncertainties. An automated re-binning procedure iteratively merges together the m_T^{tot} bins until there are no bins with statistical uncertainties $> 90\%$, and all bins are populated.

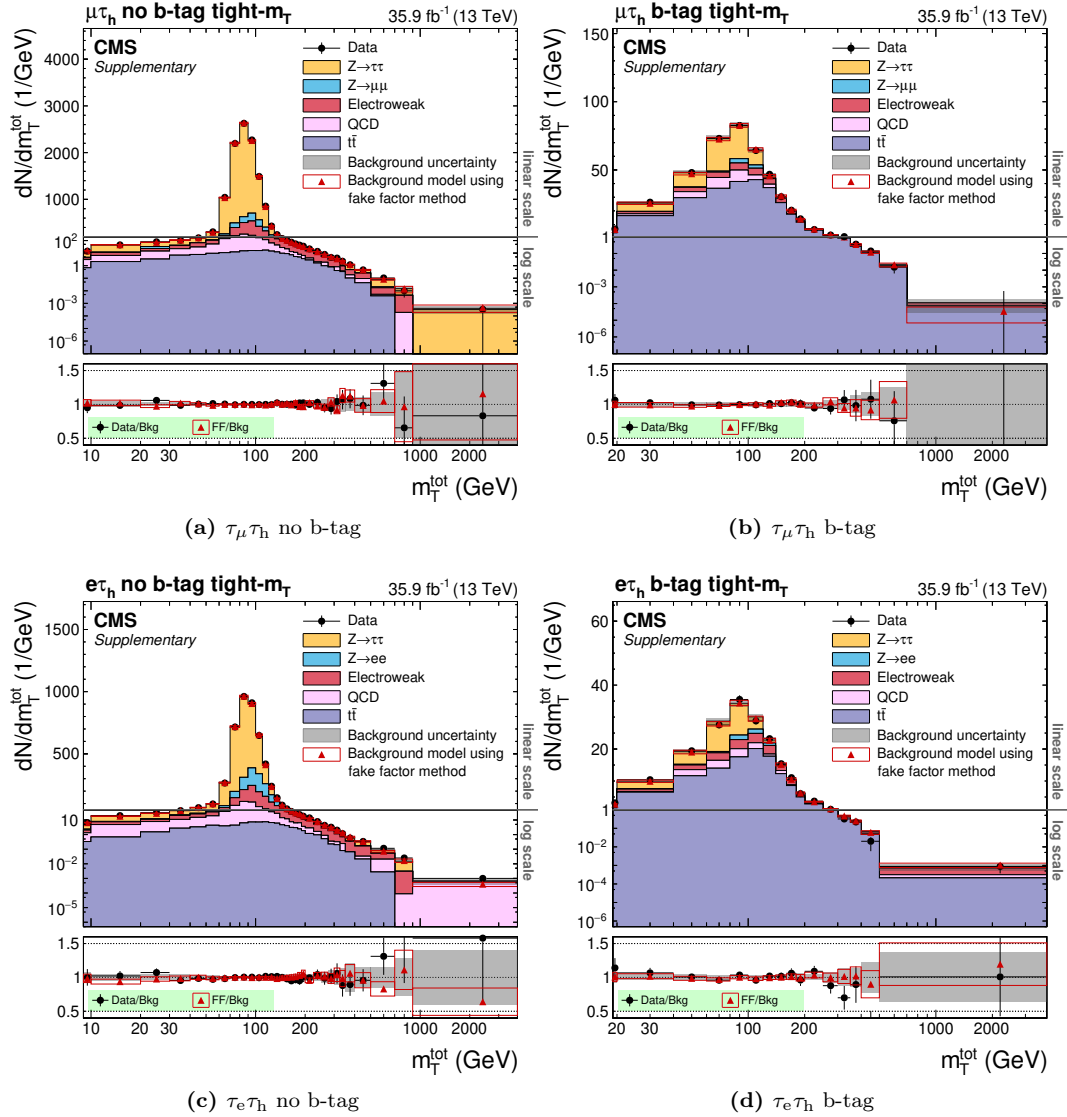


Figure 4.21: Observed and expected distributions of the m_T^{tot} variable in the no b-tag (left) and b-tag (right) categories using the classical background method described in the text. The distributions are shown for the $\tau_\mu\tau_h$ channel in the upper row and the $\tau_e\tau_h$ channel in the lower row. Only the most sensitive tight- m_T subcategories are shown in all cases. The uncertainty band shown in shaded grey includes both systematic and statistical uncertainties. The expected background distributions and uncertainties using the default background methods are shown by the red triangles and boxes respectively [1].

The m_T^{tot} distributions after the simultaneous fit to data are shown in Figure 4.21. Only

the most sensitive tight- m_T subcategories are shown for both channels. The histograms labelled "Electroweak" include the W +jets background in addition to the single top and diboson processes. The QCD background is shown by the pink histograms. The templates for all other processes producing $\text{jet} \rightarrow \tau_h$ fakes are combined with the corresponding non-jet $\rightarrow \tau_h$ histograms. The expected background and uncertainties estimated using the default background methods (i.e the FF method) are shown by the red triangles and boxes. In all distributions the background expectations are found to agree well with both the data and the FF predictions.

Comparisons of the expected and observed upper limits for the $gg\phi$ and $bb\phi$ processes are shown in Figure 4.22, where the limits are shown separately for the $\tau_e\tau_\mu$ and $\tau_\mu\tau_h$ channels. The solid black and solid red lines show the observed and expected limits obtained when using the FF method, and the green and yellow bands show the 68% and 95% confidence intervals. The observed and expected limits obtained when using the alternative background methods are shown by the solid blue lines and dashed blue lines respectively. The differences between the observed and expected limits are found to follow similar trends for both background methods.

For the $\tau_\mu\tau_h$ channel, the FF method is observed to outperform the classical method for masses up to ≈ 800 GeV. This is a result of the smaller systematic uncertainties for the FF method compared to the classical. The relative size of the uncertainties is illustrated in Figure 4.23 for the no b-tag category in the $\tau_\mu\tau_h$ channel. The uncertainties are observed to be up to 50% smaller in the region $m_T^{\text{tot}} \lesssim 300$ GeV, where the bulk of the events for $m_\phi \leq 800$ GeV signals reside. For the higher mass region, the classical method slightly outperforms the FF method. This is a result of the classical method predicting a slightly larger signal acceptance in the high m_T^{tot} bins. The signal acceptance for the FF method is reduced by a downwards pull on the p_T dependent "high- p_T tau ID efficiency" uncertainty (described in Section 4.5), which results in fewer expected signal events for the high mass points, and consequently less stringent limits. The signal/background predictions by both methods are however entirely consistent with each other within the uncertainties.

For the $\tau_e\tau_h$ channel, the FF method outperforms the classical method across almost all of the mass range. The same conclusions can be drawn for the $\lesssim 800$ GeV region as for the $\tau_\mu\tau_h$ channel. For the high mass points, the reason for the improvement in the limits when using the FF method is that several of the high m_T^{tot} bins are merged for the classical method by the automated re-binning procedure described previously. This coarser binning reduces the sensitivity significantly for the large signal masses. When the same m_T^{tot} binning is applied for the FF method the limits for the two methods are more comparable. In this case the high mass limits for the classical method are even slightly more stringent for the same reason discussed for the $\tau_\mu\tau_h$ channel.

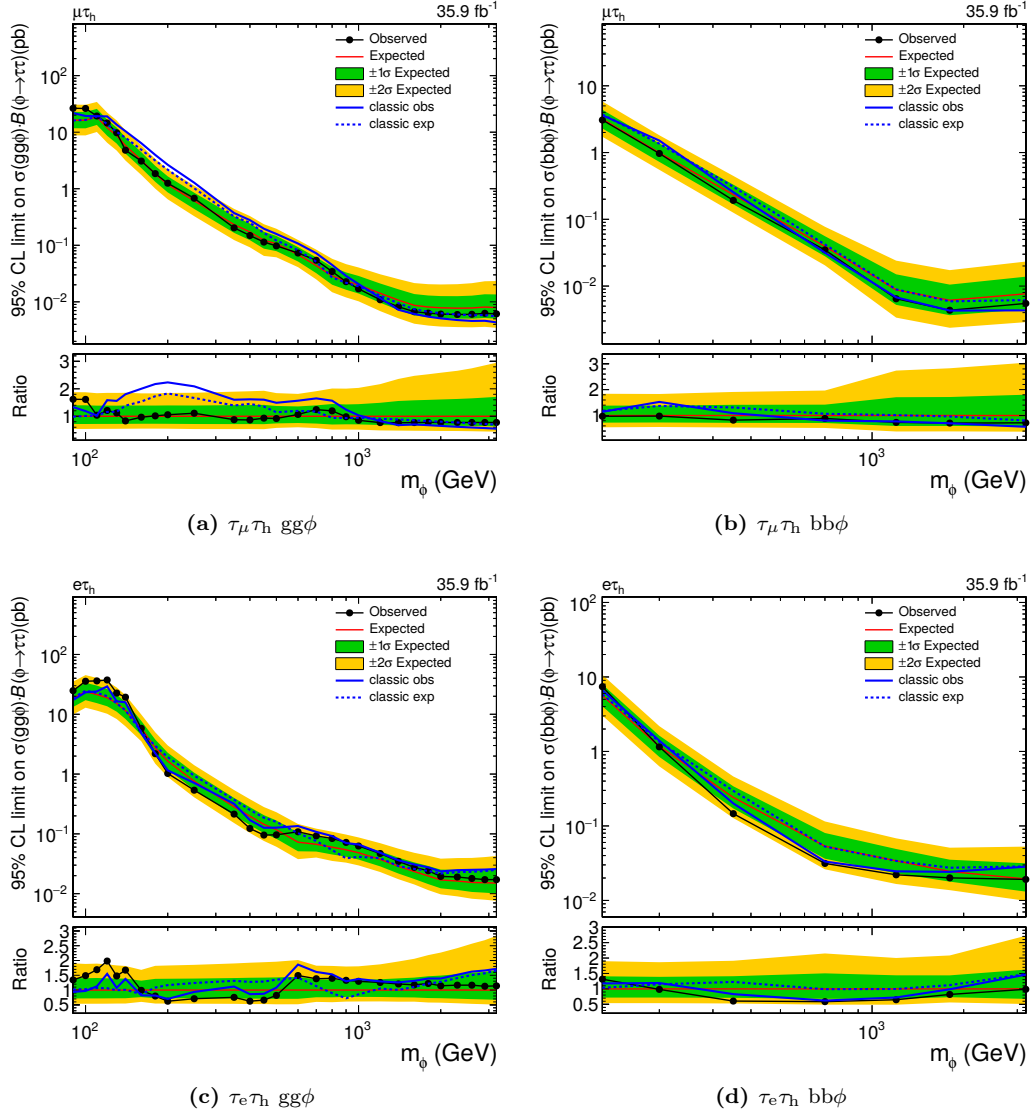


Figure 4.22: Comparisons of the observed and expected limits obtained using the FF and classical background methods. The limits are shown for both $gg\phi$ (left) and $bb\phi$ (right) processes, and separately for the $\tau_\mu\tau_h$ (upper) and $\tau_e\tau_h$ (lower) channels. The observed and expected limits obtained using the FF method are shown by the solid black and red lines, and the green and yellow bands show the 68% and 95% confidence intervals. The expected and observed limits obtained using the classical methods are shown by the solid blue and dashed blue lines.

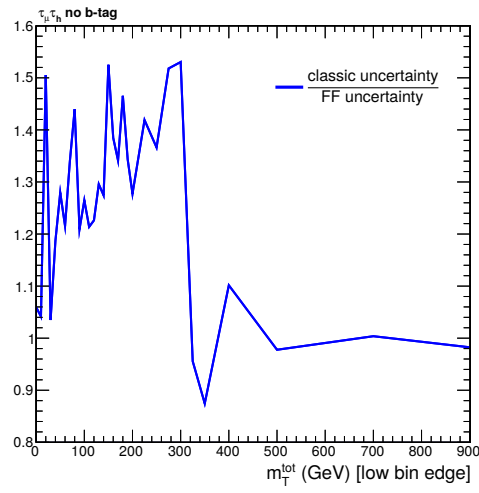


Figure 4.23: The ratio of the post-fit uncertainties on the m_T^{tot} background template obtained using the classical and FF methods. The ratio is shown for the no b-tag tight- m_T category in the $\tau_\mu \tau_h$ channel.

Chapter 5

Measurement of the Higgs boson $\mathcal{CP}$ properties

This chapter describes the measurement of the $\mathcal{CP}$ properties of the Higgs coupling to top quarks. As discussed in Chapter 1, these properties can be determined using angular correlations between jets originating from GF-induced $H+2j$ production. The Higgs bosons are reconstructed via their decay into ditau pairs, although as noted previously any Higgs decay final states may be used for this measurement. The four most sensitive ditau final states are considered, namely the $\tau_e\tau_\mu$, $\tau_e\tau_h$, $\tau_\mu\tau_h$, and $\tau_h\tau_h$ channels. This analysis utilises the datasets collected by CMS in 2016 and 2017, which correspond to a total integrated luminosity of 71.8 fb^{-1} . Due to different detector conditions, object reconstruction algorithms, and selection thresholds, the 2016 and 2017 datasets are analysed separately, and a combination is performed to determine the final results. The results presented include the measurement of the mixing angle α_{hgg} defined in Equation 1.35, and the GF cross section times branching ratio relative to the SM expectation. The results are also interpreted in terms of the coupling strengths of the top quark to the scalar and pseudoscalar Higgs fields.

5.1 Signal modelling

The generation of the GF signal samples is performed at NLO precision using `Madgraph5_aMC@NLO`, as implemented in Ref. [176]. Inclusive samples are generated with 0, 1, and 2 outgoing partons and merged according to the FxFx scheme. To increase the statistics of the $\mathcal{CP}$ -sensitive dijet events, additional exclusive samples with 2 outgoing partons were also simulated. The inclusive and exclusive samples are then stitched together by the application of event weights, such that the parton multiplicity fractions are preserved. As will be discussed below, samples are only required for three scenarios: the pure $\mathcal{CP}$ -even scenario, the pure $\mathcal{CP}$ -odd scenario, and a so called maximum-mixing scenario - where the amount of $\mathcal{CP}$ -violation is maximal.

For a generic $\mathcal{CP}$ -scenario, the cross section relative to the SM cross section, σ_{SM} , can be expressed as:

$$\mu_{\text{ggh}} = \frac{\sigma}{\sigma_{\text{SM}}} = \kappa_t^2 + \frac{\sigma_{\text{PS}}}{\sigma_{\text{SM}}} \cdot \tilde{\kappa}_t^2, \quad (5.1)$$

where σ_{PS} is the pseudoscalar cross section when $(|\kappa_t|, |\tilde{\kappa}_t|) = (0, 1)$ which is equal to $(\frac{3}{2})^2 \times \sigma_{\text{SM}}$. The larger pseudoscalar cross section is a consequence of the relative sizes of the form factors defined in Equation 1.34, since $|a_3| = \frac{3}{2} \cdot |a_2|$ for $\kappa_t = \tilde{\kappa}_t$. The differential distributions can also be formed from the superposition of pure $\mathcal{CP}$ -even ("SM"), pure $\mathcal{CP}$ -odd ("PS"), and interference

("int") distributions according to

$$\frac{d\sigma}{dx} = \kappa_t^2 \cdot \frac{d\sigma_{\text{SM}}}{dx} + \tilde{\kappa}_t^2 \cdot \frac{d\sigma_{\text{PS}}}{dx} + 2\kappa_t\tilde{\kappa}_t \cdot \frac{d\sigma_{\text{int}}}{dx}. \quad (5.2)$$

It is convenient to re-parametrise the above formula in terms of μ_{ggh} and the mixing angle α_{hgg} defined in Equation 1.35, where α_{hgg} can be written in terms of the $\tilde{\kappa}_t$, κ_t , and the two cross sections:

$$\tan \alpha_{\text{hgg}} = \frac{a_3}{a_2} = \sqrt{\frac{\sigma_{\text{PS}}}{\sigma_{\text{SM}}}} \cdot \frac{\tilde{\kappa}_t}{\kappa_t}. \quad (5.3)$$

After the re-parametrisation Equation 5.2 becomes

$$\frac{d\sigma}{dx} = \mu_{\text{ggh}} \cdot \left(c_\alpha^2 \cdot \frac{d\sigma_{\text{SM}}}{dx} + s_\alpha^2 \cdot \frac{\sigma_{\text{SM}}}{\sigma_{\text{PS}}} \cdot \frac{d\sigma_{\text{PS}}}{dx} + 2c_\alpha s_\alpha \cdot \sqrt{\frac{\sigma_{\text{SM}}}{\sigma_{\text{PS}}}} \cdot \frac{d\sigma_{\text{int}}}{dx} \right), \quad (5.4)$$

where c_α and s_α are used as shorthand for $\cos \alpha_{\text{hgg}}$ and $\sin \alpha_{\text{hgg}}$. The pure $\mathcal{CP}$ -even and pure $\mathcal{CP}$ -odd distributions are generated with $\alpha_{\text{hgg}} = 0^\circ$ and $\alpha_{\text{hgg}} = 90^\circ$ respectively. The interference distribution is obtained by subtracting the $\mathcal{CP}$ -even and $\mathcal{CP}$ -odd contributions from the maximum-mixing ("MM") sample, which is generated with $\alpha_{\text{hgg}} = 45^\circ$. Substituting $\frac{d\sigma_{\text{MM}}}{dx}$ for $\frac{d\sigma_{\text{int}}}{dx}$ we get:

$$\frac{d\sigma}{dx} = \mu_{\text{ggh}} \cdot \left((c_\alpha^2 - c_\alpha s_\alpha) \cdot \frac{d\sigma_{\text{SM}}}{dx} + (s_\alpha^2 - c_\alpha s_\alpha) \cdot \frac{\sigma_{\text{SM}}}{\sigma_{\text{PS}}} \cdot \frac{d\sigma_{\text{PS}}}{dx} + 2c_\alpha s_\alpha \cdot \frac{\sigma_{\text{SM}}}{\sigma_{\text{MM}}} \cdot \frac{d\sigma_{\text{MM}}}{dx} \right) \quad (5.5)$$

By inspecting this formula we can see that whatever σ_{PS} and σ_{MM} we scale the $\mathcal{CP}$ -odd and maximum-mixing distributions to is completely arbitrary, since the $\frac{d\sigma_i}{dx}$ distributions are always normalised to σ_{SM} by the $\frac{\sigma_{\text{SM}}}{\sigma_i}$ terms. For simplicity we choose $\sigma_{\text{MM}} = \sigma_{\text{PS}} = \sigma_{\text{SM}}$ in which case the above equation simplifies to

$$\frac{d\sigma}{dx} = \mu_{\text{ggh}} \cdot \left((c_\alpha^2 - c_\alpha s_\alpha) \cdot \frac{d\sigma_{\text{SM}}}{dx} + (s_\alpha^2 - c_\alpha s_\alpha) \cdot \frac{d\sigma_{\text{PS}}}{dx} + 2c_\alpha s_\alpha \cdot \frac{d\sigma_{\text{MM}}}{dx} \right). \quad (5.6)$$

On further inspection of the formulae we also see that variations in α_{hgg} affect only the differential distributions while leaving the inclusive yield unchanged. This parametrisation therefore allows us to separate effects that modify the cross section from those that change only the differential variables. A measurement of a deviation in the former would provide only indirect evidence of a non-SM $\mathcal{CP}$ -state, as there are many other BSM effects that could lead to an enhancement or suppression of the cross section without the need for $\mathcal{CP}$ -odd couplings. The differential information on the other hand can provide direct evidence of a non-SM $\mathcal{CP}$ -state, where any value of $\alpha_{\text{hgg}} \neq 0^\circ$ would be non-SM and $0^\circ < |\alpha_{\text{hgg}}| < 90^\circ$ would imply $\mathcal{CP}$ violation.

The signal is scaled to the cross section times branching ratio calculated in Ref. [30]. The cross section is computed at next-to-next-to-next-to-leading order (N3LO) QCD and NLO electroweak accuracies, including the effects of finite quark masses (t, b, and c) up to NLO.

5.1.0.1 Generator-level comparisons

As discussed in Chapter 1, the variable most sensitive to the Higgs $\mathcal{CP}$ -state is the η ordered azimuthal angular separation between the two highest p_T jets, $\Delta\phi_{jj}$. To compute this variable the two highest p_T jets are selected and $\Delta\phi_{jj}$ is given by

$$\Delta\phi_{jj} = \phi(j_1) - \phi(j_2), \text{ with } \eta(j_1) < \eta(j_2). \quad (5.7)$$

Figure 5.1 compares the generator-level $\Delta\phi_{jj}$ distributions for three $\mathcal{CP}$ -states: $\alpha = 0^\circ$ ($\mathcal{CP}$ -even), $\alpha = 90^\circ$ ($\mathcal{CP}$ -odd), and $\alpha = 45^\circ$ (maximally $\mathcal{CP}$ violating), applying different cuts on the invariant

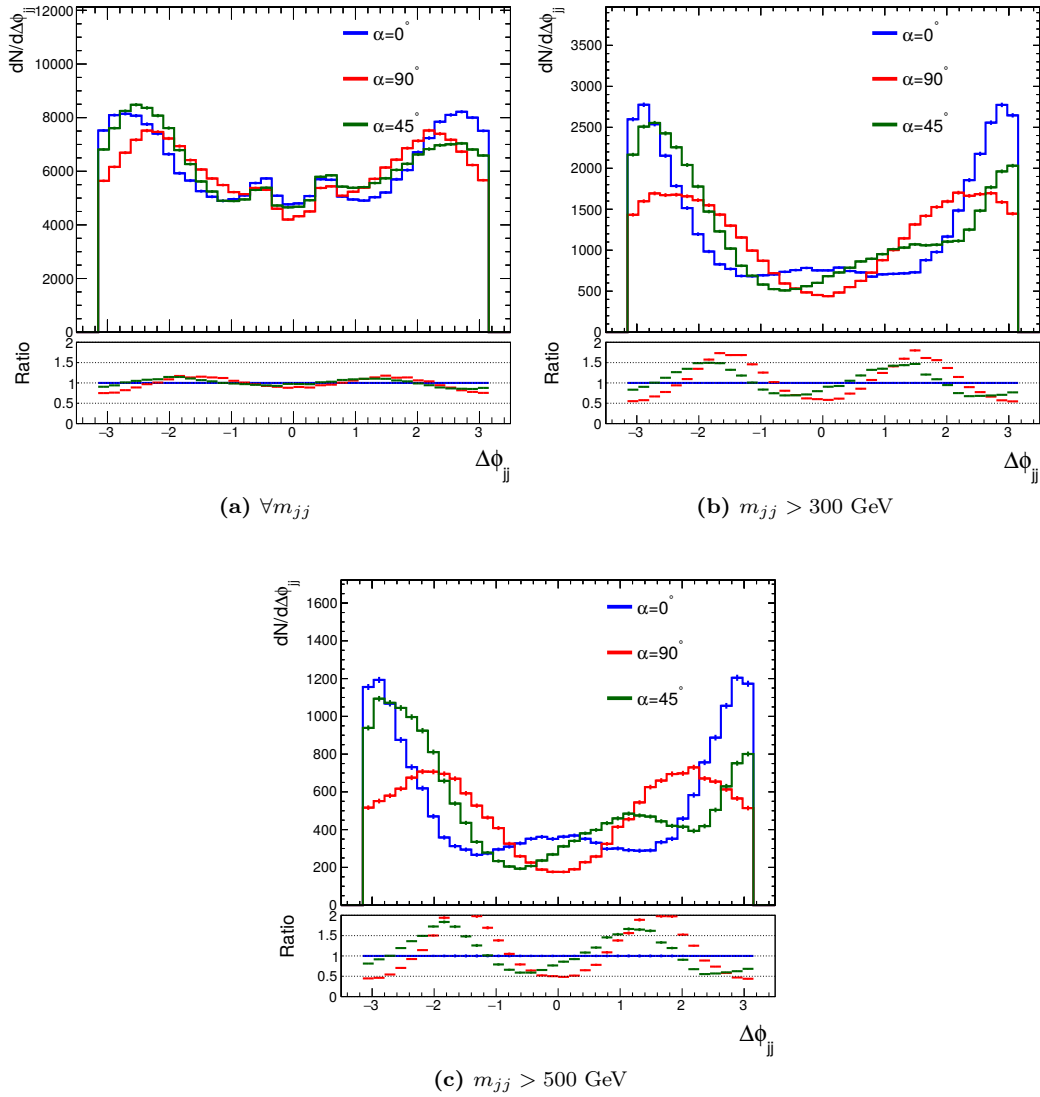


Figure 5.1: Generator-level comparisons of $\alpha = 0^\circ$ (blue), $\alpha = 90^\circ$ (red), and $\alpha = 45^\circ$ (green), are shown for the $\Delta\phi_{jj}$ variable. Comparisons are shown for inclusive dijet events (a) and with additional cuts of $m_{jj} > 300 \text{ GeV}$ (b) and $m_{jj} > 500 \text{ GeV}$ (c) applied.

mass of the jets, m_{jj} . Inset (a) shows the $\Delta\phi_{jj}$ distribution with no additional selections applied. Although some separation between the different $\mathcal{CP}$ -states is visible this is washed out by the presence of events arising from processes such as gluon-splitting which are not sensitive to the Higgs $\mathcal{CP}$ state. These uninteresting events can be removed by applying VBF-like selections such as a cut on m_{jj} . Insets (b) and (c) compare the $\Delta\phi_{jj}$ distributions after applying cuts of $m_{jj} > 300 \text{ GeV}$ and $m_{jj} > 500 \text{ GeV}$ respectively, the separation is observed to improve significantly after applying tighter m_{jj} cuts.

Figure 5.2 shows comparisons of several other variables. The Higgs transverse momentum, p_T^H , shown in inset (a) shows no significant differences between different $\mathcal{CP}$ -states. Insets (b)–(d) show the p_T distributions for the three leading jets; differences between the $\mathcal{CP}$ -states are visible here especially in the high p_T bins. The differences in the p_T distributions for the leading jet are relatively small $\lesssim 5\%$ but are more noticeable for the sub-leading and sub-sub-leading jets, where the yields in the high p_T bins for a $\mathcal{CP}$ -odd Higgs can be up to 20% lower than for a $\mathcal{CP}$ -even Higgs.

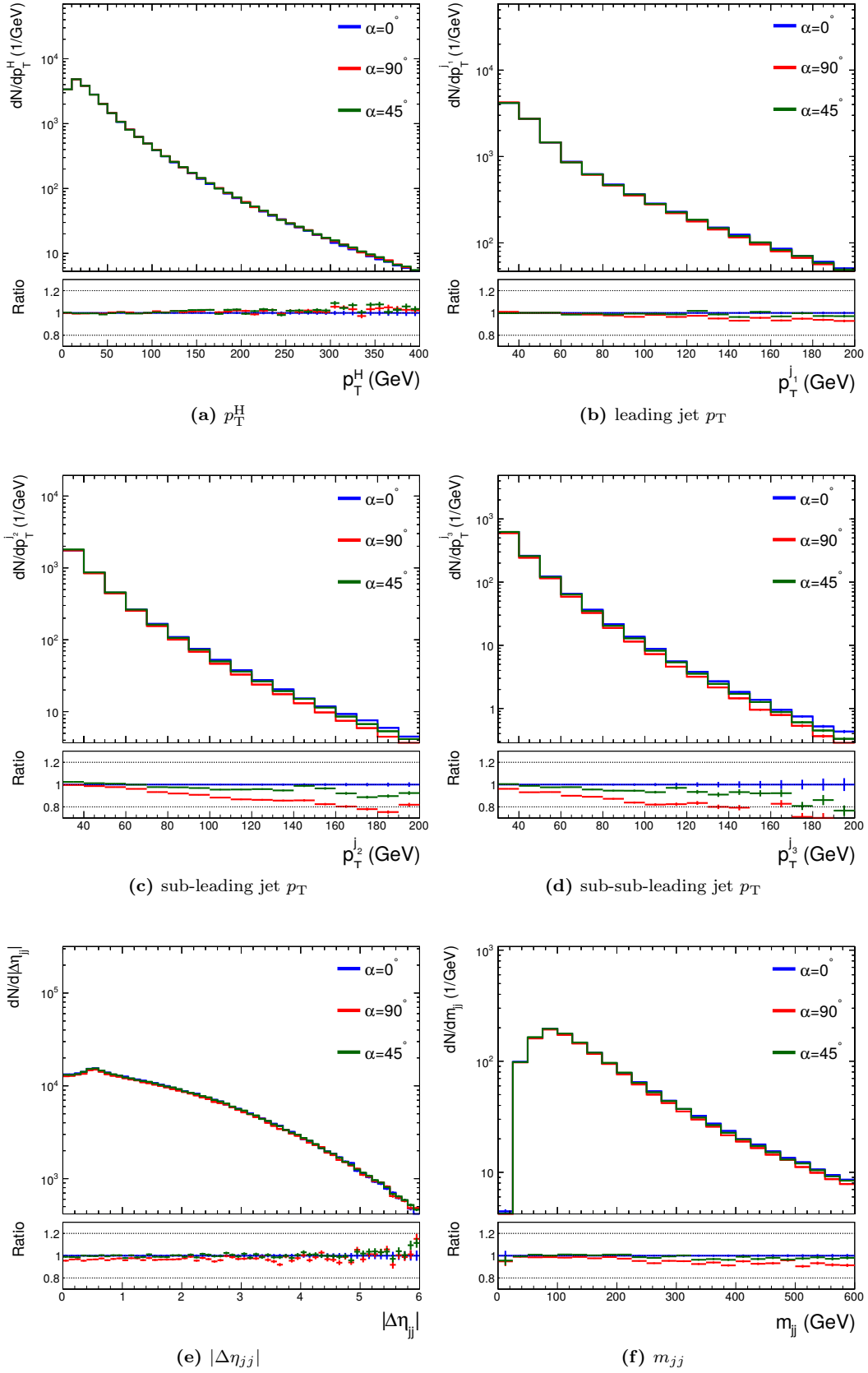


Figure 5.2: Generator-level comparisons of $\alpha = 0^\circ$ (blue), $\alpha = 90^\circ$ (red), and $\alpha = 45^\circ$ (green), are shown for several variables.

The pseudorapidity separation between the two jets, $|\Delta\eta_{jj}|$, shown in inset (e) is very similar for the different $\mathcal{CP}$ -states. Finally, the m_{jj} distribution, shown in inset (f), shows some differences between the different $\mathcal{CP}$ -states in the high m_{jj} bins - this is likely a by-product of the jet p_T variations, since m_{jj} and p_T^j are correlated.

5.1.0.2 Top quark mass effects

The signal samples are produced using `Madgraph5_aMC@NLO` in an effective field theory (EFT) approach, meaning the mass of the top quark is assumed to be large ($m_{\text{top}} = \infty$). This approximation starts to break down for Higgs events with $p_T^H > m_{\text{top}}$, where the differential cross section, $\frac{d\sigma}{dp_T^H}$, in the EFT becomes large compared the full theory. The `HNNLO` tool [177] can be used to compute the GF cross section up to NNLO accuracy using either the EFT approach or with quark mass effects included up to NLO. It is also possible to compute the differential cross section $\frac{d\sigma}{dp_T^H}$ up to NLO accuracy. The differential cross sections obtained at NLO with `HNNLO` with and without top quark mass effects included are shown in Figure 5.3. It is clear from Figure 5.3 that the EFT approach holds well below $p_T^H \approx 200$ GeV but rapidly starts to break down for more boosted events. The fraction of such boosted events is increased by the application of even the loosest possible p_T cuts on the tau decay products driven by the trigger thresholds, plus additional requirements of $m_{jj} > 300$ GeV and $N_{\text{jets}} \geq 2$ to select $\mathcal{CP}$ sensitive events. For the $\tau_h\tau_h$ channel $\approx 45\%$ of such events have $p_T^H > 200$ GeV. For the $\tau_l\tau_h$ and $\tau_e\tau_\mu$ channels the fraction of events with $p_T^H > 200$ GeV is $\approx 35\%$, which although smaller is still significant.

A correction is applied to the GF MC as a function of p_T^H to account for the top quark mass effects. This is derived using `HNNLO` at NLO accuracy. The correction is defined as:

$$\text{correction} = C \cdot \frac{\frac{d\sigma^{\text{full}}}{dp_T^H}}{\frac{d\sigma^{\text{EFT}}}{dp_T^H}}, \quad (5.8)$$

where $\frac{d\sigma^{\text{full}}}{dp_T^H}$ is the differential cross section taking into account top quark mass effects, $\frac{d\sigma^{\text{EFT}}}{dp_T^H}$ is the differential cross section in the EFT, and $C = 1.006$ is a constant factor derived to leave the inclusive cross section unchanged by the correction.

5.1.1 Theoretical uncertainties

This section describes the theoretical uncertainties applied to the GF signal process in the analysis. Theoretical uncertainties on the GF signal are split into two contributions, one that affects only the inclusive event yield (i.e the uncertainties on the cross section and branching ratio), and one that affects the shapes of variable distributions but keeps the inclusive event yield constant.

5.1.1.1 Cross section and branching ratio uncertainties

The recommended uncertainties on the GF cross section and $H \rightarrow \tau\tau$ branching ratio provided in Ref. [30] are applied. The uncertainties on the cross section include a 3.9% uncertainty due to variations of the μ_R and μ_F scales and a 3.2% uncertainty from PDFs+ α_s variations. The uncertainty on the $H \rightarrow \tau\tau$ branching-fraction includes a 1.7% uncertainty due to missing higher order corrections, a 0.99% uncertainty on the quark masses, and a 0.62% uncertainty on α_s .

5.1.1.2 Shape/acceptance uncertainties

The theoretical uncertainties applied in the analysis include: μ_R/μ_F scale uncertainties, parton-shower and UE uncertainties, an uncertainty due to the treatment of the finite quark mass effects,

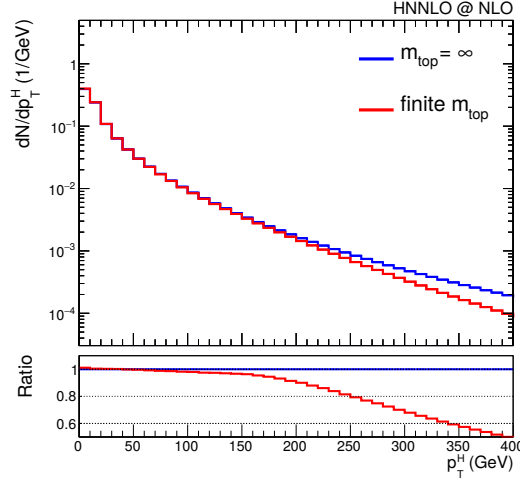


Figure 5.3: Higgs p_T distributions generated at NLO using HNNLO using EFT and taking the top quark mass into account. The ratio in the bottom inset shows the correction that is applied to the GF MC.

and jet bin-migration uncertainties. Uncertainties due to the $\text{PDF} + \alpha_s$ variations were considered but found to be negligible ($< 1\%$) compared to the other contributions.

μ_R/μ_F scale uncertainties

The standard procedure for estimating scale uncertainties is to vary the μ_R and μ_F scales used in the matrix-element computation independently by factors of 2 and $\frac{1}{2}$ times the nominal values under the constraint $\frac{1}{4} < \frac{\mu_R}{\mu_F} < 4$, and take the envelope of the variations as the uncertainty [30]. For each (μ_R, μ_F) choice the inclusive event yield is scaled to the nominal yield, such that the uncertainty on the cross section described in the previous section is not counted twice. Figure 5.4 shows the effect of the scale variations on two important variables for event selection and categorisation, p_T^H and N_{jets} . The envelope of the variations is found to be defined by the scale choices $\mu_R = \mu_F = 2$ and $\mu_R = \mu_F = \frac{1}{2}$, and therefore these scale choices are used to define the uncertainty variations.

Parton-shower/UE uncertainties

The parton shower uncertainty is estimated by varying the μ_R scale used in the parton shower generation by factors of 2 and $\frac{1}{2}$. The μ_F scale may also be varied but the effect is found to be small compared to the μ_R variations. The UE uncertainty is determined by varying several of the PYTHIA parameters according to the uncertainties on the CUETP8M1 and CP5 tunes [124, 125]. Figure 5.5 shows comparisons of the N_{jets} and p_T^H distributions for the varied UE tunes (labelled "UE up/down") and varied μ_R scale (labelled "PS up/down"). The variation in the UE tunes is found to affect mainly the N_{jets} distribution and has a negligible effect on p_T^H . The variation of the μ_R scale on the other hand affects both the N_{jets} and p_T^H distributions. The parton-shower uncertainty is therefore parametrised as a function of N_{jets} and p_T^H , whereas the UE uncertainty is parametrised only as a function of N_{jets} .

Finite top mass uncertainty

The uncertainty on the top quark mass correction is estimated by varying the μ_R and μ_F scales from the nominal values of $\mu_R = \mu_F = m_H$ to $\mu_R = \mu_F = 2m_H$ and $\mu_R = \mu_F = \frac{1}{2}m_H$. The scales

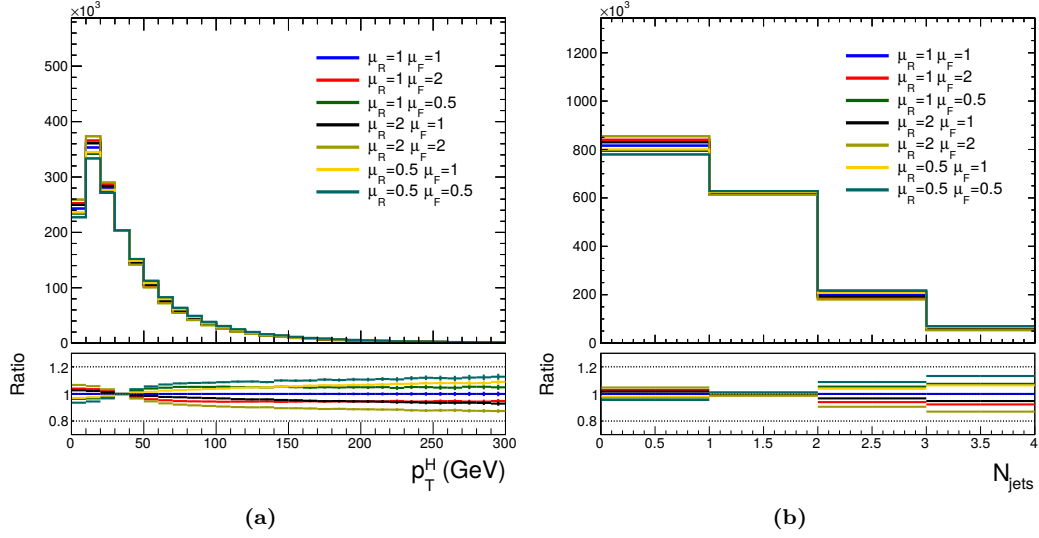


Figure 5.4: Comparison of the p_T^H (left) and N_{jets} (right) distributions for different choices of the μ_R and μ_F scales.

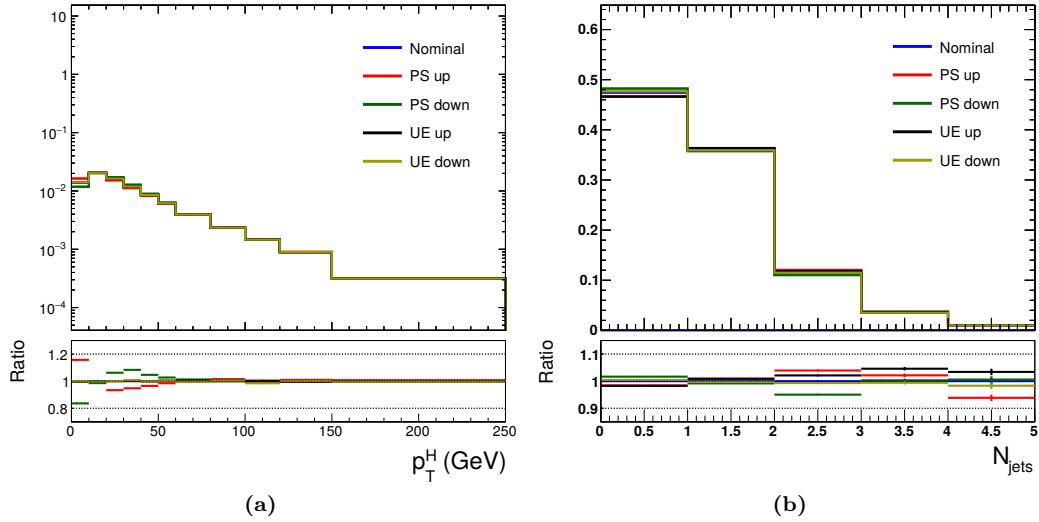


Figure 5.5: Comparison of the p_T^H (left) and N_{jets} (right) distributions for varied μ_R scales in the parton-shower generation (PS up = $\mu_R = 2$, PS down = $\mu_R = \frac{1}{2}$) and for varied UE tunes (UE up/down).

used in both the numerator and denominator are varied simultaneously such that the systematic uncertainties partially cancel. Figure 5.6 shows the size of the uncertainty as a function of p_T^H . The uncertainty grows with p_T^H but is $< 1\%$ for values of p_T^H up to 400 GeV. As this uncertainty is small compared to the other theoretical uncertainties on the p_T^H spectrum it is subsequently neglected.

The top quark mass reweighting assumes that the GF loop is dominated by the top quark contribution as is the case in the SM. While the current experimental evidence supports this hypothesis there is still room within the experimental uncertainties for a non-SM loop structure. In the case where the particle in the loop is heavy ($\gg m_{\text{top}}$), the resulting p_T^H distribution would mimic the EFT prediction. To account for the presence of heavy particles in the loop, the difference

between the reweighted MC and the unweighted MC (i.e the EFT prediction) is taken as an uncertainty.

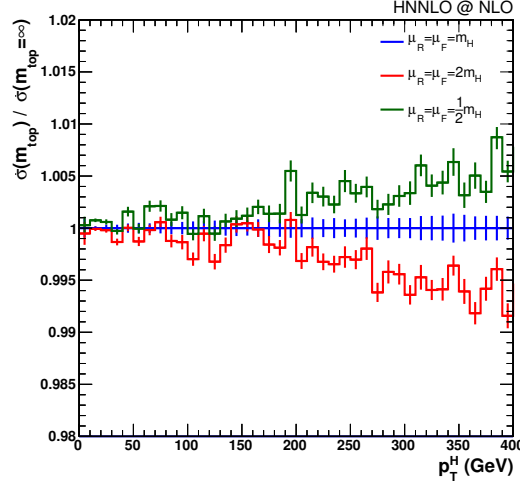


Figure 5.6: The uncertainty on the finite top quark mass correction as a function of p_T^H , estimated by varying the μ_R and μ_F scales by factors of 2 and $\frac{1}{2}$.

Jet-bin migration uncertainties

Many Higgs analyses categorise events according to the jet multiplicities. This helps to reduce backgrounds and separate different Higgs production modes. Dividing the events in this way makes the exclusive $N_{\text{jets}} = X$ cross sections sensitive to soft and collinear emissions. Conventional μ_R/μ_F scale variations may underestimate the perturbative uncertainties on these exclusive cross sections, as they account for yield uncertainties on the N_{jet} bins but underestimate the migration uncertainties between them [30, 178, 179]. The jet-bin migration uncertainties provided in Table 20 in Ref. [30] are applied to account for this effect. The uncertainties relevant for this analysis are those that affect the migration between the 0 and 1 jet boundary ($\Delta_{\text{cut}}^{0/1}$) and the 1 and 2 jet boundary ($\Delta_{\text{cut}}^{1/2}$), which are summarised in Table 5.1.

Table 5.1: The jet-bin migration uncertainties relevant to the analysis taken from Table 20 in Ref. [30], where $\Delta_{\text{cut}}^{i/j}$ is the uncertainty due to migrations across the boundary between the $N_{\text{jets}} = i$ and $N_{\text{jets}} = j$ bins. The signs before the uncertainties are used to indicate correlations across jet bins.

jet bin	$\Delta_{\text{cut}}^{0/1}$	$\Delta_{\text{cut}}^{1/2}$
$N_{\text{jets}} = 0$	4.1%	—
$N_{\text{jets}} = 1$	−8.4%	7.2%
$N_{\text{jets}} \geq 2$	−3.6%	−14.5%

5.2 Event selection

The kinematic cuts on the leptons are driven by the online trigger thresholds, where typically the offline p_T cut is 1 GeV higher than the online selection. Hadronic taus are required to have p_T at least 3–5 GeV higher than the trigger threshold depending on the channel, but often the p_T cuts are increased to help reject $\text{jet} \rightarrow \tau_h$ fakes. To reject leptons and hadronic tau candidates

originating from PU interactions, the particles are required to be associated with the primary vertex by applying cuts on the transverse and/or longitudinal impact parameters, d_{xy} and d_z . Electrons and muons are required to have $d_{xy} < 0.045$ cm and $d_z < 0.2$ cm, whereas hadronic taus must have $d_z < 0.2$ cm.

5.2.1 Event selection in the $\tau_\mu\tau_h$ channel

The events were selected online using a combination of the single-muon and muon+tau triggers. The offline selection requires an oppositely charged muon and hadronic tau pair which are separated by $\Delta R > 0.5$. These offline objects must also be matched to the corresponding trigger objects with $\Delta R < 0.5$. The selected muon must pass the *medium* muon ID, have $I_{\text{rel}} < 0.15$, and $|\eta| < 2.1$. The muon is required have p_T 1 GeV higher than the online selection. For the single-muon triggers this corresponds to a cut of 23 (25) GeV in 2016 (2017), and for the muon+tau triggers the cut is 20 (21) GeV in 2016 (2017).

The hadronic tau must be reconstructed by the HPS algorithm, pass the *tight* MVA-isolation, the *very-loose* working point of anti-electron discriminant, and the *tight* working point of the anti-muon discriminant. The τ_h candidate must additionally have $|\eta| < 2.3$ and $p_T > 30$ GeV.

Even after requiring the τ_h candidate to pass the *tight* MVA-isolation, there still remains a significant contribution to the background from W+jets events with a jet $\rightarrow \tau_h$ fake. As illustrated in Figure 3.23, the $m_T(l, E_T^{\text{miss}})$ variable gives good separation between W+jets events and other processes. The $m_T(\mu, E_T^{\text{miss}})$ is therefore required to be less than 50 GeV to suppress this background. There is also a smaller but non-negligible background in this channel from $t\bar{t}$, which is reduced by vetoing events with b-jets.

5.2.2 Event selection in the $\tau_e\tau_h$ channel

In 2016 the events were selected online using a single-electron trigger. For 2017 an electron+tau trigger is also used in combination with the single-electron trigger. The offline selection requires an oppositely charged electron and hadronic tau pair which are separated by $\Delta R > 0.5$. These offline objects must also be matched to the corresponding trigger objects with $\Delta R < 0.5$. In 2016 (2017) the electron must pass the *tight* (*loose*) working-point of the electron ID, have $I_{\text{rel}} < 0.10$ ($I_{\text{rel}} < 0.15$), and $|\eta| < 2.1$. The electron is required have p_T 1 GeV higher than the online selection. For the single-electron triggers this corresponds to a cut of 25 (28) GeV in 2016 (2017), and for the electron+tau trigger used in 2017 the cut is 25 GeV.

The hadronic tau must be reconstructed by the HPS algorithm, pass the *tight* MVA-isolation, the *tight* working point of anti-electron discriminant, and the *loose* working point of the anti-muon discriminant. The τ_h candidate must also have $|\eta| < 2.3$ and $p_T > 30$ GeV. Hadronic taus passing the tau leg of the electron+tau cross trigger in 2017 are additionally required to have $p_T > 35$ GeV.

As for the $\tau_\mu\tau_h$ channel the W+jets and $t\bar{t}$ backgrounds are reduced by requiring $m_T(e, E_T^{\text{miss}}) < 50$ GeV and vetoing events with b-jets.

5.2.3 Event selection in the $\tau_e\tau_\mu$ channel

The events were selected online using the two electron+muon triggers described in Section 3.2. The offline selection requires an oppositely charged electron and muon pair which are separated by $\Delta R > 0.3$. The offline leptons must also be matched to the corresponding trigger objects with

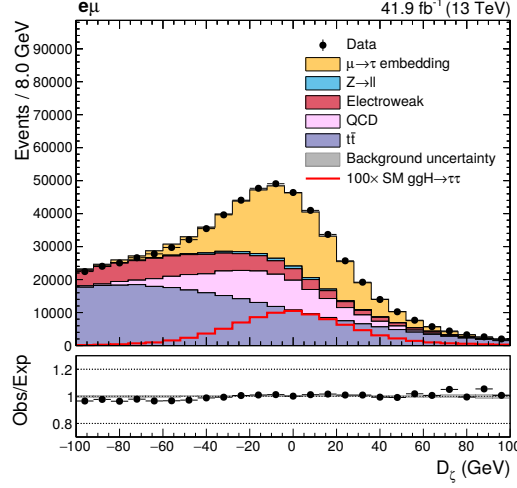


Figure 5.7: The D_ζ distribution in the $\tau_e \tau_\mu$ final state. The GF Higgs production process shown by the solid red line is scaled by a factor of 100. The uncertainty band shown in shaded grey includes only the statistical uncertainties.

$\Delta R < 0.5$ The electron is required to have $p_T > 13$ GeV, $I_{\text{rel}} < 0.15$, and pass the *tight* (*loose*) working point of the electron ID in 2016 (2017). The muon is required to have $p_T > 10$ GeV, $I_{\text{rel}} < 0.2$, and pass the *medium* muon ID. Because the two electron+muon triggers apply different p_T cuts on the leptons the electron/muon that is matched to the high p_T leg of each trigger must pass an additional offline p_T cut of 24 GeV.

The background from $t\bar{t}$ events is large in this channel, especially in regions of phase-space which contain multiple jets - exactly the region in which we measure $\mathcal{CP}$. Two additional requirements are therefore used to reduce this background. Firstly, events with b-jets are vetoed. Secondly, a cut is applied on the D_ζ variable [180] which is defined as

$$D_\zeta = \vec{E}_T^{\text{miss}} \cdot \hat{\zeta} - 0.85 \cdot (\vec{p}_T^e + \vec{p}_T^\mu) \cdot \hat{\zeta}, \quad (5.9)$$

where $\hat{\zeta}$ is the unit vector that bisects the transverse directions of the visible decay products, $\vec{p}_T^e$ and $\vec{p}_T^\mu$. The D_ζ variable separates processes producing neutrinos from tau decays, which tend to be aligned with the visible decay products, from processes where the neutrinos are back-to-back with the visible decay products such as $t\bar{t}$ and W +jets. As illustrated in Figure 5.7, processes with genuine tau leptons such as $Z/\gamma^* \rightarrow \tau\tau$ and $H \rightarrow \tau\tau$ tend to give positive values of D_ζ , whereas backgrounds such as $t\bar{t}$ tend to have negative values. D_ζ is required to be at least -35 GeV, although this cut is tightened in some categories where the $t\bar{t}$ background is larger as will be described in Section 5.3.

5.2.4 Event selection in the $\tau_h \tau_h$ channel

The events were selected online using the double-tau triggers. The offline selection requires two oppositely charged τ_h candidates which are separated by $\Delta R > 0.5$. Both τ_h candidates must pass the HPS algorithm, the *tight* MVA-isolation, the *very-loose* anti-electron discriminant, and the *loose* anti-muon discriminant. The leading τ_h is required to have $p_T > 50$ GeV and the sub-leading τ_h must have $p_T > 40$ GeV, where the latter threshold is driven by the trigger and the former is employed to reduce $\text{jet} \rightarrow \tau_h$ fakes.

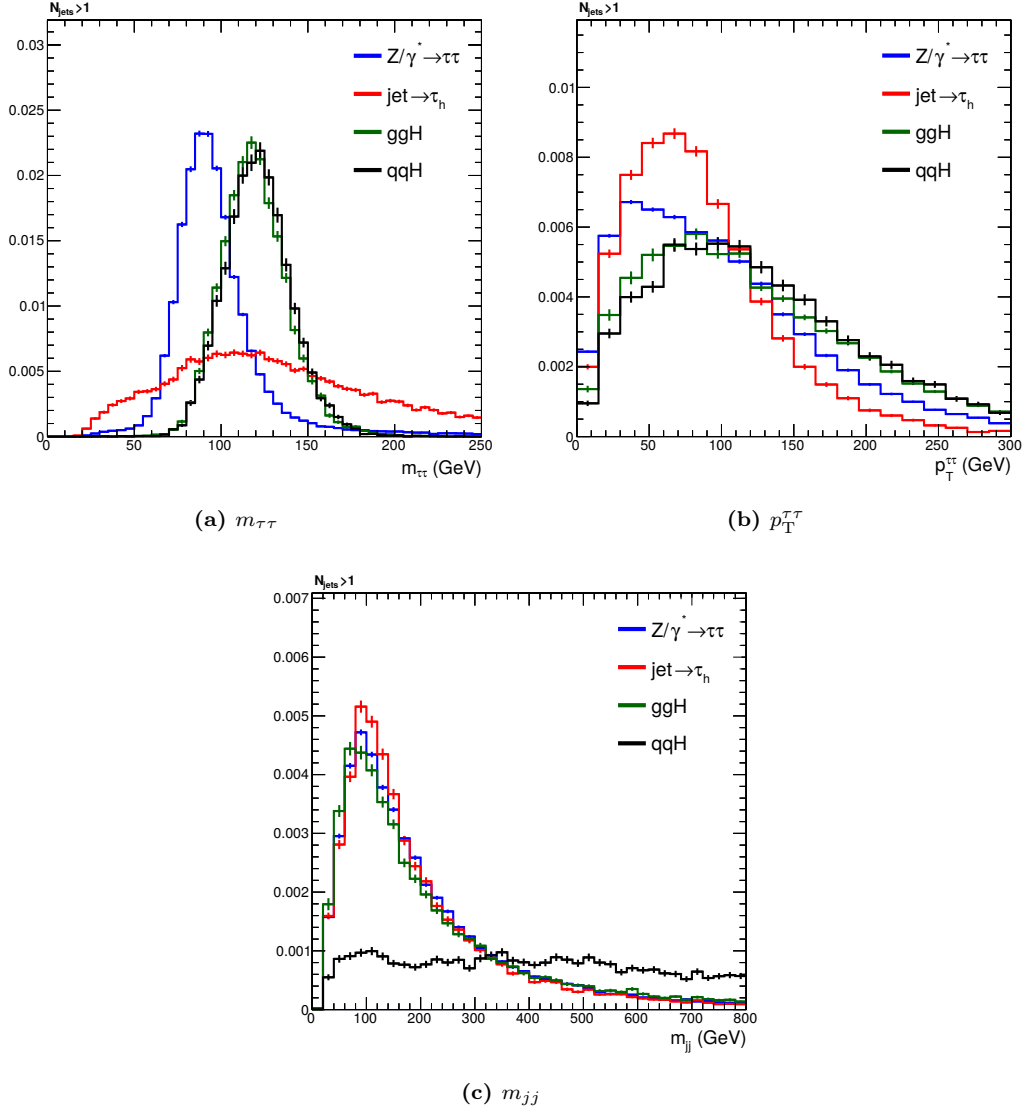


Figure 5.8: Comparisons of the normalised $m_{\tau\tau}$ (a), $p_T^{\tau\tau}$ (b), and m_{jj} (c) distributions for the $Z/\gamma^* \rightarrow \tau\tau$, $\text{jet} \rightarrow \tau_h$, $\text{GF H} \rightarrow \tau\tau$ ("ggH"), and $\text{VBF H} \rightarrow \tau\tau$ ("qqH") processes in the $\tau_\mu\tau_h$ final state.

5.2.5 Lepton vetoes and pair selection

The lepton vetoes and pair selection procedure described in Section 4.2.4 are also applied. This reduces the $Z/\gamma^* \rightarrow ll$ and diboson backgrounds and ensures that all channels are orthogonal.

5.3 Event categorisation

The event categorisation is loosely based on the categories used in the CMS $\text{H} \rightarrow \tau\tau$ discovery analysis described in Ref. [2]. For this analysis events were split into three categories: 0-jet, boosted, and VBF. The 0-jet category contains all events with $N_{\text{jets}} = 0$. The VBF category requires $N_{\text{jets}} \geq 2$, with additional selections on m_{jj} , $\Delta\eta_{jj}$, $p_T^{\tau\tau}$, and/or $p_T^{\tau_h}$ applied depending on the channel. The boosted category then contains all remaining events which are predominantly $N_{\text{jets}} = 1$ events, with the exception of some events with $N_{\text{jets}} \geq 2$ failing the other VBF category selections.

As the name suggests, the VBF category was specifically designed to be sensitive to VBF Higgs production and is therefore sub-optimal for GF Higgs production of interest here. This category was therefore suitably modified and renamed as the "dijet" category as described below. The 0-jet and boosted categories on the other hand were optimised to give sensitivity for GF and are therefore largely unchanged for this analysis.

Before discussing the procedures used to optimise the category selections the most powerful variables for signal and background separation will first be reviewed. The main backgrounds in the analysis are $Z/\gamma^* \rightarrow \tau\tau$, processes producing $\text{jet} \rightarrow \tau_h$, VBF Higgs production, and $t\bar{t}$. The variable that gives the best separation between Higgs and background events is the ditau mass, $m_{\tau\tau}$. The ditau mass is reconstructed using the **SVfit** algorithm [181], which combines the 4-vectors of the visible decay products with the $\vec{E}_T^{\text{miss}}$ using a Dynamical Likelihood Method [182, 183]. This method gives a more accurate estimation of the parent boson mass compared to using the 4-vectors of the visible decay products alone. The $m_{\tau\tau}$ resolution is about 15%, 20%, and 25% for $\tau_h\tau_h$, $\tau_l\tau_h$, and $\tau_e\tau_\mu$ events respectively, where the resolution worsens for the events with leptonic tau decays due to the presence of additional neutrinos in the final state. As illustrated in Figure 5.8 (a), the $m_{\tau\tau}$ distribution for the $H \rightarrow \tau\tau$ and $Z/\gamma^* \rightarrow \tau\tau$ processes forms a peaked structure that is maximal at around 125 GeV and 90 GeV respectively. For the other background processes no distinct peak is observed, and events are distributed according to a falling distribution. The $m_{\tau\tau}$ tends to be the most powerful variable for separating signal and background and is therefore used as a discriminating variable rather than to define categories.

Another useful variable to reduce the background is the reconstructed Higgs p_T , $p_T^{\tau\tau}$, which is computed as the vector sum of the transverse momenta of the visible tau decay products and the E_T^{miss} . As illustrated in Figure 5.8 (b), the $p_T^{\tau\tau}$ tends to be softer for the non-Higgs background processes compared to the signal.

The m_{jj} variable gives good separation between VBF and GF, where the former tends to take larger values. The m_{jj} distributions are shown in Figure 5.8 (c). At first glance it looks like applying a cut on this variable is a sensible choice to reduce the VBF background, however, an additional complication is that the high m_{jj} GF events are more sensitive to the Higgs $\mathcal{CP}$ -state (as discussed in Section 5.1.0.1). This deters the use of hard cuts on m_{jj} to reject VBF events.

The dijet category requires at least two jets in the event. Of the events passing this selection the ones with low $m_{jj} < 300$ GeV are found to have almost no $\mathcal{CP}$ sensitivity and so a minimum cut of $m_{jj} > 300$ GeV is applied. These low m_{jj} events aren't discarded but are instead included in the boosted category. The relative sensitivity after applying different cuts on $p_T^{\tau\tau}$ and m_{jj} can be estimated by comparing the $\frac{S}{\sqrt{B}}$. A range of cuts is scanned with the $\frac{S}{\sqrt{B}}$ calculated for each point, where the maximum of the scan indicates the optimal cut for separating signal and background. Figure 5.9 shows these scans for $p_T^{\tau\tau}$ and m_{jj} . The lower insets of Figure 5.9 also show the ratio of GF to VBF events as a function of the applied cut. The m_{jj} scan shows that in terms of separating GF from the backgrounds m_{jj} is not a useful variable, as the maximum of the scan is found at 0 GeV. The fraction of VBF events does however increase as a function of m_{jj} . The $p_T^{\tau\tau}$ on the other hand shows that a $p_T^{\tau\tau}$ cut around 100 GeV can help slightly improve the sensitivity, but this improvement is only $\mathcal{O}(5\%)$. This is because there is a sizeable fraction of low $p_T^{\tau\tau}$ GF events, even though on average they tend to be more boosted than the background.

Rather than apply a hard cut on $p_T^{\tau\tau}$, the use of categories may be more optimal as it does not discard the low $p_T^{\tau\tau}$ events. The use of two $p_T^{\tau\tau}$ based categories is considered: one containing

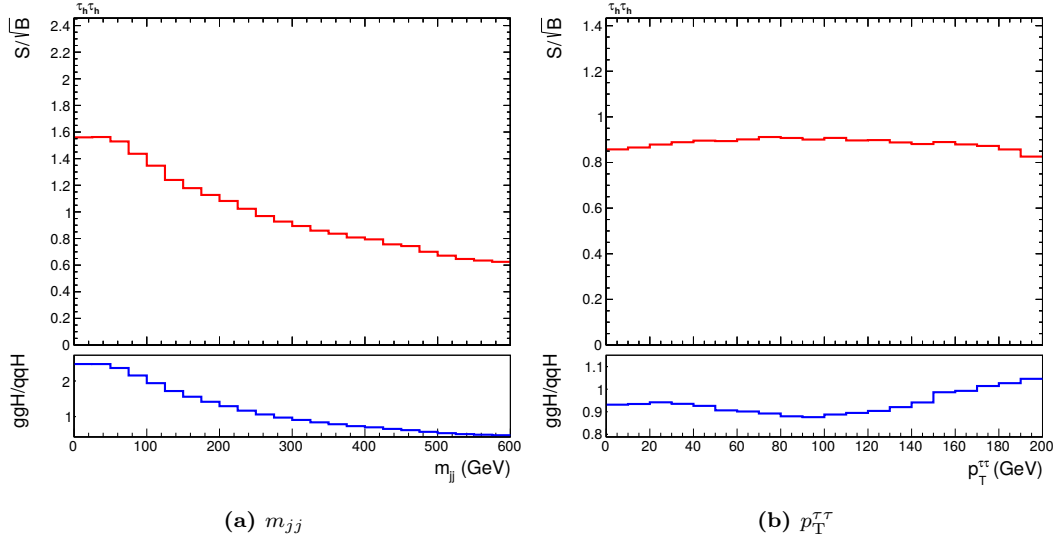


Figure 5.9: $\frac{S}{\sqrt{B}}$ scans of m_{jj} (left) and $p_T^{\tau\tau}$ (right) in the $\tau_h\tau_h$ channel. For both scans at least two jets are required and a cut of $100 < m_{\tau\tau} < 150$ GeV is applied. For the $p_T^{\tau\tau}$ scan an additional cut of $m_{jj} > 300$ GeV is also applied. The lower inset also shows the ratio of the GF yield to the VBF yield.

all events with $p_T^{\tau\tau} < X$, and the other with all events with $p_T^{\tau\tau} > X$. A range of X values is scanned to find the most optimal splitting for the categories. For this test the asymptotic median significance (AMS) [174], which is defined as

$$\text{AMS} = \sqrt{2 \left((S + B) \cdot \log \left(1 + \frac{S}{B} \right) - S \right)}, \quad (5.10)$$

is used as the figure of merit. This variable tends to give a better estimate of the sensitivity in the low statistics limit compared to $\frac{S}{\sqrt{B}}$. The AMS is calculated for the two categories and then added in quadrature to compute the "total" AMS. Figure 5.10 shows these scans for each of the four channels. The optimum $p_T^{\tau\tau}$ splitting is found to be similar for each channel and around 150–200 GeV. The improvement in the sensitivity between using one inclusive category and using two $p_T^{\tau\tau}$ is $\mathcal{O}(10 - 20\%)$.

The use of the m_{jj} variable is more complicated as it does not bring any improvement in the signal vs background differentiation but the high m_{jj} events are more sensitive to $\mathcal{CP}$. To estimate the sensitivity in this case we run the full analysis chain and fit for α_{hgg} using Asimov data (this procedure will be described in the following sections). The pure $\mathcal{CP}$ -even vs pure $\mathcal{CP}$ -odd significance is used as the figure of merit. Three options are considered:

- Option 1: Apply a hard cut on the $m_{jj} = [300, 400, 500]$ GeV and use no $p_T^{\tau\tau}$ categories.
- Option 2: Apply a hard cut on the $m_{jj} = [300, 400, 500]$ GeV and use $p_T^{\tau\tau}$ categories with $p_T^{\tau\tau}$ splitting = $[100, 150, 200]$ GeV.
- Option 3: Categorise events based on both $p_T^{\tau\tau}$ and m_{jj} giving four categories in total. The $p_T^{\tau\tau}$ splittings tested are $[100, 150, 200]$ GeV, and the m_{jj} categories are defined as $300 < m_{jj} < X$ GeV and $m_{jj} > X$, where $X = [400, 500, 600]$ GeV.

The values in square brackets indicate the values of the cuts/category-splits that are tested. For practical reasons the number of such cuts/splits are kept minimal, and the $p_T^{\tau\tau}$ splits are chosen

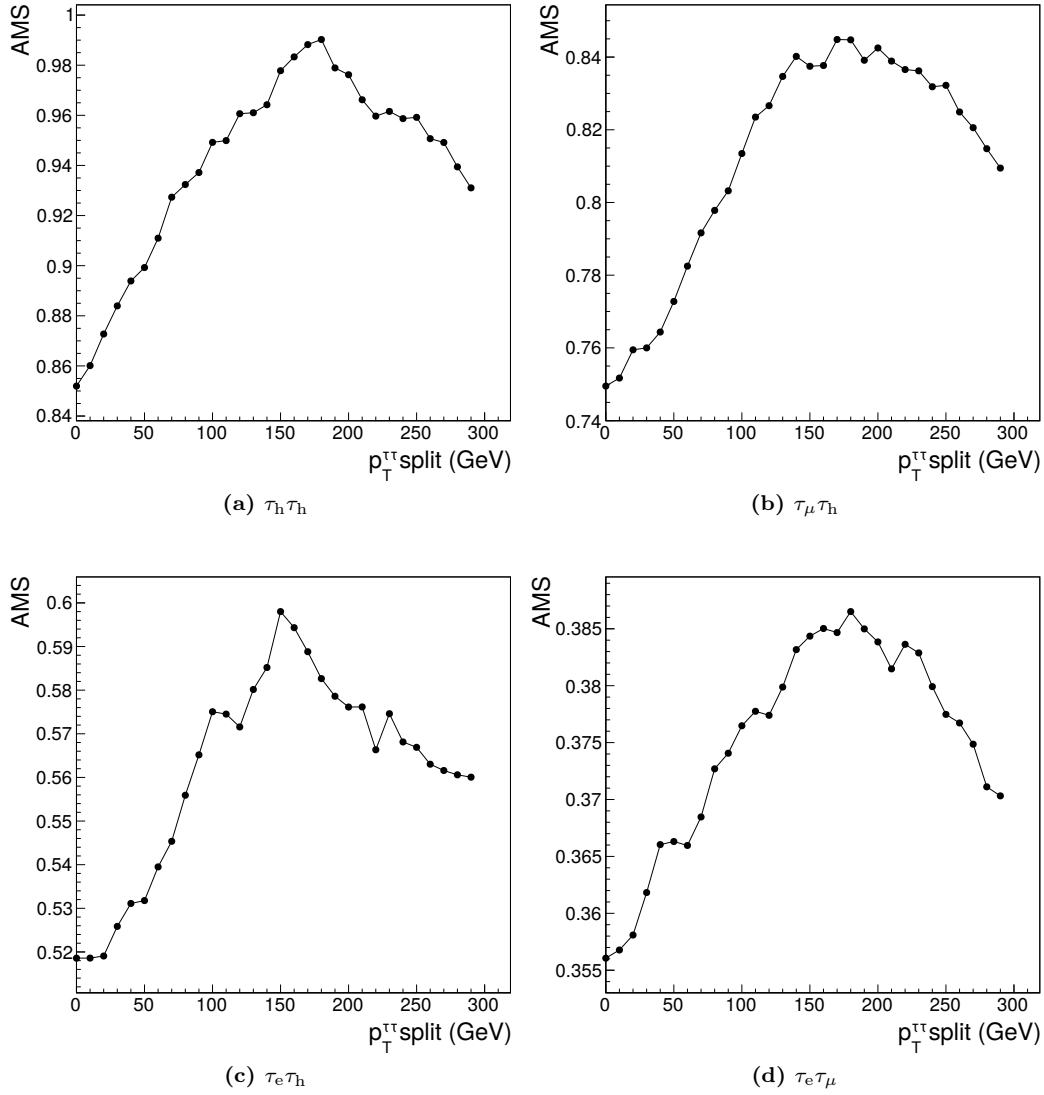


Figure 5.10: Scans of the combined AMS for different $p_T^{\tau\tau}$ category definitions. Each value of $p_T^{\tau\tau}$ shown on the x-axis indicates the split into two categories with $p_T^{\tau\tau} < X$ and $p_T^{\tau\tau} \geq X$. The AMS values computed for each category are added in quadrature to compute the total AMS which is shown on the y-axis.

based on the AMS study described above. The best results for each option are summarised below:

- Option 1: The optimal m_{jj} cut is found to be 300 GeV. This gives an expected significance of 0.92σ .
- Option 2: The optimal $p_T^{\tau\tau}$ category splitting and m_{jj} cut are found to be 200 and 300 GeV respectively. This gives an expected significance of 1.06σ .
- Option 3: The optimal $p_T^{\tau\tau}$ and m_{jj} category-splits are found to be 150 and 500 GeV respectively. This gives an expected significance of 1.26σ .

These studies show that using a combination of both $p_T^{\tau\tau}$ and m_{jj} categories gives the best sensitivity to α_{hgg} . The use of the m_{jj} based categories is observed to give $\mathcal{O}(20\%)$ improvement over applying only a hard m_{jj} cut. This categorisation is therefore adopted for the analysis.

The dijet categorisation is summarised as follows:

- Dijet events are selected by requiring $m_{jj} > 300$ GeV and $N_{\text{jets}} \geq 2$.
- The dijet events are split into "loose- m_{jj} " and "tight- m_{jj} " categories, which contain all events with $m_{jj} < 500$ GeV and $m_{jj} > 500$ GeV respectively.
- These categories are further split into two $p_T^{\tau\tau}$ categories containing events with $p_T^{\tau\tau} < 150$ GeV and $p_T^{\tau\tau} > 150$ GeV. The latter are additionally named "boosted" categories.

This gives four categories in total: "loose- m_{jj} ", "tight- m_{jj} ", "loose- m_{jj} boosted", and "tight- m_{jj} boosted". For the $\tau_e\tau_\mu$ channel an additional cut of $D_\zeta > -10$ GeV is used to suppress the large $t\bar{t}$ background in the dijet categories. Figure 5.11 gives a graphical representation of all the categories used in the analysis.

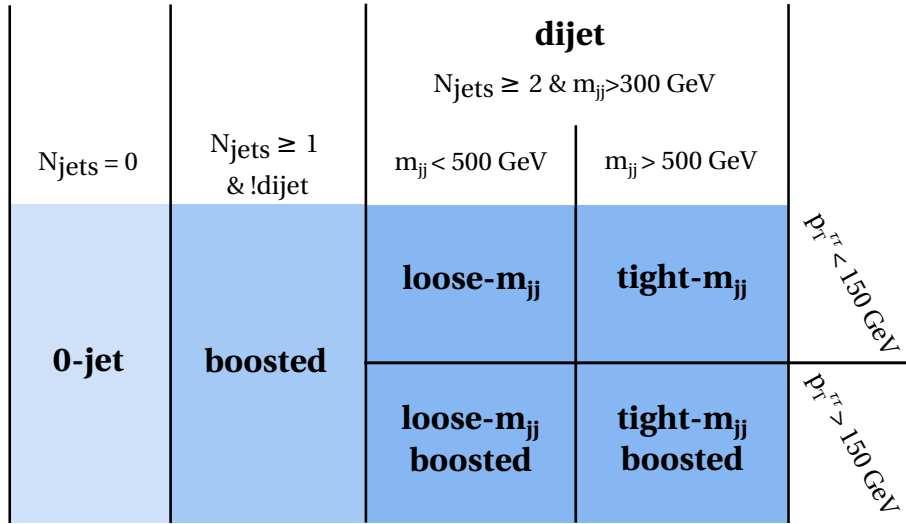


Figure 5.11: Diagram illustrating the categorisation of events.

5.4 Discriminating variables

The analysis uses a combination of one dimensional and two dimension discriminants to extract the signal and measure the Higgs $\mathcal{CP}$ properties. The 2D observables are "unrolled", meaning the bins in the 2D histograms are mapped onto a 1D histogram. When performing the fit to data to extract the results these "pseudo"-1D histograms are treated in exactly the same way as a regular 1D discriminant. Figure 5.12 illustrates the unrolling procedure. The choice of discriminant depends on the category as described below.

The 0-jet category uses the $m_{\tau\tau}$ as a 1D observable. This allows for differentiation between the Higgs signal, where most events typically have $100 < m_{\tau\tau} < 150$ GeV, and background processes. The background distributions tend to either peak at lower values of $m_{\tau\tau}$ (e.g $Z/\gamma^* \rightarrow \tau\tau$), or form more of a falling continuum rather than a distinct peak, with more events in the high $m_{\tau\tau}$ tail (e.g $\text{jet} \rightarrow \tau_h$).

For the boosted category, a 2D observable using $m_{\tau\tau}$ and $p_T^{\tau\tau}$ as the two dimensions is utilised. The $m_{\tau\tau}$ dimension provides signal and background differentiation as described above. The inclusion of the $p_T^{\tau\tau}$ dimension also provides additional discriminating power against background processes, which tend to have low $p_T^{\tau\tau}$ compared to the signal.

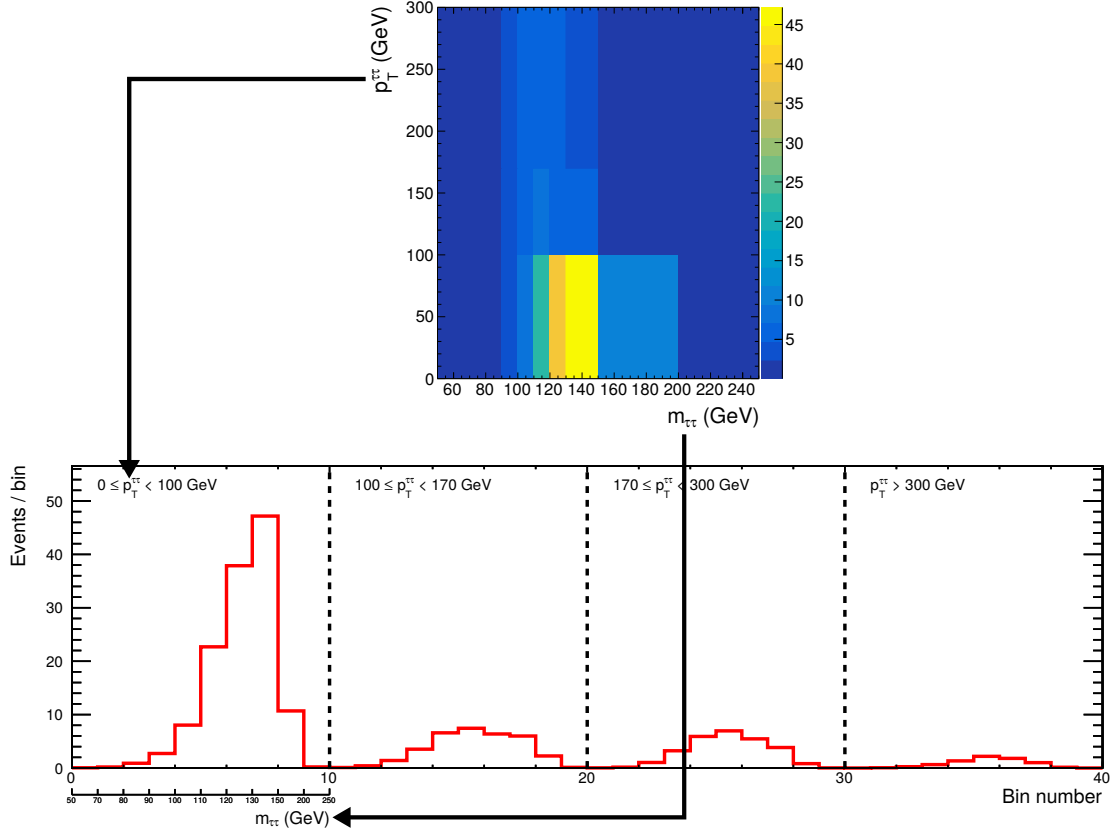


Figure 5.12: Illustration of the unrolling procedure whereby a 2D histogram is mapped onto a 1D histogram. In the unrolled plot the dashed lines and upper labels indicate the y-axis variable and binning. The small axis to the bottom left of plot indicates the x-axis variable and binning.

The discriminant used for the dijet categories needs to be able to separate the Higgs signal from the backgrounds and differentiate between different Higgs $\mathcal{CP}$ states. For these reasons a 2D discriminate is formed using $m_{\tau\tau}$ and $\Delta\phi_{jj}$ as the two dimensions. The former is used to differentiate signal from background, whereas the latter is included primarily to measure the Higgs $\mathcal{CP}$ properties - but can also provide some additional differentiation between signal and background.

5.5 Background methods

The non-Higgs backgrounds that are considered are Drell-Yan $Z/\gamma^* + \text{jets}$, QCD-induced $W + \text{jets}$, electroweak production of $W + \text{jets}$ and Z/γ^* , QCD-multijet, $t\bar{t}$, diboson, and single top.

The choice of background methods used for the analysis is driven by the need to achieve a good description of the $\Delta\phi_{jj}$ variable used for signal extraction, and to maximise the statistical power of the background templates. The latter is important as the dijet selections needed to isolate $\mathcal{CP}$ -sensitive events reduce the background yields significantly. This motivates the use of the data-driven background estimation methods described in Section 3.5.

All processes where both the $\tau_{l/h}$ candidates originate from the decays of real tau leptons are modelled by the embedding method as described in Section 3.5.2. This includes $Z/\gamma^* \rightarrow \tau\tau$, $t\bar{t} \rightarrow \tau\tau$, diboson $\rightarrow \tau\tau$, and single top production in the tW channel. In the $\tau_h\tau_h$ and $\tau_l\tau_h$ channels backgrounds in which any tau candidate originates from a $\text{jet} \rightarrow \tau_h$ fake are modelled by the FF

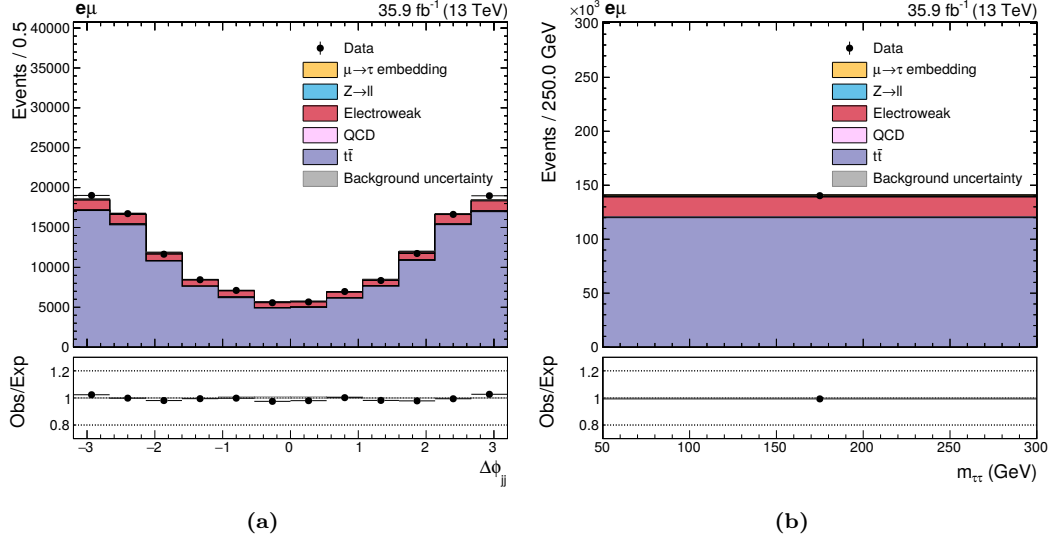


Figure 5.13: Distributions illustrating the modelling of the $t\bar{t}$ background in a control region. The control region is selected by requiring $D_\zeta < -50$ GeV and $N_{jets} \geq 1$. The left inset shows the data/background agreement for the $\Delta\phi_{jj}$ variable, where the additional requirement $N_{jets} \geq 2$ is imposed. The right inset shows the 1-bin control region that is included in the final fit to constrain the $t\bar{t}$ background. The uncertainty band shown in shaded grey includes only statistical uncertainties.

method as described in Section 3.5.3. This includes mainly QCD and W+jets events but has also a smaller contribution from the other background processes such as $t\bar{t}$. In the $\tau_e\tau_\mu$ channel the QCD is modelled using the method described in Section 3.5.5. Altogether the use of these methods means that $\approx 90\%$ of the background events are modelled from data-driven techniques.

The remaining processes are modelled from simulation. The largest background this includes is $t\bar{t}$, which is a minor contribution in the $\tau_l\tau_h$ and $\tau_h\tau_h$ channels but can be large in the $\tau_e\tau_\mu$ channel. The modelling of the $t\bar{t}$ background is validated in a control region enriched in $t\bar{t}$ events. This control region is selected in the $\tau_e\tau_\mu$ final state by requiring $D_\zeta < -50$ GeV and $N_{jets} \geq 1$. The agreement with the data is illustrated in Figure 5.13. The left inset shows the distribution of the $\Delta\phi_{jj}$ variable where excellent agreement with the data is observed. In the right inset the control region events have been grouped into a single $m_{\tau\tau}$ bin ($50 < m_{\tau\tau} < 300$ GeV) to demonstrate the agreement of the yield with the data. This one bin control region is included in the final fit to constrain the normalisation of the $t\bar{t}$ background.

5.5.1 Generator matching

It is necessary to split the MC processes depending on simulated objects from which the τ_l and τ_h candidates originate. This ensures that backgrounds that are already covered by the embedding or FF methods can be removed and are therefore not counted twice. The generator matching procedure is described in Section 4.4.1.

5.5.2 Data vs background comparisons

The description of the data by the chosen background methods is illustrated in Figures 5.14–5.15, which show the pre-fit distributions of the $m_{\tau\tau}$ and $\Delta\phi_{jj}$ variables. These variables are important for signal extraction and measurement of the Higgs $\mathcal{CP}$ -state. The uncertainty band shown in

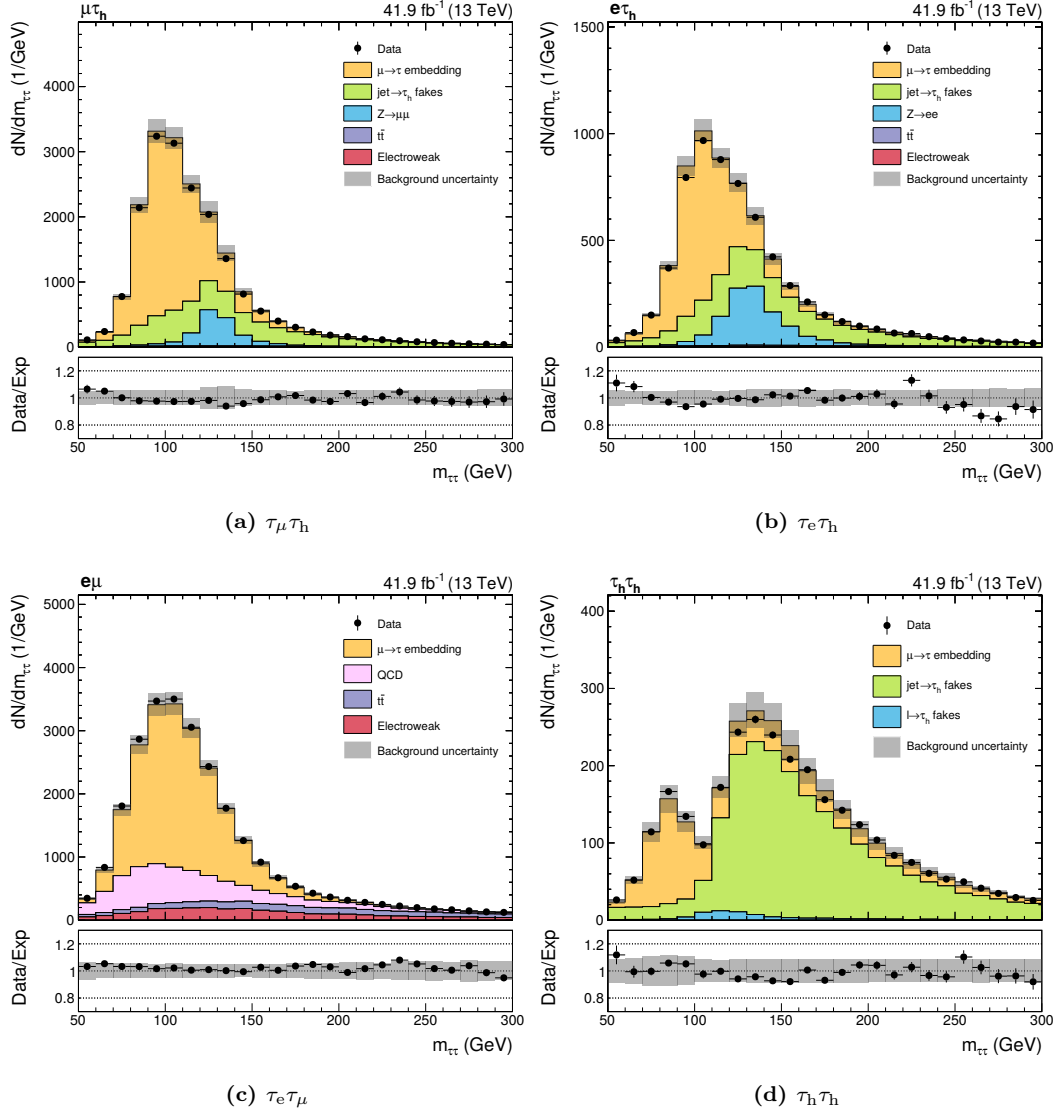


Figure 5.14: Distributions of $m_{\tau\tau}$ in the $\tau_\mu\tau_h$, $\tau_e\tau_h$, $\tau_e\tau_\mu$, and $\tau_h\tau_h$ final states. The uncertainty band shown in shaded grey includes the full uncertainty model.

shaded grey includes the full systematic uncertainty model. Good agreement with the data is observed.

5.5.3 Modelling of Higgs backgrounds

The background due to non-GF induced $H \rightarrow \tau\tau$ events includes the VBF, W-associated (VH), and Z-associated (ZH) production processes. In principle a non-SM $\mathcal{CP}$ -state would result in different cross sections and kinematic distributions for these processes. However, the lack of a tree-level $\mathcal{CP}$ -odd coupling to WW or ZZ suppresses the effect of the $\mathcal{CP}$ -odd component, provided the loop-induced effective coupling is small compared to the tree-level $\mathcal{CP}$ -even coupling. In such scenarios the kinematic distributions would resemble almost exactly the SM distributions even if the $\mathcal{CP}$ -odd component was large, and different $\mathcal{CP}$ -states would be manifest only as modifications to the cross sections and branching ratios [52]. This is the case in most weakly interacting models, as the production rates for $\mathcal{CP}$ -odd scalars are typically suppressed by three orders of magnitude or

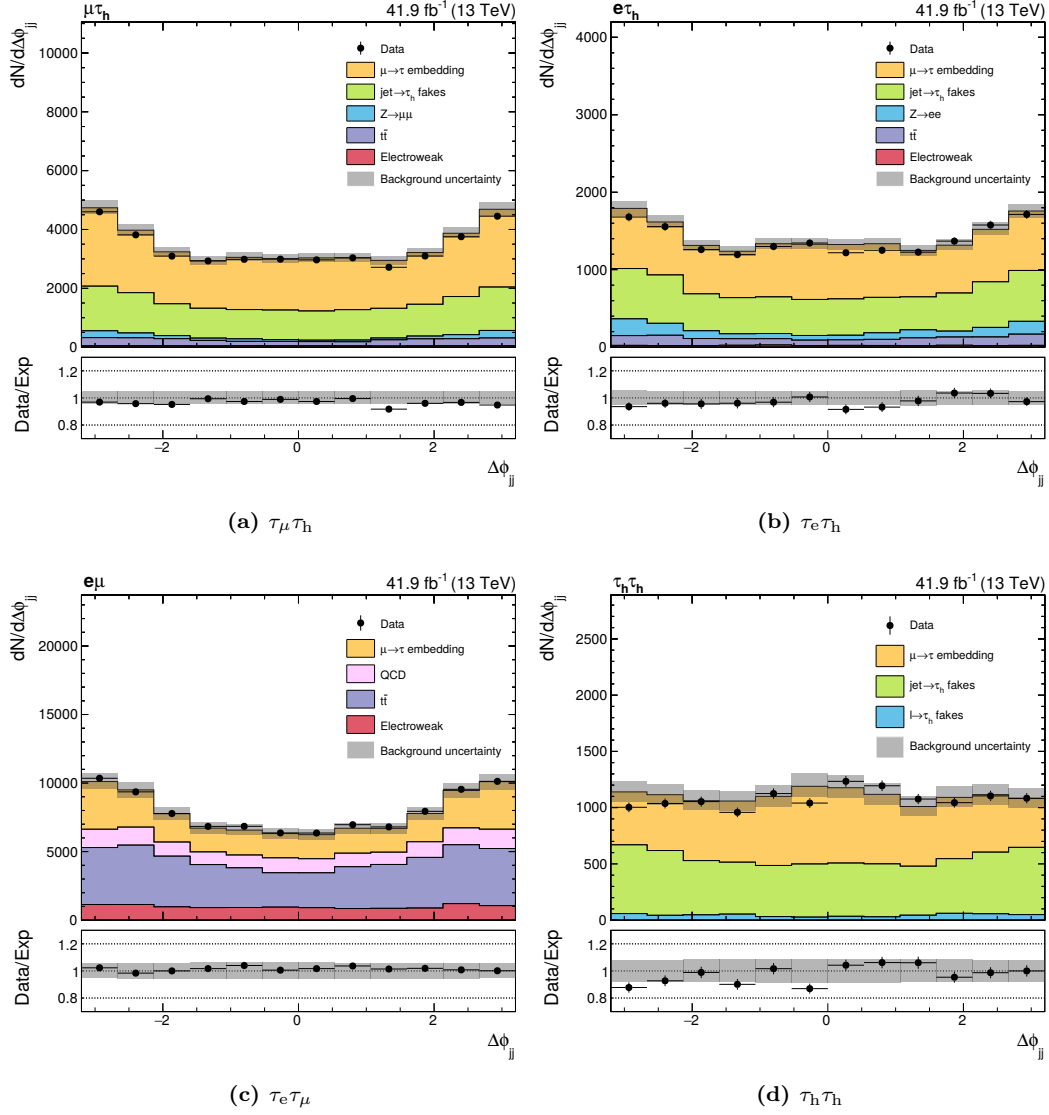


Figure 5.15: Distributions of $\Delta\phi_{jj}$ in the $\tau_\mu\tau_h$, $\tau_e\tau_h$, $\tau_e\tau_\mu$, and $\tau_h\tau_h$ final states. The uncertainty band shown in shaded grey includes the full uncertainty model.

more [52]. For this reason, the VBF and VH kinematic distributions are taken to be the same as the SM predictions in the analysis. The VBF and VH cross sections are allowed to float freely in the fit using a single parameter μ_V , thus assuming $\mu_{\text{VBF}} = \mu_{\text{WH}} = \mu_{\text{ZH}} = \mu_V$. The VBF, WH, and ZH MC samples are generated at NLO accuracy using POWHEG.

An alternative treatment of the VH and VBF backgrounds is also used to cross-check the assumptions described above. Adopting the parametrisations in Ref. [50] a generic VH/VBF state can be described by the form factors: a_1 , a_2^{WW} , a_2^{ZZ} , a_3^{WW} , and a_3^{ZZ} ; where a_1 , a_2^i , and a_3^i describe the SM coupling, the anomalous $\mathcal{CP}$ -even couplings, and the anomalous $\mathcal{CP}$ -odd couplings respectively. MC samples are generated using JHUGen 7.0.2 [184–187] at LO accuracy, neglecting the effects of $\mathcal{CP}$ -even anomalous HVV couplings ($a_2^i = 0$), and assuming the same $\mathcal{CP}$ -odd anomalous coupling of Higgs boson to ZZ and WW ($a_3^{\text{WW}} = a_3^{\text{ZZ}} = a_3$). As for the GF modelling, only three samples are required for each process: a pure $\mathcal{CP}$ -even sample, a pure $\mathcal{CP}$ -odd sample, and a maximum-mixing sample. These samples are then combined to form distributions for any generic

$\mathcal{CP}$ -state using an analogous formula to Equation 5.5 except replacing μ_{ggh} and α_{hgg} with μ_V and α_{hvv} respectively. The α_{hvv} mixing angle in this case is defined as

$$\tan \alpha_{\text{hvv}} = \sqrt{\frac{\sigma_3}{\sigma_{\text{SM}}}} \cdot \frac{a_3}{a_1}. \quad (5.11)$$

The cross check is performed for two scenarios:

1. With α_{hvv} allowed to float completely freely in the fit ($-90^\circ < \alpha_{\text{hvv}} \leq 90^\circ$). This effectively allows for so-called strongly interacting models [52], where the size of the loop-induced $\mathcal{CP}$ -odd component can be comparable to the $\mathcal{CP}$ -even contribution. However, it should be noted that these models also predict light electroweakly interacting states, and are therefore disfavoured by the lack of observations of such particles at the LHC [188].
2. With α_{hvv} limited to a small range to reflect the $\mathcal{O}(10^{-3})$ suppression in the $\mathcal{CP}$ -odd production rate for weakly interacting models. The range is given by $\tan \alpha_{\text{hvv}} \approx \sqrt{10^{-3}} \rightarrow -2^\circ \leq \alpha_{\text{hvv}} \leq 2^\circ$.

5.6 Systematic uncertainties

Two types of systematic uncertainties are considered in the analysis: normalisation uncertainties and shape uncertainties. The normalisation uncertainties affect only the yields of the signal and background templates, whereas shape uncertainties affect both the normalisations and shapes of the discriminating variable distributions. A summary of the systematic uncertainties is given in this section. This is followed by a discussion of the treatment of correlation between the MC and embedding uncertainties, and the correlations between the uncertainties in 2016 and 2017.

5.6.1 Normalisation uncertainties

The uncertainties that affect only the normalisation of the background and signal processes are described below, with the exception of theoretical uncertainties on the signal normalisation which are described in Section 5.1.1.

- **Luminosity**

The luminosity uncertainty amounts to 2.5% on the data collected during 2016 data-taking and 2.3% on the data collected in 2017. The luminosity uncertainty is applied to all backgrounds and signal processes where the yield is estimated from simulation.

- **Electron, muon and tau ID efficiency**

The combined uncertainties on electron and muon identification and isolation efficiencies are 2%. For hadronic taus the uncertainty on the tau identification efficiency is 5% (3%) in 2016 (2017). This is broken down into a part that is correlated between channels with magnitude of 4.5% (2.7%) and a part that is uncorrelated between channels with a magnitude of 2% (1.2%) for 2016 (2017). The uncorrelated part accounts for the different anti-muon/electron discriminator working-points used for each channel. Since two hadronic taus are required in the $\tau_h \tau_h$ channel the magnitude of correlated and uncorrelated uncertainty is 9% (5.5%) and 4% (2.4%) respectively for 2016 (2017). The embedded samples also have an additional uncertainty applied to cover the embedding $\rightarrow$ MC tracking efficiency correction described in Section 3.4.7 with magnitude of 2% (per-tau).

- **Trigger efficiency**

The uncertainty on the trigger efficiencies amounts to 2% for the $\tau_e \tau_h$ and $\tau_\mu \tau_h$ channels, 4% in the $\tau_e \tau_\mu$ channel, and 10% for the $\tau_h \tau_h$ channel.

- **B-tagging efficiency**

For the $\tau_e\tau_h$, $\tau_\mu\tau_h$ and $\tau_e\tau_\mu$ channels, which contain a veto on events containing b-jets, an uncertainty on the $t\bar{t}$, diboson, and single-top yields due to the b-tagging efficiency is applied. The uncertainties on the b-tagging efficiency corrections are given as a function of p_T and η . These are propagated by varying the b-tagging scale factors within their uncertainties. The magnitude of this uncertainty ranges from 1–9%. The uncertainties on the event yields for other processes are found to be negligible ($\approx 0.1\%$).

- **Background normalisation uncertainties**

- The uncertainty on the normalisation of the embedded samples amounts to 4%. This originates from the uncertainty on the measured muon identification and trigger efficiencies used to scale the embedded samples.
- A 4% uncertainty is applied to the $Z/\gamma^* \rightarrow ll$ process to account for the uncertainty on the Drell-Yan cross section.
- A 6% uncertainty is applied to the $t\bar{t}$ process to account for the uncertainty on the cross section.
- A 5% uncertainty is applied to the diboson and single-top processes to account for the uncertainties on the production cross sections.
- The uncertainties on the VBF Higgs production cross section recommended in Ref. [30] are applied. This includes a 0.4% uncertainty from variations of QCD scales and a 2.1% uncertainty from PDFs+ α_s variations. The uncertainties on the $H \rightarrow \tau\tau$ branching ratio described in Section 5.1.1 are also applied.
- The uncertainties on the W(Z)-associated Higgs production cross sections recommended in Ref. [30] are applied. This includes a 0.7% (3.8%) uncertainty from the QCD scale variations and a 1.9% (1.6%) uncertainty from PDFs+ α_s variations. The uncertainties on the $H \rightarrow \tau\tau$ branching ratio described in Section 5.1.1 are also applied.
- For the QCD background in the $\tau_e\tau_\mu$ channel, the uncertainty on the extrapolation from the anti-isolated side-band region used to measure the $\dot{R}_{\text{QCD}}$ ratios to the signal region is found to have a negligible effect on the shapes of the templates and is therefore included as a normalisation uncertainty. The magnitude of this uncertainty ranges from 3–9% depending on the event category.

- **jet $\rightarrow e/\mu$ fake-rate**

A 20% uncertainty is applied to the W+jets yield in the $\tau_e\tau_\mu$ channel which contains at least one lepton originating from a jet-fake. For all other processes the fraction of events with such fakes is small and therefore this uncertainty is negligible.

- **$l \rightarrow \tau_h$ fake-rate**

The uncertainty on the $l \rightarrow \tau_h$ fake-rate affects mainly the $Z/\gamma^* \rightarrow ll$ process. The uncertainties on the $Z/\gamma^* \rightarrow ll$ yield amount to 27%, 11%, and 23% in the $\tau_\mu\tau_h$, $\tau_e\tau_h$, and $\tau_h\tau_h$ channels respectively.

5.6.2 Shape uncertainties

The uncertainties that affect both the shapes and the yields of the background and signal processes are summarised below. Theoretical uncertainties on the GF signal are described in Section 5.1.1.

- **Hadronic tau energy scale**

Three uncorrelated uncertainties each that affect only one of the three tau decay modes used

in the analysis are applied. In 2016 each uncertainty varies the hadronic tau energy scale by 1.2%. In 2017 the size of the uncertainty is 0.8% and 0.9% for the 1-prong and 3-prong decay modes respectively. Shifts in the energy scales are also propagated to the E_T^{miss} .

- **Electron energy scale**

The uncertainty on the electron energy scale amount to 1% for electron in the barrel detector region and 2.5% in the endcaps. Shifts in the energy scales are also propagated to the E_T^{miss} .

- **$e \rightarrow \tau_h$ energy scale**

Two uncorrelated uncertainties each that affect only one of the 1-prong tau decay modes are applied to the $Z/\gamma^* \rightarrow ee$ process in the $\tau_e \tau_h$ channel. These uncertainties vary the energy scale correction by 3% for both the 1-prong and 1-prong+1- π^0 decay modes in 2016. In 2017 the size of the uncertainty is 1.4% and 1.8% for the 1-prong and 1-prong+1- π^0 decay mode respectively. Shifts in the energy scales are also propagated to the E_T^{miss} .

- **$\mu \rightarrow \tau_h$ energy scale**

Two uncorrelated uncertainties each that affect only one of the 1-prong tau decay modes are applied to the $Z/\gamma^* \rightarrow \mu\mu$ process in the $\tau_\mu \tau_h$ channel. These uncertainties vary the energy scale correction by 1.5% for both the 1-prong and 1-prong+1- π^0 decay modes in 2016. In 2017 the size of the uncertainty is 0.3% and 1.0% for the 1-prong and 1-prong+1- π^0 decay mode respectively. Shifts in the energy scales are also propagated to the E_T^{miss} .

- **Jet energy scale uncertainty**

Uncertainties on the jet energy corrections are provided as a function of jet p_T and η . There exists 27 (28) independent sources of jet energy scale uncertainties for 2016 (2017), some of which are specific to particular regions of the detector. These sources are grouped together according to the regions of the detector that they affect to form three independent uncertainties: one group that contains all uncertainties affecting only the central detector ($0 < |\eta| < 3$), one group that contains all uncertainties affecting only the HF detector ($|\eta| > 3$), and one group containing all uncertainties that affect the whole detector ($|\eta| > 0$). The uncertainties are grouped together using the formula

$$\text{uncertainty} = \sqrt{\sum s_i^2}, \quad (5.12)$$

where s_i is the uncertainty for a source i . For all samples without E_T^{miss} recoil corrections applied, namely $t\bar{t}$, diboson, and single-top samples, the jet energy scale shifts are also propagated to the E_T^{miss} .

- **E_T^{miss} unclustered energy uncertainty**

The E_T^{miss} unclustered energy uncertainty is applied to all MC processes that do not have recoil corrections applied, namely $t\bar{t}$, diboson, and single-top samples.

- **E_T^{miss} recoil correction uncertainties.**

The uncertainties on the E_T^{miss} recoil corrections are estimated by varying the uncertainties on the recoil correction parameters within their uncertainties.

- **Non- $Z/\gamma^* \rightarrow \tau\tau$ contribution to embedded samples**

As described in Section 3.5.2 the embedded samples provide an estimate for all processes producing a pair of genuine tau leptons and therefore contain a small amount of $t\bar{t}$ and diboson events. These processes produce tau leptons predominantly from two $W \rightarrow \tau\nu$ decays and the kinematics of these taus can differ slightly from those produced by $Z/\gamma^* \rightarrow \tau\tau$ decays - as is assumed during the embedding procedure. A shape uncertainty defined by adding/subtracting 10% of the diboson and $t\bar{t}$ contribution to the embedded templates using MC is used to cover any resulting bias.

- **Top p_T reweighting**

An uncertainty between no and twice the correction on the reweighting is applied to $t\bar{t}$ events.

- **Drell-Yan p_T /mass reweighting**

An uncertainty between no and twice the correction on the reweighting is applied to $Z/\gamma^* \rightarrow l\bar{l}$ events in all channels.

- **Fake-factor uncertainties**

Uncertainties on the shapes of the $j \rightarrow \tau_h$ templates come from several sources:

- Statistical uncertainty on the FF measurements, given by the fit uncertainty to the p_T dependence.
- Systematic uncertainties related to the non-closure correction, as a function of the visible mass.
- Systematic uncertainties on the DR→AR extrapolation corrections.
- Systematic uncertainties on the fractions of W +jets and $t\bar{t}$ events with one jet $\rightarrow \tau_h$ in the $\tau_h\tau_h$ channel. This is necessary as the fake-factors for these processes must be scaled by a factor of 2 for this channel as described in Section 3.5.3. For the $\tau_l\tau_h$ channels the uncertainties on the fractions are negligible compared to the other systematic uncertainties.
- An uncertainty due to the subtraction of the genuine τ_h contribution using simulation.

- **$\tau_e\tau_\mu$ QCD uncertainties**

- The uncertainties on the linear fit parameters used to measure the $\dot{R}_{\text{QCD}}$ ratios are propagated to the QCD templates as uncorrelated uncertainties.
- The uncertainty on the non-QCD background contribution to the same-sign data, which is subtracted using simulation, is estimated by varying the subtracted component by $\pm 10\%$.

- **Finite statistics in background templates**

Limitations in the statistical sizes of the signal and background templates are accounted for using the "Barlow-Beeston lite" method, which assigns a single nuisance parameter per bin that scales the total bin yield [131, 132]. With this approach, the maximum likelihood estimator for each of these nuisance parameters is independent from the others, and can be solved analytically instead of fit during the minimisation. This method gives a significant reduction in the minimisation time and does not change the expected sensitivity to α_{hgg} .

5.6.3 Correlations between MC and embedding uncertainties

MC and embedded events have several common systematic uncertainties relating to the visible tau decay products. These uncertainties include: the lepton/ τ_h identification and isolation efficiency uncertainties, the trigger efficiency uncertainties, and the electron/ τ_h energy scale uncertainties. As tau leptons are simulated in a similar way for both samples, uncertainties relating to these quantities may have correlations which need to be taken into account. Selection variables typically depend not only on the simulated objects but also on other nearby objects present in the event. For the embedded samples these additional objects either come from data or may be missing altogether, if a quantity is computed in an empty detector environment. This means that the uncertainties are likely only partially correlated, and a suitable choice of the correlation needs to be adopted.

The trigger in the embedded samples is simulated in an empty detector environment, this results in significantly different trigger efficiencies and corrections compared to the MC. The trigger

efficiency uncertainties are therefore treated as being uncorrelated. The other identification and isolation efficiencies are typically more similar for MC and embedding, and are consequently treated as being 50% correlated. The electron and τ_h energy scales are also found to differ by up to 1%; although the τ_h energy scale is corrected to MC as described Section 3.4.9. The energy scale uncertainties are therefore treated as being 50% correlated.

5.6.4 Correlations between uncertainties in 2016 and 2017

The uncertainties in 2016 and 2017 are treated as being either uncorrelated, fully correlated, or partially correlated. Uncertainties that are uncorrelated are typically statistical in nature, whereas fully correlated uncertainties have a negligible statistical component. Uncertainties with small statistical components may still be decorrelated, however, based on arguments relating to, for example, differences in detector conditions between years, differences between analysis selections/algorithms, and differences in the measurements of efficiencies/corrections for each year. The partially correlated uncertainties typically have both statistical and non-statistical components of similar magnitudes, and therefore a 50% correlation is usually assumed. Table 5.2 summarises the correlations used for each uncertainty, with the justification for each choice given in the third column.

Table 5.2: Summary of the correlations between uncertainties in 2016 and 2017.

Uncertainty	Correlation	Motivation / notes
Luminosity	Uncorrelated	Based on the recommendations in Ref. [189].
e ID/isolation efficiency	Fully correlated	The statistical uncertainties on the measured scale-factors are negligible.
μ ID/isolation efficiency	Fully correlated	The statistical uncertainties on the measured scale-factors are negligible.
τ_h ID/isolation efficiency	50% correlated	The statistical and non-statistical components of the uncertainties have similar magnitudes.
τ_h trigger efficiencies	50% correlated	The statistical and non-statistical components of the uncertainties have similar magnitudes.
e/ μ trigger efficiencies	Uncorrelated	Several of the online selections and algorithms were changed between 2016 and 2017.
B-tagging efficiency	Uncorrelated	The pixel detector was replaced between 2016 and 2017.
Cross sections and branching ratios	Fully correlated	The cross sections and branching ratios are the same for both years.
Embedding yield	Uncorrelated	Different filters and algorithms used in the triggers used to select dimuon events.
$\tau_e\tau_\mu \dot{R}_{\text{QCD}}$ extrapolation	50% correlated	Similar side-band regions are used to measure $\dot{R}_{\text{QCD}}$ in both years, but using different electron ID working points. Sizes of isolated vs anti-isolated $\dot{R}_{\text{QCD}}$ in validation region are slightly different.
$\tau_e\tau_\mu \dot{R}_{\text{QCD}}$ fit uncertainties	Uncorrelated	Fit uncertainties are statistical in nature.

Non-QCD subtraction for $\tau_e\tau_\mu$ QCD estimation	50% correlated	Some uncertainties affecting the subtracted yield are (partially) correlated e.g ID/isolation efficiency, and others are uncorrelated e.g luminosity.
jet $\rightarrow$ fake-rate	50% correlated	Electron ID working point changed between 2016 and 2017.
$l \rightarrow \tau_h$ fake-rate	50% correlated	The statistical and non-statistical components of the uncertainties have similar magnitudes.
τ_h energy scale	50% correlated	The statistical and non-statistical components of the uncertainty have similar magnitudes.
e energy scale	Fully correlated	The statistical component of the uncertainty is negligible.
$l \rightarrow \tau_h$ energy scale	Uncorrelated	Uncertainties on corrections are predominantly statistical. Size of corrections are different in 2016 and 2017.
Jet energy scale	Partially correlated	Each of the 27 (28) uncertainty sources used in 2016 (2017) have a recommendation for the correlation, which are either 0%, 50%, or 100% correlated (based on the relative sizes of the statistical and non-statistical components). These correlations are used to split the grouped uncertainties into correlated and uncorrelated parts using: correlated = $\sqrt{\sum \rho_i s_i^2}$, and uncorrelated = $\sqrt{\sum (1 - \rho_i) s_i^2}$, where ρ_i is the correlation for source i with uncertainty s_i .
E_T^{miss} uncertainties	Uncorrelated	The statistical components of the uncertainties are much larger than the non-statistical components.
non- Z/γ^* contribution to embedding	Fully correlated	Any bias in modelling taus from $W \rightarrow \tau\nu$ decays should be the same for both 2016 and 2017.
t p_T reweighting	Fully correlated	Same weights applied for both 2016 and 2017.
Z/γ^* p_T reweighting	Fully correlated	Same procedure used to derive weights for 2016 and 2017. Similar data/MC differences for unweighted dimuon events.
Fake-factor p_T fit uncertainties	Uncorrelated	Fit uncertainties are statistical in nature.
Fake-factor correction uncertainties	50% correlated	Corrections derived using the same procedures in 2016 and 2017 but derived from different data/MC events. The size of the corrections are slightly different in 2016/2017.
Fake-factor real τ_h subtraction	50% correlated	Some uncertainties affecting the real τ_h yield are (partially) correlated e.g tau ID efficiency, and others are uncorrelated e.g luminosity.
Fake-factor fractions	50% correlated	All non-QCD components estimated using MC, as for the above some uncertainties affecting these yields are uncorrelated and others are correlated.

Bin-by-Bin statistical uncertainties	Fully uncorrelated	Signal and background templates are statistically independent.
GF UE uncertainty	Uncorrelated	Different UE tunes used in 2016 and 2017.
GF μ_R, μ_F uncertainty	Fully correlated	Same matrix element computations in 2016 and 2017.
GF parton shower uncertainty	Fully correlated	Same parton shower generator in 2016 and 2017.
GF finite top quark mass reweighting	Fully correlated	Same corrections applied in 2016 and 2017.
GF jet migration uncertainties	Fully correlated	Same uncertainties applied in both years.

5.7 Measurement procedure and statistical interpretation

For generic Higgs $\mathcal{CP}$ models, both the cross section and $H \rightarrow \tau\tau$ branching ratio, $\mathcal{B}^{\tau\tau}$, may be modified with respect to the SM values. The μ_{ggh} and μ_V parameters introduced previously scale the GF and VBF/VH cross sections respectively. In principle we could introduce an additional rate parameter $\mu^{\tau\tau}$ to scale the $H \rightarrow \tau\tau$ branching ratio simultaneously for the GF, VBF, and VH processes, however, the analysis is not able to disentangle modifications to the cross sections and branching ratios. For example a simultaneous excess in the GF, VBF, and VH yields could be equally interpreted as: an enhanced $\mu^{\tau\tau}$, a coincident enhancement of μ_{ggh} and μ_V , or some combination of the two. Therefore, only two rate parameters are needed: one to scale the GF $\sigma \times \mathcal{B}^{\tau\tau}$, and the other to simultaneously scale the VBF and VH $\sigma \times \mathcal{B}^{\tau\tau}$'s. These parameters are suitably named $\mu_{\text{ggh}}^{\tau\tau}$ and $\mu_V^{\tau\tau}$, respectively, where $\mu_i^{\tau\tau} = \mu_i \cdot \mu^{\tau\tau}$. The $\mu_V^{\tau\tau}$ parameter is profiled in the fit meaning it is allowed to float freely like a nuisance parameter, but is bounded such that it may take only positive values.

A simultaneous maximum likelihood fit including all categories and the $t\bar{t}$ control region is used to extract the measured values of α_{hgg} and $\mu_{\text{ggh}}^{\tau\tau}$. The 2016 and 2017 analyses are combined taking into account the correlations between systematic uncertainties described in Section 5.6.4. The likelihood function given in Equation 4.3 is modified so that the signal expectation in each bin s_i becomes a function of α_{hgg} ,

$$\mathcal{L}(\text{data} | \mu_{\text{ggh}}^{\tau\tau}, \alpha_{\text{hgg}}, \theta) = \prod_i \text{Poisson}(n_i | \mu_{\text{ggh}} \cdot s_i(\alpha_{\text{hgg}}, \theta) + b_i(\theta)) \cdot p(\tilde{\theta} | \theta). \quad (5.13)$$

The best fit values of α_{hgg} and $\mu_{\text{ggh}}^{\tau\tau}$ are determined by maximising the $\mathcal{L}$. When fitting for α_{hgg} the parameter $\mu_{\text{ggh}}^{\tau\tau}$ is profiled with the constraint $\mu_{\text{ggh}}^{\tau\tau} \geq 0$. Likewise, when fitting for $\mu_{\text{ggh}}^{\tau\tau}$ the α_{hgg} parameter is profiled between $-90^\circ \leq \alpha_{\text{hgg}} \leq 90^\circ$. The confidence intervals are determined from the one dimensional negative log-likelihood defined as

$$-2\Delta \ln \mathcal{L} = -2 \cdot (\ln \mathcal{L}(x) - \ln \mathcal{L}(\hat{x})), \quad (5.14)$$

where x is the parameter of interest and $\hat{x}$ is its best fit value. The 68% and 95% confidence intervals are defined by $-2\Delta \ln \mathcal{L} = 1.00$ and 3.84 respectively.

The likelihood function may also be re-parametrised as a function of κ_t and $\tilde{\kappa}_t$ using Equations 5.1 and 5.3, $\mathcal{L}(\text{data} | \mu_{\text{ggh}}^{\tau\tau}, \alpha_{\text{hgg}}, \theta) \rightarrow \mathcal{L}(\text{data} | \kappa_t, \tilde{\kappa}_t, \mu^{\tau\tau}, \theta)$. A two dimensional likelihood scan of

κ_t and $\tilde{\kappa}_t$ can then be performed. As discussed above the analysis is not able to disentangle $\mu^{\tau\tau}$ and μ_{ggh} and consequently has no sensitivity to κ_t and $\tilde{\kappa}_t$ unless some value for $\mu^{\tau\tau}$ is assumed. The value of $\mu^{\tau\tau}$ is therefore fixed to unity meaning $\mathcal{B}^{\tau\tau}$ is assumed to be the SM value. The two dimensional negative log-likelihood defined in Equation 4.11 is then used to construct the 68% and 95% confidence regions, which are given by $-2\Delta \ln \mathcal{L}_{2D} = 2.28$ and 5.99 respectively. Despite the choice of $\mu^{\tau\tau} = 1$ it is easy to interpret the results for a generic model with $\mu^{\tau\tau} \neq 1$ by scaling the best fit values of κ_t and $\tilde{\kappa}_t$ and confidence regions by $1/\sqrt{\mu^{\tau\tau}}$.

For all one and two dimensional scans the expected best-fit values and confidence intervals/regions are computed using the Asimov dataset.

5.8 Results

To check the validity of the model, several goodness of fit tests are performed before the final simultaneous fit to data. The full set of systematic uncertainties and correlations are taken into account when performing these checks. Tests are chosen that are sensitive to local deviations from the model in individual bins, and correlated deviations across several bins. These tests all indicate a good compatibility between the data and the model.

The discriminating variable distributions after the simultaneous fit to the data are shown in Figures 5.16–5.20 for the 0-jet, boosted, and dijet tight- m_{jj} boosted categories. For Figures 5.17–5.20 the unrolled 2D discriminants described in Section 5.4 are shown. In these figures the small axis to the bottom left indicates the binning used for the x variables. The binning of the y variables is indicated by the broken lines, where the labels above the histograms to the right of these lines specify the range of each y bin. The GF signal shown by the solid red line is the expected distribution when α_{hgg} and $\mu_{\text{ggh}}^{\tau\tau}$ are set to their best-fit values of $\alpha_{\text{hgg}} = -59^\circ$ and $\mu_{\text{ggh}}^{\tau\tau} = 0.68$. The solid green line shows the expected GF signal for $\alpha_{\text{hgg}} = 0^\circ$ and $\mu_{\text{ggh}}^{\tau\tau} = 0.68$ for comparison. The background histograms are stacked. The contribution labelled "Electroweak" in the $\tau_l\tau_h$ and $\tau_e\tau_\mu$ channels include the diboson, single top, and, for the $\tau_e\tau_\mu$ channel, W+jets backgrounds. The background labelled $l \rightarrow \tau_h$ in the $\tau_h\tau_h$ channel includes $t\bar{t}$, single top, diboson, and Z/γ^* processes where one of the τ_h candidates is faked by a lepton and the other candidate is either also a $l \rightarrow \tau_h$ fake or a genuine τ_h . The VBF and VH background are combined into the solid blue histogram and scaled to the best-fit value of $\mu_V^{\tau\tau} = 0.64$.

In the fit to the data only 84 of the 2678 nuisance parameters are pulled by more than one standard deviation, and none by more than two standard deviations. This small number of pulls greater than one standard deviation is perfectly reasonable given the large number of nuisance parameters that are included.

The 0-jet and boosted categories have no direct sensitivity to α_{hgg} but their inclusion in the final fit provides sensitivity to the measurement of $\mu_{\text{ggh}}^{\tau\tau}$ and helps to constrain the background uncertainties, many of which are common across categories. As a consequence of these effects the expected uncertainty on α_{hgg} is reduced by $\approx 20\%$ by the inclusion of these categories in the fit. Likewise, for the low/high $m_{\tau\tau}$ bins ($m_{\tau\tau} < 100$ GeV / $m_{\tau\tau} > 150$ GeV) in the dijet categories the expected signal yields are very small compared to the background, but their inclusion also helps to constrain the backgrounds in the more signal-sensitive bins with $m_{\tau\tau} \approx 125$ GeV.

The left and right insets of Figure 5.21 show the scans of the negative log likelihood $-2\Delta \ln \mathcal{L}$ as a function of α_{hgg} and $\mu_{\text{ggh}}^{\tau\tau}$ respectively. The scans shown by the solid blue lines are the result of

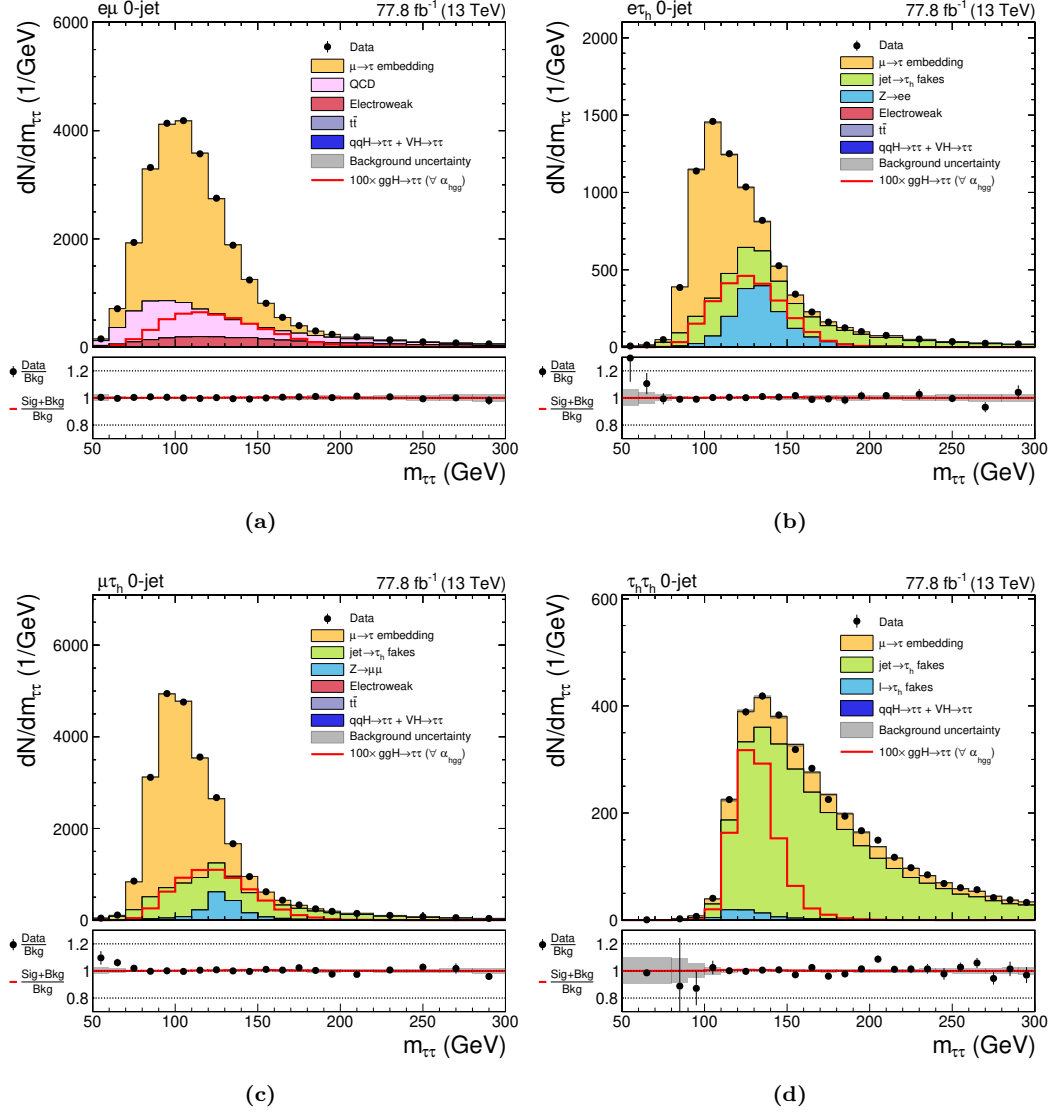


Figure 5.16: Observed and predicted $m_{\tau\tau}$ distributions in the 0-jet category after the simultaneous fit to data, where the histograms for the 2016 and 2017 analyses have been combined. The uncertainty band shown in shaded grey includes all sources of systematic uncertainties. Insets (a)–(d) show the distributions for the $\tau_e\tau_\mu$, $\tau_e\tau_h$, $\tau_\mu\tau_h$, and $\tau_h\tau_h$ channels respectively.

the fit to data, whereas the dashed blue lines show the equivalent scans using the Asimov data set. The grey lines are used to indicate the 68% and 95% CLs. The best-fit value of α_{hgg} is $-59^\circ_{-50^\circ}^{+47^\circ}$, which is a modest 1.16σ deviation from the expectation and is therefore consistent with the SM. The best-fit value of $\mu_{ggh}^{\tau\tau}$ is $0.68^{+0.32}_{-0.28}$, which is again consistent with the SM expectation.

The uncertainty on α_{hgg} can be decomposed into statistical, experimental and theoretical systematic uncertainties. The statistical uncertainty is determined by fixing all systematic uncertainty parameters to their best-fit values and repeating the scan of α_{hgg} . An additional scan of α_{hgg} is performed where only the theoretical systematic uncertainties are fixed and all others are allowed to float. Figure 5.22 shows both of these scans compared to the full scan. The three uncertainties determined in each of the three scans are used to estimate the theoretical ("theory") and experimental ("syst") systematic uncertainties using $\sigma_{\text{tot}}^2 = \sigma_{\text{stat}}^2 + \sigma_{\text{syst}}^2 + \sigma_{\text{theory}}^2$. The decomposed uncertainties are $\alpha_{hgg} = -59^\circ_{-39^\circ}^{+40^\circ}$ (stat) $_{-26^\circ}^{+21^\circ}$ (syst) $_{-17^\circ}^{+14^\circ}$ (theory). The total uncertainty

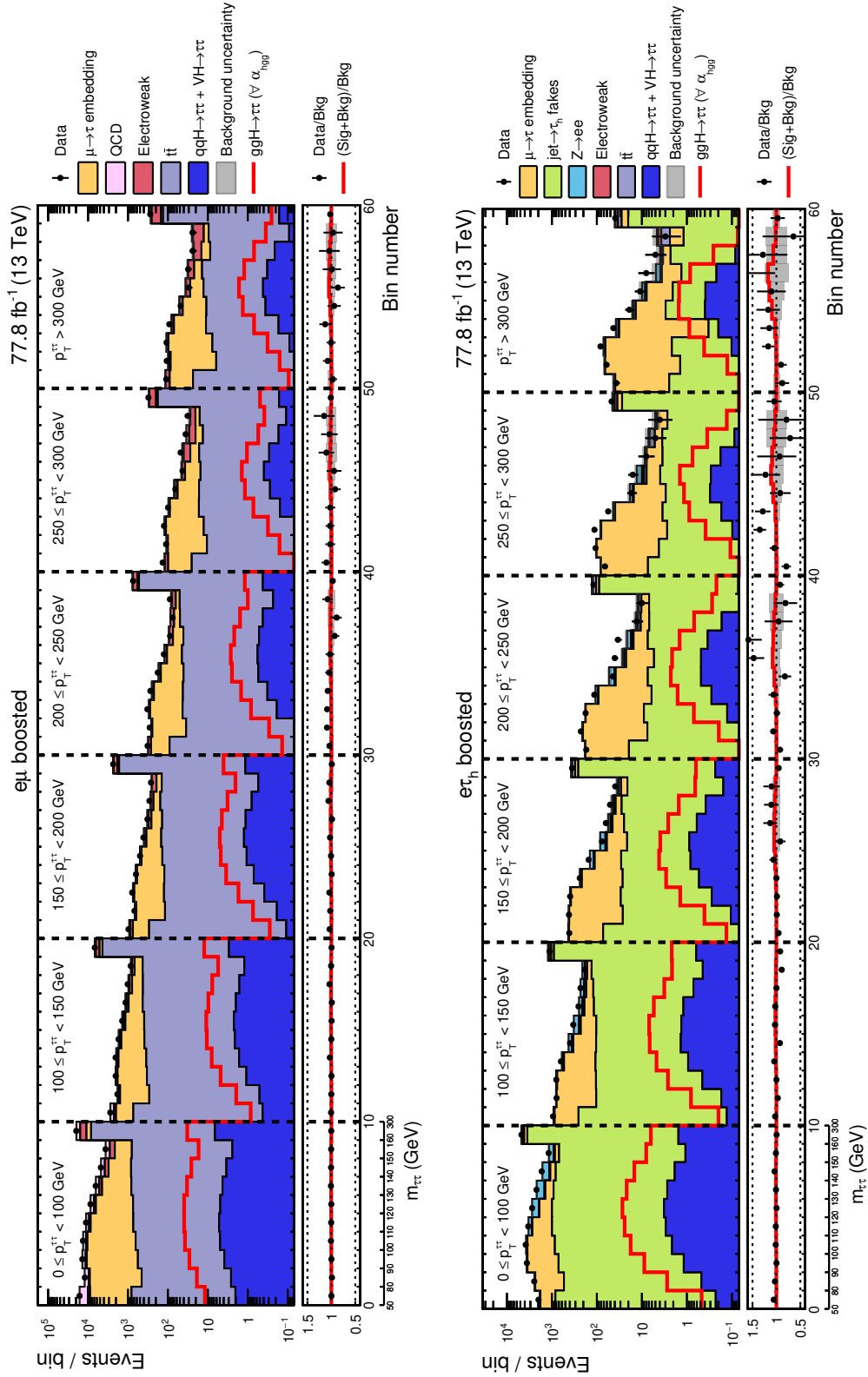


Figure 5.17: Observed and predicted 2D distributions in the boosted category after the simultaneous fit to data, where the histograms for the 2016 and 2017 analyses have been combined. The uncertainty band shown in shaded grey includes all sources of systematic uncertainties. The distributions are shown for the $\tau_e \tau_\mu$ (upper) and $\tau_e \tau_h$ (lower) channels.

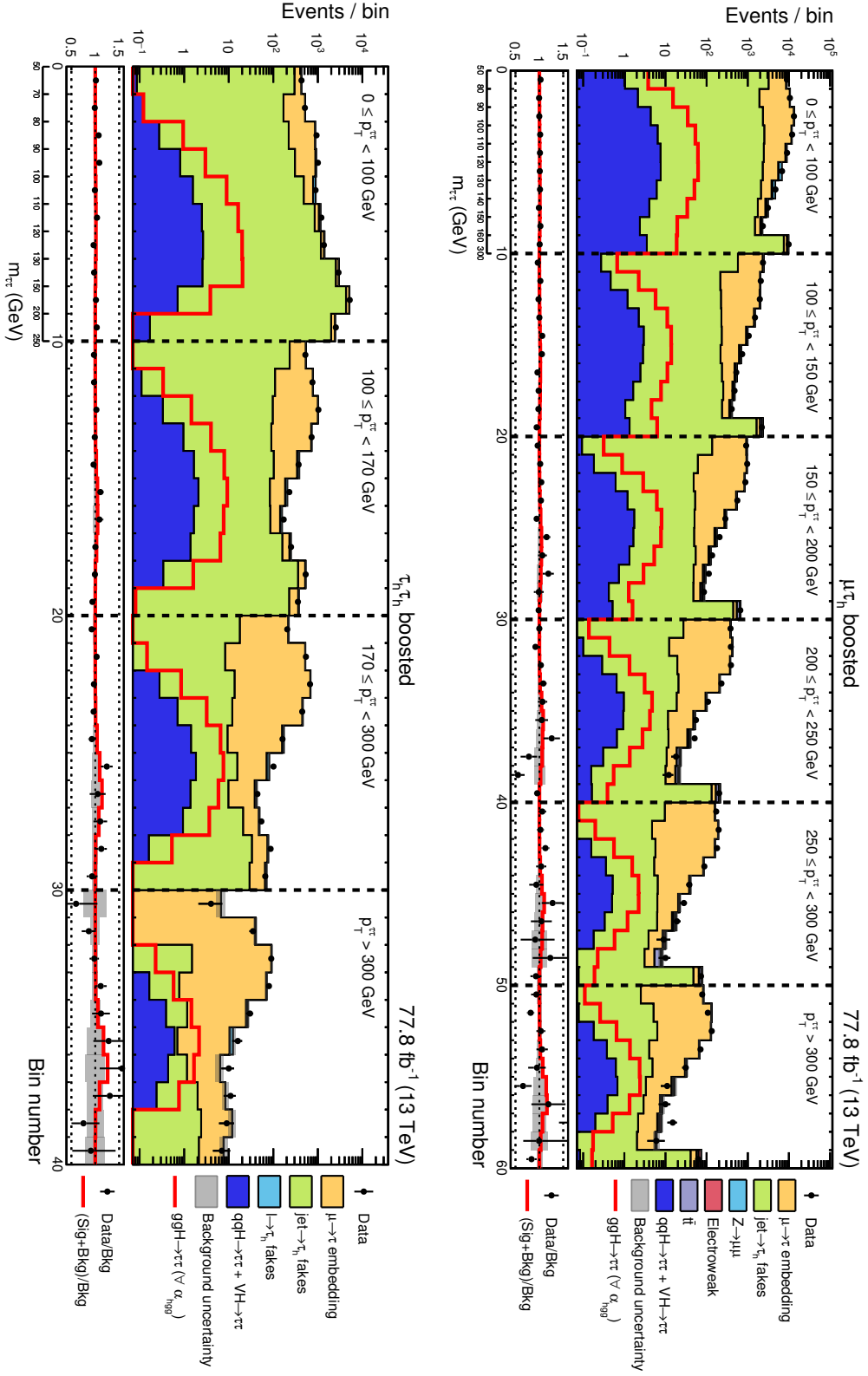


Figure 5.18: Observed and predicted 2D distributions in the boosted category after the simultaneous fit to data, where the histograms for the 2016 and 2017 analyses have been combined. The uncertainty band shown in shaded grey includes all sources of systematic uncertainties. The distributions are shown for the $\tau_\mu \tau_\eta$ (upper) and $\tau_\eta \tau_\eta$ (lower) channels.

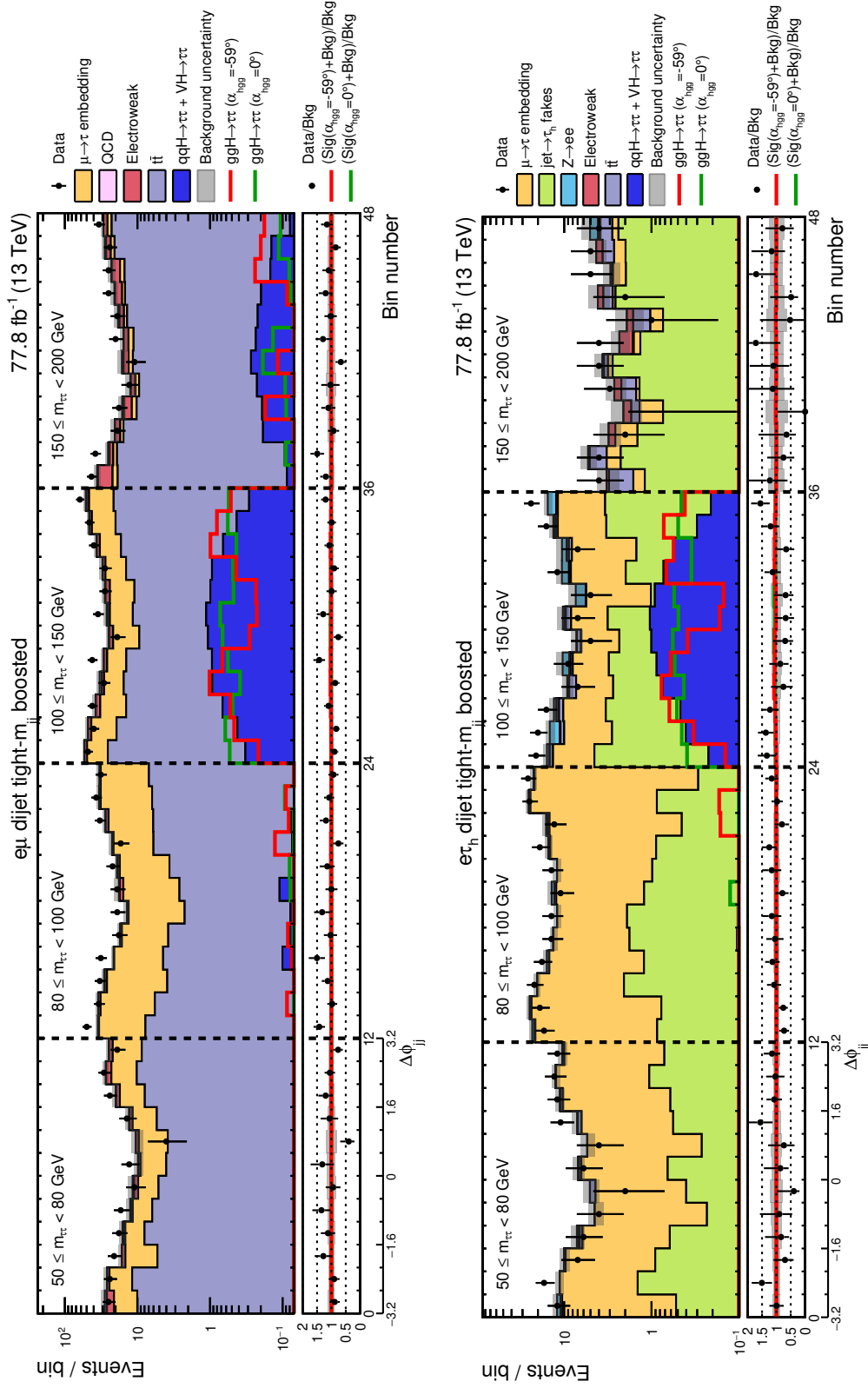


Figure 5.19: Observed and predicted 2D distributions in the dijet tight- m_{jj} boosted category after the simultaneous fit to data, where the histograms for the 2016 and 2017 analyses have been combined. The uncertainty band shown in shaded gray includes all sources of systematic uncertainties. The distributions are shown for the $\tau_e \tau_\mu$ (upper) and $\tau_e \tau_h$ (lower) channels.

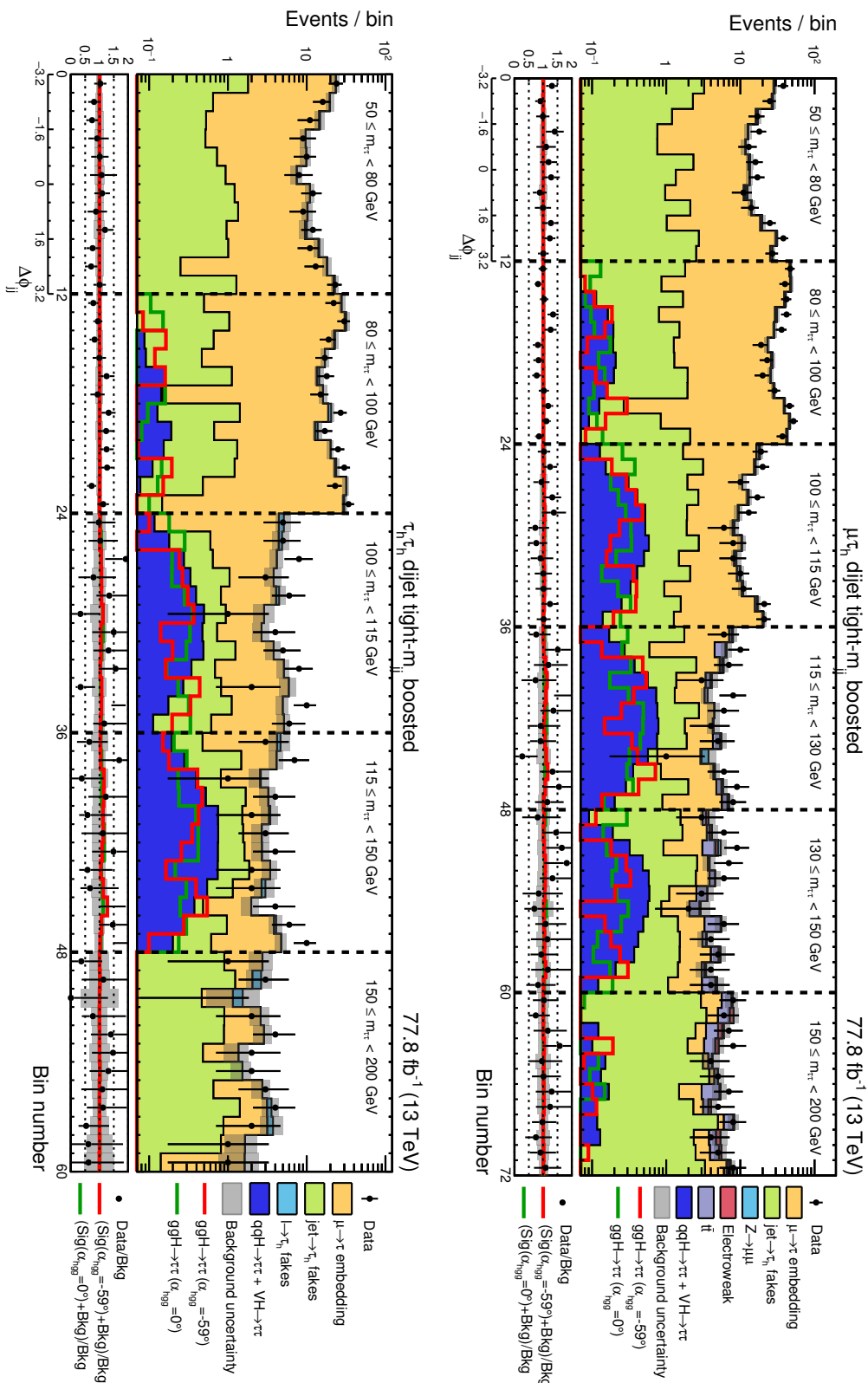


Figure 5.20: Observed and predicted 2D distributions in the dijet tight- n_{jj} boosted category after the simultaneous fit to data, where the histograms for the 2016 and 2017 analyses have been combined. The uncertainty band shown in shaded grey includes all sources of systematic uncertainties. The distributions are shown for the $\tau_\mu\tau_h$ (upper) and $\tau_h\tau_h$ (lower) channels.

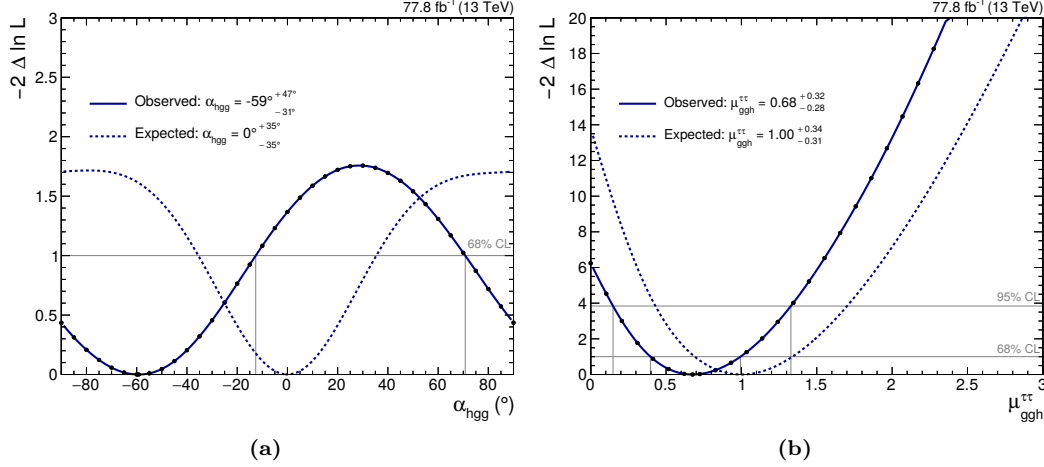


Figure 5.21: Scans of the negative log-likelihood, $-2\Delta\ln\mathcal{L}$, as a function of α_{hgg} (left) and $\mu_{\text{ggh}}^{\tau\tau}$ (right). The 68% and 95% CL intervals are indicated by the grey lines.

is found to be dominated by the statistical uncertainty. The contribution from the theoretical systematic uncertainties is also found to be smaller than the experimental contribution. The dominant experimental uncertainties are due to limitations in statistical sizes of the signal and background templates and the jet energy scale corrections. The dominant theoretical uncertainty is due to the $\mu_{\text{R}}/\mu_{\text{F}}$ scale variations.

Figure 5.23 shows the results of the two dimensional scan of κ_t and $\tilde{\kappa}_t$, where the left inset shows the observed scan and the right inset shows the expected scan based on Asimov data. The best-fit values of κ_t and $\tilde{\kappa}_t$ are $0.42^{+0.41}_{-0.68}$ and $-0.47^{+0.36}_{-0.16}$ respectively. The 68% and 95% CL regions are shown by the coloured bands. The best-fit values are slightly outside of the 68% region but within the 95% region and therefore the measured values of κ_t and $\tilde{\kappa}_t$ are consistent with the SM expectation. The sensitivity to κ_t and $\tilde{\kappa}_t$ comes partly from the rate information ($\mu_{\text{ggh}}^{\tau\tau}$), and partly from the differential information (α_{hgg}). A measurement of κ_t and $\tilde{\kappa}_t$ based only on the former would give confidence region boundaries that are exactly elliptical in shape, since the rate information has no power to differentiate between deviations in κ_t and $\tilde{\kappa}_t$. The inclusion of the differential information on the other hand does allow κ_t and $\tilde{\kappa}_t$ to be disentangled, which consequently causes the narrowing of the confidence bands at $(\kappa_t, \tilde{\kappa}_t) \approx \pm(0.7, 0.2)$ in the observed scan and at $\kappa_t \approx 0$ in the expected scan. The scans are also symmetric about $y=x$, this is because the three terms in Equation 5.2 depend only on the products κ_t^2 , $\tilde{\kappa}_t^2$, and $\kappa_t \cdot \tilde{\kappa}_t$ and therefore the differential observables are only sensitive to the relative sign of κ_t and $\tilde{\kappa}_t$. The best-fit values of $(\kappa_t, \tilde{\kappa}_t) = \pm(0.42, -0.47)$ are thus equally likely.

5.8.1 Cross check allowing for anomalous HVV couplings

Now we turn to the cross check using an alternative treatment of the VBF and VH backgrounds described in Section 5.5.3. As discussed previously the VBF and VH samples in this case are modelled using JHUGen at LO accuracy. The results obtained using these samples can differ slightly from those obtained in the previous section due to statistical differences in the samples, and due to the LO vs NLO descriptions of the variables [190]. The change in the sensitivity to α_{hgg} when floating α_{hvv} should therefore be evaluated with respect to freezing $\alpha_{\text{hvv}} = 0$ in the JHUGen samples.

The scans of α_{hgg} and $\mu_{\text{ggh}}^{\tau\tau}$ for $\alpha_{\text{hvv}} = 0$ are shown in Figure 5.24. Compared to the equivalent

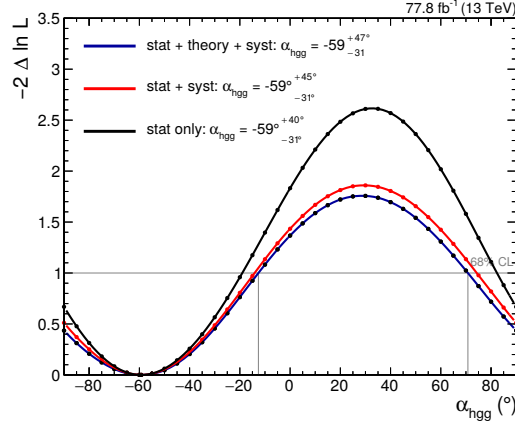


Figure 5.22: Scans of the negative log-likelihood, $-2\Delta\ln\mathcal{L}$, as a function of α_{hgg} . The "stat + theory + syst" scan corresponds to that of Figure 5.21 (a), where all systematic uncertainty parameters are profiled. For the "stat only" scan all systematic uncertainty parameters are fixed to their best-fit values. Finally, for the "stat + syst" scan the theoretical systematic uncertainties are fixed to their best-fit values and all other parameters are allowed to float. The 68% CL interval is indicated by the grey lines.

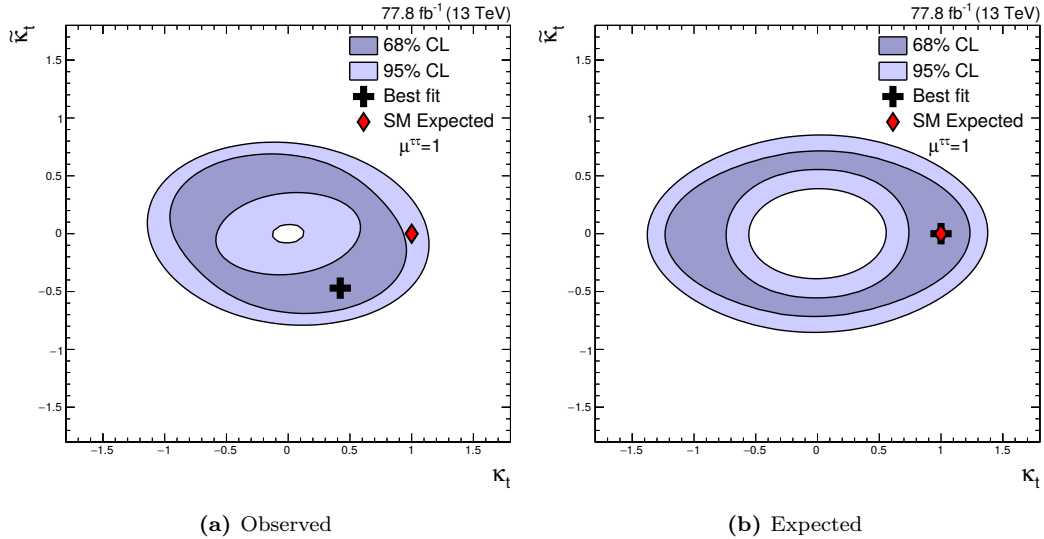


Figure 5.23: Two dimensional scans of the negative log-likelihood, $-2\Delta\ln\mathcal{L}$, as a function κ_t and $\tilde{\kappa}_t$. The 68% and 95% CL regions are indicated by the coloured bands. The left plot shows the scan from the fit to data, whereas the right plot shows the expectation using the Asimov dataset.

scans in Figure 5.21 the difference in the best-fit values of α_{hgg} is 3° which is small compared to the uncertainties. The best-fit value of $\mu_{\text{ggh}}^{\tau\tau}$ is also 12% smaller in this case, the size of this shift is again small when we consider the size of the uncertainties and the strong correlation between the measurements of $\mu_{\text{ggh}}^{\tau\tau}$ and $\mu_V^{\tau\tau}$.

The scans of α_{hgg} allowing α_{hvv} to float between $-2^\circ \leq \alpha_{\text{hvv}} \leq 2^\circ$ and between $-90^\circ \leq \alpha_{\text{hvv}} \leq 90^\circ$ are shown in Figure 5.25. In both cases there is no significant change to the expected uncertainty on α_{hgg} . The observed best-fit value is shifted by $\approx 2^\circ$ degrees for the first case, but this is overall a small effect compared to the size of the uncertainty. A significant shift of $\approx 20^\circ$ is observed when α_{hvv} is allowed to float completely freely. Figure 5.26 shows the likelihood scan

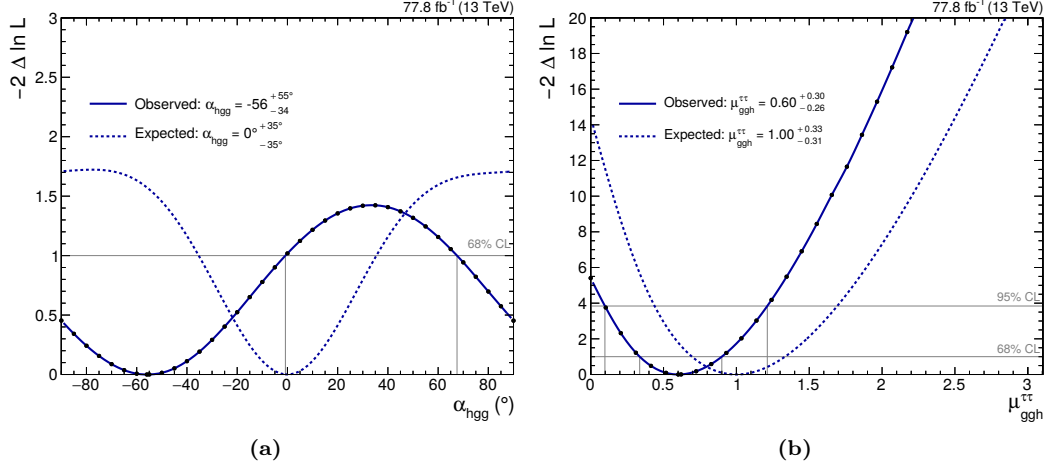


Figure 5.24: Scans of the negative log-likelihood, $-2\Delta\ln\mathcal{L}$, as a function of α_{hgg} (left) and $\mu_{ggh}^{\tau\tau}$ (right) using the LO JHUGen samples to model the VBF and VH processes with α_{hvv} fixed to 0° . The 68% and 95% CL intervals are indicated by the grey lines.

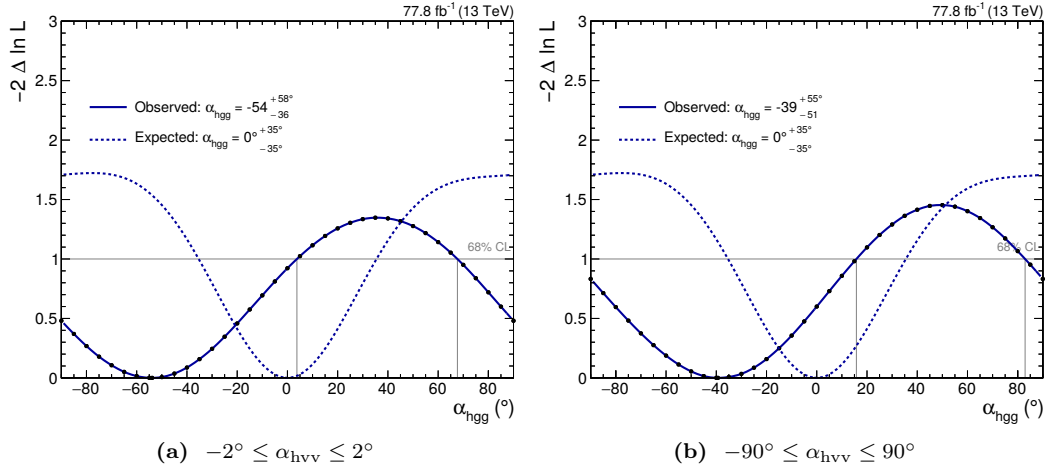


Figure 5.25: Scans of the negative log-likelihood, $-2\Delta\ln\mathcal{L}$, as a function of α_{hgg} using the LO JHUGen samples to model the VBF and VH processes with α_{hvv} floated within the ranges $-2^\circ \leq \alpha_{hvv} \leq 2^\circ$ (left) and $-90^\circ \leq \alpha_{hvv} \leq 90^\circ$ (right). The 68% CL interval is indicated by the grey line.

using α_{hvv} as the parameter of interest and floating all other parameters. The best fit value of α_{hvv} is $-37^{+29}_{-36}^\circ$. Clearly there is some slight deviation in the data which gives non-zero values of α_{hgg} and/or α_{hvv} . When α_{hvv} is fixed a best-fit value of $\alpha_{hgg} = -56^\circ$ is obtained, whereas when both α_{hvv} and α_{hgg} are floated, a smaller value of $\alpha_{hgg} = -39^\circ$ is preferred in combination with an additional pull of $\alpha_{hvv} = -37^\circ$.

The conclusion of this cross-check is that if one wishes to use the scans of α_{hgg} to test models in which α_{hvv} is large then the effects of anomalous HVV couplings should be taken into account. However, this test has also demonstrated that for weakly interacting models where $|\alpha_{hvv}| < 2^\circ$ we currently have no sensitivity to the effects of anomalous HVV couplings, and these can be safely neglected. When testing such models care should be taken in letting α_{hvv} float completely freely without good justification, as this can have a significant effect on the observed measurement of

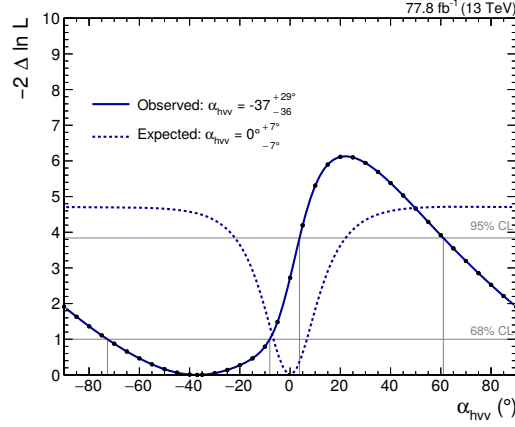


Figure 5.26: Scan of the negative log-likelihood, $-2\Delta\ln\mathcal{L}$, as a function of α_{hvv} using the LO JHUGen samples to model the VBF and VH processes. All other parameters are profiled in the fit. The 68% and 95% CL intervals are indicated by the grey lines.

α_{hgg} despite the expected results showing no indications of such an effect.

5.9 Future measurements and combinations

The analysis presented in this chapter will be extended to include also the dataset collected by CMS in 2018. This dataset corresponds to an additional integrated luminosity of 58.8 fb^{-1} . Based on projections to the 2016+2017+2018 integrated luminosity the uncertainty on α_{hgg} can be reduced to $\sim 26^\circ$ by the inclusion of these data. The sensitivity can be further improved by combining results for several Higgs decay channels. A combination of the $H \rightarrow \tau\tau$ and $H \rightarrow ZZ$ channels is planned by the autumn of 2019. The $H \rightarrow ZZ$ final state is expected to be less sensitive than $H \rightarrow \tau\tau$, but early studies suggest that combining the two channels should allow for the exclusion of large mixing angle scenarios ($\alpha_{hgg} \gtrsim 80^\circ$) at the 95% CL. A measurement of α_{hgg} in the $H \rightarrow \gamma\gamma$ final state is also expected by early 2020. The sensitivity for this channel is expected to be more comparable to the $H \rightarrow \tau\tau$ channel. Under the assumption that the $H \rightarrow ZZ$ and $H \rightarrow \gamma\gamma$ measurements collectively have about the same sensitivity to α_{hgg} as the $H \rightarrow \tau\tau$ result, CMS should be able to exclude scenarios with $\alpha_{hgg} \gtrsim 45^\circ$ using the Run 2 dataset.

Chapter 6

Development of L1 triggers targeting VBF

Higgs analyses targeting the ditau final state often suffer from low trigger efficiencies. This is due in part to the neutrinos produced in tau decays, which carry off a fraction of the energy resulting in soft τ_l and τ_h candidates. Hadronic tau triggers also typically impose relatively tight p_T and isolation requirements, which are needed to reduce the large background from QCD multi-jet events.

The characteristic jet topology of the VBF process is often exploited to reduce backgrounds. The VBF process typically results in two high- p_T jets, at least one of which is located in the forward detector regions ($|\eta| > 5$). These jet pairs usually have a large pseudorapidity separation and high invariant mass. The most recent CMS publication in the $H \rightarrow \tau\tau$ final state [2] found that VBF-tagged categories gave the best sensitivity to the inclusive $\sigma \times \mathcal{B}^{\tau\tau}$ measurement. As discussed in Chapter 1, measurements of the $\mathcal{CP}$ properties of the $H\tau\tau$ vertex are well motivated for constraining $\mathcal{CP}$ -violating Higgs scenarios. These measurements may in principle use Higgs bosons produced by any means, which includes the VBF process. It also was shown in the previous chapter that VBF-like GF event are sensitive to the $\mathcal{CP}$ properties of the $Ht\bar{t}$ vertex. Therefore, increasing the trigger efficiency for VBF (and VBF-like) events could potentially improve the sensitivity of the $\mathcal{CP}$ measurements of both $Ht\bar{t}$ and $H\tau\tau$ vertices.

The VBF topology has also been exploited by several HLT algorithms to increase signal efficiencies while maintaining reasonable event rates. The targeting of VBF events by the L1 trigger has however proved more difficult, as it has previously not been possible to implement complex multi-object variables such as m_{jj} in the L1 hardware. The upgrade of the L1 trigger during the end of year technical stop in 2015/2016 (as described in Section 2.8) opens up many more possibilities for L1 algorithms. Most notably, it is now possible to compute multiple object variables such as m_{jj} . This brings renewed hope that L1 triggers may be able to target the VBF production mode. The development of such triggers will be described in this chapter

6.1 Optimisation of the triggers

The main goals for the development of L1 VBF triggers are:

- To increase the VBF event yield for the two most sensitive ditau final states: $\tau_h\tau_h$ and $\tau_\mu\tau_h$. In particular for the $\tau_h\tau_h$ final state, which is expected to be the most important channel for $\mathcal{CP}$ measurements of the $H\tau\tau$ vertex.
- To select events as inclusively as possible. The triggers should target general features of the VBF topology, preferably without putting any requirements on the Higgs decay products. This means that the triggers can be of benefit to multiple analyses.

- To primarily target new regions of phase space not covered by the current L1 algorithms.

The second point is especially important, as the current L1 trigger must reduce the total event rate entering the HLT to 100 kHz. A trigger that is of benefit to multiple analyses can justify a larger share of the this total rate.

With the above goals in mind, the primary objective was to develop a L1 trigger that could target events using only the topology of the VBF jets. This is beneficial for $H \rightarrow \tau\tau$ analyses as it allows the low τ_h/τ_l p_T regions of the phase-space not currently recorded by CMS to be accessed. As this trigger selects VBF events in an inclusive way, it may also be useful for other analyses and therefore justify a larger share of the L1 rate. This trigger will be referred to in the following sections as the "dijet trigger". A second trigger was also developed requiring a muon in addition to the VBF jets. This trigger is designed to increase the signal acceptance for VBF events with a muon in the final state (e.g $H \rightarrow \tau_\mu\tau_h$). The muon requirement significantly reduces the L1 event rate allowing for looser selections on the jets. This trigger will be referred to in the following sections as the "muon+dijet trigger".

The main variables that were considered for the trigger optimisation studies are:

- The dijet m_{jj} . This is computed for all combinations of dijet pairs with $p_T > 30$ GeV. At least one pair is required to have m_{jj} larger than the given cut.
- The p_T of the leading L1 jet. The jet that meets this requirement does not necessarily have to be one of the jets that passes the m_{jj} selection. Keeping this jet independent from the m_{jj} selections also increases the signal efficiency for final states with jets, such as $H \rightarrow \tau_h\tau_h$ and $H \rightarrow b\bar{b}$, as one of the τ_h/b -jets can also pass the p_T requirement.
- The p_T of the sub-leading L1 jet. As above this jet does not necessarily have to be one of the jets that passes the m_{jj} selection.
- For the muon+dijet trigger, the p_T of the leading L1 muon is also included as a variable.

The efficiency turn-on curves measured in MC are shown for each of the considered variables in Figure 6.1. Figure 6.1 (b) shows the efficiency for a τ_h candidate to pass the L1 jet p_T requirement. The efficiency is observed to be considerably higher for τ_h candidates compared to jets.

The optimisation of the trigger selections requires estimations of the L1 rate and signal efficiencies. The signal efficiencies are estimated from MC, by measuring the fraction of events passing a combination of trigger selections. The sample of MC events used for this purpose depends on the trigger being studied. A sample of $H \rightarrow \tau_h\tau_h$ events was used to optimise the dijet trigger, whereas $H \rightarrow \tau_\mu\tau_h$ events were used to optimise the muon+dijet trigger.

The rate is estimated using a sample of "zero bias" events, which are events recorded by CMS without any trigger selections applied. The rate is computed as

$$\text{rate} = \epsilon N_b f_{\text{rev}}, \quad (6.1)$$

where ϵ is the fraction of zero bias events passing a given selection, N_b is the number of bunches per beam, and $f_{\text{rev}} = 11.245$ kHz is the LHC revolution frequency. The initial trigger optimisation used zero bias data collected during 2015 data-taking. This data was collected in different conditions to the design specification, with an average instantaneous luminosity $L = 0.15 \times 10^{-34} \text{ cm}^2\text{s}^{-1}$. The estimated rate must therefore be extrapolated to the design specification of $L' = 10^{-34} \text{ cm}^2\text{s}^{-1}$ using

$$\text{rate}' = \frac{L'}{L} \text{rate}. \quad (6.2)$$

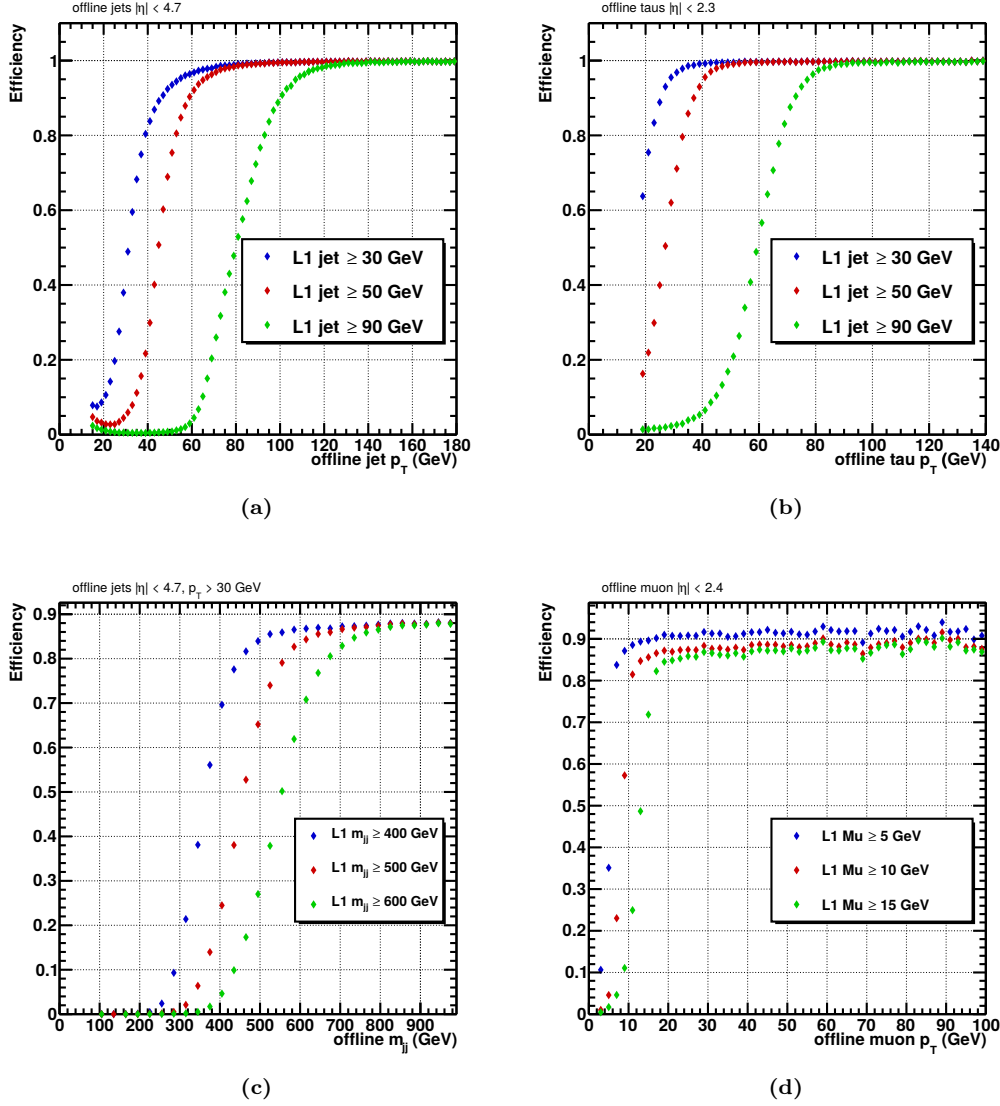


Figure 6.1: The L1 trigger efficiencies for: jets to pass L1 jet p_T cuts (a), taus to pass L1 jet p_T cuts (b), dijets to pass L1 m_{jj} cuts (c), and muons to pass L1 muon p_T cuts (d), are shown as a function of the corresponding offline variables. For each inset the efficiencies are shown for three different L1 cuts, where the magnitudes of these cuts are indicated in the legends.

This extrapolation thus assumes that the rate scales linearly with the luminosity. This assumption does not account for cases where there is a significant PU-dependence, which can cause a non-linear increase in the rates. However, this is usually a reasonable approximation as long as the luminosity per bunch (and thus the PU) for the recorded data is approximately similar to the target conditions. The zero bias data was recorded with $N_b = 517$ bunches compared to the target conditions with $N'_b = 2808$ meaning the luminosity per bunch is ~ 0.8 times the target, and therefore this approximation should hold reasonably well. After the initial optimisation studies the rate was estimated using a different method, to better take into account the extrapolation to higher instantaneous luminosities. This method will be described in Section 6.2.

The optimisation of the trigger selections starts by first performing a coarse scan of the phase space, applying many different combinations of cuts. When performing these scans a target rate

must be chosen. Target rates of 5 kHz and 1 kHz were used for the optimisation of the dijet and muon+dijet triggers respectively. The rate and signal efficiency is calculated for each combination of cuts. All combinations that give a rate less than the target rate are then ranked according to their signal efficiencies. There are usually several combinations of cuts that give approximately similar rates and efficiencies. The best candidates are then chosen based on two criteria:

- The triggers use a minimal number of selections. This is advantageous as it will simplify the measurements of efficiencies and scale factors at a later date.
- The trigger has high efficiency for other Higgs final states. This is studied using samples of $H \rightarrow b\bar{b}$ and $H \rightarrow \text{invisible}$ events.

Based on these two points, the optimal dijet trigger was found to use cuts on three variables: m_{jj} , the p_T of the leading jet, and the p_T of the sub-leading jet. The optimal m_{jj} cut and p_T cuts on the (sub)leading jet were found to be $m_{jj} > 580$ GeV and $p_T > 90(30)$ GeV. The optimal muon+dijet trigger was found to require two jets with $p_T > 30$ GeV and use cuts on m_{jj} and the p_T of the leading muon, where cuts of $m_{jj} > 360$ GeV and $p_T > 6$ GeV were found to give the best performance. The exact thresholds for these cuts were further tuned based on the rate study described in the next section.

6.2 Rate measurements

It became clear during 2016 data-taking that the LHC instantaneous luminosity would likely exceed the target luminosity, reaching luminosities up to $1.5 \times 10^{34} \text{ cm}^2\text{s}^{-1}$. The trigger thresholds were therefore re-tuned to cope with these conditions. Two versions of the triggers were developed, one designed to operate at $L = 1.2 \times 10^{34} \text{ cm}^2\text{s}^{-1}$ and the other at $L = 1.5 \times 10^{34} \text{ cm}^2\text{s}^{-1}$. As before, the trigger thresholds were chosen to give rates of about 5 kHz and 1 kHz for the dijet and muon+dijet triggers respectively.

The L1 rate measurements for the candidate triggers were also performed using a different procedure to the one described in the previous section, to better account for the extrapolation to higher instantaneous luminosities. The rate is measured in bins of the instantaneous luminosity, which varies throughout data-taking, and fitted functions are used to extrapolate the rate measurements to the target luminosities. Zero bias data collected during 2016 data-taking was also used for these measurements. This data was recorded with conditions more similar to those expected for 2017 data-taking.

After re-optimising the trigger thresholds for the higher luminosity conditions, the m_{jj} cut used for the dijet triggers was increased to 620 GeV. The p_T threshold for the leading jet was kept at 90 GeV for the $L = 1.2 \times 10^{34} \text{ cm}^2\text{s}^{-1}$ version of the trigger, but increased to 97 GeV for the $L = 1.5 \times 10^{34} \text{ cm}^2\text{s}^{-1}$ version. These triggers are named "DoubleJet_90_30_Mj30j30_620" and "DoubleJet_97_30_Mj30j30_620", respectively. For muon+dijet triggers, the m_{jj} and muon p_T cuts were increased to cope with the higher luminosity. Both versions of the trigger apply a $m_{jj} > 400$ GeV cut. The $L = 1.2 \times 10^{34} \text{ cm}^2\text{s}^{-1}$ version of the trigger requires the muon to have $p_T > 8$ GeV, which is increased to 9 GeV for the $L = 1.5 \times 10^{34} \text{ cm}^2\text{s}^{-1}$ version. These triggers are named "DoubleJet30_Mj30j30_400_Mu8" and "DoubleJet30_Mj30j30_400_Mu9", respectively.

The rates measured as a function of the instantaneous luminosity are shown in Figure 6.2. For the dijet triggers, the data is found to be well described by linear functions. This shows that these algorithms do not have a significant PU dependence. Conversely, the muon+dijet triggers

are observed to have a slight non-linear increase in the rates, indicating that these algorithms are affected by the increasing PU. Quadratic functions are used to extrapolate the rates in this case. The rates evaluated at the target luminosities are also shown at the bottom right of the figures.

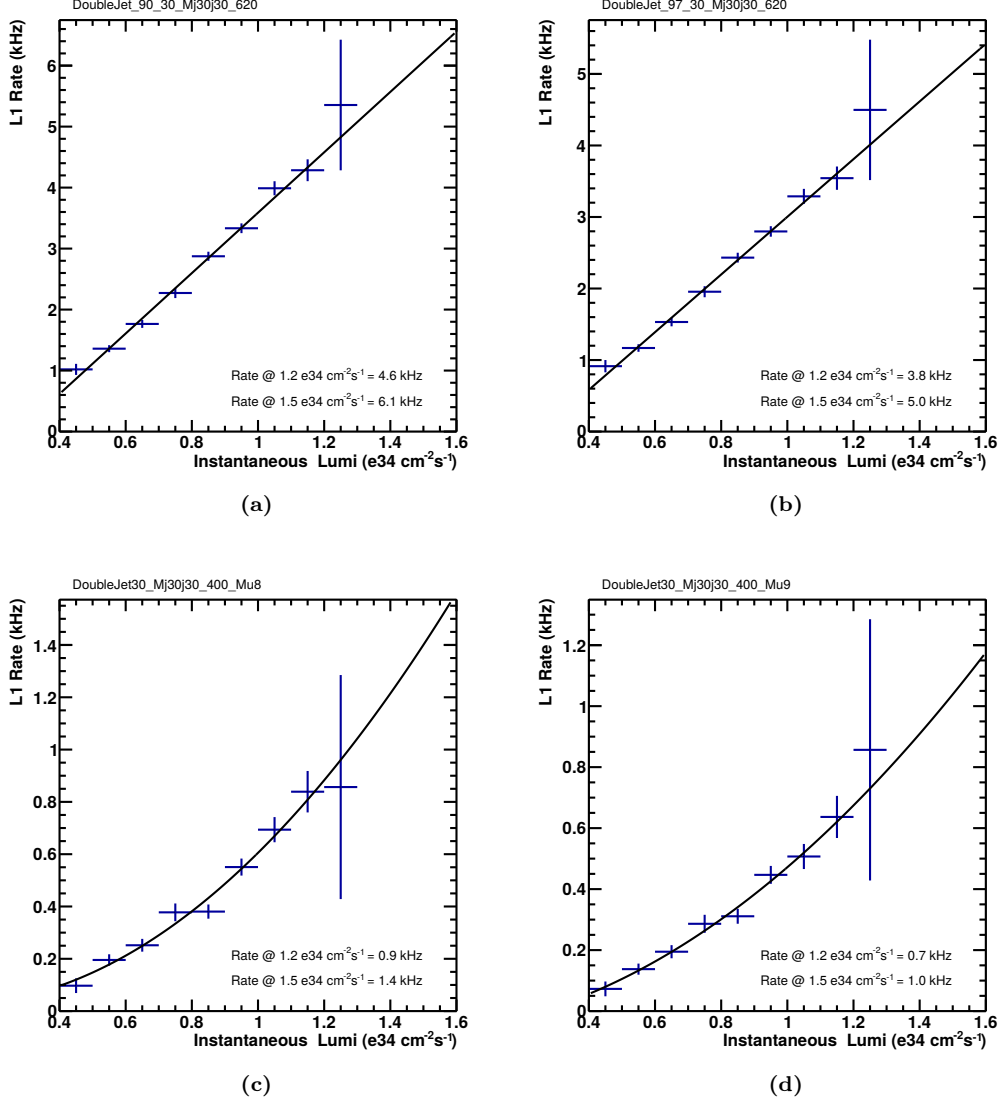


Figure 6.2: The L1 rate as a function of the instantaneous luminosity is shown for the dijet triggers (upper) and muon + dijet triggers (lower). The dijet triggers require a pair of $p_T > 30$ GeV jets with $m_{jj} > 620$ GeV. An additional jet with $p_T > 90$ (97) GeV is also required for the trigger shown in the left (right) inset. The muon+dijet triggers requires a pair of $p_T > 30$ GeV jets with $m_{jj} > 400$ in addition to a muon. The muon is required to have $p_T > 8$ (9) GeV for the trigger shown in the left (right) inset.

6.3 Performance studies

The expected signal efficiencies for the dijet and the muon+dijet triggers are shown in Tables 6.1 and 6.2. The L1 rates evaluated for an instantaneous luminosity of $1.5 \times 10^{34} \text{ cm}^{-2} \text{ s}^{-1}$ are also shown. The dijet triggers have efficiencies up to 24% for $H \rightarrow \tau_h \tau_h$ events. The dijet efficiencies are shown also for the $H \rightarrow \text{invisible}$ and $H \rightarrow b\bar{b}$ processes - two final states that often suffer from low trigger efficiencies. The efficiencies for the $H \rightarrow \text{invisible}$ process are slightly lower than for

$H \rightarrow \tau_h \tau_h$, with 18% of the events passing the lowest threshold trigger. The $H \rightarrow \tau_h \tau_h$ process has a larger efficiency as one of the hadronic taus may be reconstructed as the 90/97 GeV jet. For the same reason, the $H \rightarrow b\bar{b}$ process has the largest signal efficiency, with 30% of the events passing the lowest threshold trigger. In this case the $b\bar{b}$ pair hadronise to form two jets, either of which can pass the 90/97 GeV p_T requirement. These jets tend to have larger p_T than the τ_h candidates from $H \rightarrow \tau\tau$ decays and thus have larger efficiencies.

For the muon+dijet triggers the lowest threshold algorithm has an efficiency of 29% for $H \rightarrow \tau_\mu \tau_h$ events. This is larger than the efficiency for $H \rightarrow \tau_h \tau_h$ events to pass the dijet trigger due to the looser cuts applied on the m_{jj} and jet p_T .

The signal efficiencies for the VBF triggers are comparable to the efficiencies of the other triggers that are usually employed to trigger these events. Much of the phase space targeted by the VBF triggers is also not covered by the other L1 algorithms, and therefore their usage can give significant gains in the signal acceptance, for example:

- $H \rightarrow \tau_h \tau_h$ events are usually selected using a double-tau trigger requiring the presence of two isolated taus at L1. For 2016 data-taking, both L1 taus were required to have $p_T > 28$ GeV and $|\eta| < 2.1$. The efficiency for this double-tau trigger is 33%, which is comparable to the 24% efficiency for the VBF dijet trigger. Using both of these triggers in combination also gives an $\approx 40\%$ gain in the signal acceptance.
- $H \rightarrow \tau_\mu \tau_h$ events are usually selected using a single-muon trigger. The L1 trigger used in 2016 required the muon to have $p_T > 20$ GeV and $|\eta| < 2.1$. The efficiency for this trigger is 41%, which again is comparable to the 29% efficiency for the muon+dijet VBF trigger. Using both of these triggers in combination gives an $\approx 30\%$ gain in signal acceptance.
- The $H \rightarrow$ invisible final state relies on E_T^{miss} triggers, which often suffer from high thresholds. In 2016 the E_T^{miss} threshold used was 80 GeV which has a signal efficiency of 25%. Using both triggers in combination in this case gives an $\approx 30\%$ increase in the signal efficiency.

Table 6.1: The efficiencies of the L1 dijet triggers for three Higgs final states and the rates evaluated for an instantaneous luminosity of $1.5 \times 10^{34} \text{ cm}^2\text{s}^{-1}$.

Trigger	Rate (kHz)	Efficiency		
		$H \rightarrow \tau_h \tau_h$	$H \rightarrow \text{invis.}$	$H \rightarrow b\bar{b}$
DoubleJet_90_30_Mj30j30_620	6.1	24%	18%	30%
DoubleJet_97_30_Mj30j30_620	5.0	22%	17%	27%

Table 6.2: The efficiencies of the L1 muon+dijet triggers for the $H \rightarrow \tau_\mu \tau_h$ final state and the rates evaluated for an instantaneous luminosity of $1.5 \times 10^{34} \text{ cm}^2\text{s}^{-1}$.

Trigger	Rate (kHz)	Efficiency
		$H \rightarrow \tau_\mu \tau_h$
DoubleJet30_Mj30j30_400_Mu8	1.4	29%
DoubleJet30_Mj30j30_400_Mu9	1.0	28%

The dijet and muon+dijet triggers were included in the L1 menu for 2017 and 2018 data-taking. However, no HLT path was developed to pick up the latter and so no data was collected with this trigger. Several HLT paths were developed to pick up events collected by the dijet trigger.

Of relevance to the $H \rightarrow \tau\tau$ analysis, a HLT path requiring the presence of two τ_h candidates with $p_T > 20$ GeV and two jets with $m_{jj} > 650$ GeV was included. This path also required the (sub)leading jet to have $p_T > 115$ (40) GeV.

The potential improvement from using this HLT path is studied using MC simulating 2017 conditions. This study considered the increase in the VBF signal and $Z/\gamma^* \rightarrow \tau\tau$ background events from using the VBF trigger in combination with the double-tau trigger. The signal and background distributions are obtained for events that are triggered exclusively using the double-tau trigger, and for events that are triggered using both the VBF and double-tau triggers. A set of typical offline $\tau_h\tau_h$ selections used to target the VBF phase-space were also applied - the cuts were chosen to be similar to those used for previous CMS results such as in Ref. [191]. These offline selections require both taus to have $p_T > 20$ GeV and pass the *tight* MVA-isolation, the *loose* anti-muon discriminant, and the *very-loose* anti-electron discriminant. The events are also required to contain at least two jets with $m_{jj} > 500$ GeV, and events with additional leptons are vetoed. An additional 40 GeV p_T cut is applied to τ_h candidates passing the double-tau trigger. Events passing the VBF trigger must contain two jets with $p_T > 120$ and 45 GeV and $m_{jj} > 700$ GeV. These selections are used to approximate the typical offline cuts that may be required to select events above the trigger turn-on regions - these selections are often necessary as modelling the efficiencies accurately below the turn-ons can be difficult.

The signal and background distributions of the invariant mass of the visible decay products, $m_{\tau\tau}^{\text{vis}}$, are shown in Figure 6.3. The distributions are shown for events triggered using the double-tau trigger exclusively, and events triggered using a combination of the VBF and double-tau triggers. The lower inset of Figure 6.3 also shows the $\frac{S}{\sqrt{B}}$ for both trigger strategies. The use of the VBF trigger in combination with the double-tau trigger is observed to give a significant improvement in the $\frac{S}{\sqrt{B}}$ for most $m_{\tau\tau}^{\text{vis}}$ bins. The double-tau trigger is found to have a signal efficiency of 19%, which is increased to 30% by the VBF trigger. This corresponds to a 58% gain in the signal yield. The double-tau trigger efficiency for $Z/\gamma^* \rightarrow \tau\tau$ events is 11%, which is increased to 20% by the VBF trigger. This corresponds to a 82% increase in the background yield. The reason the background yield increases more than the signal is due to the softer p_T distributions for τ_h candidates from $Z/\gamma^* \rightarrow \tau\tau$ events (mainly due to m_Z being smaller than m_H). Despite the larger increase in the background, the total $\frac{S}{\sqrt{B}}$ is increased from 0.96 to 1.11 by the inclusion of the VBF trigger, which corresponds to a $\approx 16\%$ gain in the sensitivity.

It should be noted that this study may give a slightly over-optimistic view for the following reasons:

1. The VBF trigger was not included for the first $\approx 15 \text{ fb}^{-1}$ of data-taking in 2017.
2. While the study does consider the additional offline cuts that may be required when using the VBF trigger, these thresholds may need to be increased depending on how accurately the trigger efficiency can be modelled close to the turn-ons.
3. The LHC instantaneous luminosity exceeded the design luminosity of the trigger. For example, in 2018 data was collected at a peak luminosity of $\approx 2 \times 10^{-34} \text{ cm}^2\text{s}^{-1}$. As such, the L1 thresholds had to be increased during data-taking. The exact thresholds used depend on the run period, but the p_T selections were as large as 115 (40) GeV for the (sub)leading jet. The m_{jj} cut was not increased during data-taking. However, tighter online selections also had to be used for the double-tau trigger and, as it is not clear a priori which trigger was most affected, it is difficult to conclude if the relative gains will be the same throughout 2017/2018.

6.4 Other uses of the L1 triggers

The dijet trigger has also been employed by several other analyses, these includes:

- The $H \rightarrow \text{invisible}$ analysis. The analysis sensitivity suffers quite heavily from high thresholds used in the E_T^{miss} triggers. A HLT path seeded by the L1 dijet trigger, requiring two jets and E_T^{miss} was developed. The E_T^{miss} thresholds used were softer than those used in the usual E_T^{miss} triggers, and as such the signal acceptance is increased by the inclusion of this trigger.
- Supersymmetry searches for VBF production of electroweakinos. The decay chains of the electroweakinos include soft leptons and E_T^{miss} . The leptons are usually too soft to pass the single- and double-lepton trigger selections and so the previous strategy relied on jets/ E_T^{miss} and muon+ H_T triggers (H_T is defined as the scalar p_T sum of a collection of jets). These triggers suffer from high $p_T/E_T^{\text{miss}}/H_T$ thresholds and only consider jets up to $|\eta| < 3$, which is problematic for VBF events as at least one jet usually has $|\eta| > 3$. The inclusion of the L1 dijet trigger as a seed for a dedicated HLT path requiring jets, a soft muon, and E_T^{miss} , was able to increase the signal efficiency for this analysis.
- The search for heavy Higgs bosons decaying to two SM Higgs bosons in the $H \rightarrow hh \rightarrow \tau\tau b\bar{b}$ final state. The VBF production mode is one of the signal processes considered for this analysis. This analysis uses the HLT path developed for the $H \rightarrow \tau\tau$ analysis described above. The increase in signal efficiency can be large for this process as the jets from the $h \rightarrow b\bar{b}$ decay can also pass the high- p_T jet requirement.

The wide range of uses found for this trigger motivates the continued development and usage of such algorithms in the L1 trigger.

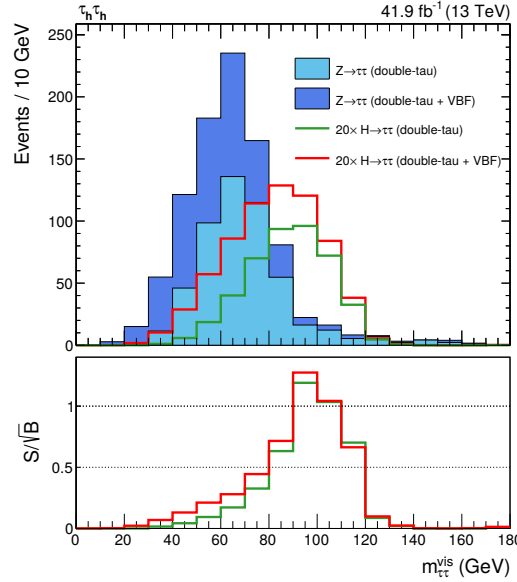


Figure 6.3: Upper inset: The $m_{\tau\tau}^{\text{vis}}$ distributions in the $\tau_h\tau_h$ channel are shown for signal and background events passing the double-tau trigger, and events passing a combination of the double-tau and VBF triggers. The signal is scaled by a factor of 20. Both signal and background histograms are shown overlaid rather than stacked. For the $Z/\gamma^* \rightarrow \tau\tau$ background, the events passing the double-tau trigger and the combination of the double-tau and VBF triggers are shown by the light blue and dark blue histograms respectively. For the $H \rightarrow \tau\tau$ signal, the events passing the double-tau trigger and the combination of the double-tau and VBF trigger are shown by the green and red lines respectively. Lower inset: The $\frac{S}{\sqrt{B}}$ is shown for events selected using the double-tau trigger (green) and events selected using both double-tau and VBF triggers (red).

Chapter 7

Conclusions

This thesis has presented two analyses of $\sqrt{s} = 13$ TeV proton-proton collision data collected by CMS using final states with tau leptons. The purpose of these analyses was to probe the Higgs sector, which can be achieved either by searching for additional Higgs bosons or by measuring properties of the observed 125 GeV state. Both of these options were explored in the preceding chapters.

The search for additional Higgs bosons decaying to ditau pairs was performed using 35.9 fb⁻¹ of data collected in 2016. This analysis uses a categorisation based on the number of b-tagged jets to target the GF and $b\bar{b}$ associated processes, which are the dominant production modes at the LHC. This search is well motivated for constraining the MSSM parameter space, as the couplings of the H and A bosons to tau leptons are enhanced for large values of $\tan\beta$. No statistically significant deviation with respect to the background-only model was observed and so exclusion limits are set for both model-independence and model-dependent interpretations. For the model-independent interpretation, limits on the cross section times branching ratios for the two processes are set at the 95% CL. For the model-dependent interpretation, exclusion limits are set in the m_A - $\tan\beta$ plane for two MSSM benchmark scenarios: $m_h^{\text{mod+}}$ and hMSSM. These results were published in Ref. [1]. A further publication is also planned in 2020 including the datasets collected in 2017 and 2018, corresponding to an additional integrated luminosity of 100.7 fb⁻¹.

A measurement of the $\mathcal{CP}$ properties of the Higgs Yukawa coupling to top quarks is also presented. This analysis combines the datasets collected in 2016 and 2017, which corresponds to an integrated luminosity of 77.8 fb⁻¹. This constitutes the first measurement of the $\mathcal{CP}$ properties of a fermionic Higgs coupling. Such measurements are well motivated as the $\mathcal{CP}$ -even and $\mathcal{CP}$ -odd components of a mixed $\mathcal{CP}$ -state can couple to fermions with comparable strengths, and therefore $\mathcal{CP}$ -violating effects can be large. The analysis adopts the parametrisation of Ref. [57] which quantifies the level of $\mathcal{CP}$ -violation in terms of a mixing angle α_{hgg} . The best-fit value of α_{hgg} is measured to be $-59^{+47}_{-50}^\circ$.

The inclusion of the dataset collected in 2018 which corresponds to 58.8 fb⁻¹ of additional data is expected to improve the sensitivity to α_{hgg} by $\mathcal{O}(30\%)$. The current analysis also considered only $H \rightarrow \tau\tau$ decays, which are expected to have the best sensitivity to the dijet phase space. It should be stressed however that the measurement can in principle be performed using any Higgs final state, and a combination including several Higgs channels will improve the overall sensitivity. A publication of the α_{hgg} measurement is planned after the dataset collected in 2018 has been analysed. This publication will also include a combination with the $H \rightarrow ZZ$ channel, which is expected to exclude scenarios with $\alpha_{\text{hgg}} \gtrsim 80^\circ$ at the 95% CL. A combination of the $H \rightarrow \tau\tau$,

$H \rightarrow ZZ$, and $H \rightarrow \gamma\gamma$ analyses is also planned by early 2020. A combination between these three channels will likely be able to exclude scenarios with $\alpha_{\text{hgg}} \gtrsim 45^\circ$ at the 95% CL.

Finally, the development of L1 triggers designed to target the VBF production process was discussed. A trigger that relied only on features of the two tagging jets was included in the L1 trigger menu for 2017 and 2018 data-taking. The use of this trigger was shown to increase the signal yield for $\tau_h\tau_h$ events significantly. This channel is expected to be the most important for measuring the $\mathcal{CP}$ -properties of the $H\tau\tau$ vertex, which complements the measurements of the $Ht\bar{t}$ vertex discussed previously. Such measurements are expected to reach appreciable sensitivities with the full Run 2 dataset.

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List of Acronyms

2HDM two Higgs doublet model	31
AMS asymptotic median significance	133
AR application-regions	82
BSM beyond the standard model	19
CHS charged hadron subtraction	55
CL confidence level	31
CSC cathode strip chambers	42
CSV combined secondary vertex	57
CTF Combinatorial Track Finder	47
DR determination-regions	82
DT drift tube	42
EB ECAL barrel	40
ECAL electromagnetic calorimeter	39
EE ECAL endcaps	40
EFT effective field theory	125
FF fake-factor	82
GF gluon fusion	22
GSF Gaussian sum filter	50
GUT grand unified theory	24
HB hadron barrel	41
HCAL hadronic calorimeter	41
HE hadron endcap	41
HF hadron forward	41
HLT high-level trigger	45

HO hadron outer	41
HPS hadrons plus strips	58
IVF inclusive vertex finding	57
KF Kalman filter	47
L1 Level-1	44
LEP Large Electron-Positron	29
LHC Large Hadron Collider	22
LO leading order	62
LS1 Long Shutdown 1	35
MC Monte Carlo	62
MPF missing transverse energy projection fraction	56
MSSM minimal supersymmetric standard model	19
N3LO next-to-next-to-next-to-leading order	122
NLO next-to-leading order	62
NNLO next-to-next-to-leading order	62
PDFs parton distribution functions	62
pdfs probability density functions	107
PF particle flow	48
PSB Proton Synchrotron Booster	35
PU pile-up	38
QCD quantum chromodynamics	19
QFT quantum field theory	19
RF radio frequency	35
RPC resistive plate chambers	42
SC supercluster	52
SM standard model	19
SPS Super Proton Synchrotron	35
SR signal-regions	82
SUSY supersymmetry	24
TEC tracker endcaps	39
TIB tracker inner barrel	38

TID tracker inner disks	38
TOB tracker outer barrel	39
UE underlying event	55
VBF vector boson fusion	22
VEV vacuum expectation value	21
VR validation-region	88
WLCG Worldwide LHC Computing Grid	45