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SEARCH FOR EXOTIC HIGGS BOSON DECAYS TO
FOUR LEPTONS WITH THE ATLAS DETECTOR

BY
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DISSERTATION

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Abstract

Searches for exotic Higgs boson decays are well motivated by various theoretical models, as well as the constraints from Higgs boson measurement results. The focus is on exotic Higgs boson decays to a pair of beyond-the-Standard-Model spin-0 particles a . This dissertation conducts two analyses for exotic Higgs bosons decaying into leptons, which are $H \rightarrow 2a \rightarrow 4\mu$ and $H \rightarrow 2a \rightarrow 4\tau$, and target the mass ranges of $1 < m_a < 15$ GeV and $15 < m_a < 60$ GeV. Various data-driven techniques are developed and used to estimate the background that cannot be well modeled with the standard simulation in ATLAS. The searches use proton-proton collision data collected with the ATLAS detector at the LHC with an integrated luminosity of 139 fb^{-1} and a center of mass energy $\sqrt{s} = 13$ TeV. No significant excess has been observed. Therefore, this dissertation summarizes the 95% confidence-level upper limits of the branching ratio of exotic Higgs decays to a pair of beyond-the-Standard-Model light scalars in the lepton final states.

For the time passed by.

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List of Abbreviations

2HDM two-Higgs-doublet model. vi, 28, 32, 33, 34, 35, 36

2HDM+S two-Higgs-doublet plus singlet model. vi, ix, xi, 28, 34, 35, 36, 37, 76, 95, 124, 132

BCS theory Bardeen-Cooper-Schrieffer theory. 8

BDT boosted decision tree. 71, 72, 108, 119

Br branching ratio. ix, x, xi, 1, 25, 31, 32, 35, 36, 37, 70, 76, 93, 94, 95, 100, 103, 104, 119, 124, 132

BRT boosted regression tree. 71

BSM beyond the Standard Model. 1, 23, 28, 31, 32, 40, 120

BW Big Wheel. 53

CB combined muon. 67, 69, 73, 119

CERN Organisation européenne pour la recherche nucléaire(European Organization for Nuclear Research).
40, 56, 63

CL confidence level. vii, x, xi, 16, 31, 34, 76, 88, 92, 94, 95, 119, 124, 132

CM center of mass. ix, 1, 10, 25, 40, 43, 44, 45, 60, 132

CP combination of charge conjugation symmetry and parity symmetry. 10

CR control region. xi, 80, 89, 100, 101, 102, 108, 109, 110, 112, 113, 114, 116, 117, 118, 119

CSC cathode strip chambers. 54, 56, 68

CT calorimeter-tagged muon. 68

CTP central trigger processor. 56

DIS deep inelastic scattering. ix, 17

DM dark matter. vi, 27, 31, 38

EMB LAr electromagnetic barrel. 51, 52, 53

EMCal electromagnetic calorimeter. 51, 52, 53, 65, 69, 72

EMEC LAr electromagnetic endcap calorimeter. 51, 52, 53

EWK electroweak. vi, ix, 9, 11, 13, 14, 19, 20, 22, 27, 80, 103, 104, 110, 116

FCal forward calorimeter. 51, 52, 53

FCNC flavor-changing neutral currents. 34

FF fake factor. vii, viii, xi, 1, 103, 104, 106, 108, 109, 110, 111, 110, 111, 112, 113, 114, 116, 117, 118, 132

ggF gluon-gluon fusion. x, 1, 23, 32, 76, 94, 98, 99, 124

GRL Good Run List. 57, 58

GUT Grand Unification Theory. 28

HCal hadronic calorimeter. 51, 52, 53, 69, 71, 72

HEC LAr hadronic endcap calorimeter. 51, 52, 53

HL-LHC High-Luminosity LHC. vi, ix, 45, 58, 132

HLT high-level trigger. 56

hMSSM habemus minimal supersymmetric standard model. ix, 34

IBL inner b-layer. 65

ID inner detector. vi, ix, 49, 50, 65, 67, 68, 69, 72, 87

ip interaction point. 44, 48, 58, 60, 68, 71

IP impact parameter. 49, 65, 72, 78, 80, 82, 84, 85, 100, 110

isol isolation. viii, x, 66, 67, 69, 77, 78, 80, 82, 84, 85, 87, 100, 108, 110

JES jet energy scale. [72](#)

JVT jet vertex tagger. [72](#)

L1 first-level trigger. [xiv](#), [56](#), [58](#)

L1A [L1](#) accpet. [56](#)

L1Calo [L1](#) calorimeter trigger. [56](#)

L1Muon [L1](#) muon trigger. [56](#), [58](#)

LAr liquid argon. [xiii](#), [52](#), [53](#), [57](#), [65](#)

LEP Large Electron-Positron Collider. [14](#)

LH likelihood-based. [65](#), [67](#), [71](#)

LHC Large Hadron Collider. [vi](#), [viii](#), [ix](#), [x](#), [1](#), [4](#), [16](#), [17](#), [20](#), [25](#), [26](#), [28](#), [31](#), [32](#), [34](#), [40](#), [43](#), [44](#), [45](#), [48](#), [56](#), [57](#), [58](#), [60](#), [65](#), [73](#), [94](#), [98](#), [102](#), [124](#), [132](#)

Linac2 Linear accelerator 2. [40](#), [43](#)

LO leading order. [104](#)

LQ leptoquark. [120](#)

LS long shutdown. [40](#)

MC Monte Carlo method. [45](#), [57](#), [60](#), [67](#), [69](#), [73](#), [79](#), [87](#), [88](#), [92](#), [103](#), [106](#), [109](#), [110](#), [112](#), [116](#), [117](#), [118](#), [132](#)

MDT monitored drift tubes. [54](#), [56](#), [68](#), [69](#)

ME extrapolated muon. [68](#), [69](#)

MicroMegas Micro-MEsh Gaseous Structure. [vii](#), [viii](#), [x](#), [40](#), [58](#), [60](#), [62](#), [63](#)

MS muon spectrometer. [vi](#), [40](#), [53](#), [54](#), [56](#), [58](#), [67](#), [68](#), [69](#), [87](#)

MSSM minimal supersymmetric standard model. [34](#)

MVA multivariate analysis. [65](#)

NLO next-to-leading order. [19](#), [76](#), [99](#), [104](#)

NMSSM next-to-minimal supersymmetric standard model. [34](#), [37](#)

NNLO next next leading orders. [ix](#), [16](#)

NP nuisance parameters. [vii](#), [89](#), [90](#)

NSW New Small Wheel. [vi](#), [vii](#), [x](#), [40](#), [58](#), [60](#), [62](#)

OF opposite flavor. [110](#)

OLR overlap removal. [71](#)

OS opposite-sign. [xi](#), [77](#), [98](#), [102](#), [108](#), [109](#), [110](#), [112](#), [113](#), [114](#), [116](#), [117](#)

p.d.f. probability distribution function. [65](#), [91](#), [92](#)

PDF Parton Distribution Functions. [vi](#), [ix](#), [15](#), [16](#), [17](#), [76](#), [88](#)

PDG particle data group. [22](#)

PLR profile likelihood ratio. [90](#), [91](#), [92](#)

PNGB pseudo Nambu-Goldstone boson. [34](#)

POI parameter of interests. [88](#), [90](#)

pp proton-proton. [ix](#), [1](#), [43](#), [44](#), [45](#), [56](#), [60](#), [65](#), [98](#), [132](#)

PS Proton Synchotron. [40](#), [43](#)

PSB Proton Synchotron Booster. [40](#), [43](#)

QCD quantum chromodynamics. [vi](#), [ix](#), [4](#), [7](#), [9](#), [14](#), [17](#), [18](#), [19](#), [20](#), [88](#), [98](#)

QED quantum electrodynamics. [vi](#), [4](#), [7](#), [9](#), [10](#), [11](#), [17](#), [18](#), [19](#)

QFT quantum field theory. [3](#), [4](#), [9](#), [19](#), [27](#)

ROC receiver operating characteristic. [xi](#), [71](#), [120](#)

RPC resistive plane chambers. [54](#), [56](#)

SCT semiconductor Tracker. [49](#), [50](#), [57](#)

sf scale factor. 87, 117, 118

SF same flavor. 102, 108, 110

SM Standard Model. vi, viii, ix, xi, 1, 3, 4, 5, 7, 8, 20, 22, 23, 25, 26, 27, 28, 31, 32, 31, 32, 33, 34, 35, 36, 37, 38, 72, 76, 79, 80, 87, 89, 92, 94, 98, 102, 103, 110, 111, 119, 124, 132

SPS Super Proton Synchotron. 40, 43

SR signal region. viii, x, xi, 80, 88, 89, 93, 94, 95, 100, 101, 102, 104, 116, 117, 118, 119

SS same-sign. 98, 100, 102, 101, 103, 111, 112, 113, 116, 118, 124, 132

SSB Spontaneous Symmetry Breaking. vi, 5, 7, 8, 14, 20

ST segment-taged muon. 68

sTGC small-strip TGC. 58

SUSY Supersymmetry. ix, 28, 34

SW Small Wheel. 53, 58

TDAQ trigger and data acquisition. vi, x, 56

TGC thin gap chambers. xvi, 54, 56, 58

Tile Tile calorimeter. 51, 52, 53, 57

TRT transition radiation tracker. 50

TST track-based soft term. viii, 73

TTVA track-to-vertex association. 66, 69, 87

v.e.v vacuum expectation value. 7, 8, 22, 23, 33, 36

VBF vector-boson fusion. 32

VR validation region. xi, 88, 100, 101, 112, 113, 114, 116, 117, 118, 119

Chapter 1

Introduction

The [Standard Model](#) is the most successful achievement in particle physics, but also known as an incomplete model which does not account for many unexplained questions. As experimentalists, there are usually two approaches to look for evidence of [beyond the Standard Model](#) physics. One is to measure the known processes with high precision, and compare to the prediction of [Standard Model](#). Another approach is to search for direct evidence of new physics. The search for exotic Higgs decays is the latter approach. Since the Higgs boson discovery in 2012 [1] [2], the properties of the Higgs boson have been studied by many efforts. The combination of different measurements of the Higgs boson properties reveals that the Higgs boson may have exotic decays, as much as the [branching ratio](#) of $\text{Br}(H \rightarrow \text{BSM}) \leq 21\%$ [3]. As one of the simplest extensions of the [SM](#), many theoretical models predict Higgs bosons decays to a new spin zero particle a , which serves as the portal linking the [Standard Model](#) and new physics.

This dissertation presents two searches for exotic Higgs decays to leptons, which are $H \rightarrow 2a \rightarrow 4\mu$ and $H \rightarrow 2a \rightarrow 4\tau$, targeting the mass range of $1 < m_a < 15$ GeV and $15 < m_a < 60$ GeV respectively. These analyses use [proton-proton](#) collision data collected from 2015 to 2018 with the ATLAS detector at [center of mass](#) energy $\sqrt{s} = 13$ TeV and integrated luminosity of $L_{\text{int.}} = 139 \text{ fb}^{-1}$. Both analyses focus on the [gluon-gluon fusion](#) production mode of the Higgs boson which is the primary production mode at the [Large Hadron Collider](#). Processes with four leptons in the final state are rare in the [Standard Model](#), which means a small amount background and high sensitivity. The major challenge in these analyses is the estimate of [Standard Model](#) processes with misidentified objects, so called fakes. These processes are usually hard to model in a simulation, suffering from low statistics and unreliable modeling of the detector response. Hence, various data-driven techniques are used for the fake estimation. In the $H \rightarrow 2a \rightarrow 4\mu$ analysis, a dedicated template method is developed to model the fake muons from semileptonically decaying hadrons. And in the $H \rightarrow 2a \rightarrow 4\tau$ analysis, an inclusive [fake factor](#) method is used to estimate fakes with complex compositions.

The thesis is organized with the following structure. Chapter 2 introduces the [Standard Model](#) following the history of its establishment. Chapter 3 describes the motivations for the searches for exotic Higgs decays. The detector is described in Chapter 4 with studies on the detector upgrades. Chapter 5 describes the method

to reconstruct and identify for the physics objects used in the analyses. The analyses of $H \rightarrow 2a \rightarrow 4\mu$ and $H \rightarrow 2a \rightarrow 4\tau$ are presented in Chapter 6 and Chapter 7 respectively. Chapter 8 discusses the interpretation of their results, and also summarizes the other exotic Higgs decay searches in ATLAS. And finally, the conclusions are provided in Chapter 9.

Chapter 2

A brief history of Standard Model

Particle physics is one of the most important branches of modern physics, which is to study and understand the nature of most fundamental constitutions of matter and radiation. The main research objects of particle physics are elementary particles, or fundamental particles, which are subatomic particles presumably with no substructure. Kickstarted by J. J. Thomson discovering the first elementary particle, the *electron* through the famous cathode ray experiment in 1897 [4], particle physics theorists and experimentalists have been working for generations to develop a model that describes all the fundamental particles and their interactions (see Figure 2.1). This model is now known as **Standard Model (SM)**.

Standard Model is a renormalizable **quantum field theory (QFT)** framework. According to **QFT**, every particle is an excitation of a specific field in the spacetime [9]. And due to the natural of the *spin-statistics theorem* [10], particles are broadly classified by their spins. Fermions, which are the particles that constitute the matter, are particles with half-integer spin. Bosons, the particles carrying the interactions, are those with integral spin.

Fermions come in two varieties, *quarks* and *leptons*. The quarks are typically represented in the following matrix: $\begin{pmatrix} u & c & t \\ d & s & b \end{pmatrix}$, where the six known species of quarks are: *up* (u), *down* (d), *charm* (c), *strange* (s), *top* (t), and *bottom* (b). In this matrix, each column is a grouping called a *generation*, ordered by the masses of the constituent quarks. And the individual 2×1 columns are called *doublets*. The quarks in the top row are called *up-type* quarks with electric charge $Q = +\frac{2}{3}$, and those in the bottom row are called *down-type* quarks with electric charge $Q = -\frac{1}{3}$. Similarly to the quarks, there are six known leptons arranged in a matrix as: $\begin{pmatrix} \nu_e & \nu_\mu & \nu_\tau \\ e & \mu & \tau \end{pmatrix}$, where the components are: *electron* (e), *muon* (μ), *tau* (τ), each has electric charge $Q = -1$, and the associated *neutrinos* (ν) with charge $Q = 0$. Every generation of leptons carries a specific *flavor*, which plays an important role in the *weak interactions*.

Bosons are categorized as *vectors* and *scalars*, with spin-0 and spin-1 respectively. The known vectors are *photon* (γ), *gluon* (g), $W^\pm$ *boson* and Z *boson*. Vector bosons are also known as *gauge bosons*, which means that each of them is corresponding to a *gauge symmetry* and mediating a fundamental force, as shown in

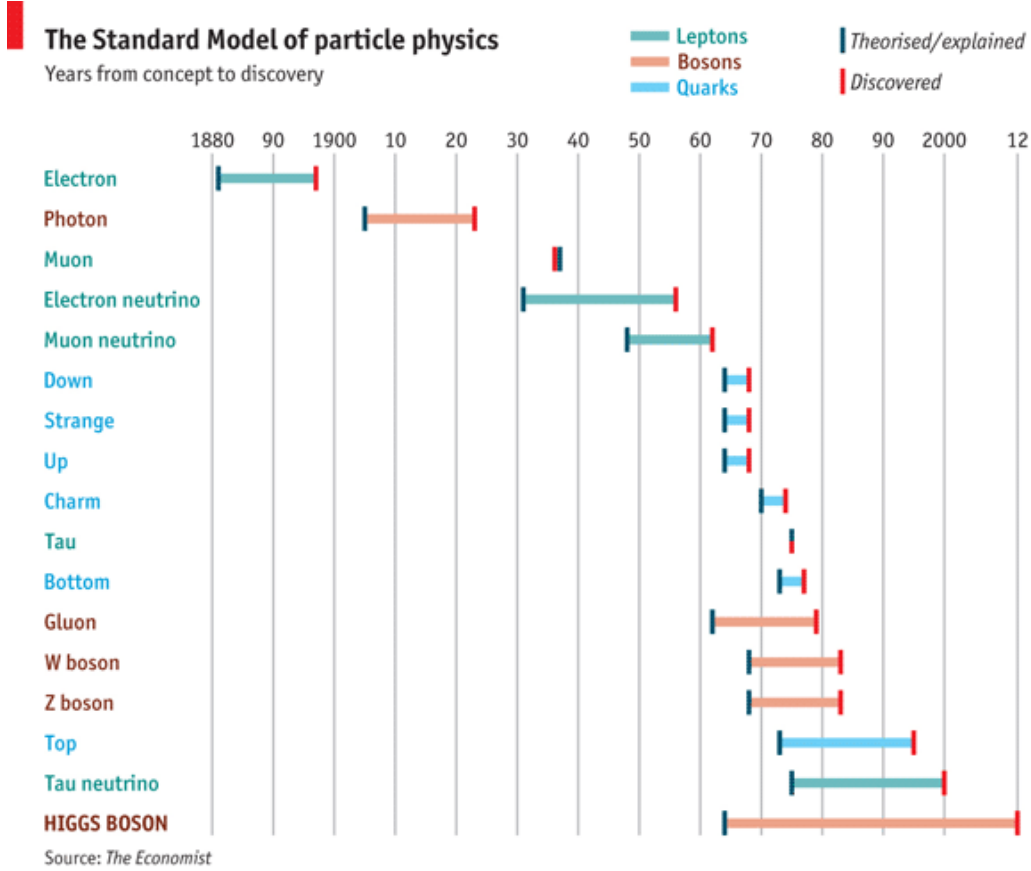


Figure 2.1: A demonstration of timeline summarizing the history of particle physics [5], starting from the discovery of electron in late 19th century, to the most recent discovery of Higgs boson, which was predicted in the 60s of 20th century [6] and eventually observed in 2012 [7, 8]. This diagram also shows how the theoretical and experimental physicists work together closely and contribute to the development of the SM.

Table 2.1. Photons are responsible for the *electromagnetic* force and interacting with all charged particles. Gluons are the particles that bind the quarks together to form *hadrons* and mediate the *strong* force. And the $W^\pm$ and Z bosons are the carriers of the *weak* force. At present, the only scalar boson has been found is the *Higgs boson* [7] [8]. This boson is the observable excitation of the *Higgs field*, which is responsible for providing masses to $W^\pm$ and Z bosons and all fermions [6].

Figure 2.2 summarizes the properties of all known fundamental particles of SM. The precision of the SM have been tested by many large particle experiments, such as AGS, SLAC, E288, PETRA, UA1, D0, CDF, DONUT, and the most recent experiments at the *Large Hadron Collider (LHC)*. All of these experiments measure the cross-sections for different processes and show excellent agreements to the predictions from the SM. However, there are still many other puzzles that the SM needs to be reconciled. This chapter will provide a brief walk through of the SM from a historical approach. Section 2.1 will illustrate the principles of symmetry and symmetry breaking in QFT. Section 2.2 will and Section 2.3 will describe the two major

Interaction	Acted on	Experiencing particles	Mediating particles	Strength at 10^{-18} m
Strong	Color charge	Quarks and gluons	Gluons	25
Electromagnetic	Electric charge	Charged particles	Photons	1
Weak	Flavor	Quarks and leptons	$W^\pm/Z$ bosons	0.8
Gravitational	Mass	Non-massless particles	Graviton	10^{-14}

Table 2.1: The summary of the four known fundamental interactions and their properties [11]. The strength are shown relative to the strength of the electromagnetic force for two up quarks separated by 10^{-18} , which is a typical situation for quarks inside a proton. As can be seen, even though the gravitational force has not been regulated in SM yet, its strength is very tiny and cannot give any impact on the measurement of the SM.

pieces of the SM, the quantum electrodynamics (QED) and the quantum chromodynamics (QCD). Finally, Section 2.4 will introduce the Higgs mechanism and the Higgs field.

2.1 Symmetry and Spontaneous Symmetry Breaking

2.1.1 Global and local symmetry

The dynamic properties of a physics system are often described by *symmetry principles*. In the mathematical language, a symmetry is a feature of the Lagrangian that preserves invariance under some mathematical transformation, or so called operation. For a very simple example, in a classic system, the *Newton's equation* $\mathbf{F} = m \frac{d^2 \mathbf{x}}{dt^2}$ keeps unchanged after the *Galilean transformation* $\mathbf{x} \rightarrow \mathbf{x} + \mathbf{v}t$ [13], where $\mathbf{v}$ corresponds to the relative motion of the observer. Similarly, the *Maxwell's equation* [14] keeps invariant under the *Lorentz transformation* [15]. Mathematically, the symmetries are described by *group theory*, like the Galilean invariance is corresponding to *Galilean group*, and the Lorentz invariance is introduced by *Lorentz group* [16].

The symmetries in physics are broadly classified as *local symmetry* and *global symmetry*. Globally symmetry refers to the transformations that applies to all the points in spacetime, such as the Lorentz transformations. According to *Noether's theorem*, each global symmetry has a corresponding conservation law [17], such as Lorentz spacetime translation invariance corresponds to the conservation of energy and momentum, and the Lorentz rotation invariance corresponds to the angular momentum conservation. Unlike the globally symmetry that appears as invariance under the change of the global spacetime coordinate, local symmetry is independent of the spacetime coordinates and may be recognized as a redundancy in the Lagrangian expressions. For example, the one-dimensional Lagrangian of an object with classic gravitational potential can be written as $L = \frac{1}{2}m\dot{x}^2 - mg(h + x)$. The choice of zero potential energy surface h can effect the expression of the Lagrangian, but preserve the results of the *Euler-Lagrange equation* $\frac{\partial L}{\partial x} - \frac{d}{dt}(\frac{\partial L}{\partial \dot{x}}) = 0$ as simple as $\ddot{x} = g$. More explicitly, in electromagnetism, the classic Lagrangian $L = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} + \frac{1}{c}j^\mu A_\mu$ (the

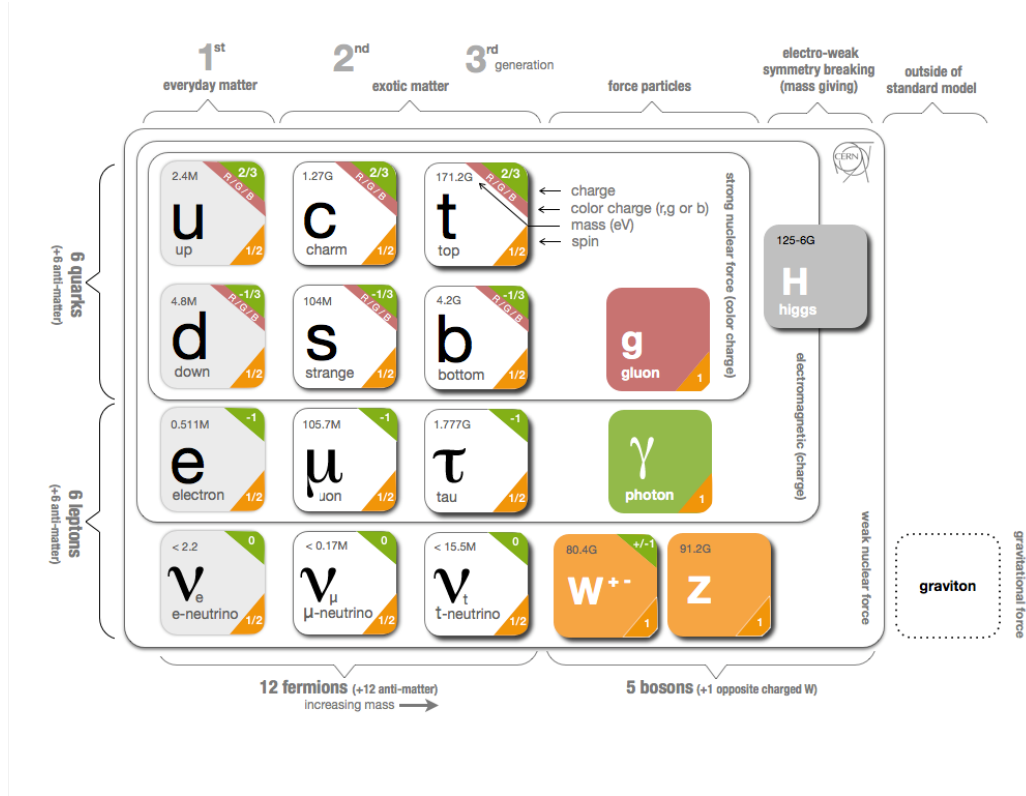


Figure 2.2: A diagram of the [Standard Model](#) of particle physics [12]. Shown are the three generations of six quarks and six leptons, which have the spin of $\frac{1}{2}$. The discovered bosons are shown as: gluon, photon, W/Z boson, and the Higgs boson, as well as the undiscovered graviton, which is a predicted mediator of gravitational force not yet covered by the [Standard Model](#).

covariant form widely used in classic field theory) varies with different choices of the vector potential, from $A_\mu \rightarrow A_\mu + \partial_\mu \lambda$, but the Euler-Lagrangian equation maintains the same and always gives the Maxwell's equations [18]. In 1940s, H. Weyl firstly introduce the idea of *Eichinvarianz* (in German) or translated as *gauge invariance* [19], which means that the local symmetry can also be interpreted as an extrapolation of the relativistic principle in *the theory of general relativity* [20]. A more elegant demonstration is explained by W. Pauli as a phase transform by U(1) gauge group [21].

The gauge invariance was once emphasized during the boom of the quantum mechanics. In 1928, the first relativistic wave equation is derived by P. Dirac, which is known as the *Dirac equation* [10], simply a covariant version of the *Schrödinger equation*. According to the Dirac equation, the *Lagrangian density* ($L = \int \mathcal{L} d^3\mathbf{x}$) of a massive fermion is written as

$$\mathcal{L}_{\text{Dirac}} = \bar{\psi}(i\partial - m)\psi \quad (2.1)$$

where $\partial = \gamma^\mu \partial_\mu$ is the *Feynman notation*, γ^μ are four *Dirac matrices* satisfying the anti-commutation

relation $\{\gamma^\mu, \gamma^\nu\} = \gamma^\mu \gamma^\nu + \gamma^\nu \gamma^\mu$, ψ is the *Dirac spinor*, and m is the fermion mass. This Lagrangian strictly follows the Lorentz invariance, but no longer keeps invariant under the U(1) gauge transform, which takes $\psi \rightarrow \psi' = e^{i\alpha(x^\mu)}\psi$ and the scalar parameter $\alpha(x^\mu)$ is spacetime dependent. To solve this, the derivative is modified to be the *covariant derivative*, as

$$\partial_\mu \rightarrow D_\mu = \partial_\mu + igA_\mu \quad (2.2)$$

where g is an arbitrary real coefficient and A_μ is a vector field satisfying $\nabla \cdot \mathbf{A} = 0$. This property of A_μ allows it can be written as an arbitrary form of $A_\mu \rightarrow A_\mu + \partial_\mu \alpha(x^\mu)$, where $\alpha(x^\mu)$ here is an arbitrary spacetime dependent scalar. But if one realize that $\nabla \cdot \mathbf{A} = 0$ is exactly the form of *Coulomb gauge* and take $g = e$ as the electronic charge, then the vector field A_μ could be interpreted as the electromagnetic four-potential. [QED](#) is now mostly derived from scratch (soon can be seen in Section 2.2)!

Inspired by the correlation between the U(1) gauge invariance and [QED](#), more effects have been taken into the development of gauge theories. In 1954, C. N. Yang and R. Mills extrapolate the Abelian gauge theory to non-Abelian gauge theory based on SU(N) group, which is known as the *Yang-Mills theory* [22]. The [SM](#) is effectively a combination of three gauge theories [23]:

$$\underbrace{\text{SU}_C(3)}_{\text{QCD}} \otimes \underbrace{\text{SU}_L(2) \otimes \text{U}_Y(1)}_{\text{QED}} \quad (2.3)$$

where the footnote $\text{SU}_C(3)$ is to avoid the confusion from the non-gauge flavor theory SU(3), $\text{SU}_L(2)$ represents the gauge group of vector bosons mediating the weak interactions, and $\text{U}_Y(1)$ represents the gauge group of weak hypercharge [24, 25].

2.1.2 Spontaneous Symmetry Breaking

[Spontaneous Symmetry Breaking](#) (SSB) is a situation that the invariance property of the symmetry is no longer preserved in the lowest energy configuration, usually referring to the [vacuum expectation value](#) (v.e.v). This is not a complicated concept but also commonly seen in daily classic systems. A famous example is to consider dropping a pencil perpendicularly to a flat ground surface. This is a system with perfect rotational symmetry around the z axis. But there are infinite number of ground state solution that the pencil lying horizontally. None these ground state have the rotational symmetry any more, and this is so called the symmetry of the physics system is “spontaneously broken”.

In the language of field theory, we can build a constant massive complex scalar field $\Phi = \frac{1}{\sqrt{2}}(\phi_1 + i\phi_2)$

with mass m . The Lagrangian density of this field could be easily written as¹:

$$\mathcal{L} = \partial_\mu \Phi^\dagger \partial^\mu \Phi - m^2 \Phi^\dagger \Phi \quad (2.4)$$

Since the mass m is real, $m^2 > 0$ and the **v.e.v** is a trivial solution with both real and complex part of Φ $\phi_1 = \phi_2 = 0$. Take Equation 2.4 as an analog but flip the sign in front of m^2 . Now the Lagrangian density becomes unstable². To get the stable form, one can rewrite Equation 2.4 by introducing a potential³

$$V(\Phi^\dagger \Phi) = \frac{1}{2\phi_0^2} m^2 (\Phi^\dagger \Phi - \phi_0)^2 + \text{const.} \quad (2.5)$$

and the Lagrangian density now is:

$$\mathcal{L} = \partial_\mu \Phi^\dagger \partial^\mu \Phi - V(\Phi^\dagger \Phi) \quad (2.6)$$

with minimum $\Phi^\dagger \Phi = \phi_0^2$, which means the system has infinite number of **v.e.v** as long as $\|\Phi\|^2 = \phi_0^2$. Notice in Equation 2.6 the U(1) global symmetry is preserved in the Lagrangian density, under the transformation of $\Phi \rightarrow \Phi' = e^{-i\theta} \Phi$. But when one picks any choice for the **v.e.v** $(\phi_0, 0)$, the global U(1) symmetry is lost. This is when **SSB** occurs in this example.

A nice interpretation is to rewrite Φ in terms of ϕ_0 and introduce two new real scalars field χ and ψ , such that

$$\Phi = \phi_0 + \frac{1}{\sqrt{2}}(\chi + i\psi) \quad (2.7)$$

Now the Lagrangian density could be expanded into a free dynamic term and an interaction term, as $\mathcal{L} = \mathcal{L}_{\text{free}} + \mathcal{L}_{\text{int}}$, where

$$\mathcal{L}_{\text{free}} = \frac{1}{2} \partial_\mu \chi \partial^\mu \chi - m^2 \chi^2 + \frac{1}{2} \partial_\mu \psi \partial^\mu \psi \quad (2.8)$$

In this free dynamic term of Lagrangian, the χ field can be interpreted as a massive scalar field, with mass $\frac{m}{\sqrt{2}}$. ψ will be a massless scalar, which is known *Nambu-Goldstone boson* [26]. In the theory of **SSB**, any

¹Notice here, that for the scalar field we do not need to implement the covariant derivation because the U(1) gauge invariance is already preserved.

²This is just a toy model for demonstration, not a proper interpretation of imaginary masses. And here m is just a parameter, without any physical meaning of mass!

³A common shorthand is to define $\lambda = \frac{m}{\sqrt{2}\phi_0}$. To reduce the confusion, we keep use the original parameterization in this dissertation.

broken global symmetry will arise a massless spin-0 particle. For example, in condensed matter physics, the *cooper pair* is effectively a Nambu-Goldstone boson that explains the superconductivity, known as the *Bardeen-Cooper-Schrieffer theory* (BCS theory) [27]. In SM, the Nambu-Goldstone bosons give masses to other gauge bosons ($W^\pm/Z$ bosons) and can cancel each other by *fine tuning*, remaining a massive scalar field which is the Higgs field [6, 28]. More details can be found in Section 2.4.

2.2 Quantum electrodynamics and electroweak interactions

2.2.1 Development of quantum electrodynamics

Maybe due to the fact that electromagnetic interaction is one of the oldest discovered fundamental interactions, *quantum electrodynamics* is the first successful formulation in the framework of QFT. QED is a U(1) gauge theory, and later inspired the development of other gauge theories such as QCD. The most intuitive way to build QED is to rewrite the Dirac's equation into a covariant form. Following the discussions in Section 2.1.1, we implement the covariant derivative and choose the electronic charge $e = g$ in Equation 2.2, Equation 2.1 now takes the form as:

$$\begin{aligned}\mathcal{L} &= \bar{\psi}(i\not{D} - m)\psi \\ &= \bar{\psi}(i\not{\partial} - m)\psi - e\bar{\psi}\not{A}\psi\end{aligned}\tag{2.9}$$

where the first term $\mathcal{L}_{\text{Dirac}} = \bar{\psi}(i\not{\partial} - m)\psi$ keeps the same as the free kinematic term for a massive fermion field introduced by Dirac's equation, and the second term shows up as the interaction term $\mathcal{L}_{\text{int}} = -e\bar{\psi}\not{A}\psi$ between the spinor field ψ and the electromagnetic field A_μ . But here A_μ field is not yet a dynamic field since a kinematic term is missing. This can be easily fixed by invoking the expression of the classical electrodynamics [29]:

$$\mathcal{L}_{\text{EM}} = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu}\tag{2.10}$$

where $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ is the electromagnetic field tensor. It is also noticed that the U(1) gauge symmetry is naturally maintained in $\mathcal{L}_{\text{EM}}$. Now we obtain the QED Lagrangian:

$$\begin{aligned}\mathcal{L}_{\text{QED}} &= \bar{\psi}(i\not{D} - m)\psi - \frac{1}{4}F^{\mu\nu}F_{\mu\nu} \\ &= \mathcal{L}_{\text{Dirac}} + \mathcal{L}_{\text{EM}} + \mathcal{L}_{\text{int}}\end{aligned}\tag{2.11}$$

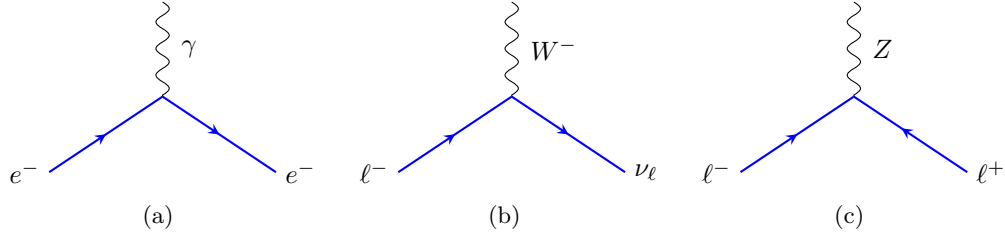


Figure 2.3: Fundamental vertices of **electroweak** interactions. As a convention, the horizontal axis is corresponding to time. (a) shows the interactions between charged fermions and photon, corresponding to a process of an electron emitting or absorbing a photon. (b) and (c) are vertices for weak interactions mediated by a W^- and Z boson respectively.

R. Feynman invented a fancy way to visualize Equation 2.11, which is later called *Feynman diagrams* and widely used in particle physics [30]. A Feynman diagram has lines with different shapes representing different types of particles. Usually the straight line represents a spinor fermion, and a squiggly line represents a vector boson. The vertices represent the interaction between these particles. Thanks to Feynman diagrams, all the **QED** phenomena can be reduced to a vertex structure shown in Figure 2.3a, every this vertex will introduce a coupling constant $ie\gamma^\mu$ from $\mathcal{L}_{\text{int}}$. But one should always keep in mind that Feynman diagrams are not only an illustration of a physics process, but also mathematically equal to the *matrix element* $\mathcal{M}$ by implementing the *Feynman rules*. The matrix element is used in the calculation of cross section of physics process, which is an observable that can be directly be measured in experiments. For example, in a typical $2 \rightarrow 2$ scattering process, the differential cross section in the **center of mass (CM)** frame is expressed as following: [31]

$$\left(\frac{d\sigma}{d\Omega}\right)_{\text{CM}} = \frac{1}{64\pi^2 E_{\text{CM}}^2} |\mathcal{M}|^2 \quad (2.12)$$

2.2.2 Discovery of weak interaction and parity violation

The weak interaction had not been realized as a fundamental interaction until a deeper look into the β decay, which is a contradict to energy conservation. In order to solve this puzzle, W. Pauli proposed a new neutral particle with half spins [32], which is named as *neutrino* (in Italian means “little neutral one”) by E. Fermi. Fermi also develop a theory that explains the beta decay as a 4-Fermi interaction model with Lagrangian density $\mathcal{L}_{\text{Fermi}} = G_F \bar{\psi}_p \psi_n \bar{\psi}_e \psi_\nu$, where G_F is the coupling constant $G_F = 1.166 \times 10^{-5} \text{ GeV}^{-2}$ [33]. This theory also successfully describe the process of a muon decay $\mu^- \rightarrow e^- + \bar{\nu}_e + \nu_\mu$. These interactions have unique properties, such as flavor changing, and very small coupling strength in large distance, and eventually called as weak interaction because it is “weaker” than electromagnetic interaction (but actually the coupling strength of weak and electromagnetic interactions are similar in small scale, see in Table 2.1).

However, more experiments in 1950s revealed problems that cannot be explained by Fermi's theory. One of the most significant problem is the violation of parity conservation in weak interactions, questioned by T. D. Lee and C. N. Yang in 1956 [34] and tested by C. S Wu in 1957 [35]. Another big constrain is that Fermi's weak theory is non-renormalizable, since G_F has a dimension of inverse power of mass [GeV⁻²]. First successful attempt to fix these problems is to distinguish the helicity of particles, that the weak interaction only act on left-handed particles and right-handed anti-particles [36]. But this progress fails again to explain the [combination of charge conjugation symmetry and parity symmetry \(CP\)](#) violation in K meson decays [37].

A successful description of weak interaction is finally achieved in 1964 by S. Glashow, S. Weinberg and A. Salam, described in the language of $SU_L(2) \otimes U_Y(1)$ gauge theory [24, 25, 38]. Besides, this theory also harmonizes [QED](#) in the same framework, known as the [electroweak \(EWK\)](#) unification. More details can be seen in Section 2.2.3. As a gauge field (photon) is predicted in the $U(1)$ gauge theory, three gauge fields are introduced in $SU_L(2) \otimes U_Y(1)$, predicting the existence of $W^\pm$ and Z bosons. This theory also posts that the left- and right-handed components of the fermion fields transform differently. More explicitly, the left handed components will be arranged in a *doublet*, such as $\begin{pmatrix} \nu_e \\ e \end{pmatrix}_L$ or $\begin{pmatrix} u \\ d \end{pmatrix}_L$, and transform under $SU_L(2)$ gauge symmetry. The right-handed components such as e_R, u_R are written as singlet. An intuitive illustration is to consider the following Lagrangian density(to be simple, only consisting electrons and electron neutrinos):

$$\mathcal{L}_{\text{Weak}} = \bar{\psi}_L(i\cancel{\partial} - \mathbf{M})\psi_L + \bar{e}_R(i\cancel{\partial} - m_e)e_R \quad (2.13)$$

where $\psi_L = \begin{pmatrix} \nu_e \\ e \end{pmatrix}_L$ and $\mathbf{M} = \begin{pmatrix} m_\nu & 0 \\ 0 & m_e \end{pmatrix}$ is the mass matrix. Similarly to [QED](#), the Lagrangian should be modified to preserve the invariance under the $SU_L(2)$ gauge symmetry transformation

$$\psi_L \rightarrow \psi'_L = e^{i\frac{g}{2}\boldsymbol{\alpha}(x^\mu) \cdot \boldsymbol{\sigma}} \psi_L \quad (2.14)$$

where g is *weak charge*, as an analogous to electric charge. $\boldsymbol{\alpha}$ is the vector phase angle in $SU(2)$ space. And $\boldsymbol{\sigma}$ is the *Pauli matrices*, also the generators of $SU(2)$ group. Rewrite the covariant derivative as:

$$\partial_\mu \rightarrow D_\mu = \partial_\mu - i\frac{g}{2}\mathbf{W}_\mu \cdot \boldsymbol{\sigma} \quad (2.15)$$

where $\mathbf{W}_\mu = \begin{pmatrix} W_\mu^1 & W_\mu^2 & W_\mu^3 \end{pmatrix}$. Similarly to A_μ , $\mathbf{W}_\mu$ follows the gauge transformation as

$$\mathbf{W}_\mu \rightarrow \mathbf{W}'_\mu = \mathbf{W}_\mu + \partial_\mu \boldsymbol{\alpha}(x^\mu) + g \mathbf{W}_\mu \times \boldsymbol{\alpha}(x^\mu) \quad (2.16)$$

Now the Lagrangian density could be rewritten as:

$$\begin{aligned} \mathcal{L} &= \bar{\psi}_L(i\cancel{\partial} - \mathbf{M})\psi_L + \bar{e}_R(i\cancel{\partial} - m_e)e_R + \frac{g}{2}\bar{\psi}_L\gamma^\mu\mathbf{W}_\mu \cdot \boldsymbol{\sigma}\psi_L - \frac{1}{8}\text{Tr } \mathbf{W}_{\mu\nu}\mathbf{W}^{\mu\nu} \\ &= \mathcal{L}_{\text{Weak}} + \mathcal{L}_{\text{int}} + \mathcal{L}_{\text{dyn}} \end{aligned} \quad (2.17)$$

where the tensor $\mathbf{W}_{\mu\nu}$ is an analog of electromagnetic tensor $F_{\mu\nu}$ defined as following:

$$\mathbf{W}_{\mu\nu} = \left(\partial_\mu + \frac{ig}{2}\mathbf{W}_\mu\right)\mathbf{W}_\nu - \left(\partial_\nu + \frac{ig}{2}\mathbf{W}_\nu\right)\mathbf{W}_\mu \quad (2.18)$$

The interaction term could be reorganized by defining the upper and lower operators $W_\mu^\pm = \frac{1}{\sqrt{2}}(W_\mu^1 \mp W_\mu^2)$:

$$\mathcal{L}_{\text{int}} = \frac{g}{2}(\bar{\nu}_e\gamma^\mu W_\mu^- e_L + \bar{e}_L\gamma^\mu W_\mu^+ \nu_e) + \frac{g}{2}(\bar{\nu}_e\gamma^\mu W_\mu^3 \nu_e - \bar{e}_L\gamma^\mu W_\mu^3 e_L) \quad (2.19)$$

Because of the nice features of Pauli matrices that $\text{Tr}(\sigma^i\sigma^j) = 2\delta^{ij}$, the dynamic Lagrangian can be written in terms of upper and lower operator as:

$$\mathcal{L}_{\text{dyn}} = -\frac{1}{4}W_{\mu\nu}^3 W^{3\mu\nu} - \frac{1}{2}W_{\mu\nu}^- W^{+\mu\nu} \quad (2.20)$$

The interpretation of Equation 2.19 is as following: the first term is corresponding to the *charged current* interactions converting electrons to neutrinos and vice versa, which is regarded as the charged spin-1 vector boson field $W^\pm$. The second term predicts a *neutral current* interaction through a neutral vector field, whereby an electron will remain an electron, and a neutrino will remain a neutrino. It is also noticeable that the Lagrangian in Equation 2.19 is invariant under the transformation of

$$f \rightarrow f' = e^{i\frac{g'}{2}Y\chi(x^\mu)}f \quad (2.21)$$

where f refers to either right- and left- handed electron or neutrino fields, $\chi(x^\mu)$ is a local phase rotation. Y is the *hypercharge*, defined by the *Gall-Mann-Nishijima formula* $Y = 2(Q - I^3)$ [39, 40, 41], where Q is the electric charge and I^3 is the *weak isospin quantum number*. More explicitly:

- For right-handed particles, $I^3 = 0$.
- For left-handed neutrinos and up-type quarks $I^3 = +\frac{1}{2}$.
- For left-handed charged leptons and down-type quarks $I^3 = -\frac{1}{2}$.

This symmetry is called $U_Y(1)$ symmetry.

2.2.3 Electroweak unification

In order to unify electromagnetic interaction and weak interaction, the covariant derivative in Equation 2.15 is modified to take account of $\chi(x^\mu)$ (from Equation 2.21) by introducing a spin-1 gauge field B_μ , as

$$\partial_\mu \rightarrow D_\mu = \partial_\mu - i\frac{g}{2}\mathbf{W}_\mu \cdot \boldsymbol{\sigma} - i\frac{g'}{2}YB_\mu \quad (2.22)$$

Such that the Lagrangian density in Equation 2.17 becomes:

$$\begin{aligned} \mathcal{L} &= \bar{\psi}_L(i\not{\partial} - \mathbf{M})\psi_L + \bar{e}_R(i\not{\partial} - m_e)e_R + \bar{\psi}_L(\frac{g}{2}\gamma^\mu \mathbf{W}_\mu \cdot \boldsymbol{\sigma} + \frac{g'Y}{2}\gamma^\mu B_\mu)\psi_L + \bar{e}_R\frac{g'Y}{2}\gamma^\mu B_\mu e_R \\ &\quad - \frac{1}{4}\mathbf{B}_{\mu\nu}\mathbf{B}^{\mu\nu} - \frac{1}{4}W_{\mu\nu}^3W^{3\mu\nu} - \frac{1}{2}W_{\mu\nu}^-W^{+\mu\nu} \\ &= \mathcal{L}_{\text{EWK}} + \mathcal{L}_{\text{int}} + \mathcal{L}_{\text{dyn}} \end{aligned} \quad (2.23)$$

where the tensor of $\mathbf{B}_{\mu\nu}$ is simply as:

$$\mathbf{B}_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu \quad (2.24)$$

The interaction term $\mathcal{L}_{\text{int}}$ could be expanded in terms of upper and lower operators:

$$\begin{aligned} \mathcal{L}_{\text{int}} &= \frac{1}{2}(\bar{\nu}_e\gamma^\mu(g'YB_\mu + gW_\mu^3)\nu_e + \bar{e}_L\gamma^\mu(g'YB_\mu - gW_\mu^3)e_L) \\ &\quad + \bar{e}_R\frac{g'}{2}Y\gamma^\mu B_\mu e_R \\ &\quad + \frac{g}{\sqrt{2}}(\bar{\nu}_e\gamma^\mu W_\mu^- e_L + \bar{e}_L\gamma^\mu W_\mu^+ \nu_e) \end{aligned} \quad (2.25)$$

The first term of Equation 2.26 indicates that the neutral weak current is corresponding to a gauge field defined as

$$Z_\mu = \frac{g'YB_\mu + gW_\mu^3}{\sqrt{g^2 + g'^2Y^2}} \quad (2.26)$$

The electromagnetic field is orthogonal to the neutral electroweak field and redefined as

$$A_\mu = \frac{gB_\mu - g'YW_\mu^3}{\sqrt{g^2 + g'^2}} \quad (2.27)$$

And since Z_μ and A_μ are from a rotation of B_μ and W_μ^3 , a *weak mixing angle* θ_W could be defined as

$$\cos \theta_W = \frac{g}{\sqrt{g^2 + g'^2}} \quad (2.28)$$

Hence the the electromagnetic field can be written in terms of W_μ^3 and B_μ as

$$A_\mu = \sin \theta_W W_\mu^3 + \cos \theta_W B_\mu \quad (2.29)$$

Now the interaction part of Lagrangian density $\mathcal{L}_{\text{int}}$ could be expanded in terms of Z_μ and A_μ :

$$\begin{aligned} \mathcal{L}_{\text{int}} = & \frac{1}{2\sqrt{g'^2 + g^2}} ((g^2 + g'^2) \bar{\nu}_e \not{Z} \nu_e - 2gg' \bar{e} \not{A} e + (g'^2 - g^2) \bar{e}_L \not{Z} e_L + 2g'^2 \bar{e}_R \not{Z} e_R) \\ & + \frac{g}{\sqrt{2}} (\bar{\nu}_R \gamma^\mu W_\mu^- e_L + \bar{e}_L \gamma^\mu W_\mu^+ \nu_e) \end{aligned} \quad (2.30)$$

Now we can reintroduce the fundamental vertices in [EWK](#) interactions as shown in [Figure 2.3](#). The electromagnetic interaction is mediated by photon γ and only acts on charged fermions, as [Figure 2.3a](#). The charged current weak interaction is mediated by a charged vector boson $W^\pm$, and only acts between the left-handed fermions and their partners, as [Figure 2.3b](#). And the neutral current weak interaction will be acting on both left- and right-handed fermions, mediated by Z boson, as shown in [Figure 2.3c](#).

However, $W^\pm$ and Z bosons have not been observed directly for a long time. This is because in the [EWK](#) theory, all three vector bosons should be massless, such like photon. Until 1983, the $W^\pm$ and Z bosons are observed in UA1 and UA2 experiments [[42](#), [43](#)], with masses $m_W = 80.379 \pm 0.012$ GeV and $m_Z = 91.1876 \pm 0.0021$ GeV⁴. In order to explain the mass problem of $W^\pm$ /Z bosons, the [SSB](#) mechanism is introduced ([Section 2.1.2](#)) and will be further described in [Section 2.4](#).

⁴These values are coming for the later precision measurement in Tevatron experiment at Fermilab and [Large Electron-Positron Collider \(LEP\)](#) experiments at CERN.

2.3 Quantum chromodynamics and strong interaction

2.3.1 Establishment of the quark model

Since the discovery of proton and neutron with α particle scattering experiments in early 20th century, a new theoretical model is needed to describe a force that binds the protons and neutrons in atomic nucleus. The confusion is to distinguish whether this force is correlated to the radioactive decay (the weak force). In 1935, M. Yukawa was the first person to speculate the difference and establish an effective description of a strong short-range force by exchanging a massive spin-0 particle between protons and neutrons [44], where the coupling part of the Lagrangian is

$$\mathcal{L}_{\text{Yukawa}} = -\lambda\bar{\psi}\Gamma\phi\psi \quad (2.31)$$

Here ϕ is representing a new scalar field. λ is the Yukawa coupling constant. Γ is defined as $\Gamma = 1$ for scalar, $\Gamma = i\gamma^5$ for pseudoscalar. By calculating the matrix element [45], the classical Yukawa potential could be derived as $V(r) = -\frac{\lambda^2}{4\pi}e^{-mr}$, where $m \approx 100$ MeV for the scale of a nucleon. This particle is called *meson*, and eventually observed after 11 years since the theoretical prediction, which is names as pion with ± 1 or 0 electrical charge ($\pi^\pm$ and π^0).

However, the improvements of experimental equipment (mainly of the invention of bubble chambers and spark chambers) led to a bewildering variety of particles discoveries from 1950s to 1960s. Experimentalists found many new spin-0 particles, similar to pion but heavier, such as η , K^- , $K^\pm$, K^0 , $\bar{K}^0$, as well as new spin- $\frac{1}{2}$ particles heavier than protons and neutrons, categorized as *baryons*. In order to organize the “zoo” of these hundreds of new particles at hand, a classification scheme for mesons and baryons (collectively called *hadrons*) is introduced by M. Gall-man in 1964, so called the *eightfold way* [46]. In this model, mesons are baryons are not fundamental particles, instead are composed of by three types of quarks, which are fermions but with fraction electric charge. Therefore mesons could be explained as bound state of quarks doublets, like charged pion $\pi^+(u\bar{d})$, whereas the baryons are bound state of quarks triplets, like proton $p(uud)$ and neutron $n(udd)$. The success of this model is proven by the observation of a predicted new baryon Ω^- with $-\frac{3}{2}$ spin [47]. In order to satisfy the Pauli exclusion principles, Gell-man also introduced a new gauge boson, called *gluon* that binding the quarks together, carrying a *color charge* of red, blue and green, which is an analog of electric charge $\pm$ [48].

2.3.2 Parton model and Parton Distribution Functions

Based on the idea of static quark model, that hadrons have substructure, in 1969 R. Feynman established Parton Model [49] to calculate the scattering cross sections and the structure functions for the nucleons. In this model, the hadrons are considered as a generic composition of many point-like constituents called *partons*. One simple example is to think of a proton at high energy. At first, the proton is composed by the quark triplet (uud), called as *valence quarks*. Then these valence and gluons can also produce an arbitrary numbers of lower energy partons, which are gluons and virtual quark-antiquark pairs called sea quarks.

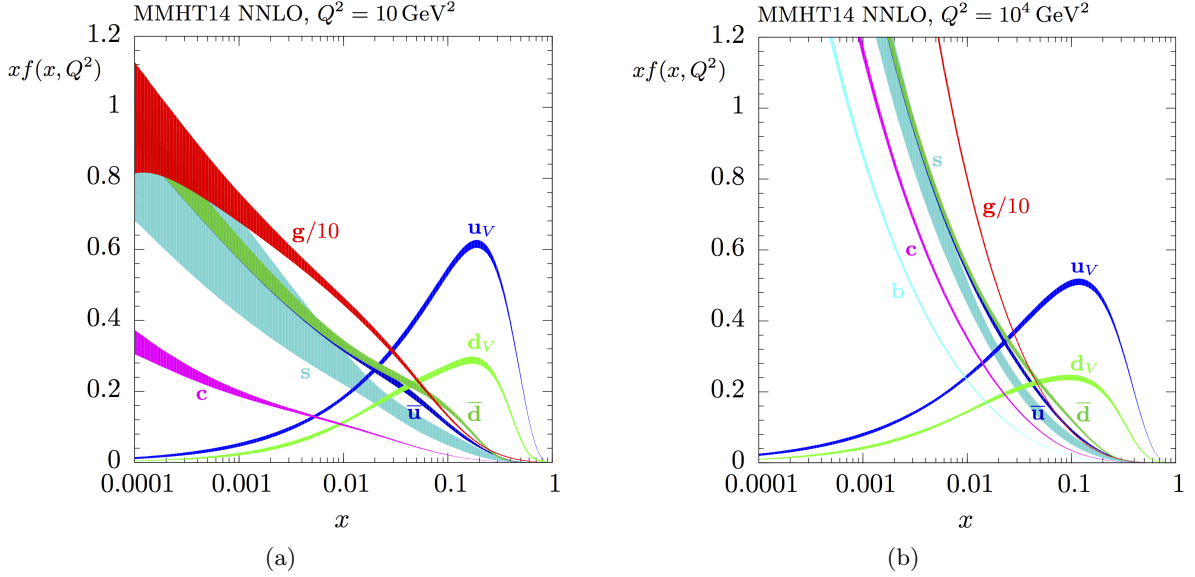


Figure 2.4: Examples of PDF at LHC with NNLO corrections, provided by MMHT2014 with NNLO corrections at (a) $Q^2=10 \text{ GeV}^2$ and (b) $Q^2=10^4 \text{ GeV}^2$ [50]. The colored bands are uncertainties with 68% confidence level. The function of $xf(x, Q^2)$ is plotted as a function of x for different partons, where the gluon parton is scaled by a factor of 10.

The probability density of finding a parton i in the given hadron with the given momentum is defined as the Parton Distribution Functions (PDF) $f_i(x, Q^2)$. Here Q^2 is the energy scale of the collision, and x is the momentum fraction of the proton that the interacting parton holds. For each hadron, the PDF should satisfy:

$$\int_0^1 x \sum_i f_i(x, Q^2) dx = 1 \quad (2.32)$$

Figure 2.4 shows an example of PDF, where the valence quarks generally dominate at low energy Q^2 , while other virtual partons are more likely to participate process with higher energy. And gluons g dominant in lower x .

In order to calculate the hadron collision process analytically, the concept of *factorization* [51] is needed, which is actually an approximation theorem by defining a cutoff scale Q_F , above which will be treated as collinear radiation, and below which can be absorbed into the PDF. Effectively, it is a separation between the perturbative phenomenon and the non-perturbative part in calculations [52]. Physically, the perturbative part is corresponding to a hard-scattering process, and the non-perturbative part is corresponding to a long-distant component. For example, the total cross section for a high energy collision process $\ell + p \rightarrow \ell' + X$ (see the Feynman diagram in Figure 2.5a) can be written as following [53]:

$$\sigma(\ell + p \rightarrow \ell' + X) = \sum_{i=q,\bar{q},g} C_i(x, Q^2 | Q^2 > Q_F^2) \otimes f_i(x, Q^2 | Q^2 < Q_F^2) \quad (2.33)$$

The term of C_i is corresponding to the perturbative coefficient function as a series in terms of the coupling constant α_s for the energy scale above the cutoff scale, whereas the part of low energy below Q_F emission is included in the term of PDF $f_i(x, Q^2)$.

This theory is soon applied in the **deep inelastic scattering (DIS)** experiments by J. Bjorken and E. Paschos in the same year at SLAC [54]. In this experiments, the scaling behavior is observed, which means the properties of probed hadrons in high-energy scattering is determined of by dimensionless kinematic quantities, instead of the absolute energy of the experiments (as shown in Figure 2.5b). This is the first evidence that protons have a substructure, which meets predictions of the quark model. But the achievement of this experiments is more fruitful. The results suggest the hints of asymptotic freedom in strong interaction, which largely boosted the establishment of **QCD** theory.

It is also important to mention that the parton model is still widely used in modern hadron collision experiments like **LHC**. In **LHC** Run2, the global PDFs NNPDF3.0, MMTH14 and CT14 are the latest used for works including this dissertation, and of which the uncertainties will be one of the dominant sources for proton-proton collision cross sections.

2.3.3 Quantum chromodynamics

Even though **quantum chromodynamics (QCD)** could be understood as an analog of well-developed **QED**, the theory has not been established for a long time. This is because the suggested symmetry of **QCD** is $SU_C(3)$, which is a non-Abelian group. **QCD** is the first successful use of Yang-Mills theory. In **QCD**, there

are three colors (red, green blue) for each quark, which can be written as triplet $\mathbf{q} = \begin{pmatrix} q_r \\ q_g \\ q_b \end{pmatrix}$, where each q_c corresponding to a quark Dirac spinor. In order to maintain Lagrangian density invariant after a local

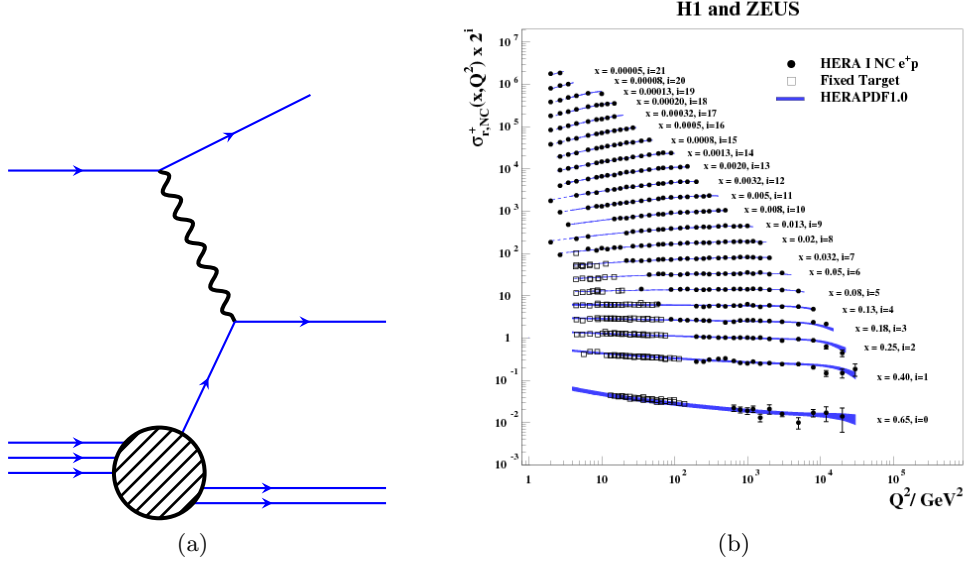


Figure 2.5: (a) Feynman diagram of $\ell + p \rightarrow \ell' + X$ deep inelastic scattering. As a convention, the horizontal axis is corresponding to time. (b) Inclusive cross section of e^+p DIS as a function of Q^2 , combining the data from HERA I NC and fixed target results [55]. It can be clearly seen that the cross section is independent of Q^2 but only correlated to the dimensionless variable x .

gauge transform of

$$q \rightarrow q'_c = e^{i \sum_{\alpha=1}^8 \chi^n(x_\mu) \lambda_\alpha} q \quad (2.34)$$

where λ_α are SU(3) generators (called *Gell-mann matrices* [56]), the covariant derivative should be written as

$$\partial_\mu \rightarrow D_\mu = \partial_\mu + i \frac{g_s}{2} G_\mu^\alpha \lambda_\alpha \quad (2.35)$$

As the $W_\mu^{1,2,3}$ is the QED gauge field, $G_\mu^{\alpha=1\dots 8}$ is the gauge field (gluon field) for QCD and the index α is traversing all eight possible combinations of the three colors, collectively written as G_μ . Now we can write down to QCD Lagrangian density as

$$\begin{aligned} \mathcal{L}_{\text{QCD}} &= \sum_f \bar{q}_f (i \not{D} - m_f) q_f - \frac{1}{4} \mathbb{G}^{\mu\nu} \mathbb{G}_{\mu\nu} \\ &= \sum_f \bar{q}_f (i \not{\partial} - m_f) q_f - \frac{g_s}{2} \sum_f \bar{q}_f \gamma^\mu \lambda G_\mu q_f - \frac{1}{4} \mathbb{G}^{\mu\nu} \mathbb{G}_{\mu\nu} \\ &= \mathcal{L}_{\text{free}} + \mathcal{L}_{\text{int}} + \mathcal{L}_{\text{dyn}} \end{aligned} \quad (2.36)$$

Here $\mathbb{G}$ is the gluon field tensor, defined as

$$\mathbb{G}_{\mu\nu}^\alpha = \partial_\mu G_\nu^\alpha - \partial_\nu G_\mu^\alpha - g_s f^{\alpha\beta\gamma} G_\mu^\beta G_\nu^\gamma \quad (2.37)$$

where $f^{\alpha\beta\gamma}$ is the structure function of SU(3). Figure 2.6 shows a set of Feynman diagrams describing some basic processes of QCD.

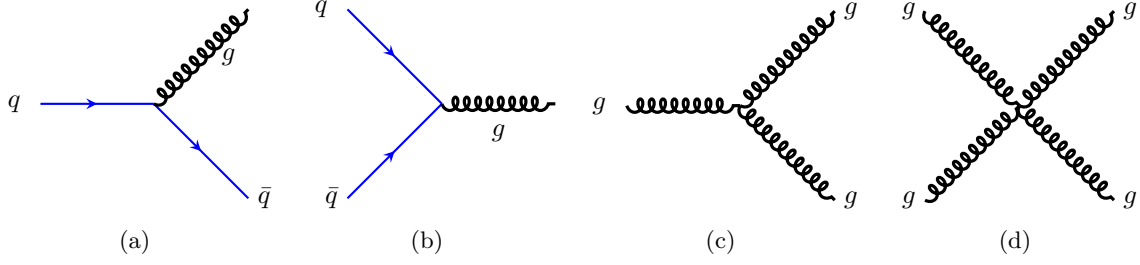


Figure 2.6: Some fundamental Feynman diagram of the lowest order QCD processes, including (a) a gluon radiation, (b) quark anti-quark annihilation, (c) gluon splitting and (d) gluon scattering. As a convention, the horizontal axis is corresponding to time.

As can be seen, QCD shares many common place with QED. But there are several critical difference make the strong interaction very different from electroweak. One of the most significant differences between QCD from QED is the coupling constant. The coupling constant of QED could be simply derived from the $\mathcal{L}_{\text{int}}$ of Equation 2.9 as $\alpha_{\text{QED}} = \frac{e^2}{4\pi} \approx \frac{1}{137}$ (in natural units), which is well known as the *fine structure constant*. But the for QCD the constant $\alpha_{\text{QCD}} > 1$, hence the calculation of next-to-leading order (NLO) terms is a headache, which makes the perturbation theory no longer work.

In QFT, because of the existence of quantum fluctuation effect, the interaction vertices will be corrected by the virtual particle-antiparticle pairs. In QED, this procedure is called *vacuum polarization*, which can be interpreted as *loops* in the context of Feynman diagrams⁵. Therefore, the measured coupling constant will be different from the bare coupling constant from Lagrangian, and dependent on the energy scale Q^2 of the measurement. This effect is called *running coupling constant*, and can be described by the β function defined as following:

$$\beta(g) = Q^2 \frac{\partial g}{\partial(Q^2)} = \frac{\partial g}{\partial \log(Q^2)} \quad (2.38)$$

For example, the β function of QED with one loop correction is $\beta(e) = \frac{e^3}{12\pi^2}$, and the coupling constant

⁵Vacuum polarization has been experimentally observed in 1997 [57]

could be corrected as [58]

$$\alpha_{\text{QED}}^{\text{1-loop}}(Q^2) = \frac{e^2}{4\pi - \frac{e^2}{3\pi} \log\left(\frac{Q^2}{4m^2}\right)} \quad (2.39)$$

The correction of α_{QCD} is derived by D. Gross, F. Wilczek and H. D. Politzer based on Yang-Mills theory as following [59, 60]:

$$\beta(g_s) = -\left(11 - \frac{n_s}{3} - \frac{2n_f}{3}\right) \frac{g_s^3}{16\pi^2} \quad (2.40)$$

$$\alpha_{\text{QCD}}(Q^2; \Lambda^2) = \frac{g_s^2}{4\pi - \frac{g_s^2}{4\pi} \frac{1}{n_c} (2n_f - 11n_c) \log\left(\frac{Q^2}{\Lambda^2}\right)} \quad (2.41)$$

where $n_f = 6$ and $n_c = 3$ for the number of flavor and number of color for [SM](#), but could be generalized to any n_c and n_f . $\Lambda^2 \approx 220 \text{ MeV}$ is the energy cutoff for [QCD](#). The β function in Equation 2.40 is smaller than zero, which dictates that the coupling strength for strong interaction decreases with increasing of the energy scale. This is known as the *asymptotic freedom* where charged particles barely interact with each other at small distance [61].

For those with large distances (on the order of femtometers) and low energy scales $Q^2 \ll \Lambda^2$, the perturbative calculation blows up, and a non-perturbative approach needed to calculate the interactions, and confirms the effectiveness of factorization approximation. Experimentally, this is known as *color confinement*, which means that at low energies and large distances, the quarks and gluons cannot be observed individually, instead they combine to form colorless hadrons. Even with more energy, it is favorable to produce a new quark-antiquark pair from vacuum, rather than putting more energy to separate the two particles.

In high energy hadron colliders such as [LHC](#), the boosted quarks or gluons that flying out of the incident hadron with large amount of energy will create colorless bound states of hadrons. This process is called *hadronization* and refer to the transition of colored partons to colorless hadrons. Furthermore, these partons can continue to radiate energetic gluons (Bremsstrahlung) and hadronized to heavy hadrons. These hadrons will further decay into collimated hadrons. This avalanche process called *showering*. A demonstration can be find in Figure 2.7. In particle detectors such as ATLAS, the ensemble of these colorless hadrons is called a *jet*.

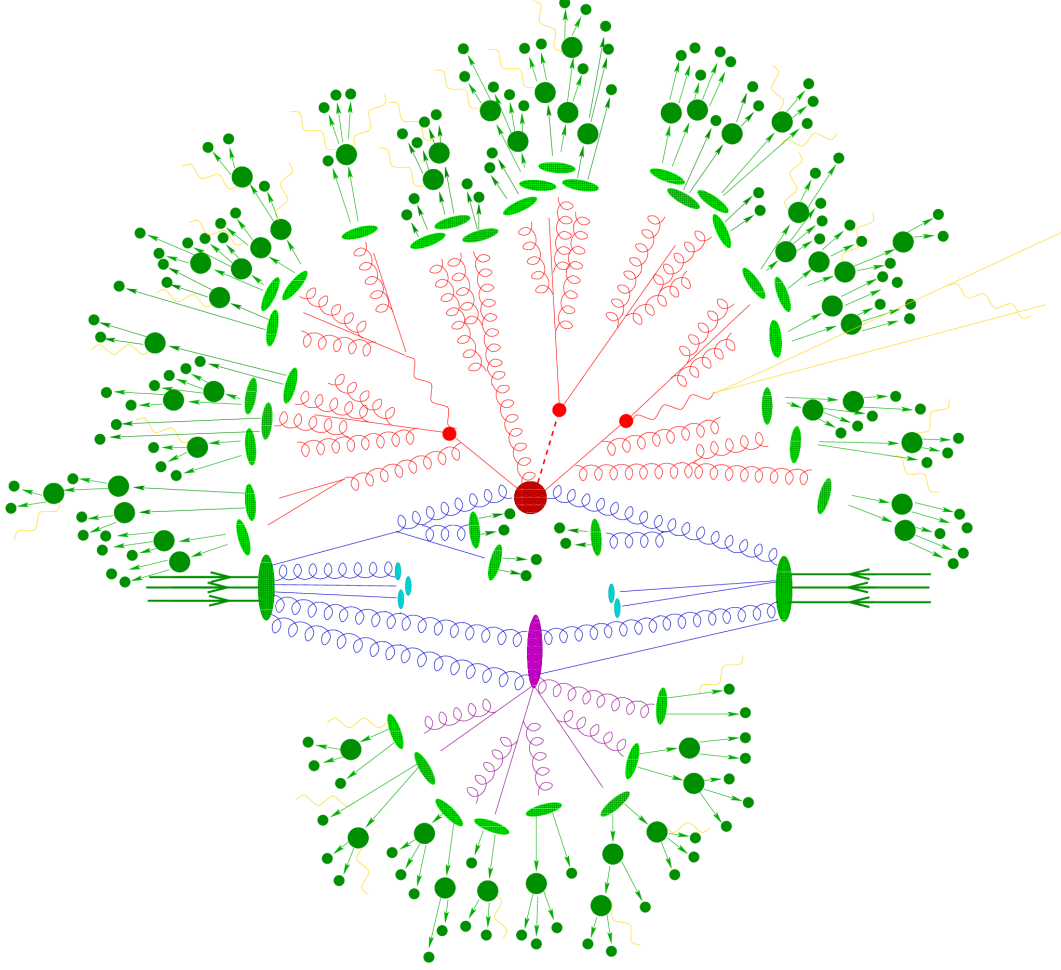


Figure 2.7: Sketch of a hadron-hadron collision as simulated by a Monte-Carlo event generator [62]. The red blob in the center represents the hard collision, surrounded by a tree-like structure representing Bremsstrahlung as simulated by parton showers. The purple blob indicates a secondary hard scattering event. Parton-to-hadron transitions are represented by light green blobs, dark green blobs indicate hadron decays, while yellow lines signal soft photon radiation.

2.4 Higgs mechanism

2.4.1 Higgs mechanism for $W^\pm/Z$ masses

As mentioned in Section 2.2.3, the masses of the **EWK** gauge bosons are massless. To explain the contradiction from the observations, a **SSB** mechanism is introduced by P. Higgs, F. Englert and R. Brout [6, 28], so called the *Higgs mechanism*, which is a broken symmetry in $SU_L(2) \otimes SU_C(3)$. To illustrate we can consider a doublet of complex scalar field, $\Phi(x^\mu) = \frac{1}{\sqrt{2}} \begin{pmatrix} \phi_1(x^\mu) + i\phi_2(x^\mu) \\ \phi_3(x^\mu) + i\phi_4(x^\mu) \end{pmatrix}$, just as we did in Section 2.1.2 but has four degrees of freedom. Now we can write down the Lagrangian density of this scalar field in a covariant

form corresponding to Equation 2.6:

$$\mathcal{L}_\Phi = (D_\mu \Phi)^\dagger D^\mu \Phi - V(\Phi^\dagger \Phi) \quad (2.42)$$

where D_μ is the [EWK](#) covariant derivation defined in Equation 2.22. $V(\Phi^\dagger \Phi)$ is the same as Equation 2.5. The ground state can be manually set as $\phi_1 = \phi_2 = \phi_4 = 0$ and $\phi_3 = \phi_0$. Approximately an excited state close to the ground state can be expressed as $\Phi(x^\mu)' = \begin{pmatrix} 0 \\ \phi_0 + \frac{1}{\sqrt{2}}(h(x^\mu) + i\chi(x^\mu)) \end{pmatrix}$ where $h(x^\mu)$ and $\chi(x^\mu)$ are real, this is known as the *unitary gauge*. Notice that $\chi(x^\mu)$ plays the role as the massless Nambu-Goldstone boson in Section 2.1.2, and we effectively ignore them since they are not physically observable. The massive part $h(x^\mu)$ is the *Higgs field*, plugging into Equation 2.42 we obtain:

$$\begin{aligned} \mathcal{L}_\Phi = & \frac{1}{2} \partial_\mu h \partial^\mu h + \frac{g^2}{2} W_\mu^- W^{+\mu} \left(\phi_0 + \frac{h}{\sqrt{2}} \right)^2 \\ & + \frac{1}{4} (g^2 + g'^2) Z_\mu Z^\mu \left(\phi_0 + \frac{h}{\sqrt{2}} \right)^2 - m^2 h^2 + \frac{m^2}{\sqrt{2}\phi_0} h^3 + \frac{m^2}{8\phi_0^2} h^4 \end{aligned} \quad (2.43)$$

The last three term is called *Higgs potential*, which takes the shape of a Mexican hat as seen in Figure 2.8:

$$V(h) = m^2 h^2 - \frac{m^2}{\sqrt{2}\phi_0} h^3 - \frac{m^2}{8\phi_0^2} h^4 \quad (2.44)$$

But more clearly, the field of $W^\pm/Z$ and h are massive, with

$$m_{W^\pm} = \frac{g}{\sqrt{2}} \phi_0 \quad (2.45)$$

$$m_Z = \sqrt{\frac{g^2 + g'^2}{2}} \phi_0 \quad (2.46)$$

$$m_H = \sqrt{2}m \quad (2.47)$$

And the electromagnetic field A_μ remains as massless. It is also easy to figure that the weak mixing angel is exactly the mass ratio between Z and $W^\pm$ bosons, as defined in Equation 2.28, also reported in [particle data group \(PDG\)](#) as well [63].

It also can be noticed that the Higgs [v.e.v](#) ϕ_0 is determined by the masses of W and Z bosons, but Higgs boson mass is a free parameter. From experimental observations, all of these masses are measured, including the mass of Higgs boson which is observed in 2012 by ATLAS and CMS collaborations [7, 8], $m_H \approx 125$ GeV.

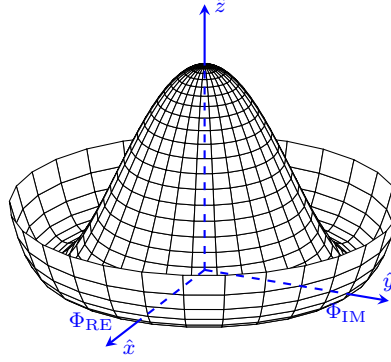


Figure 2.8: The shape of the two-dimensional Higgs potential $V(\Phi^\dagger\Phi)$. As can be seen, the local minimum has infinite numbers of states.

2.4.2 Higgs mechanism for SM fermions

To illustrate the Higgs mechanism acting on fermions, we can go back to the free Dirac Lagrangian density in Equation 2.1. As introduced in Section 2.2.2, the left- and right-handed fermions interact differently. We can use a projection operator $\frac{1 \mp \gamma^5}{2}$ to split the left- and right-handed part of the spinor field ψ as:

$$\psi_L = \frac{1 - \gamma^5}{2} \psi \quad (2.48)$$

$$\psi_R = \frac{1 + \gamma^5}{2} \psi \quad (2.49)$$

Hence the kinematic part and mass part of Equation 2.1 can be written as:

$$i\bar{\psi}\not{\partial}\psi = i\bar{\psi}_L\not{\partial}\psi_L + i\bar{\psi}_R\not{\partial}\psi_R \quad (2.50)$$

$$-m\bar{\psi}\psi = -m\bar{\psi}_L\psi_R - m\bar{\psi}_R\psi_L \quad (2.51)$$

Consider that the covariant derivative D_μ also acts differently on left- and right-handed parts, as can be seen in Equation 2.23, or more explicitly as:

$$D_\mu\psi_L = \left(\partial_\mu - \frac{ig}{2}\mathbf{W}_\mu \cdot \boldsymbol{\sigma} - \frac{ig'}{2}YB_\mu \right) \psi_L \quad (2.52)$$

$$D_\mu\psi_R = \left(\partial_\mu - \frac{ig'}{2}YB_\mu \right) \psi_R \quad (2.53)$$

Now we can rewrite Equation 2.4.2 as $\partial_\mu \rightarrow D_\mu$ as:

$$\mathcal{L}_{\text{fermion}} = i\bar{\psi}_L\not{D}\psi_L + i\bar{\psi}_R\not{D}\psi_R \quad (2.54)$$

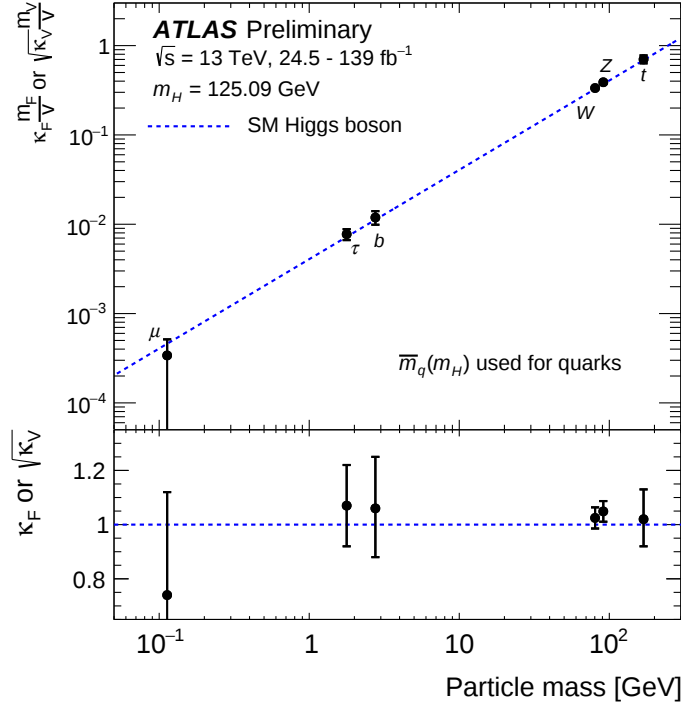


Figure 2.9: Reduced coupling strength modifiers κ_F for fermions ($F=t, b, \tau, \mu$) and $\sqrt{\kappa_V}$ for weak gauge bosons ($V=W, Z$) as a function of their masses m_F and m_V , respectively, and the v.e.v of the Higgs field $v = \frac{\phi_0}{\sqrt{2}} = 246$ GeV. The SM prediction for both cases is also shown (dotted line). The couplings modifiers κ_F and κ_V are measured assuming no BSM contributions to the Higgs boson decays, and the SM structure of loop processes such as ggF , $H \rightarrow \gamma\gamma$ and $H \rightarrow gg$. The lower inset shows the ratios of the values to their SM predictions [64, 65].

which can be seen that the Lagrangian density is gauge invariant, but the masses of the fermions are zero.

Here is the Higgs mechanism again introduced. Since Higgs field is a scalar, the coupling to fermions the same as the Yukawa coupling in Equation 2.31. We can write down the additional term of Yukawa as:

$$\mathcal{L}_{\text{Yukawa}} = -\lambda \bar{\psi} \Phi \psi = -\lambda (\bar{\psi}_L \Phi \psi_R + \bar{\psi}_R \Phi \psi_L) \quad (2.55)$$

Applying the unitary gauge of the Higgs field as Equation 2.47, $\mathcal{L}_{\text{Yukawa}}$ can be expressed in terms of each fermion f_i , as:

$$\mathcal{L}_{\text{Yukawa}} = \sum_i \frac{\lambda_i \phi_0}{\sqrt{2}} (f_{iL}^\dagger f_{iR} + f_{iR}^\dagger f_{iL}) \quad (2.56)$$

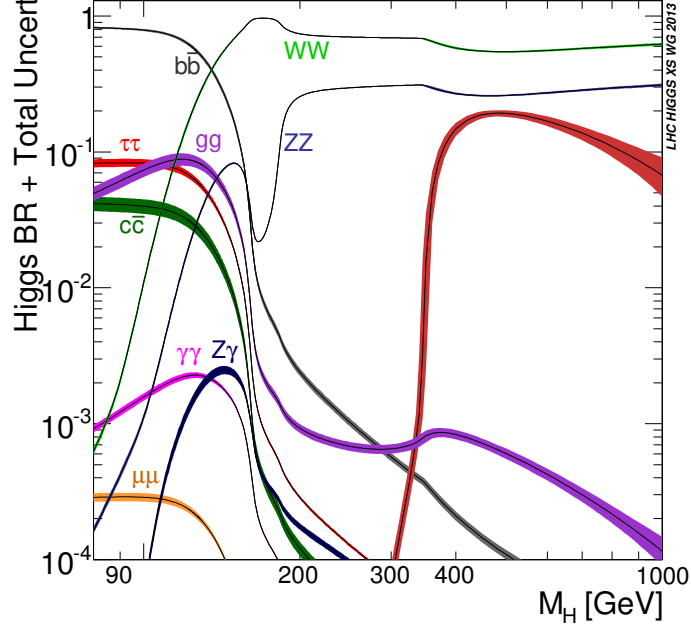


Figure 2.10: Higgs [branching ratio](#) and their uncertainties as a function of Higgs mass with [center of mass](#) energy $\sqrt{s} = 13$ TeV [66], which corresponds the scenario of [LHC Run 2](#).

It is easy to see that $\mathcal{L}_{\text{Yukawa}}$ takes the same form as Equation, where the mass of each fermion would be:

$$m_i = \frac{\lambda_i \phi_0}{\sqrt{2}} \quad (2.57)$$

One the other hand, the coupling strength of Higgs field to fermions is proportional to their masses. According to this mechanism, the production and decay modes of Higgs boson in high energy colliders is understood. Figure 2.9 shows the most recent results Higgs coupling measurement in ATLAS, where the Higgs coupling strength between the fermions κ_F and vector boson κ_V are:

$$\kappa_F = \frac{m_F}{\sqrt{2}\phi_0} = 2\lambda_F \quad (2.58)$$

$$\kappa_V = \frac{m_V^2}{2\phi_0^2} \quad (2.59)$$

which indicates the [branching ratio](#) (Br) of the [SM](#) Higgs decays. However, the [SM](#) Higgs boson can also decay to massless vector bosons such as gluons and photons through loops processes of fermions or $W^\pm/Z$ bosons [67]. The final calculated branching ratio of Higgs decays as a function of the mass of Higgs boson m_H can be find in Figure 2.10.

Chapter 3

Motivations for searches of exotic Higgs decays

In the past decades, [Standard Model](#) has been tested by many experiments, and shown to be robust. For instance, [Figure 3.1](#) shows the most recent results of cross section measurements of several processes at [LHC](#) comparing the experimented measurement to the theoretical predictions of [SM](#). Furthermore, the properties of the Higgs boson have been measured and compared to [SM](#) predictions.

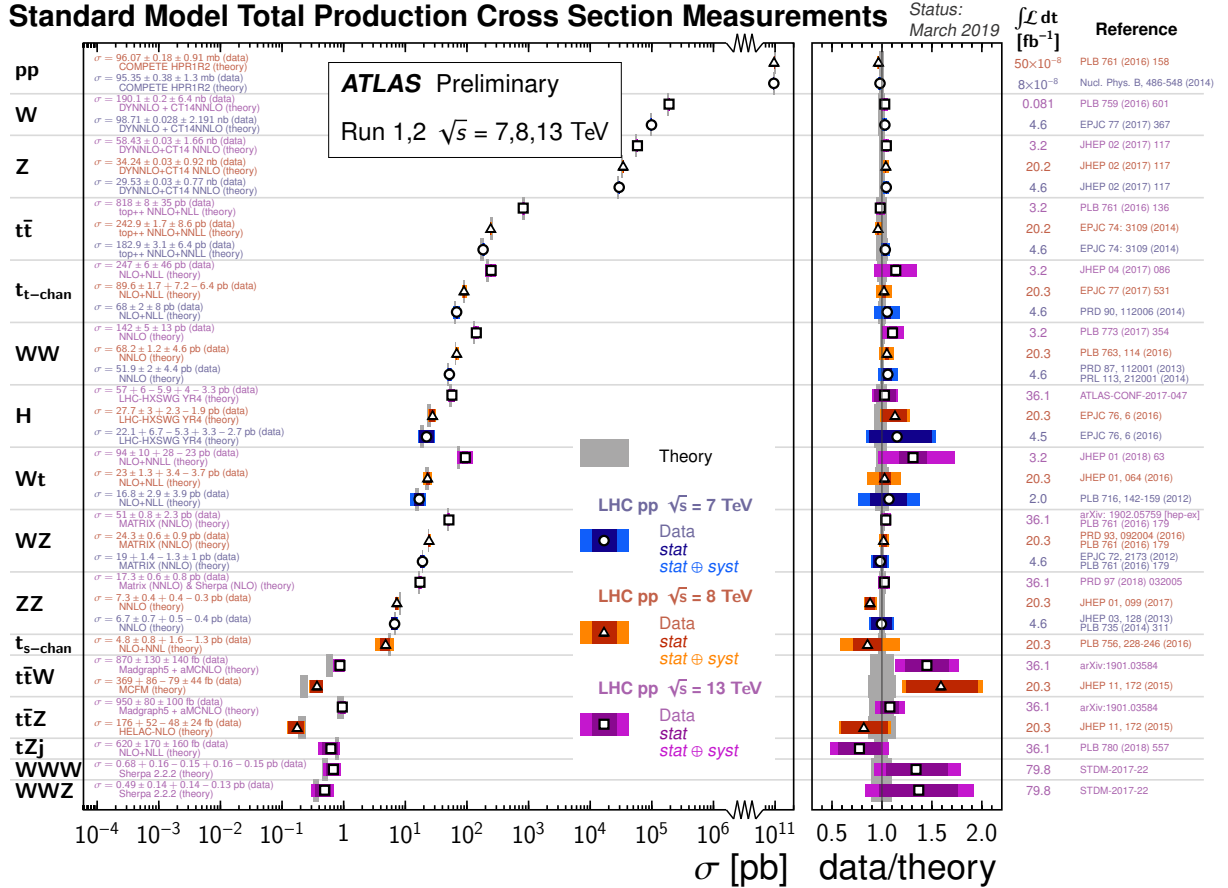


Figure 3.1: Summary of several Standard Model total production cross section measurements in ATLAS [68], corrected for leptonic branching fractions, compared to the corresponding theoretical expectations and ratio with respect to best theory. The Run1 and Run2 data are combined in this plot.

However, there are still a long list questions unanswered or contradict to [SM](#), including but not limited to:

- The question of neutrino mass. Neutrinos do not have a right-handed partner in [SM](#), which means they are massless under Higgs mechanism (see Section 2.4.2). However, many experiments show evidence of neutrino oscillation [69], indicating that the mass of neutrinos are non-zero.
- The matter-antimatter asymmetry problem. The observed universe is dominated by baryonic matter instead of anti-matter. The [SM](#) does not provide an explanation for why this should be so.
- The naturalness question (also known as *hierarchy problem*): the [electroweak](#) scale $\Lambda_{\text{EWK}} \approx 10^2$ GeV is much smaller than the Planck scale $\Lambda_{\text{Planck}} \approx 10^9$ GeV. According to [QFT](#), the mass of Higgs boson is corrected by loop level contributions from particles that couple to the Higgs field. Explicitly, the observable Higgs mass could be written as:

$$m_{\text{H}}^2 = (m_{\text{H}}^0)^2 + \Delta m_{\text{H}}^2 \approx 125 \text{ GeV} \quad (3.1)$$

where m_{H}^0 is the bare mass of the Higgs boson, and the correction term Δm_{H}^2 is:

$$\Delta m_{\text{H}}^2 = -\frac{\lambda_f^2}{8\pi^2} \Lambda_{\text{UV}}^2 + \dots \quad (3.2)$$

where $\Lambda_{\text{UV}} = \Lambda_{\text{Planck}}$ is the cutoff scale up to which the [SM](#) is assumed to be valid. Considering the mass spectrum of fermions, the largest contribution in Equation 3.2 is from the top quark since $m_t \approx 175$ GeV (as shown in Figure 3.2). The term Δm_{H}^2 takes the magnitude $\Lambda_{\text{Planck}}^2 \approx 10^{19}$, which would need to be corrected that the bare mass parameter of the same magnitude to keep the corrected m_{H} at the [EWK](#) scale of 10^2 GeV. Unless there is some new physics in this scale, the process of “fine tuning” process is unnatural.

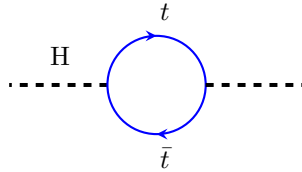


Figure 3.2: Feynman diagram of a one-loop correction to the Higgs mass m_{H} .

- Explanation for [dark matter](#) (DM). Many astrophysical observations [70, 71] provide compelling evidence for the presence of [dark matter](#). The estimated amount of the dark matter is five and a half

times larger than the ordinary matter, yet it is not included in the SM.

- As one of the fundamental forces, the gravitational force is not included in the SM.

The motivation of many BSM theories is to solve these limitations, such as Grand Unification Theory (GUT) [72] and Supersymmetry (SUSY) [73]. There is a large effort underway to search for experimental evidence for new physics as proposed by these theories, especially at the energy frontier such as the LHC physics. So far there is no clear evidence of new physics found. Figure 3.3 and Figure 3.4 shows a summary of the mass limits of most recent searches for new physics signatures in ATLAS.

However, there are still many open topics that have a large potential to discover new physics. One possibility that has not been explored on much is the searches for BSM Higgs decays, also known as exotic Higgs decays, which will be explained in Section 3.1. Exotic Higgs decays are the main motivation for the work performed in this dissertation. Section 3.2 will introduce the two-Higgs-doublet plus singlet model (2HDM+S) model, which is a simple extension of 2HDM. The 2HDM+S model is used as a benchmark for the phenomenology in exotic Higgs decays. Finally, other motivations will be covered in Section 3.3.

ATLAS SUSY Searches* - 95% CL Lower Limits

October 2019

ATLAS Preliminary

$\sqrt{s} = 13 \text{ TeV}$

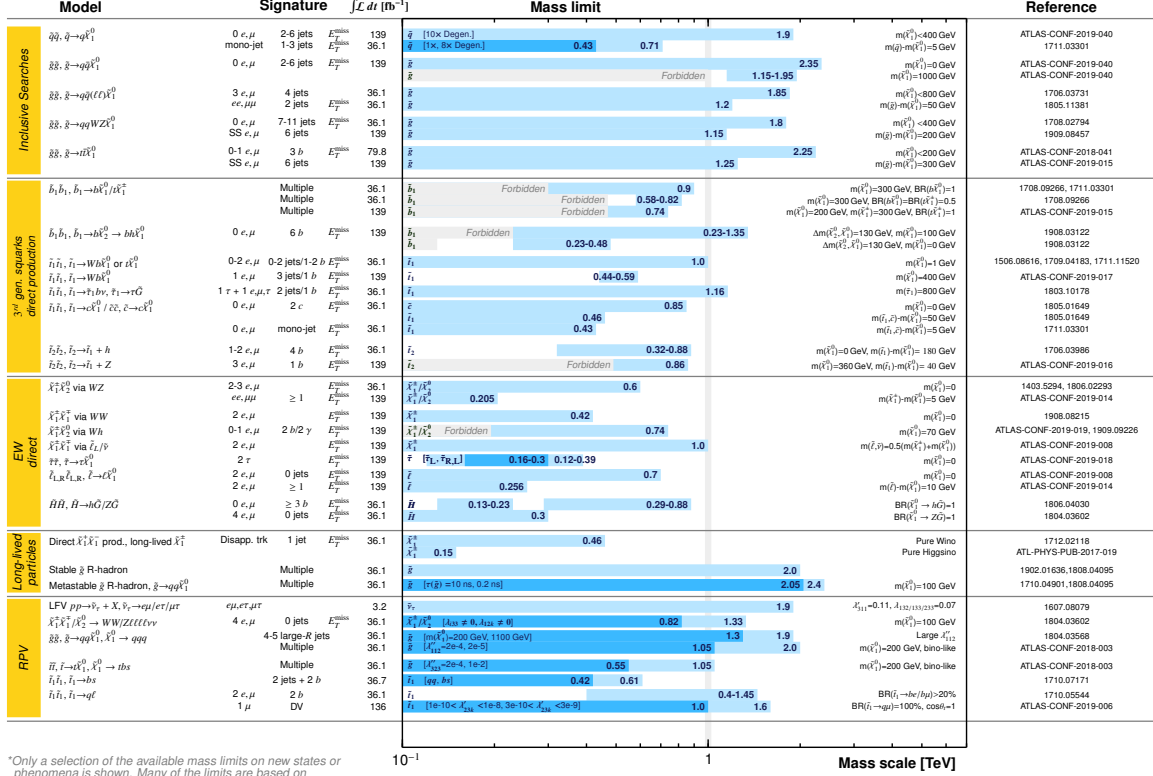


Figure 3.3: Mass reach of the ATLAS searches for **Supersymmetry**. A representative selection of the available search results is shown [74]. Results are quoted for the nominal cross section in both a region of near-maximal mass reach and a demonstrative alternative scenario, in order to display the range in model space of search sensitivity. Some limits depend on additional assumptions on the mass of the intermediate states, as described in the references provided in the plot. In some cases these additional dependencies are indicated by darker bands showing different model parameters.

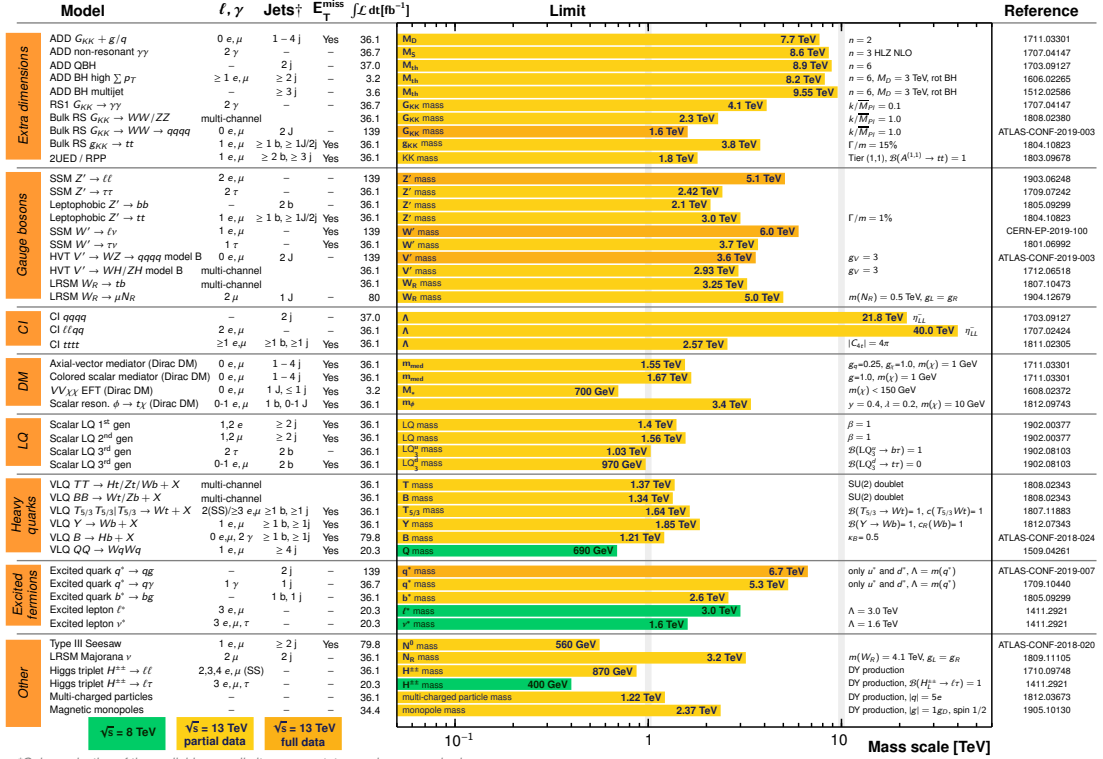
ATLAS Exotics Searches* - 95% CL Upper Exclusion Limits

Status: May 2019

ATLAS Preliminary

$$\int \mathcal{L} dt = (3.2 - 139) \text{ fb}^{-1}$$

$$\sqrt{s} = 8, 13 \text{ TeV}$$



*Only a selection of the available mass limits on new states or phenomena is shown.

[†]Small-radius (large-radius) jets are denoted by the letter j (J).

Figure 3.4: Reach of ATLAS searches for new phenomena other than SUSY [75]. Only a representative selection of the available results is shown. Green bands indicate 8 TeV data results; yellow (orange) bands indicate 13 TeV data results with partial (full) data set.

3.1 Indirect evidence of exotic Higgs decays

The Higgs boson may play an essential role connecting the SM and BSM physics. Searches for *exotic Higgs decays* are a particularly rich and fruitful way to seek evidence of new physics, especially due to the following reasons:

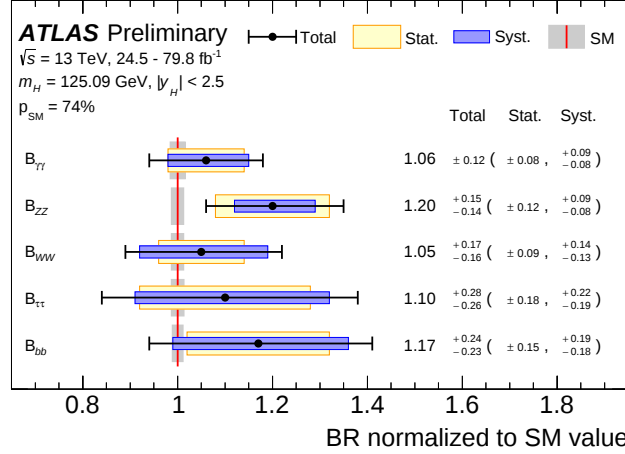


Figure 3.5: Branching ratios for $H \rightarrow \gamma\gamma$, $H \rightarrow ZZ^*$, $H \rightarrow WW^*$, $H \rightarrow \tau\bar{\tau}$ and $H \rightarrow b\bar{b}$ normalized to their SM predictions, measured under SM assumptions for the Higgs boson production processes [3]. The black error bars, blue boxes and yellow boxes show the total, systematic, and statistical uncertainties in the measurements, respectively. The blue bands indicate the theory uncertainties on the predictions.

- The width of the SM Higgs boson is extremely narrow, $\Gamma_H \approx 4.07$ MeV. The reason is that the $\gamma\gamma$, gg channels are suppressed by loop factors, and the decays to WW and ZZ are suppressed by multi-body phase space. For the fermion final states, the dominant decay mode of $H \rightarrow b\bar{b}$ is controlled by the Yukawa coupling to b quarks $\lambda_b = \frac{\sqrt{2}m_b}{\phi_0} \approx 0.018$ (derivation can be seen in Equation 2.59). This indicates that even a small coupling to another light state can easily open a sizable decay mode¹ [76, 77].
- Combining the most recent measurements of the 125 GeV Higgs boson in multiple SM channels, it is possible to constrain $\text{Br}_{\text{BSM}} \lesssim 21\%$ at 95% confidence level [3]. The branching ratio of Higgs decays to available final states compared to the SM prediction is shown in Figure 3.5. Branching ratios of $\mathcal{O}(10\%)$ into exotic decay modes are still allowed, and remain reasonable target for direct searches using the data set already collected and from the future LHC. Table 3.1 summarizes the expected number of exotic Higgs decay events in different LHC run scenarios [78].

- The Higgs can provide one of a few “portals” that allow SM matter to interact with hidden-sector matter that is not charged under SM forces, such as dark matter. It is possible to construct a singlet

¹The branching ratio of exotic Higgs decays can simply be calculated as $\text{Br}(H \rightarrow XX) = \frac{\Gamma_X}{\Gamma_H}$.

Higgs production mode	$\sqrt{s} = 7 \text{ TeV}, 5 \text{ fb}^{-1}$		$\sqrt{s} = 8 \text{ TeV}, 20 \text{ fb}^{-1}$		$\sqrt{s} = 14 \text{ TeV}, 130 \text{ fb}^{-1}$	
	$\sigma(\text{pb})$	$N_{\text{event}}^{\text{Br(BSM)}=10\%}$	$\sigma(\text{pb})$	$N_{\text{event}}^{\text{Br(BSM)}=10\%}$	$\sigma(\text{pb})$	$N_{\text{event}}^{\text{Br(BSM)}=10\%}$
ggF	15.13	7600	19.27	38500	49.85	1.5×10^6
VBF	1.22	610	1.58	3200	4.18	125000
$HW^\pm$	0.58	290	0.70	1400	1.5	45000
$HW^\pm(\ell^\pm\nu)$	0.58×0.21	62	0.70×0.21	300	1.5×0.21	9600
HZ	0.34	170	0.42	830	0.88	26500
$HZ(\ell^+\ell^-)$	0.34×0.067	11	0.42×0.067	56	0.88×0.067	1800
$t\bar{t}H$	0.086	43	0.13	260	0.61	18300

Table 3.1: The number of exotic Higgs decays in existing [LHC](#) data different run scenarios, assuming the [SM](#) production cross section of a 125 GeV Higgs boson and a [branching ratio](#) of $\text{Br(BSM)} = 10\%$ for various Higgs production modes [\[78\]](#).

scalar field s coupling to the [SM](#) Higgs field through the *Higgs portal* with an additional term in the Lagrangian:

$$\Delta\mathcal{L} = \zeta s^2 \left(\psi_0 + \frac{h}{\sqrt{2}} \right)^2 \quad (3.3)$$

of which the details are introduced in [Section 3.2](#). [Figure 3.7](#) shows the sensitivities of exotic Higgs decay for different coupling constants ζ . Comparing to [Table 3.1](#), it can be seen that most cases are reachable in the [LHC](#) energy and luminosity scales.

3.2 Theoretical frameworks

3.2.1 Two-Higgs-doublet model

As introduced in [Section 2.4](#), the [SM](#) Higgs boson is a complex scalar doublet. However, there are no constraints for the numbers of Higgs doublets, namely it is very easy to extend the number of Higgs doublets. The simplest model is the [two-Higgs-doublet model](#) ([2HDM](#)) proposed in 1973 by T.D. Lee [\[79\]](#). This model assumes two $SU(2)$ doublet spin-0 fields:

$$\Phi_1 = \begin{pmatrix} \phi_1^+ \\ \phi_1^0 \end{pmatrix} \quad \Phi_2 = \begin{pmatrix} \phi_2^+ \\ \phi_2^0 \end{pmatrix} \quad (3.4)$$

where $\phi_1^+, \phi_1^0, \phi_2^+, \phi_2^0$ are complex fields.

The scalar potential can be written as

$$\begin{aligned}
V_{\text{2HDM}} = & m_{11}^2 \Phi_1^\dagger \Phi_1 + m_{22}^2 \Phi_2^\dagger \Phi_2 - [m_{12}^2 \Phi_1^\dagger \Phi_2 + \text{h.c.}] + \frac{1}{2} \lambda_1 (\Phi_1^\dagger \Phi_2)^2 + \frac{1}{2} (\Phi_2^\dagger \Phi_2)^2 \\
& + \lambda_3 (\Phi_1^\dagger \Phi_1) (\Phi_2^\dagger \Phi_2) + \lambda_4 (\Phi_1^\dagger \Phi_2) (\Phi_2^\dagger \Phi_1) + [\frac{1}{2} (\Phi_1^\dagger \Phi_2)^2 + \text{h.c.}]
\end{aligned} \tag{3.5}$$

where all masses and coupling parameters can be chosen to be real. The minimum of the scalar potential should preserve the $U(1)$ gauge symmetry of the SM, such that the scalar fields develop the following [vacuum expectation value](#):

$$\langle \Phi_1 \rangle = \begin{pmatrix} 0 \\ v_1 \end{pmatrix} \qquad \langle \Phi_2 \rangle = \begin{pmatrix} 0 \\ v_2 \end{pmatrix} \tag{3.6}$$

Using the minimum conditions, the two mass parameters m_{11} and m_{22} can be expressed in terms of v_1 and v_2 . Therefore, with $v = \sqrt{v_1^2 + v_2^2} \approx 246$ GeV and $\tan \beta = v_1/v_2$, the [two-Higgs-doublet model \(2HDM\)](#) can be expanded around the potential minimum in terms of this component fields as

$$\Phi_1 = \begin{pmatrix} \phi_1^+ \\ \frac{1}{2}(v \cos \beta + \phi_1^0) \end{pmatrix} \qquad \Phi_2 = \begin{pmatrix} \phi_2^+ \\ \frac{1}{2}(v \sin \beta + \phi_2^0) \end{pmatrix} \tag{3.7}$$

From the original 8 scalar degrees of freedom, 3 Goldstone bosons ($G^\pm$ and G) are absorbed by the $W^\pm$ and Z bosons. The remaining $8 - 3 = 5$ degrees of freedom from the physical Higgs states of the model includes: two CP-even scalars (h and H with masses m_h and m_H respectively and $m_H \geq m_h$), one CP-odd pseudoscalar (A), and a pair of charged Higgs bosons ($H^\pm$) [80].

It is convenient to express the scalar doublet fields in the Higgs basis [81] as

$$H_1 = \begin{pmatrix} H_1^+ \\ s H_1^0 \end{pmatrix} = \Phi_1 \cos \beta + \Phi_2 \sin \beta \qquad H_2 = \begin{pmatrix} H_2^+ \\ H_2^0 \end{pmatrix} = -\Phi_1 \sin \beta + \Phi_2 \cos \beta \tag{3.8}$$

such that [vacuum expectation value](#) of these fields is $\langle H_1^0 \rangle = \frac{v}{\sqrt{2}}$ and $\langle H_2^0 \rangle = 0$.

The Higgs boson coupling to fermions are described again by Yukawa interactions as:

$$\mathcal{L}_{\text{Yukawa}} = - \sum_{i=1}^2 [\bar{f}_L \tilde{\Phi}_i y_i^u u_R + \bar{f}_L \Phi_i y_i^d d_R + \bar{f}_L \tilde{\Phi}_i y_i^e e_R + \text{h.c.}] \quad (3.9)$$

where $\tilde{\Phi} = i\sigma_2 \Phi$ and $y_i^{u,d,e}$ are two 3×3 Yukawa matrices in the flavor space of each Higgs doublet ($i = 1, 2$), and fermion classes $f = u, d, e$ for up-type quarks, down-type quarks and leptons, respectively.

The absence of tree-level [flavor-changing neutral currents \(FCNC\)](#) is guaranteed by the Glashow-Weinberg condition [82] that all fermions of a given representation receive their masses by renormalizable Yukawa couplings to a single Higgs doublet, in which case the tree-level couplings of neutral Higgs bosons are diagonal in the mass eigenbasis. This condition could be satisfied under four discrete assignments, where by convention up-type quarks are always taken to couple to Φ_2 :

- **Type-I** $y_1^{u,d,e} = 0$, which means all fermions couple to one doublet. In this type of model, one Higgs doublet provides masses to all fermions, which is similar to the [SM Higgs](#).
- **Type-II** $y_1^u = y_2^d = y_2^e = 0$, which means up-type quarks couple to one doublet and the down-type quarks and leptons couple to the other. In this type of model, one Higgs doublet provides masses to up-type quarks and the other to down-type quarks and leptons, which is similar to the [minimal supersymmetric standard model \(MSSM\)](#)
- **Type-III** $y_1^u = y_1^d = y_2^e = 0$, which means quarks couple to one doublet and leptons to the other. In this type of model, one Higgs doublet provides masses to quarks and the other to leptons.
- **Type-IV** $y_1^u = y_1^e = y_2^d = 0$, which means up-type quarks and leptons couple to one doublet and down-type quarks couple to the other. In this type of model, one doublet provides masses to up-type quarks and leptons, and the other to down-type quarks.

Many analyses have been performed to search for additional Higgs bosons in ATLAS. One of these models is the [habemus minimal supersymmetric standard model \(hMSSM\)](#) [83, 84], which is a model where the lighter Higgs boson h has a mass of approximately 125 GeV. In addition, the non-observation of superparticles at [LHC](#) indicates the [SUSY](#)-breaking scale M_S is greater than 1 TeV. The [hMSSM](#) model has two free parameters $\tan \beta$ and the mass of the heavier Higgs m_A . Figure 3.6 [85] summarizes the most recent results of searches for these signatures in ATLAS, where the light shaded or hashed regions indicate the observed exclusions.

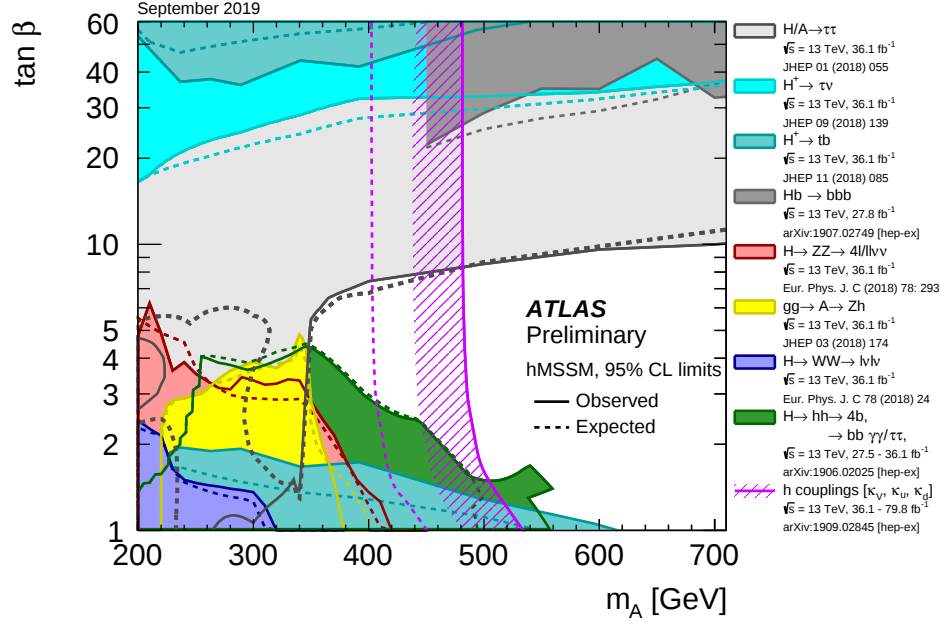


Figure 3.6: Regions of the $[m_A, \tan \beta]$ plane excluded in the $hMSSM$ model via direct searches for heavy Higgs bosons and fits to the measured rates of observed Higgs boson production and decays [85]. Limits are quoted at 95% CL and are indicated for the data (solid lines) and the expectation for the SM Higgs sector (dashed lines). The light shaded or hashed regions indicate the observed exclusions.

3.2.2 Two-Higgs-doublet plus singlet model

As could be seen in Figure 3.6, a wide region of phase space has been covered for the additional Higgses with mass > 125 GeV. Unfortunately no significant evidence of discovery has been found. More extensions of the 2HDM have been proposed. One of the simplest models is 2HDM+S, which adds a pseudo Nambu-Goldstone boson (PNGB) of a new approximate global symmetry. This PNGB could be significantly lighter than the other spin-0 particles. A well known example of a model that includes such a light pseudoscalar (a) is the next-to-minimal supersymmetric standard model (NMSSM) [86]. Since in this case the amount of symmetry breaking turns out to be proportional to soft breaking trilinear terms, the mass of a can be less than half of the SM Higgs mass. The most distinctive consequence of this model are exotic decays of Higgs $h \rightarrow aa$ for $m_a < \frac{m_h}{2}$, as well as $h \rightarrow Za$ for $m_a < m_h - m_Z$. In this work, we will focus on the case of $h \rightarrow aa$ for $m_a < \frac{m_h}{2}$.

3.2.3 Pseudoscalar a coupling to SM Higgs

In 2HDM+S, a complex singlet S is added to the the 2HDM model:

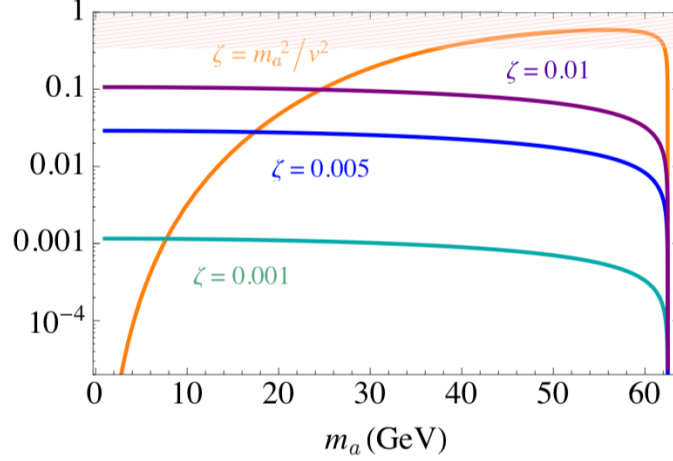


Figure 3.7: Exotic Higgs branching ratio to a pseudoscalar a versus the mass of the pseudoscalar m_a [78].

$$S = \frac{1}{\sqrt{2}}(S_{\text{Re}} + S_{\text{Im}}) \quad (3.10)$$

The complex field S only couples to the two Higgs doublets $H_{1,2}$ but has no direct Yukawa couplings, acquiring all of its coupling to SM fermions through the mixing with the two Higgs doublets. In this setup, there are totally two pseudoscalar states in the 2HDM+S model, one is mostly A from the two doublets, and one that is mostly S_{Im} , which could be written as the mostly-singlet-like pseudoscalar [78]:

$$a = \cos \theta_a S_{\text{Im}} + \sin \theta_a A \quad (3.11)$$

where $\theta_a \ll 1$.

Two terms in the effective Lagrangian give rise to $h \rightarrow aa$ decays [78]:

$$\begin{aligned} \mathcal{L} &\supset g_{hAA} hAA + \lambda_S |S^2|^2 \\ &= g_{hAA} \sin^2 \theta_a haa + 4\lambda_S v_s \zeta \cos^2 \theta_a haa \end{aligned} \quad (3.12)$$

where v_s is the singlet vacuum expectation value and ζ determines the singlet scalar content of the SM-like Higgs. The first term by itself can easily give rise to $\text{Br}(h \rightarrow aa) \approx 10\%$ if $g_{hAA} \approx v$ and $\theta_s \approx 0.1$, see Figure 3.7. The second term indicates that $\text{Br}(h \rightarrow aa)$ and $\text{Br}(h \rightarrow Za)$ can be independently adjusted.

3.2.4 Pseudoscalar a decays to SM particles

The decay of a to SM fermions proceeds via the couplings to A multiplied by $\sin\theta_a$. Therefore, once the type of 2HDM has been specified, the couplings between a to SM particles are predicted in the model. Figure 3.8 3.9 3.10 3.11 shows the results of $\text{Br}(a \rightarrow XX)$ [78], which could be summarized as follows:

- **Type-I:** since all fermions couple to H_2 , the branching ratios are independent of $\tan\beta$. The pseudoscalar a coupling to all fermions are proportional to those of the SM Higgs, as seen in Figure 3.8.

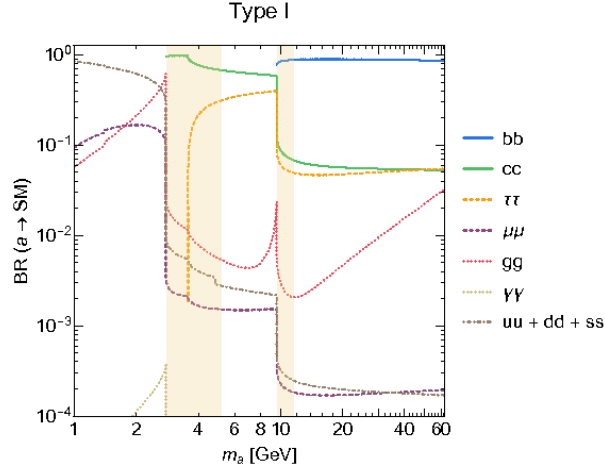


Figure 3.8: Branching ratio of a singlet-like pseudoscalar in the 2HDM+S for Type-I Yukawa couplings [78]. Decays to quarkonia likely invalidate the calculations in the shaded regions, where branching ratio is independent of $\tan\eta$

- **Type-II:** the exotic decay branching ratios are those of NMSSM models. Unlike Type-I models, they are now dependent on $\tan\beta$ with decays to down-type fermions suppressed or enhanced for $\tan\beta$ smaller or greater than 1, respectively. In this type of model, the pseudoscalar a corresponds to the R-symmetry limit of the NMSSM, as seen in Figure 3.9.
- **Type-III:** the branching ratios are dependent on $\tan\beta$. For $\tan\beta > 1$, pseudoscalar decays to leptons are enhanced. When $\tan\beta > 1$, this type of model suggests searching for 4τ over the entire mass range above the $b\bar{b}$ threshold, as seen in Figure 3.10.
- **Type-IV:** the branching ratio are $\tan\beta$ dependent. Compared to NMSSM, when $\tan\beta < 1$ the pseudoscalar decays to up-type quarks and leptons can be enhanced. As seen in Figure 3.11.

The coupling between the pseudoscalar a and other SM particles are Yukawa-like couplings, which means that the branching ratio is dependent on the masses of the decay products. The dominant decay mode will

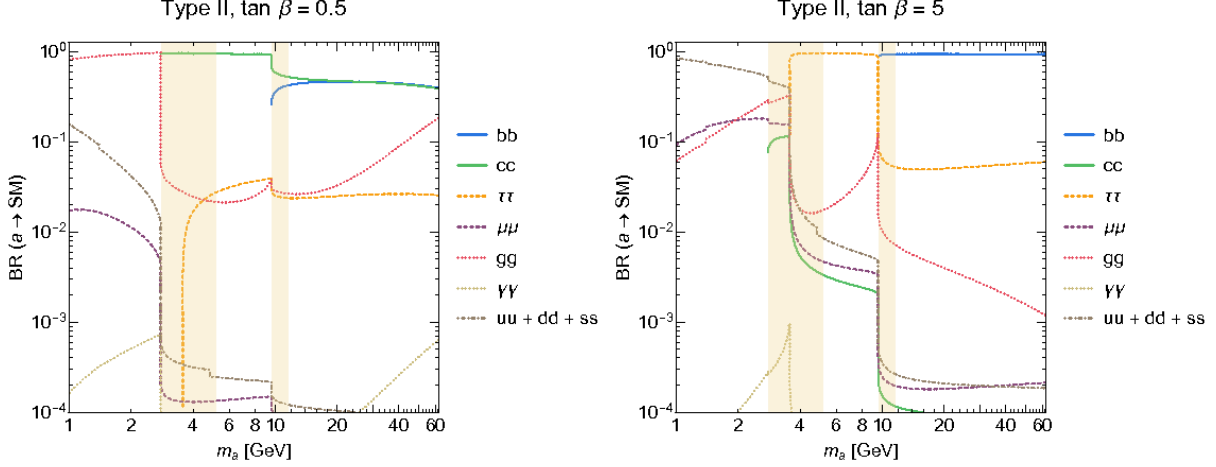


Figure 3.9: Branching ratio of a singlet-like pseudoscalar in the 2HDM+S for Type-II Yukawa couplings [78]. Decays to quarkonia likely invalidate the calculations in the shaded regions, where $\tan \beta = 0.5$ (left) or $\tan \beta = 5$ (right).

be $a \rightarrow ff$ where f is the heaviest fermion satisfying $m_a \geq m_f$. Work done in this dissertation is guided by this theoretical framework, where an analysis for $H \rightarrow 2a \rightarrow 4\mu$ is aiming sensitivities at small m_a from 1 to 15 GeV, and $H \rightarrow 2a \rightarrow 4\tau$ analysis expects good sensitivities for higher m_a from 15 to 60 GeV.

3.3 Pseudoscalar portal to dark matter

A pseudoscalar a can also provide a possible portals between dark matter and SM particles. Assuming DM to be a Dirac fermion χ with mass m_χ , the coupling to a pseudoscalar mediator a would be given by

$$\begin{aligned}\mathcal{L}_{\text{DM}} &= y_\chi a \bar{\chi} i \gamma^5 \chi \\ &= y_\chi (\cos \theta_a S_{\text{Im}} + \sin \theta_a A) \bar{\chi} i \gamma^5 \chi\end{aligned}\tag{3.13}$$

The Fermi collaboration has published limits on DM annihilation into final states containing photons [87]. By analyzing data taken from the Fermi Gamma Ray Space Telescope, an excess of gamma rays of energy $\approx 1 - 3$ GeV was found in the region of the galactic center. This excess's spectrum has been fit by DM annihilation to a number of final states, such as 10 GeV DM annihilating to $\tau^+ \tau^-$ [88] and 30 GeV to $b\bar{b}$ [89]. This possibility suggests searching for the cases when a light pseudoscalar plays the role as the mediator between DM and SM fermions [90].

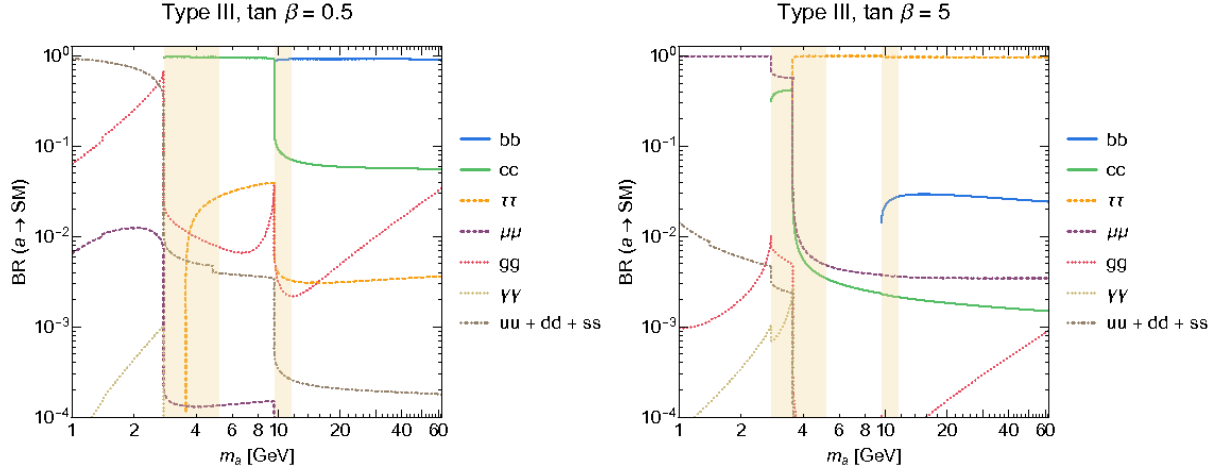


Figure 3.10: Branching ratio of a singlet-like pseudoscalar in the 2HDM+S for Type-III Yukawa couplings [78]. Decays to quarkonia likely invalidate the calculations in the shaded regions, where $\tan \beta = 0.5$ (left) or $\tan \beta = 5$ (right).

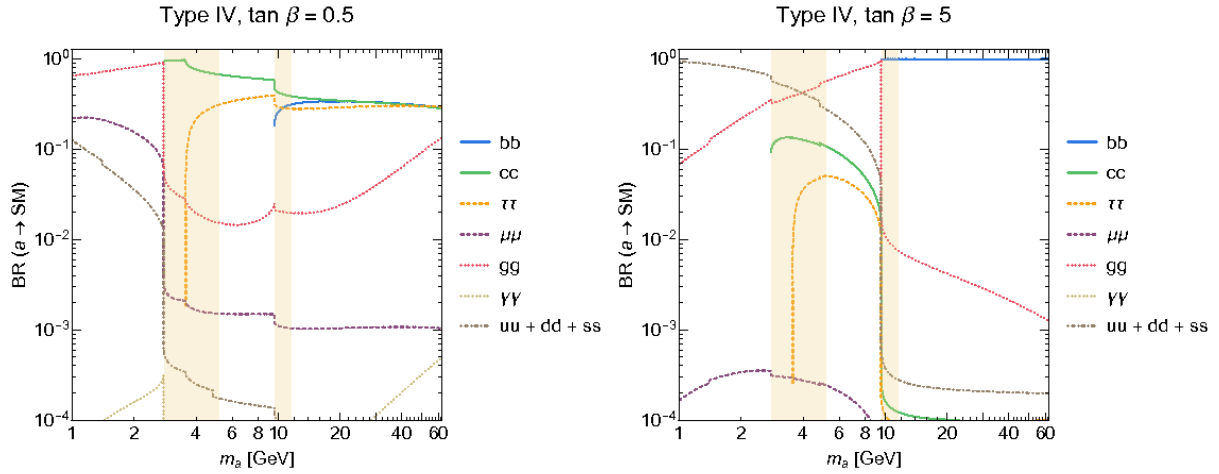


Figure 3.11: Branching ratio of a singlet-like pseudoscalar in the 2HDM+S for Type-IV Yukawa couplings [78]. Decays to quarkonia likely invalidate the calculations in the shaded regions, where $\tan \beta = 0.5$ (left) or $\tan \beta = 5$ (right).

Chapter 4

Large Hadron Collider and the ATLAS experiment

This chapter provides a brief introduction about [Large Hadron Collider \(LHC\)](#), the world's largest and most powerful particle accelerator. More details about the technical design, construction and operation of [LHC](#) could be found in [\[91\]](#).

The rest of this chapter will provide a review of the ATLAS detector, which mainly covers the ATLAS detector instrumentation used in this work. The final part of this chapter describes the studies about the upgrade of ATLAS detector in [LS2](#), which mainly focus on my simulation studies of the [Micro-MEsh Gaseous Structure \(MicroMegas\)](#) in [New Small Wheel \(NSW\)](#) in the [muon spectrometer \(MS\)](#).

4.1 Overview of the Large Hadron Collider

The [Large Hadron Collider \(LHC\)](#) [\[91\]](#) is a super-conducting particle accelerator with perimeter of 27 km allocated 100 m underground at the border of France and Switzerland, belong to [Organisation européenne pour la recherche nucléaire](#)(European Organization for Nuclear Research) ([CERN](#)). [Figure 4.1](#) shows a diagram of [LHC](#) outlook, as well as the major experiments operating on the ring of the collider: [ATLAS](#) [\[93\]](#), [LHCb](#) [\[94\]](#), [ALICE](#) [\[95\]](#) and [CMS](#) [\[96\]](#). Proton beams are accelerated by the superconductor magnets which provide a strong magnetic field of 8.3T, then collided in opposite directions at locations of these four major experiments. The designed aim for [LHC](#) is to reveal the physics [beyond the Standard Model](#) by colliding the energetic proton beams in a [center of mass \(CM\)](#) energy up to $\sqrt{s} = 14$ TeV.

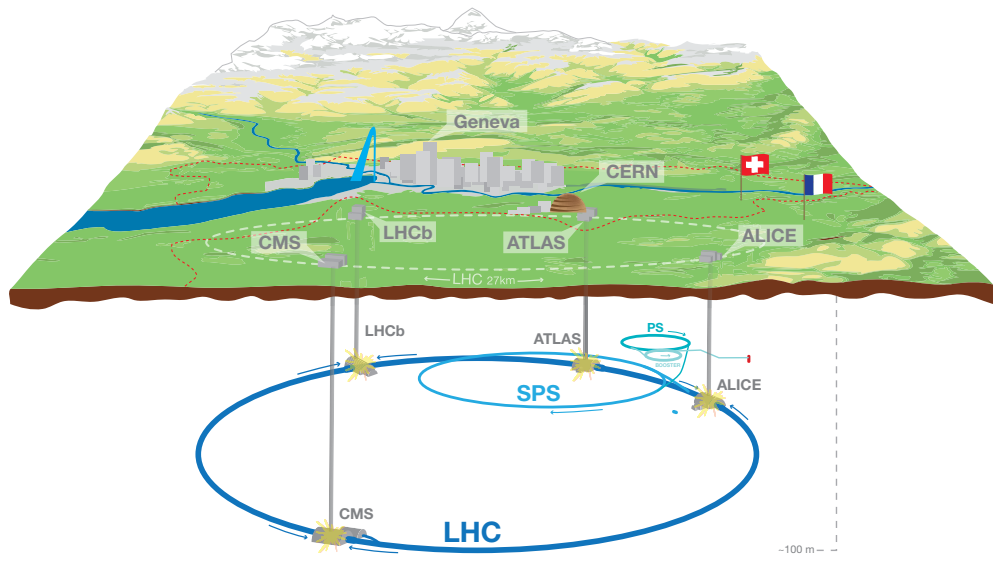


Figure 4.1: An overall view of LHC, sitting on the border between France and Switzerland, near the city of Geneva [92]. Four experiments located on the ring of the collider: ATLAS [93], LHCb [94], ALICE [95] and CMS [96].

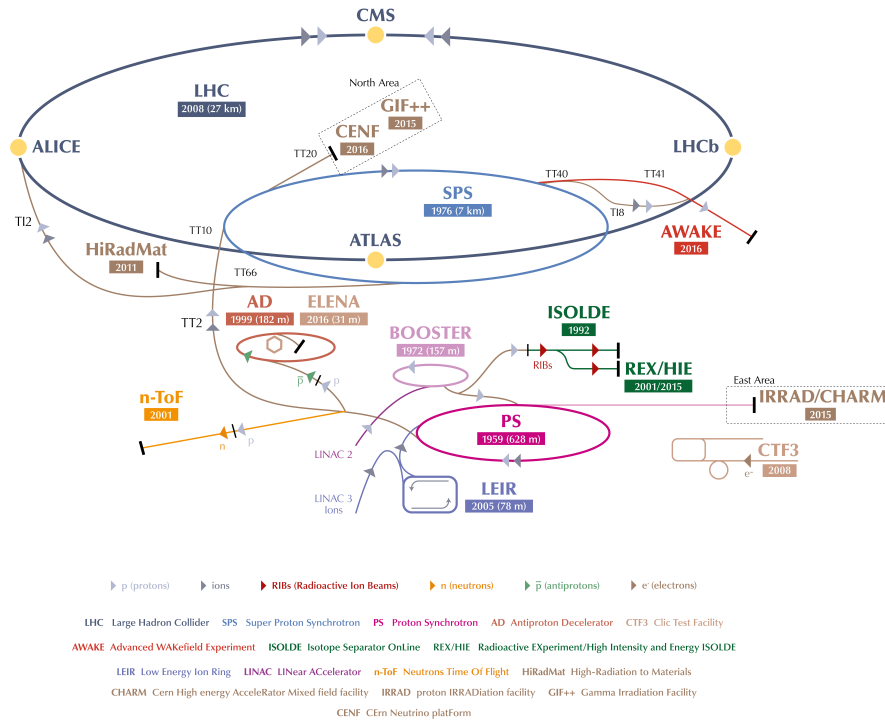


Figure 4.2: The LHC is the last ring in a complex chain of particle accelerator [97]. The smaller machines are used in a chain to help boost the particles to their final energies and provide beams to a whole set of smaller experiments, include Linac2, PSB, PS and SPS, and eventually arrive the ring of LHC.

The **LHC** is the last ring in a complex chain of particle accelerator. The smaller machines are used in a chain to help boost the particles to their final energies and provide beams to a whole set of smaller experiments, as seen in Figure 4.2. The proton beams are firstly injected to **Linear accelerator 2 (Linac2)** [98] and accelerated to energy of 50 MeV. After that, the proton beams travel to the **Proton Synchrotron Booster (PSB)** which accelerates them to 1.4 GeV. After **PSB**, the proton beams are injected to **Proton Synchrotron (PS)** and reach the energy of 25 GeV. The **SPS** then accelerates the proton beams up to 450 GeV. Eventually, the proton beams are injected to the ring of **LHC**.

4.1.1 Parameters of **LHC** operations

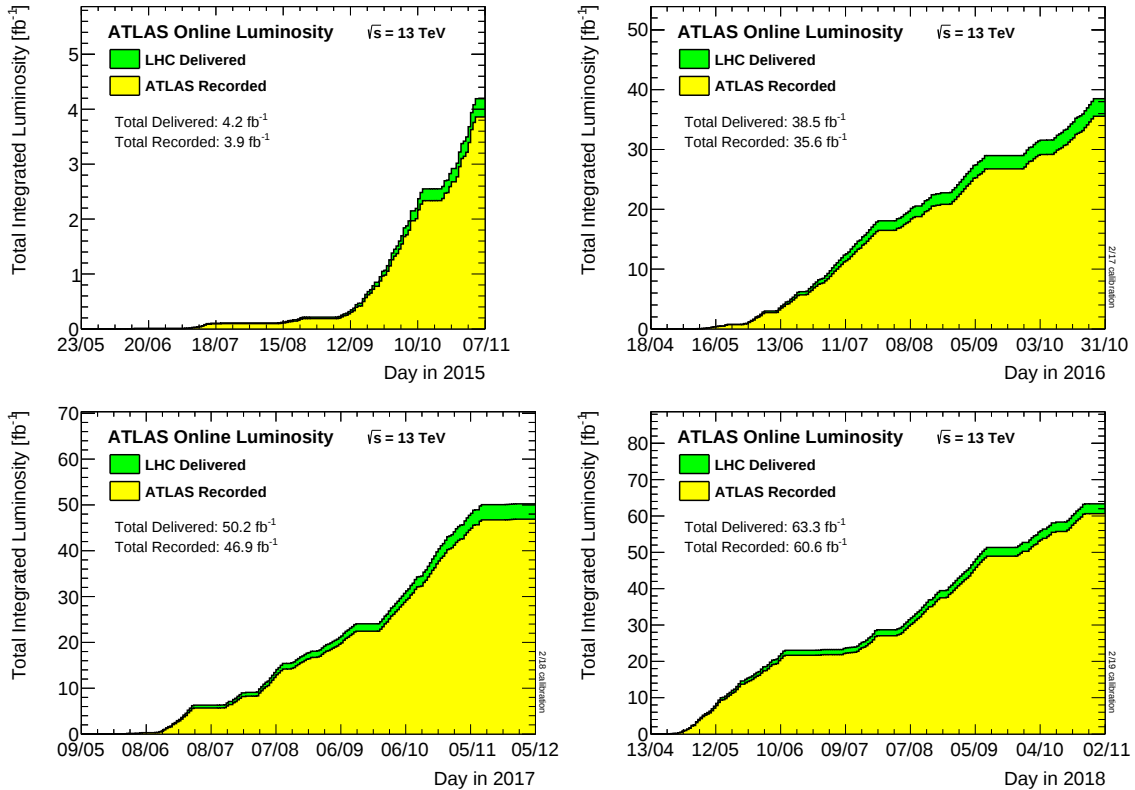


Figure 4.3: Cumulative luminosity versus time delivered to (green) and recorded by ATLAS (yellow) during stable beams for **pp** collisions at 13 TeV **center of mass** in 2015 (top left), 2016 (top right), 2017(bottom left) and 2018(bottom right) [99]. The delivered luminosity accounts for luminosity delivered from the start of stable beams until the **LHC** requests ATLAS to put the detector in a safe standby mode to allow for a beam dump or beam studies.

The frequency of bunch crossing at **LHC** is parameterized as instantaneous luminosity $L_{\text{inst.}}$. The integrated luminosity $L_{\text{int.}}$ is the parameter that is proportional to the total number of bunch collisions recorded

in the data taking period, which could be written as the integral as:

$$L_{\text{int.}} = \int L_{\text{inst.}} dt \quad (4.1)$$

For a given integrated luminosity $L_{\text{int.}}$, the number a certain physics process N with cross section σ the could be expressed as $N = \sigma L_{\text{int.}}$.

The instantaneous luminosity $L_{\text{inst.}}$ is defined as [91]:

$$L_{\text{inst.}} = \frac{N_p^2 n_p f_{\text{rev.}} \gamma_r}{4\pi \epsilon_n \beta^*} F \quad (4.2)$$

where N_p is the numbers of particles per bunch, n_p is the number of bunches per proton beam, $f_{\text{rev.}}$ is the revolution frequency, γ_r is the relativistic γ -factor, ϵ_n is the normalized transverse beam emittance, and β^* is the β function that describes the transverse size of the particle beam at the [interaction point \(ip\)](#).

In Run2 of [LHC](#), the integrated luminosity of each year is shown in Figure 4.3. The total L_{int} is 140.3 fb^{-1} as shown in Figure 4.4, which is used in this thesis. F is the geometrical correction factor that takes into account the angle of the bunch crossing.

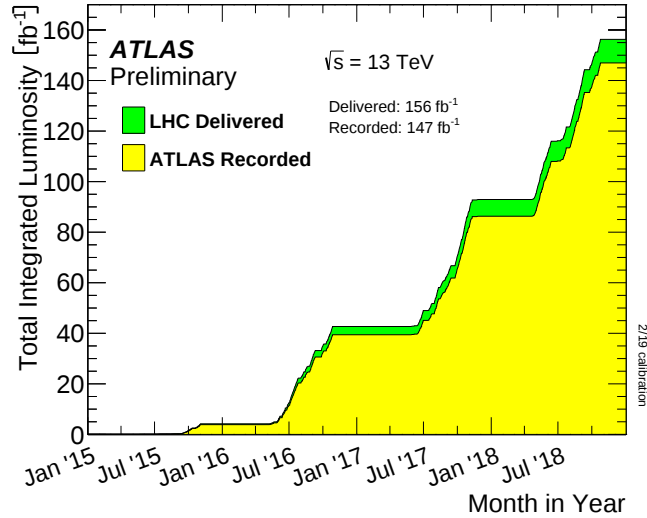


Figure 4.4: Cumulative luminosity versus time delivered to ATLAS (green) and recorded by ATLAS (yellow) during stable beams for [pp](#) collisions at 13 TeV [center of mass](#) energy in [LHC](#) Run 2 [99].

Another significant parameter is the number of particle interactions per bunch crossing μ . The multiple [proton-proton \(pp\)](#) interactions in each bunch crossing is usually known as **pile-up** and comes from two effects:

- **in-time** pile-up, which means multiple [proton-proton](#) interactions in the same bunch crossing.

- **out-of-time** pile-up, which refers to the effects from previous **proton-proton** interactions that before the bunch crossing that is currently being recorded. This effect is mainly due to the delay from the response of the detector electronics.

The time average pile-up is defines as $\langle\mu\rangle$. In **LHC** run2, $\langle\mu\rangle \approx 33.7$, as shown in Figure 4.5.

The effect of pile-up is simulated in the **MC** generations. When running on **MC**, the random run and lumiblock numbers are applied which are provided by a reweighting calculated based on the pile-ups [100].

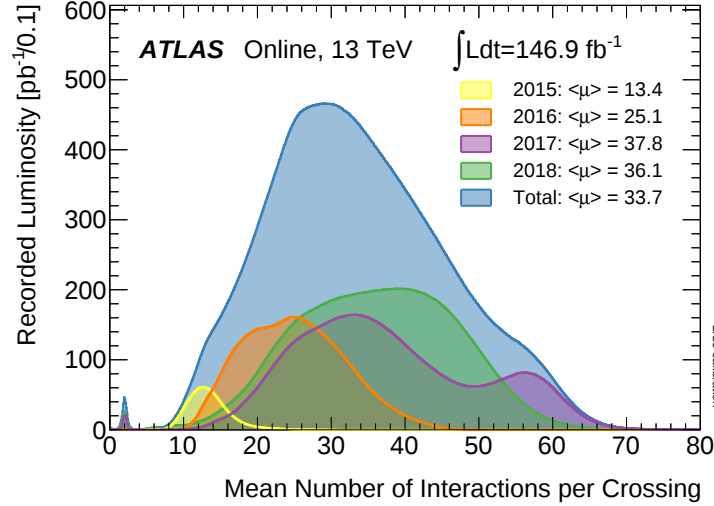


Figure 4.5: The mean number of interactions per bunch crossing for **pp** collision data at **center of mass** energy $\sqrt{s} = 13$ TeV during the **LHC** run2 [99]. The distributions for each year are weighted by the corresponding luminosity.

4.1.2 LHC and HL-LHC upgrade plans

The operation of **LHC** is planned as several run period. The first period Run 1 starts from 2011, with **center of mass** energy of $\sqrt{s} = 7$ TeV at 2011 and later $\sqrt{s} = 8$ TeV at 2012. The **LHC** Run 2 is during 2015 to 2018, with **center of mass** energy of $\sqrt{s} = 13$ TeV. The future Run 3 is planned to start from 2021, with **center of mass** energy increased to $\sqrt{s} = 14$ TeV. And finally after a set of detector and accelerator upgrades, **LHC** will run with its maximum designed **center of mass** energy of $\sqrt{s} = 14$ TeV with the luminosity increased by 5 to 7 times the current one, which is known as **High-Luminosity LHC (HL-LHC)**. Figure 4.6 shows the whole plan of **LHC** up to **HL-LHC**.

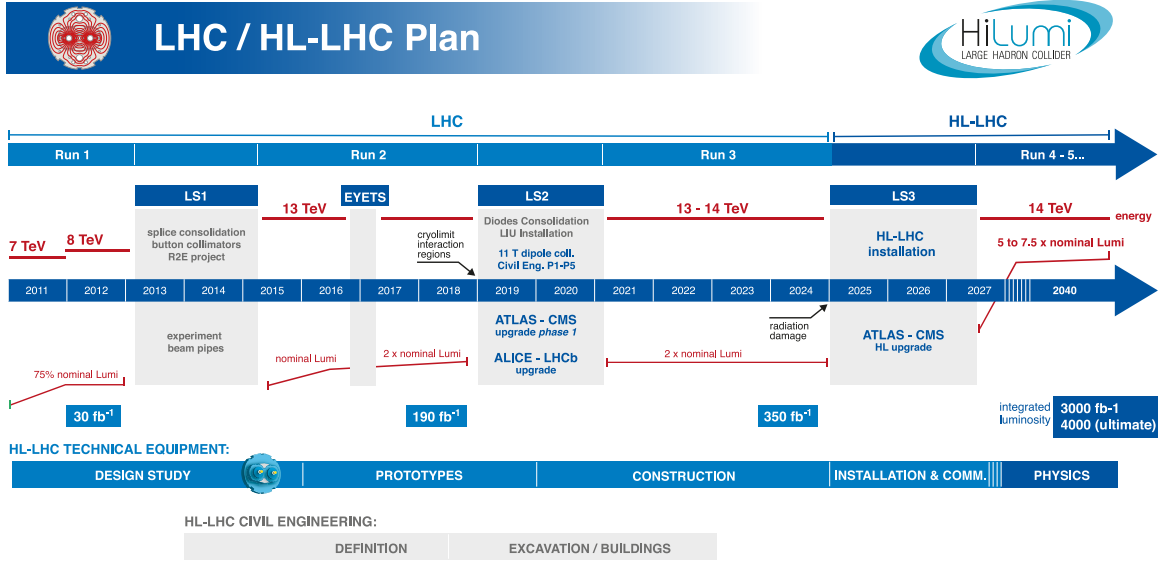


Figure 4.6: Schedule of LHC and HL-LHC from 2011 to 2040 [101].

4.2 ATLAS detector

The ATLAS detector is one of the most complex particle detectors which have ever been today, with 100 million electronic channels and over 3000 km of cabling. The detector is located 100 m underground at the LHC Point 1, centered around the LHC beam pipe. Figure 4.7 shows a cut-away view of the ATLAS detector, which is 25 m high and 44 m long in dimensions.

This section describes the detector details which are relevant to the work of this thesis. In Section 4.2.1, the detector coordinate system will be setup and used in later sections. Section 4.2.2, Section 4.2.3 and Section 4.2.4 will describe details in different detector systems. And finally Section 4.2.5 will introduce the trigger system in ATLAS.

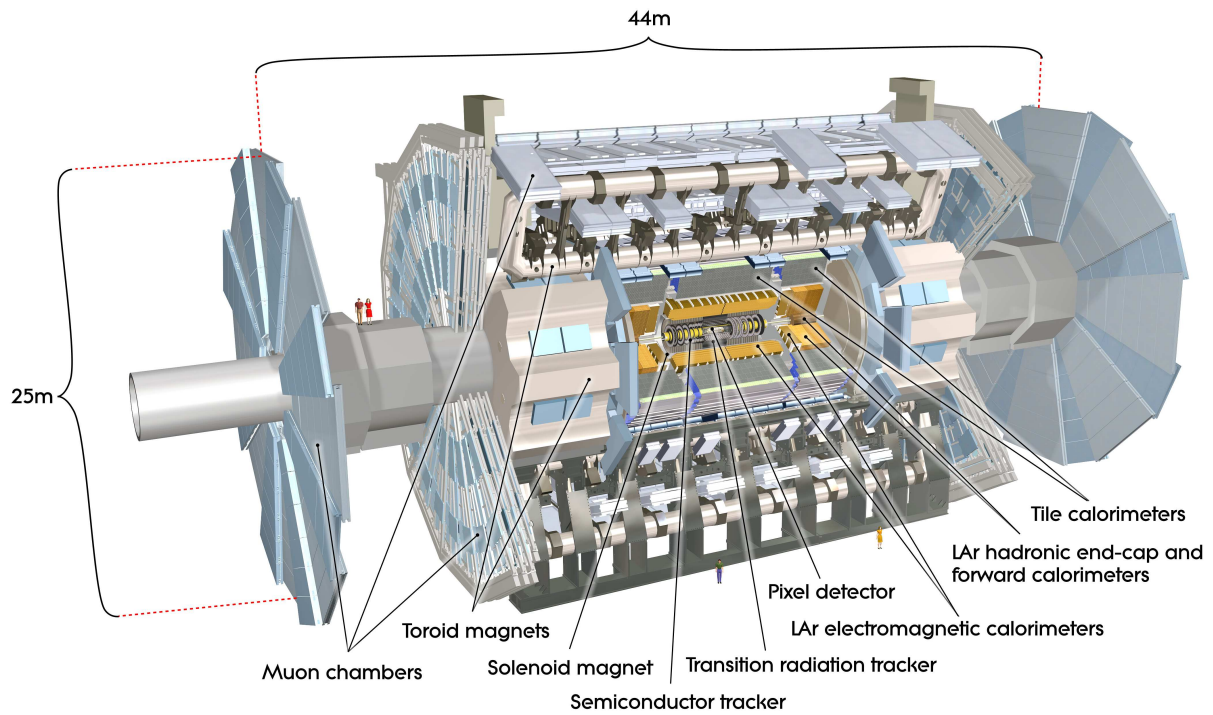


Figure 4.7: Cut-away view of the ATLAS detector [93]. The dimensions of the detector are 25 m in height and 44 m in length. The overall weight of the detector is approximately 7000 tonnes.

4.2.1 Coordinate system of ATLAS detector

ATLAS uses a right-handed coordinate system, where the origin is at the nominal [interaction point](#) (ip) in the center of the detector: the positive x -axis is defined as the direction pointing towards the center of the LHC ring, the positive y -axis is pointing upwards, and the beam direction is the z -axis is defined as the beam direction, as seen in Figure 4.8.

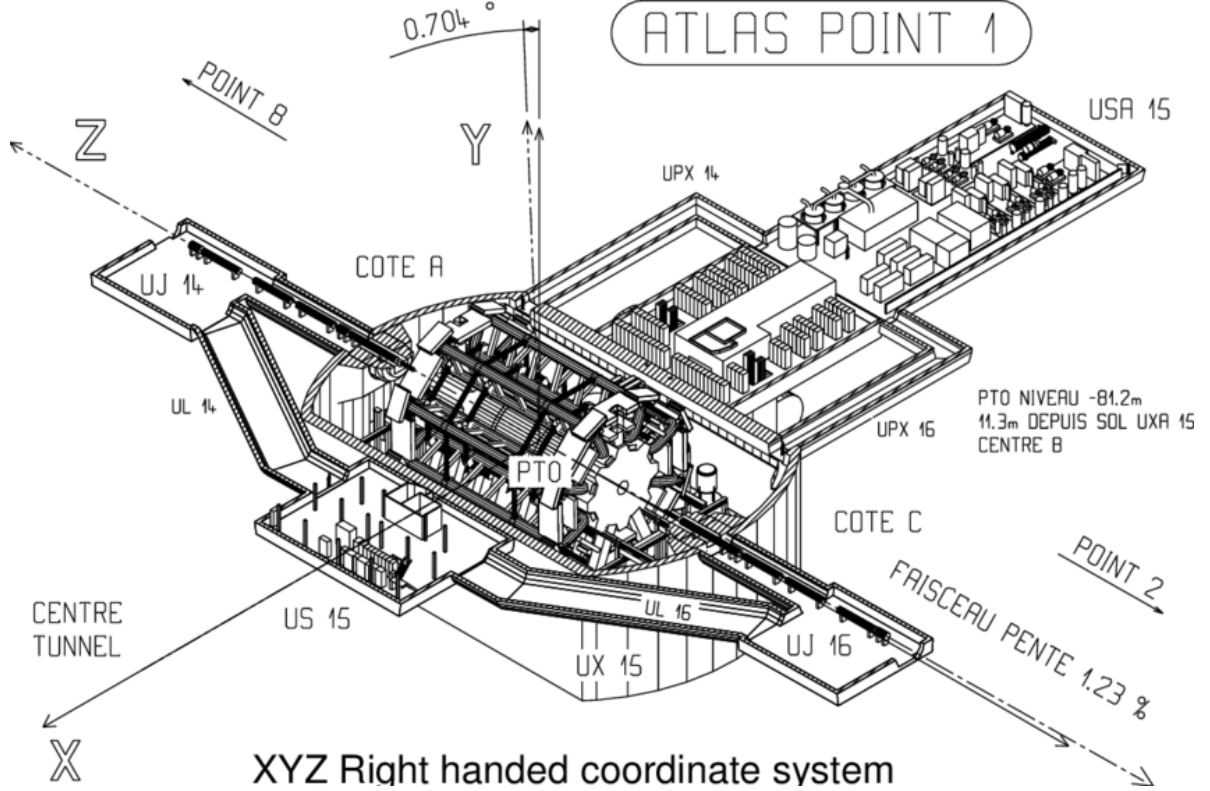


Figure 4.8: Illustration of the ATLAS Detector oriented in the global coordinate system [102].

The $x - y$ plane is perpendicular to the beam direction, and is referred as the transverse plane. A cylindrical coordinate system is used in the transverse plane, labeling as (r, ϕ) , where ϕ is the azimuthal angle about the z -axis.

The 3-dimensional momentum of physics objects measured in ATLAS are usually described by $\vec{p} = (p_T, p_z)$, where p_T is known as the transverse momentum, which is the momenta projection in the transverse plane. Due to the fact of momentum conservation, the vector sum of the transverse momentum is ideally zero: $\sum p_T = 0$.

Another kinematic property of the physics objects is the pseudo-rapidity η , which is defined in terms of

the polar angle θ as

$$\eta = -\ln \tan \frac{\theta}{2} \quad (4.3)$$

. When $\eta = 0$, the direction of the particle is pointing perpendicular to the beam axis, whereas large values of $|\eta|$ mean that the direction is close to the beam pipe, which is pointing to the end-cap region of the detector. The pseudo rapidity is an approximation of the rapidity y defined as

$$y = \frac{1}{2} \ln \frac{E + p_z}{E - p_z} \approx \eta - \frac{\cos \theta}{2} \left(\frac{m}{p_T} \right)^2 \quad (4.4)$$

, where E is the energy of the particle, and p_z is the particle momentum projection along the beam direction, and m is the mass of the object.

Since the detector has an onion-like structure, the most important information about the position of an object is (η, ϕ) . The distance between the two objects in $\eta - \phi$ plane could be expressed by ΔR as

$$\Delta R = \sqrt{\eta^2 + \phi^2} \quad (4.5)$$

4.2.2 Tracking in the inner detector

The **inner detector** (ID) [103, 104] is build around the beam pipe with a cylindrical geometry as sown in Figure 4.9. The **inner detector** is designed to be reconstruct tracks of charged particles tracks with excellent momentum resolution in a 2 T solenoid magnet, for particles with $p_T \geq 500$ MeV and $|\eta| < 2.5$.

The ATLAS **inner detector** contains three specific sub-components, as seen in Figure 4.10.

Pixel detector

The pixel detector [105] is the closest part to the beam pipe, which is composed of 80 million pixels covering 1.7 m^2 of $|\eta| < 2.5$ region. Each pixel has an area of $50 \mu\text{m} \times 400 \mu\text{m} = 20\,000 \mu\text{m}^2$. The resolution of each pixel is $14 \mu\text{m}$ in ϕ direction and $115 \mu\text{m}$ in z direction. The barrel part of ATLAS detector is covered by 3 pixel layers with 1456 pixel modules, each containing 46080 readout channels. In the end-cap area, there are three pixel disks with 144 modules and 6.6 million readout channels. Each barrel layer provides one measurement for each charged particle track, to reconstruct not only tracks, but also primary and secondary vertices.

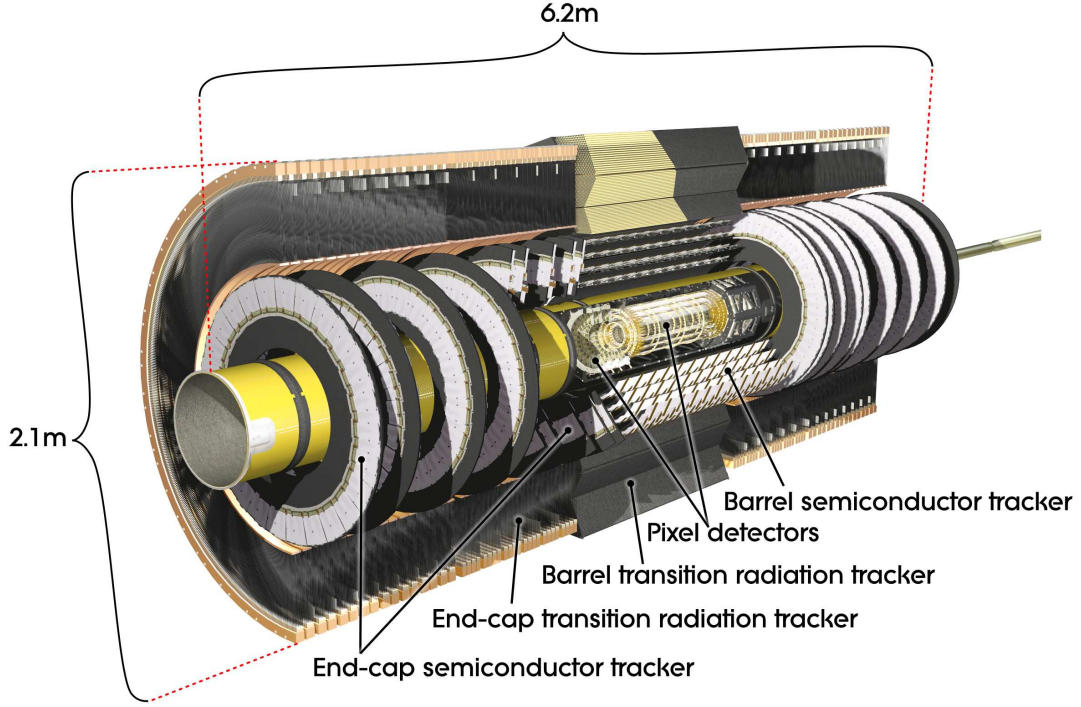


Figure 4.9: Cut-away view of the ATLAS inner detector [93].

Semiconductor Tracker

Semiconductor Tracker (SCT) [106] is a silicon micro-strip tracker, consisting of 4088 two-sided modules with over 6 million strips. The SCT covers the region of $|\eta| < 2.5$ surrounding the pixel detectors with an area of 63 m^2 . The strips are placed with a distance of $80 \mu\text{m}$ and rotated by 50 mrad with respect to each other. The position resolution of each SCT strip is $17 \mu\text{m}$ in the transverse plane and $580 \mu\text{m}$ along z -axis. For each track of the charged particle, the SCT will provide 4 to 9 precision measurement and contribute to the measurement of momentum, **impact parameter** and vertex identification, combined with the pixel detector.

Transition radiation tracker

Transition radiation tracker (TRT) [107] is made of 350000 straw drift tubes covering the volume of 12 m^3 volume for $|\eta| < 2.0$. Each drift tube has a diameter of 4 mm and length of 144 cm (37 cm) in the barrel(end-cap) region. TRT has a resolution of $130 \mu\text{m}$ in ϕ direction. There are totally 73 layers of tubes in the barrel region and 160 planes in the end-cap region, which provides transition radiation tracking for charge particle identification. Each charged particle will travel through at least 36 tubes, providing the measurement of the charge collection time in the tubes. The charged particle with a lower mass tends to emit more transition

radiation, hence TRT will provide strong discrimination between charged leptons and charged hadrons.

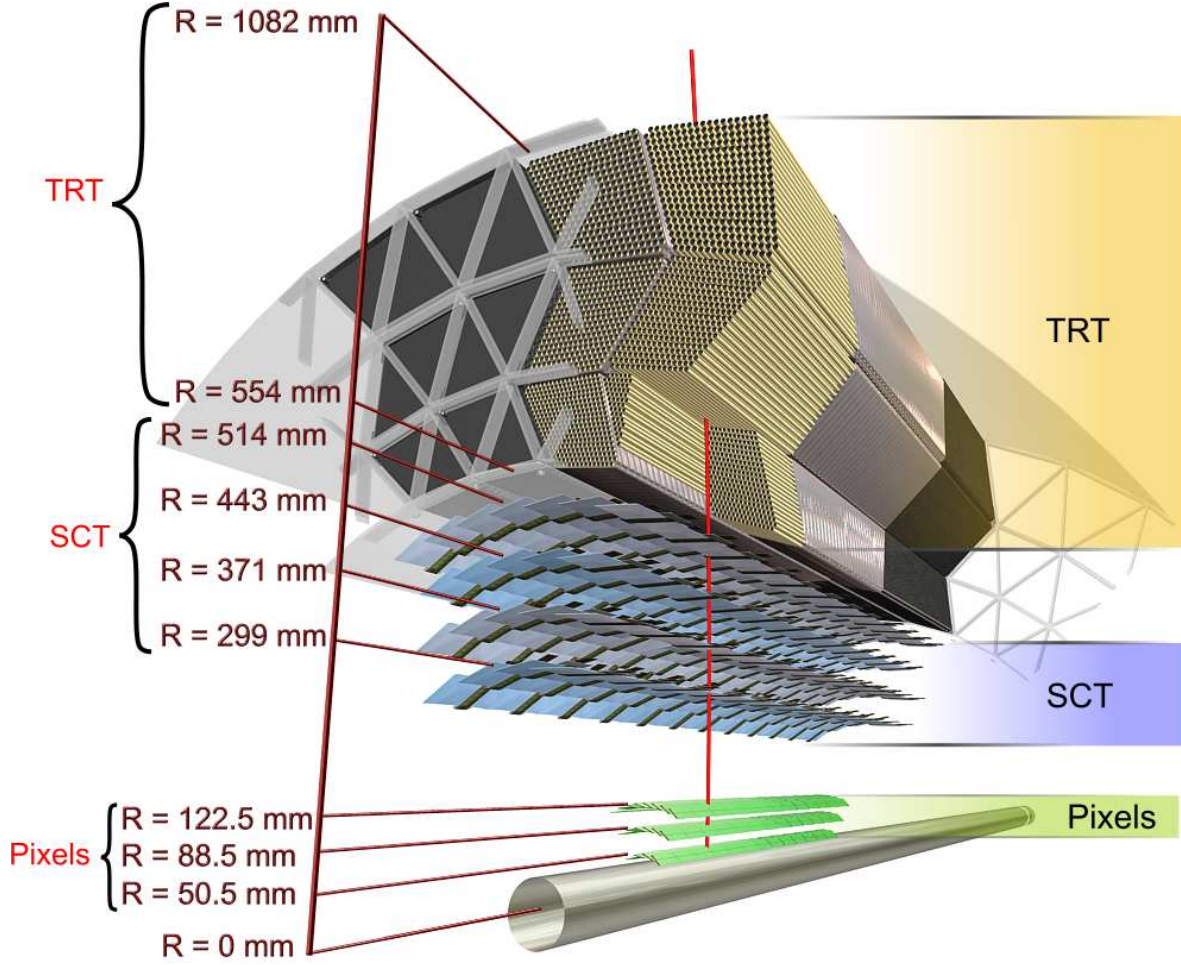


Figure 4.10: Drawing showing the sensors and structural elements traversed by a charged track of 10 GeV p_T in the barrel ID ($\eta = 0.3$) [93]. The track traverses successively the beryllium beam-pipe, the three cylindrical silicon-pixel layers with individual sensor elements of $50 \times 400 \mu\text{m}^2$, the four cylindrical double layers (one axial and one with a stereo angle of 40 mrad) of barrel SCT of pitch $80 \mu\text{m}$, and approximately 36 axial straws of 4 mm diameter contained in the barrel TRT modules within their support structure.

4.2.3 Calorimeter system

In ATLAS, the calorimeter system consists of two types of calorimeters: **hadronic calorimeter (HCal)** and **electromagnetic calorimeter (EMCal)** [108]. The HCal is to measure the energy of particles interacting via strong nuclear force (gluons and quarks), whereas the EMCal is designed to measure the particles interacting through electromagnetic force (mainly electrons and photons). The calorimeters system covers the range of $|\eta| < 4.9$ in the detector, and provides sufficient containment for both electromagnetic and hadronic showers. Both calorimeters are sampling calorimeters, which use the “active” material to provide

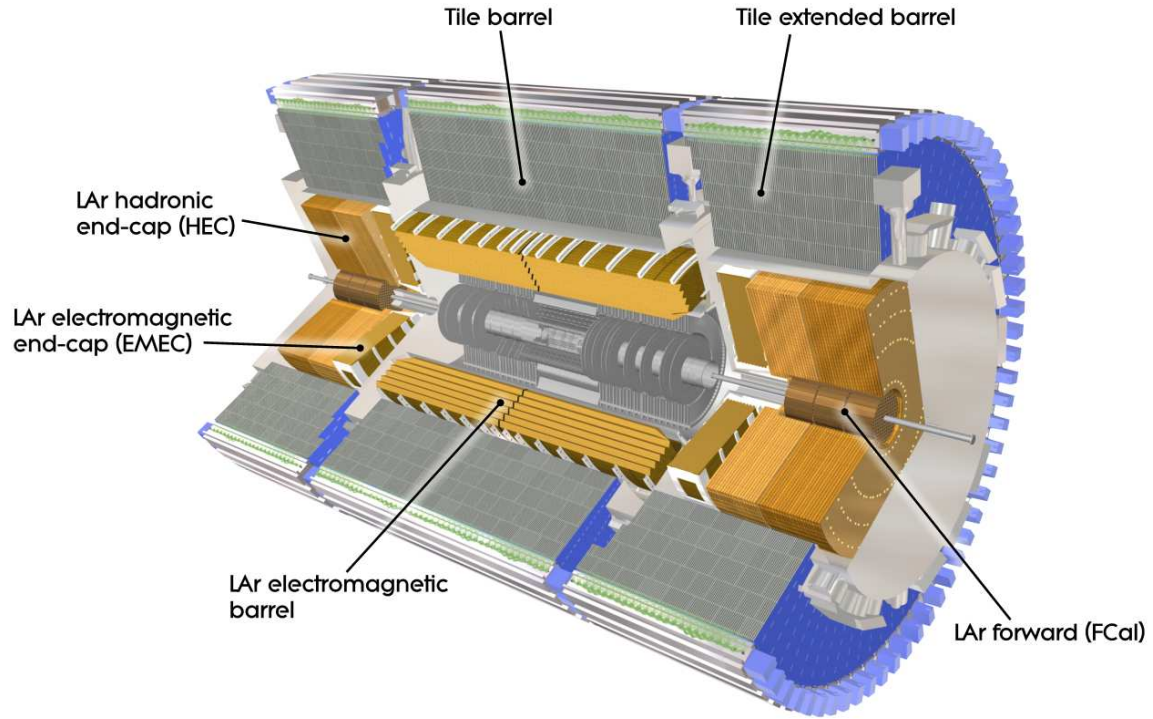


Figure 4.11: Cut-away view of the ATLAS calorimeter system [93], including EMB, EMEC, Tile, HEC and FCal.

the detectable signal which is different from a dense “absorber” material and reducing the particle energy. Therefore, only a fraction of energy is measurable by detector sensors, and needs a calibration to measure the calorimeter energy by studying the sensor response.

Figure 4.11 shows a cut-away view of the ATLAS calorimeter system, and Table 4.1 summarizes the different detectors used in ATLAS calorimeters.

One important characteristics of the calorimeter is the energy resolution. In ATLAS, the energy resolution of EMCal is [63]:

$$\frac{\sigma(E)}{E} = \frac{10\%}{\sqrt{E}} \oplus \frac{0.3}{E} \oplus 0.4\% \quad (4.6)$$

where the first term is the stochastic term, the second term represents the noise and the third term is a constant uncertainty.

Detector	Calorimeter	Material		Coverage	Number of channels
		Absorber	Active		
EMB	EMCal	lead	liquid argon	$ \eta < 1.475$	99712
EMEC	EMCal	lead	liquid argon	$1.375 < \eta < 3.2$	62208
Tile	HCal	steel	scintillator	$ \eta < 1.0$	5760
	HCal	steel	scintillator	$0.8 < \eta < 1.7$	4092
HEC	HCal	copper	liquid argon	$1.5 < \eta < 3.2$	5632
FCal	EMCal	copper	liquid argon	$3.2 < \eta < 4.9$	1008
	HCal	tungsten	liquid argon		754
Total					179166

Table 4.1: Summary of the sampling calorimeters, with the corresponding (absorber and active) materials, $|\eta|$ coverage and number of channels.

Electromagnetic calorimeter

ATLAS EMCal is formed by LAr electromagnetic barrel (EMB) covering the barrel part, and two LAr electromagnetic endcap calorimeter (EMEC) located at the end-cap region. EMB covers the region of $|\eta| < 1.475$. There are two EMEC, where the inner one covers $1.375 < |\eta| < 2.5$ and the outer one covers $2.5 < |\eta| < 3.2$. Between EMB and EMEC, there is a transition region that degrades the performance known as the “cracked region” for $1.37 < |\eta| < 1.52$. Most ATLAS analyses involving photons and electrons veto the objects in this region to exclude misidentified electrons or photons.

EMB and EMEC have three radial layers. The first sampling layer is the strips, which are finely segmented in $\Delta\eta = 0.0031$ with 8 strips in front of each cell. The strips layer measure the fine information which provide the discrimination for electron/photon showers from pions. The second layer is made of many fine distributed square cells, which collect the largest fraction of energy in electromagnetic showers. And the last layer collects the tail end of the electromagnetic shower. Figure 4.12 shows a sketch of a barrel module of EMCal with its three-layer structure.

Hadronic calorimeter

The ATLAS HCal is composed by Tile calorimeter (Tile) [109], LAr hadronic endcap calorimeter (HEC) and forward calorimeter (FCal). The Tile calorimeter is a sampling calorimeter using steel as the absorber and scintillator as the active material, covering the region of $|\eta| < 1.7$, located behind EMB. The HEC uses LAr as the active material and copper as the absorber, locating behind EMEC and covering $1.5 < |\eta| < 3.2$. Finally FCal extends the hadronic calorimeter range by covering $3.1 < |\eta| < 4.9$ with much coarser granularity.

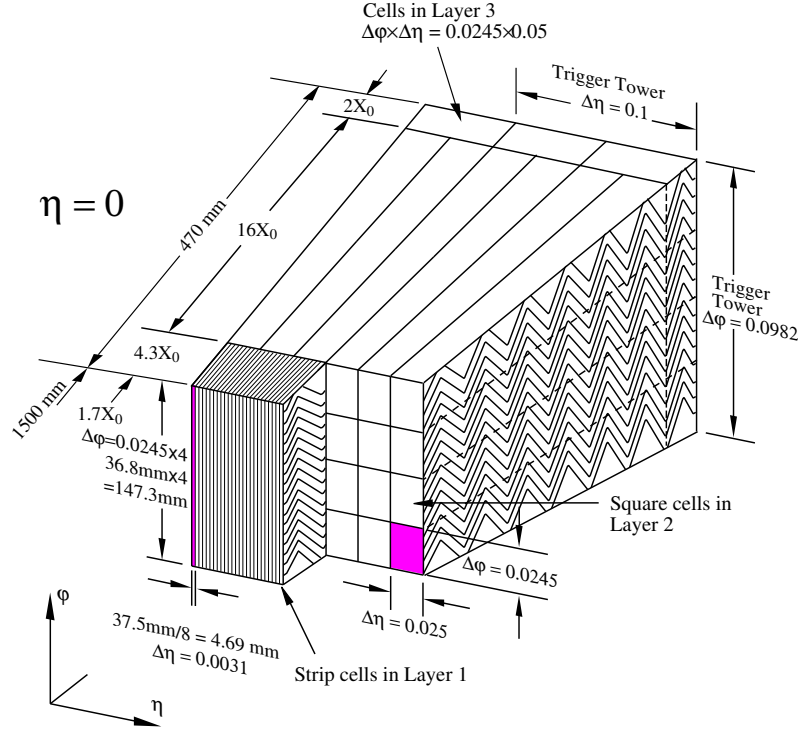


Figure 4.12: Sketch of a barrel module where the different layers are clearly visible with the ganging of electrodes in ϕ [93]. The granularity in η and ϕ of the cells of each of the three layers and of the trigger towers is also shown.

4.2.4 Muon spectrometer

Muon spectrometer (MS) [110] is the most outer part in ATLAS, as shown in Figure 4.13. Due to the nature of heavy mass, muons have a much smaller energy loss in Bremsstrahlung. Therefore muons can travel through the calorimeters. The ATLAS **muon spectrometer** system is based on the magnetic deflection of muon tracks in large superconducting toroid magnets. The muon tracks will be measured and used to calculate the momentum. The barrel area is covered by three cylindrical shells around the beam axis, called “stations”, with a radius of 5 m, 7.5 m and 10 m. The end-cap region of **muon spectrometer** contains three wheels perpendicular to the beam axis, named based on their radius, as one **Small Wheel (SW)** and two **Big Wheel (BW)**.

The ATLAS **muon spectrometer** is composed by **monitored drift tubes (MDT)** [111], **cathode strip chambers (CSC)** [112], **resistive plane chambers (RPC)** [113] and **thin gap chambers (TGC)** [114].

Monitored drift tubes

Monitored drift tubes (MDT) [111] are aluminum tubes with a diameter of 30 mm and thickness of 400 μm . Each tube is filled by gas which is a mixture of 93% Ar and 7% CO_2 . When a muon muon pass through

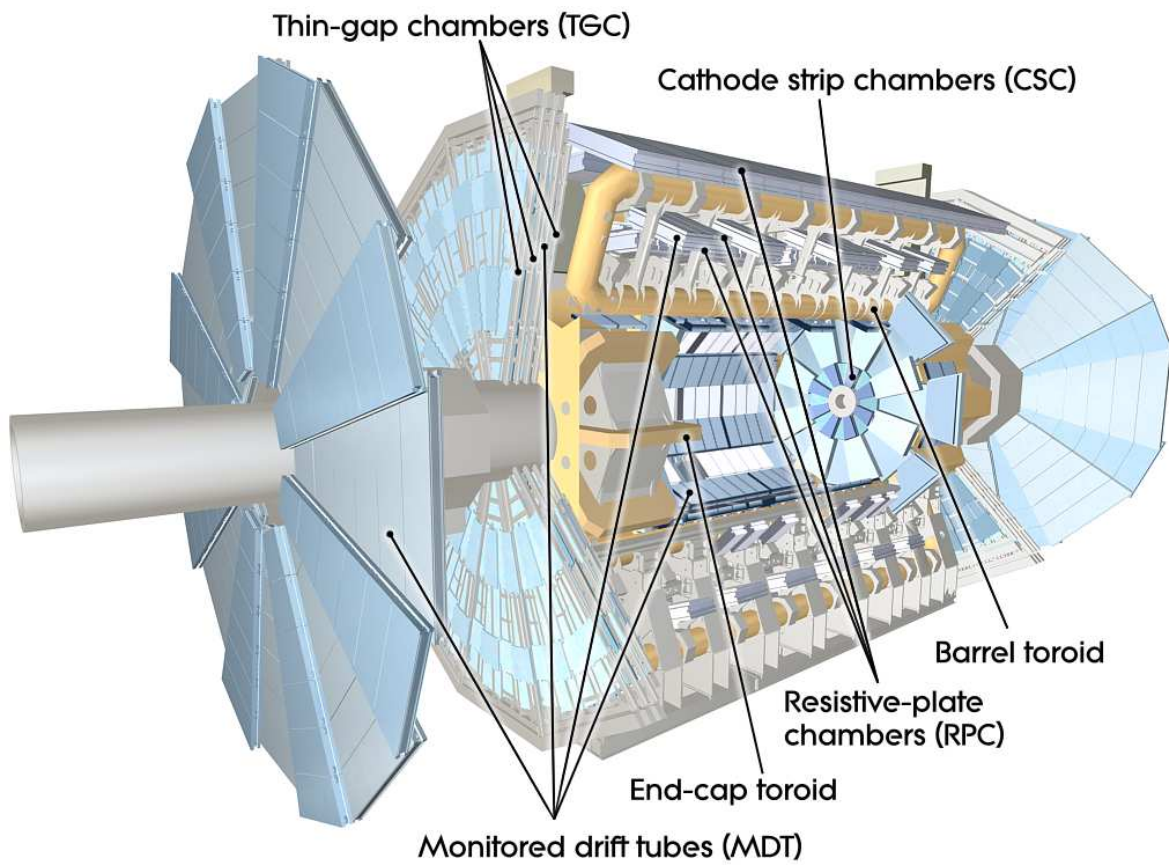


Figure 4.13: Cut-away view of the ATLAS muon system.

the tube, the gases will be ionized and produce an electron avalanches collected by a tungsten wire which is placed in the center of the tube. The resolution achieved by a single tube can be as good as 80 μm , whereas the resolution of the total chamber can reach as good as 35 μm .

Cathode strip chambers

Cathode strip chambers (CSC) [112] is used in the forward region, this is because of a much higher particle rate in the end-cap region and the performance of MDT is strongly suffered. CSC are multi-wire proportional chambers with multiple anode wires exposing in a gas mixture of 80% of Ar and 20% of CO_2 . CSC have a fast response with good spatial resolution, thus used to measure the charge deposition with a resolution of 40 μm in r -direction and 5 mm in ϕ -direction.

Resistive plane chambers

Resistive plane chambers (RPC) [112] are chambers filled with a gas mixture of $\text{C}_2\text{H}_2\text{F}_4$ and a small fraction of resistive component SF_6 , contained in two resistive parallel plates of bakelite. The passed muons will ionize the gas and the signal will be amplified by the electric field, which will be further readout by metallic strips. RPC are mainly used as triggers, where the chambers in the middle station triggers for low p_{T} muons, and the chambers in outer station are used for high p_{T} muon triggers.

Thin gap chambers

Similarly to CSC, thin gap chambers (TGC) [114] are used in the muon spectrometer end-cap region to improve the performance with very high particle rate. TGC are multi-wire proportional chambers, filled by a gas mixture of CO_2 and $\text{n-C}_5\text{H}_{12}$, with a similar structure as CSC but a higher granularity. TGC are used as trigger chambers, but also provide secondary measurement in azimuthal direction. TGC and RPC are working together to provide a trigger efficiency above 99% in the nominal LHC luminosity.

4.2.5 Trigger and data acquisition system

In ATLAS, trigger and data acquisition (TDAQ) [115] system is a crucial component of the experiment, responsible for select events of interest at a recording rate of ≈ 1 kHz from up to 40 MHz of proton-proton collisions, which equals to 25 ns per bunch crossing of LHC.

Figure 4.14 shows an overview workflow of ATLAS TDAQ system in LHC run2. The TDAQ system consists of a hardware-based first-level trigger (L1) and a software-based high-level trigger (HLT). The L1 trigger decision is produced from central trigger processor (CTP), which received inputs from L1 calorimeter

trigger ([L1Calo](#)) and [L1 muon trigger](#) ([L1Muon](#)), and also protect front-end readout buffers from overflowing. This effect is also known as “dead-time”, which is either the amount of time needed to allow readout windows to process data, or the empty time window of the downstream front-end buffers. The decision by [L1CTP](#) is called [L1 accpet](#) ([L1A](#)), of which the maximum rate is up to 100 kHz. Events are buffered in the readout system, and later processed by [HLT](#), which has an acceptance rate up to 1 kHz. The events passing the [HLT](#) decisions will be transferred to a local Tier-0 computing facility based at [CERN](#) for offline reconstructions. At the end of the day, the ATLAS detector will only save one event out of 40000 produced. Each stored event has a size of 1 – 2 MB, which means the readout system will write 1 – 2 GB/s.

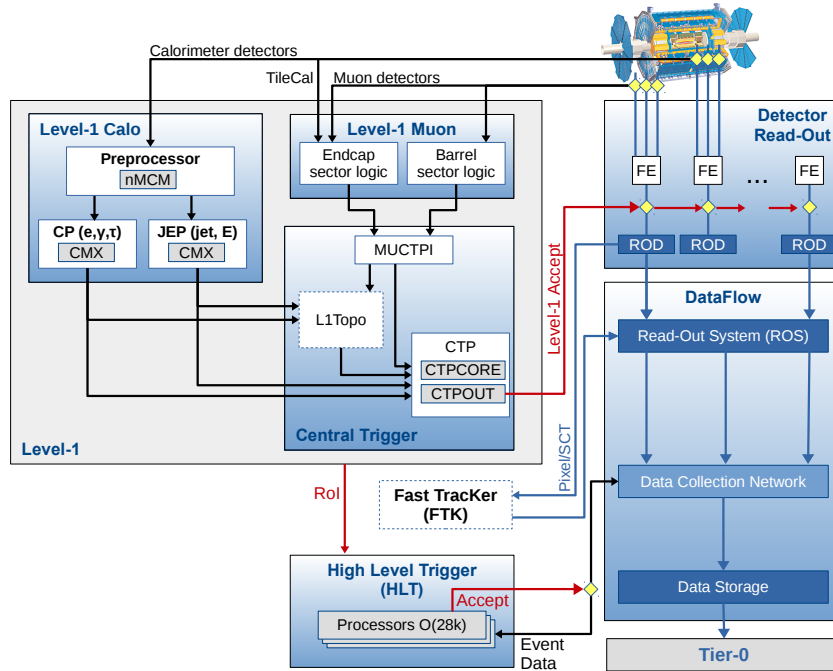


Figure 4.14: The ATLAS TDAQ system in run2 with emphasis on the components relevant for triggering [115].

ATLAS provides an XML file containing a list of events in data which have passed the data quality criteria. The corrupted events that occurs errors in [Tile](#), [LAr](#), [SCT](#) will be vetoed. This XML file is called [Good Run List](#) (GRL). In this analysis, full [LHC](#) Run-2 data will be used, and four different [GRL](#) are used, for 2015 data, 2016 data, 2017 data and 2018 data.

Trigger requirement will be implemented in both data and MC. Table 4.2 listed all unprescaled muon triggers used in ATLAS through out full Run2 of LHC. The lowest p_T thresholds for the single-lepton triggers ranged from 24 GeV to 26 GeV. Dielectron (dimuon) trigger thresholds ranged from 2×12 GeV (2×10 GeV) to 2×17 GeV (22, 8 GeV). Trielectron (trimuon) triggers had thresholds of 17, 9, 9 GeV (3×6 GeV). The triggers are only used when they are unprescaled (a separate GRL is maintained for each trigger, recording which lumiblocks the trigger was unprescaled for). When running on MC, the random run and lumiblock numbers are applied which are provided by a reweighting calculated based on the pile-ups [100].

4.3 ATLAS NSW upgrade

As mentioned in Section 4.2.4, the end-cap region of muon spectrometer usually have a high particle rate. This case will be more extreme with the increase of the instantaneous luminosity of LHC, where the MS end-cap region will suffer from a large rate with fake muons, as well as the lost of efficiency of L1 muon trigger. Therefore, Small Wheel in ATLAS muon spectrometer will be updated to New Small Wheel (NSW) [116] in the Phase-I upgrade of the ATLAS detector.

The proposed NSW detector system is designed to meet the requirement of high rate of particles for LHC Run 3 and HL-LHC. Figure 4.15 shows the layout of NSW, which consists of 16 detector planes in two multi-layers. Each multi layer comprises four small-strip TGC (sTGC) and four Micro-MESh Gaseous Structure (MicroMegas) detector planes. The trigger system in NSW will also be updated, which implemented look-up-tables to reject the muon tracks which are not orienting from interaction point (ip) [117].

This section describes some simulate studies to understand the MicroMegas performance at NSW. Section 4.3.2 presents the expected particle hit rate at MicroMegas in LHC Run 3 scenario, and Section 4.3.3 will present a GEANT4 simulation study for MicroMegas sensitivity for neutral particles, such as photons and neutrons.

4.3.1 Introduction of MicroMegas detector

A schematic of a MicroMegas detector is shown in Figure 4.16. The gas-filled chamber is split into two main regions: the **drift gap**, which is the region between the drift cathode and the woven mesh with an electric field of 0.6 kV/cm; and the **amplification gap**, which is the region of 39 kV/cm between the mesh and the strips. The gas is chosen to be a mixture of a noble and a poly-atomic gas. When a charged particle crossed the detector, if it is energetic enough, it ionizes the atoms of the gas and creates ion-electron pairs in the drift gap. The primary electrons from ionization further drift towards the mesh, while the ions towards the drift

ℓ	Trigger name	2015	2016	2017	2018
e	HLT_e24_lhmedium_L1EM20VH	✓			
	HLT_e24_lhmedium_iloose_L1EM18VH				
	HLT_e24_lhtight_nod0_ivarloose		✓		
	HLT_e26_lhtight_nod0				✓
	HLT_e26_lhtight_nod0_ivarloose		✓	✓	✓
	HLT_e60_lhmedium	✓	✓		
	HLT_e60_lhmedium_nod0		✓	✓	✓
	HLT_e60_medium				
ee	HLT_2e12_lhloose_L12EM10VH	✓			
	HLT_2e15_lhvloose_nod0_L12EM13VH		✓		
	HLT_2e17_lhvloose_nod0		✓		
	HLT_2e17_lhvloose_nod0_L12EM15VHI			✓	✓
	HLT_2e24_lhvloose_nod0			✓	✓
μ	HLT_mu20_iloose_L1MU15	✓			
	HLT_mu26_ivarmedium		✓	✓	✓
	HLT_mu50	✓	✓	✓	✓
	HLT_mu60_0eta105_msonly	✓	✓	✓	✓
$\mu\mu$	HLT_mu18_mu8noL1	✓			
	HLT_mu22_mu8noL1		✓	✓	✓
	HLT_2mu10	✓	✓		
	HLT_2mu14		✓	✓	✓
$e\mu$	HLT_e17_lhloose_mu14	✓			
	HLT_e17_lhloose_nod0_mu14		✓	✓	✓
	HLT_e24_lhmedium_nod0_L1EM20VHI_mu8noL1		✓		
	HLT_e26_lhmedium_nod0_mu8noL1			✓	✓
	HLT_e7_lhmedium_mu24	✓			
	HLT_e7_lhmedium_nod0_mu24		✓	✓	✓
$\mu\mu\mu$	HLT_mu18_2mu4noL1	✓			
	HLT_mu20_2mu4noL1		✓	✓	✓
	HLT_3mu4			✓	
	HLT_3mu6	✓	✓	✓	✓
	HLT_3mu6_msonly	✓	✓	✓	✓

Table 4.2: Triggers used in the 2015–8 analysis. Triggers are only used when they are unrescaled (a separate [GRL](#) is maintained for each trigger that records which lumiblocks the trigger was unrescaled). This maximizes the integrated luminosity available for analysis.

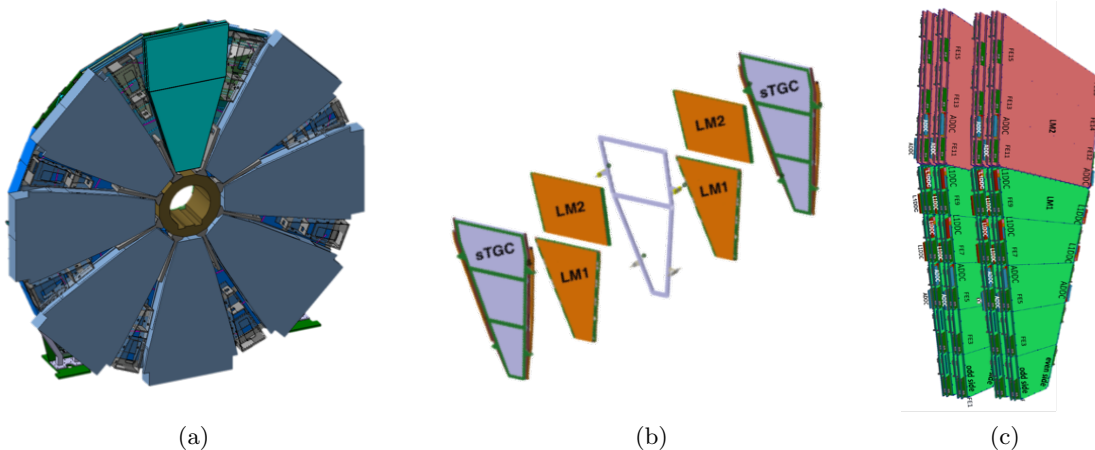


Figure 4.15: (a) Layout of the ATLAS NSW, consisting of 8 large wedge sectors and 8 small wedge sectors. (b) Arrangement of one NSW sector, consisting two quadruplets of sTGC and two quadruplets of MicroMegas and one supporting frame, where each MicroMegas wedge is divided into two modules. (c) Structure of one MicroMegas wedge, consisting two quadruplets. The module 1 is in green and module 2 is in red.

plane. The electric field allows the primary electron to gain energy and to be capable of producing further ionization, which results in an avalanche. This electron multiplication is collected by the readout strips as the final signal. To withstand the harsh radiation environment of the ATLAS NSW, carbon resistive strips with large resistivity are introduced.

4.3.2 Simulation for NSW MicroMegas particle hit rate

This analysis is to understand the expected background rates based on MicroMegas NSW simulations [118]. Following MC samples are used, normalized to the product of the corresponding cross sections:

- Ordinary pp events at center of mass energy $\sqrt{s} = 14$ TeV.
- Neutral particles with $|\eta| > 6$ from the ordinary pp events, simulated with long time scales. These are the non-collision processes with the long-lived neutral particles bouncing in the detector cavern, known as “cavern background”.
- Single muon events with energy of 50 GeV, as the “signal” events.

Figure 4.17 shows the expected particle hit rate with the LHC Run3 where the instantaneous luminosity $L_{\text{inst.}} = 7.0 \times 10^{34}$ Hz/cm², as a function of the radius of NSW. The highest rate is as high as 23 kHz/cm², which is at the inner most region (with large value of $|\eta|$). The major background is from fake muons, as

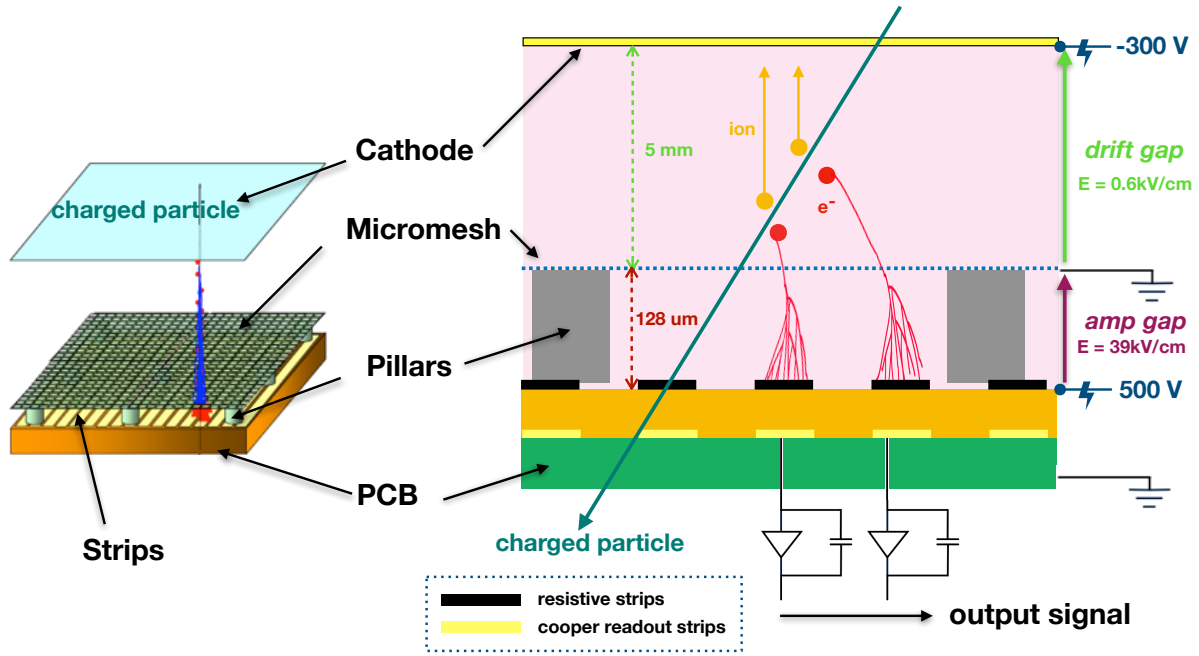


Figure 4.16: A schematic of the configuration of a [MicroMegas](#) detector gas gap, not drawn to scale. The right figure shows the side view of the left one. The gas gap region (pink in the right figure) is divided into two regions by the micro-mesh: the drift gap with a thickness of 5 mm and the amplification (amp) gap with a thickness of 128 μm. The micro-mesh is held in place by non-conductive pillars. The additional layer of resistive strips are placed above the readout ones. The cathode and the resistive strips are connected to a voltage of -300 V and 500 V respectively, while the micro-mesh is grounded. Therefore electric fields of 0.6 kV/cm and 39 kV/cm are produced in the drift gap and amplification gap respectively. A charged particle ionizes the gas and the electron drifts towards the micro-mesh. The increased electric field of the amplification gap causes the primary electrons to generate an avalanche and the final charge is collected by the readout strips.

well as from cavern background. These rates will be effectively reduced by the upgraded trigger system for NSW, to be as low as 0.4 kHz/cm^2 and maintains a high efficiency of muon selections [119].

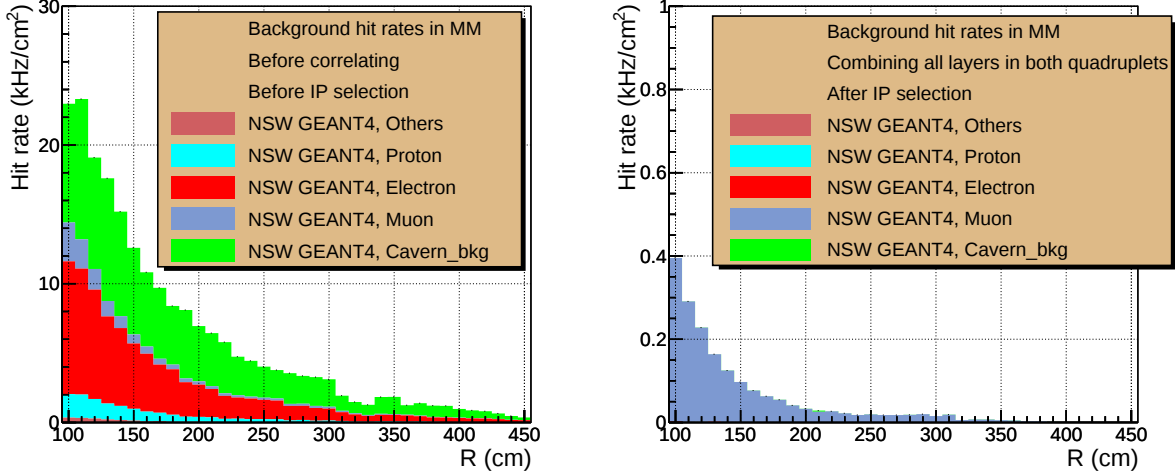


Figure 4.17: The prediction of background particle rate in ATLAS NSW MicroMegas with LHC run3 scenario [119], as a function of the distance to the center of the beam pipe R . In the left plot, no decision of NSW trigger is made, either from the correlation in the multi-layers of MicroMegas, or the pointing direction to the ip, where in the right plot those information are used to reject most of the fakes muons, with a high purity of muons as 99%.

4.3.3 NSW MicroMegas sensitivity to neutral particles

As seen in Section 4.3.2, there is a large proportion of non-collision neutral particles contributing to the background in NSW MicroMegas. These particles are mainly low-energy photons and neutrons, which are long-lived particles that bounce around in the detector cavern and enter the end-cap region. In order to ensure MicroMegas detector is capable of withstanding the radiation environment, it is important to understand how the MicroMegas detector will respond to these neutral particles.

In order to better simulate the response of MicroMegas with respect to neutral particles, the following additional details are being included in the GEANT4 simulation procedure:

- The full material information of MicroMegas chamber, including all the density and elemental components obtained from the fabricators.
- The detailed micro-mesh structure, which is not included in standard ATLAS MicroMegas NSW simulations [118].
- SHIELDING physics list [120], which gives a good representation of low energy neutron physics. This

Attenuation nominal	γ flux $m^{-2}s^{-1}$	received flux $m^{-2}s^{-1}$	Sensitivity	
			GIF++ measured	GEANT4 simulated
10	4.96×10^6	1.80×10^4	3.62×10^{-3}	3.44×10^{-3}
100	6.00×10^5	2.50×10^3	4.16×10^{-3}	3.88×10^{-3}

Table 4.3: Comparison of the integrated [MicroMegas](#) sensitivity to photons, between the GIF++ measurements [122] and GEANT4 simulation [121].

physics list is not included in standard ATLAS [NSW](#) simulations [118] due to the high cost of processing time and memory.

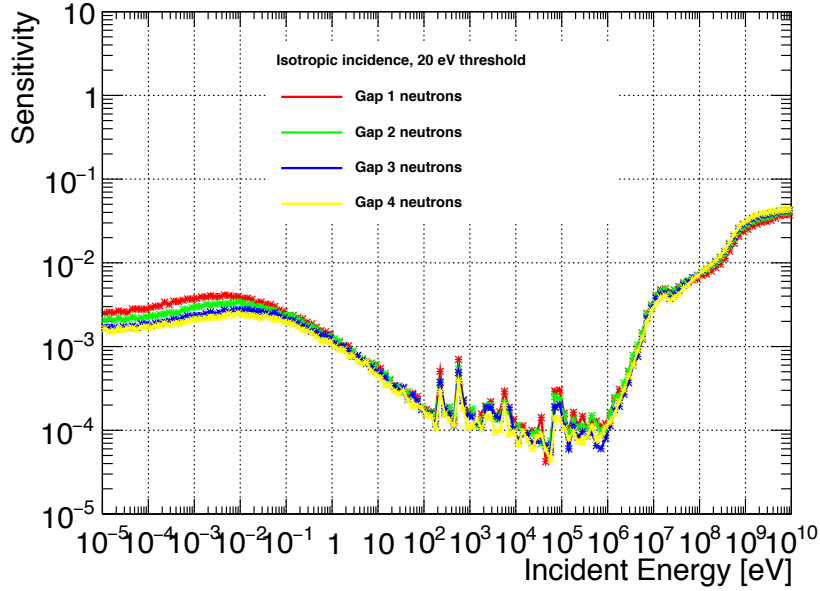


Figure 4.18: The sensitivities of the four gas gap in the same [MicroMegas](#) quadruplet with isotropic incident direction for neutrons [121].

The response of [MicroMegas](#) detector is represented by sensitivities, with is the ratio of particles that generate signal above the digital threshold 20 eV, with respect to the total number of particles initiated from the particle gun. Figure 4.18 and Figure 4.19 show the [MicroMegas](#) sensitivity to photons and neutrons respectively, as a function of the incident energy, in different gas gaps within the same [MicroMegas](#) quadruplet. In Figure 4.18, the spikes are due to the resonant neutron capture cross sections of some atoms, which exceeded the elastic scattering cross sections.

The simulated photon sensitivities are convoluted with the incident spectra with attenuation factor of 10 and 100 respectively. This results is compared to [CERN](#) GIF++ experiment [122], with the relative disagreement of 5.0% and 18.8% respectively, as seen in Table 4.3

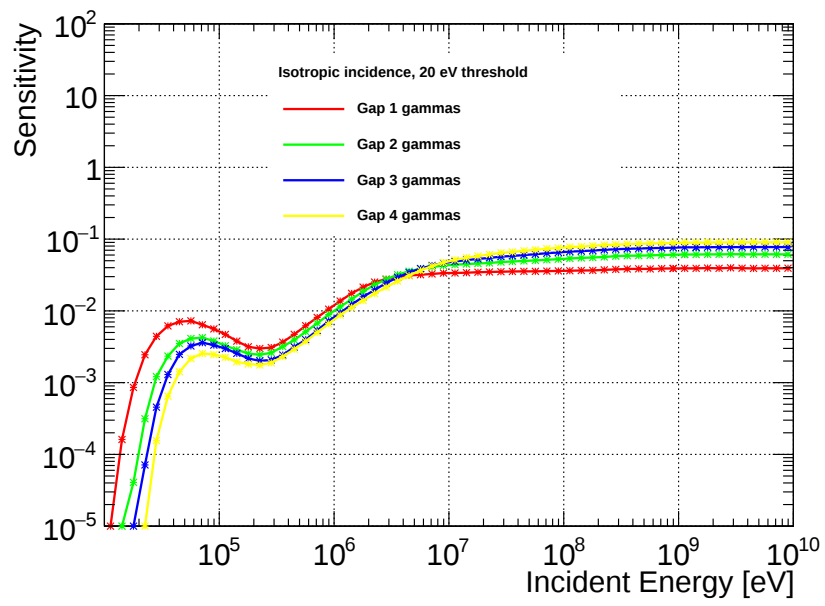


Figure 4.19: The sensitivities of the four gas gap in the same [MicroMegas](#) quadruplet with isotropic incident direction for photons [121].

Chapter 5

Physics objects

In particle physics experiments at the [LHC](#), one of the main goal is to reconstruct and measure the outgoing particles produced in [proton-proton](#) collisions. After an event is accepted by the trigger system, the physics objects of interests such as electrons, muons and jets are reconstructed from the detector digital signals, with the combination of different sub-parts. [Figure 5.1](#) shows a slice view of the detector reconstructing various of objects with different components of the detector.

This chapter will describe the detailed method of standard ATLAS algorithm for physics object reconstructions and identifications used in this thesis: electrons ([Section 5.1](#)), muons ([Section 5.2](#)), jets ([Section 5.3](#)) including τ jets ([Section 5.3.1](#)) and b-jets ([Section 5.3.2](#)), and missing transverse momentum ([Section 5.4](#)).

5.1 Electrons

Electrons are mainly reconstructed based on the combination of the electromagnetic clusters in the [LAr electromagnetic calorimeter](#) and the tracks in [inner detector](#). The [EMCal](#) is divided into a grid of 3×5 towers with a size of $\Delta\eta \times \Delta\phi = 0.025 \times 0.025$. The algorithm scans over all towers and finds a local maxima which is used to seed cluster reconstruction. The clusters are associated to the tracks in the [ID](#) with the primary interaction vertex. The properties of the matched tracks are further used to identify the cluster as being consistent with a prompt electron from hard scattering, or an electron from photon conversion [[123](#)]. The photon conversion usually takes place in the detector material. One of the important properties to distinguish the electron from conversions is the [impact parameter \(IP\)](#) d_0/σ_{d_0} , which is the distance of closest approach of the track to the measurement plane, and $z_0 \sin \theta$ where z_0 is the distance along the beam line between the point of the d_0 measurement and the beam spot position, and θ is the polar angle of the track. Furthermore, the [inner b-layer \(IBL\)](#) provides the number of hits in the innermost pixel layer to reduce the amount of fake electrons from photon conversions.

The identification of electrons is provided as electron qualities. The identification algorithm used in run2

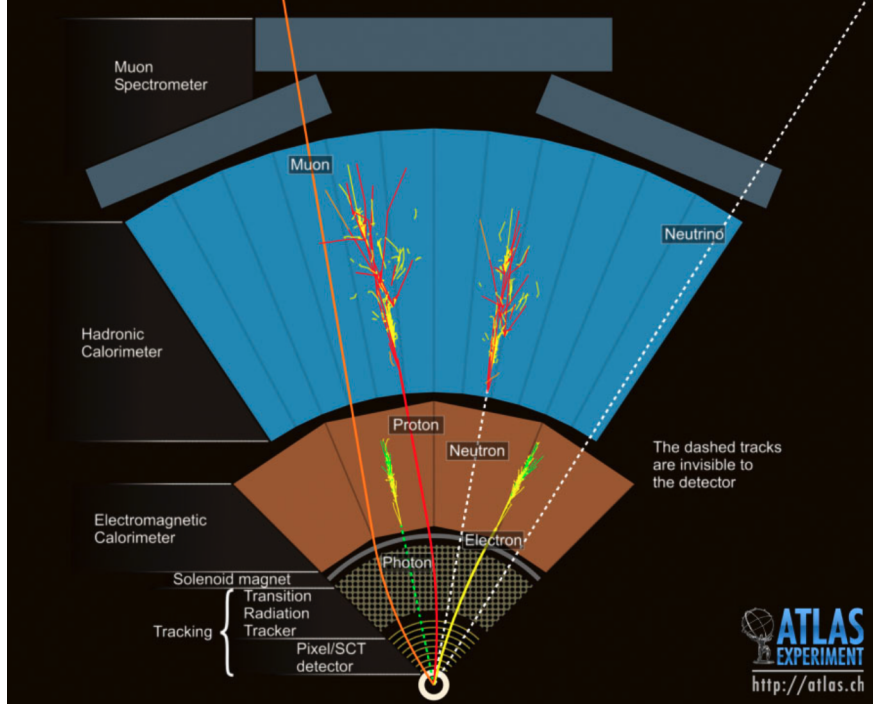


Figure 5.1: A slice of the ATLAS detector depicting the interaction of various of particles with different components of the detector [93].

analysis is a **likelihood-based (LH)** method, which uses a **multivariate analysis (MVA)** that uses signal and background **probability distribution functions** of discriminating variables of electron candidates [124]. Three levels of identification working points are provided: **Loose**, **Medium** and **Tight**, whose efficiencies are shown in Figure 5.2.

In addition to the identification criteria, **isolation (isol)** criteria are used to distinguish between the electrons produced in the hard-scattering and the ones from heavy flavor decays or misidentified jets. The **isolation** variables are calculated using the following variables:

- $E_T^{\text{topocone20}}$, which is a calo-based variable defined as the scalar sum of the transverse energy of topological clusters within a cone size of $\Delta R = 0.2$ around the electron candidate, and the contribution of E_T of the electron candidate is subtracted.
- $p_T^{\text{varcone20}}$, which is a track-based variable defined as the scalar sum of the transverse momentum of tracks within a cone size of $\Delta R = 0.2$, apart from the electron candidate track itself. The tracks must be associated with the same primary vertex of the candidate, known as **track-to-vertex association (TTVA)** tracks.

In the work of this thesis, the electron **isolation** operating working points are defined by fixed upper cuts

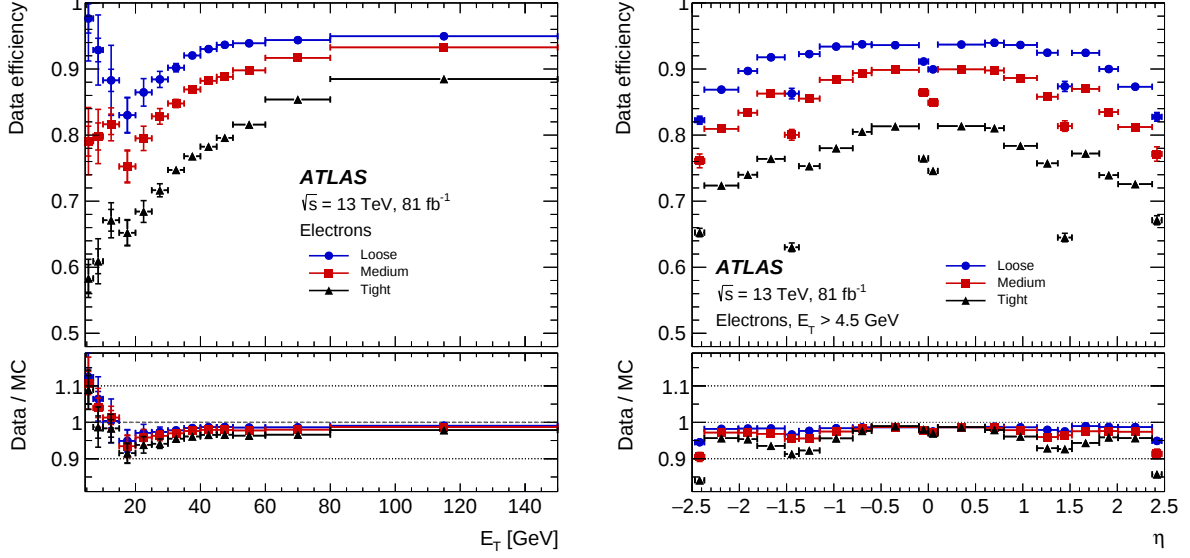


Figure 5.2: The electron identification efficiency in $Z \rightarrow ee$ events in data as a function of transverse energy E_T (left) and as a function of η (right) for the Loose, Medium and Tight operating points [124]. The efficiencies are obtained by applying data-to-simulation efficiency ratios measured in $J/\Psi \rightarrow ee$ and $Z \rightarrow ee$ events to $Z \rightarrow ee$ simulation. For both plots, the bottom panel shows the data-to-simulation ratios.

on these quantities with respect to the E_T or p_T of the electron candidate. Table 5.1 summarizes the cuts of these **isol** operation points, and Figure 5.3 shows the efficiency of these working point for electrons from inclusive $Z \rightarrow ee$ events.

Working point	Calorimeter isolation	Track isolation
Gradient	$\epsilon = 0.1143 \times p_T + 92.14\%$ (with E_T^{cone20})	$\epsilon = 0.1143 \times p_T + 92.14\%$ (with $p_T^{\text{varcone20}}$)
FCHighPtCaloOnly	$E_T^{\text{cone20}} < \max(0.015 \times p_T, 3.5 \text{ GeV})$	-
FCLoose	$E_T^{\text{cone20}}/p_T < 0.20$	$p_T^{\text{varcone20}}/p_T < 0.15$
FCTight	$E_T^{\text{cone20}}/p_T < 0.06$	$p_T^{\text{varcone20}}/p_T < 0.06$

Table 5.1: Definition of the electron **isolation** working points and **isolation** efficiency ϵ [124]. In the **Gradient** working point definition, the unit of p_T is GeV. All working points use a cone size of $\Delta R = 0.2$ for calorimeter **isolation** and $\Delta R_{\text{max}} = 0.2$ for track **isolation**.

5.2 Muons

Muons are one of the cleanest objects to reconstruct and identify in ATLAS detector. As muons travels through the entire detector, the tracks are reconstructed by a combination of the **inner detector**, the calorimeters and the **muon spectrometer**. In ATLAS, there are four different types of muons defined based on which sub-detectors are used:

- **Combined muon (CB)**: tracks are first reconstructed independently in the **ID** and the **MS**. Then

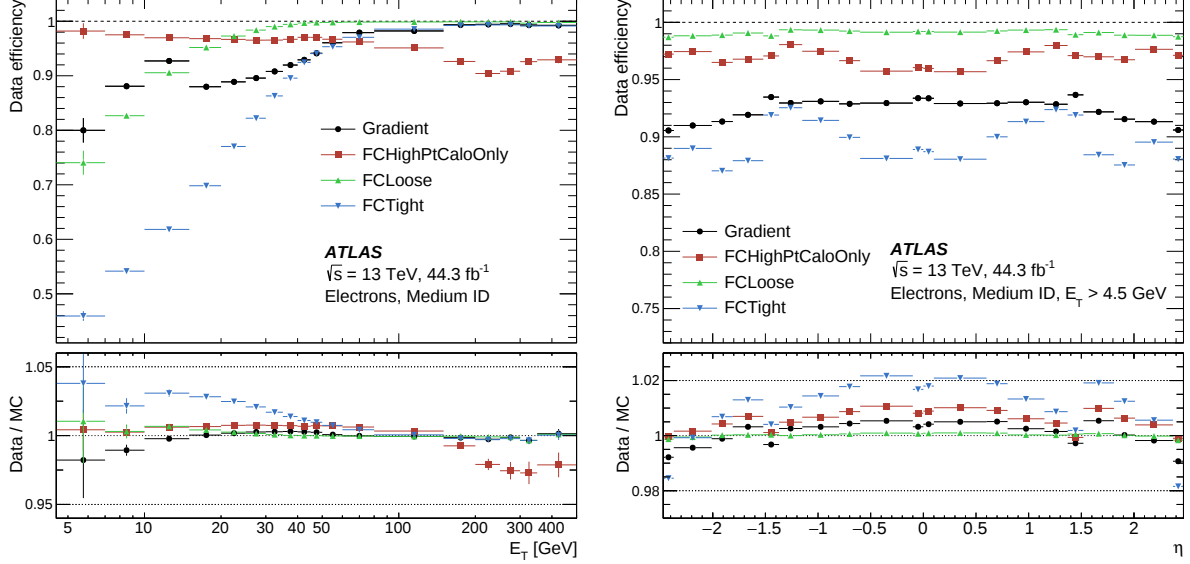


Figure 5.3: Efficiency of the different [isolation](#) working points for electrons from inclusive $Z \rightarrow ee$ events as a function of the electron E_T (left), electron η (right) [124]. The electrons are required to fulfill the [Medium](#) selection from the [likelihood-based](#) electron identification. The lower panel shows the ratio of the efficiencies measured in data and in [MC](#) simulations. The total uncertainty is shown, including the statistical and systematic components.

the tracks are combined by a global fit. Most muon track candidates follow an outside-in pattern recognition (namely the muon is first reconstructed in [MS](#) and extrapolate to match tracks in [ID](#)), but the inside-out pattern is also used for complementarity. This is the dominant method of muon reconstruction in ATLAS.

- **Segment-taged muon (ST)**: a track in the [ID](#) is classified as a muon if it is associated to a track segment in the [MDT](#) or [CSC](#). These muons typically have low p_T and cannot travel through the entire [MS](#).
- **Calorimeter-taged muon (CT)**: a track in the [ID](#) is classified as a muon if it is associated with an energy deposit in the calorimeter system that is expected from a minimum ionizing particle. There are the muons which ionize in the calorimeter system and do not travel to the [MS](#). These muons usually have lower purity.
- **Extrapolated muon (ME)**: the muon trajectory is reconstructed based only on the [MS](#) track in at least two layers. The ones which are not pointing to [interaction point](#) are rejected. These muons are the ones outside the coverage of the [ID](#) with $2.5 < |\eta| < 2.7$.

To reject prompt muons from pion and kaon decays, different qualities of muon identification are defined, mainly based on the tracks from each portion of the sub-detectors:

- **Loose:** all kinds of reconstructed muons are used, except for **CT** and **ST** muons which are restricted to $|\eta| < 0.1$. This working point is designed to maximize the identification efficiency.
- **Medium:** the default selection for muons in ATLAS, which minimizes the systematic uncertainties associated with muon reconstruction and calibration. Only **ME** and **CB** tracks are used in this selection, which are required to have at least 3 hits on at least two layers of **MDT** except the $|\eta| < 0.1$ region where tracks with at least one precision layer but no more than one precision hole layer are allowed. Moreover, a loose selection on the compatibility between the **ID** and **MS** momentum measurements is applied, in order to suppress the contamination due to hadrons mid-identified as muons.
- **Tight:** the working point with maximized purity of muons. Only **CB** muons with at least 2 stations and satisfying the **Medium** selection criteria are considered.
- **VeryLoose** a collection of all reconstructed muons fulfilling the hit requirements on the associated **ID** track.

Figure 5.5 shows the efficiency of the different identification working points using a **MC** $t\bar{t}$ sample as a function of the muon p_T .

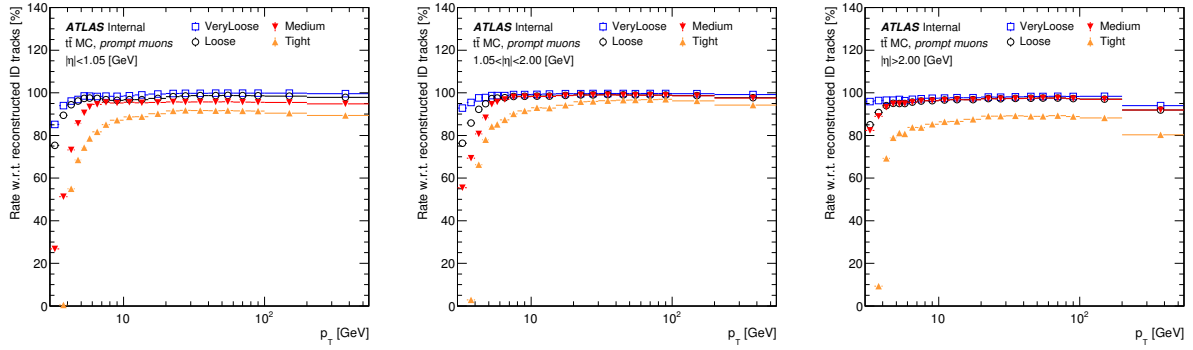


Figure 5.4: Muon identification efficiency as function of p_T of reconstructed muons with respect to inner detector tracks for prompt muons from $t\bar{t}$ **MC** for muons with different $|\eta|$ [125].

As the case of electrons, **isol** requirements are also applied for muons to reject the muons from heavy flavor hadron decays. The variables used for the **isol** working points are similarly to the ones for the electrons defined in Section 5.1, except using a $p_{T,TTVA}^{\text{varcone}}$ variable with the size of the cone of $\Delta R = 0.3$. Table 5.2 summarizes the definitions of the muon **isol** working points, and the efficiencies are shown in Figure 5.5.

Working point	track isolation	calorimeter isolation
FCTightTrackOnly	$p_{T,TTVA}^{\text{varcone30}}/p_T < 0.06$	-
FCLoose	$p_{T,TTVA}^{\text{varcone30}}/p_T < 0.15$	$E_T^{\text{topocone20}}/p_T < 0.30$
FCTight	$p_{T,TTVA}^{\text{varcone30}}/p_T < 0.04$	$E_T^{\text{topocone20}}/p_T < 0.15$
FixedCutHighPtTrackOnly	$p_T^{\text{varcone20}} < 1.25 \text{ GeV}$	-
FCTightTrackOnly_FixedRad	$\max(p_{T,TTVA}^{\text{varcone30}}, p_{T,TTVA}^{\text{varcone20}})/p_T < 0.06$	-
FCLoose_FixedRad	$\max(p_{T,TTVA}^{\text{varcone30}}, p_{T,TTVA}^{\text{varcone20}})/p_T < 0.15$	$E_T^{\text{topocone20}}/p_T < 0.30$
FCTight_FixedRad	$\max(p_{T,TTVA}^{\text{varcone30}}, p_{T,TTVA}^{\text{varcone20}})/p_T < 0.04$	$E_T^{\text{topocone20}}/p_T < 0.15$

Table 5.2: Muon **isolation** working points definitions [126].

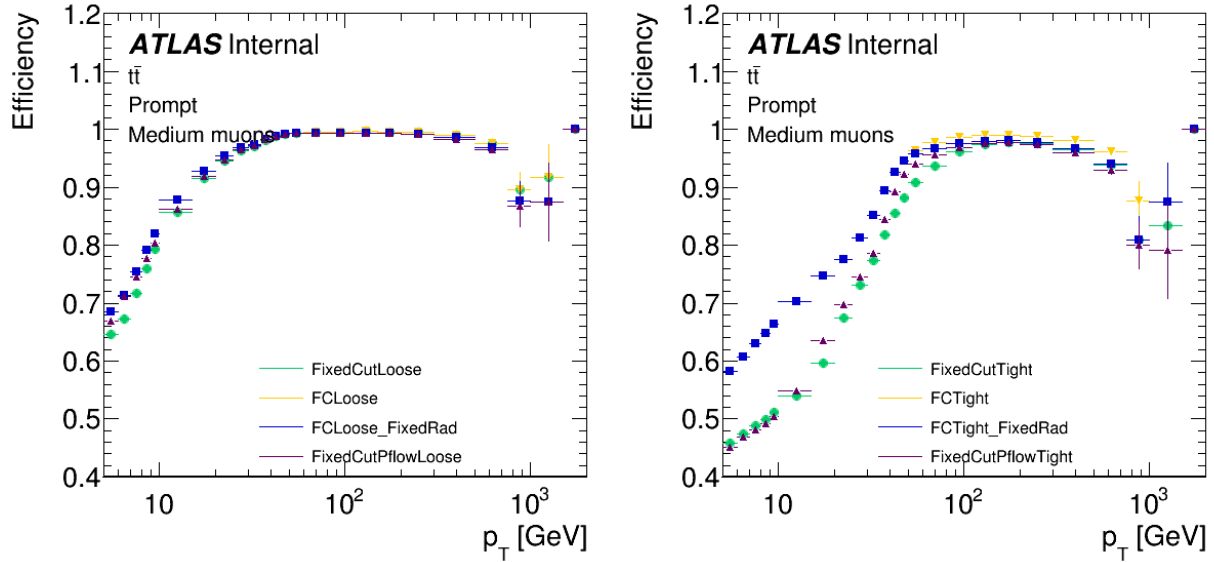


Figure 5.5: Efficiency of **Loose**(left) and **Tight**(right) **isolation** working points as a function of p_T for prompt muons from $t\bar{t}$ MC, with respect to muons passing **Medium** identification working point [126].

5.3 Jets

In particle physics, one big challenge is to measure the hadronic final state and reconstruct quark and gluons. As described in Section 2.3, colored particles cannot be directly observed in the detector, instead, a collection of quarks and gluons will be smeared in the showering process. The solution is to build these objects as a jet, which is a collimated spray of particle showers originating fragmentation and hadronization of quarks and gluons with energy depositions in the **HCal** and the **EMCal**. In the work of this thesis, jets are reconstruct using the anti- k_T algorithm [127]. This section introduces the reconstruction, calibration and identification of two different types of jets used in this thesis: τ jets in Section 5.3.1 and b -jet in Section 5.3.2.

5.3.1 τ jets

With a mass of 1.777 GeV and a proper decay length of 87 μm , τ leptons decay either leptonically ($\tau \rightarrow \ell \nu_\ell \nu_\tau, \ell = e, \mu$) with $\text{Br} \approx 35\%$ or hadronically ($\tau \rightarrow \text{hadrons} + \nu_\tau$) with $\text{Br} \approx 65\%$. These decays usually happen before the τ lepton reaches the active regions of ATLAS detector. Therefore, leptonically-decaying τ s are usually just considered the same as electrons or muons. In this section, only hadronic τ decays are considered. The hadronic decay products contain one or three charged π in 72% and 22% of all cases respectively.

τ candidates are seeded by jets, which are reconstructed using the anti- k_T algorithm [127] with a distance parameter of $R = 0.4$. Three-dimensional clusters of **hadronic calorimeter** cells are calibrated, known as **topoclusters**. Jets seeding τ candidates are required to have $p_T > 10$ GeV and $|\eta| < 2.5$, where the ones within the calorimeter crack region $1.37 < |\eta| < 1.52$ are vetoed. A τ vertex is chosen as the candidate track vertex with the largest fraction of momentum from tracks associated with the jet within $\Delta R < 0.2$. All these tracks must have $p_T > 1$ GeV and an **interaction point** requirement of $|d_0| < 1$ mm and $|z_0 \sin \theta| < 1.5$ mm. The energy of the τ jet candidate is obtained through a dedicated **boosted regression tree** (BRT) algorithm [128].

The identification algorithm of τ jets is designed to reject backgrounds from quark- and gluon-initiated jets. The identification uses **boosted decision tree** (BDT) based methods, with the input variables summarized in Table 5.3. Several different working points are provided corresponding to different τ jet identification efficiency values, labeled as **Loose**, **Medium** and **Tight**. For each working point, requirements on the BDT score are determined as a function of $\tau_{\text{had-vis}} p_T$, in order to achieve a constant value for the combined reconstruction and identification signal efficiency, as shown in Figure 5.6. Signal efficiencies are respectively 0.6, 0.55 and 0.45 for the 1-prong **Loose**, **Medium** and **Tight** working points. The corresponding 3-prong working points have efficiencies of 0.5, 0.4 and 0.3 [129].

This BDT based identification algorithm is not sufficient to discriminate electrons reconstructed as τ jets. Therefore another **likelihood-based overlap removal** (OLR) algorithm is provided for the discrimination between hadronic τ jet and electrons, with a geometrical matching within $\Delta R < 0.4$ between the τ jet candidate and electrons with $p_T > 5$ GeV. The likelihood score is parameterized by η^{track} and $p_T(\tau)$. Reconstructed 1-prong τ jet candidate will be rejected if matched to an electron candidate of a large likelihood score. Figure 5.7 shows the discriminating power of this **LHOLR** method in terms of the **receiver operating characteristic** (ROC) curve in different bins of η^{track} . The nominal working point is chosen corresponding to a τ identification efficiency of 95% [129].

Variable	Offline	
	1-track	3-track
f_{cent}	•	•
$f_{\text{leadtrack}}^{-1}$	•	•
	•	•
$ S_{\text{leadtrack}} $	•	
$f_{\text{iso}}^{\text{track}}$	•	
ΔR_{Max}		•
$S_{\text{T}}^{\text{flight}}$		•
m_{track}		•
$f_{\text{EM}}^{\text{track-HAD}}$	•	•
$f_{\text{track}}^{\text{EM}}$	•	•
$m_{\text{EM+track}}$	•	•
$p_{\text{T}}^{\text{EM+track}}/p_{\text{T}}$	•	•

Table 5.3: Discriminating variables used as input to the tau identification algorithm at offline reconstruction and at trigger level, for 1-track and 3-track candidates [129].

5.3.2 b -jet tagging

b -tagged used in this thesis described as a veto to reduce the contribution of heavy flavor backgrounds such as $t\bar{t}$ for $H \rightarrow 2a \rightarrow 4\tau$ analysis, as in Chapter 7.

b -jets are reconstructed from three-dimensional topoclusters in [HCal](#) and [EMCal](#) using the anti- k_{T} algorithm [127] with radius parameter $R = 0.4$. Each topocluster is firstly calibrated to the electromagnetic scale response, then corrected by the [jet energy scale](#) (JES) derived from $\sqrt{s} = 14$ TeV data and simulations [130]. Further selections are applied to reject jets originated from pile-up interactions by the [jet vertex tagger](#) (JVT) [131], which is a multivariate algorithm using the tracks matched to the jet.

The b -hadrons are characterized by a longer lifetime $c\tau \approx 450 \mu\text{m}$, hence decay at a longer distance of 3 – 5 mm away from the primary vertex. Hence, the principle method to identify a b -jet is to reconstruct a secondary vertex. In ATLAS, the b -tagging algorithm is based on a [boosted decision tree](#) (BDT), called MV2c10, using the information of:

- [impact parameter](#) of [inner detector](#) tracks matched to the jet.
- existence of a displaced secondary vertex.
- reconstructed flight path of b - and c -hadrons inside the jet cone.

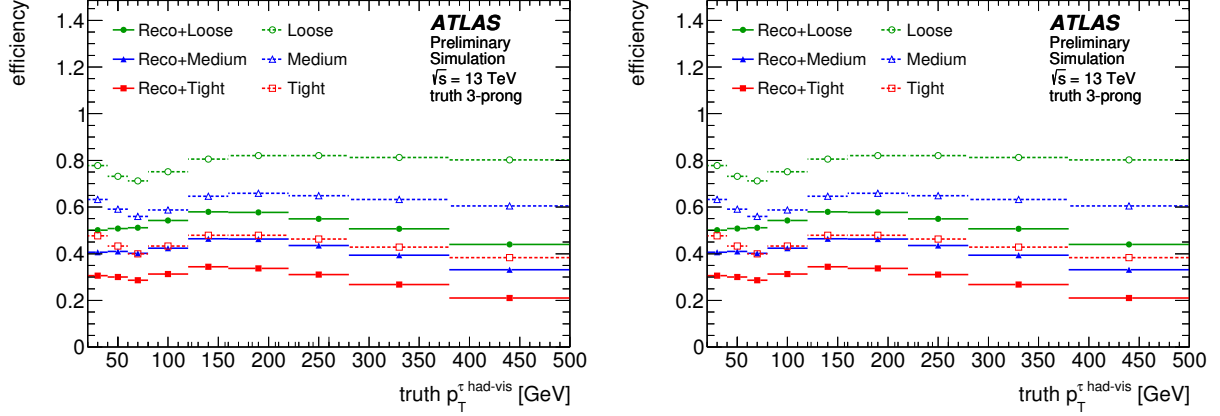


Figure 5.6: Efficiency for $\tau_{\text{had-vis}}$ identification (open symbols) and combined reconstruction and identification efficiency (full symbols) as a function of the average number of interactions per event, for 1-track (left) and 3-track (right) $\tau_{\text{had-vis}}$ candidates [129].

$p_T > 0.4 \text{ GeV}$
$ \eta < 2.5$ (ATLAS detector coverage for the ID tracking volume)
$\left \frac{d_0}{\sigma(d_0)} \right < 2$
$ z_0 \sin(\theta) < 3.0 \text{ mm}$

Table 5.4: Track selections for the TST [134]

5.4 Missing transverse energy

For objects such as SM neutrinos that do not interact with detector, cannot be reconstructed and identified directly. The commonly used object is the missing transverse energy E_T^{miss} , which is defined as [132]:

$$E_T^{\text{miss}} = - \left[\sum_{e, \gamma, \tau, \mu, \text{jet}} p_T + \sum_{\text{soft}} p_T(\text{track}) \right] \quad (5.1)$$

where the last “sort” term is **track-based soft term** (TST). The TST is defined as the sum of track p_T from tracks extrapolated from the primary vertex which are not within $\Delta R < 0.05$ of an electron or photon, $\Delta R < 0.2$ of a τ jet, or matched with a jet or a **combined muon**. Additional requirement for tracks used in the TST are summarized in Table 5.4

The performance of E_T^{miss} is shown in Figure 5.8, which is studied by $W \rightarrow e\nu$ analysis using data collected in 2015 of LHC run2 [133].

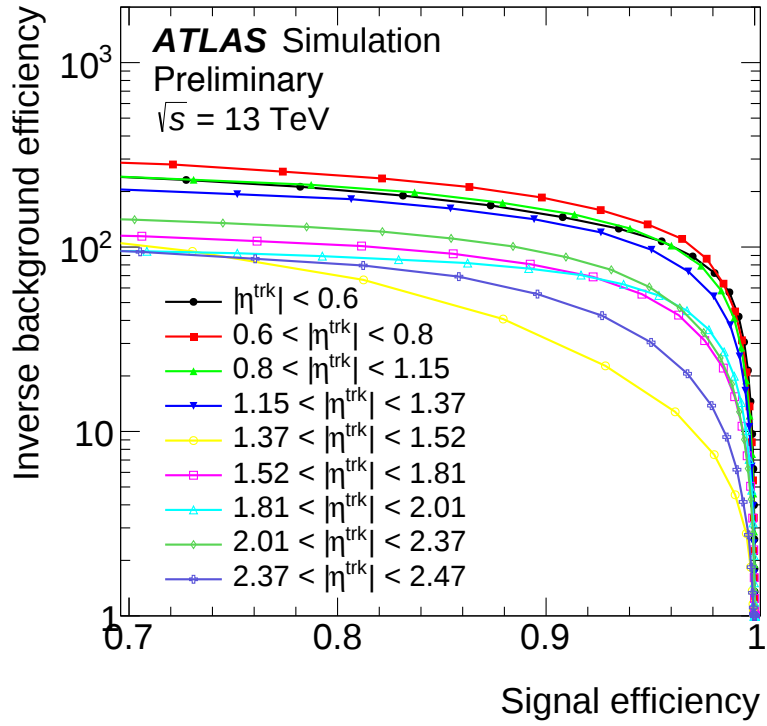


Figure 5.7: Inverse of electron mis-identification efficiency as a function of the tau identification efficiency (namely ROC curve) for different bins of η^{track} [129].

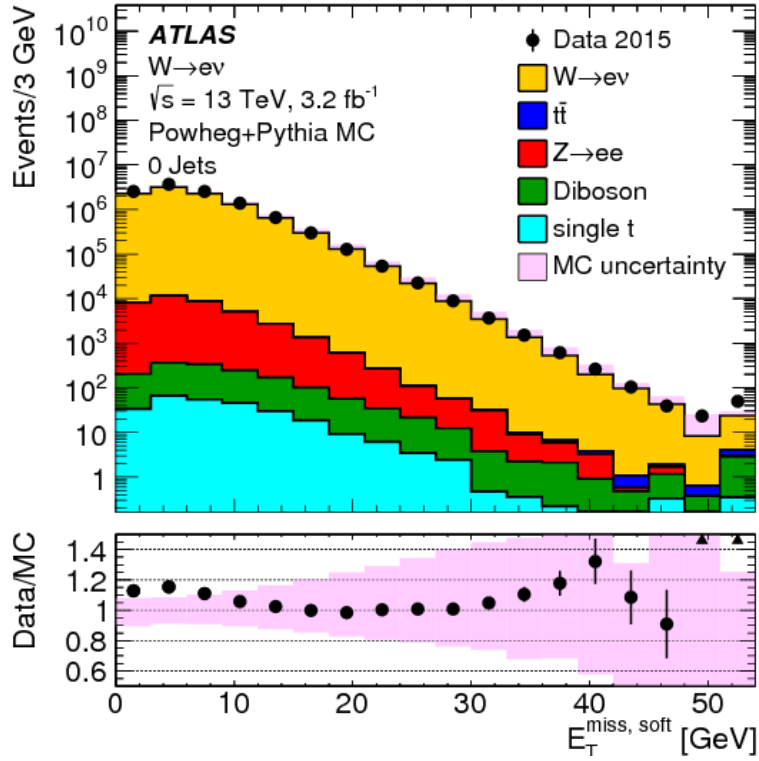


Figure 5.8: Distributions of E_T^{miss} with $N_{\text{jet}} = 0$ events with $W \rightarrow e\nu$ in data [133]. The expectation from MC simulation is superimposed and includes all relevant background final states passing the event selection. The shaded areas indicate the total uncertainty for MC simulations, including the overall statistical uncertainty combined with systematic uncertainties comprising contributions from the electron, jet, and the track-based soft term.

Chapter 6

Search for exotic Higgs decays to 4μ

This chapter presents a search for an exotic Higgs decays into a pair of light bosons (presumed to be pseudoscalars a) in a four-muon (μ) final state: $H \rightarrow 2a \rightarrow 4\mu$. In 2017, this analysis was published with 36.1 fb^{-1} of data collected by the ATLAS detector from 2015 to 2016 [135]. No significant excess of events above Standard Model background predictions is observed. This chapter discusses the work using 140 fb^{-1} of data collected by ATLAS experiments in full run2 of LHC from 2015 to 2018, with a stronger upper limit set at 95% [confidence level](#) on the [branching ratio](#) of exotic Higgs decays $\text{Br}(H \rightarrow aa \rightarrow 4\mu)$.

6.1 Signal models and features

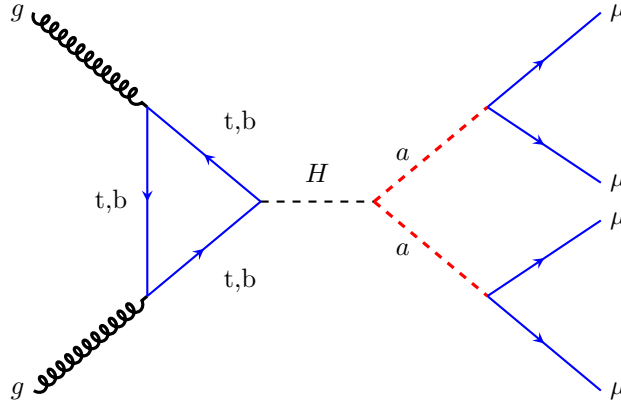


Figure 6.1: The Feynman diagram of the $H \rightarrow 2a \rightarrow 4\mu$ signal model, where the Higgs is produced through [ggF](#).

Signal samples of [ggF](#) Higgs boson production are generated with [POWHEG-BOX](#) [136, 137, 138] using the CT10 [PDF](#) set at [next-to-leading order \(NLO\)](#)[139]. The Higgs boson mass is assumed to be $m_H = 125 \text{ GeV}$. The [SM](#) Higgs is then replaced with a neutral scalar Higgs a from the [2HDM+S](#), and this Higgs is decayed to 4μ via two pseudoscalar bosons by the [PYTHIA8](#) generator [140], at the series of mass points of $m_a = 0.5, 1, 2, 2.5, 4, 6, 8, 10$ and 15 GeV . The a boson is performed in the narrow-width approximation and the coupling to the muons is assumed to be that of a pseudoscalar. Figure 6.1 shows the Feynman diagram

of the signal process.

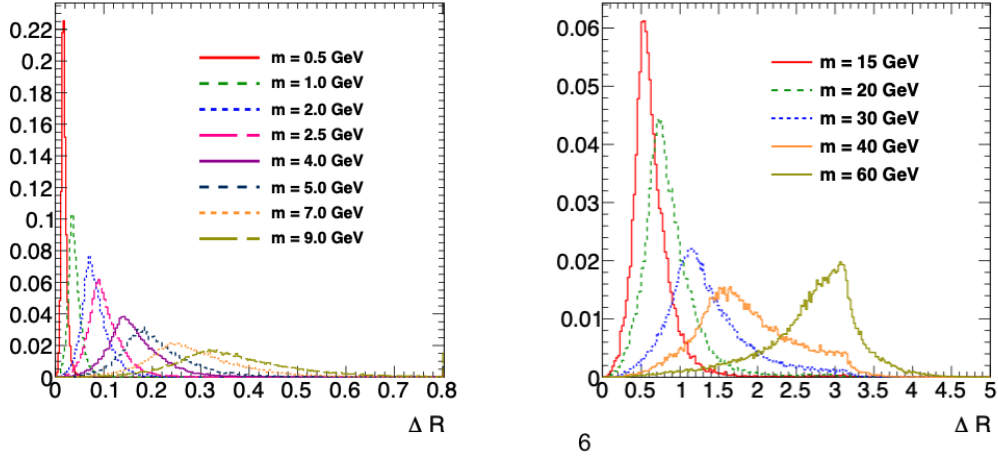


Figure 6.2: The separation between muons in $a \rightarrow \mu\mu$ decays for a mass range of $m_a < 10$ GeV (left) and $m_a > 15$ GeV (right). For the low masses, the muon pairs become very close together.

Figure 6.2 highlights a key feature of the search strategy for these signatures. The decay products are close by, particularly for the very low masses. The usual calculations of the muon *isolation* variables are spoiled by the track or energy deposition from the near-by muon. Hence, a strategy is used to correct the muon *isolation* where the affect from any close-by muon within the radius of *isolation* is subtracted. Figure 6.3 shows the improvement in this correction algorithm. The *isolation* requirement in this analysis all has this correction implemented.

6.2 Event selections

In this analysis, muons are required to pass a baseline selection of $p_T > 5$ GeV, or $p_T > 15$ GeV if the muon is calo-tagged, $|\eta| < 2.5$, and the **Loose** identification working point. The impact parameters have to satisfy $z_0 \sin\theta < 0.5$ mm and $d_0 < 1$ mm to reject background muons such as cosmic-ray muons. Muons can be combined, standalone, calo-tagged or silicon-associated forward. But events with less than three combined muons will be rejected.

Considering the fact that the final states are composed by four muons which decays from two same parents, the primary observables for this analysis are the invariant masses of the two muon pairs. Here we define m_{12} and m_{34} as the invariant masses of the two muons that make up a quadruplet, with the defining constraint that $m_{12} > m_{34}$. Thus m_{12} identifies the primary pair and m_{34} the secondary pair. Alternative pairings of *opposite-sign* (OS) muons can be formed. The invariant masses of these alternative pairings are denoted m_{14} and m_{23} . Explicitly, the positively charged muon from the primary pair is paired with the

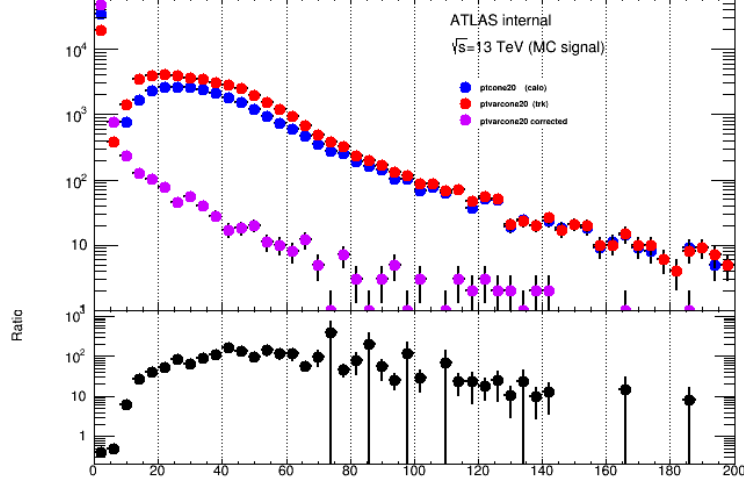


Figure 6.3: The p_T^{cone20} isolation variable distribution for muons from $H \rightarrow 2a \rightarrow 4\mu$, where $m_a = 5$ GeV. The different colors are representing: (blue) The scalar sum of transverse momentum in a cone of $\Delta R < 0.3$, (red) deviates from p_T^{cone20} above 50 GeV to optimize the performance of muons with higher momentum, (purple) $p_T^{\text{varcone20}}$ with a close-by correction. The ratio plot on the bottom is showing the comparison before and after the close-by correction.

negatively charged lepton from the secondary pair to create m_{14} , and the positively charged muon from the secondary pair is paired with the negatively charged muon from the primary pair to create m_{32} .

For events with more than four muons, only one quadruplet is selected through the quadruplet selection ranking process. The quadruplet with the smallest $\Delta m_{\mu\mu} = m_{12} - m_{34}$ is selected, considering the parents should have the same masses ideally.

After a quadruplet has been selected from the event, further requirements on the muons of the quadruplet are applied, including the isolation, muon identification quality and impact parameter significance requirement:

Isolation cut: the FCLoose_FixedRad isol working point is used for muons, which applies cuts on track-based and calo-based isolation variables p_T^{cone20} , p_T^{cone30} , and $E_T^{\text{topocone20}}$. Here the close-by correction is implemented as described in Section 6.1. Only the other three muons in the quadruplet are used for this correction.

Higgs boson mass window: $m_{4\mu}$ must be in the window $120 \text{ GeV} < m_{4\mu} < 130 \text{ GeV}$.

Di-muon mass window: $0.88 < m_{12,34} < 20 \text{ GeV}$. The lower bound comes from the limitation of the heavy flavor data-driven background estimation (see Section 6.3.2). The range in $m_{\mu\mu}$ considered is limited by the ability to reconstruct overlapping muons at low mass. At high mass, another search

extends from 20 GeV to the Z mass.

Quarkonia veto: J/Ψ and Υ will provide charm quarks and anti-quarks which can decay semi-leptonically to a muon pair. Therefore our event is rejected if either

$$(m_{J/\Psi} - 0.25 \text{ GeV}) < m_{12,34,14,23} < (m_{\Psi(2S)} + 0.30 \text{ GeV})$$

$$(m_{\Upsilon(1S)} - 0.70 \text{ GeV}) < m_{12,34,14,23} < (m_{\Upsilon(3S)} + 0.75 \text{ GeV})$$

Compatibility requirement: $\frac{m_{34}}{m_{12}} > 0.85$, which shapes a “wedge” on the m_{12} vs m_{34} plane. The distribution of $\frac{m_{34}}{m_{12}}$ of signal samples can be seen in Figure 6.4.

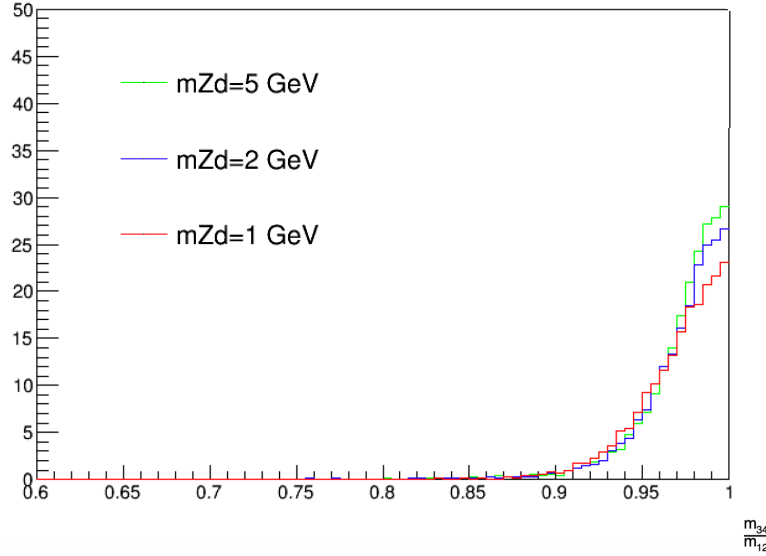


Figure 6.4: Scan of $\frac{m_{34}}{m_{12}}$ to define the requirement for the signal region for $m_a = 1, 2, 5 \text{ GeV}$.

6.3 Background estimation

A few [SM](#) processes can provide four muons in the final states, and are the dominant sources of background. There is also a contribution from multijet events, particularly decays of heavy-flavor jets that produce muons, such as events with two or four b -quarks in the final state. A leading part of the $b\bar{b}$ contribution comes from double semi-leptonic decays, where the b -quark decays to a muon and a c -quark which further decays into another muon and light hadrons. Resonances produced in the b -quark decay chain (i.e. $\omega, \rho, \phi, J/\Psi$) that result in pairs of muons contribute to the $b\bar{b}$ background. Events with $b\bar{b}b\bar{b}$ may satisfying the signal region requirements if each of the b -quarks decays semi-leptonically.

The prompt muon contribution is estimated from MC simulation, whereas the heavy-flavor background is estimated using a fully data-driven method.

6.3.1 Prompt background simulation

The following SM processes are considered:

- $ZZ^* \rightarrow 4\mu$: Non-resonant SM production of a pair of Z bosons, each decaying to a pair of muons.
- $H \rightarrow ZZ^* \rightarrow 4\mu$: Higgs boson production with subsequent decays to four leptons.
- **Triboson production**: Higher-order electroweak (EWK) processes with couplings of α^6 , which lead to four muons in the final state with two additional particles.

6.3.2 Heavy-flavor background estimation

A method to estimate the $b\bar{b}$ and $b\bar{b}b\bar{b}$ backgrounds has been developed using fully data-driven inputs and only relying on simulation to account for contributions from EWK backgrounds. The basic idea behind the method is to derive a shape for the background prediction for the four-muon process using the product of two dimuon spectra. This method works because the $b\bar{b}$ background is dominated by events where both b-quarks decay to pairs of muons ($2\mu + X$) via two semileptonic decays or resonances, e.g. ω , ρ , ϕ , J/Ψ , so the modeling of the pairs can be studied separately using a higher statistics sample of dimuon + single muon events, where the single muon represents a b with a single semi-leptonic decay $b \rightarrow \mu + X$.

The smaller contribution coming from $b\bar{b}b\bar{b}$ events where each b-quark decays semi-leptonically and produces a muon, is also accounted by this data-driven method. All the requirements on the muons, including isolation and d_0 significance (referred to as IP hereafter), are applied. One important challenge is that in this analysis a Higgs boson mass requirement $120 \text{ GeV} < m_{4\mu} < 130 \text{ GeV}$ is used, which introduces correlations between the dimuon pairs. These correlations associated to the event kinematics are studied in data using a sample enriched in $b\bar{b}$ events, where the isolation and IP requirements are inverted. This is possible because the kinematic constraints from the Higgs mass requirement and the muon identification requirements are uncorrelated.

Modeling of the 4-muon background shape

In the case of the $b\bar{b}$ background, the four muons are produced as two independent muon pairs. Therefore a high statistics sample comprised of pairs of muons is used to derive a shape for each pair of muons. A

small contribution to this shape comes from the cases where two muons originate in independent b -decays, and the shape of the $b\bar{b}b\bar{b}$ background is also accounted for.

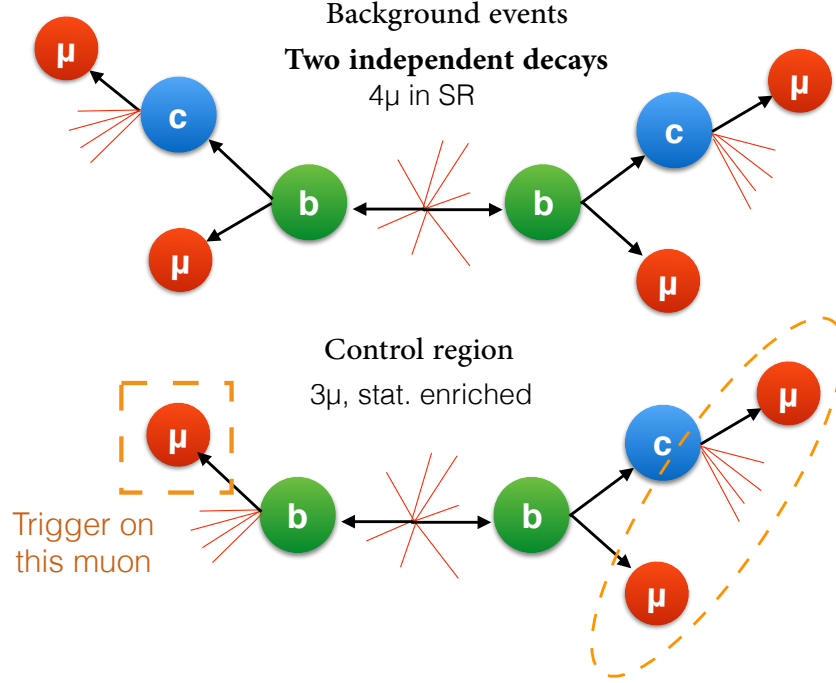


Figure 6.5: Sketches showing the basic idea of the data-driven method for heavy flavor background in $H \rightarrow 4\mu$ analysis. On the top is a $b\bar{b}$ processes that provide four muons and enters the SR. The bottom sketch demonstrates $b\bar{b}$ processes which contribute to three muons and enter the CR and enriched by statistics. In the bottom sketch, the single muon on the left hand side provides the muon trigger, and the invariant mass of the di-muon pair on the right hand side will be used to construct the 1-D template.

It should be emphasized that the same quality and identification requirements on the muons are used in the background estimation method as in the analysis. The background estimation procedure is based on constructing a $m_{\mu\mu}$ distribution (removing the quarkonia resonances) in a sample of three muons. Two selections are used in order to account for different muon momentum thresholds:

- **High- p_T :** Require 3 muons, a pair with thresholds $p_T \geq 20$ GeV and $p_T \geq 10$ GeV, and an additional muon with threshold $p_T \geq 5$ GeV. The muons in the pair are required to be matched to the dimuon trigger.
- **Low- p_T :** Require 3 muons, a pair with thresholds $p_T \geq 5$ GeV and $p_T \geq 5$ GeV, and an additional muon with threshold $p_T \geq 25$ GeV. The additional muon is required to be matched to the single muon trigger.

Figure 6.5 is a sketch demonstrating the overall idea of the data-driven method. It should be noted that in simulation, $> 97\%$ of signal events satisfying the analysis selection criteria are comprised of this combi-

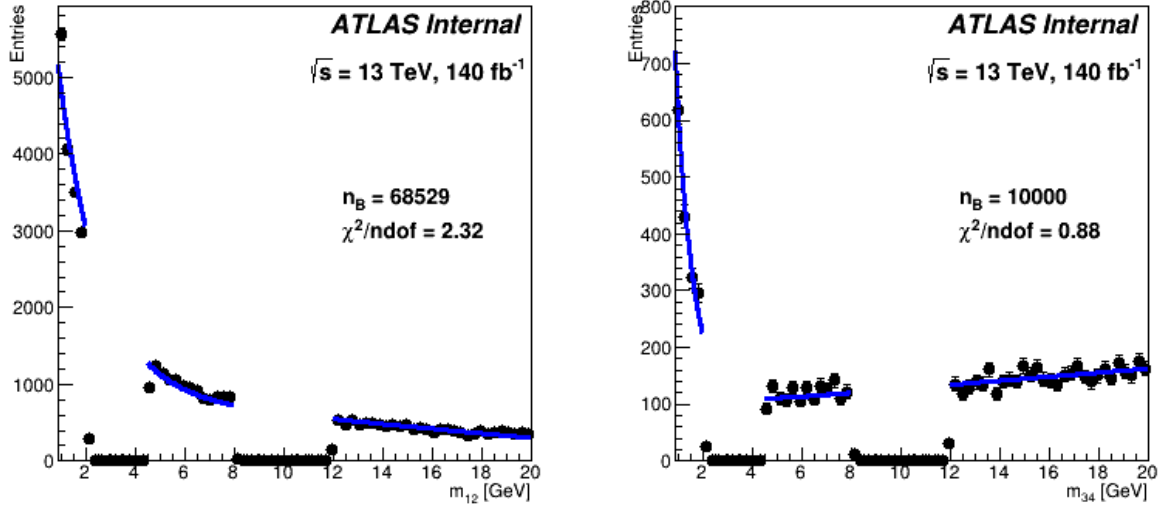


Figure 6.6: Dimuon mass distributions for the (left) high- p_T and (right) low- p_T pair configurations. The fits used to model the shapes of these dimuon pairs are also shown.

nation of pairs, high p_T and low p_T pair. Focusing on this combination provides an acceptable background determination within the associated uncertainties of the method. Figure 6.6 shows the distributions of muon pairs for the high p_T (m_{12}) and low p_T (m_{34}) configurations. These distributions are fit using an exponential function combined with a first degree polynomial, showing reasonable modeling of this background shape. The two dimuon templates are multiplied together, resulting in a 4-muon template, which is demonstrated in Figure 6.7. The results of the Cartesian product are shown in Figure 6.8.

Modeling of the kinematics from the Higgs boson mass constraint

In this analysis, the two muon pairs are not independent of each other, mainly because of the requirement that the mass $m_{4\mu}$ is compatible with the Higgs boson mass. In order to include the effect of this requirement, an efficiency map for the Higgs boson mass requirement is derived as a function of the two muon pairs m_{12} and m_{34} using a sample enriched in $b\bar{b}$ events. Events are selected with four muons satisfying the full analysis selection except for the [isolation](#) and [IP](#) requirements on the muons which are inverted. The distributions of m_{12} vs. m_{34} are fit before and after applying the Higgs boson mass constraint. The fit function used is a Gaussian convoluted with a polynomial of second degree, in two dimensions:

$$f(m_{12}, m_{34}) = (G(\mu_{x1}, \sigma_{x1}; m_{12}) + \text{Poly}(\text{par}_{x1}, \text{par}_{x2}; m_{12})) \times (G(\mu_{y1}, \sigma_{y1}; m_{34}) + \text{Poly}(\text{par}_{y1}, \text{par}_{y2}; m_{34}))$$

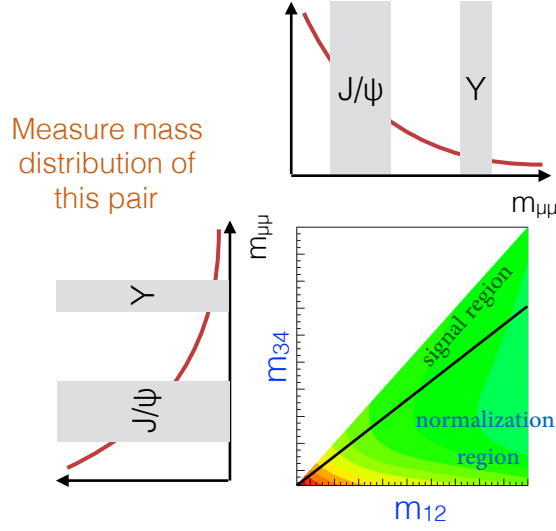


Figure 6.7: Sketch demonstrating the construction of the 2-D template, the two templates on the top and right are measured from the low- p_T and high- p_T templates. The region covered by the gray-colored boxed are the quark-onia veto. The final 2-D template is divided by the compatibility cut of $\frac{m_{12}}{m_{34}} < 0.85$, where the diagonal wedge region is the signal region, and the remaining triangle will be used for normalization.

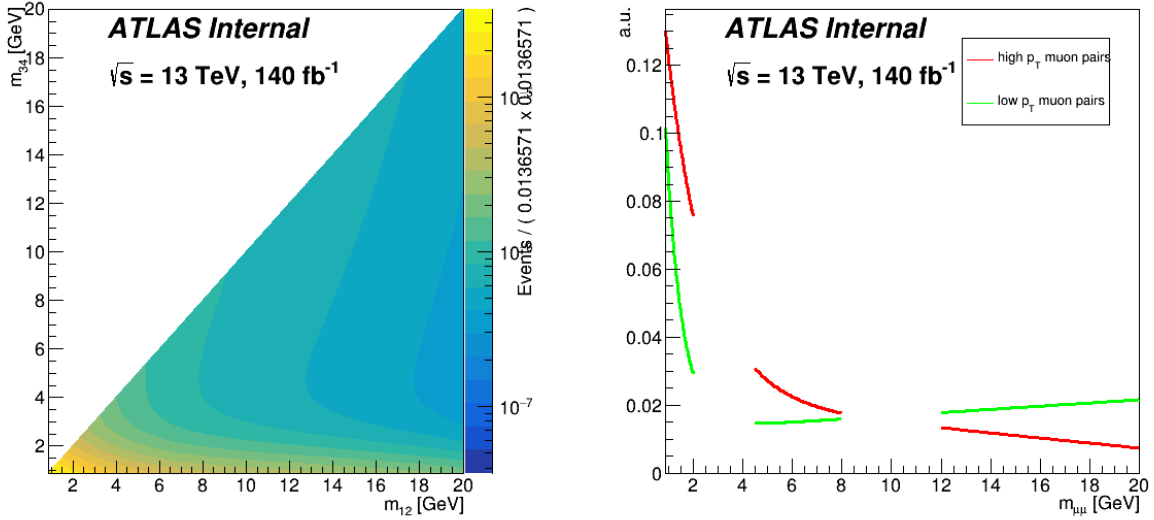


Figure 6.8: (left) The product of a high- p_T and a low- p_T dimuon template, resulting in a 4-muon template. (right) The projections of the predictions on the two axes from the 4-muon template are shown. Note that only the region where m_{12} is larger than m_{34} is used in this analysis.

The left and right columns of Figure 6.9 show the results of the fits before and after applying the Higgs boson mass constraints respectively. The resulting efficiency map, defined as the ratio of the distribution after and before applying the Higgs mass constraint ratio distribution is shown in the central plot of Figure 6.10. This efficiency map is applied to the shape template derived previously (left plot of Figure 6.8 or 6.10), to produce the final shape template for the four-muon $b\bar{b}$ background, shown in the right plot of the same figure.

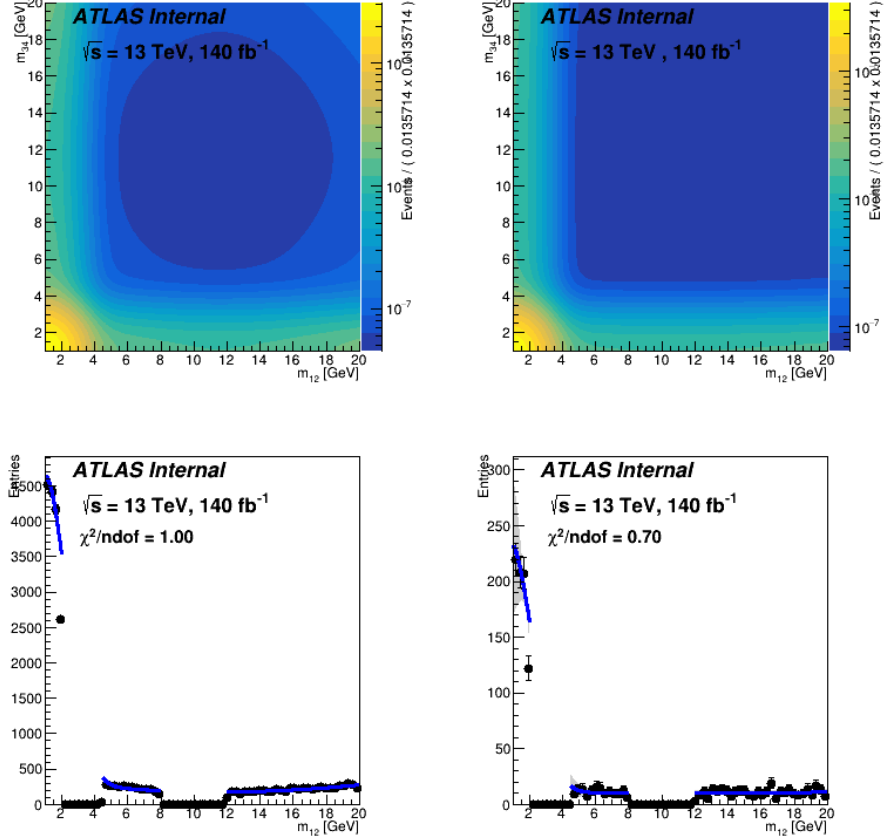


Figure 6.9: Before and after the Higgs mass constraint. The map of m_{12} vs. m_{34} derived from fitting events satisfying the full analysis selection criteria with the exception of the *isolation* and *impact parameter* requirements which are inverted. Note that only the region where m_{12} is larger than m_{34} is used in this analysis. The right plots are for the case before applying the Higgs boson mass constraint, and the left two plots are for the case that invariant mass of 4 muons inside the Higgs boson mass window.

Normalization of the background prediction

The final template shape (shown on the right of Figure 6.10) covers both the signal region (the region defined by the compatibility cut $|m_{12} - m_{34}| < 0.15m_{12}$, which we call *region A*) as well as the control region defined

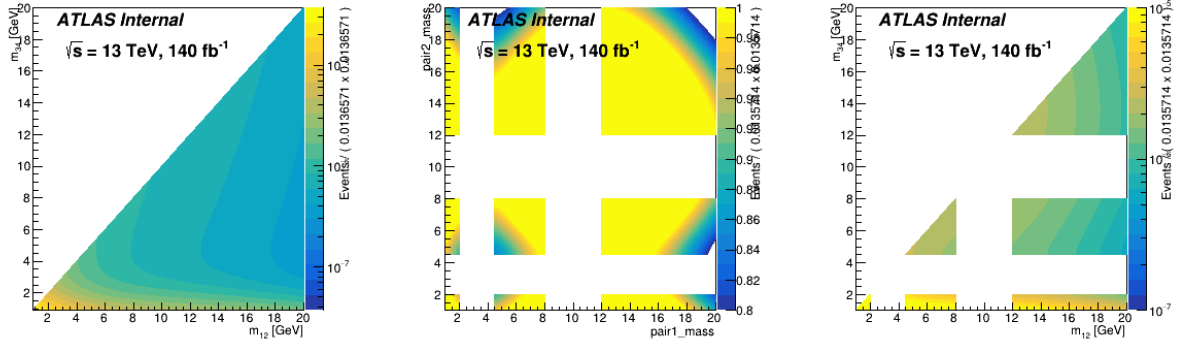


Figure 6.10: (left) Initial prediction for the shape template derived from the product of two dimuon pairs (see Fig. 6.8). (center) Efficiency map of m_{12} vs. m_{34} derived as the ratio of the distributions obtained after and before applying the Higgs boson mass requirement, in a regions with **isolation** and **impact parameter** requirements inverted. (right) Final prediction for the shape template derived from the product of two dimuon pairs and corrected for the impact of the Higgs boson requirement.

by inverting the compatibility cut (which we call *region B*). The template is to be normalized to data in this control region, and then the prediction in the signal region can be extracted from this normalized template. The ratio of region A to region B given by the template is 0.285.

Unfortunately, due to the limited data set size, there are no data events in region B. Therefore the yield in region B is estimated by counting events in another control region defined by inverting the Higgs mass window requirement¹ (region C) and applying a transfer factor to extrapolate into region B. The transfer factor is measured in two complementary regions (region D and E), which correspond to region B and C except for an inversion of the **impact parameter** and **isolation** requirements. These various regions are shown pictorially in Figure 6.11.

The event counts in the various regions are shown in Figure 6.12, and are also listed in Table 6.1.

Region	Definition	Events	Prediction
E	inversed isol IP AND inversed $120 < m_{4\mu} < 130$ GeV	4910	—
D	inversed isol IP AND $m_{34}/m_{12} < 0.85$	362	—
C	inversed $120 < m_{4\ell} < 130$ GeV	14	—
B	$m_{34}/m_{12} < 0.85$	0	$14 \frac{362}{4910} = 1.03$
A	Signal Region	0	$1.03 \times 0.404 = 0.417$

Table 6.1: Event counts and predictions in each region. When the Higgs boson mass window cut is inverted, the region $80 < m_{4\ell} < 100$ GeV is also excluded. The variable **isol IP** is the combination of **isolation** and **impact parameter** requirements. The 0.404 transfer factor from the integral of region B to region A comes from the template.

The final prediction in the signal region A is obtained by scaling the prediction in region B by the ratio

¹The range $80 < m_{4\ell} < 100$ GeV is also excluded, to limit contributions from single Z bosons with a Dalitz decay.

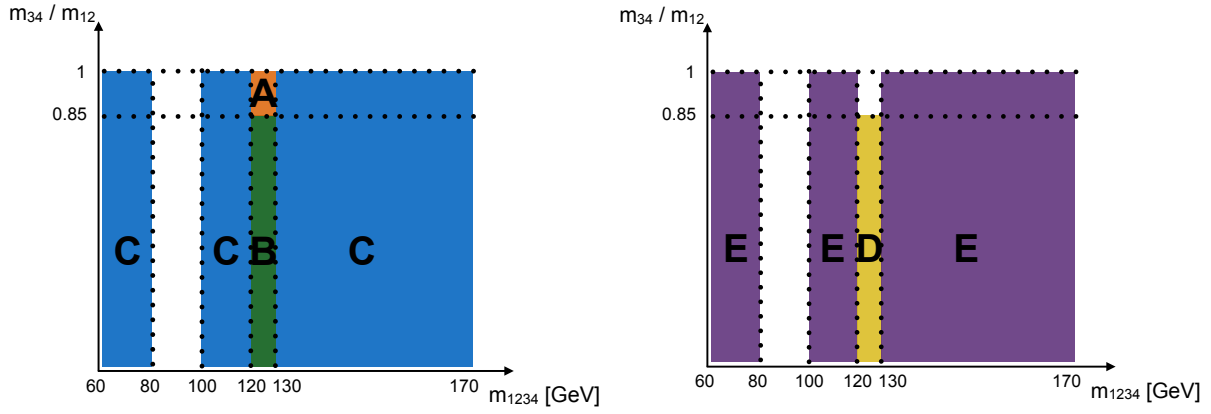


Figure 6.11: Definition of regions used in the ABCD background estimate. *isolation* and *impact parameter* requirements are inverted on the D and E regions, defined on the right plot.

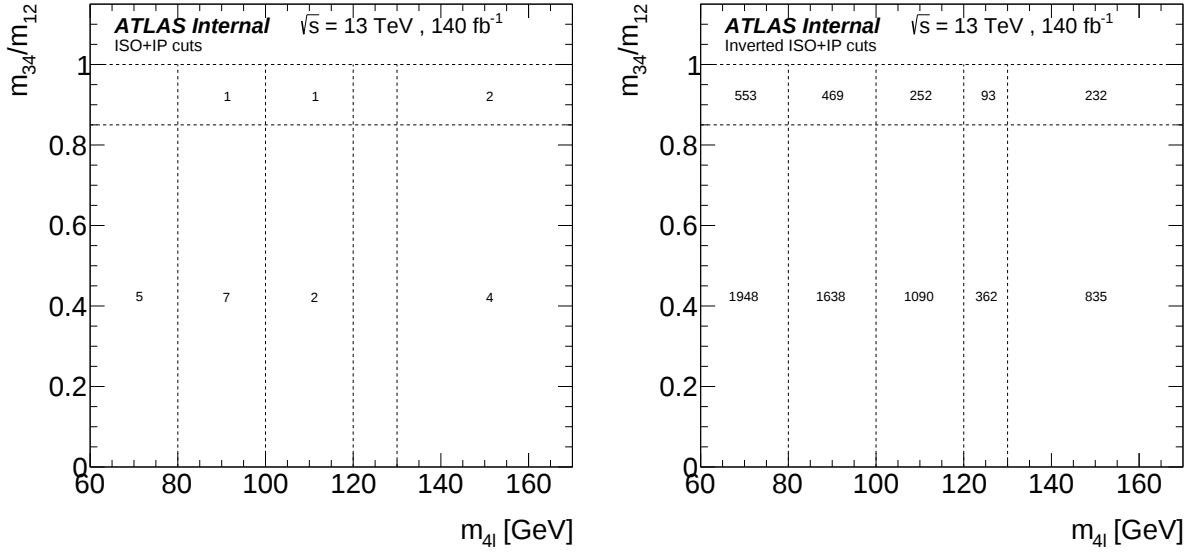


Figure 6.12: Inputs in the BCD regions used in the ABCD background estimate. Note that the signal region, satisfying the muon identification requirements on isolation and IP as well as the kinematic requirements $m_{34}/m_{12} > 0.85$ and $120 < m_{4\mu} < 130$ GeV, is blinded in the figure on the left.

between the areas corresponding to regions A and B in the template. This ratio is equal to 0.285. The final prediction for the total background is $N = 0.43 \pm 0.08$ and is shown in Figure 6.14 with yellow color. The derivation of the uncertainty on this background prediction is described next.

6.4 Uncertainties on the background estimation

In the low mass region, the four muon final states is very rare after vetoing the quarkonia in SM. Hence the analysis is mainly affected by a high relative statistical error, and the uncertainties from the data-driven method of the heavy-flavor background modeling. In this section, the estimation of different sources of uncertainties from the background is explained.

6.4.1 Experimental uncertainties

Scale factor (sf), which are parameterized as a function of object kinematics, quality and other characteristics are used to correct MC simulation with respect to data. Uncertainties associated with muon track-to-vertex association (TTVA), reconstruction and isolation (isol) efficiencies are considered. For $p_T < 10$ GeV, the muon reconstruction efficiency corrections are extrapolated. Their systematic variation of scale factor under different conditions and statistical constrains due to the number of events used to study the scale factor uncertainties are also considered. Table 6.2 shows the break down of experimental systematic for the prompt background categories.

SM process	Relative Uncertainty [%]		
	$H \rightarrow 4\ell$	$ZZ \rightarrow 4\ell$	$VVV \rightarrow 4\ell + 2X$
MC statistical uncertainty	± 2.2	± 16.9	± 0.20
Muon reconstruction stat.	± 0.26	± 0.32	± 0.16
Muon reconstruction stat. at low p_T	± 0.26	± 0.32	± 0.16
Muon reconstruction syst.	$+1.25$ -1.24	± 1.00	± 0.80
Muon reconstruction syst. at low p_T	± 0.70	± 0.71	< 0.01
Muon isol stat.	± 0.55	± 0.65	< 0.01
Muon isol syst.	$+1.52$ -1.51	$+1.46$ -1.45	± 0.80
Muon TTVA stat.	± 0.15	± 0.13	± 0.08
Muon TTVA syst.	± 0.16	± 0.18	± 0.30
Pileup scale factor	—	—	—
Track p_T smearing in the ID	$+0.57$ -0.06	$+1.29$ 0.00	< 0.01
Track p_T smearing in the MS	$+0.37$ -0.42	$+0.47$ -1.29	< 0.01
Muon calibration scale	$+0.09$ -0.38	$+0.46$ -1.29	< 0.01

Table 6.2: Relative uncertainties on the prompt background for the $H \rightarrow 4\mu$ analysis.

6.4.2 Theoretical uncertainties

The sources of theoretical uncertainties on prompt [SM](#) background processes can affect the normalization and shape of the [MC](#) simulation. Uncertainties on the cross-section affect the normalization of [MC](#) samples, while an imperfect simulation will lead to an imprecise shape of kinematic variables.

The choice of [PDF](#) sets affects the results of [MC](#) simulations. Therefore, [PDF](#) parameters are varied. The signal [MC](#) was generated with a set of 100 [PDF](#) variations. In the event generation stage, each [PDF](#) set results in a different weight that is later used as the event weight in a truth analysis. The systematic uncertainty on the acceptance is then assessed as 15.87% of the yields.

The [QCD](#) scale uncertainties on the choice of renormalization and factorization scales are also considered. Six pairwise variations of renormalization μ_r and factorization μ_f scales were considered: $(\mu_r, \mu_f) = (0.5, 0.5), (0.5, 1.0), (1.0, 0.5), (1.0, 2.0), (2.0, 1.0)$ and $(2.0, 2.0)$. The procedure of is the same as it is for the [PDF](#) set uncertainties. The systematic uncertainty on the acceptance is then assessed as the largest variation above and below the nominal.

6.4.3 Uncertainties of heavy-flavor background estimate

Several sources of uncertainty are considered for the heavy-flavor data-driven estimation. The uncertainties associated with the shape of the background estimate are derived by varying each parameter in the model up and down by 1σ and combined to form an envelope around the central value. The envelope is used as an uncertainty on the shape of the background in terms of m_{12} vs. m_{34} . The uncertainty on the background normalization is derived by propagating the statistical uncertainties in the various regions used in the “ABCD” estimate. There is an additional uncertainty on the ratio of events in the template outside of the lepton compatibility window, for $|m_{12} - m_{34}|/m_{12} > 0.15$, where the normalization is applied, and inside, for $|m_{12} - m_{34}|/m_{12} < 0.15$, where the prediction is derived. The uncertainty on the ratio is derived by varying the templates within their uncertainties, recalculating the ratios, and studying the range of variations.

6.5 Statistical interpretation

In ATLAS, a frequentist statistical analysis is used to interpret the results of a search. Typically for the work in this dissertation, the [parameter of interests](#) (POI) in the statistical analysis is the signal strength $\mu = \frac{\sigma_{\text{obs.}}}{\sigma_{\text{theory}}}$, defined as a scale factor of to total number of signal events predicted by the new physics model. In ATLAS, the data in the [signal region](#) must be **blinded** during the design of the research strategy. The results need to first be validated in a background-only fit assuming $\mu = 0$. If there are no significant

deviations in the modeling of data in the [validation region](#), the data in the [signal region](#) will be unblinded. Finally, if there is no significant excess observed in the [signal region](#), exclusion limits will be set using CL_s method described in Section 6.5.4.

6.5.1 Introduction of the general likelihood function

The likelihood function L described the probability of an observation when the physics model is defined. The expected number of events, also known as the expected **yields** y could be generally written as:

$$y = \mu s + b \quad (6.1)$$

where b is the number of background events expected from [SM](#) processes, s is the number of events predicted by the signal model, and μ is the signal strength.

The likelihood function $L(n, \mu, b, \boldsymbol{\theta})$ is a function of the number of data events n , the signal strength μ , expected yields of background b and a set of [nuisance parameters](#) (NP) $\boldsymbol{\theta} = (\theta_0, \theta_1, \dots, \theta_m)$, for m systematic uncertainties. Each θ_i is a [nuisance parameters](#) that continuously interpolates the variation from nominal value for a certain systematic error for $\pm 1\sigma$. The systematic errors must be well controlled otherwise will kill an experimental observation. For instance, if one of the dominant systematic uncertainties θ_k is significantly larger than the rest, the significance would become as

$$\sigma_{\text{sig.}}(\theta_k) = \frac{s}{\sqrt{b(1 + \theta_k^2 b)}} \quad (6.2)$$

Once $\theta_k \gg 1$, a significant observation will be strongly suppressed and very difficult to achieve.

The expression of the likelihood function $L(n, \mu, \boldsymbol{\theta})$ in the most general case of an analysis could be written as the product of a set of Poissonian distributions $P(k|\lambda)$ of event yields in the [signal region](#) (SR) and [control region](#) (CR) [141]:

$$L(n, \mu, b, \boldsymbol{\theta}) = \prod_i P(n_{\text{SR}} | \lambda_{\text{SR}}(\mu, b_{\text{SR}}, \boldsymbol{\theta})) \times \prod_i P(n_{\text{CR}} | \lambda_{\text{CR}}(\mu, b_{\text{CR}}, \boldsymbol{\theta})) \times C_{\text{syst.}}(\boldsymbol{\theta}) \quad (6.3)$$

where i is corresponding to index of binning, and the Poissonian distribution could be generally written as

$$P(k|\lambda) = \frac{\lambda^k e^{-\lambda}}{k!} \quad (6.4)$$

and the λ function is

$$\lambda(\mu, b, \boldsymbol{\theta}) = \mu s + b(1 + \sum_i \theta_i^2 b) \quad (6.5)$$

And finally, the last factor $C_{\text{syst.}}$ is corresponding to the probability density function describing the variations of $\boldsymbol{\theta}$ around $\boldsymbol{\theta}$, where Gaussian function is most commonly used. Therefore, $C_{\text{syst.}}$ could be expressed as

$$C_{\text{syst.}}(\boldsymbol{\theta}^0, \boldsymbol{\theta}) = \prod_i \frac{1}{\sqrt{2}} e^{-\frac{1}{2}(\theta_i^0 - \theta_i)^2} \quad (6.6)$$

where i here is referring to the index of each systematic uncertainties.

6.5.2 Fitting of the nuisance parameters

Once the analysis strategy is defined, the likelihood function $L(n, \mu, b, \boldsymbol{\theta})$ is regressed as a function of μ and $\boldsymbol{\theta}$. Therefore, the likelihood function $L(\mu, \boldsymbol{\theta})$ could be maximized to fit μ and $\boldsymbol{\theta}$ as $L(\hat{\mu}, \hat{\boldsymbol{\theta}})$. The fitting of $H \rightarrow 4\mu$ is handled by `HISTFACTORY` [141] and the fitting in $H \rightarrow 4\tau$ is done by `HISTFITTER` [142]. The pre-fit central value for $\boldsymbol{\theta}$ is usually set to be 0, and the range of variation is set to be $\pm 1\sigma$.

After the fit, the **nuisance parameters** can be affected by the results of fitting as:

- **Pull:** the post-fit **NP** central value non zero. Investigation will be needed if large pull is observed.
- **Constrain:** the post-fit **NP** error is reduced from the pre-fit, indicating that the assigned uncertainty is too large. This case usually needs checking if this is legitimate or coming from a model issue.

6.5.3 Statistical hypothesis testing

The likelihood function L can be also maximized when μ is fixed, and the fitting is only applied on $\boldsymbol{\theta}$, written as $L(\mu, \hat{\boldsymbol{\theta}})$. A test statistics q_μ can be defined as

$$q_\mu = -2 \ln \frac{L(\mu, \hat{\boldsymbol{\theta}})}{L(\mu, \hat{\hat{\boldsymbol{\theta}}})} \quad (6.7)$$

The defined test statistics is called **profile likelihood ratio (PLR)**, since the **NPs** are profiled in order to maximize the likelihood when setting the **parameter of interests**. According to Wilk's theorem [143], **PLR** also follows a χ^2 distribution, profiling “builds in” the effect of **NPs**, hence could be treated as a function of **POI** only.

In the language of statistic hypothesis testing, the signal+background model is considered as the alternative hypothesis H_1 , corresponding to $\mu = 1$, and the background only model is the null hypothesis H_0 , corresponding to $\mu = 0$. There will be two types of error in statistical hypothesis testing:

- **Type-I error** $\text{Prob}(\text{reject } H_0 | H_0)$, which means the rejection of a true null hypothesis.
- **Type-II error** $\text{Prob}(\text{reject } H_0 | H_1)$, which means the failure to reject a false null hypothesis.

The power of an hypothesis test is the probability to reject the null hypothesis when it is indeed wrong. The power of a test will increase as the rate of type-II error decreases. The PLR is known as the most powerful test statistics to reject a null hypothesis [144], which means to minimize the probability of Type-II error for a given level of Type-I error.

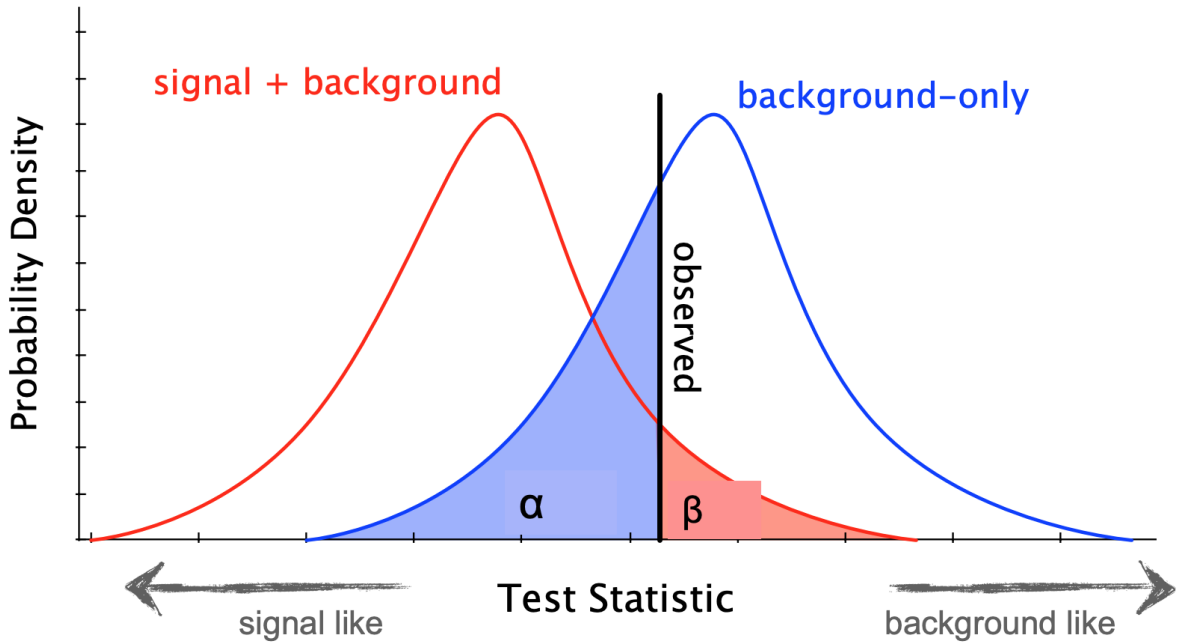


Figure 6.13: An illustration of the p-value calculation, where the blue curve is corresponding to the p.d.f. of background only hypothesis H_0 , and the red curve is corresponding to the signal+background hypothesis H_1 . The black line is the observed value of q_μ^{obs} .

The observed p-value is a measure of the compatibility of the data with the tested hypothesis, which is defined as

$$\alpha = p_0 = \int_{-\infty}^{q_\mu^{\text{obs}}} (q_\mu | H_0) dq_\mu \quad (6.8)$$

$$\beta = p_1 = \int_{q_\mu^{\text{obs}}}^{\infty} f(q_\mu | H_1) dq_\mu \quad (6.9)$$

where $f(q_\mu|\mu')$ is the **probability distribution function** (p.d.f.) of q_μ for signal strength $\mu = 0$, as seen in the Figure 6.13.

There are usually two methods to calculate $f(q_\mu|\mu')$:

- **Asymptotic approach:** if the distribution of $\hat{\mu}$ is Gaussian, the p.d.f. of q_μ could be written as an asymptotic approach [143]:

$$f(q_\mu|\mu') = \frac{1}{2\sqrt{q_\mu}} \frac{1}{\sqrt{2\pi}} \times \left[\exp\left(-\frac{1}{2}\left(\sqrt{q_\mu} + \frac{\mu - \mu'}{\sigma}\right)^2\right) + \exp\left(-\frac{1}{2}\left(\sqrt{q_\mu} - \frac{\mu - \mu'}{\sigma}\right)^2\right) \right] \quad (6.10)$$

where μ' and σ the mean value and standard deviation of μ . In ATLAS, a pseudo data set will be used to approximate σ , known as Asimov data set [145]. Asimov data set is constructed by a pseudo data set defined to be equal to the estimated value with zero signal strength, namely background only. Using the asymptotic approach, the Asimov data set can gives the median results and bands immediately.

- **Toy model:** explicitly build p.d.f. by generating “toy” (pseudo) experiments assuming a specific value of μ and $\hat{\theta}$. The procedure starts by randomizing both main measurement and auxiliary measurement via MC, fitting the model twice for the numerator and denominator of **profile likelihood ratio**, and finally evaluating the PLR and adding to the histogram. This procedure can be very time consuming, but in some cases with tiny statistics can be useful.

6.5.4 CL_s method

If there is no significant excess over the expected background from SM, the data is included in an exclusion fit to derive on-side upper limit at 95% **confidence level** (CL). The limits are calculated by CL_s prescription [146]. To illustrate how CL_s method works, we can use Figure 6.13 as an example. Since the y-axis is the p.d.f., the integral of each red and blue curves equal to 1.

In order to place limits on new physics and solve the problem of data with a downward fluctuation with respect to data, a new quantifier called CL_s is defined as

$$\text{CL}_s = \frac{p_1}{1 - p_0} = \frac{\text{CL}_{s+b}}{\text{CL}_b} \quad (6.11)$$

If $\text{CL}_s < 0.05$, this could be interpreted as the signal+background hypothesis is excluded as 95% **confidence level**.

6.6 Expected upper limits with 139 fb⁻¹

The total background expectation in the $H \rightarrow 4\mu$ analysis is shown in Figure 6.14 as a function of the average di-muon pair mass $m_{\mu\mu} = \frac{m_{12}+m_{34}}{2}$. The error bands shown include the systematic uncertainties of the data-driven method on the heavy-flavor background estimate, and the statistical uncertainty. The statistical uncertainty is the dominant source of error in this analysis.

The total expected yields per background category are summarized in Table 6.3. The distribution of the average muon pair mass $\langle m_{\mu\mu} \rangle = \frac{m_{12}+m_{34}}{2}$ is shown in Figure 6.14, where the error band includes the systematic uncertainty on the heavy-flavor background estimation and the statistical uncertainty.

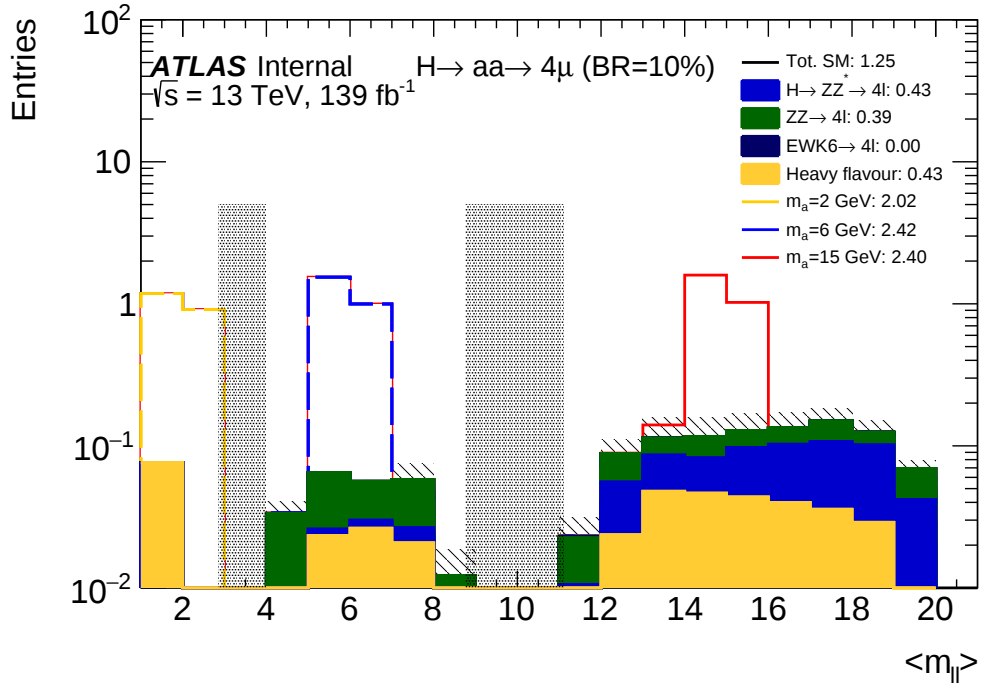


Figure 6.14: Total background expectation in [signal region](#) as a function of $\langle m_{\mu\mu} \rangle$. The error band includes the systematic uncertainty on the heavy-flavor background estimate and the statistical uncertainty.

Process	Expected yields
$H \rightarrow ZZ \rightarrow 4\mu$	0.43 ± 0.01
$ZZ \rightarrow 4\mu$	0.39 ± 0.02
$VVV \rightarrow 4\ell + 2X$	0.04 ± 0.00
Fakes	0.43 ± 0.08

Table 6.3: Total expected background yields in the low mass signal region for the full Run-2 luminosity.

The fitting of maximum likelihood function $L(\mu, \hat{\theta})$ is calculated based on the asymptotic model, following

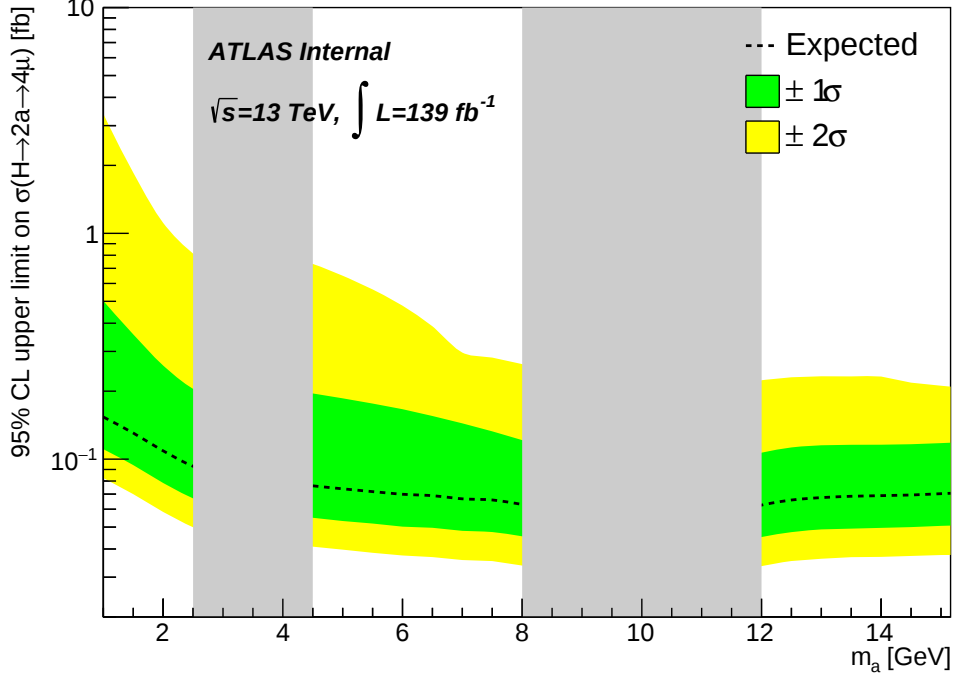


Figure 6.15: The 95% [confidence level](#) upper limit bound on the cross section $\sigma(H \rightarrow 2a \rightarrow 4\mu)$ using the [LHC run2](#) data set for $m_H = 125$ GeV.

the procedure in Chapter 6.5. The upper limits of the cross section $\sigma_{\text{Br}(H \rightarrow 2a \rightarrow 4\mu)}$ can be derived as:

$$\sigma(H \rightarrow 2a \rightarrow 4\mu) = \frac{\mu N_{\text{signal, SR}}}{L_{\text{int.}}} \quad (6.12)$$

were $N_{\text{signal, SR}}$ is the expected number of signal events in the [SR](#) assuming 100% of $H \rightarrow 2a \rightarrow 4\mu$, and $L_{\text{int.}}$ is the integrated luminosity of the used data set which is 139 fb^{-1} , as shown in Figure 6.16. The upper limits on the branching ratio of $\text{Br}(H \rightarrow 2a \rightarrow 4\mu)$ can be calculated as:

$$\text{Br}(H \rightarrow 2a \rightarrow 4\mu) = \frac{\sigma(H \rightarrow 2a \rightarrow 4\mu)}{\sigma(H_{\text{SM}})} \quad (6.13)$$

where $\sigma(H_{\text{SM}})$ is the cross section of [SM](#) via [ggF](#) $\sigma(H_{\text{SM}}) = 48.61 \text{ pb}$ [147].

6.7 Results with 36.1 fb^{-1}

The exactly same analysis was performed on data collected at the ATLAS detector in 2015 to 2016 with an integrated luminosity $L_{\text{int.}} = 36.1 \text{ fb}^{-1}$. The $\langle m_{\mu\mu} \rangle$ distribution in the [SR](#) is shown in Figure 6.17. No

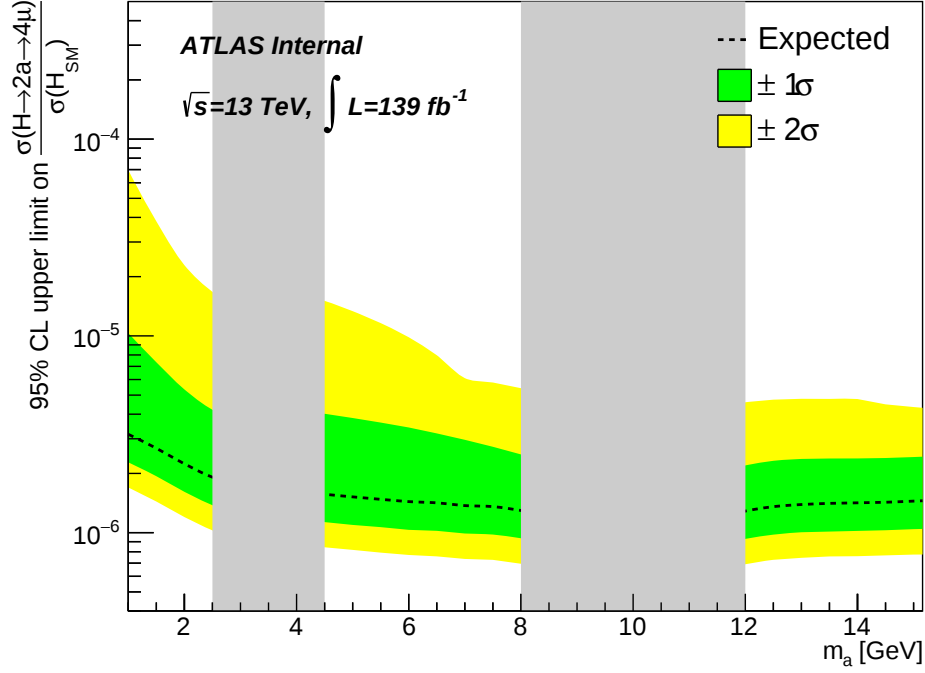


Figure 6.16: The expected 95% [confidence level](#) upper limit bound on the cross section $\text{Br}(H \rightarrow 2a \rightarrow 4\mu)$ using the [LHC](#) run2 data set for $m_H = 125$ GeV.

events are observed to pass the [SR](#) selection, where the total background prediction of [SM](#) is 0.4 ± 0.1 . The two-dimensional panel m_{34} versus m_{12} distribution is shown in Figure [6.18](#), where the crossed dots are 16 data events fails the Higgs mass window selection $120 < m_{4\mu} < 130$ GeV. This is comparable with a prediction of 15 ± 2 events from non-resonant [SM](#) $ZZ^* \rightarrow 4\mu$ process.

The observed limit for $\text{Br}(H \rightarrow 2a)$ is shown in Figure [6.19](#), where the region of $m_a > 15$ GeV is covered by an another analysis dedicated for higher mass [[135](#)], assuming Type-II [2HDM+S](#) model with $\tan \beta = 5$, which have no sensitivity in $2HDM + s$ model but more sensitive for the interpretation of exotic Higgs decays to a pair of vector bosons. More interpretations of the results will be discussed in Chapter [9](#).

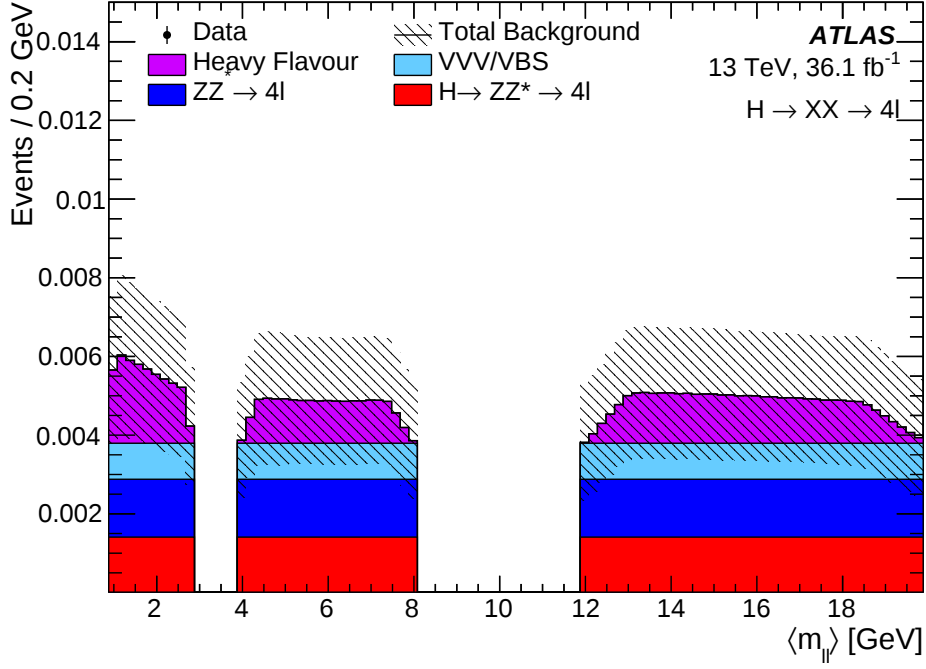


Figure 6.17: Distribution of $\langle m_{\mu\mu} \rangle = \frac{1}{2}(m_{12} + m_{34})$ for the events in the SR of $H \rightarrow 2a \rightarrow 4\mu$ analysis, using the data collected the ATLAS detector with an integrated luminosity of $L_{\text{int.}} = 36.1 \text{ fb}^{-1}$.

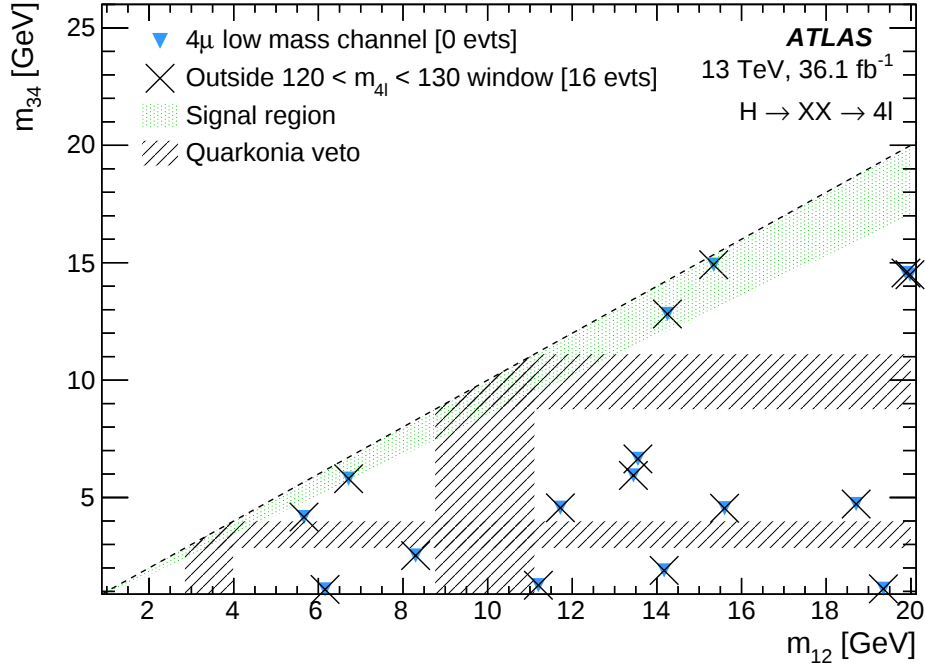


Figure 6.18: Distribution of m_{12} vs m_{34} for events selected in SR of $H \rightarrow 2a \rightarrow 4\mu$ analysis, using the data collected the ATLAS detector with an integrated luminosity of $L_{\text{int.}} = 36.1 \text{ fb}^{-1}$. The crossed events are the ones fails the Higgs mass window requirement $120 < m_{4\mu} < 130 \text{ GeV}$. The events outside the green shaded region are the events fails the compatibility requirement $m_{34}/m_{12} > 0.85$.

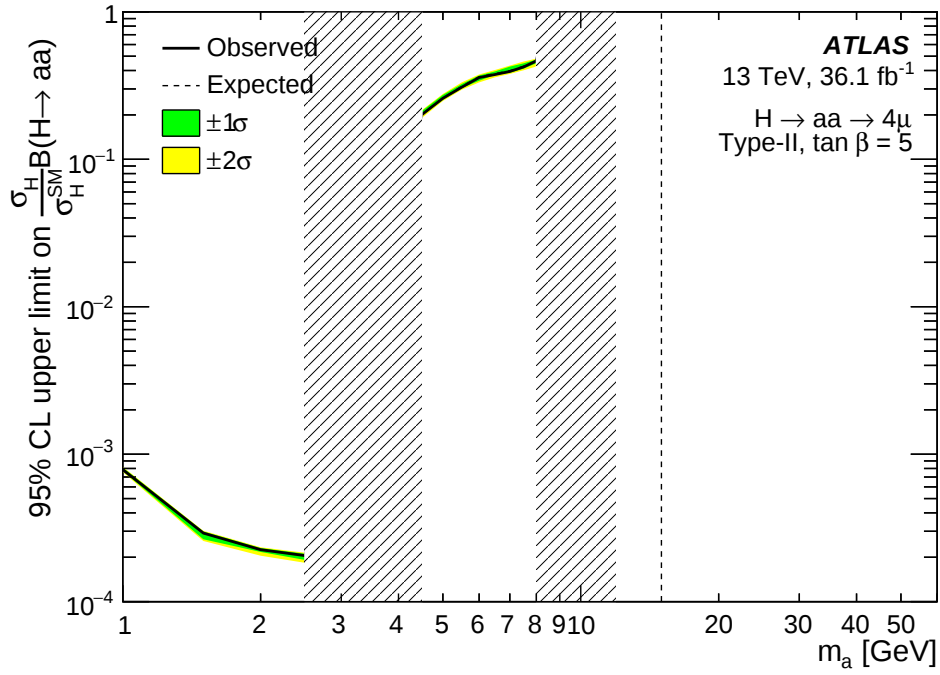


Figure 6.19: Upper limit at 95% CL on the branching ratio of $H \rightarrow 2a$ using the data collected the ATLAS detector with an integrated luminosity of $L_{\text{int.}} = 36.1 \text{ fb}^{-1}$. The benchmark model used here is $2HDM + s$ type-II model with $\tan \beta = 5$.

Chapter 7

Search for exotic Higgs decays to 4τ

This chapter presents a search for exotic Higgs decays into a pair of light bosons (presumed to be pseudoscalars a) in a four- τ final state: $H \rightarrow 2a \rightarrow 4\tau$. In this analysis, the main strategy is to take advantage of the unique properties of the τ leptons, as which can decay both leptonically and hadronically. This analysis focuses on the final state with two **same-sign** (SS) τ decays to leptons, and two **SS** hadronic τ jets. The **SS** selection effectively reduces most of the **SM** background, such as Drell-Yan production. The dominant background is from objects misidentified as leptons or τ jets, called fakes. A data-driven inclusive fake-factor method is used to estimate the fake background.

This analysis performed for the first in ATLAS. The full **LHC** run 2 data set from 2015 to 2018 is used.

7.1 Signal models and features

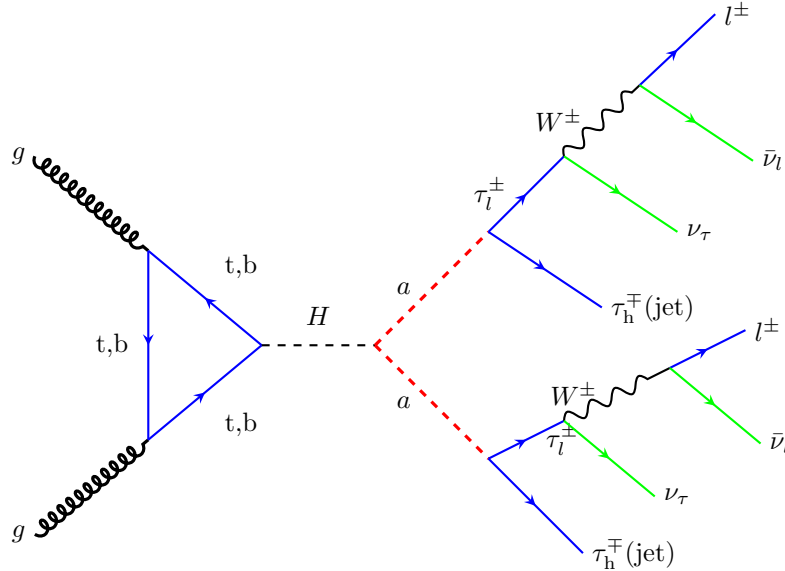


Figure 7.1: The Feynman diagram of the $H \rightarrow 2a \rightarrow 4\tau$ signal model, where the Higgs is produced through **ggF**. The final state requires two **SS** leptons (from τ_l) and two **SS** τ_h jets. Requiring from the same parent a , one hadronically decayed τ_h and another leptonically decayed τ_l . Six neutrinos are produced in this process in total.

The τ lepton is the heaviest lepton in the SM, and can decay not only to lighter leptons (e and μ) but also to hadrons. Therefore, with four τ leptons, it is possible to select specific final states where two τ leptons decay leptonically (τ_e) and two decay hadronically (τ_h). In particular, this search requires that both τ_l have **same-sign** (SS), and both τ_h have **opposite-sign** (OS) to the selected leptons. This requirement can significantly suppress backgrounds, which in a **proton-proton** collider largely consist of **OS** lepton pairs, such $t\bar{t}$, Drell-Yan production of Z bosons, and **QCD** multijet background events with semileptonic decays. Figure 7.1 shows the Feynman diagram of the signal in the $H \rightarrow 4\tau$ analysis.

As in the case of the $H \rightarrow 4\mu$ analysis, the angular separation ΔR between the hadronic τ_h and the lepton from the leptonically decaying τ_l is one of the crucial features of the $H \rightarrow 4\tau$ analysis. Figure 7.2 shows the ΔR distribution for different m_a , from 4 to 60 GeV. Signals with low m_a have very small angular separation between the leptonically decayed τ and hadronically decayed τ , within the size of a jet $R = 0.4$ reconstructed using the anti- k_T4 algorithm. Since the τ jet and the lepton overlapping, the usual τ jet identification algorithm no longer works (as described in Section 5.3.1). Since it is heavily dependent on signatures such as the number of charged tracks (prongs) and the total charge of the jet.

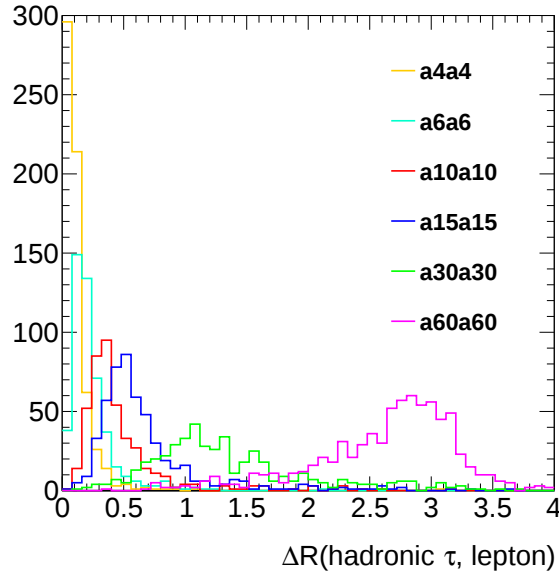


Figure 7.2: The separation between the lepton (e or μ) from a τ -decay and a hadronically-decaying τ for a range of masses m_a (similar behavior is expected for Z_d) in truth level. For very low masses, the pairs become very close together.

Figure 7.3 shows two sketch of some topologies of the signal production. The analysis only targets signals with $m_a > 15$ GeV.

Signal samples of **gluon-gluon fusion** Higgs boson production are generated with POWHEG-BOX [136, 137, 138] using the CT10 PDF set [139] at **next-to-leading order** (NLO), with the same settings as described

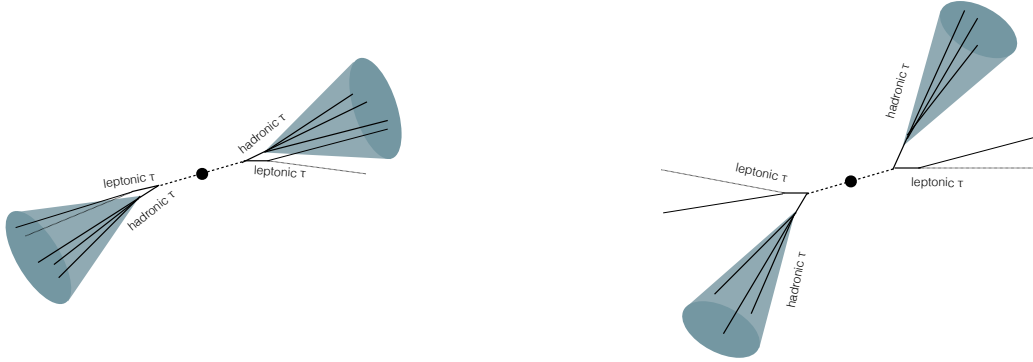


Figure 7.3: Sketches depicting different topologies for the signal. (Left) The case when the lepton and τ jet are overlapping, where the usual identification of τ jets no longer works, due to the influence of the additional charged track from the lepton inside the τ jet cone. (Right) The case when the leptons and τ jets are not overlapping, or resolved, in this case the usual τ jet identification still affords good performance.

in Section 6.1.

A further filter is applied at the generator level to select events with two hadronic τ and two leptonic τ in the final state. The efficiency of this filter is measured to be 30.1%, consistent with τ lepton [branching ratio](#). In this selection, no lepton sign requirement is applied, therefore two thirds of the generated signal pass the signal region selection. In order to get better statistics, we further apply a lepton p_T cut of 3 GeV and a jet p_T cut of 12.5 GeV. The efficiencies of the filter for different signal mass hypothesis are shown in Figure 7.4.

7.2 Object and event selection

In ATLAS, the lowest p_T threshold of unprescaled τ jet triggers is too high for the signal of $H \rightarrow 2a \rightarrow 4\tau$. Single- or di-lepton unprescaled triggers are used.

Electron candidates are required to have $p_T > 7$ GeV and be within the fiducial region $|\eta| < 2.47$, but not within the transition region between the barrel and the end-cap ($1.37 < |\eta| < 1.52$). Selected electrons are required to satisfy the **Medium** likelihood identification working point and the **FCTight isolation** requirement. Since this analysis has a [same-sign](#) requirement, the electrons are required to satisfy the **Tight** electron charge identification working point.

Only combined muons are used in this analysis. Muon candidates are required to have $p_T > 5$ GeV and $|\eta| < 2.5$. Selected muons must satisfy the **Medium** muon identification working point. The [impact parameter](#)

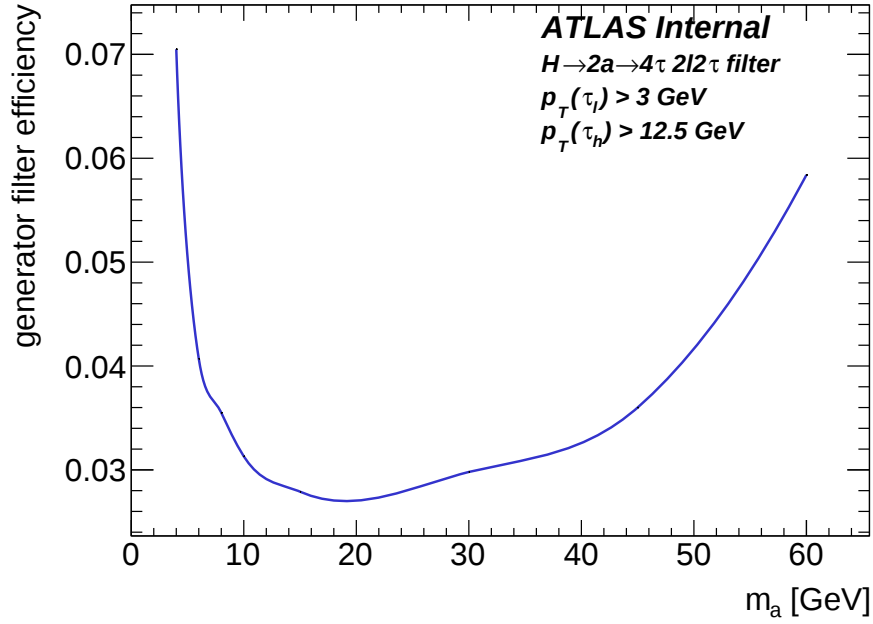


Figure 7.4: The efficiencies of the signal generator filter of $H \rightarrow 2a \rightarrow 4\tau$ selecting events with 2 leptonically-decayed τ and 2 hadronically-decayed τ with $p_T(\tau_{\text{lep}}) > 3 \text{ GeV}$ and τ with $p_T(\tau_{\text{had}}) > 12.5 \text{ GeV}$, as a function of mass m_a .

significance is required to $d_0/\sigma_{d_0} < 4$. This is optimized for muons coming from a τ decay, due to a relatively longer life time of the τ lepton.

Tau jets are expected to contain either one or three charged tracks, corresponding in a narrow jet with low track multiplicity. τ_h candidates are seeded by jets reconstructed with the anti- k_T algorithm [148], which is required with a radius parameter of $R = 0.4$. Selected τ_h must have $p_T > 20 \text{ GeV}$, $|\eta| < 2.5$ but not within $1.37 < |\eta| < 1.52$, and absolute charge equal to 1. The identification working point for a τ jet candidate is **Medium**, of which the identification is described in Section 5.3.1. In order to reduce background from processes involving b -jets, the event must not contain any b -tagged jets. The b -veto algorithm could be found in Section 5.3.2. The missing transverse energy reconstruction is described in Section 5.4.

Event-level selection criteria are applied to reconstruct the two-body and four-body kinematic quantities. Two possible pairings are considered between the leptons and τ jets. For these two pairings, the combination with a smaller scalar sum of the two pair's visible mass is selected as the best pairing.

A requirement that the total visible mass is smaller than 125 GeV is applied to reject events not coming from the decay of a Higgs boson. The m_{T2} is calculated by a bisection-based asymmetric method [149], as shown in Figure 7.6. The m_{T2} value has good discrimination power for different signal masses. To the low statistics, the final limit setting is based on a cut-and-count fit.

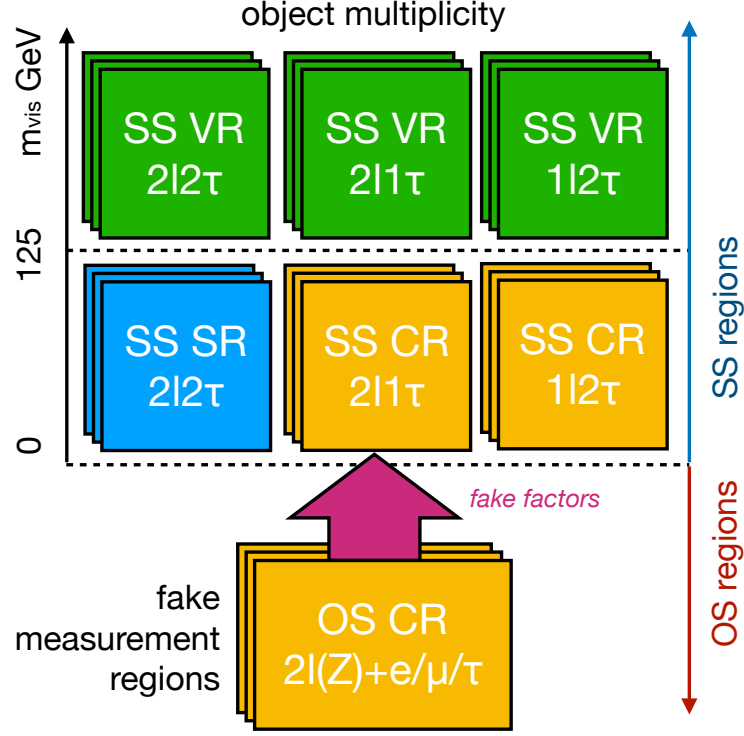


Figure 7.5: A diagrammatic overview of the definitions of SR, CR and VR.

In addition to the **signal region** (SR), two regions are designed as data **control region** (CR), which has either one lepton or one τ jet not reconstructed. These two high statistics CR have small signal contamination. Finally, the estimation of background is cross check in three **validation region** (VR), with an inverted requirement on $m_{\text{vis}} > 125$ GeV.

In order to estimate the contributions from background processes with multiple misidentified objects, a fake factor method is used (see Section 7.3.2). The fake factors are measured in a Z+jets **control region** with Z candidates reconstructed from **OSSF** pairs of leptons. One additional electron, muon or τ jet is also requested.

Figure 7.5 summarizes the definition of these regions. All regions are orthogonal to each other, with the lepton or τ jet multiplicity, m_{vis} or OS/SS requirements.

7.3 Background estimation

Due to the requirement of **SS** leptons and **SS** τ jets, most of **SM** processes with a large cross section in **LHC** do not enter the regions of interests. However, there are many **SM** processes with misidentified objects with

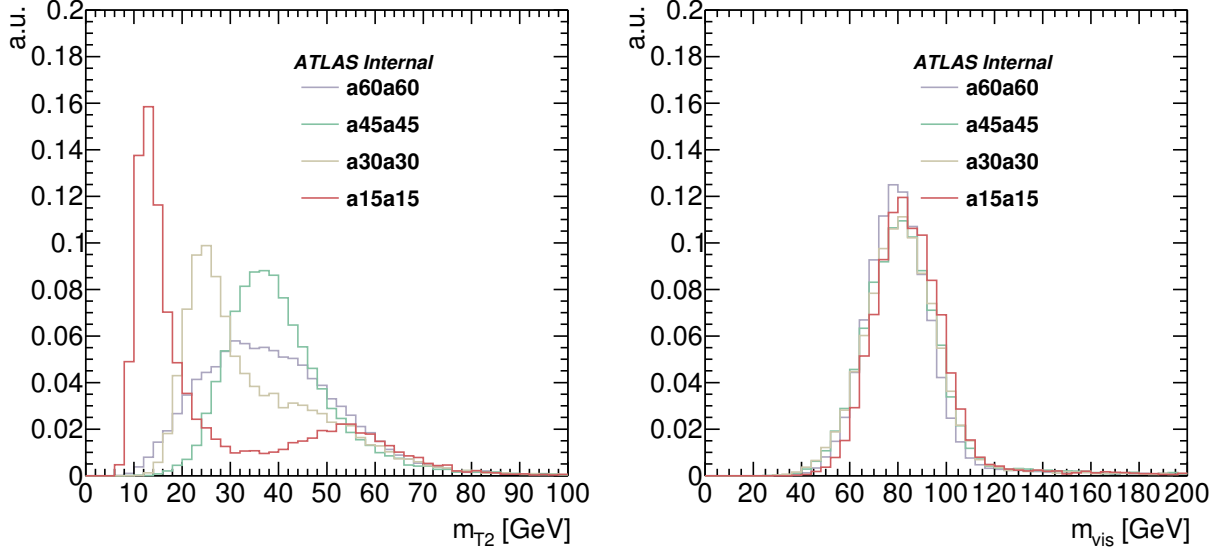


Figure 7.6: The m_{T2} (left) and m_{vis} (right) distribution of $H \rightarrow aa \rightarrow 4\tau$ with masses from 15 to 60 GeV in reconstructed level. The m_{T2} is calculated by a bisection-based asymmetric method [149]. The m_{T2} value has shown good discrimination power for different signal masses. The plots are produced from reconstructed signal events requiring two **SS** leptons and two **SS** τ jets. All distributions are normalized to 1.

a final state as the **SR** selection, which are known as fakes. Even though the misidentification occurs at a very low rate in ATLAS, the fakes become an important contribution if the cross section of the certain process is large enough. Table 7.1 lists the processes which might enter the **signal region** for the $H \rightarrow 4\tau$ analysis, as well as their theoretical cross sections in **LHC** at $\sqrt{s} = 13$ TeV. Usually, the more fake objects, the less likely to have the contribution from the given process. But considering the scale differences of the cross sections is very large, the contribution of multiple fakes is still significant.

As can be seen, due to the complexity of the different types of fakes, it is hard to use **MC** to estimate, because of the statistics, as well as the unreliable simulation. Hence, we use an inclusive data-driven **fake factor (FF)** method to estimate the fake contribution, which is introduced in Section 7.3.2. This method measures the efficiencies of fakes from an inclusive data sample, and derive the systematic uncertainties from differences in fake sources.

7.3.1 Prompt background simulation

The following **SM** processes are considered:

- $ZZ^* \rightarrow 4\tau$: Non-resonant **SM** production of a pair of Z bosons, and no additional final state particles in the matrix element-level process.
- **Triboson production**: Higher-order **electroweak (EWK)** processes with couplings of α^6 , which lead

SM process	$\sigma_{\text{SM}} \times \text{Br}(\text{pb})$ [150]	fake process	
$Z \rightarrow e^+e^- + \text{jets}/\gamma^*$	1.98×10^3	$2\ell(\text{SS})2\tau(\text{SS})$ $2\ell(\text{SS})1\tau$ $1\ell2\tau(\text{SS})$	jet $\rightarrow \tau$, jet/ $\gamma^* \rightarrow e/\mu$, $e \rightarrow \tau$ jet/ $\gamma^* \rightarrow e/\mu$, jet $\rightarrow \tau$ or $e \rightarrow \tau$ $e \rightarrow \tau$, jet $\rightarrow \tau$
$Z \rightarrow \mu^+\mu^- + \text{jets}/\gamma^*$	1.98×10^3	$2\ell(\text{SS})2\tau(\text{SS})$ $2\ell(\text{SS})1\tau$ $1\ell2\tau(\text{SS})$	(jet $\rightarrow \tau$) ² , jet/ $\gamma^* \rightarrow e/\mu$ jet/ $\gamma^* \rightarrow e/\mu$, jet $\rightarrow \tau$ (jet $\rightarrow \tau$) ²
$Z \rightarrow \tau_{\text{had}}^+\tau_{\text{lep}}^- + \text{jets}/\gamma^*$	9.01×10^2	$2\ell(\text{SS})2\tau(\text{SS})$ $2\ell(\text{SS})1\tau$ $1\ell2\tau(\text{SS})$	jet $\rightarrow \tau$, jet $\rightarrow \tau$ jet/ $\gamma^* \rightarrow e/\mu$ jet $\rightarrow \tau$
$W \rightarrow e/\mu/\tau_{\text{lep}} + \text{jets}/\gamma^*$	2.11×10^4	$2\ell(\text{SS})2\tau(\text{SS})$ $2\ell(\text{SS})1\tau$ $1\ell2\tau(\text{SS})$	(jet $\rightarrow \tau$) ² , jet/ $\gamma^* \rightarrow e/\mu$ jet $\rightarrow \tau$, jet/ $\gamma^* \rightarrow e/\mu$ (jet $\rightarrow \tau$) ² or $e \rightarrow \tau$, jet $\rightarrow \tau$
$W \rightarrow \tau_{\text{had}} + \text{jets}/\gamma^*$	1.37×10^4	$2\ell(\text{SS})2\tau(\text{SS})$ $2\ell(\text{SS})1\tau$ $1\ell2\tau(\text{SS})$	jet $\rightarrow \tau$, (jet/ $\gamma^* \rightarrow e/\mu$) ² (jet/ $\gamma^* \rightarrow e/\mu$) ² jet $\rightarrow \tau$, jet/ $\gamma^* \rightarrow e/\mu$
$t\bar{t} \rightarrow \ell^+\ell^- + \text{jets}/\gamma^*$	3.78	$2\ell(\text{SS})2\tau(\text{SS})$ $2\ell(\text{SS})1\tau$ $1\ell2\tau(\text{SS})$	(jet $\rightarrow \tau$) ² , jet/ $\gamma^* \rightarrow e/\mu$ jet/ $\gamma^* \rightarrow e/\mu$, jet $\rightarrow \tau$ (jet $\rightarrow \tau$) ²
$t\bar{t} \rightarrow \tau_{\text{had}}\tau_{\text{lep}} + \text{jets}/\gamma^*$	0.43	$2\ell(\text{SS})2\tau(\text{SS})$ $2\ell(\text{SS})1\tau$ $1\ell2\tau(\text{SS})$	(jet $\rightarrow \tau$) ² , jet/ $\gamma^* \rightarrow e/\mu$ jet/ $\gamma^* \rightarrow e/\mu$, jet $\rightarrow \tau$ (jet $\rightarrow \tau$) ²
$W^+W^- \rightarrow \ell^+\ell^- + \text{jets}/\gamma^*$	0.60	$2\ell(\text{SS})2\tau(\text{SS})$ $2\ell(\text{SS})1\tau$ $1\ell2\tau(\text{SS})$	(jet $\rightarrow \tau$) ² , jet/ $\gamma^* \rightarrow e/\mu$ jet/ $\gamma^* \rightarrow e/\mu$, jet $\rightarrow \tau$ (jet $\rightarrow \tau$) ²
$W^+W^- \rightarrow \tau_{\text{had}}^+\tau_{\text{lep}}^- + \text{jets}/\gamma^*$	0.07	$2\ell(\text{SS})2\tau(\text{SS})$ $2\ell(\text{SS})1\tau$ $1\ell2\tau(\text{SS})$	jet $\rightarrow \tau$, jet $\rightarrow \tau$ jet/ $\gamma^* \rightarrow e/\mu$ jet $\rightarrow \tau$
$WZ \rightarrow \tau_{\text{had}}^+ + \tau_{\text{had}}^+\tau_{\text{lep}}^- + \text{jet}/\gamma^*$	1.94×10^{-3}	$2\ell(\text{SS})2\tau(\text{SS})$	jet/ $\gamma^* \rightarrow e/\mu$

Table 7.1: A list of some SM processes contributing to the regions of interests with different kinds of fakes. The cross section and branching ratio is take from Reference [150].

to four muons in the final state with two additional particles.

- **ttV**: production of $t\bar{t}$ pairs produced in association with [electroweak](#) vector bosons W and Z. This process is modeled by samples generated at [LO](#) using MADGRAPH5_AMC@NLO v2.2.2 and showered with HERWIG++ [151].

The process of $H \rightarrow ZZ^* \rightarrow 4\tau$ is not considered, because of the small cross section and small [branching ratio](#).

7.3.2 Introduction to the inclusive fake factor method

A inclusive [fake factor](#) method is used to estimate backgrounds in the 4τ search.

The total number of objects in the [signal region](#) (SR) N_{tot} passing a “loose” selection L can be expressed

as a sum of the number of real objects N_R and fake objects misidentified by the detector N_F . On the other hand, it can also be considered as a sum of the number of objects that pass the “tight” selection N_T and the number of objects that fail $N_{L'}$, where L' signifies passing the loose requirement L but failing the tight selection.

$$N_{\text{tot}} = N_R + N_F = N_T + N_{L'}.$$

This can be written as a 2×2 matrix:

$$\begin{bmatrix} N_T \\ N_{L'} \end{bmatrix} = \begin{bmatrix} r & f \\ 1-r & 1-f \end{bmatrix} \begin{bmatrix} N_R \\ N_F \end{bmatrix}, \quad (7.1)$$

where the real rate r (usually called efficiency ϵ) is the “tight” selection efficiency for a real object $r = N_T^{\text{real}}/N_R$, and the fake rate f is the rate of a fake object passing the “tight” selection criteria $f = N_T^{\text{fake}}/N_F$.

The matrix equation (7.1) may be expanded as:

$$\begin{aligned} N_T &= rN_R + fN_F = N_T^{\text{real}} + N_T^{\text{fake}}, \\ N_{L'} &= (1-r)N_R + (1-f)N_F = (N_R - N_T^{\text{real}}) + (N_F - N_T^{\text{fake}}). \end{aligned}$$

In practice, we know N_T and $N_{L'}$ but not N_R and N_F . Therefore we invert the matrix to get N_R and N_F in terms of N_T and $N_{L'}$:

$$\begin{bmatrix} N_R \\ N_F \end{bmatrix} = \frac{1}{r(1-f) - f(1-r)} \begin{bmatrix} 1-f & -f \\ r-1 & r \end{bmatrix} \begin{bmatrix} N_T \\ N_{L'} \end{bmatrix},$$

giving

$$\begin{aligned} N_R &= \frac{1}{r(1-f) - f(1-r)} [(1-f)N_T + (-f)N_{L'}], \\ N_F &= \frac{1}{r(1-f) - f(1-r)} [(r-1)N_T + rN_{L'}] = \frac{1}{r(1-f) - f(1-r)} [r(N_T + N_{L'}) - N_T]. \end{aligned}$$

In cases of interest, the rate r (efficiency) is close to one. Approximating $r \approx 1$ gives:

$$\{N_F\}^{\text{approx}} = \frac{1}{1-f} N_{L'}.$$

We know that the number of fakes in the Tight region is by definition $N_T^{\text{fake}} = f N_F$. Therefore, in the ‘tight’ region, the fake estimate is

$$\{N_T^{\text{fake}}\}^{\text{approx}} = \frac{f}{1-f} N_{L'}.$$

From the matrix equation 7.1 we know $N_{L'} = (1-r)N_R + (1-f)N_F$. Substituting for $N_{L'}$ yields

$$\begin{aligned} \{N_T^{\text{fake}}\}^{\text{approx}} &= \frac{f}{1-f} [(1-r)N_R + (1-f)N_F] \\ &= f N_F + \frac{f}{1-f} (1-r)N_R. \end{aligned}$$

The left-hand side of this expression is the approximate value of the fake estimate, and the first term on the right hand side is the exact expression of N_T^{fake} , by definition. Therefore, we can get the accurate fake estimate as:

$$\begin{aligned} N_T^{\text{fake}} &= \{N_T^{\text{fake}}\}^{\text{approx}} - \frac{f}{1-f} (1-r)N_R \\ &= \frac{f}{1-f} N_{L'} - \frac{f}{1-f} (1-r)N_R. \end{aligned}$$

The interpretation of the second term $(1-r)N_R$ is the real objects in the L' (Loose but not Tight) region. It can be easily estimated by [Monte Carlo method](#) simulation. Then the fake estimation in the ‘tight’ region is obtained as a data driven term subtracted by a [MC](#) term:

$$N_T^{\text{fake}} = \left\{ \frac{f}{1-f} N_{L'} \right\}_{\text{data}} - \left\{ \frac{f}{1-f} N_{L'} \right\}_{\text{truth matched MC}}.$$

The coefficient $F \equiv \frac{f}{1-f}$ can be directly measured from the ratio of the number of objects passing the ‘tight’ selection to the number failing. This [fake factor](#) F is then parametrized as a function of the object properties $F(p_T, \eta, \dots)$.

Now we extend the [fake factor](#) method to two objects.

$$\begin{bmatrix} N_{TT} \\ N_{TL'} \\ N_{L'T} \\ N_{L'L'} \end{bmatrix} = \begin{bmatrix} r_1 r_2 & r_1 f_2 & f_1 r_2 & f_1 f_2 \\ r_1(1-r_2) & r_1(1-f_2) & f_1(1-r_2) & f_1(1-f_2) \\ (1-r_1)r_2 & (1-r_1)f_2 & (1-f_1)r_2 & (1-f_1)f_2 \\ (1-r_1)(1-r_2) & (1-r_1)(1-f_2) & (1-f_1)(1-r_2) & (1-f_1)(1-f_2) \end{bmatrix} \begin{bmatrix} N_{RR} \\ N_{RF} \\ N_{FR} \\ N_{FF} \end{bmatrix}$$

Taking the inverse of this matrix we get:

$$\begin{bmatrix} N_{RR} \\ N_{RF} \\ N_{FR} \\ N_{FF} \end{bmatrix} = \frac{1}{(r-f)^2} \begin{bmatrix} (1-f)^2 & (f-1)f & f(f-1) & f^2 \\ (f-1)(1-r) & (1-f)r & f(1-r) & -rf \\ (r-1)(1-f) & (1-r)f & r(1-f) & -rf \\ (1-r)^2 & (r-1)r & r(r-1) & r^2 \end{bmatrix} \begin{bmatrix} N_{TT} \\ N_{TL'} \\ N_{L'T} \\ N_{L'L'} \end{bmatrix}.$$

If $r \approx 1$, this can be simplified as:

$$\begin{bmatrix} N_{RR} \\ N_{RF} \\ N_{FR} \\ N_{FF} \end{bmatrix} = \frac{1}{(1-f)^2} \begin{bmatrix} (1-f)^2 & (f-1)f & f(f-1) & f^2 \\ 0 & 1-f & 0 & -f \\ 0 & 0 & 1-f & -f \\ 0 & 0 & 0 & 1 \end{bmatrix} \begin{bmatrix} N_{TT} \\ N_{TL'} \\ N_{L'T} \\ N_{L'L'} \end{bmatrix}.$$

The fake estimate we want in the TT (Tight & Tight) region is:

$$N_{TT}^{\text{fake}} = rfN_{RF} + frN_{FR} + f^2N_{FF},$$

where

$$\begin{aligned} N_{RF} &= \frac{1}{1-f}N_{TL'} - \frac{f}{(1-f)^2}N_{L'L'}, \\ N_{FR} &= \frac{1}{1-f}N_{L'T} - \frac{f}{(1-f)^2}N_{LL'}, \\ N_{FF} &= \frac{1}{(1-f)^2}N_{L'L'}. \end{aligned} \tag{7.2}$$

Substituting these terms into the fake estimate gives:

$$\{N_{TT}^{\text{fake}}\}^{\text{approx}} = \frac{f}{1-f} N_{TL'} + \frac{f}{1-f} N_{L'T} - \left(\frac{f}{1-f}\right)^2 N_{L'L'}.$$

Using the same trick as in the single-object case, we get

$$N_{TT}^{\text{fake}} = \{F_2 N_{TL'} + F_1 N_{L'T} - F_1 F_2 N_{L'L'}\}_{\text{data}} - \{F_2 N_{TL'} + F_1 N_{L'T} - F_1 F_2 N_{L'L'}\}_{\text{truth matched MC}}.$$

Since the **FF** should be parametrized as a function of the properties of the object, this expression is rewritten as a formula of reweights of each event:

$$N_{TT}^{\text{fake}} = \left[\sum_{TL'} F_2 + \sum_{L'T} F_1 - \sum_{L'L'} F_1 F_2 \right]_{\text{data}} - \left[\sum_{TL'} F_2 + \sum_{L'T} F_1 - \sum_{L'L'} F_1 F_2 \right]_{\text{truth matched MC}}. \quad (7.3)$$

Extrapolating this expression for the three-object case gives:

$$N_{TTT}^{\text{fake}} = \left[\sum_{L'TT} F_1 + \sum_{TL'T} F_2 + \sum_{TTL'} F_3 - \sum_{L'L'T} F_1 F_2 - \sum_{L'TL'} F_1 F_3 - \sum_{TL'L'} F_2 F_3 + \sum_{L'L'L'} F_1 F_2 F_3 \right]_{\text{data}} \quad (7.4)$$

$$- \left[\sum_{L'TT} F_1 + \sum_{TL'T} F_2 + \sum_{TTL'} F_3 - \sum_{L'L'T} F_1 F_2 - \sum_{L'TL'} F_1 F_3 - \sum_{TL'L'} F_2 F_3 + \sum_{L'L'L'} F_1 F_2 F_3 \right]_{\text{truth matched MC}}. \quad (7.5)$$

7.3.3 Fake background estimate

τ jet **fake factor** measurement

The **fake factors** are measured using a tag-and-probe technique, with Z + jets events. The Z + jets **control region** (**CR**) selects events with a Z boson candidate, the “tag”. This method requires two **OSSF** leptons (two muons or two electrons), passing the **Tight** identification requirements and the **FCTight isolation** requirement. The transverse momenta of the leading and subleading leptons must be $p_T > 20$ GeV and $p_T > 15$ GeV, respectively, and their **isol** must satisfy $|d_0|/\sigma_{d_0} < 2.0$. The di-lepton invariant mass must be consistent with that of the Z boson ($|m_{ll} - m_Z| < 10$ GeV). In order to reduce the contribution from heavy-flavor jets, events with a b -tagged jet at the 85% efficiency working point are rejected. A missing transverse momentum cut of $E_T^{\text{miss}} < 20$ GeV is applied to reducing the contributions from the WZ di-boson process, hence reduce the uncertainty from the prompt process subtraction.

The “probe” τ jet must pass the “loose” requirement, which is $p_T > 20$ GeV, $|\eta| < 2.5$ but not within

$1.37 < |\eta| < 1.52$, one or three charged tracks, absolute charge equals to 1 and pass the **BDT** identification working point with 99.5% efficiency. The “tight” selection is the same as described in Section 7.2, which is also the same as the “loose” requirement except using the **Medium** identification working point, as described in Section 5.3.1.

The **fake factor** of a τ jet F could be expressed as:

$$F = \frac{\text{pass “tight” selection}}{\text{pass “loose” but fail “tight” selection}}.$$

In order to take into account the differing pile-up in different data-taking periods, the fake rate and **fake factors** are derived separately for each period, 2015–16, 2017 and 2018. Z+jets **MC** samples produced via SHERPA2.0 are used to cross check the major component in this **control region**. Figure 7.7 shows the kinematics of **MC** and data distributions with the requirement of a Z boson candidate, with “loose” requirement on the “probe” τ jet.

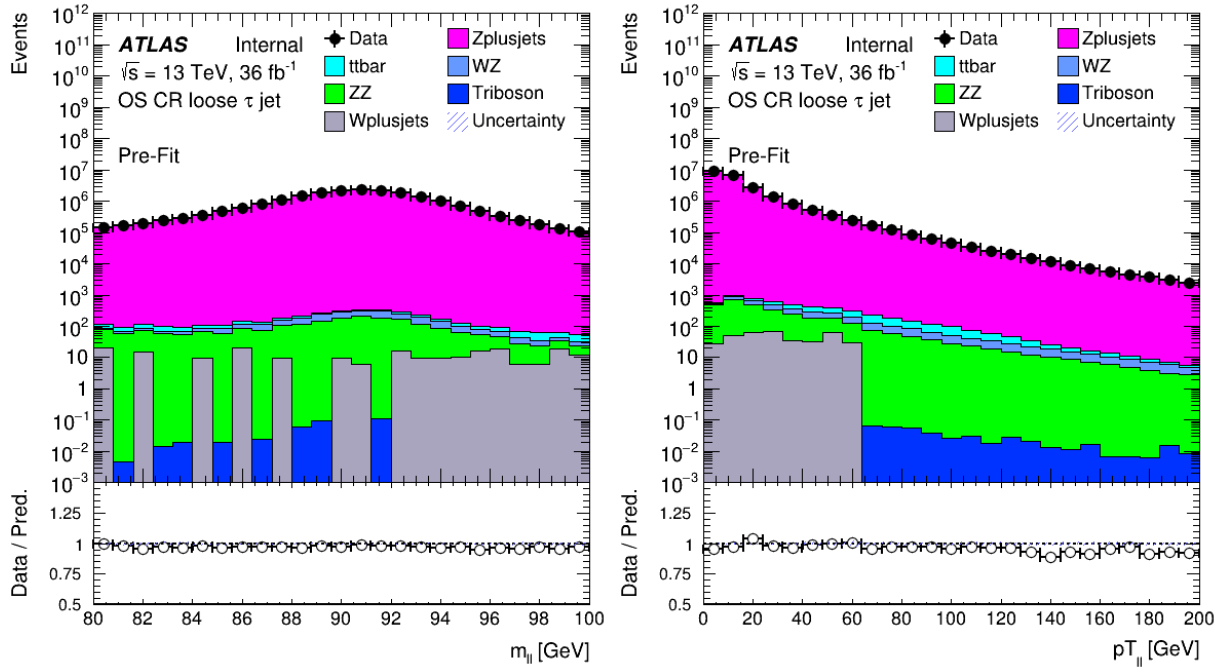


Figure 7.7: Kinematic quantities of the “tag” di-lepton pair in **OS CR** for $H \rightarrow 4\tau$, shown with data collected in 2015-16 compared to the corresponding **MC** simulation, requiring a Z candidate, $E_T^{\text{miss}} < 20$ GeV and a “loose” τ jet. These plots shows good agreement between data and simulation, and the major component are coming from Z+jets as expected.

The tau jet **fake factors** is parametrized as a function of number of charged tracks (1- or 3-prongs). The results of the τ jet **fake factor** measurement is shown in Figure 7.10. The error bars include the statistical

uncertainties and systematic uncertainties of the τ jet FF measurement, as explained in Section 7.4.

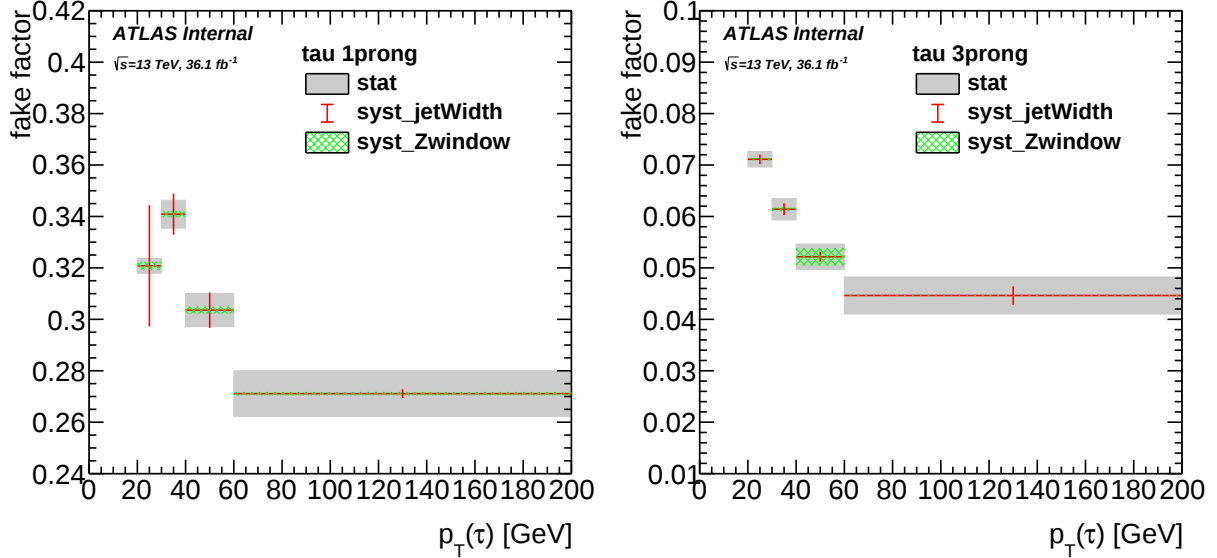


Figure 7.8: Measured τ jet fake factor from the OS CR, including both statistical and systematic uncertainties, separately for (left) one-prong and (right) three-prong τ decays, using the 2015-16 data.

Lepton fake factor measurement

The strategy to measure the fake factor for electrons and muons uses the same tag-and-probe technique as the τ jet fake factor measurement. However, due to the fact that the “tag” lepton pair could be mixed with the “probe” leptons, the OS CR is adapted to optimize the whole procedure.

The “tag” OSSF di-lepton pair is required to have $p_T > 30$ GeV, with an isolation requirement of $E_T^{\text{topocone20}} < 0.01p_T$ and $p_T^{\text{varcone30}} < 0.01p_T$ for both electrons and muons. The impact parameter requirement is $d_0/\sigma_{d_0} < 3$ for muons and $d_0/\sigma_{d_0} < 5$ for electrons. The invariant mass of the OSSF “tag” lepton pairs must be within a Z mass window of $|m_{ll} - m_Z| < 10$ GeV.

In order to avoid the cases where a prompt lepton from a real Z boson is misidentified as a “probe” lepton, the “probe” lepton and the “tag” leptons have to be opposite flavor (OF). Namely, only consider the cases of $2e + \mu$ and $2\mu + e$ are considered, for the muon and electron fake factors respectively. The selection on the “probe” lepton is summarized in Table 7.2.

As described in Section 7.3.2, the fake factor is defined as:

$$F = \frac{N_{\text{“Tight”}}^{\text{data}} - N_{\text{“Tight”}}^{\text{prompt}}}{N_{\text{“Loose”}}^{\text{data}} - N_{\text{“Loose”}}^{\text{prompt}}} \quad (7.6)$$

where N^{data} is the number of events with a probe lepton satisfying the “Loose” or “Tight” requirements,

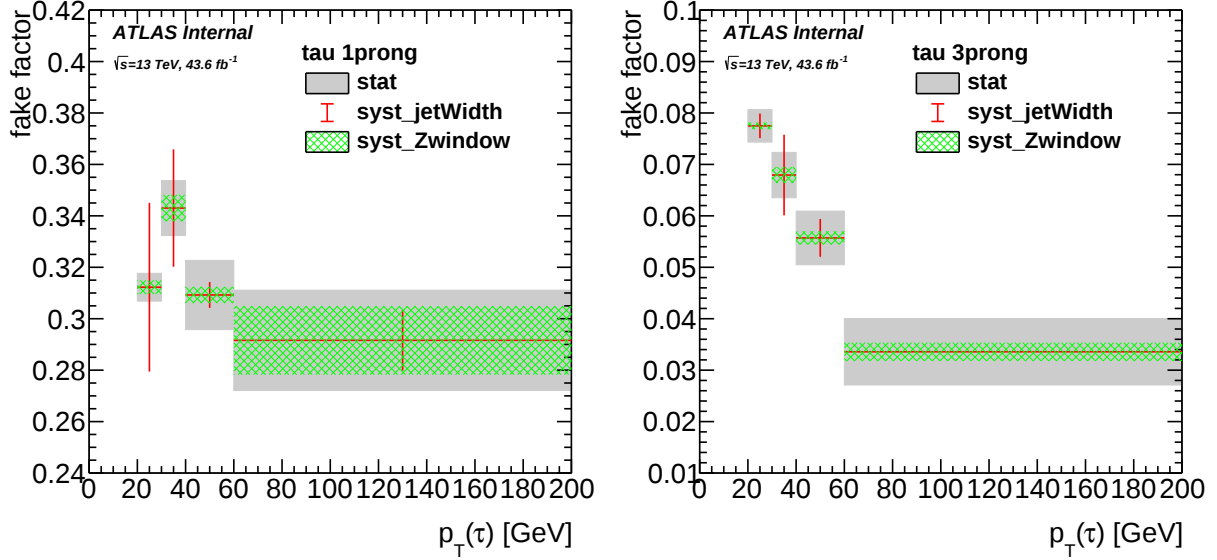


Figure 7.9: Measured τ jet fake factor from the OS CR, including both statistical and systematic uncertainties, separately for (left) one-prong and (right) three-prong τ decays, using the 2017 data.

Electron		Muon	
Loose	Tight	Loose	Tight
$p_T > 7$ GeV		$p_T > 5$ GeV	
-		$ \eta < 2.7$	
$z_0 < 0.5$		$z_0 < 0.5$	
LooseAndBLayerLH	MediumLH	Medium	
$d_0/\sigma_{d_0} < 4$		$d_0/\sigma_{d_0} < 4$	$d_0/\sigma_{d_0} < 7$
-	FixedCutTight isol	-	FixedCutTight isol

Table 7.2: The requirements of “probe” leptons for lepton fake factor measurements.

and N^{prompt} is the number of events with a probe lepton produced from hard scattering, which determined from MC. In this analysis, these processes include SM diboson production of ZZ, WZ, and triboson EWK production, as well as more rare processes of $t\bar{t}$ associated with a vector boson $t\bar{t}W$ and $t\bar{t}Z$.

7.4 Uncertainties on the background estimation

As in the $H \rightarrow 4\mu$ analysis, the main source of the uncertainties on the background estimate are the statistical uncertainties, due to the fact that the final state of two SS τ jets and two SS leptons is very rare in the SM. The systematic uncertainties of the fake factor measurements are also studied and are explained in this section. The experimental uncertainties, such as the object reconstruction and the identification are also considered.

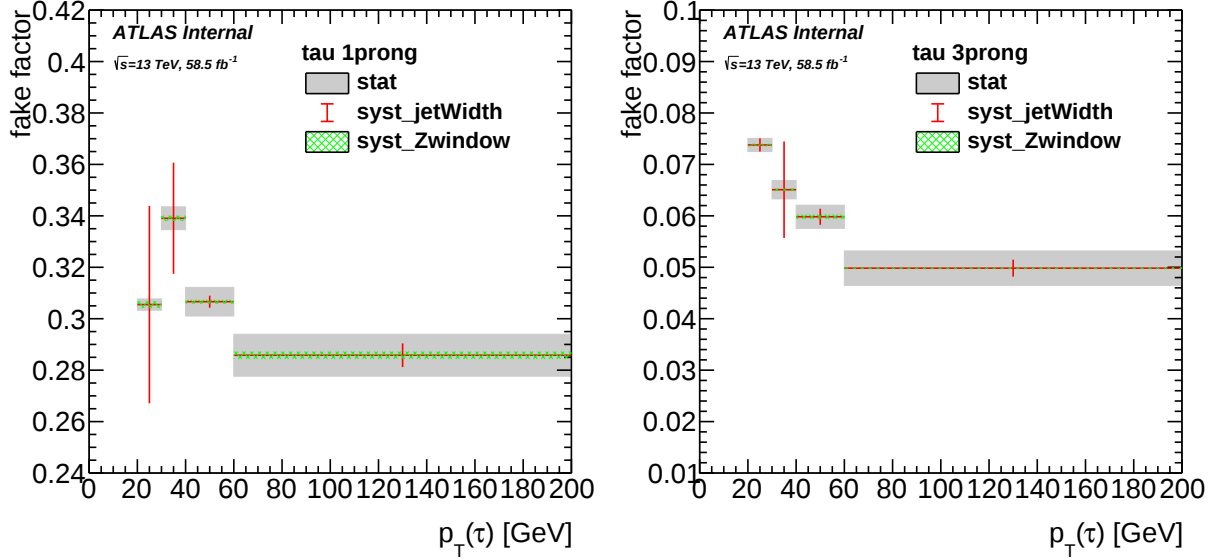


Figure 7.10: Measured τ jet fake factor from the OS CR, including both statistical and systematic uncertainties, separately for (left) one-prong and (right) three-prong τ decays, using the 2018 data.

7.4.1 Systematic uncertainties on the τ jet FF measurement

The systematic uncertainties on the tau jet fake factor measurements are obtained from the following sources:

- **Fake origin composition difference, parametrized as jet width:** Due to the different characteristics of light-flavored quark and gluon jets, the fake factor can depend on the quark/gluon composition of the sample used. So the extent to which this differs between the OS CR (where the fake factor is measured), and the other regions is used as a systematic uncertainty. Figure 7.11 shows the fraction of fake τ jets that originate from light quarks and gluons as a function of $p_T(\tau)$, as obtained from MC truth level information from SHERPA3.0 Z + jets simulation. These fractions are similar in the measurement region (OS CR) and the fake τ jet enriched validation region (SS VR with one lepton and two τ jets).

The τ jet width as a quantity sensitive to the quark-gluon fraction, which is shown in Figure 7.12. The τ jet width is calculated from the tracks associated to the τ jet:

$$\text{jet width}(\tau) = \frac{\sum_i p_T(\text{trk}^i) \times \Delta R(\text{trk}^i, \text{center})}{p_T(\tau)}$$

where i is the i^{th} track associated to the τ jet. The gluon jets have more radiation and so tend to be wider than light-quark jets. The distributions of τ jet widths in the measurement region OS CR and

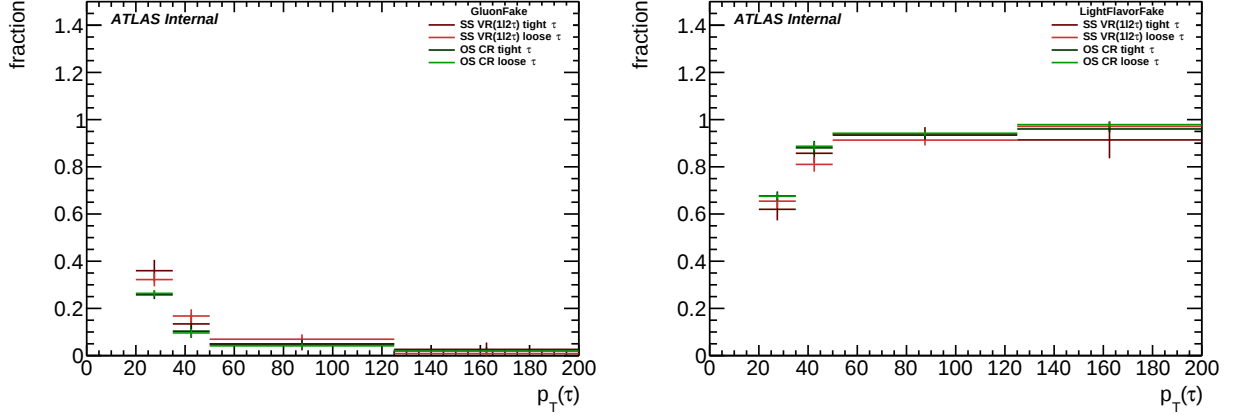


Figure 7.11: Fake τ jet fraction as a function of $p_T(\tau)$, obtained from MC simulation. The fraction of the probe τ jet truth origin, either from a gluon (left) or light-flavored quark (right), with “loose” (lighter color) or “tight” (darker color) requirement in the OS CR (Green) and the SS VR (Red) with one lepton and two τ jets. The fraction coming from heavy flavor quark is very small, because of the b-veto selection.

validation region are shown in Figures 7.13 and 7.14 for 1-prong and 3-prong τ jets respectively.

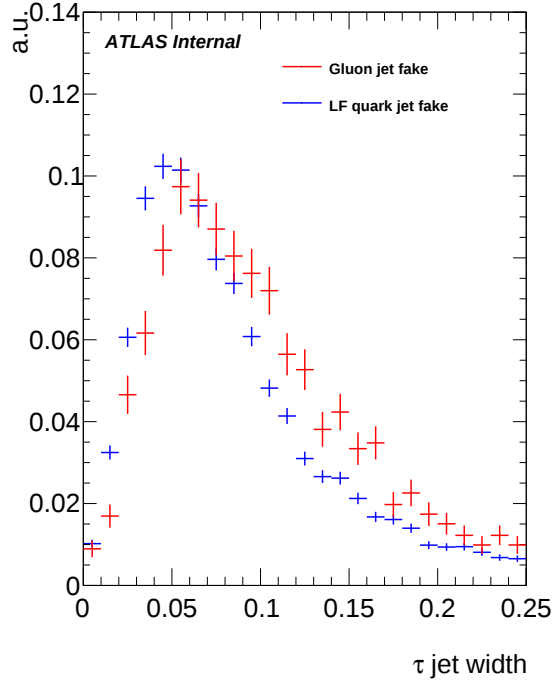


Figure 7.12: Normalized τ jet width distribution from light flavor quark (red) fakes and gluon (blue) jet fakes. Using a SHERPA Z(ee)+jets MC sample with heavy flavor jets vetoed in generator level.

A re-weighting factor as a function of $p_T(\tau)$ and jet width is obtained by dividing the normalized two-

dimensional distribution of $p_T(\tau)$ vs jet width in the **SS VR** to that in the **OS CR**. This re-weighting factor is further applied to the **OS CR** to calculate the **fake factor** variation associated to the τ jet width, which will correspond to the systematic uncertainty due to the differing quark/gluon fraction between the measurement region (**OS CR**) and the **SS validation region**.

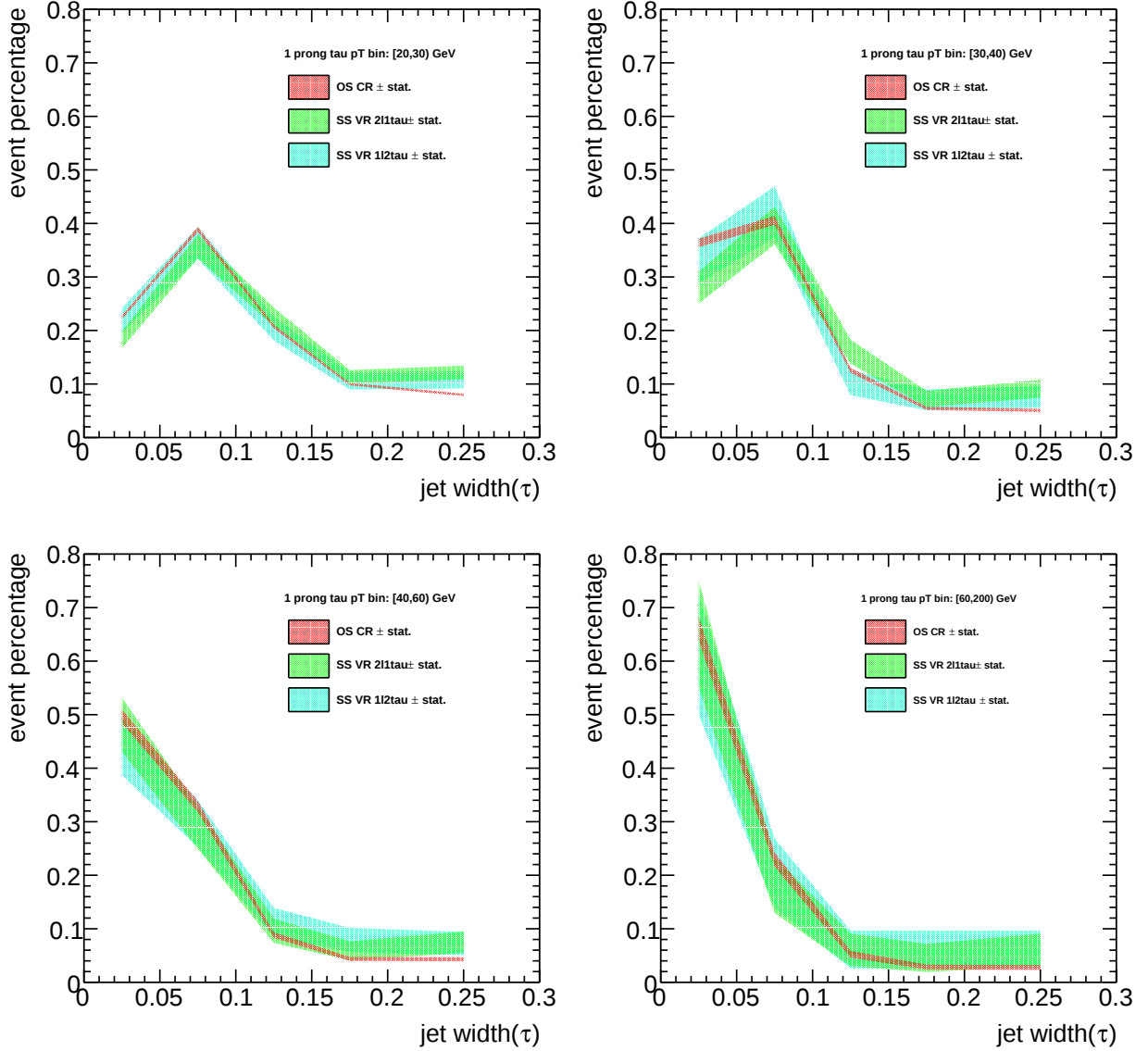


Figure 7.13: τ jet width distribution, normalized in each p_T bin, obtained from 2015–6 data for one-prong τ jets. Comparison among the measurement region **OS CR** and **validation region** is shown with the corresponding statistical uncertainties.

- **The Z tag mass requirement:** The measured **fake factor** depends on the chosen of Z boson mass window. It is measured with the Z boson mass window widths differed, from 20 GeV (± 10 GeV) to

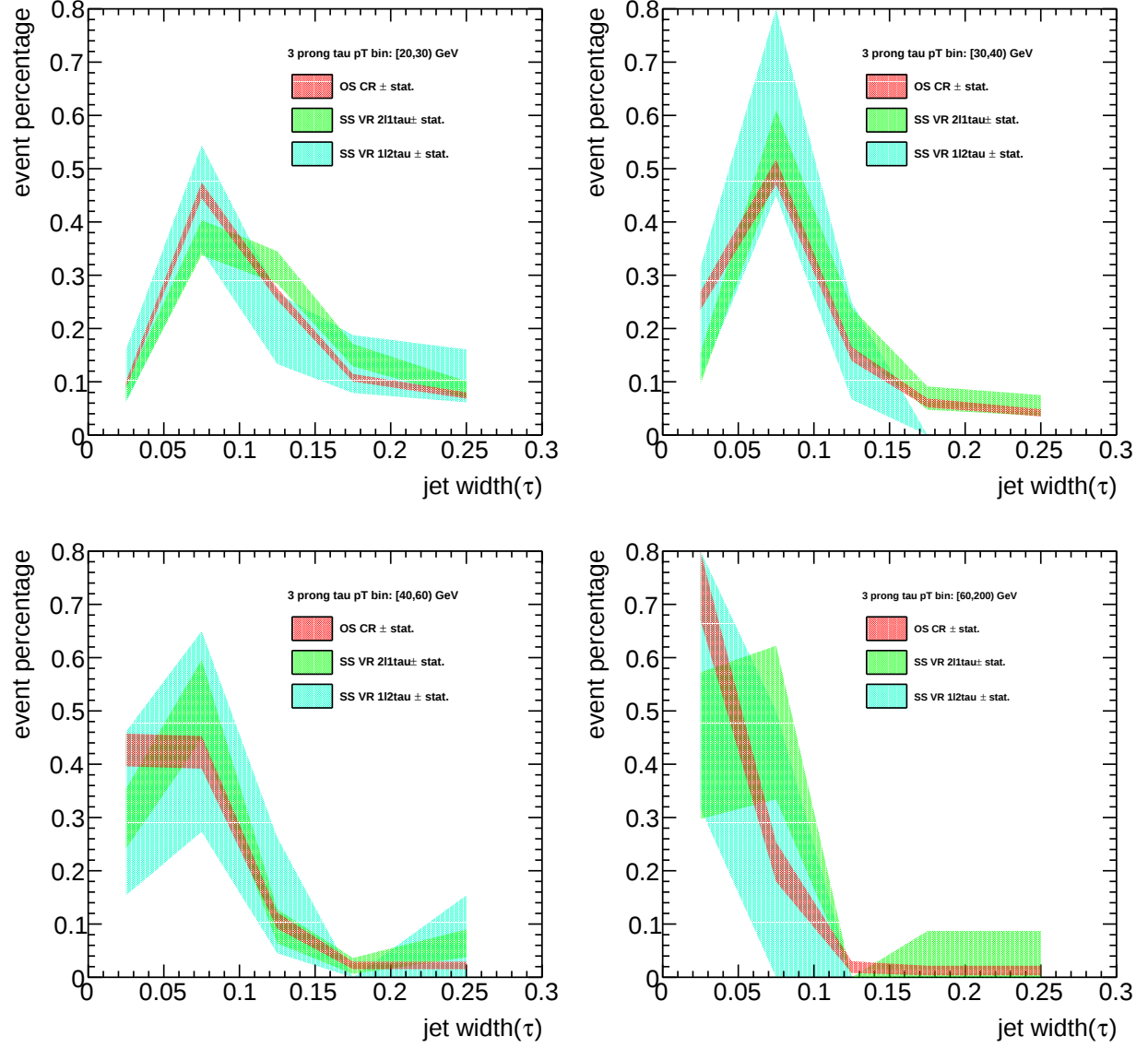


Figure 7.14: τ jet width distribution, normalized in each p_T bin, obtained from 2015–6 data for three-prong τ jets. Comparison among the measurement region OS CR and validation region is shown with the corresponding statistical uncertainties.

10 GeV (± 5 GeV) , and the variation taken as a systematic uncertainty.

- **Statistical uncertainty in the OS CR:** The uncertainty from finite statistics in the measurement region is propagated as a systematic uncertainty when the measured **fake factor** is applied in the **same-sign validation region** and **signal region**.
- **Prompt process modeling:** When the prompt process is subtracted, a 20% variation is applied on the normalization factors of the **electroweak** processes. This variation is meant to include the theoretical uncertainties in the calculation of **electroweak** cross sections due to variations in the factorization and renormalization scales.

7.4.2 Systematic uncertainties of lepton **FF** measurement

The systematic uncertainties on the lepton **fake factor** is measured using the following criteria:

- The statistical constraint in the **OSCR** region for the measurement. Based on the expression for the lepton **FF** in Equation 7.6, the statistical uncertainty on the measurement is:

$$\sigma_{\text{stat}} = \sqrt{q_1^2 + q_2^2} \quad (7.7)$$

where

$$\begin{aligned} q_1 &= \sqrt{p_1 + p_2} \times \frac{F}{(1 + F)^2} \\ q_2 &= \sqrt{p_3 + p_4} \times \frac{F}{(1 + F)^2} \end{aligned} \quad (7.8)$$

in which

$$\begin{aligned}
p_1 &= \left(\frac{\sqrt{N_{\text{"Loose"}}^{\text{data}} - N_{\text{"Tight"}}^{\text{data}}}}{N_{\text{"Loose"}}^{\text{prompt}} - N_{\text{"Tight"}}^{\text{prompt}}} \right)^2 \\
p_2 &= \left(\frac{\sqrt{N_{\text{"Tight"}}^{\text{data}}}}{N_{\text{"Tight"}}^{\text{prompt}}} \right)^2 \\
p_3 &= \left(\frac{\sqrt{\sum_{\text{prompt MC}} (\sigma_{\text{"Loose"}}^{\text{prompt}})^2 - \sum_{\text{prompt MC}} (\sigma_{\text{"Tight"}}^{\text{prompt}})^2}}{N_{\text{"Loose"}}^{\text{prompt}} - N_{\text{"Tight"}}^{\text{prompt}}} \right)^2 \\
p_4 &= \left(\frac{\sqrt{\sum_{\text{prompt MC}} (\sigma_{\text{"Tight"}}^{\text{prompt}})^2}}{N_{\text{"Tight"}}^{\text{prompt}}} \right)^2
\end{aligned} \tag{7.9}$$

The above calculation is derived from the propagation of statistical uncertainties from data and the **MC** samples using for subtraction.

- The uncertainties from the prompt subtraction process. In the step of Equation 7.5, the last term is obtained from the prompt **MC** samples. The uncertainty on the cross sections for these **MC** processes can be impactful on the measurements of the **fake factor**. Hence, we set an uncertainty of 20% on the cross sections of all the subtracted **MC** samples.
- The truth composition of a fake electron or muon in the **OSCR** is not necessarily the same as the region in which the **FF** is applied, such as the **VR** and **SR**. The main origin of a truth electron can be either from a photon conversion $\gamma^* \rightarrow e^+e^-$ or a misidentified quark jet. The main source of a muon can be either a light flavor quark jet or a semi-leptonically decaying heavy flavor jet. For an electron, in order to enrich the sample in fake electrons from photon conversions, a sample of reconstructed electrons seeded by either photons or electrons is used. For muons, the d_0/σ_{d_0} cut is varied for the "Loose" muon probe, from $d_0/\sigma_{d_0} < 7$ (nominal cut) to $d_0/\sigma_{d_0} < 4$, in order to vary the contributions from light flavor jet fakes.

7.4.3 Experimental systematic uncertainties

Experimental systematic uncertainties associated to τ jets arise from the τ calibration, reconstruction, identification and the efficiency of overlapping removal with respect to electrons. The **scale factors** are

measured from the difference between data and MC for the $Z \rightarrow \tau\tau$ process with a tag-and-probe technique.

The b -jet veto uncertainties are calculated by reversing the truth b -jet identification efficiency, using the scale factors of the mv2c10 b -tagger obtained from Reference [152]. The final weight implemented on each MC simulated sample with truth b -jets is:

$$\sum_{\text{iterating over truth } b\text{-jets}} \left(\frac{1 - \epsilon_b \text{sf}_b}{1 - \epsilon_b} \right)$$

where $\epsilon_b = 85\%$ is the efficiency of the chosen b -tagging working point for the b -veto.

The systematic uncertainties associated with the measurement of E_T^{miss} are derived for the response as well as for the resolution. They depend on the composition of the hard terms and on the magnitude of the corresponding soft term, as described in Section 5.4. The extraction of the systematic uncertainties for the reconstructed E_T^{miss} is based on data-to-MC comparisons of spectra of observables measuring the contribution of the softer terms with respect to the overall E_T^{miss} [133]. The systematic uncertainty on the E_T^{miss} resolution is determined as the root mean square of E_T^{miss} projections on x and y directions.

Electron identification, reconstruction and isolation efficiency uncertainties are considered, implemented by scale factors measured from the data-to-MC comparison of $Z \rightarrow ee$ events. The kinematic smearing and calibration uncertainties of electrons are also considered. For electron calibration uncertainties, a transfer factor is extracted from detector raw hit value to energy. A smearing algorithm is used to simulate the detector resolution for MC samples by pseudo randomly modifying the kinematics of the objects according to an assumed probability distribution.

The systematic uncertainties of muon reconstructions and energy measurements are obtained using the same procedure as in Section 6.4.1.

7.5 Validation of the background estimate in the $\mu\mu$ channel

In order to verify that the measured inclusive FF in Section 7.3.3 and Section 7.3.3 can well describe the fake component in the interested SS regions, two SS VR with $m_{\text{vis}} > 125$ GeV are designed to validate the background estimation. The regions are split by lepton flavor. This section concentrates on the $\mu\mu$ channel.

A maximum likelihood function fit is performed. The regions that participate in the fit in the following regions:

- The SS signal region ($2\tau 2\mu$) with the Asimov data set.
- The SS control region ($2\tau 2\mu$) inverting the mass cut $m_{\text{vis}} > 125$ GeV

- The [SS control region](#) ($1\tau 2\mu$) with is enriched in statistics.

The [CR](#) are fitted with the data of 139 fb^{-1} . The results of each region before and after fit are summarized in Figures [7.15](#).

7.6 Result in the $\mu\mu$ channel

The final estimation of the [signal region](#) in the $\mu\mu$ channel of the $H \rightarrow 4\tau$ analysis is summarized in Table [7.3](#). The total yield estimation of background in the [SR](#) is 0.24 ± 0.12 . 1 data event is observed from the 139 fb^{-1} dataset collected at the ATLAS detector, which is consistent with the [SM](#) prediction, considering the ± 1 Poisson error in the data.

Process	Expected yields
$ZZ \rightarrow 4\tau$	0.11 ± 0.07
Fakes	0.13 ± 0.10

Table 7.3: Total expected background yields in the $\mu\mu$ channel of $H \rightarrow 4\tau$ in the [signal region](#) for the full Run-2 luminosity.

The upper limit of $\text{Br}(H \rightarrow 2a \rightarrow 4\tau)$ in the $\mu\mu$ channel at 95% [CL](#) is set, as Figure [7.16](#) shows. The further interpretation with regard of the exclusion limits of the $2HDM + s$ model is discussed in Chapter [8](#).

7.7 Suggestion for future searches

As discussed in Section [7.1](#), the sensitivity of the $H \rightarrow 4\tau$ decreases when the mass of the pseudoscalar is lower than 15 GeV. This is because the angular separation ΔR between the hadronic τ_h and the lepton from the leptonically decayed τ_l will be less than 0.4, which is the radius of the reconstructed τ jet based on Anti- k_T algorithm with $R = 0.4$. The current τ identification is using a [BDT](#) based algorithm, which dependent on the signatures such as numbers of charged tracks and the total charge. Hence, once an additional track inside the jet cone, this identification is no longer reliable.

This section presents some preliminary studies on the new identification algorithm of these overlapped di-tau objects. In the sake of simplify the problem, we consider a leptonic tau decays to muons overlapping with a hadronic τ , since a track of muon is much easier to distinguish comparing to an electron.

The di-tau identification uses non-isolated [combined muon](#) (CB) tracks as a seed, and looks for a reconstructed τ jet within $\Delta R < 1.0$ vicinity. The τ jet identification variables will be recalculated with a removal of the muon impact, and then feed the [BDT](#) score. The identification efficiency as a function of ΔR

between the muon track and the center of the τ jet cone is shown in Figure 7.17, using the signal samples of $m_a = 10, 30$ GeV. It can be seen that the standard tau BDT maintains a reasonable efficiency in the region of $\Delta R > 0.4$, but totally no identification power in the region of $\Delta R < 0.4$, whereas the new di-tau identification keeps the efficiency in a level of 60 – 80% in all ranges of ΔR .

Figure 7.18 shows the ROC curves of the new di-tau identification algorithm. In order to select a sample with the same kinematics and final state expect with a quark jet instead of τ jet, the background sample is $H \rightarrow \text{LQ} \rightarrow c\mu c\mu$, where leptoquark (LQ) is BSM particle decaying to a lepton and a quark [153], where the mass is set to be $m_{\text{LQ}} = 12$ GeV. It can be seen that the discrimination power of the new di-tau identification remains the same as the standard algorithm in the case of $\Delta R > 0.4$, where the muons and τ jet are still separated. But once $\Delta R < 0.4$, the new di-tau identification algorithm has an overwhelmingly better performance.

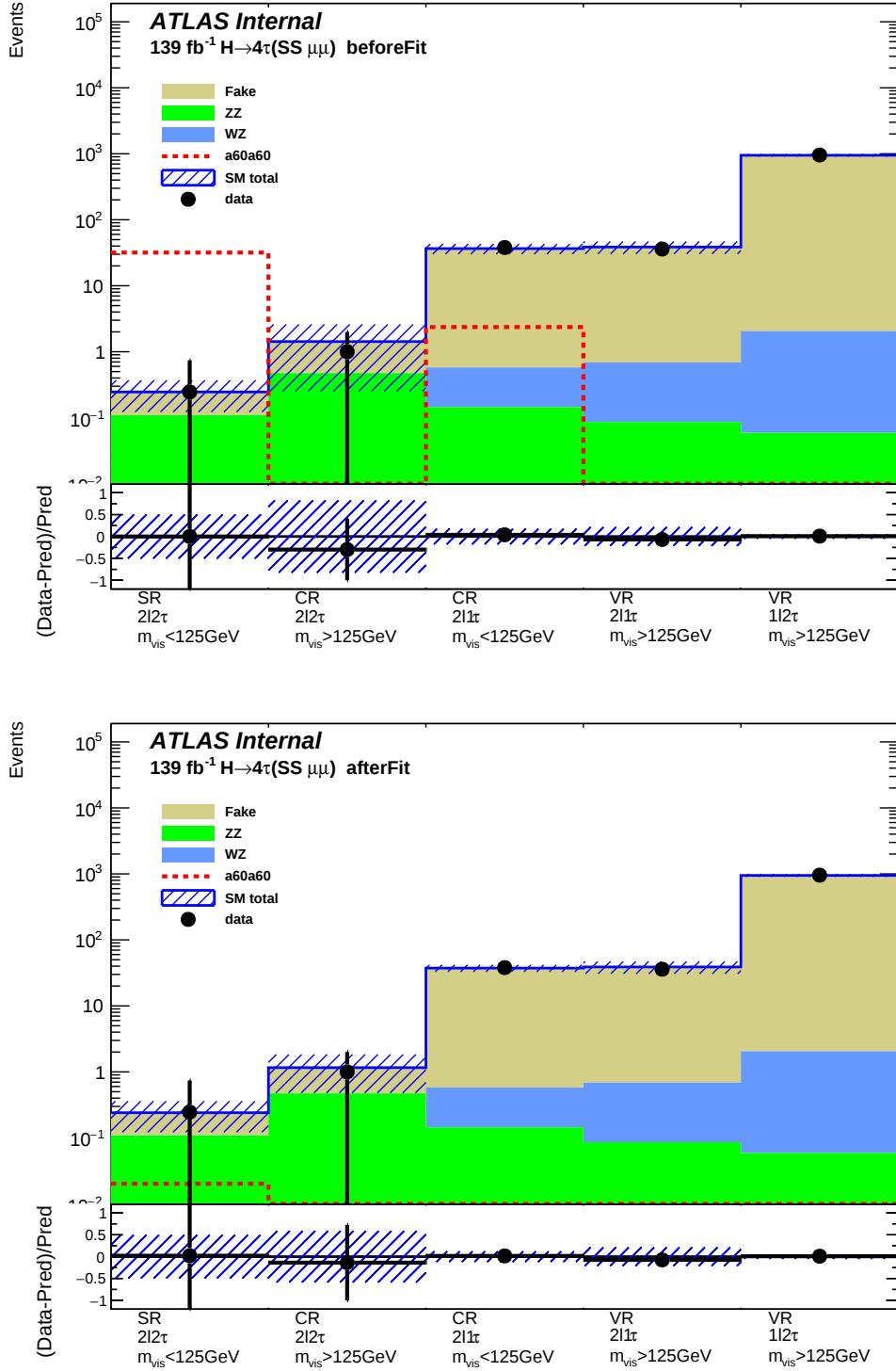


Figure 7.15: Before(up) and after(down) the background-only fit on the **control** regions and the extrapolation to the **validation** region of the cut-and-count analysis. All uncertainties are included in the uncertainty band. The lower bands shows the relative disagreement between the data and the prediction.

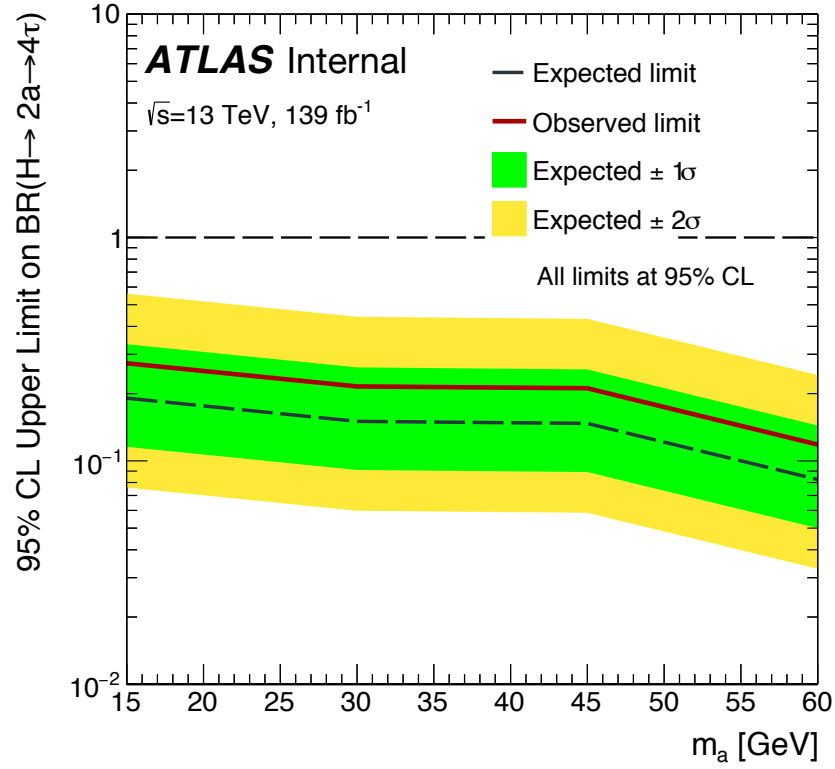


Figure 7.16: The 95% CL upper limit on the branching ratio of $\text{Br}(H \rightarrow 2a \rightarrow 4\tau)$ in the $\mu\mu$ channel. The expected limit is shown in dashed line and the observed limit is in red solid line.

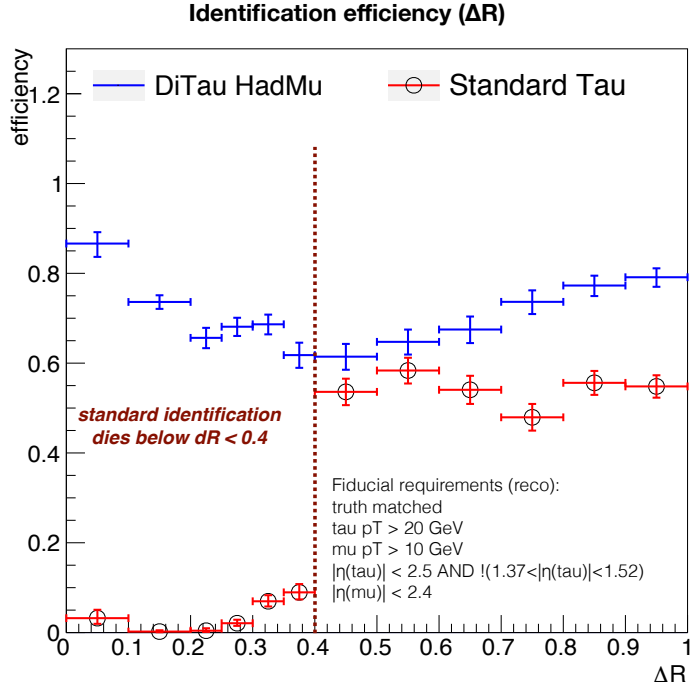


Figure 7.17: The efficiency as a function of $\Delta(\tau, \mu)R$ between the muon track and the center of the τ jet for the new ditau identification algorithm and the standard τ jet algorithm.

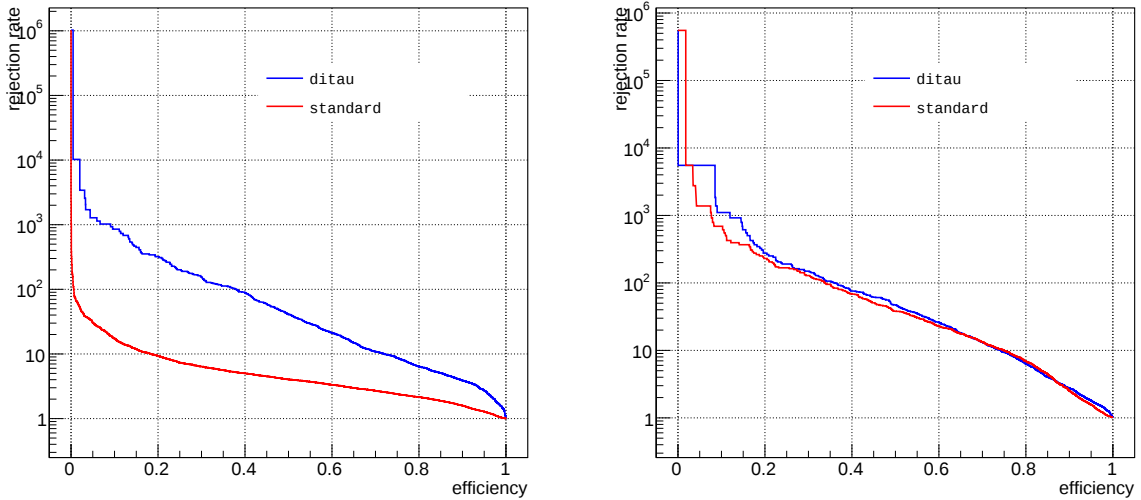


Figure 7.18: The ROC curves for $\Delta R < 0.4(\tau, \mu)$ (left) and $\Delta R(\tau, \mu) > 0.4$ (right) for the new ditau identification algorithm and the standard τ jet algorithm.

Chapter 8

Discussions and interpretations

In this chapter, we will discuss about the results of the two analyses presented in this thesis and their interpretations, using the 2HDM+S model described in Section 3.2.

Based on the results of the analyses of $H \rightarrow 2a \rightarrow 4\mu$ and $H \rightarrow 2a \rightarrow 4\tau$ (SS $\mu\mu$ channel), no significant excess of data over the expected SM background has been observed. A strong upper limits of the $\text{Br}(H \rightarrow 2a \rightarrow 4\mu/4\tau)$ have been set at a 95% CL. The interaction between the pseudoscalar a with the SM particles are discussed in Section 3.2, hence an upper limit of $\text{Br}(H \rightarrow 2a)$ could be derived, assuming the Higgs boson production section is the same as the predicted ggF Higgs cross section in the SM.

The following figures show the ATLAS summary of the observed and expected 95% CL upper limits on $\frac{\sigma_h}{\sigma_{\text{SM}}} \times \text{Br}(H \rightarrow aa)$ in 2HDM+S model with different types of interactions between the two Higgs doublets and SM fermions. Figures 8.1- 8.10 shows the upper limits of $\frac{\sigma_h}{\sigma_{\text{SM}}} \times \text{Br}(H \rightarrow aa)$ as a function of the pseudoscalar mass m_a with different $\tan\beta$ values. Figures 8.11-8.13 show the upper limits of $\frac{\sigma_h}{\sigma_{\text{SM}}} \times \text{Br}(H \rightarrow aa)$ as a function of $\tan\beta$ assuming $m_a = 40$ GeV.

The upper limits gained in the works from this thesis is shown. The main constrain of $H \rightarrow 4\mu$ is statistics. Using the full LHC run2 data set improved by about a factor of 4 on limit strength compared to the one with data from 2015 and 2016. The $H \rightarrow 4\tau$ is a new search in ATLAS, which is able to access the phase space with a strong complementarity to the others. In the 2HDM+S Type-III model, the coupling between a and τ lepton is dominant, hence strong limits have been set from $H \rightarrow 4\tau$ analysis.

As can be seen, according to the Yukawa-like coupling, the $H \rightarrow 4\mu$ analysis gives good sensitivity in the low mass region. The $H \rightarrow 4\tau$ analysis has the strong limit exclusion for 2HDM+S Type-III model, as the coupling between the a and the third generation of lepton is preferred especially for large $\tan\beta$.

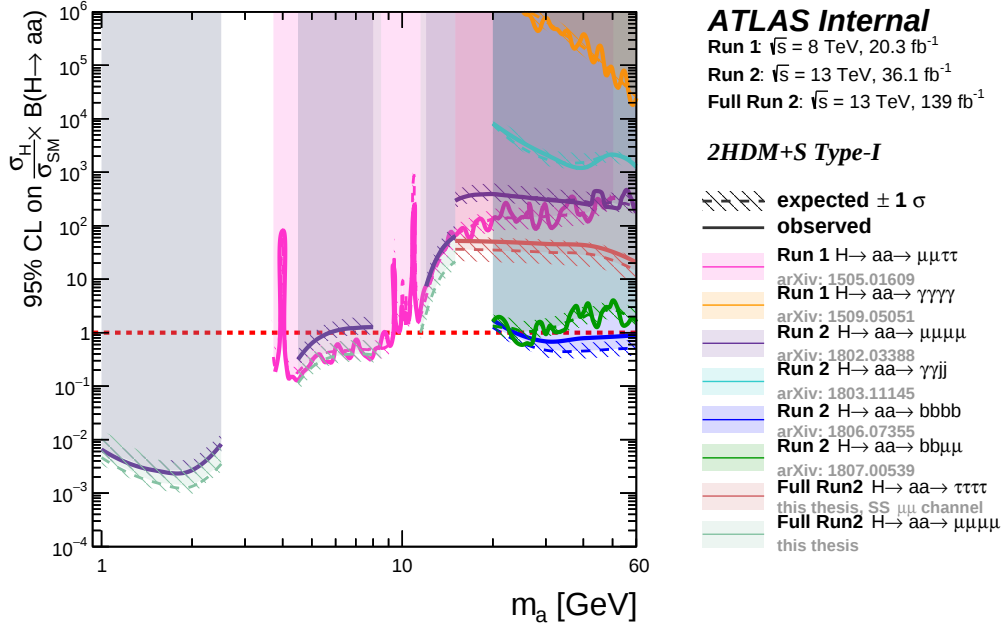


Figure 8.1: Observed and expected 95% CL upper limits on $\frac{\sigma_H}{\sigma_{SM}} \times \text{Br}(H \rightarrow aa)$ in the 2HDM+S type-I scenario. The branching fractions of the pseudoscalar boson to SM particles are computed following the prescriptions in Reference [154].

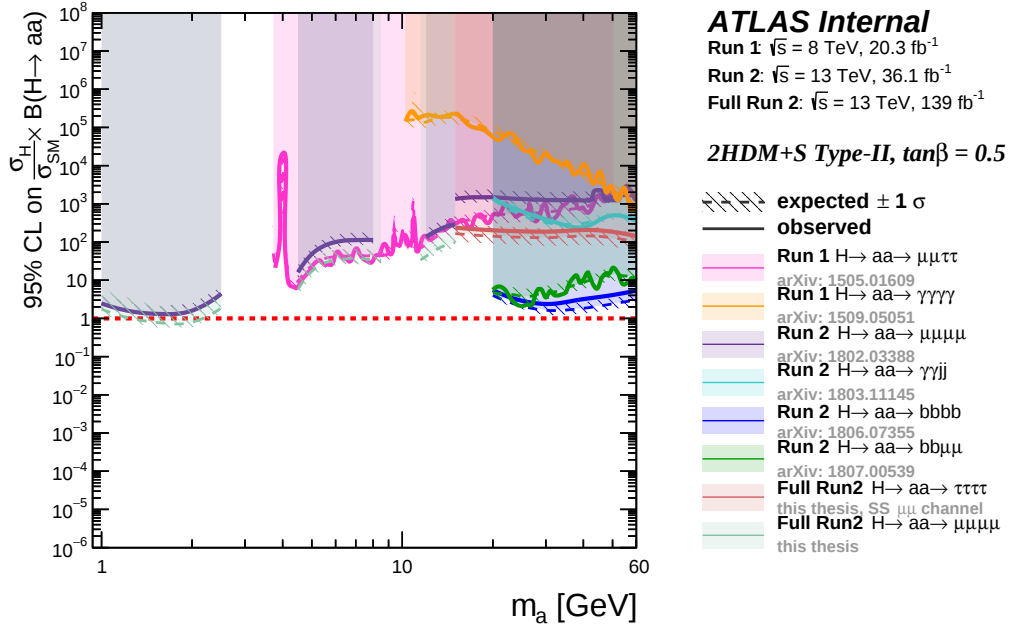


Figure 8.2: Observed and expected 95% CL upper limits on $\frac{\sigma_H}{\sigma_{SM}} \times \text{Br}(H \rightarrow aa)$ in the 2HDM+S type-II $\tan\beta = 0.5$ scenario. The branching fractions of the pseudoscalar boson to SM particles are computed following the prescriptions in Reference [154].

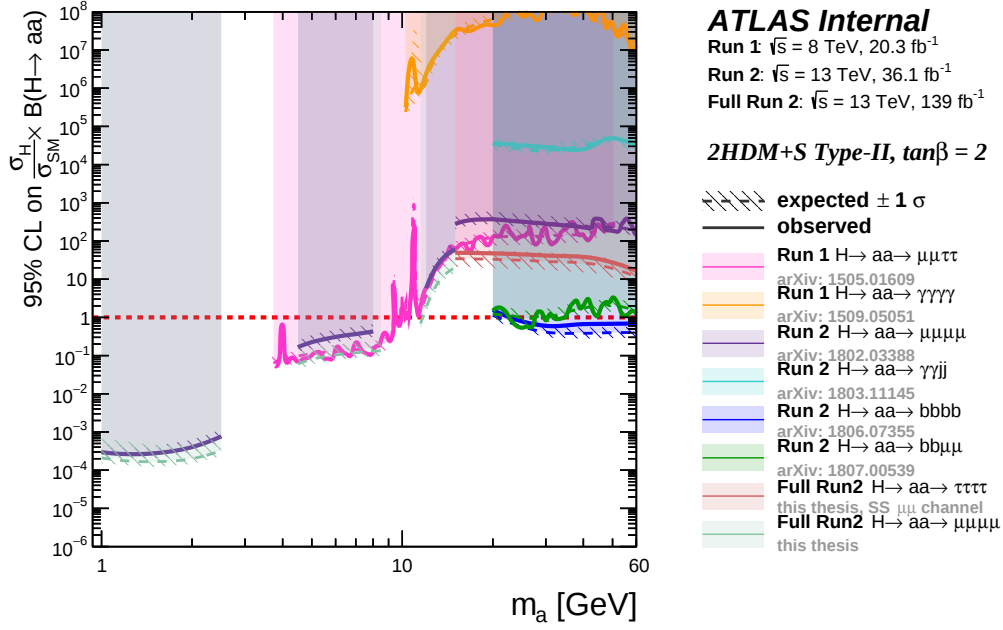


Figure 8.3: Observed and expected 95% CL upper limits on $\frac{\sigma_H}{\sigma_{SM}} \times \text{Br}(H \rightarrow aa)$ in the 2HDM+S type-II $\tan\beta = 2.0$ scenario. The branching fractions of the pseudoscalar boson to SM particles are computed following the prescriptions in Reference [154].

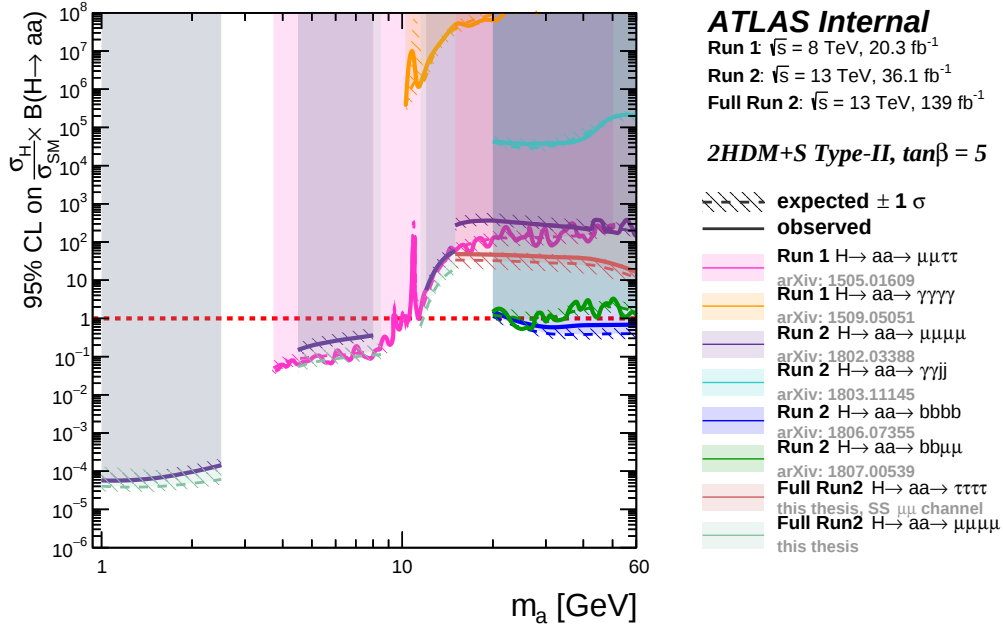


Figure 8.4: Observed and expected 95% CL upper limits on $\frac{\sigma_H}{\sigma_{SM}} \times \text{Br}(H \rightarrow aa)$ in the 2HDM+S type-II $\tan\beta = 5.0$ scenario. The branching fractions of the pseudoscalar boson to SM particles are computed following the prescriptions in Reference [154].

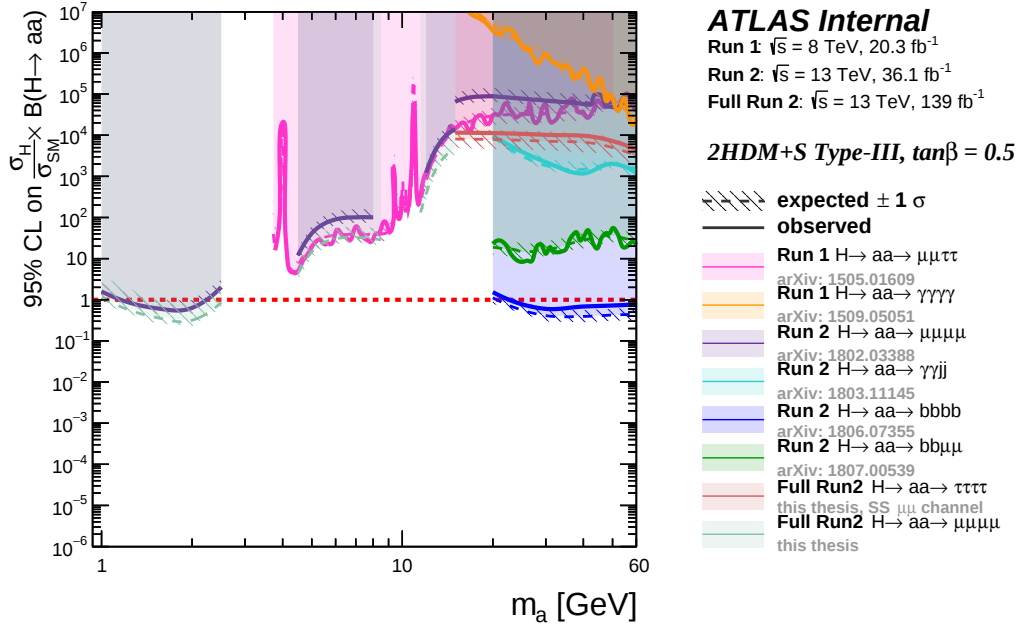


Figure 8.5: Observed and expected 95% CL upper limits on $\frac{\sigma_H}{\sigma_{SM}} \times \text{Br}(H \rightarrow aa)$ in the 2HDM+S type-III $\tan\beta = 0.5$ scenario. The branching fractions of the pseudoscalar boson to SM particles are computed following the prescriptions in Reference [154].

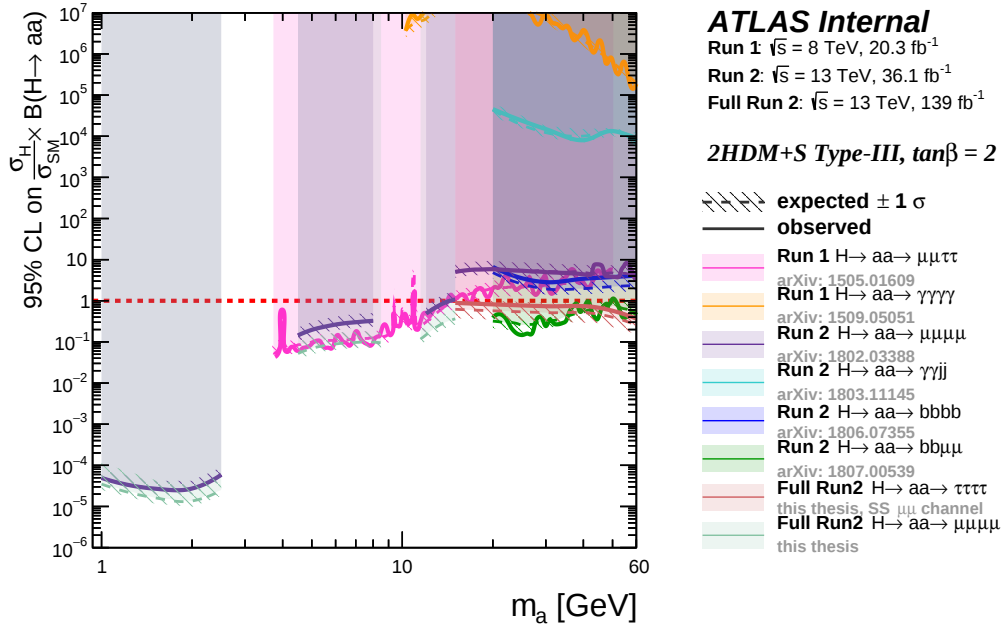


Figure 8.6: Observed and expected 95% CL upper limits on $\frac{\sigma_H}{\sigma_{SM}} \times \text{Br}(H \rightarrow aa)$ in the 2HDM+S type-III $\tan\beta = 2.0$ scenario. The branching fractions of the pseudoscalar boson to SM particles are computed following the prescriptions in Reference [154].

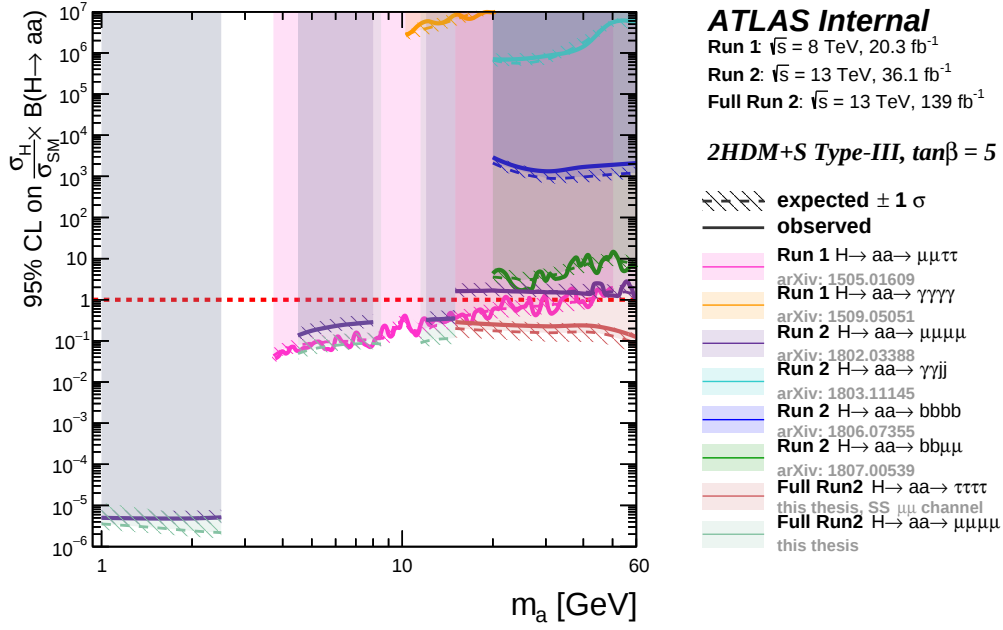


Figure 8.7: Observed and expected 95% CL upper limits on $\frac{\sigma_H}{\sigma_{SM}} \times \text{Br}(H \rightarrow aa)$ in the 2HDM+S type-III $\tan\beta = 5.0$ scenario. The branching fractions of the pseudoscalar boson to SM particles are computed following the prescriptions in Reference [154].

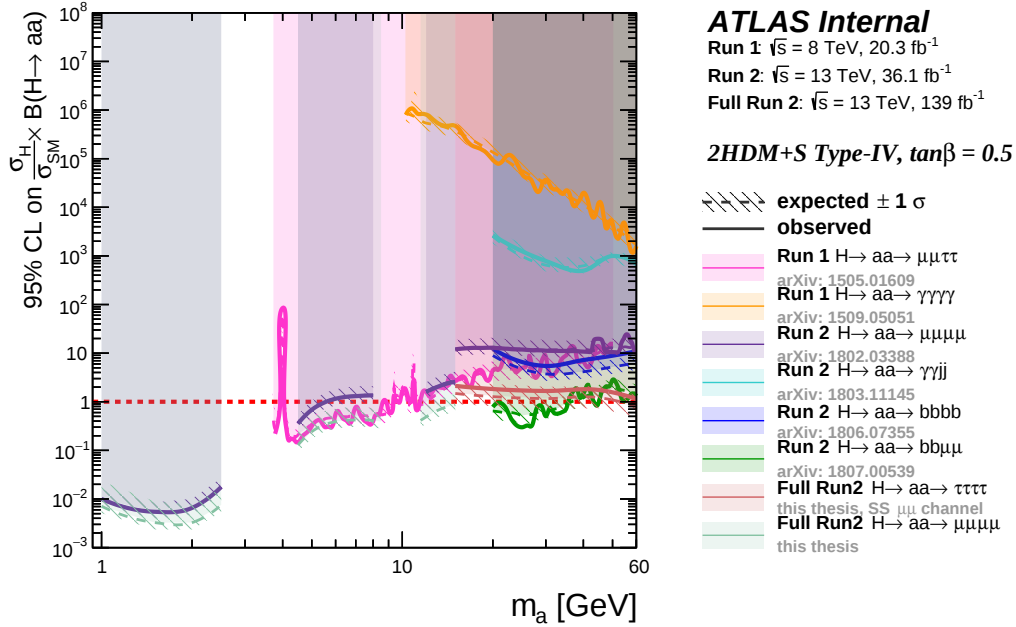


Figure 8.8: Observed and expected 95% CL upper limits on $\frac{\sigma_H}{\sigma_{SM}} \times \text{Br}(H \rightarrow aa)$ in the 2HDM+S type-IV $\tan\beta = 0.5$ scenario. The branching fractions of the pseudoscalar boson to SM particles are computed following the prescriptions in Reference [154].

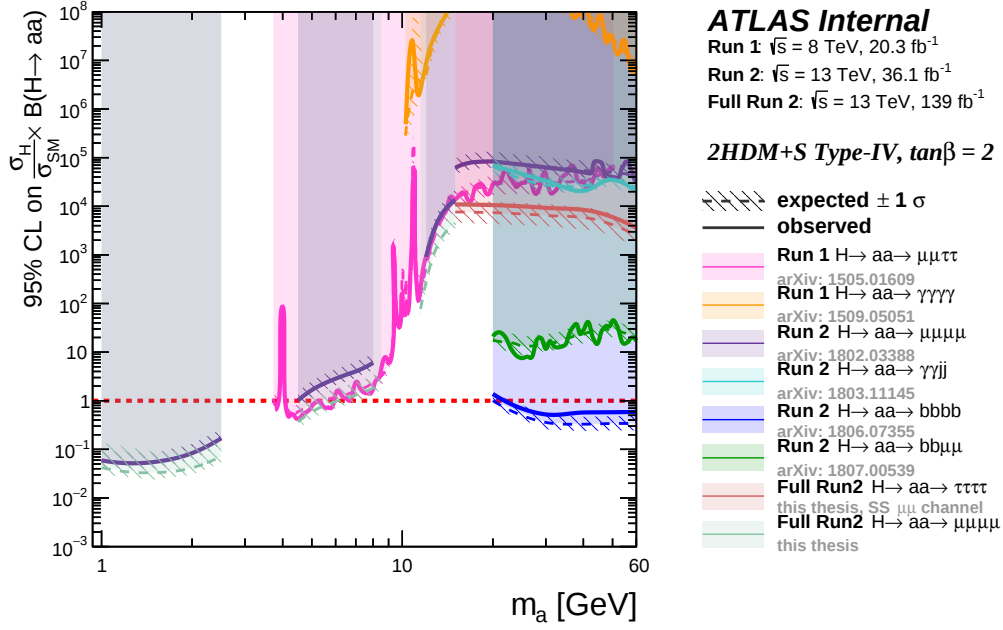


Figure 8.9: Observed and expected 95% CL upper limits on $\frac{\sigma_H}{\sigma_{SM}} \times \text{Br}(H \rightarrow \gamma\gamma)$ in the 2HDM+S type-IV $\tan\beta = 2.0$ scenario. The branching fractions of the pseudoscalar boson to SM particles are computed following the prescriptions in Reference [154].

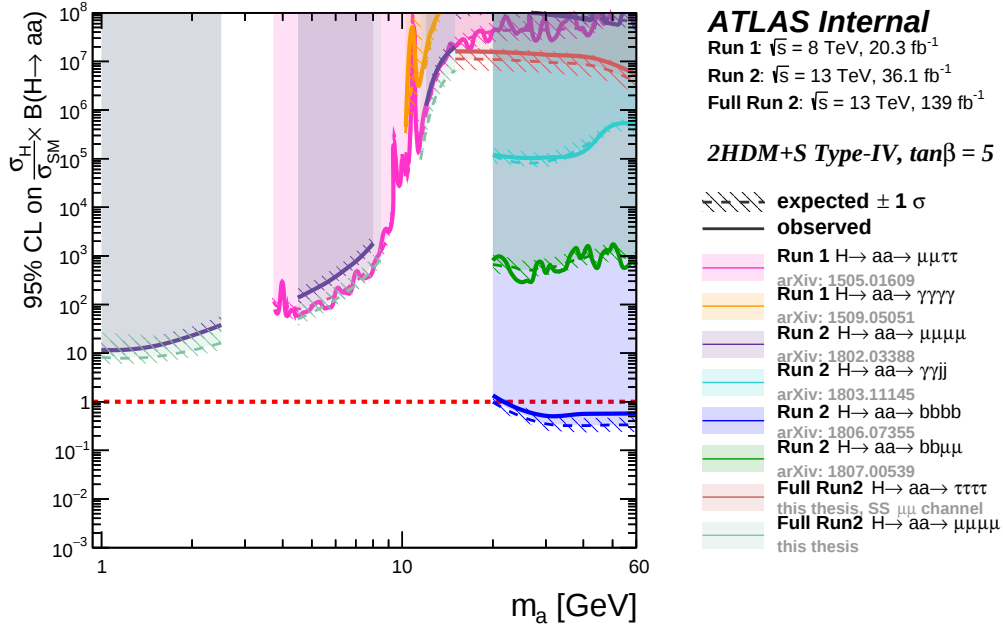


Figure 8.10: Observed and expected 95% CL upper limits on $\frac{\sigma_H}{\sigma_{SM}} \times \text{Br}(H \rightarrow \gamma\gamma)$ in the 2HDM+S type-IV $\tan\beta = 5.0$ scenario. The branching fractions of the pseudoscalar boson to SM particles are computed following the prescriptions in Reference [154].

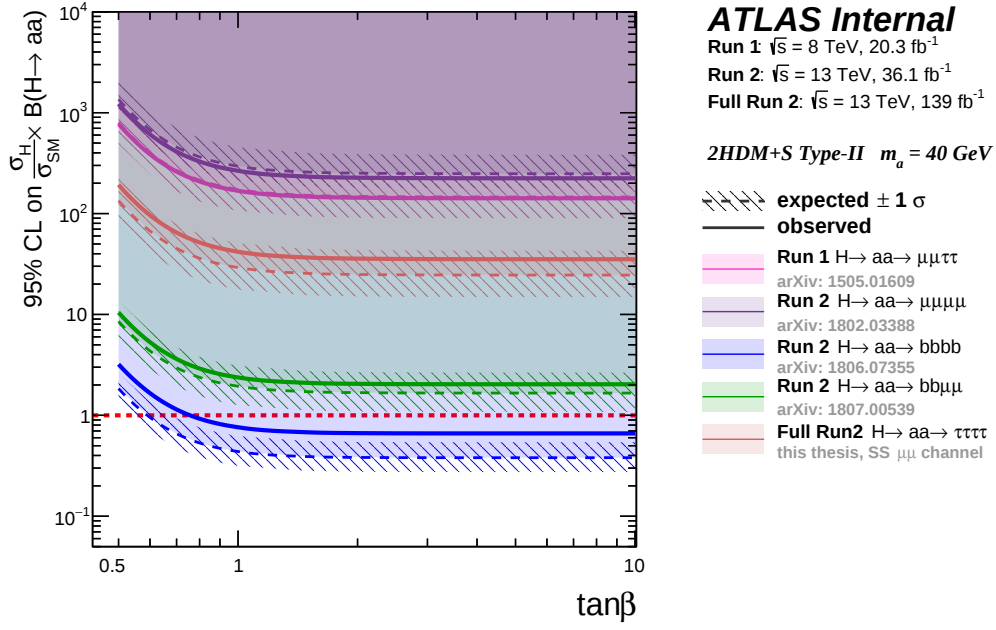


Figure 8.11: Observed and expected 95% CL upper limits on $\frac{\sigma_H}{\sigma_{SM}} \times \text{Br}(H \rightarrow aa)$ in the 2HDM+S type-II scenario for different $\tan \beta$ values for a fixed pseudoscalar mass $m_a = 40$ GeV. The branching fractions of the pseudoscalar boson to SM particles are computed following the prescriptions in Reference [154].

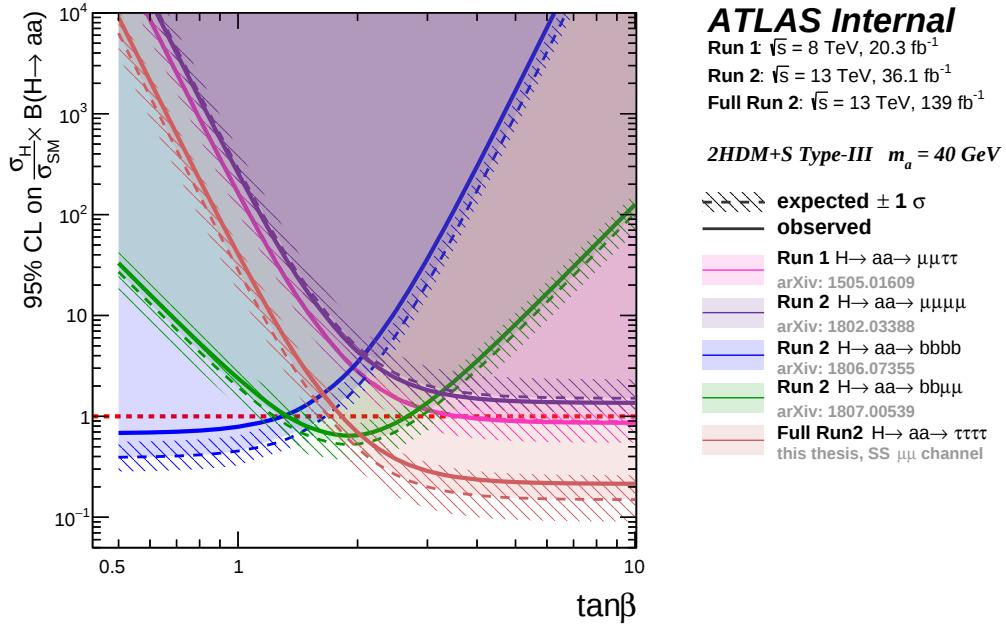


Figure 8.12: Observed and expected 95% CL upper limits on $\frac{\sigma_H}{\sigma_{SM}} \times \text{Br}(H \rightarrow aa)$ in the 2HDM+S type-III scenario for different $\tan \beta$ values for a fixed pseudoscalar mass $m_a = 40$ GeV. The branching fractions of the pseudoscalar boson to SM particles are computed following the prescriptions in Reference [154].

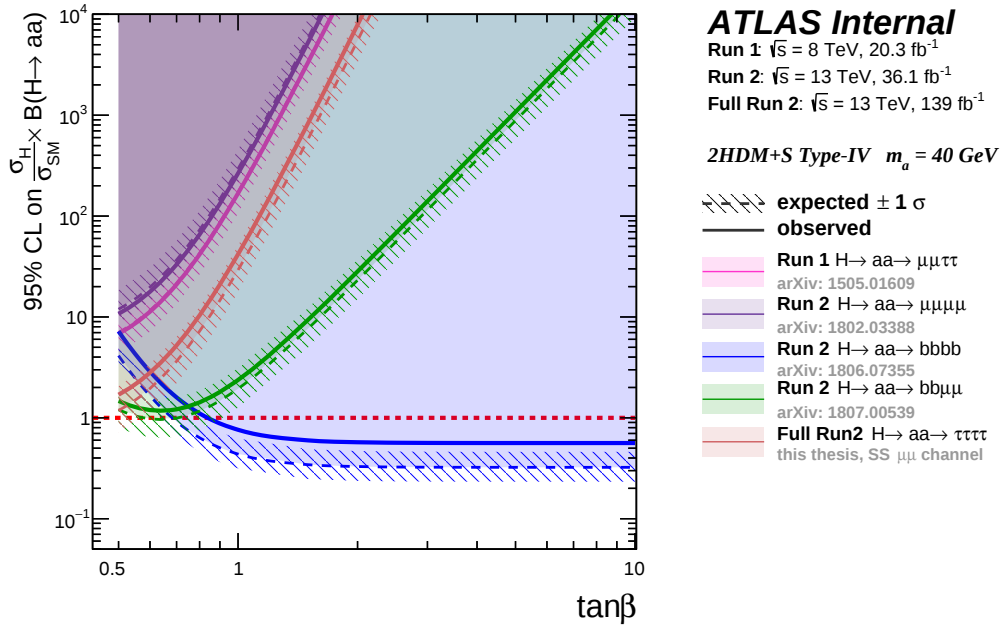


Figure 8.13: Observed and expected 95% CL upper limits on $\frac{\sigma_H}{\sigma_{SM}} \times \text{Br}(H \rightarrow aa)$ in the 2HDM+S type-IV scenario for different $\tan\beta$ values for a fixed pseudoscalar mass $m_a = 40$ GeV. The branching fractions of the pseudoscalar boson to SM particles are computed following the prescriptions in Reference [154].

Chapter 9

Conclusions

The searches for exotic Higgs decays are motivated by many unanswered questions in the [Standard Model](#). The most recent measurements can only constrain the upper limit of exotic Higgs decays to 21% at 95% [CL](#), which means a tempting potential of the discovery of new physics. This dissertation presents searches for exotic Higgs decays in two different final states, a search of $H \rightarrow 2a \rightarrow 4\mu$ aiming for $1 < m_a < 15$ GeV and a search for $H \rightarrow 2a \rightarrow 4\tau$ aiming for $15 < m_a < 60$ GeV, performed with data collected in [proton-proton](#) collisions at a [center of mass](#) energy of $\sqrt{s} = 13$ TeV by the ATLAS detector in the full run2 of the [LHC](#). No significant excess of data over the expected background has been observed. Therefore, these analyses are able to set stringent exclusion upper limits on the [branching ratio](#) of exotic Higgs decays to a pair of pseudoscalars using the [2HDM+S](#) model.

In these searches, one major challenge is to model the background from misidentified objects. Utilizing these data-driven methods can overcome the shortcomings in typical [MC](#) simulations, such as low statistics and non-reliable detector response in the simulation model, reducing systematic uncertainties, and increasing the sensitivity of the analysis. Different data-driven techniques are utilized. In the $H \rightarrow 4\mu$ analysis a dedicated template method is developed for the estimation of fake muons from heavy flavor quark decays. The $H \rightarrow 4\tau$ analysis is a [same-sign](#) search, which means low contamination from [SM](#) backgrounds. Due to complex composition of different [SM](#) fake processes, the inclusive [fake factor](#) method is implemented.

Another difficulty in exotic Higgs decay searches is that the low mass of the signal pseudoscalars will give a small angular separation between the final state particles. Therefore, the performance of standard reconstruction and identification is no longer effective. In the analysis of $H \rightarrow 2a \rightarrow 4\mu$, the calculation of muon isolation variables are optimized to remove close-by muons. The $H \rightarrow 2a \rightarrow 4\tau$ analyses targets higher mass, and a preliminary strategy targeting the low mass region is proposed using dedicated di-tau identification algorithm.

In ATLAS, many exotic Higgs decay searches have been performed looking for a signatures in a variety of final states. A summary of these analyses and their corresponding interpretations for different assumptions in the [2HDM+S](#) model is discussed. Despite the absence of a direct observation of exotic Higgs decays,

strong upper limits on $\text{Br}(H \rightarrow aa)$ are made by these analyses. Among them, the $H \rightarrow 2a \rightarrow 4\mu$ analysis gives strongest limits for $m_a < 10$ GeV, and the $H \rightarrow 2a \rightarrow 4\tau$ analysis gives the best limits on the 2HDM+S Type-III model. The efforts of ATLAS and other experiments will continue to hunt for the direct evidence of exotic Higgs decays. Since most of these analyses suffer from low statistics, the future run of the LHC and the High-Luminosity LHC will benefit searches for exotic Higgs decays. We are looking forward to future explorations!

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