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**New Physics from Warped compact Extra Dimensions:
from model building to colliders signals.**

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Resumo

No Modelo Padrão que descreve a física das partículas elementares e suas interações o campo de Higgs pode ser imaginado como um campo composto formado por uma força forte ainda desconhecida. Tal hipótese é bastante atrativa para completar o Modelo Padrão a altas energias. Problemas como hierarquia e naturalidade podem ser mais facilmente evitados. No contexto de uma força forte porém métodos de cálculo baseados em expansões perturbativas não têm mais validade. Uma alternativa para entender as propriedades básicas desse tipo de teorias é trabalhar em termos de teorias de gravitação com dimensões extras.

Nesta tese nos focamos no caso de uma dimensão espacial extra. Características genéricas desse tipo de cenário são a existência de partículas de gravidade massivas, associadas com a métrica cinco-dimensional que acopla com o Modelo Padrão para matéria, levando a assinaturas diretas em colisores de partículas (como o LHC no CERN). Tais partículas de gravidade se acoplam com o setor de Higgs. A descoberta do bóson de Higgs abriu um novo campo de investigação para sua detecção direta, no estado final com dois bósons de Higgs.

Nos usamos técnicas de Monte Carlo para estudar as estratégias de análise que levariam a um melhor reconhecimento de novas ressonâncias que decaem em pares de bósons de Higgs em colisores hadrônicos, que podem ser interpretadas como partículas de gravidade massivas. Finalmente apresentamos as buscas experimentais por tais ressonâncias realizadas no contexto do experimento CMS com dados retirados do segundo run do LHC (com uma energia de centro de massa de 8 TeV).

Palavras chave: Modelo padrão, Setor eletrofraco, Gravidade quântica.

Áreas do conhecimento: Física de altas energias, Física de partículas.

Abstract

The Higgs field of the Standard Model theory for elementary particles and interactions can be realized as a composite state from an underlying strong sector. Such hypothesis is very attractive as an ultraviolet completion of the Standard Model since it solves the hierarchy and avoids naturalness problems. The standard perturbative methods cannot be used in the context of strongly interacting theories, however those can be broadly described in terms of extra dimensional models of gravity.

We focus on the case of one additional Warped compact Extra Dimension (WED). The generic signatures of this scenario are the manifestation of heavy gravity particles, associated with the five dimensional metric, that couples with the Standard Model matter leading to direct collider signatures.

The heavy gravity particles couples to the Higgs sector. The higgs discovery had opened a new investigation channel to LHC direct detection that is the di-higgs final state.

We use Monte Carlo techniques to study the analysis strategies that would lead to a best recognition of new resonances decaying to a pair of higgses in hadron colliders, that can be interpreted as the gravity particles. We finally present resonance searches performed with data taken by the CMS experiment on the 8 TeV LHC run. The results are interpreted as the gravity particles signatures in the WED context.

Key words: Standard Model, Electroweak sector, Quantum gravity.

Áreas do conhecimento: High energy physics, Particle physics.

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Note: I swear I wrote the acknowledgements before to have the announcement of the members of the jury!

Contents

Introduction	2
I Warped compact Extra Dimensional models	9
I.1 The Higgs mechanism	15
I.2 The Standard Model fields	16
I.3 Gravity particles	21
I.4 Gravity particles couplings to matter	30
I.5 Phenomenology of gravity particles at a proton collider	35
II Phenomenology of resonant double higgs production at LHC	44
II.1 The LHC and its general purpose experiments	47
II.2 Jets	51
II.3 Topology of processes like $X \rightarrow Y Y \rightarrow z z z z$	53
II.4 $pp \rightarrow X \rightarrow h(b\bar{b})h(b\bar{b})$	55
II.5 $pp \rightarrow X jj \rightarrow h(b\bar{b})h(b\bar{b}) jj$	67
III Searches for resonant double higgs production on CMS detector	74
III.1 The CMS detector	77
III.2 $pp \rightarrow X \rightarrow h(\gamma\gamma)h(b\bar{b})$ at CMS	83
III.3 $pp \rightarrow X \rightarrow h(b\bar{b})h(b\bar{b})$ at CMS	106
III.4 Interpretation of the di-higgs LHC results	112
Conclusions	116
Appendices	119
A Signals of strongly coupled new physics in theories with a light higgs boson	120
A.1 About perturbative unitarity in a higgs-less scenario	123
A.2 Partial unitarization of $2 \rightarrow n$ ($n > 2$) processes	124
A.3 Cross section sensitivity	126
A.4 Probing anomalous couplings at LHC in WBF di-higgs production	130
B Bulk Standard Model fields	131
B.1 Bulk gauge bosons	131
C Effects of higgs-radion mixing in LHC production	134
D Counting experiment	136
E Classification for the limits	137
References	139

Introduction

The Standard Model of particle physics (SM) is the unified description of electromagnetic and weak forces by means of the $SU(2)_L \otimes U(1)_Y$ local gauge symmetry. The spontaneous breaking of this symmetry leads to successful predictions, like the existence of a neutral weak current mediated by the Z boson [1]. The structure of the Electroweak Symmetry Breaking (EWSB) in the SM is very precisely tested [2, 3].

The mechanism necessary to trigger EWSB in the SM (known as Higgs mechanism) [4] requires as ingredient the existence of a scalar field (the so called Higgs doublet) and predicts the existence of a physical scalar (h) particle whose mass (m_h) lies in the same scale than the carriers of the weak interactions (m_Z and m_W) [5]. The five known elementary interactions of nature are the Higgs force, the electromagnetic force, the weak force (responsible by radioactivity), and the strong force (that binds quarks together).

After decades of theoretical and experimental successes the SM experienced its zenith with the discovery of the higgs boson - its ultimate prediction - by the CMS and ATLAS collaborations [6, 7].

Open questions on top of the SM however remain: The Higgs mechanism is a description of spontaneous symmetry breaking rather than an understanding of its dynamics. The description of the higgs particle as a fundamental scalar with mass at the electroweak scale poses naturalness problems.

The quantum effects induced by the higgs potential and the higgs couplings to top quarks introduces quadratic divergences to m_h value, pushing it to very high values ($\sim$ to the cutoff of the theory). To stabilize the higgs mass at ~ 125 GeV a very fine tuning between SM and new physics corrections expected to appear in the Planck scale* $\sim (10^{19})$ is required. What is known as the hierarchy problem. This fine tuning suggests that new physics beyond the Standard Model should appear before the Planck scale.

The underlying Beyond Standard Model (BSM) physics can interact with the SM particles either weakly or strongly. Examples of weakly coupled BSM theories are the Supersymmetric extensions of the SM in which a new set of particles are introduced to solve the hierarchy and naturalness problem.

In the case in which the SM is completed by a strongly interacting theory the Higgs doublet can be realized as a composite state formed by the components of this strong sector, like the π mesons in Quantum Chromodynamics (QCD) are composite states of light quarks. The failure of the perturbative methods in strongly interacting theories limit their predictability. The low energy effects of the underlying strong sector however can be exploited by means of effective constructions of particle dynamics [8].

The Large Hadron Collider (LHC), the same machine that allowed the higgs boson discovery by the CMS and ATLAS collaborations will be active for at least one more decade with the purpose of revealing the nature of EWSB. The discovery potential for this machine need to be profited at maximum. From a practical point of view New Physics Beyond the Standard Model (BSM) can manifest itself at colliders in two ways:

*The scale where quantum gravity effects becomes strong.

Direct detection: BSM scenarios generically predict the existence of additional resonances with mass within the collider reach that can be directly detected. The additional particles typically are connected with the electroweak sector of the SM coupling mostly with the Electroweak bosons (W, Z and higgs) and top-quarks.

Indirect detection: The couplings between the Standard Model particles can be modified by the presence of new forces or heavy states. This kind of effects are present for example in models where the Higgs is interpreted as a composite state of some strong dynamics. The deviations of the SM couplings that could arise from the simpler composite Higgs case that afford a light higgs boson are already very constrained [9, 10].

In this thesis we will be interested on direct detection of BSM particles that can arise from strongly interacting BSM theories. The predictions of new particles expected from strongly coupled theories however are parallel to the predictions of constructions of weakly coupled theories propagating in compact extra spacial dimensions [11, 12]. The references [13, 14] worked out examples where the special case where the chiral lagrangians (that commonly appear in theories where Higgs is a light pseudo goldstone boson of a strong sector global symmetry) is successfully described by considering a theory where gravity propagates in a Warped compact Extra Dimension. The approach we adopt in this thesis is to study the direct physics predictions from specific BSM constructions based on the existence of a Warped compact Extra Dimension (WED) [15]. The knowledge of the higgs boson properties is used as a tool for the searches. We use this setup to discuss benchmark points for direct searches for spin-0 and spin-2 resonances at LHC.

This document is organized in three chapters that can be read in an independent way. These chapters are inter-connected, each one acting as one ingredient to the others. The class of models based on WED is presented in chapter (I). In this chapter we mention how this construction solves the hierarchy problem and also comment on the behaviour of matter fields on the WED background. We describe how a quantum description of the five dimensional space-time metric would manifest itself in our world as a tower of massive states with spin-0 and spin-2. We discuss the model parameters and mention the validity of the model predictions. We study the hypothesis behind the coupling of those particles with SM matter and calculate its decays and production on LHC.

In chapter (II) we study the LHC sensitivity to direct detection of a resonance on the di-higgs channel using as benchmark the gravity particles introduced on first chapter. Both inclusive production and production by weak boson fusion are explored. We study the case where both higgses decay to a pair of b -quarks, we exploit the resonance mass hypothesis that defines the regimes where different search techniques should be used.

We predict the LHC sensitivity for the detection of a di-higgs resonance produced inclusively by emulating experimental quantities like b -tag and pileup, using conservative estimations based in the LHC running conditions and performances of CMS/Atlas detectors during the Run I (2010-2012). In the study of the resonant di-higgs production of a BSM resonance produced by weak boson fusion we intend also to understand better the experimental challenges in the search for new physics in the LHC environment projected

for the next decade. This study is in progress, we show however the motivations and preliminary results.

Chapter (III) contains the discussions on the first experimental searches considering higgs bosons as final state of BSM physics performed within CMS experiment. We look for an inclusively produced resonance that decays into a di-higgs pair. We describe the basic experimental elements important to perform the searches for a di-higgs resonance in the four b -quarks final state in different mass ranges. Finally, we present the CMS search for a di-higgs resonance in the two photons and two b -quarks final state, we compare the sensitivity reached by the experiment with the predictions of the models described on the first chapter.

For the good flow of information we decided to add to the thesis our most recent works and the works in progress. In appendix (A) we study indirect BSM signatures related with modifications of the Higgs sector that are testable at LHC and are not yet very constrained. We quickly discuss the minimal model for the low effective description of a strongly interaction theory and the anomalous couplings induced by this construction in the effective Higgs sector. We study the collider effects of the anomalous couplings on the rates for multi-boson (longitudinally polarized weak bosons and/or higgs bosons) LHC production. We also present the preliminary studies we had performed on the effect of these anomalous couplings on di-higgs LHC production in the weak boson fusion mode.

Note about nomenclature: To avoid confusion we write Higgs with capital letter to refer to the mechanism and doublet while higgs with non capital letter stands for the neutral boson. Following similar logic we write Graviton (Radion) to refer to the five dimensional field, while graviton (radion) refers to the effective four dimensional particle. The weak bosons W and Z are collectively denoted using the letter V .

Acronyms and jargon's

(B)SM = (Beyond the) Standard Model *of elementary particles*

NP = New Physics

WED = Warped *compact* Extra *spatial* Dimensions

RS = Randall-Sundrum = *theory described in* [16]

KK = Kaluza-Klein

EWPT = ElectroWeak Precision Tests

EMT = Energy Momentum Tensor

(N)MSSM = (Next to) Minimal Supersymmetric Standard Model

RGE = Renormalization Group Equation

NDA = Naive Dimensional Analysis

2HDM = Two Higgs Doublet Model

WBF = Weak Boson Fusion

BKG = Background

PDF = Parton Distribution Function

MC = Monte Carlo

(N)LO = (Next to) Leading Order

ISR = Initial State Radiation

FSR = Final State Radiation

CA = Cambridge-Aachen (*jet clustering algorithm*)

AkT = Anti-kT (*jet clustering algorithm*)

LHC = Large Hadron Collider

CMS = Compact Muon Solenoid

ATLAS = A Toroidal LHC ApparatuS

BDT = Boosted Decision Tree

MVA = Multi VARIative

PF = Particle Flow

RD = Run Dependent

WP = Working Point

HLT = High Level Trigger

CiC = Cuts in Categories

CSV = Combined Secondary Vertex JER = Jet Energy Scale

JES = Jet Energy Resolution

PER = Photon Energy Scale

PES = Photon Energy Resolution

Contributions made public during thesis period

1. **Hidden sector effects on double higgs production near threshold at the LHC,**
R.Rosenfeld and A.O.
Published in: Phys. Lett. B **702**, 201 (2011); **arXiv:** 1009.4497 hep-ph.
2. **Multi Higgs and Vector boson production beyond the Standard Model.**
A Belyaev, A.O. , R.Rosenfeld and M.C.Thomas,
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3. **Scale-invariant resonance tagging in multijet events and new physics in Higgs pair production.**
M.Gouzevitch, A.O., J.Rojo, R.Rosenfeld, G.Salam and V.Sanz,
Published in: JHEP **1307**, 148 (2013); **arXiv:** 1303.6636 hep-ph.
4. **Gravity particles from Warped Extra Dimensions, a review.**
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A.O.
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5. **LES HOUCHES 2013: PHYSICS AT TEV COLLIDERS**
NEW PHYSICS WORKING GROUP REPORT. Two contributions:
 - **Resonant new physics in vector boson fusion double Higgs production at the LHC.**
A. Belyaev, A.O., O. Bondu, A. Massironi, V. Sanz, R. Rosenfeld
 - **Non-resonant new physics in vector boson fusion double Higgs production at the LHC,**
R. Contino, A.O., O. Bondu, A. Massironi, J. Rojo, M. Gouzevich.**arXiv:** 1405.1617 hep-ph.
6. **Search for new diboson resonances in semi-leptonic final states with boosted topology.**
The CMS collaboration.
cds entry: <https://cds.cern.ch/record/1696467>
Submitted to JHEP; **arXiv:** 1405.3447 hep-exp.
7. **Search for the resonant production of two Higgs bosons in the final state with two photons and two bottom quarks**
The CMS collaboration.
cds entry: <http://cds.cern.ch/record/1697512>

8. **Search for heavy resonances in the W/Z-tagged dijet mass spectrum in pp collisions at 8 TeV**

The CMS collaboration.

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Projects motivated during thesis period

In (approximate) order of current progress:

1. **Resonant new physics in vector boson fusion double Higgs production in HL-LHC enviroment.**
A. Belyaev, A.O., O. Bondu, V. Sanz, R. Rosenfeld, A. Santos.
2. **Search for BSM non-resonant production of two Higgs bosons in the final state with two photons and two bottom quarks in 8 TeV data.**
within the CMS collaboration.
3. **Gravity particles from Warped Extra Dimensions, a review.**
Part II - radion.
A.O.
4. **Non-resonant new physics in vector boson fusion double Higgs production in HL-LHC enviroment.**
O. Bondu, M. Gouzevich, A.O.
5. **Resonant new physics in WW production in vector boson fusion production mode HL-LHC enviroment.**
P. Govoni, L. Brianza, S. Pigazzini, A. Santos, A.O.
6. **Resonant new physics in vector boson fusion double weak boson production at the LHC.**
within the CMS collaboration.
7. **Scale-invariant resonance tagging in multijet events and new physics in semi-hadronic Higgs pair production as a tool to increase the range of the collider searches.**
A.O.
8. **Beyond Standard Model axial and vectorial spin-1 resonances from minimal composite model.**
C. Savoy, A.O., O. Eboli, E. Bertuzzo.
9. **Minimal realization of the higgs doublet as a composite state of KK fermions.**
E. Ponton, A.O, R. Rosenfeld.

I Warped compact Extra Dimensional models



Figure covering this chapter: Allusion to compact extra dimensions with a color treated photo of a spiral stair.

From the five known elementary forces of nature the electromagnetic, weak, strong and Higgs interactions are described by Standard Model of elementary particles, while gravitational interaction remains a odd character - the (existence of a) quantum description of gravitational field is still a big field of research. The scale of gravity interactions is defined by the weakness of the potential gravitational interaction between two objects M_1 and M_2 at a given distance r :

$$V(r; M_1, M_2) = \frac{1}{M_{Pl}^2} \frac{M_1 M_2}{r} \quad (\text{I.1})$$

The mass scale that weakens the gravitational interaction is the Planck mass, to simplify notation on the following developments it is usual to define the reduced Planck mass as $\bar{M}_{Pl} = M_{Pl}/\sqrt{8\pi}$ GeV, the value of reduced Planck mass is $\bar{M}_{Pl} \sim 2.4 \times 10^{18}$ GeV while the natural mass scale for the mediators of weak force is $M_{weak} \sim 100$ GeV.

In the course of last century it was realized that to assume the existence of compact extra spatial dimensions where gravitational interaction could propagate modifies the gravity behavior at short distances. Therefore M_{Pl} would not anymore be a fundamental scale of nature and the n -dimensional gravitational scale approaches the n -dimensional weak scale, justifying the weakness of gravity interactions with respect to the remaining forces of nature we experience on our four dimensional effective world [17, 18, 19, 20, 21].

In the latest part of last century the references [22, 23] studied the hypothesis of the flat compact extra dimensions only accessible by gravity interaction. Following reference [24] if one consider the existence on n flat compact extra dimensions with radius r_c and linear metric the hierarchy problem would be solved if $r_c \sim 10^{\frac{30}{n}-17}$ cm, in order to avoid bounds coming from experimental tests of gravity we need to consider $n > 1$.

A proposal of consider a Warped compact Extra Dimension (WED) with the focus of these works was to show a mechanism where the effective difference between $\bar{M}_{Pl}$ and M_{weak} is an outcome of the existence of one finite additional spatial dimension with non-trivial metric was made in [16]. In this work the extra dimension is compactified on a circle. The symmetry of the warp factor however allows the identification of opposite sides of the circle in such a way the extra dimension can be described as a line segment between two four dimensional manifolds known as branes. The five dimensional space defined between the two branes is referred as the bulk. In figure (I.1) we show an scheme of the compactification hypothesis. This theory setup is commonly referred as RS scenario.

In RS scenario the Higgs doublet is localized on the TeV brane such the warp factor dilutes the hierarchy problem. Quantum fluctuations of the space-time metric are interpreted as particles, the fluctuations around the 3+1 infinite part of metric correspond to the five dimensional Graviton field (spin-2) while the fluctuations around the finite size of the spatial extra dimension manifest itself as a spin-0 field, common called Radion [26].

The finiteness of the extra spatial dimension makes the four dimensional effective theory (our world) to contain excitation modes, acting as heavy resonances, known as

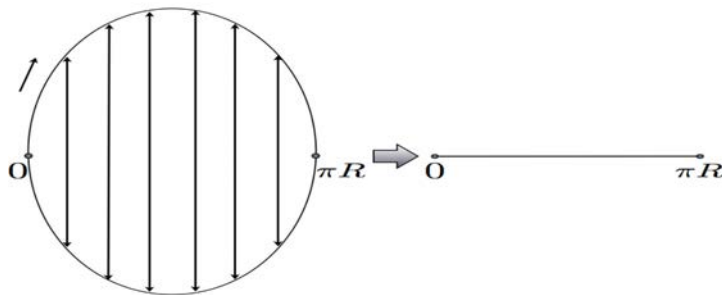


Figure I.1: Scheme for the compactification of the compact extra dimension on RS theory, each point along the circumference represents one four dimensional brane. The Planck ($\phi = 0$) and TeV ($\phi = \pi$) branes represent the four dimensional boundaries of the extra dimension ϕ compactified on a interval $\phi \in [0, \pi]$ [25]. In the figure the circumference radius is denoted by R in our text the circumference radius is denoted by r_c .

Kaluza-Klein (KK) modes*. The mass properties of the KK modes are directly calculable from the space-time geometry. The mass of the first KK mode of Graviton field (spin-2) is expected to be on TeV range if the size and curvature of the extra dimension is chosen such that the Planck-TeV hierarchy is solved on the full theory.

In the original WED proposal [16] the SM matter is localized on the same brane of the Higgs doublet. The consequences of allowing matter fields to propagate along the extra dimension under the space-time setup started to be investigated just some time after this publication. This hypothesis opened a big window for particle phenomenology: to promote the SM matter fields to be 5-dimensional introduces heavy KK modes for all the matter fields on the effective four dimensional theory [32, 33]. Indirect consequences of the above models were studied by many authors, in particular [34, 35, 36, 37]. The full classical action reads:

$$S = S_{Bulk} + S_{TeV} + S_{Planck} + S_{Matter} , \quad (\text{I.2})$$

where S_{Bulk} is the bulk gravitational action and S_{Matter} is the action matter fields, the pieces of gravitational action confined on the branes are denoted $S_{TeV/Planck}$.

We use the indices $M, N = 1, \dots, 5$ with the sign convention $(-, +, +, +, +)$ for the space-time directions. If we denote the five dimensional space-time metric as g_{XY} its determinant will be denoted by $g = -\det[g]$ and therefore the unit of five dimensional volume reads $(-\int d^5x \sqrt{g})$.

The gravitational part of the action can be written as the brane terms:

$$S_{TeV/Planck} = - \int d^4x \sqrt{g(\phi = 0, \pi)} \Lambda_{TeV/Planck} , \quad (\text{I.3a})$$

*WED constructions can be also realized as a product of finite number of copies of space-time gauge with hierarchical gauge couplings, known as sites [27, 28, 29], where a tower of excitation states is also expected. This scenario is known as *Dimensional Deconstruction* (or, by its old fashioned name *moose models* [30, 31]).

and the bulk Einstein-Hilbert action:

$$S_{Bulk} = \int d^4x \int_{-\pi}^{\pi} d\phi \sqrt{g} (-\Lambda_{bulk} + 2M_5^3 R), \quad (I.3b)$$

where the Λ 's are constant vacuum energy densities living on the branes/bulk (also known as tensions), R is the Ricci scalar, M_5 is the five dimensional Planck mass. See for example [38] and references within.

The five dimensional Einstein-Hilbert equations for gravity dynamics, derived from the action, are:

$$\sqrt{g} \left(R_{MN} - \frac{1}{2} g_{MN} R \right) = T_{MN}, \quad (I.4)$$

where the five dimensional Ricci tensor and scalar are denoted by R_{MN} and $R \equiv R^M_M$ and T_{MN} is the five dimensional energy momentum tensor of the universe. As first approximation we will neglect matter contribution to the energy momentum tensor.

The most general solution to solve the classical Einstein equations of motion maintaining four dimensional Poincaré invariance is of the form:

$$ds^2 = e^{-2\sigma(\phi)} \eta_{\mu\nu} dx^\mu dx^\nu + r_c^2 d\phi^2, \quad (I.5)$$

where the indices $\mu, \nu = 1, \dots, 4$ are used for the four dimensional directions and r_c is the radius of the extra dimension (denoted as R in figure I.1). Using the metric antzats (I.5) the Ricci tensor is given by:

$$R_{\mu\nu} = e^{-2k\sigma(\phi)} (4\sigma(\phi)'^2 + \sigma(\phi)'') \eta_{\mu\nu}, \quad (I.6a)$$

$$R_{55} = -4\sigma(\phi)'^2 - 4\sigma(\phi)'' . \quad (I.6b)$$

The primes denotes derivatives with respect to the ϕ variable. We neglect the matter back-reaction on the Einstein gravity equations and assume $\Lambda_{bulk} = \Lambda_{Planck} = -\Lambda_{TeV} = \Lambda$ resulting that [16]:

$$\sigma(\phi) = r_c |\phi| \sqrt{\frac{-\Lambda}{24M_5^2}} \equiv r_c |\phi| k. \quad (I.7)$$

The function $\sigma(\phi)$ is linear on the extra dimension and acts as a exponential warp factor along the compact extra dimension. Literature refers to the factor $k \equiv \sqrt{\frac{-\Lambda}{24M_5^2}}$ as the curvature factor. This form strongly depends on the relation between the negative cosmological constant and the other tensions.

Integrating out the extra dimension we find the four dimensional effective gravity action, and consequently the reduced four dimensional Planck mass:

$$\overline{M}_{Pl}^2 = 2M_5^3 r_c \int_0^\pi d\phi e^{-2\sigma(\phi)} = \frac{M_5^3}{k} (1 - e^{-2\pi k r_c}) \quad (I.8)$$

Usually the curvature factor k is assumed to be of the order of the five dimensional Planck scale (M_5). In this way any value of $kr_c > 1$ does not produce strong hierarchy between the mass parameters of the theory (k, M_5 and $\bar{M}_{Pl}$).

Both k and r_c are free parameters of the theory, the scale up to the gravity interaction is tested constrains the value of the compactification radius r_c . In order to avoid handle with extremely large and small numbers we use the combinations $\tilde{k} \equiv \frac{k}{\bar{M}_{Pl}}$ and kr_c as basic parameters of the theory. In the next section we discuss the relation of $r_c k$ value and the hierarchy of weak and gravitational mass scales. From now on we use the orbifold condition thus we use the ϕ variable between 0 and π and therefore suppress the modulus notation.

This chapter is divided in five sections. The first section motivates the model by showing how the hierarchy problem is sooth when considering a localized Higgs field. We them comment on the behaviors of fermions and gauge bosons in the WED background in section (I.2). In section (I.3) we derive the radion and KK-graviton physics from the metric perturbations. Section (I.4) is dedicated to the gravity particles couplings to the SM matter fields. Finally in section (I.5) we calculate the necessary ingredients to study the direct detection of those particles at proton colliders.

I.1 The Higgs mechanism

The Higgs sector is added to the theory just like the SM construction: as a $SU(2)_L \times U(1)_Y$ scalar doublet, we also admit the existence of a quartic potential.

If we confine the Higgs doublet (H) to live on the TeV brane its five dimensional action is:

$$S_H = -r_c \int d^4x d\phi \sqrt{g} [\partial_\mu H^\dagger \partial^\mu H - \lambda(|H|^2 + v_0^2)^2] \delta(\phi r_c - \pi r_c), \quad (\text{I.9})$$

when we integrate out the extra dimension (ϕ) we arrive at

$$S_H = - \int d^4x [e^{-2\pi k r_c} \partial_\mu H^\dagger \partial^\mu H - e^{-4\pi k r_c} \lambda(|H|^2 + v_0^2)^2]. \quad (\text{I.10})$$

We make the kinetic term canonical redefining $H \rightarrow e^{-\pi k r_c} H$. The 4 dimensional Higgs action reads on equation (I.11).

$$S_H = - \int d^4x [\partial_\mu H^\dagger \partial^\mu H - \lambda(|H|^2 + e^{-2\pi k r_c} v_0^2)^2]. \quad (\text{I.11})$$

We define the four dimensional vacuum expectation value of the Higgs doublet as:

$$v \equiv e^{-\pi k r_c} v_0 \quad (\text{I.12})$$

If we assume the five dimensional Higgs vev v_0 of the order of the 5D fundamental mass scale M_5 the separation between Planck (gravity) and EW (higgs mass) scale is assured when $k r_c \sim 11$. We will not address on this thesis the hypothesis where also the Higgs field is considered as a bulk field, scenario exploited for example on [39].

Other constructions admit the Higgs doublet could be a composite state of the heavy KK fermions, bounded by an additional strong force [40, 41, 42, 43], or even by excitations of QCD [44, 45]. Alternative ways to introduce of the higgs field on the WED construction also has been considered. For example, the higgs field could be originated from the components of the higher dimensional gauge fields, corresponding to the extra dimension [46, 47, 48].

The Planck-EW hierarchy reduction by the exponential space-time warp factor, represented by equation (I.12), is the most celebrated feature of Warped Extra Dimensional scenarios.

I.2 The Standard Model fields

For definitiveness we will refer to WED case where matter fields allowed to propagate on the extra dimension as bulk scenario and to the case that considers all the SM fields to be confined on the TeV brane as RS1 scenario.

The five dimensional equations of motion for bulk fields are solved performing a separation of variables on the fields five dimensional wave function that separates the ϕ and x_μ behaviors as follows:

$$X(x^\mu, \phi) = \sum_{n=0}^{\infty} X^{(n)}(x^\mu) f_X^{(n)}(\phi). \quad (\text{I.13})$$

Each four dimensional expansion mode $X^{(n)}$ is also known as the n -th KK mode while each fifth dimensional part $f_X^{(n)}(\phi)$ is commonly called as the n -th KK mode *profile*. A nice and complete guide to the behavior of bulk fields on bulk WED scenario can be found on [25].

In the simplest bulk WED construction only the SM gauge symmetry $SU(2)_L \times U(1)_Y$ is enlarged to the bulk. The five dimensional action for matter fields on the bulk model reads [33]:

$$S_{Matter} = - \int d^5x \sqrt{g} \left(\frac{1}{4} F_{MN} F^{MN} + i \bar{\Psi} \gamma^M D_M \Psi + i M_f \bar{\Psi} \Psi \right), \quad (\text{I.14})$$

where $F_{MN} = \partial_M A_N - \partial_N A_M + ig_5 [A_M, A_N]$ represents the SM gauge fields field strengths, Ψ the SM fermion fields and D_M is the covariant derivative. The bulk fermion fields admit a five dimensional mass term M_f . The masses of the KK modes of each bulk fields are geometrically generated. Usually literature calculates the masses of the KK spectrum ignoring the back-reaction of bulk matter on space-time metric*.

The masses of the zero modes however, are free of constrains. The zero modes are naturally taken as massless and correspond to the SM fields before EWSB. Deviations of the presented scenario, that holds not so strict structure for KK-particles couplings and masses can be modelled on WED construction by means of brane terms of the introduction of bulk fields with strong metric back-reaction.

In the next sub-sections we discuss the consequences of considering the gauge fields and fermion fields to be allowed to propagate in the bulk of the extra dimension ignoring any field back-reaction in the absence of brane terms.

*Some works addresses the bulk matter space-time back-reaction [49, 50, 51]. These references find that the deformation in space-time due the presence of matter is very small and therefore the changes on the mass spectra of the KK modes because of back-reaction in the metric are also small.

I.2.1 Bulk Gauge fields

When we extend a gauge symmetry to the bulk of a extra dimension, we need to consider gauge fields on all the directions where the symmetry need to be restored. If for example we assume a five dimensional abelian local symmetry, the gauge field we introduce to restore local symmetry has five components $A_M(x, \phi) = (A_\mu(x, \phi), A_5(x, \phi))$.

The physics of bulk gauge bosons is fully determined by geometry and boundary conditions. The effect of brane terms to the physics of bulk field have been investigated for example at [52, 53], It has also been argued that the effect of brane terms can be safely absorbed in the effective coupling constants [54].

In appendix (B.1) we derive carefully the solutions of the equations of motion for bulk gauge fields in the absence of brane terms. We find that when we chose boundaries conditions for a five dimensional gauge field that mimic SM pattern*, we find the gauge boson zero mode $A_\mu^{(0)}$ to have a flat profile on the extra dimension. The masses of the weak bosons will be generated by the Higgs mechanism.

I.2.2 Bulk fermions

The five dimensional mass term M_f results on one additional degree of freedom to define fermionic profiles. We find the fermionic KK modes to be vector-like, each mode have both left and right handed chirality.

The boundary conditions assure that only one of the chiralities survive for the massless zero mode. To construct a massive four dimensional field we need to introduce two bulk fermions – to fix notation we call the bulk fermion field where left (right) handed zero mode survive as left (right) handed fermion. The zero mode fermion profiles for a right/left handed fermion reads:

$$f_{L/R}^{(0)}(\phi) = \sqrt{\frac{(1 \mp 2c_{f_{L/R}})k}{e^{(1 \mp 2c_{f_{L/R}})k\phi r_c} - 1}} e^{-(1 \mp 2c_{f_{L/R}})k\phi r_c}, \quad (\text{I.15})$$

where the $c_{f_{L,R}}$ are the arbitrary constants from the separation of variables used for each bulk fermion. To simplify the analysis of the theory free parameters literature takes for the light fermions $c_{f_L} = -c_{f_R} \equiv c_f$, with this convention $c_f < 1/2$ fixes the zero mode fermion profile to be TeV localized and $c_f > 1/2$ fixes the zero mode fermion profile to be Planck localized.

The original proposals for the embedding of fermions in the WED scenario [36] considered the symmetry $SU(2)_L \times U(1)_Y$ enlarged to the bulk. The tree level masses for the four dimensional Standard Model fermions are generated under the usual EWSB mechanism. The five dimensional Yukawa interactions are introduced [33]:

$$L_{Yukawa} = \lambda_{ij}(\bar{\Psi}_L^i \Psi_R^j H + h.c.) \quad \delta(\phi r_c - \pi r_c), \quad (\text{I.16})$$

where the constants λ_{ij} are dimensionless Yukawa constants for the fermion fields i, j .

*The massless zero mode of $A_5(x, \phi)$ to vanish.

From a schematic point of view, if we take the fermions in the diagonalized mass basis the fermion masses in the bulk WED scenario are dependent on both $\lambda_f \equiv \lambda_{ij}$ (where f is the physical fermion) and fermion localization:

$$m_f^2 = \lambda_f v^2 \left(\int d\phi \quad f_{\Psi R}(c_{fL}; \phi) f_{\Psi L}(c_{fR}; \phi) \quad \delta(\phi r_c - \pi r_c) \right). \quad (\text{I.17})$$

Assuming $c_{fL} = -c_{fR} \equiv c_f$ the electroweak four dimensional fermion mass is:

$$m_f^2 \sim \lambda_f k \left(c_f - \frac{1}{2} \right) e^{(1-2c_f)\pi k r_c}. \quad (\text{I.18})$$

If we take the constant $\lambda_f k \sim 1$ the fermion masses hierarchy is created by fixing the fermion profile localization in the extra dimension. The big hierarchy between fermion masses is exponentially reduced to a hierarchy between the value of the c_f constants. For example, to reproduce the tiny electron mass (assuming all the Yukawa constants $\lambda_f k = 1$) we need to choose $c_e = 0.64$. A complete analysis of the flavor problem in the simplest version of WED bulk scenario has been done for example on [36].

I.2.3 Electroweak Precision Tests

The choice of fermion localization parameters is a more delicate issue on the third family of quarks. The t_L and b_L are zero modes of a bulk $SU(2)_L$ doublet – holding a common profile. There are three geometrical parameters controlling the third family profile localization c_{3L} (for the the left handed doublet), c_{t_R} and c_{b_R} (for the right handed singlets). To reproduce both the top and bottom quark masses (with two orders of magnitude difference) we need one of the left or right handed fermion to have TeV localized profiles.

The fermion localizations are chosen in a way to fix the four dimensional couplings to their SM value. On the other hand the effect of heavy KK modes produces dimension 6 operators. These operators contribute to SM predictions producing limits on KK modes mass scale. The pure bosonic operators contribute at tree level to the oblique parameters S and T (as defined in [55]). The fermionic operators contribute non-obliquely contributing to processes like $Z \rightarrow b_L \bar{b}_L$. The simplest solution for the third family embedding on the WED bulk on minimal gauge group setup is to choose the t_R profile to be TeV localized and the t_L profile to be Planck localized ($c_{t_R} \rightarrow 1$ and $c_{t_L} \rightarrow 0$). This scenario leads to very stringent limits on the KK fermion masses (of the order of 10 TeV), and on the KK-gauge masses (of the order of 20 TeV) [26, 35].

The authors of references [56, 57, 58] realized that enlarging the gauge symmetry in the WED bulk to a $SU(2)_L \times SU(2)_R$ custodial symmetry causes a partial cancellation of the oblique parameters loop corrections caused by the heavy KK modes, lowering down the indirect limits on the mass scale of matter KK modes.

These scenarios makes possible to fix the profiles localization differently from the simplest scenario, specially for the third family of quarks, without hurting SM precision tests.

Enlarging custodial symmetry to the bulk open the possibility for the existence of new light fermions with exotic charge [57]. The phenomenology of such exotic particles was

studied for example by [59]. In this note we are ignoring the phenomenology of any BSM particle other than the gravity particles.

In figure (I.2) we show a schematic view of the profiles of the zero modes matter fields along the extra dimension ϕ in contrast to the KK-graviton profile.

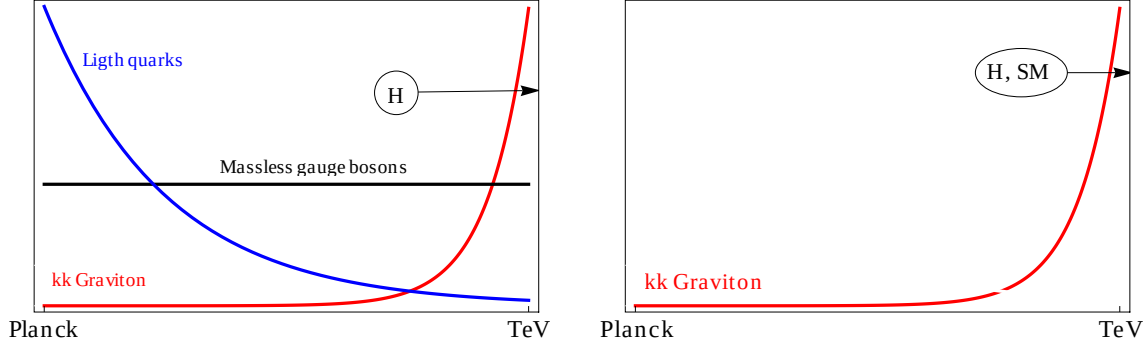


Figure I.2: Scheme of matter localization on the different WED scenarios. **Left:** Bulk scenario. **Right:** RS1 scenario. The combination of Exponential and Bessel functions makes the KK-graviton profile to be very TeV localized. To do not overload the figure we do not show the profile of zero mode Graviton, neither of the third generation of fermions.

I.2.4 Energy momentum tensor for matter fields

The Energy Momentum Tensor (EMT, also known as stress tensor) condenses information about the density and flux of energy of matter fields. This quantity had already appeared in the general equations of motion for gravity (I.4). We had neglected the contribution of matter fields to derive the dynamics of the free gravity. The couplings of the Graviton (and Radion) particles with matter however, not coincidentally, will take form by means of the EMT (and its trace). In this subsection we review the basic concepts of the physics of this object in a flat four dimensional space-time.

The EMT tensor in a space with a general background metric $g_{\mu\nu}$ is calculated from the action as:

$$T_{\mu\nu} = -\frac{2}{\sqrt{g}} \frac{\delta S_{Matter}}{\delta g^{\mu\nu}}. \quad (\text{I.19})$$

We denote $T_{\mu\nu}^i$ as the part of the four dimensional energy-momentum tensor related with the free field i and $T_{\mu}^{\mu(i)} = \eta_{\mu\nu} T^{\mu\nu(i)}$ as its trace.

Up to quadratic terms in the Higgs field the EMT of the the four dimensional Higgs doublet reads (see for example [60]):

$$T_{\mu\nu}^H = \partial_{\mu} H^{\dagger} \partial_{\nu} H - \frac{1}{2} g_{\mu\nu} (\partial_{\alpha} H^{\dagger} \partial^{\alpha} H - m_H^2 H^{\dagger} H) . \quad (\text{I.20a})$$

The Yang Mills EMT of a four dimensional gauge field A_μ and field strength $F_{\mu\nu}^A$, in unitary gauge is given by:

$$T_{\alpha\beta}^{YM}(A) = g_{\alpha\beta} \frac{1}{4} F_{\mu\nu}^A F^{A,\mu\nu} + F_\alpha^{A,\mu} F_\beta^{A,\mu}, \quad (\text{I.20b})$$

and finally,

$$T_{\mu\nu}^{\psi_{L/R}} = \frac{i}{4} (\bar{\psi} \gamma_\mu \overline{D}_\nu \psi + \bar{\psi} \gamma_\nu \overline{D}_\mu \psi) - g_{\mu\nu} \frac{i}{2} \bar{\psi} \gamma_\alpha \overline{D}^\alpha \psi, \quad (\text{I.20c})$$

is the EMT of a four dimensional free massless Dirac fermion ψ .

The space-time conserved currents in the universe can be written in terms of the EMT, the space-time conservation laws can be translated as symmetries of the EMT by means of the Noether theorem (see for example [61] for a historical introduction). In a nutshell the Noether theorem states that the conservation of energy and momentum is translated on the vanishing of EMT divergence $\partial_\mu T^{\mu\nu} = 0$. Analogously the consequence of the conservation of angular momentum in EMT is the vanishing of its rotational $\partial_\mu \wedge T^{\mu\nu} = 0$.

The last part of the Noether theorem states that classically in absence of a mass scale the divergence of the scale current $T_\nu^\mu x^\nu$ vanish. A theory with massive states is not scale invariant reflecting that the EMT for massive free fields is proportional to the field mass, the term that breaks the scale invariance.

A theory of massless fields however have no preferred mass scale classically and therefore we should expect the theory to be scale invariant. Quantum corrections to scattering amplitudes however introduces an energy dependence to the theory. This energy dependence is absorbed by the bare coupling constants by a renormalization procedure.

The renormalization of a theory consists of absorbing the divergences that appear from loop factors on the physical vertices, redefining the fields and bare coupling constants of the theory while counter terms are added to the bare Lagrangian. This procedure is systematized by the Renormalization Group Equations (RGE) formalism, a pedagogical introduction can be found in the textbooks [62, 63]. In the RGE formalism the running of the coupling constant g of a given interaction i with the energy scale μ is captured by the function $\beta_i \equiv \mu \frac{\partial g}{\partial \mu}$. The renormalized Lagrangian contains counter terms proportional to the loop induced variation of the coupling δg multiplied by the bare Lagrangian. When calculating the EMT by equation (I.19) with respect to the renormalized action the counter term related with the energy flow of the coupling will contribute to the EMT through β_i function. The anomalous scale dependence of the theory predictions are reflected in non vanishing of the trace of the EMT.

In the massless QCD and QED the running of the α_{QCD} and α_{QED} couplings contribute to their β functions. The trace of the EMT of free massless photon and gluon fields is [62]:

$$T_\mu^{\mu(g,\gamma)} = \frac{\alpha_s}{2\pi} \beta_{QCD} \sum_{colors} F_{\mu\nu}^{(g)} F^{(g),\mu\nu} + \frac{\alpha_{QED}}{2\pi} \beta_{QED} F_{\mu\nu}^{(\gamma)} F^{(\gamma),\mu\nu}. \quad (\text{I.21})$$

I.3 Gravity particles

We refer as *gravity particles* the particles resulting from the quantum fluctuations around the classical metric solution (I.7). The fluctuation modes can be decomposed into a 4D tensor $\otimes$ 4D vector $\otimes$ scalar components:

$$\delta g_{MN}(x, \phi) \propto \left(\begin{array}{c|c} h_{\mu\nu}(x, \phi) & h_{\mu,5}(x, \phi) \\ \hline h_{\mu,5}(x, \phi) & h_{55}(x, \phi) \end{array} \right) \quad (\text{I.22})$$

The equations of motion for the independent fields of (I.22) form a over constrained system. The axial gauge, defined by fixing $h_{\mu 5} = 0$, decouples the dynamics of tensor and scalar perturbations. The perturbed space-time metric in this gauge, defined with the space-time element $ds^2 \equiv g_{\mu\nu} dx^\mu dx^\nu + g_{55} d\phi^2$, is written as [64, 65, 66]:

$$ds^2 = e^{-2(\epsilon_r h_{55} + \sigma(\phi, r_c))} (\eta_{\mu\nu} + \epsilon_g h_{\mu\nu}(x, \phi)) dx^\mu dx^\nu + r_c^2 (2\epsilon_r h_{55}(x, \phi) + 1)^2 d\phi^2. \quad (\text{I.23})$$

To hold control of the order of the expansion we include the expansion parameters ϵ_g and ϵ_r on the perturbed metric. To match the canonical dimensions of the four dimensional graviton and radion the fields $h_{\mu\nu}$ and h_{55} have the canonical dimensions respectively $[mass]^{3/2}$ and $[mass]$ and therefore the ϵ_g and ϵ_r have units of $[mass]^{-3/2}$ and $[mass]^{-1}$.

The tensor fluctuations ($h_{\mu\nu}$) correspond to Graviton modes and the scalar fluctuation (h_{55}) are related with the Radion mode. We will treat the Graviton and Radion fields as free fields and will also consider independently the expansion parameters ϵ_g and ϵ_r . The dynamics and interactions of the gravity particles is found substituting the form (I.23) on the full action (I.2):

$$S_{Bulk} = S_{Graviton} + S_{radion}. \quad (\text{I.24})$$

The inverse of the metric is a power series on the Graviton and Radion fields. For example, the inverted tensorial part of the metric is [67]:

$$g^{\mu\nu} = e^{2(\epsilon_r h_{55} + \sigma(\phi, r_c))} [\eta^{\mu\nu} - \epsilon_g h^{\mu\nu} + \epsilon_g^2 h^\mu_\sigma h^{\sigma\nu} + O(\epsilon_g^3 h^3)], \quad (\text{I.25})$$

therefore the presence of the element of volume $\sqrt{g}$ in the action above will make the full theory also be an infinite series on powers of the Graviton and Radion fields. Therefore gravity is not a renormalizable field theory and we might worry about the validity of the perturbative treatment of Graviton modes as quantum fields.

In a complete theory higher curvature terms (proportional to the five dimensional Ricci scalar R) can be induced in the gravity action [16, 35]. *Naive Dimensional Analysis* (NDA) power counting states that each higher order correction induced to the bare Lagrangian would be suppressed by the theory cutoff. See for example [68, 69, 70] for a formal definition of NDA in effective field theories.

Reference [71] calculates the NDA loop factor in D dimensions. The D dimensional loop factor produces a hierarchy between five dimensional cut-off and the Planck scale of $\Lambda_D \sim 10M_D$ when $D \in [5, 11]$.

In accordance with the NDA estimative the higher curvature terms will be proportional to R/Λ_5^2 [34]. The five dimensional Ricci scalar is where $R = 20k^2$ [16]. When we use the approximation $M_5 \sim \overline{M}_{Pl}$ we find that the radiative higher curvature corrections can be neglected in the graviton description if:

$$\tilde{k} < (\sim 5). \quad (\text{I.26})$$

This inequality however is only approximative. The NDA loop factor is not based on a explicit calculation. We do not know the underlying theory and therefore we are ignorant of the factors controlling the higher order corrections.

I.3.1 Graviton modes

To find the classical free field solutions for the Graviton modes we need again to fix a gauge, we adopt the transverse-traceless gauge ($\partial_\mu h^{\mu\nu} = h^\mu_\mu = 0$) on the graviton action (I.24). The Graviton action when expanding the metric up to the order ϵ_g power is [25]:

$$S_{Graviton} = -\frac{\epsilon_g^2}{2} M_5^3 r_c \int d^4x \int_0^\pi d\phi e^{-2kr_c\phi} (\partial_\rho h_{\mu\nu} \partial^\rho h^{\mu\nu} + e^{-2kr_c\phi} \partial_5 h_{\mu\nu} \partial^5 h^{\mu\nu}) , \quad (\text{I.27})$$

which results in the equations of motion and boundary conditions:

$$\begin{cases} \square h_{\mu\nu} + e^{2kr_c\phi} \partial_5 (e^{-4kr_c\phi} \partial_5 h_{\mu\nu}) = 0 \\ (\delta h_{\mu\nu}) \partial_5 h_{\mu\nu} |_{\phi=0,\pi} \end{cases} , \quad (\text{I.28})$$

To solve the graviton equations of motion (I.28) KK expansion on graviton modes $h_{\mu\nu}^{(n)}$, separating the behavior of each mode along the extra dimension on the functions $f_h^{(n)}(\phi)$:

$$h_{\mu\nu}(x^\mu, \phi) = \sum_{n=0}^{\infty} h_{\mu\nu}^{(n)}(x^\mu) f_h^{(n)}(\phi) . \quad (\text{I.29})$$

Using the expansion (I.29) in the equation of motion (I.28) we find the four dimensional equation of motion for each mode:

$$\square h_{\mu\nu}^{(n)} = m_{Gr}^{(n)2} h_{\mu\nu}^{(n)} , \quad (\text{I.30})$$

where the $m_{Gr}^{(n)}$ are constants that appear in the separation of variables procedure. The equation (I.30) is equivalent to the equation of motion of a four dimensional spin-2 particle of mass m_n , for each mode n . The profiles obey the differential Sturm-Liouville problem (see for example [72]):

$$-\partial_\phi (e^{-4\phi\pi k} \partial_\phi f_h^{(n)}) = m_{Gr}^{(n)2} \pi^2 e^{-2\phi r_c k} f_h^{(n)} , \quad (\text{I.31a})$$

$$(\delta f_h^{(n)}) \partial_\phi f_h^{(n)} |_{\phi=0,\pi} = 0 , \quad (\text{I.31b})$$

$$\int_0^\pi d\phi e^{-2\phi r_c k} f_h^{(n)} f_h^{(m)} = \delta_{mn} . \quad (\text{I.31c})$$

The solutions of (I.31) are defined by the boundary conditions. Equation (I.31b) is satisfied imposing both Dirichlet ($f_h^{(n)}|_{\phi=0,\pi} = 0$) or Neumann ($\partial_\phi f_h^{(n)}|_{\phi=0,\pi} = 0$) conditions. We assume the zero graviton mode to be massless $m_{Gr}^{(0)} = 0$, the general solution for the zero mode graviton profile is:

$$f_h^{(0)}(\phi) = c_0^{(0)} + c_1^{(0)} e^{4\phi r_c k}. \quad (\text{I.32})$$

The only boundary condition that allows a non-massive zero mode is the Neumann one, resulting a constant massless zero mode, using the orthonormality condition (I.31c):

$$f_h^{(0)}(\phi) = c_0^{(0)} = \sqrt{\frac{2k}{1 - e^{-2k\pi r_c}}}. \quad (\text{I.33})$$

The physical zero mode graviton arises from the canonical kinetic term on the action (I.27), the profile of the canonical field is:

$$\tilde{f}_h^{(0)}(\phi) = f_h^{(0)}(\phi) e^{-kr_c \phi} = e^{-kr_c \phi} \sqrt{\frac{2k}{1 - e^{-2k\pi r_c}}} \sim e^{-kr_c \phi} \sqrt{2k}, \quad (\text{I.34})$$

therefore the profile of the four dimensional massless carrier of gravitational force have a Planck localized profile. We might worry about the universality of graviton (zero mode) interactions with matter on bulk scenario. The couplings of the graviton with fermions (zero mode) are $\propto \int d\phi e^{kr_c \phi} f_h^{(0)} f_\psi^2 \propto \int d\phi f_\psi^2 = 1$.

The general solution for massive states ($m_n \neq 0$) is given in terms of Bessel functions:

$$f_h^{(n)}(\phi) = N_h^{(n)} e^{2\phi r_c k} \left[J_2 \left(\frac{m_{Gr}^{(n)} e^{\phi r_c k}}{k} \right) + b_h^{(n)} Y_2 \left(\frac{m_{Gr}^{(n)} e^{\phi r_c k}}{k} \right) \right]. \quad (\text{I.35})$$

The orthonormality conditions constrain the values of $N_h^{(n)}$ and $b_h^{(n)}$ of each mode. The mass of each mode (m_n) is found imposing the Neumann boundaries condition. The mass of the first graviton KK mode is:

$$m_{Gr} \equiv m_{Gr}^{(1)} = x_1 \tilde{k} e^{-kr_c \pi} \overline{M}_{Pl}, \quad (\text{I.36})$$

where $x_1 = 3.83$ is the first zero of the Bessel function J_1 . If we want to use both m_{Gr} and $\tilde{k}$ as the free parameter of theory we must note the relation:

$$kr_c = \frac{1}{\pi} \log \left(x_1 \tilde{k} \frac{\overline{M}_{Pl}}{M_{Gr}} \right). \quad (\text{I.37})$$

The value of kr_c parameter is dominated by the value of $\overline{M}_{Pl}$. The dependence on the parameters M_{Gr} and $\tilde{k}$ is very mild. Following the relation of equation (I.37) the kr_c value varies between 10 and 12 if $\tilde{k}$ vary between 0.01 and 1, and M_{Gr} vary between 100 GeV and 1.5 TeV. Therefore, assuming the relation (I.8) is valid, the value of kr_c remain an acceptable to reproduce the mass hierarchy of equation (I.12) for a wide range of $\tilde{k}$ and M_{Gr} parameters.

I.3.2 Radion modes

In this section we treat the dynamics of the scalar metric fluctuations (Radion). Analogously to section (I.3.1), we replace the form (I.23) in the the Radion action, resulting in [73]:

$$S_{Radion} = \frac{1}{2} \int d^4x d\phi \sqrt{g} \left[\partial_\mu r \partial^\mu r + \frac{2M_5^3}{k} \left(1 - \frac{k}{24M_5^3} r^2 \right) R \right], \quad (I.38)$$

where the field $r(x, \phi)$ is defined to have canonical kinetic term, this field relates to the scalar metric fluctuations $h_{55}(x, \phi)$ as:

$$r(x, \phi) = \sqrt{\frac{24M_5^3}{k}} e^{-k \epsilon_r h_{55}(x, \phi) \pi}. \quad (I.39)$$

The free radion field ($r(x, \phi)$) in this theory behaves as a massless field. This field however couples with the SM particles trough expansions of the metric on S_{Matter} action. It also interacts with photons [74], slightly modifying propagation of light. The existence of a massless scalar particle with electromagnetic is experimentally disfavored, see for example [75, 76]. Also, the distance between branes (r_c) was introduced in the un-perturbed metric (I.5) as a free parameter, even though its value is an important feature on the theory. The value of r_c is important to define the hierarchy of gravity and weak scales meaning we need one mechanism that dynamically fix its value (see equation I.12).

Reference [77] proposed that the compactification radius could be dynamically generated as vacuum expectation value of a potential, at the classical level. The potential energy would be created by the existence of a bulk scalar field with specific brane potential*. Following the approach of [77] we begin with the static metric (I.5), and introduce a bulk scalar with the action:

$$S_\Phi = \int d^4x d\phi \int_0^\pi \sqrt{g} (\partial_M \Phi \partial^M \Phi - M_\Phi^2 \Phi^2) + \quad (I.40a)$$

$$+ \int d^4x \int_0^\pi \sqrt{g} [\lambda_{TeV} (\Phi^2 - v_{TeV}^2)^2 \delta(\phi - \pi) + \lambda_{Planck} (\Phi^2 - v_{Planck}^2)^2 \delta(\phi)] . \quad (I.40b)$$

We will refer to the Φ field as "background field". The brane terms (I.40b) makes the field Φ to acquire a vacuum expectation value Φ_v , determined by minimizing the action (I.40) with respect to Φ . The Φ_v general form is:

$$\Phi_v(\phi) = e^{2\sigma(\phi)} (A e^{\nu\sigma(\phi)} + B e^{-\nu\sigma(\phi)}), \quad (I.41)$$

where $\nu = \sqrt{4 + M_\Phi^2/k^2}$ and A and B are arbitrary constants, to be found with the equations of motion.

The brane terms (I.40b) dynamically generates a potential energy that is dependent on the 5th dimension component. Integrating out the fifth dimension obtain that the potential energy depends only of the compactification radius, this potential turns out to

*This generic idea can be adapted to specific models, for example [78] suggested the Radion potential could be stabilized by the condensation of bulk super-symmetric particles.

have a complicated form:

$$V(r_c) = k(\nu + 2)A^2(e^{2\nu k r_c \pi} - 1) + k(\nu - 2)B^2(1 - e^{-2\nu k r_c \pi}) \quad (\text{I.42})$$

$$+ \lambda_{TeV}(\phi)e^{-4k\pi r_c}(\Phi_v(0)^2 - v_{TeV}^2)^2 + \lambda_{Planck}(\phi)(\Phi_v(\pi)^2 - v_{Planck}^2)^2,$$

assuming $kr_c \gg 1$, we find this potential to be minimized if:

$$r_c k = \frac{4}{\pi} \frac{k^2}{M_\Phi^2} \log \frac{v_{Planck}}{v_{TeV}}. \quad (\text{I.43})$$

This picture provides a framework where the r_c value would be dynamically generated in terms of a bulk scalar field. However it does not treat the dynamics of the radion field.

In the next steps we develop the theory for Radion dynamics, we follow the work done by [26, 79] to derive the classical dynamics of radius stabilization and by [64] to comment on the dynamics of the quantum fields. We stress the key points of the theory necessary to better understand the origin of the radion mass and potential, the intermediate steps can be found in the above references.

I.3.2.1 Background field dynamics and r_c stabilization

To consider the back-reaction of the Φ field on geometry we need to include the effects of the classical bulk scalar of (I.40) on the derivation of the metric solution (I.53b). We consider the scalar action:

$$S_{scalar} \equiv S_{Radion} + S_\Phi. \quad (\text{I.44})$$

In addition to the action (I.40) we assume a bulk background field to have a potential $U(\Phi)$ with the special form:

$$U(\Phi) = \frac{1}{8} \left(\frac{\partial W(\Phi)}{\partial \Phi} \right)^2 - \frac{1}{6M_5^3} W(\Phi)^2, \quad (\text{I.45})$$

where $W(\Phi)$ is a arbitrary real function following some normal mathematical requirements (like to be continue and finite in the ϕ interval). This function is commonly known as super-potential. This form for the potential and brane terms allows intelligent manipulation of the coupled equation of motion. We can find the solution of the equations of motion derived as function of the the $W(\Phi)$:

$$\sigma' = \frac{1}{6M_5^3} W(\Phi), \quad (\text{I.46a})$$

$$\Phi' = \frac{1}{2} \frac{\partial W(\Phi)}{\partial \Phi}, \quad (\text{I.46b})$$

where the prime denotes derivative with respect to the fifth space-time component ϕ . Once the super-potential function $W(\Phi)$ is defined, both the classical metric and background field solutions are also defined.

The classical dynamics of σ and Φ dynamically fixes the size of the extra dimension that minimizes the $V(r_c)$ potential.

In general for given $W(\Phi)$ the metric solution and stabilized r_c have to be calculated numerically. Both [79, 80] also work out examples in terms of benchmark super-potentials. We recover the free field solution of equation (I.43) for a very heavy background field ($M_\Phi \gg k$).

I.3.2.2 Quantum field fluctuations and radion mass

Together with the metric quantum fluctuations (equation I.23) we also introduce quantum fluctuations ($\tilde{\Phi}$) around the vacuum expectation value of the background field Φ_v :

$$\Phi(x, \phi) = \Phi_v(\phi) + \tilde{\Phi}(x, \phi), \quad (\text{I.47})$$

If we introduce the scalar field fluctuations on the S_{Scalar} the action minimization leads to coupled equations of motion for $h_{55}(x, \phi)$ and $\tilde{\Phi}(x, \phi)$. Massaging the equations we decouple the equation that governs the $h_{55}(x, \phi)$ dynamics:

$$e^{2\sigma} \square h_{55} = 4\sigma' \frac{\Phi_v(\phi)''}{\Phi_v(\phi)'} h_{55} - 2 \frac{\Phi_v(\phi)''}{\Phi_v(\phi)'} h_{55}' - 4\sigma'' h_{55} - 2\sigma' h_{55}' + h_{55}''. \quad (\text{I.48})$$

We stress that at this point the functional forms of $\Phi_v(\phi)$ and $\sigma(\phi)$ was already defined by the classical theory (section I.3.2.1). As usual, equation (I.48) is solved performing separation of variables, we expand the Radion field in KK modes as:

$$h_{55}(x, \phi) = \sum_{n=0}^{\infty} h_{55}^{(n)}(x^\mu) f_r^{(n)}(\phi), \quad (\text{I.49})$$

the separation of variables procedure introduces one mass eigenvalue for each mode $m_r^{(n)}$. The equations that defines each mode wave function are:

$$\square h_{55}^{(n)} = m_r^{(n)2} h_{55}^{(n)}, \quad (\text{I.50a})$$

$$4\sigma' \frac{\Phi_v(\phi)''}{\Phi_v(\phi)'} f_r^{(n)} - 2 \frac{\Phi_v(\phi)''}{\Phi_v(\phi)'} f_r^{(n)'} - 4\sigma'' f_r^{(n)} - 2\sigma' f_r^{(n)'} + f_r^{(n)''} = -e^{2\sigma} m_r^{(n)2} f_r^{(n)}. \quad (\text{I.50b})$$

The lightest mode $h_{55}^{(0)}$ is related to the four dimensional radion. The zero mode radion in the case we ignore the $\tilde{\Phi}$ back-reaction will remain massless. To better understand the role of the back-reaction of the background field on the radion mass we present in the next section a derivation of the radion mass for a specific choice of functions, that allows analytical solutions.

I.3.2.3 Explicit example

In this section we present a derivation of specific model, worked out in the references [79, 26] We chose the super-potential as:

$$W(\Phi) = \frac{6k}{\kappa^2} - u\Phi^2, \quad (\text{I.51})$$

where u and κ are arbitrary constants that define the super-potential. We also assume the brane potentials have the form:

$$\lambda_{Planck/TeV} = \pm W(\Phi_{\pm}) \pm W'(\Phi_{\pm})(\Phi - \Phi_{\pm}) + \gamma_{Planck/TeV}^2 (\Phi - \Phi_{\pm}), \quad (I.52)$$

where $\Phi_{\pm}$ are the value of the background field on Planck/TeV brane and $\gamma_{Planck/TeV}$ are arbitrary constants. The classical solutions for the background field vacuum expectation value and metric warp function are:

$$\Phi_v(\phi) = v_{Planck} e^{-ur_c \phi}, \quad (I.53a)$$

$$\sigma(\phi) = k\phi r_c + \frac{\kappa^2 v_{Planck}}{12} e^{-2ur_c \phi}. \quad (I.53b)$$

We will assume a small back-reaction of the background field on warp function, meaning that $\kappa^2 v_{Planck/TeV} \ll 1$. If we define the variable z such as $dz e^{-\sigma(y(z))} = dy$, the location of the branes on this coordinated is $z_{Planck} = 0$ and $z_{TeV} = \frac{1}{k}(e^{kr_c} - 1)$. After some manipulations equation (I.50b) is reduced to*:

$$-f_r^{(n)''} + \frac{\alpha(\alpha+1)k^2}{(kz+1)^2} f_r^{(n)} = m_r^{(n)2} f_r^{(n)}, \quad (I.54)$$

where $\alpha = -(3/2 + u/k)$ and the primes denotes derivation with respect to the z variable. The boundary conditions that define the problem on the $\kappa^2 v_{Planck/TeV} \ll 1$ approximation, and after the change of variables in the explicit example are:

$$(f_r^{(n)'} + \frac{\alpha k}{kz+1} f_r^{(n)})|_{Planck, TeV} = 0. \quad (I.55)$$

The profiles solutions are constructed with $(\alpha + 1/2)^{th}$ order Bessel and Neumann functions:

$$f_r^{(n)}(z) = a_n \left(z + \frac{1}{k}\right)^{1/2} N_{\alpha+1/2}(m_n(z + 1/k)) + b_n \left(z + \frac{1}{k}\right)^{1/2} J_{\alpha+1/2}(m_n(z + 1/k)). \quad (I.56)$$

As usual the boundaries conditions determine the separation constants a_n and b_n relation, and finally the m_n values. The orthonormality relations determine the overall profiles normalization.

The non negligible scalar back-reaction on warp metric results at $\sigma(\phi)'' \neq 0$, in this case the equation that defines the radion profiles are:

$$f_r^{(n)''} - 2\sigma' f_r^{(n)'} - 4\sigma'' f_r^{(n)} + 2u f_r^{(n)'} - 4u\sigma' f_r^{(n)} = -m_r^{(n)2} e^{2\sigma} f_r^{(n)}. \quad (I.57)$$

We treat the metric back-reaction (recognized in the term σ'') as a perturbation in the parameter $l \equiv \kappa v_{Planck}/\sqrt{2}$. Using this expansion in (I.54) and keeping terms only up to

*As poetic licence we denote the profile function on ϕ and z variables with the same symbol $f_r^{(n)}(\phi) = f_r^{(n)}(\phi(z)) \rightarrow f_r^{(n)}(z)$.

l^2 we arrive to a equation for the zero mode radion as:

$$t''_{r_0} = 2(k+u)t'_{r_0} = -\tilde{m}_r^2 e^{2k\phi r_c} - \frac{4}{3}(k-u)ue^{-2u\phi r_c}, \quad (\text{I.58})$$

where the function $t_{r_0}(\phi)$, a perturbation function with respect with the free field solution, and the perturbation parameter l are defined as:

$$f_r^{(0)}(\phi) = e^{2k\phi r_c}(1 + l^2 t_{r_0}(\phi)), \quad (\text{I.59})$$

and the constant $\tilde{m}_r^2 \equiv \frac{m_r^{(0)2}}{l^2}$. The boundaries conditions are:

$$t'_{r_0} + \frac{2}{3}ue^{-2u\phi r_c}|_{Planck/TeV} = 0. \quad (\text{I.60})$$

We find the radion mass as function of the fundamental parameters of the theory:

$$m_R^2 \equiv m_r^{(0)2} = 4l^2 \frac{(2k+u)u^2}{3k} e^{-2(u+k)r_c}. \quad (\text{I.61})$$

The exact value of the radion mass in this explicit model is very sensitive to the exact value of the free parameter u of the specific super-potential (I.51). The parameter have dimension of mass, we assume it to be at the same scale of the other mass scales of the theory ($k, \overline{M}_{Pl}$). If we vary the u parameter between $3 \times k$ and k and assume $\tilde{k} = 0.2$ the radion mass varies between $m_R^2 \sim l^2 \times 0.08 \text{ GeV}$ to $m_R^2 \sim l^2 \times 7 \times 10^{16} \text{ GeV}$, where l is the expansion parameter and therefore should be small. A more precise mass estimative is not useful without a microscopic motivation of the form of the super-potential.

The radion kinetic term is derived from the scalar Ricci term of gravitational action ($\sim \int d\phi \sqrt{g} R$). If we ignore the perturbation terms on the radion profile (the terms proportional to l^2 on I.59) we arrive to the kinetic term:

$$S_{radion}^{kinetic} = \frac{6M_5^3}{k} (e^{2k\pi r_c} - 1) \int d^4x (\partial h_{55}^{(0)}(x))^2. \quad (\text{I.62})$$

The canonically normalized radion field is:

$$r(x) \equiv \sqrt{\frac{k}{6M_5^3(e^{2k\pi r_c} - 1)}} h_{55}^{(0)}(x), \quad (\text{I.63})$$

If we use the un-perturbed relation for the reduced Planck mass (I.8) we have the relation:

$$h_{55}^{(0)}(x) = \frac{1}{\sqrt{6}} \frac{1}{e^{-k\pi r_c} \overline{M}_{Pl}} r(x). \quad (\text{I.64})$$

Regardless the functional form of the perturbation t_{r_0} the extra dimensional profile function ($f_r^{(0)}(\phi)$) is TeV localized by the factor $e^{2k(\phi+\pi)r_c}$, we can extend this conclusion to general forms of the super-potential.

I.3.3 Radion - higgs mixing

The Higgs scalar can have a non-minimal coupling to gravity. The lowest level non-minimal Higgs coupling to gravity (two derivative) is known as the scalar-curvature mixing. This term was first considered in discussions about conformal invariance on curved space-time backgrounds (see for example [81]) and also motivates cosmological models of Higgs inflation (see for example [82]). The scalar-curvature mixing reads:

$$S_\xi = \xi \int d^5x \sqrt{g} H^\dagger H R \delta(\pi - \phi) \quad (\text{I.65})$$

The action (I.65) was studied in the WED theories by [83] and nice references for the study of Higgs-curvature mixing on this context are [64, 73]. The scalar-curvature mixing can be motivated from a more fundamental level, see for example [84]. This action contains a mass mixing between the (h) and radion (r) gauge eigenstates. Also, the Lagrangian defined by this action is two derivative level and it induces a single field term on the higgs boson EMT:

$$T_\mu^{\mu(h)} = 6\xi v \square h + (6\xi - 1) \partial_\mu h \partial^\mu h + 6\xi h \square h + 2m_h^2 h^2, \quad (\text{I.66})$$

therefore the radion couplings to matter (defined by equation I.68) will contain a kinetic mixing between the higgs and radion gauge eigenstates. These gauge eigenstates are related to the the kinetic and mass diagonalized states h_ξ and r_ξ (where with $m_{h_\xi} < m_{r_\xi}$) by two angles:

$$r = (\sin \theta - \sin \rho \cos \theta) h_\xi + (\cos \theta + \sin \rho \sin \theta) r_\xi, \quad (\text{I.67a})$$

$$h = \cos \rho \cos \theta h_\xi - \cos \rho \sin \theta r_\xi. \quad (\text{I.67b})$$

Those angles are calculable from the mass physical observables, the radion scale Λ_r and the ξ parameter, the explicitly expression can be find for example at [73]. We ignore the higgs-radion mixing, otherwise explicitly stated.

I.4 Gravity particles couplings to matter

The couplings of radion and KK-graviton with matter fields arise from expanding the metric perturbations (equation I.23) in the action S_{Matter} . The first order set of interactions of the radion with SM matter have the form:

$$\mathcal{L}_i = -\frac{r}{\Lambda_R} \times a_i T_\mu^{\mu(i)}(x) - \frac{x_1 \tilde{k}}{m_{Gr}} h^{\mu\nu(1)} \times d_i T_{\mu\nu}^i, \quad (\text{I.68})$$

where the factors d_i and a_i are defined to absorb the possible volume suppression on the radion coupling with the field i . The source of this parameters depend on the field hypothesis. Additional terms on the radion couplings with matter that are not proportional to energy momentum trace can be present on bulk scenario [85], we discuss those contributions along the following subsections.

In the free theory the parameter $\Lambda_R = \sqrt{6} \frac{m_{Gr}}{x_1 \tilde{k}}$, see equation (I.64). The exact relation between its parameter and the original parameters of the microscopic theory when the radion has a non-zero mass however relies on a sequence of approximations and therefore we chose to use the parameter Λ_R (the scale that suppresses all the radion couplings) as a free parameter of the theory.

I.4.1 KK-graviton couplings to matter

The four dimensional effective couplings of KK-graviton to matter are proportional to the overlap integral between the profiles of the involved fields $d_i \equiv \int d\phi f_{Gr} f_i^2$. The KK-graviton is extremely TeV localized. In the RS1 scenario all the d_i 's parameters are approximately one and therefore the strength of the couplings between KK-graviton and SM matter are democratic with each field degree of freedom. This feature will be more clear in figure (I.8).

In the bulk scenario therefore the parameter $d_i \leq 1$ indicates how much the profile of the field i is away from TeV brane. This parameter is also known as volume suppression. In the following we calculate the volume suppression of the bulk KK-graviton coupling to each SM field hypothesis.

Higgs doublet

The Higgs doublet is confined to TeV brane, therefore the volume suppression on its couplings to KK-graviton is proportional to the value of KK-graviton profile ($f_h^{(0)}$) on TeV brane:

$$d_H = f_h^{(0)}(\phi = \pi) \sim 1. \quad (\text{I.69})$$

The decay rate of KK-graviton to a pair of physical higgses is given by [86]:

$$\Gamma(h_{\mu\nu} \rightarrow hh) = \frac{m_{Gr}^3}{960\pi} \left(\frac{\tilde{k} x_1}{m_{Gr}} \right)^2 \left(1 - \frac{4m_H^2}{m_{Gr}^2} \right)^{5/2} \quad (\text{I.70})$$

Massless gauge bosons

As outlined in figure (I.2) the profile of the massless zero mode of the spin-1 bulk bosons is flat with respect to the metric. Therefore its four dimensional coupling with KK-graviton is suppressed by the factor:

$$d_g = \frac{2}{kr_c\pi} \frac{(1 - J_0(x_1))}{x_1^2 |J_2(x_1)|}. \quad (\text{I.71})$$

This coupling suppression holds for both gluons, photons and weak bosons on R_ξ gauge. The decay rates of KK-graviton to a pair of SM non-massive gauge bosons, photon (γ) and gluon (g), are given by [86]:

$$\Gamma(h_{\mu\nu} \rightarrow gg, \gamma\gamma) = 2d_g^2 \times N_{g,\gamma} \times \frac{m_{Gr}}{160\pi} \left(\tilde{k} x_1 \right)^2, \quad (\text{I.72a})$$

where the color factors are:

$$N_g = 8, N_\gamma = 1 \quad (\text{I.72b})$$

In figure (I.3) we show the dependence of the d_g parameter with the KK-graviton mass using the relation (I.37).

Figure I.3: Dependence of volume suppression of KK-graviton and bulk massless gauge boson the coupling with the KK-graviton mass.

Weak bosons

The W and Z fields couplings to KK-graviton have two components: one component related with the massless part of the weak bosons (transverse polarizations) volume suppressed by the d_g parameter and one component proportional to the weak bosons EWSB mass that is not volume suppressed. The physical weak boson fields are defined on the unitary gauge. In this gauge we write the KK-graviton couplings to the massive weak bosons as:

$$d_W T_{\mu\nu}^{WW} = - \left(m_W^2 \frac{\delta(g_{\alpha\beta} W_+^\alpha W_-^\beta)}{\delta g^{\mu\nu}} + \frac{m_Z^2}{2} \frac{\delta(g_{\alpha\beta} Z^\alpha Z^\beta)}{\delta g^{\mu\nu}} \right) - d_g T_{\mu\nu}(W^\pm, Z, \gamma)^{YM}, \quad (\text{I.73})$$

where $T_{\mu\nu}^{YM}(W^\pm, Z, \gamma)$ is the canonical Yang Mills energy momentum tensor for $SU(2)_L \times U(1)_Y$ for the massless part of the gauge fields (see equation I.20b).

The decay rates of KK-graviton to a pair of physical SM massive weak bosons are [86]:

$$\Gamma(h_{\mu\nu} \rightarrow ZZ) = \frac{m_{Gr}^3}{960\pi} \left(\frac{\tilde{k} x_1}{m_{Gr}} \right)^2 (A_Z + d_g B_Z + d_g^2 C_Z) \sqrt{1 - 4 \frac{m_i^2}{m_{Gr}^2}}, \quad (I.74)$$

$$\Gamma(h_{\mu\nu} \rightarrow WW) = \frac{m_{Gr}^3}{480\pi} \left(\frac{\tilde{k} x_1}{m_{Gr}} \right)^2 (A_W + d_g B_W + d_g^2 C_W) \sqrt{1 - 4 \frac{m_i^2}{m_{Gr}^2}}, \quad (I.75)$$

where for each $i = W, Z$:

$$A_i = \left(1 + 12 \frac{m_i^2}{m_{Gr}^2} + 56 \frac{m_i^2}{m_{Gr}^2} \right), \quad B_i = 80 \left(1 - \frac{m_i^2}{m_{Gr}^2} \right) \frac{m_i^2}{m_{Gr}^2}, \quad C_i = 12 \left(1 - 3 \frac{m_i^2}{m_{Gr}^2} + 6 \frac{m_i^4}{m_{Gr}^4} \right). \quad (I.76)$$

As usual the longitudinal degrees of freedom of the weak bosons results of the absorption of the the Goldstone bosons by the physical states. At high energies it is possible to realize analytically the relation between the Goldstones bosons and the longitudinal polarization of weak bosons [87, 88], what is known as equivalence theorem.

Figure (I.4) shows the split of the KK-graviton decay rate to weak bosons in the different weak bosons polarizations modes. In the RS1 scenario the KK-graviton decays to preferentially to transverse polarized modes while in the bulk scenario the KK-graviton decays preferentially to longitudinally polarized modes. This change of behavior is a direct consequence of the transverse modes volume suppression d_g .

Figure I.4: Separation of polarizations on decay product of $h_{\mu\nu}^{(1)} \rightarrow ZZ$. The same behavior holds to $h_{\mu\nu}^{(1)} \rightarrow WW$ decays. The continuous red line represents the portion of decay to pure longitudinal gauge boson modes (LL), the dot-dashed blue is the portion of decay to pure transverse modes (TT) and the dashed green the portion of decay to longitudinal-transverse mixed final states (LT).

Fermions

Analogously to the case of the KK-graviton couplings to weak gauge bosons, the KK-graviton couplings to fermions also have one component proportional to the Dirac energy momentum tensor of each fermion chirality that is weighted by the overlap between the profiles of KK-graviton and the massless fermions, and also other component proportional to the EWSB mass.

Schematically the coupling of KK-graviton to a fermion ψ in the fermion diagonalized mass basis (used in equation I.18) have the form:

$$d_\psi T_{\mu\nu}^\psi = d_{f_L} T_{\mu\nu}^{\psi_L} + d_{f_R} T_{\mu\nu}^{\psi_R} + d_{f_{LR}} g_{\mu\nu} \lambda_f v^2 (\bar{\psi}_R \psi_L + h.c.), \quad (\text{I.77})$$

where $d_{LR} \equiv \int f_{Gr} f_{f_R} f_{f_L}$ and $d_{L/R} \equiv \int f_{Gr} f_{f_{L/R}} f_{f_{L/R}}$. Three localization parameters are necessary to define the KK-graviton couplings to each fermion generation.

The SM leptons and first two quarks generations are very light compared with the electroweak scale, therefore we need to choose the profiles of both fermion chiralities to be Planck localized, resulting in the constants $d_{f_L} = d_{f_R} = d_{f_{LR}} \sim 0$. We therefore ignore KK-graviton couplings to bulk light fermions.

The top quark however has mass ~ 175 GeV, comparable to the weak scale and therefore its coupling to KK-gravitons can be sizeable. For practical reasons we consider the model where only the SM gauge symmetry $SU(2)_L \times U(1)_Y$ is enlarged to the bulk and the constants $d_{f_L} \sim d_{f_{LR}} \sim 0$ and $d_{f_R} \sim 1$ [89]. The decay rate of KK-graviton to a pair of top quarks with mass m_T is:

$$\Gamma(h_{\mu\nu} \rightarrow t\bar{t}) = \frac{m_{Gr}^3}{80\pi} \left(\frac{\tilde{k} x_1}{m_{Gr}} \right)^2 \left(\frac{3}{4} - \frac{7}{2} \frac{m_T^2}{m_{Gr}^2} + 2 \frac{m_T^4}{m_{Gr}^4} \right) \sqrt{1 - 4 \frac{m_T^2}{m_{Gr}^2}}. \quad (\text{I.78})$$

In constructions with an extended bulk gauge group the localization of the profiles of the third family of quarks is modified and consequently the KK-graviton couplings to the top and bottom quarks is also modified.

I.4.2 Radion couplings to matter

In this subsection we review the definition of the a_i constants that appear in equation (I.68) for each field hypothesis. We follow the references [36, 90].

Massive fields

The contributions to the radion couplings to massive fields that arises from bulk effects deserves a detailed study, the reference [85] find the effect of consider bulk matter to all the radion branching ratios that are dominant on each mass range to be very small. We ignore the bulk effects in radion coupling to massive states and consider the coupling between the radion and massive gauge bosons as:

$$a_V T_\mu^{\mu(V)}(x) = \left(\int d\phi f_H^2 f_r \right) \int d^4x \quad (m_W^2 W_{+\mu} W_-^\mu + \frac{m_Z^2}{2} Z_\mu Z^\mu), \quad (\text{I.79})$$

and the couplings of the radion to each fermion f_i , in the mass diagonalized fermion basis, as:

$$a_{f_i} T_\mu^{\mu(f_i)}(x) = \left(\int d\phi f_H^2 f_r \right) \int d^4x \quad m_{f_i} \quad (\bar{\psi}_{f_i L} \psi_{f_i R} + h.c.). \quad (\text{I.80})$$

We assume the profile of the radion field to be TeV localized (see section I.3.2.3) and therefore in both RS1 and bulk scenario the strength of the couplings between the radion and a massive SM states is:

$$\int d\phi f_H^2 f_r \sim 1 \quad (\text{I.81})$$

Massless Gauge fields

Following equation (I.68) the couplings of the radion field with the massless gluon (g) and photon (γ) are proportional to the β function (see equation I.21). In the case of bulk gauge fields we find also a generic source in addition to equation (I.68) contributing to the coupling of the massless gauge bosons:

$$S_{bulk}^{(g,\gamma)} \equiv -\frac{1}{4k\pi r_c} \int d^4x \quad \frac{r}{\Lambda_R} \left(\sum_{colors} F_{\mu\nu}^{(g)} F^{(g),\mu\nu} + F_{\mu\nu}^{(\gamma)} F^{(\gamma),\mu\nu} \right). \quad (\text{I.82})$$

This term arises from the expansion of the bulk field strength F_{MN} on space-time components. The coupling suppressions for the leading order radion couplings to both gluon and photon in the bulk scenario are:

$$a_g = \frac{\alpha_s}{2\pi\Lambda_R} \beta_{QCD} + \frac{1}{4k\pi r_c}, \quad a_\gamma = \frac{\alpha_{QED}}{2\pi\Lambda_R} \beta_{QCD} + \frac{1}{4k\pi r_c}. \quad (\text{I.83})$$

The biggest effect of the change of scenario from RS1 to bulk in hadron colliders is the constructive contribution of the bulk term (I.82) on the gluon fusion radion production.

I.5 Phenomenology of gravity particles at a proton collider

We are interested in the direct detection of the gravity particles in a proton-proton collider. Protons are objects made by three valence quarks and a sea of gluons and quark and anti-quark pairs. In an interaction with a proton the constitution of the sea depend on the energy scale of the interaction. This energy scale can be interpreted as the resolution of the hadron that is being tested and it is also known as factorization scale [91]. The functions that determine the density of quarks and gluons at given factorization scale are known as Parton Distribution Functions (PDFs), those are determined experimentally [92].

The cross section for a hard scattering process with two hadrons i, j with four momentum p_i and p_j in the initial state is calculated from a proton collision as a convolution of hard scattering cross section $\hat{\sigma}$ with the PDFs of the initial state partons at the factorization scale μ .

$$\sigma_{PP} = \sum_{i,j} \int dx_1 dx_2 f_1(x_1, \mu^2) f_2(x_2, \mu^2) \hat{\sigma}(p_i, p_j, \mu), \quad (\text{I.84})$$

where x_1 and x_2 are the fraction of energy each one of the partons carries from the initial protons and f_i is the PDF associated to each parton. At high energies the proton density is dominated by gluons.

There are four production modes for a heavy resonance in hadron colliders. Two of them are inclusive gluon fusion and quark fusion (Drell Yan). In figure (I.5) we show the diagrams for inclusive production of a resonance that couple to a pair of gluon and/or a pair of quarks.

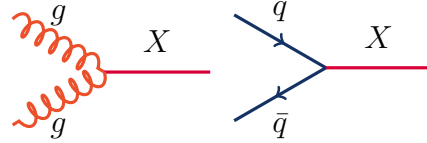


Figure I.5: Feynman diagrams for inclusive production of a X resonance that couple to a pair of gluons and/or a pair of quarks on hadronic colliders. **Left:** Gluon fusion. **Right:** Drell Yan. To bulk scenario the Drell Yan production

A heavy resonance that couples to the weak bosons can also be produced in association with a weak boson and the Weak Boson Fusion (WBF). We show the diagrams of WBF and weak boson associated production in the figure I.6.

Assuming the weak boson from associated production to be hadronic both production modes are characterized by the presence of two additional light quarks in the hard process, what at this level we will call jets. The common cuts we use di-jet system are:

$$p_{t_j} > 20 \text{ GeV}, \quad \Delta R_{jj} > 0.1, \quad (\text{I.85})$$

where p_{t_j} is the transverse momenta of each one of the jet (final state light quark) and ΔR_{jj} is the angular distance between the hard jets. Relaxing of the ΔR_{jj} cut is designed to take into account the possibility of tag the two hard jets as coming from the weak boson in the associated production.

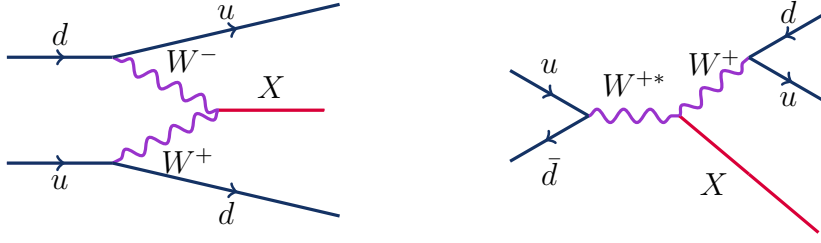


Figure I.6: Examples of Feynman diagrams for the production of a resonance X that couples to a pair weak bosons in a proton collider. **Left:** Weak boson fusion. **Right:** Weak boson associated production.

I.5.1 KK-graviton

For the calculation of RS1 KK-graviton cross section at proton colliders we had used **MadGraph5** [93]. The source model used was the RS1 model [94] from **Feynrules** [95] database. We had checked the branching ratios of the model with the analytical expressions presented here. The results for RS1 KK-graviton production cross sections also agrees with the WED implementation default on **Pythia6** [96].

The production cross sections for the bulk KK-graviton was also calculated using **MadGraph5**, we adapted the **Feynrules** model [94] to match the bulk scenario (with no KK-graviton coupling to top quarks) removing all the bulk KK-graviton couplings to fermions and rescaling the KK-graviton couplings to massless gauge bosons (both photon, gluon and transverse part of weak bosons) with the volume suppression d_g (equation I.71).

We had found agreement on the production cross section between the model mentioned above and the bulk WED scenario implemented by the authors [97] on **CalcHep** [98]. We used **CTEQ61** [99] as PDF set, considering dynamical factorization scale.

In the figure (I.7) we show the production cross section of KK-graviton on both RS1 and bulk scenario on the exclusive production, weak boson associated and WBF production* for the parameter $\tilde{k} = 0.2$. The coupling of KK-graviton to bulk light fermions is negligible, therefore we consider only gluon fusion as inclusive production of bulk KK-graviton.

When comparing the production cross sections of KK-graviton in bulk and RS1 scenario we see that for each production mode the production cross section in the bulk scenario is smaller than the to RS1 when the same compactification scale is assumed. This effect is expected in the gluon fusion production mode, since the KK-graviton coupling to gluons suffer by the volume suppression d_g .

The dominant part the KK-graviton coupling is to transversely polarized weak bosons

*To calculate processes with **MadGraph5** simulation we can only specify initial and final state. When we generate quark initiated process $pp \rightarrow jj h_{\mu\nu}^{(1)}$ (di-jet associated production) this full process will both WBF like and the weak boson associated topologies. We separate the two topologies by means of a mass cut on the final state di-jet system. The WBF cross section is defined by imposing a cut $M_{jj} > 100$ GeV in this process. The weak boson associated production is defined by the same process but under the cut $M_{jj} < 100$ GeV.

Figure I.7: Production cross sections with $\tilde{k} = 0.2$ in a proton-proton collider, in pb. The red curves correspond to the inclusive production, the green ones to WBF and the blue ones to the hadronic weak boson associated production. The dot-dashed, continuous and dashed curves are respectively for the center of mass energy of 14 TeV, 13 TeV and 8 TeV. **Right:** Bulk scenario. **Left:** RS1 scenario. The cross section values scale with $\tilde{k}^2$ on all production mechanisms.

and therefore the di-jet associated production modes (WBF and associated production) also experiences the volume suppression (d_g).

The decays modes of the resonance is also an important information for its collider phenomenology. The best search mode for a given resonance hypothesis at colliders is characterized by a compromise between the size of the branching fraction and the quantity of expected SM background events with the same final state that would fake the resonance detection.

Figure (I.8) shows the KK-graviton branching fractions to SM particles. The left part of this figure shows its branching ratios on bulk scenario, assuming the composite top scenario. The right part shows KK-graviton branching ratios on RS1 scenario. In this figure we clearly note the couplings of RS1 KK-graviton with SM particles are democratic with respect to the number of degrees of freedom of each field, while the couplings of bulk KK-graviton with SM particles are predominantly to the fields localized on TeV brane (massive bosons and higgs field). The highest branching fraction of a RS1 KK-graviton is to di-jet final states ($gg + q\bar{q}$). In the bulk scenario the highest KK-graviton branching fractions are to the massive states, heavy di-bosons and top quarks (in the composite top scenario).

The left part of figure (I.9) shows a zoom of KK-graviton branching ratios comparing two distinct scenarios for third family WED embedding: The thick curves considers a fully composite top quark [97] while thin dashed ones we ignores KK-graviton coupling with top quark. It is not a surprise to note that KK-graviton branching ratios to bosons pairs ($W/Z/H$) increases when there is no KK-graviton coupling to top quark.

The right part of figure (I.9) shows the KK-graviton total width in the discussed scenarios for the same parameter $\tilde{k} = 0.2$. Given the same geometric parameters the total width on RS1 scenario the KK-graviton is around two orders of magnitude larger than

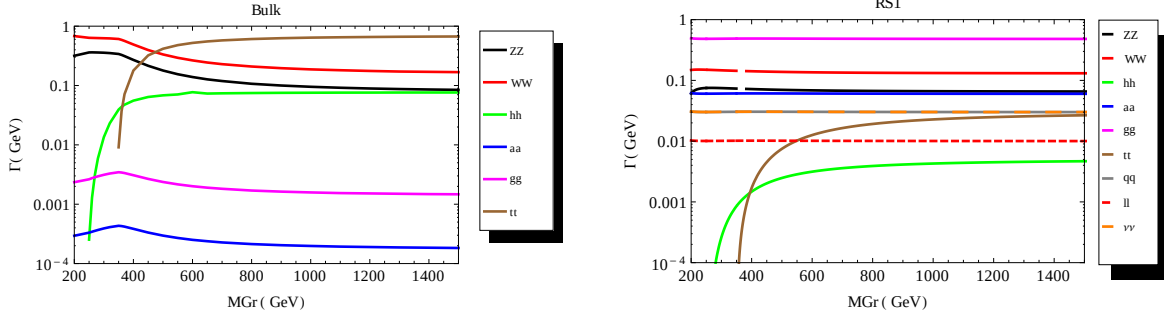


Figure I.8: KK-graviton branching fractions. **Left:** Bulk scenario assuming composite top quark. **Right** RS1 scenario. The dashed line represents the individual KK-graviton decay rates to a pair of light fermions $f\bar{f}$, where f represents both light quarks (u, d, s, c, b) or leptons (e, μ, τ). The branching ratios are independent of $\tilde{k}$ parameter.

in bulk scenario. The experimental resolution on the width of the resonance depends on the investigated channels and also of the mass hypothesis probed. The experimental searches are usually designed using signal models with negligible width with respect to experimental resolution.

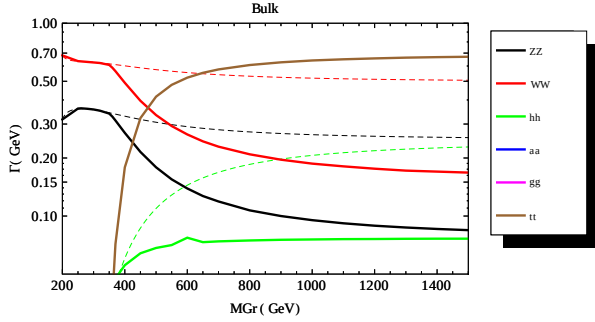


Figure I.9: **Left:** Zoom of the dominant branching ratios for bulk KK-graviton comparing two hypothesis of fermion WED embedding. The thick curves considers a fully composite top quark ([97]) and thin dashed ones completely ignores coupling with top quark. **Right:** Total width for $k/M_{Pl} = 0.2$. The green curve represents RS1 scenario and the red curves represents bulk scenario. We show two different hypothesis of treating KK-graviton couplings to bulk third family of quarks: the continuous curve considers a fully composite top quark ([97]) while dashed curve ignores KK-graviton coupling with top quark. Total width values scale with $\tilde{k}^2$ parameter.

I.5.2 Radion production

We consider the radion production in hadron colliders (and decays) defined by three independent parameters. We chose these parameters to be the radion mass M_R , the

compactification scale kr_c and the radion-higgs mixing parameter ξ . We do not want to use the strict relation between the ultraviolet mass scale of the theory Λ_R and the WED geometry and consider this value also as an external input.

Analogously to the SM higgs boson the coupling of a radion field to a pair of photons and gluons (with momenta p_1 and p_2) receives contributions from heavy fermions and weak bosons loops that are comparable to a_g and a_γ . Therefore the loop corrections are important on considering the gluon fusion production mode and the radion decay width to photons (see for example [100]).

In the bulk scenario the electroweak loop corrections to the radion coupling to massless gauge bosons are considered together with trace anomaly and bulk components in the Born level couplings (see section I.4.2). The Born level radion couplings to gluons (g) and photons (γ) are:

$$\frac{c_{iir}}{\Lambda_R} (p_1 \cdot p_2 g_{\mu\nu} - p_{1,\mu} p_{2,\nu}), \quad (\text{I.86a})$$

where $i = g, \gamma$ and the c_{iir} constants are:

$$c_{ggr} = \frac{\alpha_s}{4\pi} \left(\sum_{\text{fermions}} Q_f F_{1/2}(\tau_f) - 2 \beta_{QCD} \right) - \frac{1}{4kr_c \pi} \quad (\text{I.86b})$$

$$c_{\gamma\gamma r} = \frac{\alpha_{QED}}{4\pi} \left(\sum_{\text{fermions}} Q_f F_{1/2}(\tau_f) + F_1(\tau_w) - 2 \beta_{QED} \right) - \frac{1}{4kr_c \pi}, \quad (\text{I.86c})$$

The $F_{1/2}(\tau_f)$ correspond to the loop functions related with a fermion f with charge Q_f running in the loop and $F_1(\tau_w)$ is the loop function for a W boson running in the loop*.

Differently of what was made to KK-graviton, we only perform calculations for radion phenomenology for the bulk scenario. This point is not particularly important for phenomenology matters since the change of scenarios does not modify the coupling momentum dependence and therefore does not modify production and decay topology. Reference [85]

*The one loop level functions to higgs-like scalar coupling to massless bosons mediated by fermion and weak boson loops are [101]:

$$F_{1/2} = -\frac{2}{\tau^2} \tau [1 + (\tau - 1)f(\tau)],$$

$$F_1 = \frac{1}{\tau^2} (2 + 3\tau + 3\tau(2 - \tau)f(\tau)),$$

where $\tau_i \equiv \frac{M^2}{4m_i^2}$, M is the mass of the produced particle and m_i is the mass of the particle on the loop. The f functions depend on the ratio of masses τ_i as:

$$f(\tau) = \begin{cases} \arcsin^{-2}(\sqrt{\tau}), & \text{if } \tau \leq 1 \\ -\frac{1}{4} [\log(\frac{1+\sqrt{1+\tau-1}}{1-\sqrt{1+\tau-1}}) - i\pi]^2, & \text{if } \tau > 1 \end{cases}.$$

In the limit $\tau \rightarrow \infty$ them $F_{1/2} \rightarrow 4/3$ and $F_1 \rightarrow 7$.

note that there is no significant modification on the radion branching fractions to the dominant decay modes if instead of bulk scenario we consider RS1 scenario.

The QCD and EW corrections have a big impact the production cross sections and decays of a higgs-like particle in proton colliders. Differently from a generic higgs-like, where the scalar coupling to gluons happens exclusively mediated by the fermion loop the trace anomalies (and eventually the bulk term) are a point-like interaction between the radion and gluons.

In a first approach we ignore the fact that the QCD corrections to a process like $gg \rightarrow X$ in the case of the coupling be mediated by the a loop of a finite mass quarks are different from the QCD corrections to a point-like interaction. Under this hypothesis it is possible to explore the similarity between the radion and higgs interactions with matter and extrapolate the state-of-art to next-to-next-to-leading logarithmic accuracy (NLL) on soft gluon resummation scheme for QCD corrections and NLO EW corrections already calculated for the SM higgs (h_{SM}) cross sections to calculate radion production and decay. The references [102, 103, 104] pushed this method for the interpretation the radion parameter space using LHC8 data.

The productions modes of the higgs (h) and radion (r) particle in a proton-proton collider can be divided in two ways: gluon fusion initiated and the non-gluon fusion initiated production modes. The non gluon fusion production modes are the weak boson associated production (Wr), WBF and the associated production with a pair of top quarks ($t\bar{t}r$). The WBF and weak boson associated radion production involves its couplings to weak bosons, while the $t\bar{t}$ associated production involves the radion coupling to top quarks. The dominant part of radion couplings to massive particles have exactly the same structure than the higgs-like case. It is clear that for a non-gluon fusion production we can directly extrapolate all the QCD and EW higher order corrections calculated to a higgs-like particle production to the radion case:

$$\sigma(pp \rightarrow r + X)_{Wr, WBF} = \sigma(pp \rightarrow h_{SM} * + X)_{Wr, WBF} \frac{\Gamma(\phi \rightarrow VV)}{\Gamma(h_{SM} * \rightarrow VV)}, \quad (I.88a)$$

$$\sigma(pp \rightarrow r + X)_{t\bar{t}r} = \sigma(pp \rightarrow h_{SM} * + X)_{t\bar{t}r} \frac{\Gamma(\phi \rightarrow t\bar{t})}{\Gamma(h_{SM} * \rightarrow t\bar{t})}, \quad (I.88b)$$

where $h_{SM}*$ denotes a SM like boson with the some mass hypothesis than the radion (r), and the cross sections for $h_{SM}*$ are the quantum corrected ones. All the quantum corrections on the decay rates factorize. The ratio of decay rates can be taken as the leading order ratio.

The inclusive modes for radion LHC production is mainly gluon fusion. We naively calculate the quantum corrected gluon fusion production of radion field using the same

assumption we had done for the non-gluon fusion modes:

$$\sigma(pp \rightarrow \phi + X)_{gg} = \sigma(pp \rightarrow h_{SM} * + X)_{gg} \frac{\Gamma(\phi \rightarrow gg)}{\Gamma(h_{SM} * \rightarrow gg)} \quad (I.89)$$

In figure (I.10) we show the ratio between the cross sections for gluon fusion radion production mode on LHC8 calculated, using NLL accuracy on based on soft gluon resummation scheme [105] and our Born level calculations, using CTEQ6L as PDF set. We had used the NLO calculation for the α_s value and two loop level QCD beta function* for the Born calculation (defined by the equations I.86).

Figure I.10: **Right:** Ratio between the cross sections for gluon fusion production mode in a proton collider with center of mass energy of 8 TeV, calculated by [105] and the same gluon fusion production at Born level. The red (green) curve is for higgs (radion) production. We use $\Lambda_R = 3$, $kr_c = 11$ and ignore radion-higgs mixing.

For practical reasons we had used a different PDF set from the one used by the reference that performed the calculation [105] (this group use MRST [107] as PDF set). Therefore it is not possible to compare the exact value for the gluon fusion K-factors to higgs gluon fusion production on a proton-proton collider with 8 TeV center of mass energy. The values for the ratio σ/σ_{LO} reported in figure (I.10) lies around 1.5 and 3, compatible with [105].

The correction on the gluon fusion production of a higgs-like particle have a kink when the mass of the produced particle around 350 GeV, due to the loop function transition when the mass hypothesis of the resonance exceeds two times the top quark mass. The different slope on the radion and higgs correction is consequence of the different dependence of α_S and the on the Born radion gluon fusion production (I.86).

In figure (I.11) we show the production cross section to both bulk radion and higgs particle with mass calculated with formula (I.89), having in mind that for a $M_R > 350$ we

* The two loop level α_S constant is $\alpha_S(q) = \frac{\alpha_S}{1 + \alpha_S w(q)}$, where $w(q) = (\beta_0 + \beta_1 * \alpha_S) * \text{Log}(q/m_Z)$, $\beta_0 = \frac{1}{2\pi} (11.0 - \frac{2}{3} n_f)$, $\beta_1 = \frac{1}{8\pi^2} (102.0 - \frac{38}{3} n_f)$ and n_f is the number of fermions with mass bellow the energy scale q [106].

The two loop level β function for QCD is $\beta_{QCD}(q) = (\beta_0 + \beta_1 \alpha_S) \log\left(\frac{q}{m_Z}\right)$ [106].

may be underestimating gluon fusion cross section values. The radion gluon fusion production is more dominant over the non-gluon initiated radion production modes than SM higgs case. In appendix C we comment on the effect of mixing in the different production modes.

Figure I.11: Production cross sections for a higgs-like particle on a proton-proton collider with 8 TeV CM energy. We assume no higgs-radion mixing. Red for gluon fusion production, green for WBF production, blue for W associated production, blue for Z associated production and black for $t\bar{t}$ associated production. **Left:** SM higgs. **Right:** radion (with $\Lambda_R = 3$ TeV and $kr_c = 11$). The radion production cross section scales with Λ_R^{-2} on RS1 scenario for all production modes.

To calculate radion decay rates we use the same trick used in the above section to calculate production cross section. The quantum corrected radion decay rates to weak bosons and fermions can be straightforwardly calculated from [108] by the relation:

$$\Gamma(r \rightarrow VV, ff) = \left(\frac{\Gamma(r \rightarrow VV, ff)}{\Gamma(h_{SM}^* \rightarrow VV, ff)_{SM}} \right)_{LO} \Gamma(h_{SM}^* \rightarrow VV, ff) \quad (\text{I.90a})$$

and, for the radion decay to massless particles:

$$\Gamma(r \rightarrow gg, \gamma\gamma) = \left(\frac{\Gamma(r \rightarrow gg, \gamma\gamma)}{\Gamma(h_{SM}^* \rightarrow gg, \gamma\gamma)} \right)_{LO} \Gamma(h_{SM}^* \rightarrow gg, \gamma\gamma). \quad (\text{I.90b})$$

Similar caveats as the ones mentioned for the radion production about the quantum corrections of higgs scalar to be extended to the radion gluon fusion production can be done also for the radion decays to a pair of gluons or photons. We note however that for radion mass $> 2M_W$ its branching ratio to gluons and photons is sub-dominant and however any K-factor correction on the radion branching ratio to gluons and photons that does not exceed $O(10)$ is not relevant for phenomenology.

When kinematically allowed the radion decays to a pair of higgs bosons. The particles

in the process $r \rightarrow h h$ have no color and therefore the decay process does not receive QCD corrections. We ignore the electroweak corrections and use the leading order expression. The leading order expression for this mode is:

$$\Gamma(r \rightarrow h h)_{LO} = \frac{m_R^3}{32\pi^2 \Lambda_R^2} \left(1 + \frac{4m_H^2}{m_R^2}\right) \sqrt{1 - \frac{4m_H^2}{m_R^2}} \quad (\text{I.91})$$

The gluon pair final state dominates the branching fractions in the mass range ($m_R < 2m_W$) [85, 100]. Being larger than the rate to b-quarks what is a clear effect of the trace anomaly (see figure I.12).

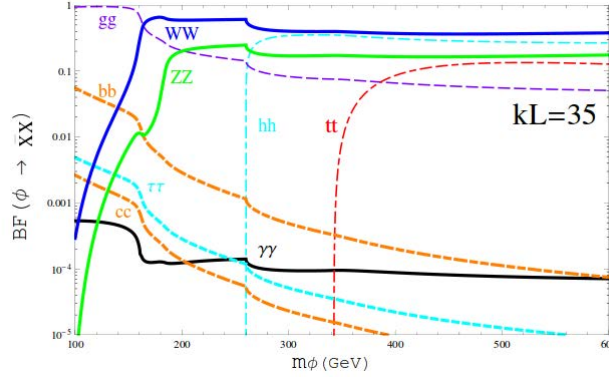


Figure I.12: Radion decay rates cross without higgs-radion mixing for a higgs boson of 130 GeV. Figure from [103].

Once the decays to the electroweak bosons (W , Z and h) becomes kinematically possible the decays for these states become the dominant part on branching ratio. In the non mixing scenario up to the 700 GeV mass hypothesis the radion width is below 0.1 GeV when the parameter $\Lambda_R = 10$ TeV [83] therefore it is valid to use narrow width approximation in cross section calculations.

II Phenomenology of resonant double higgs production at LHC

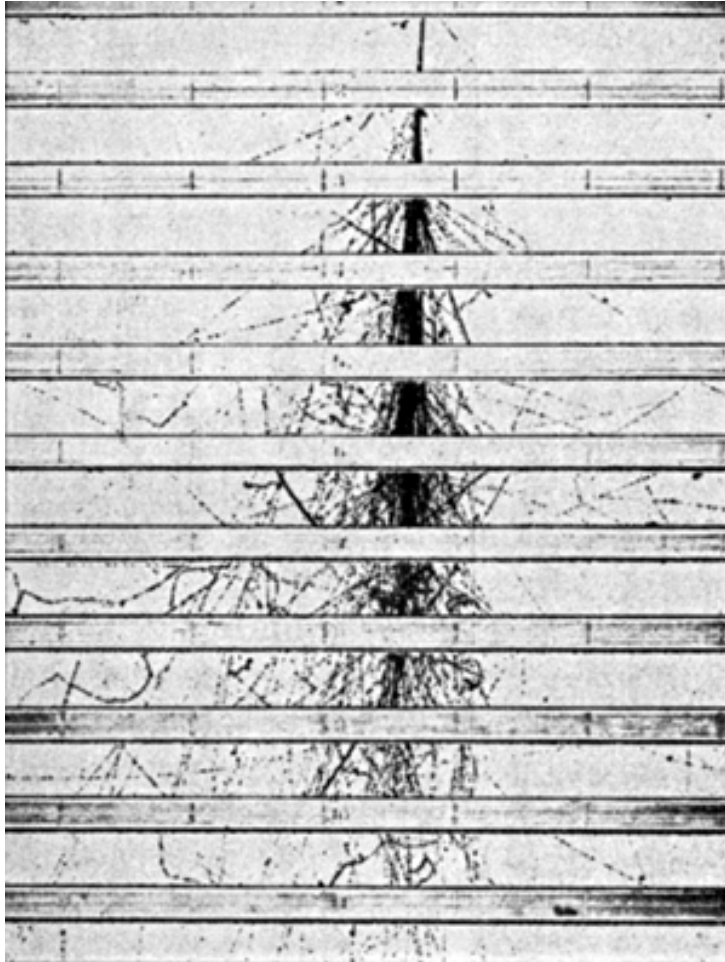


Figure covering this chapter: A jet formed by a cosmic ray passing through photographic layers, photo taken by the MIT cosmic rays group in 1933 [109].

In this chapter we discuss phenomenology of spin 0 and spin 2 resonances at LHC collider. We divide the chapter in five sections. In section (II.1) we describe the basics of the LHC project and its experiments. We next study the general topology of event processes with the decay of a heavy resonance into identical unstable resonances in section (II.3). In the following sections (II.4) and (II.4) we exploit the LHC sensitivity to a fully hadronic di-higgs resonance on inclusive and WBF production modes respectively. The goal of this work was to demonstrate the potential of fully hadronic di-higgs searches at the LHC on a wide range of new resonance mass hypotheses.

In general BSM new physics is supposed to couple to higgs sector, consequently in many SM extensions exotic heavy resonances have the higgs and weak bosons among its main decays. The collider search for heavy spin 0/2 resonances decaying to a pair of higgses benefits from the now well known higgs mass and the characterization of its decays [6, 7]. A SM higgs boson with mass 125 GeV decays preferentially to a b-quark pair (57.8%). In the table (II.1) we show the decay rates for the main channels of a di-higgs pair if we assume SM-like branching ratios for the 125 GeV higgs (as suggested by data [110]).

channel	frequency(%)
$h(b\bar{b}, c\bar{c}, gg)h(b\bar{b}, c\bar{c}, gg)$	47.86
$h(b\bar{b})h(b\bar{b})$	33.29
$h(b\bar{b}, c\bar{c}, gg)h(VV^*)$	33.40
$h(b\bar{b}, c\bar{c}, gg)h(\tau^+\tau^-)$	8.77
$h(VV^*)h(VV^*)$	5.83
$h(l^+l^-)h(VV^*)$	3.06
$h(l^+l^-)h(l^+l^-)$	0.40
$h(b\bar{b}, c\bar{c}, gg)h(\gamma\gamma)$	0.32
$h(b\bar{b}, c\bar{c}, gg)h(\mu^+\mu^-)$	0.03
$h(l^+l^-)h(\gamma\gamma)$	0.03

Table II.1: Branching ratios for a higgs pair assuming the SM like branching fractions.

The dominant decay rate of a pair of higgses is into fully hadronic mode. Searches for fully hadronic final states at hadronic colliders have as challenge the very large QCD multi-jet background. The improvement of techniques that allow the search for New Physics (NP) signal extraction in multi-jet events is an important element of the LHC physics program.

In the signal simulation the resonances was required to be produced on-shell and further decayed into a pair of (125 GeV) higgs bosons in **MadGraph5**. To simulation purposes the total width is assumed to be 1 GeV in all resonance mass spectra, which is negligible with respect to detector resolution and also consistent with the radion [64] and bulk KK-graviton [89] scenarios (see figure I.9).

II.1 The LHC and its general purpose experiments

The Large Hadron Collider (LHC) is a circular complex designed for acceleration and high energy collision of hadrons. It is installed in a 26.7 km tunnel situated in between France and Switzerland, the same tunnel used by LEP (Large Electron Positron collider) [111, 112]. The LHC and the pre-acceleration system is hosted by the *Center Europeene pour la Recherche Nucleaire* (CERN).

The majority of the LHC operation time is dedicated to accelerate and collide proton beams. Approximately one moth of each year of LHC operation is reserved to lead-lead ad proton-lead collisions, for specific Quark Gluon Plasma studies [113].

Four experiments profits from LHC collisions to take data: LHCb, ALICE, CMS and ATLAS. LHCb and ALICE are designed to the research on flavor and heavy ions physics, respectively. CMS and ATLAS are multi purpose designed. Their focus is to investigate the electroweak sector of SM and BSM and as well QCD physics. In figure (II.1) we show the schematic position of the detectors of the different experiments and main accelerator modules. A full description of LHC and the pre-acceleration system and the list and purpose of all experiments that profit of this complex can be found elsewhere, see for example [111]. In this thesis we are interested on the physics that can be extracted from proton-proton collisions recorded by the two multi-purpose ATLAS and CMS experiments.

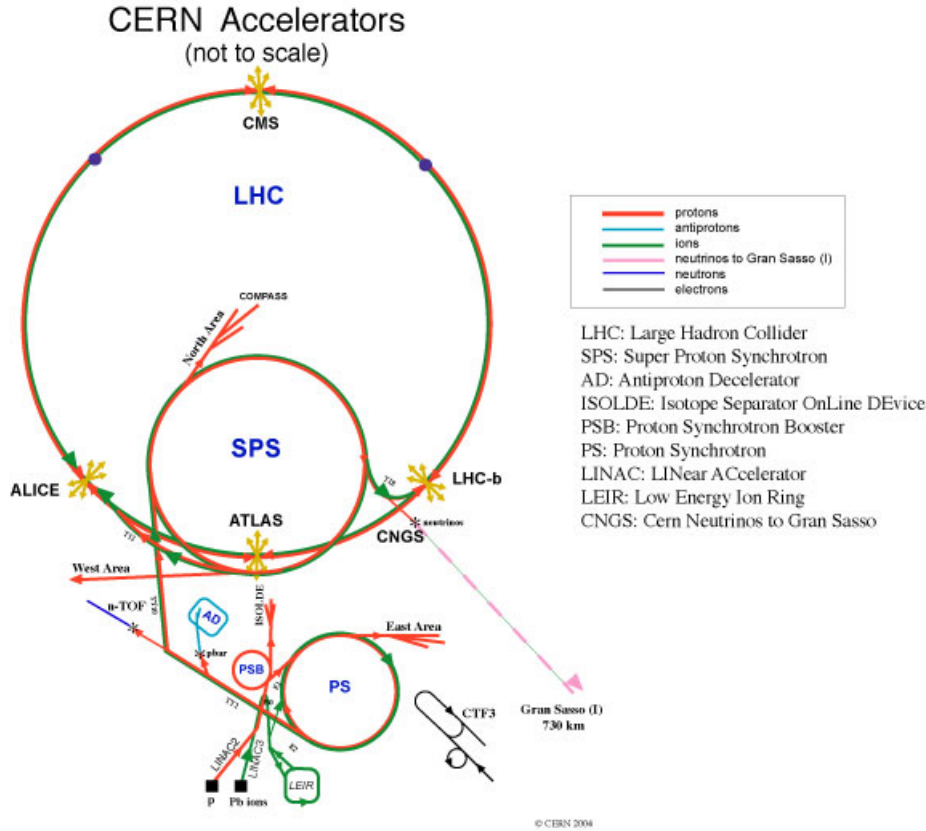


Figure II.1: Schematic view of LHC accelerator system and its four main experiments [114]. The yellow stars represents the collision points.

In the LHC the protons (or lead atoms) beams are organized in bunches. The length, gauge and spacing of the bunches are controlled during the LHC pre-acceleration system. If we consider the collision area between beams to be shaped as a perfect ellipse with transverse sections σ_x and σ_y the instantaneous luminosity can be calculated as*

$$L = \frac{f}{2\pi} \frac{N_{p1} N_{p2}}{\sqrt{2}\sigma_x \sqrt{2}\sigma_y}, \quad (\text{II.1})$$

where the N_p 's are the number of nuclei in each bunch and f is the frequency of collisions. The instantaneous luminosity relates the rate of collisions by time dN/dt with the total cross section:

$$\frac{dN}{dt} = L\sigma. \quad (\text{II.2})$$

The potential for NP discovery of the experiments is a by product of the total luminosity ($\equiv N_{events}/\text{cross section}$) produced with LHC collisions. This quantity is the probability of any event to occur on a beam collision. For example, giving a instantaneous luminosity of $O(10/\text{sec}/\text{nb})$, and a given process with cross section of 10 nb then we should expect a rate of 100 events per second.

II.1.1 LHC run conditions

The LHC was designed to reach 14 TeV of center of mass (CM) energy in performing proton-proton collisions. Due technical problems the machine had started its first run with 7 TeV CM energy started in 2010 collecting 35/pb of data. Since then the CM energy is being gradually increased. In the second run of LHC the CM energy was increased to 8 TeV by a trade off between the magnets sensitivity and the to the higgs boson discovery probability.

In the LHC conditions in the Run I (2010-2012) the bunches were separated by a distance of 50 ns[†]. Meaning a bunch crossing happens every 50 ns during this bunch crossing more than one inelastic collision may occur. For each triggered event however there is the possibility of some of the recorded radiation not to be originated from the triggered hard event.

The average instantaneous luminosity delivered for LHC to the experiments during Run 1 corresponds to 5/sec/nb[‡]. For example the cross section for (not interesting to elementary particles research) elastic proton interactions at 8 TeV CM energy is 74.7 ± 1.7 mb [115], therefore we expect $\sim 75 \times 10^6$ events per second. The kind of events we are interested have a cross section of order of pb (see for example figures C.2 and I.11).

*The formula (II.1) receives corrections considering the possible deformation on beam geometry, angle of crossing, gradient of density and emittance. This discussion however is not important in this thesis context.

[†]https://twiki.cern.ch/twiki/bin/view/CMSPublic/LumiPublicResults#Pileup_distribution

[‡] <https://twiki.cern.ch/twiki/bin/view/CMSPublic/LumiPublicResults>

The capacity for data storage of the experiments is finite. To be sensitive to $\sim$ pb cross sections out of LHC collisions one important computational part of the experiments is the design of efficient triggers systems to select a specific hard event among many background collisions. Each trigger is specially designed to select a given pattern, according with each suitable physics search.

The events that contaminate the triggered hard events are also known as pileup events. There are two kinds of pileup events: The *out of time* pileup consists of events from different crossing. The *in time* pileup consists of events originated from the same bunch crossing where the hard event was originated.

To increase the probability of a rare event to occur during a given data taking period one needs to increase the instantaneous luminosity. The instantaneous luminosity can be increased in two ways: decreasing the time spacing between the bunches (increasing out of time pileup) or increasing the bunches density (increasing in time pileup). Both methods increase the quantity of pileup events and also demand the development of better and faster trigger algorithms.

II.1.2 Future

The center of mass energy during the Run 2 period of LHC operation stage is agreed to be 13 TeV and pileup to not increase too much from what used on the last phases of Run 1. Increasing pileup (decreasing the bunches distance) increases the experimental challenge to the detectors recognize the hard event localization. One of the most important experimental challenges at high pileup are to improve the techniques for removal of underlying event QCD jets, that contaminates the hard event. The jet area methods developed by the authors [116, 117] are among the current techniques used for pileup subtraction by LHC collaborations [118, 119, 120].

During Run 3 the bunch spacing is expected to decrease to 25 ns and the pileup events are expected to double. The target for integrated luminosity to be recorded up to the end of Run 3 is around 300/fb. The center of mass energy however depends on the success of run 2 conditions, the expectations is to be able to increase CM energy up to 14 TeV.

In the table (II.2) we show an approximative time-line for the LHC operations. We had also add on this table the prospects for High Luminosity LHC (HL-LHC). The CM energy for HL-LHC is expected to be 14 TeV.

period		$\sqrt{s}$ (TeV)	$\int L$ (fb $^{-1}$)	bunch spacing	<pileup>
2009-2012	Run 1	7 and 8	5+20	50 ns	21
2013-2014		Phase 1 upgrades			
2015-2017	Run 2	13	-	25 ns	25
2018-2019		Phase 1 upgrades - cont.			
2020 $\sim$ 2021	Run 3	14?	300	-	50
$\sim$ 2022		Phase 2 upgrades			
$\sim$ 2024	HL-LHC	14	3000	-	$\sim$ 130

Table II.2: Approximate time-line for LHC operations. The numbers followed by a question mark are the optimist expectations. The references are found along the text.

II.1.3 The multi-purpose experiments

The detectors of the CMS and ATLAS experiments are constructed on a cylindrical form along the beam pipe with the interaction point is on the center of the detector (see figure II.2). This detector topology allows maximal capacity to cover all the collision products.

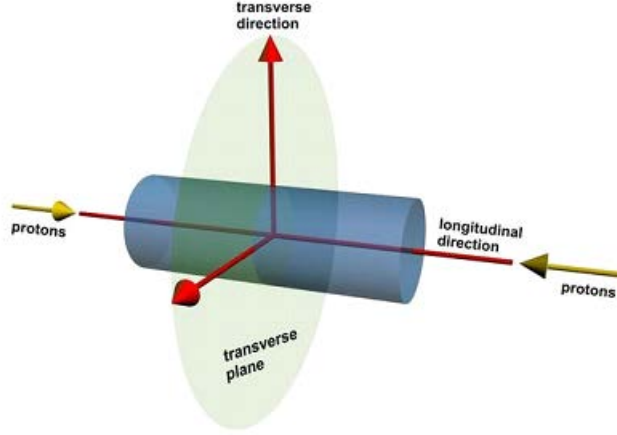


Figure II.2: Scheme for CMS and ATLAS detector format and their position with respect to LHC beam.

To study the collision products we use a coordinate system centered on the collision point. The azimuthal angle ϕ is the one that defines the plane transverse to the beam, and the polar angle θ is defined as angle between the cylinder transverse plane and the longitudinal (beam) direction. The pseudo-rapidity of objects leaving the collision point is $\eta = -\log \tan(\theta/2)$.

In parallel with the LHC time-line all the experiments schedule upgrades of their detector systems [121, 122] in order to be able to operate on the prospected LHC conditions, for repair of detectors, new installations and preparations to the new processing conditions both for at hardware and software level. At software level the increase of instantaneous luminosity is a challenge for trigger and pileup rejection methods. In chapter (III) we present a more extensive description of the CMS experiment, detector apparatus, sub-detectors coverage and software.

II.2 Jets

In the feasibility studies presented in this chapter we are interested in multi-jet final states. Jets are objects build out of the partons that are product of the interactions and are moving in the same direction. A jet is defined as a collection of particles clustered with a given algorithm. The jet clustering algorithms need to be robust against infrared and collinear radiation [123].

A jet clustering definition depend of two quantities: the distance between two sub-jets objects (d_{ij}) and the distance between one given object and the core of the jet (d_{iB}). Distance definitions can be balanced by the object p_T in a different manner:

$$d_{ij} = \min(p_{T,i}^{2p}, p_{T,j}^{2p}) \frac{\Delta_{ij}}{R} \quad (\text{II.3a})$$

$$d_{iB} = p_{T,i}^{2p} \quad (\text{II.3b})$$

where $\Delta_{ij} = (y_i - y_j)^2 + (\phi_i - \phi_j)^2$ is the angular distance between two objects, y_i is the rapidity of the parton i and R the resolution parameter that controls the extension of the jet. The different distances definitions (value of p) lead to different jet definitions. In this thesis we will make use of two of them:

- **Anti- k_t :** $p = -1$ (AkT), this means that the hardest the particle is, close it is from others. The clustering is done from the hardest particles from soft particles, these last ones do not modify the shape of the jets, while hard particles do. the most energetic sub-particles are clustered together first, this distance definition of make the jet definition more correspondent to the hard process and more resilient against soft background immersion. If no other hard particles are close by the algorithm will results on perfect cones [124].
- **Cambridge/Aachen:** $p = 0$ (C/A), the distances between objects are purely geometrical, the objects are clustered with angular ordering, allowing the formation of hierarchical substructure between clustered cores.

For nice reviews on details of the different jet definitions and their properties see [125], and also the lectures [126]. The AkT algorithm is commonly used to find resolved jets, while the C/A algorithm is more suitable for jet substructure studies [127].

The pile-up radiation is subtracted from the clustered objects using the jet-area method [116, 117]. This method is implemented in **FastJet** program [128]. After the jet area method subtraction there are still a fraction of reconstructed jets which are dominated by the pile-up energy, or are fake jets made of spurious overlap of soft particles produced in min-bias events (events triggered requesting a minimum activity related with inelastic scattering, like a minimum quantity of tracks). Particles produced in the pileup interactions can be collected by the jet clustering algorithm into objects with large enough p_T , known as a real pileup jet.

***b*-tagging**

The flavor of the jet can be inferred from the topology of jet constituents. The mass of *b*-quark is relatively large on comparison with the first and second generation of quarks (~ 5.7 GeV). The *b*-quarks decays weakly to a *c*-quark associated with a off shell *W* boson before hadronizes.

Therefore B hadrons (hadrons formed with a *b*-quark) are unstable therefore their lifetimes are usually longer than the lifetimes of unstable mesons that decay only hadronically (with the first and second generation of quarks). *B*-mesons have a flight distance of $\sim 450\mu\text{m}$ and defining a secondary vertex inside the jet. The weak decays of B hadrons have a high multiplicity of final tracks (~ 5 charged tracks).

Considering the weak decays approximately 11% of them are semi-leptonic ($\sim 20\%$ including cascade effects) [129]. Neutrinos are present in about the 35% of B hadron decays therefore the energy corrections due the presence of neutrinos in the final state are usually larger in *b*-quark initiated jets. Usually the analyses that explicitly rely on *b*-jet mass resolution construct a *b*-jet specific energy regression, see for example [130].

A *b*-tag algorithm is constructed to have as input experimental quantities based on the above considerations and output one real number between 0 and 1. Different techniques can be used to combine several experimental inputs in one unique discriminator like Multi-Variate (MVA) method, Lagrange multipliers or Neural Networks. The value of the *b*-tag discriminator allows us to define a Working Point (WP) from where the "*b*-tag" variable is considered as a boolean variable.

II.3 Topology of processes like $X \rightarrow Y Y \rightarrow z z z z$

We will exploit the topology of process chains like $X \rightarrow Y Y \rightarrow z z z z$, with X and Y being massive particles. We define the ratio:

$$r_M \equiv \frac{m_X}{2m_Y}. \quad (\text{II.4})$$

The boost of the Y particle from the rest frame of X particle to the rest frame of the Y particles depends only on the ratio of the X and Y masses, not of any absolute mass scale. If we calculate:

$$\beta_Y = |\vec{p}_Y|/E_Y = \sqrt{1 - 1/r_M^2}, \quad (\text{II.5a})$$

the boost factor is:

$$\gamma_Y = 1/\sqrt{1 - \beta_Y^2} = r_M. \quad (\text{II.5b})$$

The angular separation between two partons z coming from the decay of a boosted particle Y is given by equation (II.6b), see section 5.3 of [125]. If $(p_{T,Y} \gg m_Y)$:

$$m_Y \cong \varepsilon(1 - \varepsilon)p_{T,Y}\Delta R_{zz}^2, \quad (\text{II.6a})$$

resulting that:

$$\Delta R_{zz} \cong \frac{1}{\sqrt{\varepsilon(1 - \varepsilon)}} \frac{m_Y}{p_{T,Y}} \sim \frac{1}{\sqrt{\varepsilon(1 - \varepsilon)}} \frac{1}{r_M}, \quad (\text{II.6b})$$

where ε is the fraction of transverse momenta carried by each parton, for symmetric decays the split of momenta tend to be also symmetrical $\varepsilon \sim 1/2$. Using equation (II.6b) we estimate the partons of the the most energetic higgses (central in rapidity) start to merge on a cone of radius 0.5 if:

$$r_M \sqrt{1 - r_M^{-2}} \gtrsim 4. \quad (\text{II.7a})$$

Supposing the experimental capability of distinguish z -like objects requires those objects to be isolated by a cone defined by a 0.5 radius we should expect a signal composed by four well separated reconstructed z -like objects at $r_M \lesssim 4$ (resolved case). For $r_M \gtrsim 4$ we should be able to recognize an object composed by two merged z -like objects, what we call boosted scenario (see figure II.3), the jets formed by two prongs will be called fat jets.

Each mass regime has a best signal reconstruction solution. An important point for the fully resolved topology is the strategy to pair the z -like objects as coming from the decay of a Y . An intermediate region between boosted and resolved regimes can hold a non negligible portion of events with a mixed topology (where one of the Y particle is seen as a merged system of two z 's and the other Y is resolved). In the boosted and mixed topologies there is no z -object combinatoric problem to reconstruct the Y 's candidates. An object composed of two identical sub-objects is expected to have symmetrical internal structure.

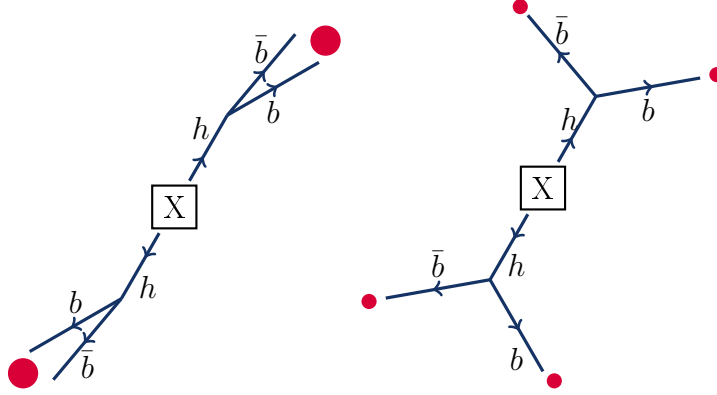


Figure II.3: Topology diagrams for the generic process $X \rightarrow 2Y \rightarrow 4z$. **Left:** boosted regime, corresponding to large values of the mass ratio $r_M = M_X/2M_Y$. **Right:** resolved regime, described by small values of r_M .

The common practice is to treat the boosted and resolved regimes with different analyses schemes, the intermediary region is eventually under-covered. In the next sections we propose a method that allow within the same analysis to interpolate resolved and boosted regime, maintaining uniform efficiency on the whole range of mass hypotheses. At parton level, and without any kinematical cut the proposed signal extraction is designed to be fully efficient on all range for m_X hypothesis. We refer to the method as scale invariant tagger.

II.4 $pp \rightarrow X \rightarrow h(b\bar{b})h(b\bar{b})$

We consider the feasibility of the LHC inclusive search of fully hadronic di-higgs resonances. The radion model is used as benchmark model for a scalar X resonance. We also check the angular sensitivity of our analysis applying it to the case where the heavy resonance is a spin 2 particle, a KK-graviton.

To validate our signal reconstruction strategy we simulated a generic process $pp \rightarrow X \rightarrow 2Y \rightarrow 4z$ using a toy MC simulation. We assume X to be produced at rest, X and Y to decay isotropically, like spin 0 resonances. In toy simulation any particle width is neglected. The main differences between the toy MC and matrix element generators are that the latter includes the full kinematics distributions for the X resonance, inclusive the proper angular distributions accordingly with spin hypothesis.

Both toy and **MadGraph5** signal events were showered and hadronized with **Pythia8**, version 8.170. At shower level events we consider Initial State Radiation (ISR), Final State Radiation (FSR) and multi-particle interaction (pileup). We refer to the showered events as being at hadron level simulation. We also study the topology of signal events without the showering process, we will refer to this level of simulation as parton level. The events after shower but without hadronization will be referred as hadron level. We analyze both parton and hadron level events using a **FastJet** routine in the toy MC events. Parton and hadron level events were clustered with anti- k_t algorithm [124] with radius parameter 0.5. Most experimental multi-jet analysis in CMS use $R = 0.5$ and in ATLAS use $R = 0.4$ [131, 132].

We do not add any kinematical cuts to study the efficiency reconstruction of parton level toy events. To study the efficiency of reconstruction of hadron level simulation however we use very mild cuts on minimum transverse momenta and maximum rapidity $|y|$ of each jet and of the sum of the transverse momenta of all the jets (H_T). We list the basic cuts in equation (II.8).

$$p_T > 25 \text{ GeV}, \quad |y| < 5, \quad H_T > 100 \text{ GeV}. \quad (\text{II.8})$$

In figure (II.4) we show the event fractions for different final state jet multiplicities for both parton and hadron level simulation.

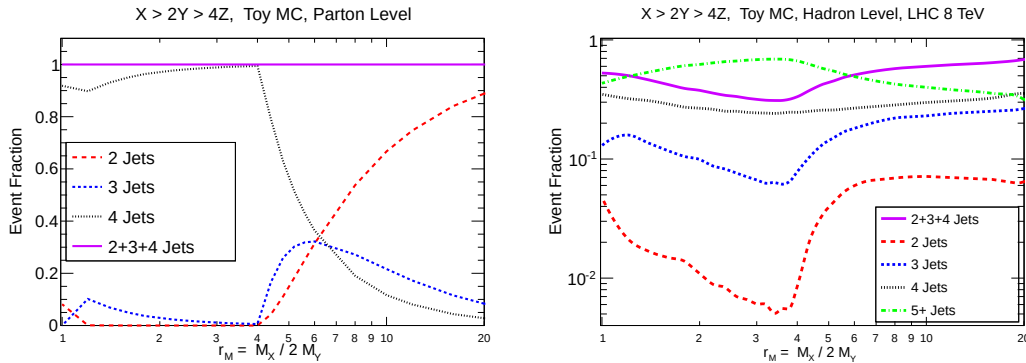


Figure II.4: **Left:** the event fraction of reconstructed jets as a function of resonances mass ratio r_M equation (II.4) for parton-level toy Monte Carlo events. **Right:** the same at hadron-level, under the basic cuts of table (II.8). By construction parton level the only allowed topologies are two-, three- and four-jet events, so their sum adds to to the total number of events.

At parton level we can identify the three different regimes we discussed before: boosted, resolved, and mixed. By construction, events can only have between two and four jets, the topology regimes are easy to identify.^{*} The separation of regimes happens around $r_M \sim 4$, as anticipated by equation (II.7a).

At hadron-level both the shower and the kinematic cuts makes the separation of regimes to be no longer obvious. The basic cuts we applied can suppress some of the jets from the hard process, while QCD radiation enhance signal jets multiplicity. Therefore we cannot classify regimes based on jet multiplicity. In the next section we detail an analysis strategy that define and interpolate kinematic regimes. Still more than 50% of the events have between two and four jets for any value of r_M .

II.4.1 Reconstruction algorithm

We base the reconstruction algorithm for event selection on the number of jets per event by the presence of jets with non-trivial substructure and high invariant mass. Each one of the two hardest jets in the event is re-clustered using the C/A algorithm. We perform the re-clustering with a big radius $R_{sj} = 1.1$. To be sure all particles coming from radiation of the hard event are included on the jet. The last step of C/A clustering for jet j is undone, resulting in sub-jets (prongs).

Hard QCD radiation (quark and gluon initiated jets) can also present jets with high invariant mass $\sim O(100)$ GeV, that contain a given number of prongs from early gluon and/or early quark radiation, faking a fat tag. A QCD splitting like $q \rightarrow qg$, for example, is mainly asymmetric. The distinction between QCD jets and jets coming from boosted heavy bosons can be done using the informations about the number of prongs and the relation between their invariant masses.

Boosted top quarks form one jet with tree hard prongs ($t \rightarrow W(q\bar{q})b$) boosted W , Z or higgs bosons form a jet with two nearly symmetric hard prongs. Therefore our signal is characterized by the presence of two symmetric hard prongs. We identify those as j_1 and j_2 , with $m_{j_1} > m_{j_2}$ and process with the BDRS mass-drop tagger [127]. The BDRS tag relies in two conditions:

1. A minimum cut as bellow assure the prongs not to be too asymmetric[†]:

$$y = \frac{\min(p_{t,j_1}, p_{t,j_2})^2}{m_j^2} \Delta R_{j_1,j_2}^2 > y_{\text{cut}}, \quad (\text{II.9a})$$

where $\Delta R_{j_1,j_2}$ is the angular separation between the two sub-jets.

2. The more massive prong is expected to be significantly lighter than the parent jet by a factor μ . This condition is known as mass-drop:

$$m_{j_1} \leq \mu \cdot m_j \quad (\text{II.9b})$$

^{*}At $r_M \sim 1$ the two higgses have very small boost, its decay products tend to be back-to-back. The probability we see some fraction of events with 2 and 3 jets due coincidences is bigger.

[†]When $p_{t,j_1}, p_{t,j_2} \gg m_j$ we have that the rapidity $y \sim \min(p_{t,j_1}, p_{t,j_2}) / \max(p_{t,j_1}, p_{t,j_2})$, see also equation (II.6a).

If one of the two conditions are not satisfied we replace j with j_1 and apply the unclustering to the new j , repeating the procedure until we find a sub-jet for which the mass-drop and asymmetry conditions are both satisfied. In this work we use ($\mu = 0.67$) and $y_{\text{cut}} = 0.09$ as in the original BDRS paper [127]. When j is reduced to a single particle with no success on the requirements the process stops without tagging the jet.

The jet candidates that survives these two requirements are fat tagged. We classify events by the quantity of tags: events with **2-tags** are supposed to belong to boosted regime. Events with **1-tag** are supposed to be on mixed regime, and events with **0-tag** are supposed to be on un-boosted regime. The two higgses candidates of the un-boosted regime are constructed from the pairs holding the minimum difference on di-jet invariant mass, among the four hardest jets.

Once the first event classification is done we apply quality requirements to recognize the events as a di-higgs event. We apply mass and angular quality requirements to tag the two higgses according with the regime. These angular requirements are important to background rejection in all the regimes. The mass resolution and general angular quality requirements we use are:

1. We require the masses of the two h candidates to be the same up to a given mass tolerance f_m , to account for experimental mass resolution, as well as mass smearing due to underlying event, hadronization, out-of-cone radiation and radiation lost on neutrinos emission:

$$\left| \frac{m_{H_1} - m_{H_2}}{\langle m_H \rangle} \right| \leq f_m, \quad (\text{II.10a})$$

where $\langle m_H \rangle$ is the average mass of the two reconstructed h resonances. In realistic analysis the mass resolution depends on mass scale. Generically the mass resolution scale like $f_m \sim M^{-1/2}$ where M the jet mass. The detector mass resolution in di-jet systems is $\sim 15\%$ for a large range of jets p_T [131, 132, 133]. We assume a fixed mass resolution tolerance of $f_m = 15\%$ to all mass hypothesis.

2. The masses of the two h candidates must lie in a mass window around m_H , where the width of the window, determined by the mass resolution of the detector.

$$m_H(1 - f_m) \leq m_{H_1}, m_{H_2} \leq m_H(1 + f_m). \quad (\text{II.10b})$$

3. We require separation in rapidity of the two h candidates to be smaller than some upper value,

$$\Delta y \equiv |y_{H_1} - y_{H_2}| \leq \Delta y_{\text{max}} \quad (\text{II.10c})$$

We will take $\Delta y_{\text{max}} = 1.3$, this value was optimized from the high mass di-jet searches at the LHC [131, 132].

The angular distribution of scalar resonances is flat in the decay angle between the higgses in the rest frame of the X resonance, common called $\cos\theta^*$. This angle can be written in term of the rapidity gap between the higgses candidates (equation II.10c) as:

$$\cos\theta^* = \frac{\exp(\Delta y) - 1}{\exp(\Delta y) + 1}, \quad (\text{II.11})$$

the cut in $\Delta y_{\max}$ directly determines the value of the angle $\cos\theta_{\max}^*$ that will lead to resonance tagging.

The expected tagging efficiency on very boosted **2 tag** case at parton level can be anticipate as the product between signal cut acceptance under the y_{cut} and the cosine of maximum polar angle a tagged higgs can reach on rest frame of the X resonance $\cos\theta_{\max}^*$:

$$\epsilon_{2\text{-tag}}^{\text{lim}} \equiv \epsilon_{2\text{-tag}}(r_M \gg 1) = \left(1 - \frac{2y_{\text{cut}}}{1 + y_{\text{cut}}}\right)^2 \cos\theta_{\max}^*. \quad (\text{II.12a})$$

In the very boosted limit:

$$\epsilon_{2\text{-tag}}(r_M \gg 1) \sim 0.4. \quad (\text{II.12b})$$

The CA jet definition used to find the fat jet prongs is not resilient with respect to soft radiation from pile-up. To improve mass resolution of boosted higgses candidates, and make the higgs-tagging more resilient we filter the components of the mass-drop tagged jet [127, 134]. The filtering process consists on re-cluster the constituents of each tagged jet with a radius $R_{\text{filt}} = \min(\Delta R_{\text{sj,sj}}/2 R_f)$, where $\Delta R_{\text{sj,sj}}$ is the angular distance between the two sub-jets after mass-drop and $R_f = 0.3$. To account for at least one QCD emission three hardest sub-jets are retained ($n_{\text{filt}} = 3$).

In the characterization of a boosted higgs candidate the BDRS mass-drop tag (equation II.9) implies the tagged jet passed rapidity and asymmetry cuts. We want to assure same conditions hold for the three different possible topologies, and therefore similar efficiencies. Special requirements (inspired on mass drop tagger) for the two jets tagged as from a resolved higgs candidate are applied also to the **1 tag** and **0 tag** samples:

1. We prevent big energetic asymmetry inside of the constituents. We require the p_T of the jets from the higgs candidate to be:

$$p_T^{H_2} \geq y_{\text{cut}} \cdot p_T^{H_1}, \quad (\text{II.13a})$$

where $p_T^{H_1}$ and $p_T^{H_2} \leq p_T^{H_1}$. This cut plays a similar role as the asymmetry requirement in the BDRS mass-drop tagger, it helps rejecting events where a soft jet arises from FSR.

2. By the same motivation we impose also a requirement for the invariant mass of each sub-jet candidate

$$\max(m_{H_1}, m_{H_2}) \leq \mu \cdot m_H, \quad (\text{II.13b})$$

where μ is the same parameter as in the BDRS mass-drop tagger.

3. The rapidity of the two constituents of the candidate ($y_{Y_{i,1}}$ and $y_{Y_{i,2}}$, with $i = 1, 2$) must be smaller than some upper value:

$$\Delta y \equiv |y_{H_1} - y_{H_2}| \leq \Delta y_{\max}^{\text{res}} \quad (\text{II.13c})$$

As it was done in fully boosted scenario (equation II.12b) we also can estimate the parton level expected signal efficiency fully resolved case ($r_M \sim 1$) given the maximum rapidity separation of the jets from each higgs decay. Considering the impact of only this cut the signal efficiency is given by:

$$\epsilon_{0\text{-tag}}(r_M \sim 1) = \left[\frac{\exp(\Delta y_{\max}^{\text{res}}) - 1}{\exp(\Delta y_{\max}^{\text{res}}) + 1} \right]^2. \quad (\text{II.14})$$

We determine $\Delta y_{\max}^{\text{res}}$ value cut demanding that the efficiencies at low r_M match the asymptotic large r_M value, resulting that:

$$\Delta y_{\max}^{\text{res}} = \ln \left(\frac{1 + \sqrt{\epsilon_{2\text{-tag}}^{\text{lim}}}}{1 - \sqrt{\epsilon_{2\text{-tag}}^{\text{lim}}}} \right) \sim 1.5, \quad (\text{II.15})$$

Given the tag elements presented we construct the analysis chain. The analysis chain has been implemented in a code based on **FastJet3**.

- If the two hardest jets in the event are found fat, we test both under the quality requirements described on last section. In case both fat jets satisfy the quality requirements of section (II.10) the event is assigned to the **2-tag**
- If one of the two fat jets do not pass quality requirements the events with fewer than three jets after cuts are discarded. We then apply (II.10) to check if the event belongs to tree-jet topology. If both candidates are tagged as higgs (one boosted and one resolved), the event is classified into the **1-tag** sample
- If the event is not assigned as **2-tag** or **1-tag** sample, we check the event can be classified as being two resolved higgses. Events with fewer than four jets passing the basic cuts are discarded. We select the jet pairing such that the combination ij and kl of the five hardest jets in the event leads to di-jet masses m_{ij} and M_{kl} that minimizes the difference $|m_{ij} - m_{kl}|$, and use this pairing to reconstruct the two higgs candidates among the five hardest jets*. If the reconstructed higgs candidates pass satisfy (II.10), the event is classified as belonging to the **0-tag** sample.

The asymmetry cut of equations (II.10c, II.13c) are an important element to higgs tagger in all the regimes of the presented analysis. Some other details of the resonance strategy however can be modified without affecting the general philosophy of the work. See [135] for a recent phenomenological comparison of different fat-taggers.

*Often one gluon from large-angle initial-state radiation can have a larger p_T than the four original b -jets. We check this choice improves the efficiency of the resolved configuration as compared to selecting only the four leading jets.

II.4.2 Signal efficiencies in toy signal events

We study the performance of events selection on toy signal. In figure (II.5) we show the efficiency of the resonance resulting of the pair tagging algorithm described above as a function of resonances mass ratio r_M . We show both the total efficiency and the breakup of the efficiencies corresponding to the **2-tag**, **1-tag** and **0-tag** samples. For the validation toy results we use the loose selection cuts of table (II.8). The impact of those loose cuts in the final efficiencies are negligible if $r_M > 1$. Therefore differences between parton and hadron-level arise uniquely from ISR and FSR. The behavior at 14 TeV center of mass is very similar.

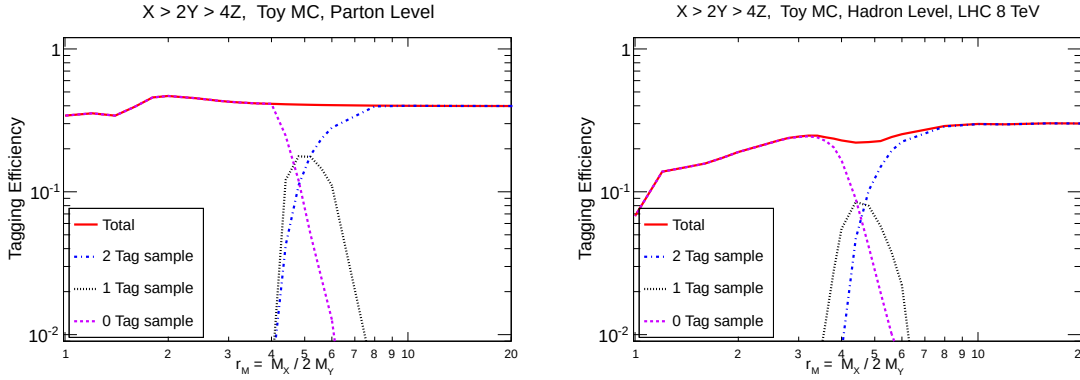


Figure II.5: The efficiency of the resonance pair tagging algorithm as a function of resonances mass ratio r_M for toy MC simulation. We show both the total efficiency and the break-up for the **2-tag**, **1-tag** and **0-tag** samples. **Left plot:** Parton level simulation, no cuts have been applied.. **Right plot:** Hadron-level simulation under the jet selection cuts of table (II.8).

By construction the total parton level reconstruction efficiency is found to be rather flat in all the mass range, between 30% and 40% for all mass values. This result agrees well with we had estimated on section (II.10). At low r_M , the **0-tag** sample dominated the total efficiency, while for large r_M the 2-tag sample takes all total efficiency. The **1-tag** sample takes a non negligible part of efficiencies at intermediate boosts.

For hadron level events the total reconstruction efficiency are lower, and the split on regimes is not trivial. The contamination from QCD radiation and pileup lower down the efficiency of the BDRS-tagger, and consequently lower down the efficiency of hadron level simulation on the boosted regime (**2-tag** sample) with respect to the parton level estimation. The relative loose of efficiency is proportional to r_M value. The **1-tag** sample and the low r_M **0-tag** sample are the most affected by the showering procedure. These samples are on the edges between regimes therefore they are the most sensitive to QCD radiation.

We obtain a reasonable flat tagging efficiency of between 20% and 30% for hadron level simulation on all topologies, except close to $r_M = 1$ (where the jets pairing is not trivial due the absence of any boost). Given the non triviality of hadron level events it is remarkable that efficiencies and the split of regimes between parton and hadron level signal is comparable and the differences are well understood.

II.4.3 Dependency on the jet radius

We test the robustness of the jet reconstruction strategy against different jet radius definition. Of course the relative fraction of **2-tags**, **1-tag** and **0-tag** events will vary with R , but, giving the scale invariance character of the pair tagging algorithm their sum should be stable. In figure (II.6) we show the total efficiency in parton and hadron-level events for three different radii, $R = 0.3, 0.5$ and 0.8 .

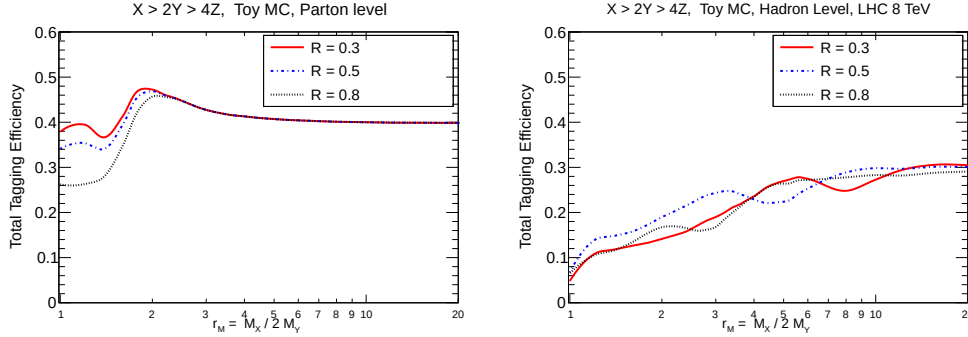


Figure II.6: The total tagging efficiency for parton-level events (left plot) and for hadron-level events (right plot) for different values of the jet radii R as a function of the mass ratio r_M . The default radius used in this paper is $R = 0.5$.

The efficiency reconstruction for the parton level events in the boosted regime is indeed R -independent, for low r_M the probability of 2 partons to end on the same jet (mis-tag) increases with the radius definition of the jet clustering. Also by the fact the higgs bosons are predominantly at rest (if small QCD radiation), in this configuration it is quite frequent to have a confusing pairing of the original jets already at parton level. On top of this, at hadron level, large-angle ISR can appear as additional jets also cause confusion on jets pairing. The total efficiency is considerably independent of the jet radius definition. It reveals the robustness of event classification to the physics characterization.

II.4.4 Gravity particles signal

We apply the pair tagging algorithm for the gravity particles benchmarks. The signal benchmarks generated by **MadGraph5**. To be more realistic we use more severe selection cuts, inspired by the typical trigger and angular acceptances of the LHC experiments [136, 137].

$$p_T > 50\text{GeV}, \quad |y| < 2.5, \quad H_T > 300 \text{ GeV}. \quad (\text{II.16})$$

The comparison of hadron level tagging efficiencies between the toy Monte Carlo events and the **MadGraph5** radion and KK-graviton events, at LHC 8 TeV as a function of the mass ratio r_M can be seen in figure (II.7). The toy MC results agrees very well with the radion events*.

*We have generated fewer mass points with **MadGraph5** than with the toy MC, hence the somewhat less smooth distributions in the former case.

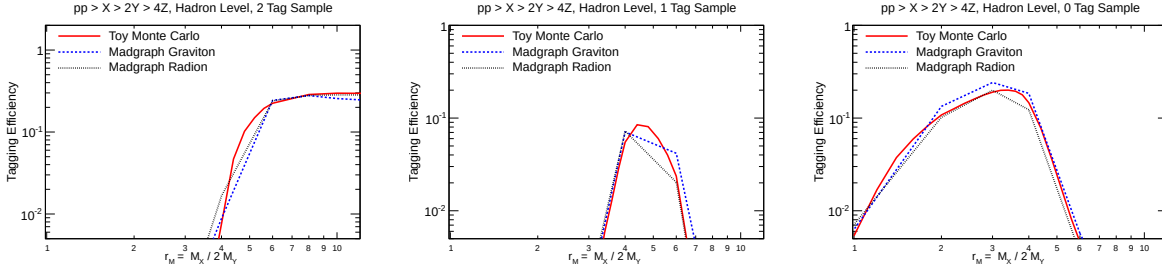


Figure II.7: Comparison between the hadron-level tagging efficiencies, at the 8 TeV LHC, for the toy Monte Carlo events and the radion and graviton MadGraph5 events. We show the total efficiencies and the breakup in different tagged samples as a function of the boost factor r_M .

We do not apply strong angular cuts along the analysis chain therefore we do not expect to be sensitive to spin. Indeed the efficiencies for the radion and graviton are very similar. The spin-zero vs. spin-two angular distributions do not lead to any large differences at the level of the reconstruction.

II.4.5 b -tagging estimation

We introduce b -tag by weighting events according to some requirements. We had adopted requirements inspired by the ATLAS and CMS capabilities [138, 139, 140, 141]: the probability of tagging a b -quark is taken to be $f_b = 0.75$, the mistag probabilities of c -quarks, $f_c = 0.10$, and of light quarks and gluons, $f_l = 0.03$. These requirements are expected to be realistic, we will see in section (III.3) that actually they are too conservative.

In our MC simulation we only shower the events but do not hadronize events. To be able to estimate the impact of b -tag in the search we choose to not decay b -quarks. The final state of the investigated channel includes four b -quarks. We obviously want to be able to estimate background rejection power gained by using this information.

We estimate b -tag performance by defining b -quark properties to be detector-taggable, both on p_T and angular range. We require only one b -tag per higgs candidate. A b -quark is considered taggable if it contains at least one b quark with $p_{T,b} \geq p_{T,b}^{\min} = 10$ GeV. A c quark with $p_{T,c} \geq p_{T,b}^{\min}$, is considered mis-taggable.

We first determine the number of b -quarks within each of the two jets of one higgs candidate. If at least one taggable b -quark is found, the b -tag efficiency is denoted by f_b . If no taggable b -quark is found in the candidate, but we find at least one c quark within mistag requirements, the b -tag efficiency is denoted by f_c . If none of the condition above are fulfilled (higgs candidate jet which contains only light quarks and gluons) the event is b -tagged with a mistag probability f_l . If the two higgs candidate's jets each contain at least one b quark, the event is assigned a weight $f_b^2 = 0.56$.

In figure (II.8) we see results for the b -tagging efficiencies on graviton mediated higgs pair production at the LHC 8 TeV. To consider b -tagging re-weighting decreases signal efficiency by $\sim 50\%$.

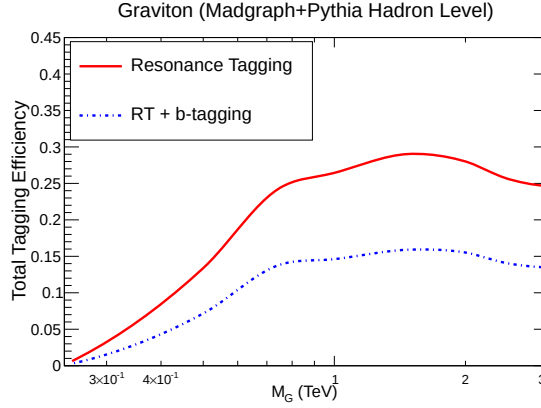


Figure II.8: The b -tagging efficiencies for graviton mediated Higgs pair production at the LHC 8 TeV. We show the efficiency of the resonance tagging (same as in the upper right plot of Fig. II.7) together with that including the b -tagging.

We could had also considered a more optimistic scenario for the b -tagging, in which each higgs candidate is required to have two b -tags. In particular CMS has demonstrated the ability to tag pairs of B -hadrons even for angular separations $\Delta R_{b\bar{b}} < 0.4$ [138], which suggests that this scenario could be viable also in the highly boosted limit where the B -hadrons are within a single AkT jet [141]. We found this requirement to be very severe for the **0 tag** signal sample.

II.4.6 Background rejection

We estimate background rejection according to the mass hypothesis m_X . As background we consider only the SM multi-jet final states, expected to be dominant. We produced a large sample of QCD multi-jets with **Pythia8** [142], starting from di-jet configurations and with the shower radiation taking care of generating the higher-order jet topologies *. The resulting hadron level events are processed through the same analysis chain as the signal events. One example of how QCD radiation can mimic the conditions for resonance tagging is when fake mass drops can be generated from a sufficiently symmetric splitting of a quark or gluon [145].

The kinematics of massless jet pair production determine that the mass of the di-jet will be given in terms of the p_T of the jets and their rapidity separation in the laboratory frame Δy by $m_X = 2p_T \cosh(\Delta y/2)$. Therefore, to properly cover all phase space the generation cut for QCD di-jets should be at least:

$$p_T^{\min} \sim \frac{m_X}{2} \frac{1}{\cosh(\Delta y_{\max}/2)}, \quad (\text{II.17})$$

for any candidate resonance mass m_X .

***Pythia8** is known to underestimate the amount of QCD multi-jet topologies by a factor up to two with respect to experimental data [131], for the accuracy requested to a feasibility study we had considered this precision to be sufficient. Similar results have been obtained with the **Alpgen** parton level event generator [143] matched to **Pythia8** using the MLM matching [144].

In order to achieve an efficient QCD multi-jet event generation we demand that the invariant mass of the system to be in a window around the hypothesis and the p_T cut:

$$m_X(\text{system}) \subset [m_X(1 - f_m), m_X(1 + f_m)] . \quad (\text{II.18})$$

In the four jet configuration, it is reasonable to require also a minimum p_T value of half of the one above. We are using a rapidity cut of $\Delta y_{\text{max}} = 1.3$, therefore we find that the Pythia8 minimum p_T in generation should be $p_T^{\text{min}} \sim m_X/5$. We had explicitly verified that the QCD di-jet cross-section is not modified if looser generation cuts are adopted. No generation cut in rapidity is applied.

In figure (II.9) we show the QCD di-jet cross section obtained from the Pythia8 multi-jet sample at LHC 14 TeV.

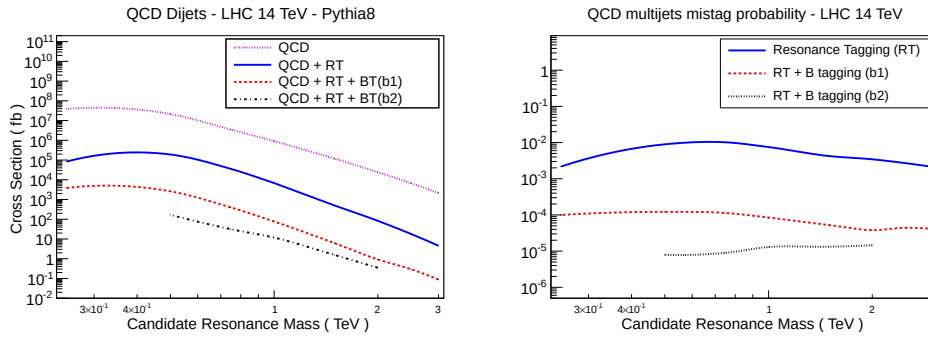


Figure II.9: **Left:** the QCD di-jet cross section before and after the events are processed through the resonance tagging algorithm, for LHC14. **Right:** Total probability of resonance tagging to mis-tag background, as a function of the mass point M . The dot-dashed lines (b2) correspond to require two b -taggable quarks on each higgs candidate.

To consider b -tagging reduces the mis-tagging probability of events in about 10^{-3} to 10^{-4} . To require two b -tags decreases one order of magnitude more the background rejection. In figure (II.10) we can see the efficiency of mistag splitted in regimes. Background events split on tag categories qualitatively similar to what was seen for the signal, also similar to signal total mistag probability is approximately scale invariant.

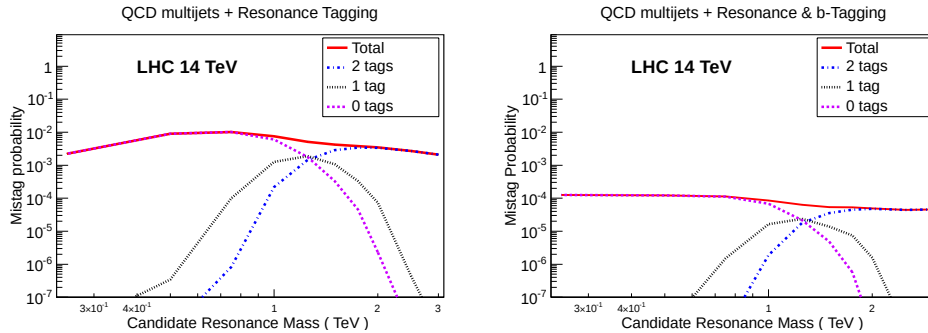


Figure II.10: Decomposition of the probability of resonance tagging to mistag background, without (left) and with (right) b -tagging, as a function of the mass point M on LHC14. The mistag probability of QCD di-jets at LHC8 is very similar.

II.4.7 LHC sensitivity

To derive model independent bounds on BSM scenarios with enhanced higgs pair production in the $4b$ final state we use a counting approach (see appendix D). We combine the results of the signal efficiencies and the multi-jet background rejection of the resonance tagging algorithm. We assume a total integrated luminosity of $\mathcal{L} = 25 \text{ fb}^{-1}$ at 8 TeV and of 500 fb^{-1} at 14 TeV. We also assume branching fraction to be SM like.

In figure (II.11) we show LHC sensitivity (95% confidence level exclusion limit) to $\sigma(pp \rightarrow X)\text{BR}(X \rightarrow HH)$ derived out of the algorithm efficiencies.

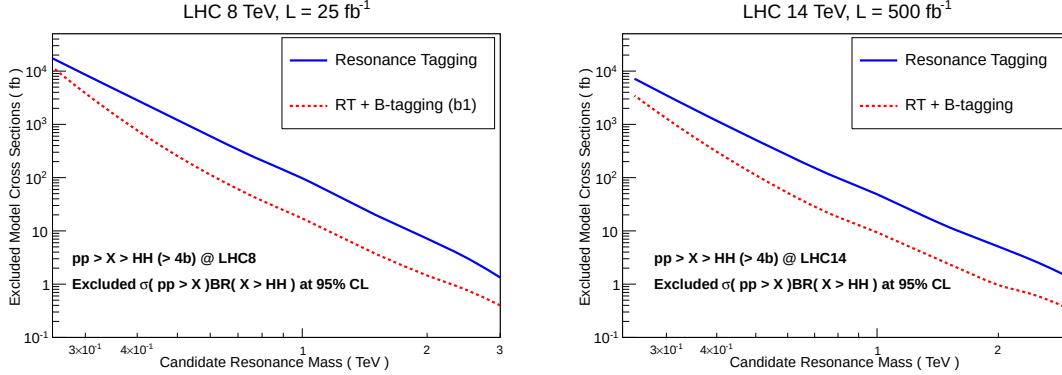


Figure II.11: The model independent exclusion ranges in $\sigma(pp \rightarrow X)\text{BR}(X \rightarrow HH)$ at the 95% confidence level for the production of a heavy resonance X which then decays into a higgs boson pair, which decay subsequently into a bottom-antibottom pair. The left plot show the exclusion ranges at LHC8 with $\mathcal{L} = 25 \text{ fb}^{-1}$ while the right plot corresponds to 14 TeV with $\mathcal{L} = 500 \text{ fb}^{-1}$.

Using the resonance tagging algorithm and our naive b -tag estimation we predict to LHC 8 TeV (14 TeV) to be sensitive to cross sections as small as 200 fb (50 fb) for $m_X \sim 500 \text{ GeV}$ at, while at higher masses, $m_X \sim 2 \text{ TeV}$, the $4b$ final state is sensitive to cross sections smaller than 1 fb at both energies. The cross section increase is the element of the improvement in exclusion reach with the center of mass energy (to 8 TeV to 14 TeV).

At high mass although, the predicted sensitivity to fully hadronic di-higgs channels is very similar on both center of mass energies investigated (8 TeV and 14 TeV). The increase in luminosity and cross section is compensated by the corresponding increase of the high mass QCD di-jet cross sections when going from 8 to 14 TeV. In the 14 TeV center of mass scenario however the increase of expected signal would allow the application of more restrictive selection to background rejection and use more involved technique.

After the publication of [146] other two references showing similar developments appeared [147, 148]. Recently the ATLAS collaboration had released results on the $hh \rightarrow b\bar{b}b\bar{b}$ channel using 19.5/fb of 8 TeV data, considering only the resolved event topology. The benchmark model it was used was a very narrow KK-graviton (bulk scenario) [149]. The spin hypothesis of the benchmark signal however is not relevant to the gluon fusion channel. The total efficiency reached for signal on request four b -tagged jets up to 1 TeV is around 10%, not too far from what we had predicted (see figure II.8). As expected the signal reconstruction efficiency falls above the 1 TeV hypothesis (as it will be shown in figure III.24).

The $t\bar{t}$ contribution is a non-negligible of the smooth background, between 10 and 20% according to the presence or not of a veto. The four body spectra shape in the falling regime is more steep for this component than the multi-jet QCD component (see section III.3.2).

For the signal extraction in a real analysis we need to define a strategy for background estimation. Both CMS and ATLAS signal strategies for resolved case defines a data driven method which extract the amount of background from the signal depleted region, so called *sideband method*. After the background estimation the signal selection efficiency drops more than 50%. The 2 b -tag signal sample is also considered.

The final expected sensitivity for a $hh \rightarrow b\bar{b}b\bar{b}$ signal obtained by ATLAS collaboration however are very similar to what we had predicted up to 1 TeV hypothesis, and worsen for bigger mass hypothesis (see figure II.11).

II.5 $pp \rightarrow X \ jj \rightarrow h(b\bar{b})h(b\bar{b}) \ jj$

In this section we investigate the resonant double higgs production by vector boson fusion mechanism. Analogously to last section we concentrate in the final state where both higgs bosons decays into b-quarks, leading to a $b\bar{b}b\bar{b}jj$ final state. In all the benchmark models of a di-higgs resonance described on section (I) the WBF resonance production is dominant over the resonance production associated with weak bosons and sub-dominant with respect to gluon fusion production mode. The hierarchy of cross sections between gluon fusion and WBF production mode is $\sim O(10)$ for the bulk KK-graviton hypothesis and $\sim O(100)$ for the RS1 and bulk radion hypothesis.

Differently from the inclusive gluon fusion production in the WBF production the topology of the events final states depends on the model hypotheses. We perform our pre-feasibility studies on the resonant WBF signal of a bulk KK-graviton resonance, assuming 1 GeV total width. For sake of final comparison on the WBF jets topology only, we perform the same simulation under RS1 scenario.

We want to be more precise than the last section and aim to study also the detector effects on the signal extraction. We target our studies to the conditions of run 2 of LHC. We also plan to study the impact of pileup increase on run 3 and HL-LHC conditions. The best we can do in this direction outside the collaborations is to use the public simulations tools compatible with have been accorded between the LHC and the experiments [150]. The main motivations to the non-experimental community to perform studies to prospect the LHC sensitivity for NP theories based on fast detector simulation is to be able to do realistic predictions for feedback theories model building and also learn the detectors capability and run conditions that can favor signal detection to feedback the upgrade studies.

For this study we had generated 20k of signal events using the parameter $\tilde{k} = 0.1$, and varying the bulk KK-graviton mass m_{Gr} from 260 to 1450 GeV. Given the PDF suppression on the heavy quark flavors, to consider only quarks with u and d flavor as initial final state is sufficient to cover the WBF signal topology.

II.5.1 Parton level study

A pre-feasibility study of the signal, with the same spirit of the last section was made to setup an analysis strategy. We looked at the topology of hard events at matrix element level only and construct one unique baseline analysis with the goal of being able to distinguish two higgses and two (WBF-like) forward jets for a wide range of resonance mass hypothesis. The higgses were decayed to b -quark pairs on *Pythia8*, therefore no generation cut is assigned to the b -jets of signal.

We clustered the "jets" with the AkT05 algorithm using *FastJet*. This procedure makes sense even at a parton level preliminary analysis since two partons that are close enough can cluster on a jet. Following [146] we based the reconstruction algorithm for event selection in the number of jets per event.

For the preliminary parton level study we consider only events with four b 's and assuming perfect b -tag performance. In the analysis flow we require at least two jets in the event after b and fat jet-tagging. The WBF jets are chosen as the pair that holds the largest invariant mass among the non-tagged jets. We need to assume that the b -quarks have the minimal transverse momentum and angular requirements to a jet to be tag-gable as a b -quark, in this section we use the following cuts:

$$p_T^b > 20 \text{ GeV}, \quad |\eta_b| < 2.5, \quad (\text{II.19})$$

the basic generation cuts we use on the WBF jets in this section are:

$$p_T^j > 10 \text{ GeV}, \quad m_{jj} > 100 \text{ GeV}, \quad |\eta_j| < 5, \quad \Delta R(jj) > 0.01. \quad (\text{II.20})$$

We follow closely section (II.4) and classify events by the number of mass-drop tagged jets (fat jets) under the same parameter definitions (equation II.9). For the fat tag in this section we require also:

$$m_{\text{fat jet}} > 120 \text{ GeV}. \quad (\text{II.21})$$

As pre-selection we used the following cuts in the WBF jets:

$$p_T^j > 30 \text{ GeV}, \quad m_{jj} > 400 \text{ GeV}, \quad |\eta_j| < 4.7. \quad (\text{II.22})$$

After this first selection we reconstruct the higgses from the remaining jets. We also take advantage of the fact that the di-higgs system is boosted by using the following cuts:

$$m_{hh} > 250 \text{ GeV}, \quad p_T^{hh} > 60 \text{ GeV}, \quad \Delta\eta(hh) < 2 \quad (\text{II.23})$$

We construct the analysis flow as:

- If the two hardest jets in the event are found fat and satisfy the di-Higgs quality requirements of equation (II.23) the event is assigned to the **2-tag**.
- If one of the two fat jets do not pass quality requirements or only one of the jets among the two hardest is fat, then events with fewer than three jets after cuts are discarded. If the event passes the selection (II.23) it is assigned to **1-tag** sample.
- If the event is not assigned as **2-tag** or **1-tag** sample, we check the event can be classified as being two resolved Higgses. Events with fewer than six jets passing the basic cuts are discarded. To select the higgses candidates we pair the four non-WBF jets i, j, k, l minimizing the mass difference $|M_{ij} - M_{kl}|$. If the two reconstructed Higgs candidates satisfy the cuts (II.23) the event is classified as belonging to the **0-tag** sample.

The QCD production of four b -quarks and two additional is the background with the largest cross section. These events were generated with Alpgen [143]. We had also considered the electroweak background producing Z bosons pairs decaying to b -quarks in association with jets, generated with MadGraph5. Finally, we had also considered the SM WBF production of higgs boson pairs with the higgs boson decaying into b -quarks. The SM di-higgs pair production does not have a resonant structure, generated with MadGraph.

Table (II.3) summarizes the different samples used, along with the expected cross-sections and number of generated events.

Process	Generator	σ (pb)	Number of events
QCD $b\bar{b}b\bar{b}jj$	AlpGen	148	500k
EWK $Z(b\bar{b})b\bar{b}jj$	MadGraph	1.46	200k
EWK $Z(b\bar{b})Z(b\bar{b})jj$	MadGraph	0.06	210k
WBF SM $H(b\bar{b})H(b\bar{b})jj$	MadGraph	2.71×10^{-4}	50k

Table II.3: Background samples and the corresponding cross sections used for the analysis.

The cuts performed at the generator level on background generation were:

$$p_T^j > 20 \text{ GeV}, \quad |\eta_j| < 5, \quad \Delta R(jj) > 0.3, \quad (\text{II.24a})$$

for the WBF jets system. For the b -jets system we use:

$$\Delta R(bj) > 0.3, \quad p_T^b > 20 \text{ GeV}, \quad |\eta_b| < 2.5, \quad \Delta R(bb) > 0.1. \quad (\text{II.24b})$$

The di-jet invariant mass cut design to favor the WBF topology ($M_{jj} > 100$) was applied only to the electroweak background, that was generated with **MadGraph5**. It is not straightforward to add such a cut on **AlpGen**. The QCD **AlpGen** background was used only for the preliminary parton level results, to be on time to participate the 2013 Les Houches proceedings.

Figure (II.12) shows the pseudo-rapidity difference among the two WBF jets of KK-graviton production on both RS1 and bulk scenario.

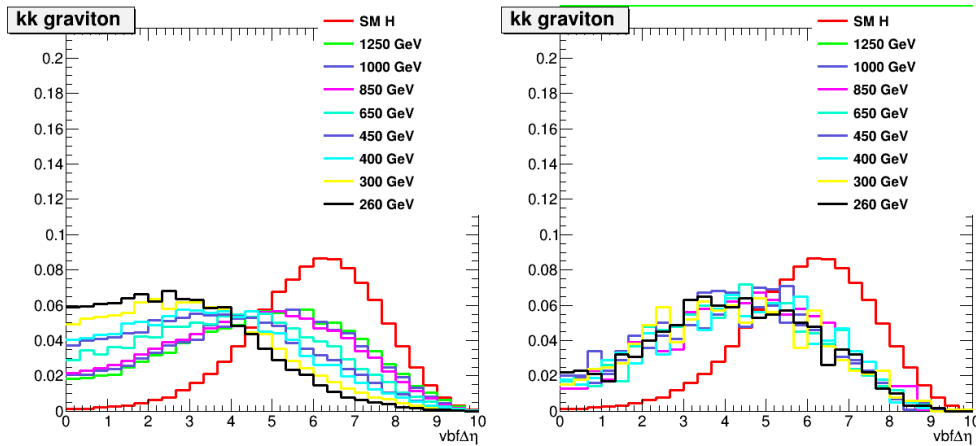


Figure II.12: Separation in pseudo rapidity between the WBF di-jet system on WBF production of KK-graviton for LHC13 after the generation cuts. **Left:** Bulk scenario. **Right:** RS1 scenario. The red curve corresponds to SM single higgs WBF production, this curve is the same on both distributions. Distributions are shape normalized.

The kinematics of the WBF jets in the signal differ from the case of the Standard Model WBF higgs production, which tend to be more central. It was already noted that the behavior of WBF jets reminiscent of the on shell production of a resonance depends of the spin of such particle [151, 152, 153].

The KK-graviton couplings with massive weak bosons have different p_T dependence in bulk and RS1 scenario. The topology of WBF jets is also sensible on the underlying model that defines the momentum dependence of couplings. There is a stronger mass dependence of WBF topology of processes involving bulk scenario KK-gravitons. This is due its couplings to be mainly with longitudinal gauge bosons, making the dominant part of coupling to be more momentum dependent than RS1 case. In this case the KK-graviton coupling to weak bosons is mainly to transverse modes, therefore the coupling is mainly not dependent of momenta and there is no mass dependence on the WBF topology.

In table (II.4) we show the impact of our analysis cuts on the WBF jets system (II.22) and on di-higgs system (II.23) in bulk KK-graviton hypothesis.

Sample	basic cuts eqs. (II.19,II.20)	jet merging (AkT5)	$M_{jj} > 400$ GeV	$M_{HH} > 250$ GeV	$p_T^{HH} > 60$ GeV	$\Delta\eta_{HH} < 2$
1450 GeV	0.53	0.49	0.44	0.44	0.41	0.35
1250 GeV	0.52	0.47	0.43	0.43	0.40	0.35
1050 GeV	0.52	0.48	0.44	0.44	0.40	0.35
850 GeV	0.51	0.47	0.42	0.42	0.39	0.34
650 GeV	0.51	0.43	0.39	0.39	0.36	0.33
450 GeV	0.54	0.38	0.34	0.34	0.32	0.30
400 GeV	0.55	0.33	0.29	0.29	0.28	0.26
300 GeV	0.58	0.17	0.14	0.14	0.13	0.12
260 GeV	0.60	0.09	0.07	0.07	0.07	0.07
SM $H(b\bar{b})H(b\bar{b}) \, jj$	0.41	0.38	0.36	0.36	0.30	0.12
$Z(b\bar{b}) \, b\bar{b} \, jj$	0.50	0.36	0.14	0.10	7.91E-02	4.55E-02
$Z(b\bar{b})Z(b\bar{b}) \, jj$	0.62	0.51	0.17	0.12	9.66E-02	6.61E-02
$b\bar{b} \, b\bar{b} \, jj$	0.70	0.20	0.11	6.73E-02	5.49E-02	4.55E-02

Table II.4: Cut flow on resonant signal and background. The displayed numbers are efficiencies.

In the second column of table (II.4) we see the effect of consider realistic AkT05 jets on the analysis flow. At very low mass ($m_{Gr} < 400$ GeV) both higgses are not very boosted. In this limit there is larger probability of jets to superpose generating a intrinsic mis-tag of events, caused by an accidental merging of the b quarks on different higgses or accidental merging with the WBF jets.

The first column this table (containing the angular and p_T cuts that allows b -tag) tell us the selection on b -quarks (II.19) reduce the signal efficiency on a bit more than 40%. The angular selection already reduces $\sim 40\%$ of the phase space for each b -quark, a further suppression is due the transverse momentum cuts. The background loss of efficiency in the first column is due the pre-selection cuts we apply to signal (II.20) super-seed the generation cuts.

The loss of efficiency of signal in low mass regime seen in table (II.4) is effect of the higher jet multiplicity environment. Apart from this effect we expect our analysis strategy to hold parton level scale invariance on most of the mass spectrum. In figure (II.13) we find signal reconstructed efficiency separated by the number of mass drop tags at parton level.

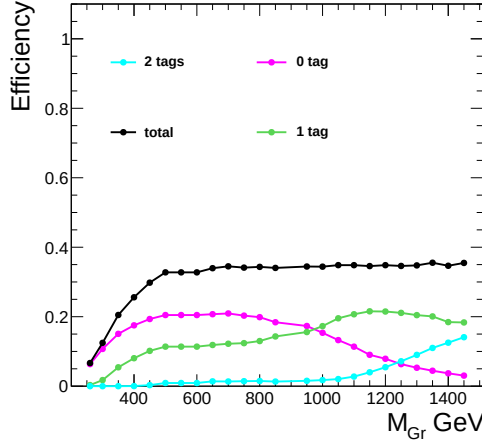


Figure II.13: Signal separation on number of mass drops after analysis flow in a parton simulation level.

In this figure (II.13) we see that for a resonance mass hypothesis larger than 1 TeV there is a systematic increase of events with two tags. The two b -quarks products of each one of the higgses start to merge in one jet. This change of behavior at 1 TeV is expected from calculating the distance radius between two partons giving the boost of their mother particle.

In table (II.5) we show the quantity of signal events expected with the luminosity expected for HL-LHC ($L = 3 \text{ ab}^{-1}$), after all selections of table II.4. Although in the public note [154] we show the values of $S/\sqrt{B}$ values we chose to suppress the explicitly mention to this quantity here. This information can be easily derived from table (II.4). We want to stress although that aim of this pre-feasibility study is to be able to sanity check the analysis flow in preparation for the more realistic results. The most powerful variables for background rejection in this channel are related with its resonant character and rely on reconstructing the resonance and the b -tag characterization of signal.

In all the probed mass spectra there is a non-negligible rate of events where we find one or two fat jets and therefore jet substructure will be an important element even in the case of a resonance with $m_X < 1 \text{ TeV}$.

To advance demonstrate the discovery potential of the channel and push the need of a higher simulation level study we also display the parton level distributions of the mass of the leading reconstruct higgs and the reconstructed di-higgs resonance in figure (II.14). By construction the reconstruction of parton level MC events is fully effective to each higgs candidate.

	$pp \rightarrow b\bar{b}b\bar{b}jj$ σ (pb)	$\sigma \times \text{eff.}$ (pb)	Nevents (3000/fb)
1450 GeV	6.91E-06	8.17E-07	2.45E+00
1250 GeV	1.86E-05	2.14E-06	6.41E+00
1050 GeV	7.87E-05	9.13E-06	2.74E+01
850 GeV	2.63E-04	3.00E-05	9.01E+01
650 GeV	1.28E-03	1.40E-04	4.20E+02
450 GeV	8.87E-03	8.80E-04	2.64E+03
400 GeV	1.41E-02	1.20E-03	3.60E+03
300 GeV	1.55E-02	6.42E-04	1.93E+03
260 GeV	5.72E-03	1.27E-04	3.80E+02
SM $H(b\bar{b})H(b\bar{b}) \, jj$	2.71E-04	3.33E-05	1.00E+02
$Z(b\bar{b}) \, b\bar{b} \, jj$	1.46E+00	6.63E-02	1.99E+05
$Z(b\bar{b})Z(b\bar{b}) \, jj$	5.96E-03	3.94E-04	1.18E+03
$b\bar{b} \, b\bar{b} \, jj$	1.48E+02	6.73E+00	2.02E+07

Table II.5: Cross section and events for HL-LHC, after all selections of Table II.4. The parameters used are $\tilde{k} = 0.1$ and naively fixing $kr_c = 11$.

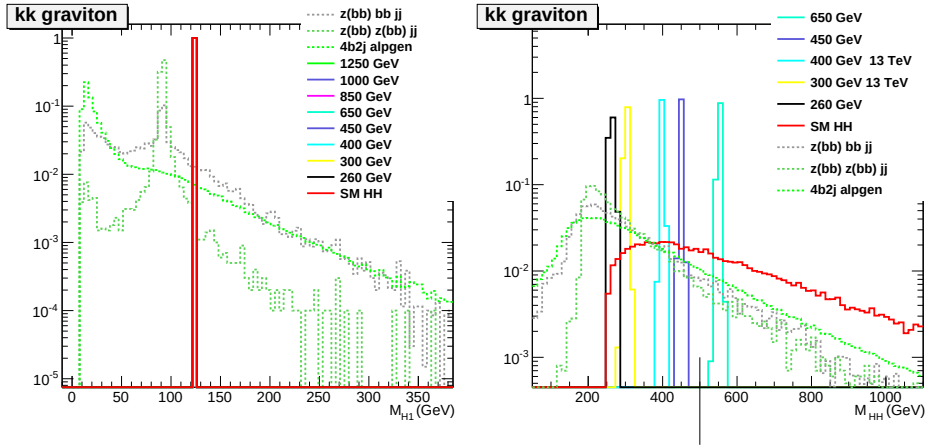


Figure II.14: Parton level distributions of reconstructed objects with no cuts applied. Distributions are shape normalized. **Left:** Leading reconstructed higgs. **Right:** Reconstructed bulk KK-graviton. Those distributions are only for illustration.

II.5.2 Preliminary considerations for detector simulation level study

Generically the ATLAS and CMS detectors are composed by a cylindrical inner tracker covered by calorimeter towers, under the influence of a strong parallel magnetic field. The input are events after showering and hadronization, the **Delphes** program includes the effects of magnetic fields on trajectories and simulates the responses of tracker calorimeter towers to each particle species, additional Gaussian smearing are also applied.

The future studies for detector upgrade performance on each one of the LHC conditions (listed on table II.2) performed by the CMS collaboration are being done based on **Delphes3**, a public fast simulation tool for fast detector simulation. The most updated CMS detector configurations to **Delphes3** can be found on CMS public twiki [122]. The ATLAS collaboration had already setup a full simulation framework for internal studies. The public physics and performance studies made on the ATLAS upgrade conditions can be found on ATLAS public twiki [121]. For the non-ATLAS community studies, the collaboration released a letter of intent [155] with the detector parameters to be used.

The prediction of particles reconstruction and tag capability are important elements of the detector simulation. The reconstruction and isolation of objects are performed with a **FastJet** based algorithm. The flavor tag of jets and leptons are modelled using smearings accordingly with detector range and expected response. The triggers paths are emulated based on the stand-alone particles. To simulate pileup environment **Delphes3** mixes random min bias events into detector response on modeling extra jets. An external modulo should be added to mixing pile-up events with the main interaction the pile-up subtraction with **FastJet** area subtraction method is done.

In our study we interface **Pythia6** and **Delphes3** using **HepMC** format. We also had implemented the **FastJet** substructure analysis acting in **Delphes** objects.

One of the practical difficulties of perform a detector level study is the computational demand on processing of the simulations. Also, the fact we have a semi realistic environment for tagging physics objects, we also need to consider background components to account fake tagging.

In the last section we had only considered the backgrounds containing four b -taggable b -quarks on signal selection. The dominant background for this search is expected to be multi-jet QCD. We had estimated at parton level that indeed, after a basic selection, the QCD production of two $b\bar{b}b\bar{b} \ jj$ is dominant when over the EW backgrounds listed on table (II.3), see table (II.5). Even if we want to keep the same requirement of have four b -tags to tag a signal to this study we need also to take into account the b -tag fake rate and consider also processes with less b -quarks on final state, like $b\bar{b} \ jj \ jj$.

The $t\bar{t}$ is a non negligible component on the inclusive $h(b\bar{b})h(b\bar{b})$ search on LHC8 [149]. The SM production of $t\bar{t}$ pairs is known to have high jet multiplicity already at LHC7 [156]. However its contribution on the $h \rightarrow b\bar{b}$ search on WBF mode can be non-negligible and need to be taken into account for a full result.

III Searches for resonant double higgs production on CMS detector

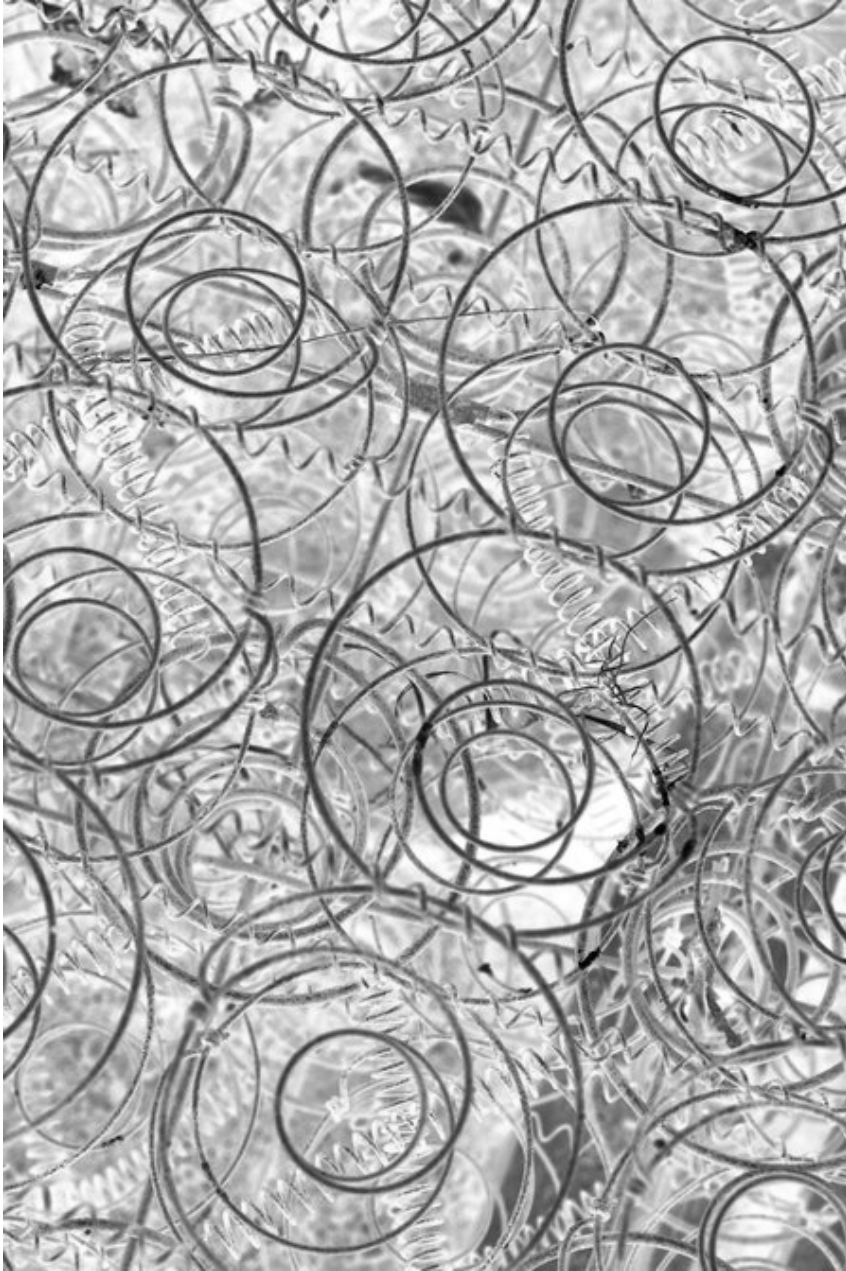


Figure covering this chapter: Allusion to the format of charged particles tracks due the presence of a magnetic field, in a color treated photograph of a mattress.

Since the celebrated (~ 125 GeV) higgs discovery [6, 7] the experimental searches for heavy resonances on LHC are being extended to channels with final states that contain higgs particles. An evidence of a signal on di-higgs final state would state the spin of the resonance to be even, following the motivations already presented along this thesis (because of spin conservation is it not possible to a spin odd particle to decay in two spin-0 particles [157]).

Before a detailed discussion of the analyses we present the CMS sub-detectors and the basics of reconstruction techniques in the section (III.1). In the section (III.2) we describe the $pp \rightarrow X \rightarrow \gamma\gamma b\bar{b}$ analysis. We then comment on the two CMS di-higgs CMS analyses for the process $pp \rightarrow X \rightarrow b\bar{b}b\bar{b}$ in the section (III.3). And finally in the section (III.4) we compare all the public di-higgs searches.

In all the studies the radion particle (as introduced on section I) was used as benchmark for the MC simulations of a mother particle decaying into pair of higgses on LHC environment. For simulation purposes only we use a constant width of 1 GeV for all the different resonance masses, what is compatible with the radion model [64] and with bulk KK-graviton [89] hypothesis. Therefore the final resolution on the mass peak in the analyses will be dominated by detector reconstruction effects. With the intent of checking the analyses dependency of spin hypothesis and showering model we had also simulated samples for a heavy bulk KK-graviton and MSSM hypothesis, we had respectively used **MadGraph5+Pythia6** and **Pythia6** for the Leading Order hard scattering simulation + showering.

The gluon fusion on shell radion production was simulated using **MadGraph5**. The events are hadronized under **Pythia6** and processed on full (**Geant4** based [158, 159]) simulation of CMS detector. The **Geant4** simulation is impressively close to reality, but at the level of precision requested we have to correct for simulation imperfections and bring the simulated objects closer to reality.

III.1 The CMS detector

The CMS experiment (Compact Muon Solenoid) is a cylinder of 28.7 meters of length and 15 meters diameter, with 14.000 tons.

The main characteristics of CMS are the central super-conducting solenoid with four layers external muon chambers (muon stations) intercalated by iron yokes. The solenoid cylinder generates 3.8 Tesla magnetic field. A very detailed description of CMS detector is found on its Technical Design Reports [160], a more recent source for technical information about CMS detector can also be found on [161].

In figure (III.1) we show a transverse slice of CMS marking its sub-detectors, the signals the different stable objects would leave on the sub-detectors are also outlined. Four layers of muon chambers are installed outside the solenoid cylinder. Inside the solenoid cylinder we have three sub detectors:

- The tracker system is composed of a silicon strip and a pixel detector. It is installed on the inner part of cylinder. The high magnetic field of the solenoid assure the tracker to be very sensitive to the charge of very energetic particles. The bend of the particles in the presence of the magnetic field guarantees a good p_T resolution. The CMS tracker is very dense and act as a pre-shower stage.
- The following slice is the electro-magnetic calorimeter (ECAL), where the total energy of electrons and photons are deposited and measured. The ECAL is made by homogeneous scintillating crystals.
- Finally, on the last layer, we have a hadronic calorimeter (HCAL), where the total energy of hadrons are deposited and measured.

The cylindrical form of the detector determines two regions to be covered by the sub-detectors barrel and end-cap. The barrel is the cylindrical region along the z -axis. The endcap is oriented perpendicularly to the z -axis, closing the barrel cylinder. The capabilities of the detection of each kind of stable particle depend on the reach of each dedicated sub-detector.

Figure (III.2) shows a longitudinal view of one quadrant of CMS and the angular range of the CMS sub-detectors. After the muon chambers very forward hadronic calorimeters are installed around the beam direction, covering a pseudo-rapidity region of $3.0 < |\eta| < 5.0$. There is a barrel-endcap transition region between $1.44 < \eta < 1.57$ in the calorimeters coverage.

III.1.1 Luminosity

The knowledge of luminosity recorded by CMS experiment is an important information to allow physics interpretations of the data events. The luminosity measured by CMS is made by the pixel detector [164]. The measure of luminosity based on very forward calorimeter activity [165] is used for estimation of systematic uncertainties and cross checks.

The calculation of luminosity based on pixel detector follows the formula (II.1), corrected for the real geometry of the bunch crossing cross section and emittance, see for example [166].

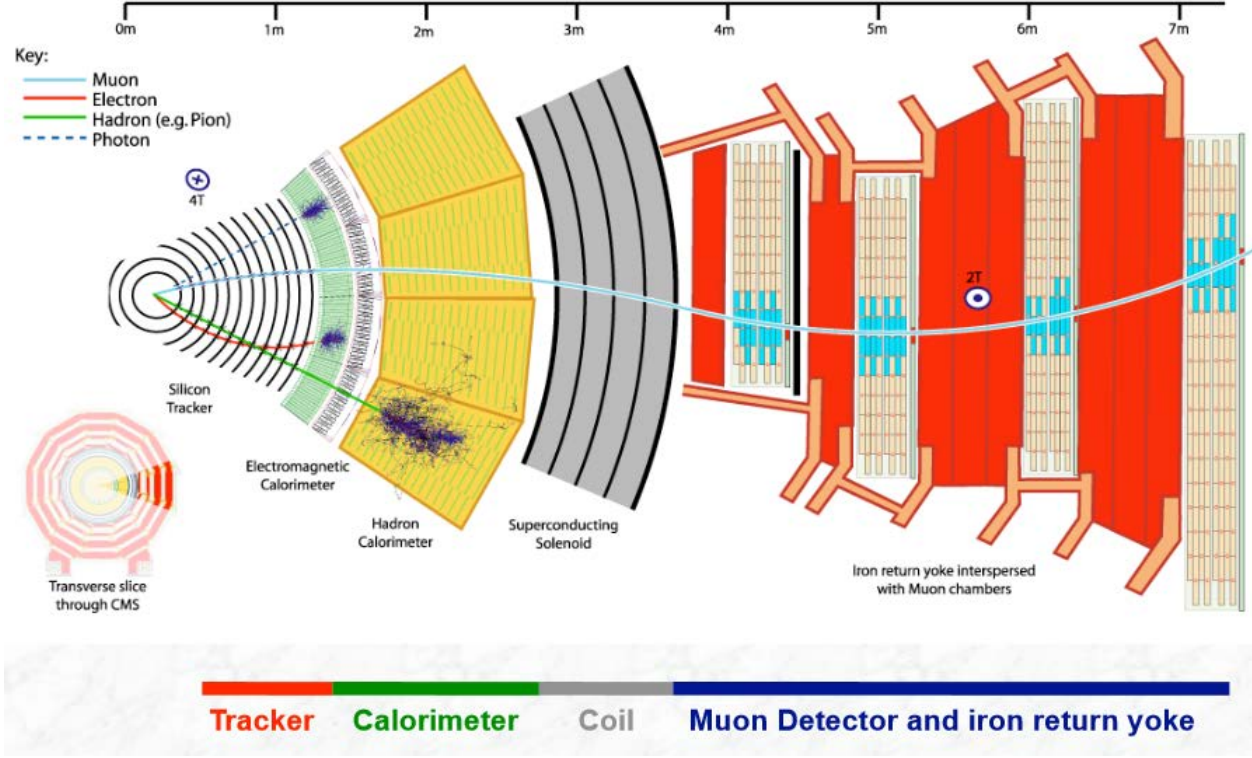


Figure III.1: Schematic view of CMS sub-detectors [162].

The relation between the number of hits and luminosity is $\frac{dN_{hits}}{dt} = \sigma_{mb}L$, where σ_{mb} is the cross section estimated from minimum bias events [167]. The geometry of the overlap between the colliding beams is estimated using the Van der Meer technique, in simple words consist of systematic scan varying the beams centrality.

The most updated results on CMS luminosity can be found on the public Twiki [168]. The central value for integrated luminosity recorded by CMS detector during the 8 TeV run is 19.6/fb, with 2.8% of systematic uncertainty [164, 168].

III.1.2 Trigger

There are two trigger levels on CMS: The Level 1 (L1) is the first processing act of the signals collected during a bunch crossing. The L1 trigger have a latency time determined from the time of processing for the particles identification. The objects identification at L1 trigger are mostly made at hardware level, based on the sub-detectors signatures and a only partial event reconstruction. To reduce the time of latency the event processing must be fast, it looks for single objects, for example for an event with a single muon or multi-jets [169].

The second trigger level is also known as High Level Trigger (HLT). This level of events trigger relies on more complete software event reconstruction, its time of latency is larger. The analysis dedicated triggers are designed at this level. The trigger perfor-

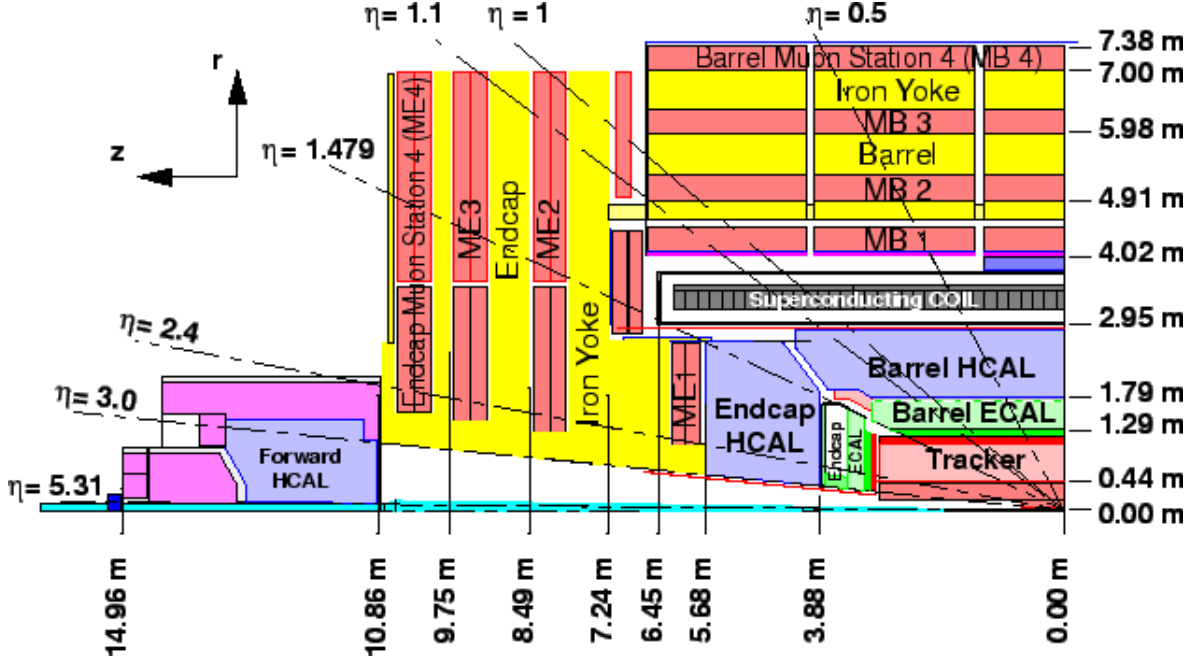


Figure III.2: Longitudinal view of one quadrant of CMS sub-detectors [163].

mances reported by the different searches are calculated with respect to HLT.

A primary vertex is assigned to each triggered event. The primary vertex is defined as the most probable vertex from where the hard event was originated. The definition of primary vertex for triggered events that contain jets and/or charged leptons is made using the tracks information. In this case the primary vertex is the one that holds the biggest $\sum p_{T_{track}}^2$ where the sum is for all the tracks attached to the vertex.

III.1.3 Particle Flow

The CMS collaboration developed an algorithm to reconstruct the physical objects combining the information from all sub-detectors. The information of the calorimeters, tracker and muon chambers are combined with an algorithm known as Particle Flow (PF) [170].

The reconstructed particles are characterized in five types: muons, electrons, photons, charged and neutral hadrons. We refer to PF objects as detector-traceable objects reconstructed with the PF algorithm.

III.1.4 Jets

All the objects resulted from clustering of PF objects are passed through a jet clustering algorithm implemented in *FastJet*. The jet energy measured by the detector needs to be calibrated for correct the experimental and theoretical expected mis-measurements.

A jet calibration is done under one specific jet definition. The AkT05 jet is very seriously calibrated, while the C/A08 is only approximatively calibrated. The di-jet cross first measurement in CMS is found in the references [171, 172].

The Jet Energy Corrections (JEC) are designed to account with detector energy mis-measurement of PF objects. Those are data driven techniques to correct the energy measured by the calorimeters using elements like the transverse energy measurements from particle flow reconstruction, expected detector response, energy balance of the event and the removal of soft radiation. The JEC are divided in corrections to the overall Jet Energy Scale (JES) and Jet Energy Resolution (JER) [173, 174, 175]. The potential lost of energy from radiation of neutrinos can only be corrected by looking on the jets substructure, this procedure is called *regression*.

After the mentioned procedures jets are considered stand-alone objects. A jet however can be tagged as pileup or for its flavor.

III.1.4.1 Pileup rejection

Pileup jets are characterized by a large fraction of charged PF candidates that are not attached to primary vertex. The fraction of charged PF candidates attached to the primary vertex is quantified by the the sum of the p_T of all the charged PF objects attached to a given vertex (β^*), balanced by the total jet p_T :

$$\beta^* \equiv \frac{\sum_{jetPUvertex} p_T^{tk}}{\sum_{jet} p_T^{tk}}. \quad (\text{III.1})$$

Fake jets are characterized also by a rather uniform distribution of the energy within the jet cone. To help the fake jet identification and rejection we define the RMS variable as the median of the distance of each constituent and the direction of the jet (ΔR) balanced by p_T :

$$RMS \equiv \frac{\sum_{constituents} p_T^2 \Delta R^2}{\sum_{constituents} p_T^2}. \quad (\text{III.2})$$

III.1.4.2 b -tag

There are more than one algorithm designed inside CMS collaboration to recognizes b -quark initiated jets [176, 177], as mentioned before the b -tag algorithm is resumed on one real valued variable. The jet b -tag algorithms used by CMS have as ingredient track information, only jets in the angular range $|\eta| < 2.5$ can be b -tagged. All the tag methods are designed to exploit the properties of a B hadron. The tag efficiencies are dependent of p_T and η range. The simplest and faster algorithms are designed to be suitable to be used on HLT, while more complex and robust algorithms can be designed to be used on the physics analyses. Three Working points (WPs) are defined from the value of the b -tag variable: loose, medium, and tight.

Two b -tag discriminators are constructed based on tracker information only: The *track counting* variable is constructed using only the impact parameter of tracks inside the jet*.

*The impact parameter of a track is calculated based in the distance between the point where the track extrapolation crosses the jet direction and the primary vertex.

The *jet probability* variable however combines the track counting with the matches of these good tracks with primary and displaced vertices [178].

The *Combined Secondary Vertex* algorithm (CSV), for example, combines all the known experimental variables that can distinguish b -jets using a likelihood ratio technique. As a rough estimative the probability of tag a b -quark originated jet, with medium CSV b -tag definition, is $\sim 60\%$. There is however $\sim 10\%$ possibility of mis-tag c -quark originated jet, and of $\sim 0.01\%$ of mis-tag light-quark originated jet with medium CSV tag [179].

III.1.5 Electromagnetic objects

The electromagnetic objects are defined as photons and electrons. The interaction of those particles with the detector material generates showers of secondary electromagnetic particles. The hard part of the electromagnetic showers are formed from hard electrons bremsstrahlung ($e + A \rightarrow e\gamma$) and electrons pair production from protons ($\gamma + A \rightarrow e^+e^-$) where A is the material atomic electronic cloud. When the energy of showering particles is smaller than a critical threshold a large part of the energy is dissipated through the ionization process. This later lead to the scintillation phenomena in ECAL that allow to detect particles in the crystal.

Electromagnetic showers usually leave signals in the ECAL, if the shower is very energetic it can also deposit energy in HCAL. Electrons (and positrons) leave tracks on the inner tracker. A reconstruction method other than the PF definition is also used by specific analyses. The characterization of electromagnetic objects is based on two complementary algorithms: the ECAL driven algorithm starts from super-clusters signals and search for compatible hits on the inner tracker. The tracker driven reconstruction technique is the contrary: it extrapolates the trajectories up to ECAL towers.

Photons and electrons can lead to very similar detector signals. The photon signals and therefore the photon reconstruction method depends on which position of the detector the photon had first interacted with material:

- If the photon converts inside the ECAL leaves no signal in the tracker system. All the photon energy is deposited on this calorimeter but since the ECAL is an one layer detector no direction can be reconstructed. In this case the photon is identified as a footprint in ECAL not attached to any track.
- The photon can also start to shower the in tracker system, induced by its high density material [180]. If the photon start to shower inside the tracker the only part of its energy is measured by the ECAL system, its trajectory however can be reconstructed.
- A very early electron-positron conversion can lead to very separated electron positron pair reaching the ECAL, so called conversion. Photon conversions are characterized by secondary tracks of an electron-positron pair with an invariant mass consistent with zero, originating from the a common vertex. In this case the photon is recognized if opposite charge tracks are fitted to the same common vertex, with no invariant mass.

Many SM processes have electromagnetic objects as final states. Those processes are background to BSM signals and also serves as calibration processes for identification [181].

The photons identification is calibrated using the SM $Z \rightarrow \mu\bar{\mu}\gamma$ process. This processes is very rare therefore the $Z \rightarrow e\bar{e}$ process is also used for photon ID calibration (see for example [182]).

More than 50% of the photons converts, the charged particles (product of the conversion) are bent on ϕ direction. To hold a good resolution on converted photons the individual ECAL crystals in barrel are organized on super-clusters that are long on ϕ direction and short on η . In endcap however the super-clusters are organized as 5×5 cubes of crystals.

Shower shape and isolation variables are used to differentiates electrons and photons.

III.1.5.1 Calorimeters shower shape

In most of the cases the energy of electromagnetic objects is fully deposited in the ECAL. The distribution of the ECAL energy deposited hints about the nature of the object that generated the shower [183].

An electromagnetic shower can also deposit some energy on HCAL. An electron or converted photon can be mis-identified with a jet. A non-converted photon photon can be mis-identified as a neutral meson. The ratio of hadronic to the electromagnetic energy in a cone of $\Delta R = 0.15$ around the super-cluster position helps to discriminate between jets and pure electromagnetic objects.

A very energetic neutral pion that decay into a photon pair can be mis-reconstructed as a single photon. The spread of energy in the crystals is an information that helps to distinguish a isolated hard photon from a pair of collimated photons from π^0 fragmentation.

The ratio of the energy in the 3×3 crystals centered at the crystal with maximum energy to the super-cluster energy ($R9$) determine purity of photon candidates. Photons that had converted in the ECAL are characterized by a high $R9$ variable, while low $R9$ photons probably had initiated shower on the inner tracker.

III.1.5.2 Isolation

The isolation is an important element to distinguish prompt photons and electrons from photons and electrons from jet hadronization. The direction from where to calculate the PF isolation of the non-converted photons is defined as the line connecting the photon super-cluster position and the primary vertex.

Instead of a detector-based isolation (based on ECAL, HCAL and tracker) in CMS the isolation of electromagnetic objects is calculated with respect to PF particles. An isolation cone of size $\Delta R < 0.3$ is used. Like it is done for the jets, the contribution of the pile-up and underlying event inside the isolation cone is estimated with an effective area method.

III.2 $pp \rightarrow X \rightarrow h(\gamma\gamma)h(b\bar{b})$ at CMS

This final state is characterized by a pair of photons and a pair of b -jets. Its prime challenge is the photon characterization. We construct the analysis on top of the $h \rightarrow \gamma\gamma$ [184] analysis expertise. The full dataset collected by the di-photon trigger in 2012 at 8 TeV, summing up to a integrated luminosity of $19.7/\text{fb} \pm 2.6\%$ was analyzed*.

The $\gamma\gamma b\bar{b}$ final state combines the large branching fraction of the $h \rightarrow b\bar{b}$ decay with the comparatively low background and good mass resolution of the $h \rightarrow \gamma\gamma$ channel. A preliminary study of this channel in the context of the CMS detector is found in the physics TDR [185].

This search is performed in the mass range 260-1100 GeV. The upper bound for the search range is defined by the transition between the two distinct jets regime and the one jet regime in which the boosted decay products of the hadronic Higgs are reconstructed into one fat-jet (estimated to be ~ 1 TeV).

In subsection (III.2.1) we list the MC backgrounds components considered along the study. In the subsections (III.2.2), (III.2.3) and (III.2.4) respectively we detail this analysis specificities on the primary vertex assignment and on the reconstruction of photons and jets objects. Finally in subsection (III.2.5) we describe the signal extraction strategies. The results are summarized in subsection (III.2.7).

III.2.1 Backgrounds

The reconstructed photons can be either prompt photons or light neutral mesons faking photons (also known as fakes). We split the simulation of this background in 3 contributions: the fake-fake contribution (simulated with multi-jet QCD samples on *Pythia6*), the fake-prompt contribution (simulated with γ + jets QCD samples on *Pythia6*), and the prompt-prompt contribution, simulated with *Sherpa*. The dominant background at the pre-selection level is the QCD production of two prompt photons. When known the K-factors are applied, following the recipe from the 2011 Hgg analysis [186].

One specificity of the *Sherpa* generator is to provide event weights to enrich the multi-jet sample with high jet multiplicity. The *Sherpa* simulation indeed proved itself to better describe the topology of two prompt photons in addition to high multiplicity of hard jets than *Pythia6*. A specific k-factor was also derived for this sample.

The SM single higgs production in all the production modes is background to this channel. Even if the cross section for the SM single higgs production is low compared with the QCD component these signals can fake the resonant signal that is looked for. The relevant samples are given in table III.2. The gluon fusion and WBF production modes was simulated with *Powheg* + *Pythia6*. While associated production (Wh , Zh and $t\bar{t}h$) was entirely generated with *Pythia6*.

*The usage of the 2011 data is not considered since they do not add anything to the sensitivity of this low background analysis while requiring an equivalent effort from the point of view of data analysis.

Sample	# events	Cross section (pb)	K-factor
$\gamma\gamma$ + jets	14404429	120.354	3.565
$\gamma\gamma + t\bar{t}$	122040	0.14584	-
$\gamma\gamma + t$	324676	0.00337	-
$\gamma\gamma + l\nu$ (FSR)	1000310	1.84	-
$\gamma\gamma + l\nu$ (ISR)	1003920	0.319	-
$Z(ll)\gamma$	6588161	132.6	-
γ + jets ($p_t = 20$ -40 GeV)	5901106	150.34	1.3
γ + jets ($p_t > 40$ GeV)	17722752	478.58	1.3
$\gamma + t\bar{t}$ +jets	1719954	14.	-
QCD multi-jets ($p_t = 30$ -40 GeV)	6047441	12208.	1.3
QCD multi-jets ($p_t > 40$ GeV)	9764546	51439.	1.3
$ll jj$ (DY)	30290538	2950.0	1.15

Table III.1: List of the non-resonant SM background samples. The horizontal lines separates the components that contains respectively two prompt photons, one prompt photon plus a fake and two fake photons.

Sample	# events	Cross section (pb)
Gluon fusion	96161	4.39×10^{-2}
WBF	99671	3.60×10^{-3}
$Wh + Zh$	100151	$1.61 \times 10^{-3} + 9.47 \times 10^{-4}$
$t\bar{t}h$	93183	2.95×10^{-4}

Table III.2: List of the Standard Model higgs $\rightarrow \gamma\gamma$ samples, with mass $m_H = 125$ GeV.

III.2.2 Primary vertex assignment

Usually the direction of a photon cannot be directly reconstructed within the CMS detector. To reconstruct the photon four-momentum, the di-photon invariant mass and a proper identification of the interaction vertex in a general procedure it is necessary to find an alternative solution other than the tracks information. The solution used for the $h \rightarrow \gamma\gamma$ analysis [187] for the identification of the primary vertex was to construct a BDT (Boosted Decision Tree) variable with the correlations between the kinematic properties of the relics of the proton-proton collision with the kinematic properties of two photons from the decay of a higgs boson.

In the $h h \rightarrow \gamma\gamma b\bar{b}$ analysis however the presence of charged particles (jets) associated to the hard interaction process is an advantage that can be used for the identification of the primary vertex. In the case of jets attached to the primary vertex the most natural solution is to select the vertex with the largest value of $\sum p_T^2$ of the event, where the sum runs over the tracks associated to the vertex.

For this study the signal events are identified by the presence of at least two jets with $p_T > 25$ GeV and one reconstructed photon pair selected according to the details that will be mentioned below. The generator level photon pair within a distance $\Delta R = \sqrt{\Delta\eta^2 + \Delta\phi^2} < 0.05$ are matched to the primary vertex. The vertex identification efficiency is defined as the probability of the $h h \rightarrow \gamma\gamma b\bar{b}$ generator level vertex to be

within a given distance from the reconstructed vertex $|z(\text{gen}) - z(ID)| < \Delta z$.

It is known from the $h \rightarrow \gamma\gamma$ analysis that for $\Delta z < 1$ cm the resolution on the di-photon invariant mass is not limited by the angle resolution [187]. For the di-photon mass reconstruction it is therefore sufficient to localize the vertex within about 1 cm.

We found the vertex identification efficiency to be above 98% from a distance $\Delta z \leq 1$ cm for a 300 GeV mass hypothesis in both vertex identification methods.

III.2.3 Photons

The $h \rightarrow \gamma\gamma$ analysis is designed to correctly characterize two photons coming from a common primary vertex, this group does not use the PF definition of a photon, but construct its own framework to deal with the detector information and recognize a di-photon [186], optimized to di-photon characterization. In this analysis we use exactly the same photon reconstruction and identification strategy*.

The optimization procedure of the photon identification cuts on the shower shape and isolation variables defines Working points (WPs) to be taken as a reference. The cuts performed on those variables are optimized differently to the four photon categories defined to SM $h \rightarrow \gamma\gamma$ analysis[187], namely barrel/endcap with high and low $R9$. The performance and validation of CiC cuts, as well as data/MC scale factors, were studied on $Z \rightarrow e\bar{e}$ and $Z \rightarrow \mu\bar{\mu}\gamma$ events. Figure (III.3) shows the two photons p_T after the basis photon selection:

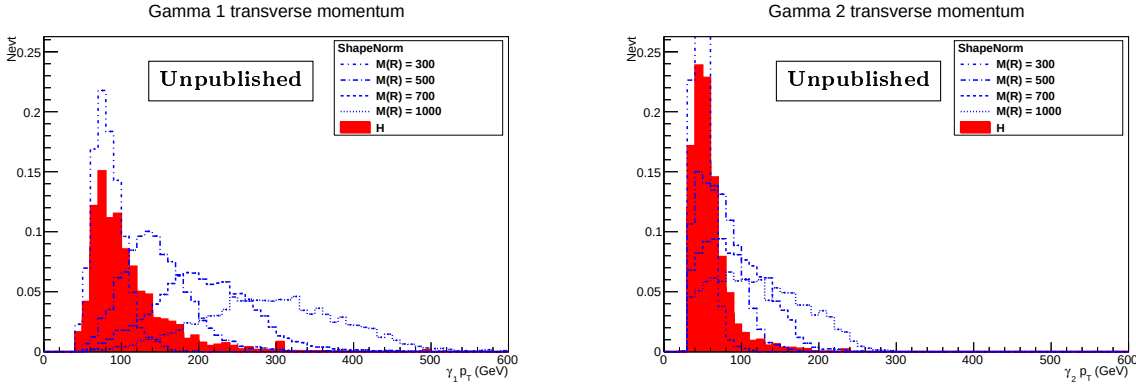


Figure III.3: Transverse momentum of the two selected photons for the radion signal at different masses compared with the single SM higgs production.

The experimental uncertainties on the photons identification and energy corrections on data and MC are computed using the goodness of Z reconstruction on $Z \rightarrow ee$ and $Z \rightarrow \mu\mu\gamma$ events. In case of a low mass resonance ($\lesssim 300$ GeV) decaying to the considered

*Two photon identification (ID) strategies exist in the SM $h \rightarrow \gamma\gamma$ analysis: one of them is cut-based and another is based in a boosted decision tree (BDT) [188]. The BDT-based photon ID uses as input the kinematics of the SM $h \rightarrow \gamma\gamma$ to construct the BDT variable. The cut-based photon ID as documented on [187] will be used.

channel the kinematics of the $h \rightarrow \gamma\gamma$ is very similar to the one from the SM single higgs and so within the p_T range of the SM Z -peak calibration point.

In figure (III.3) we also note that in the case of high mass resonance ($> \sim 500$ GeV) the two higgses are boosted and therefore each one of the photons resultant from the higgs decay is also boosted. The drawback of the photons of the investigated signal to have sizeable boost is that the p_T of the photons can be beyond the range for which the uncertainties on the energy corrections can be computed. We overcome this fact increasing the systematic uncertainty on the photon energy scale factors (as will be clarified on the systematics part).

In the SM $h \rightarrow \gamma\gamma$ analysis three independent frameworks are used for sanity checks. One unique framework however is used for nominal results. The $hh \rightarrow \gamma\gamma b\bar{b}$ analysis nominal results are constructed with the same framework used in [184], a second photon identification framework is synchronized within few % level and also used for some specific studies, like bias studies for background estimation.

III.2.4 Jets selection

The issues concerning jets specific to this analysis is the choice of the categorizations of events based on b -tag information, the jet combinatorics and the kinematic fit we apply in order to tune energy corrections on the chosen two jets to match better higgs mass.

At pre-selection level the jet candidates are requested to have $p_T > 20$ GeV and $|\eta| < 2.5$ and hold a separation from each one the two selected photons by at least $\Delta R = 0.3$.

We use the Combined secondary Vertex (CSV) tagger for b -tag characterization, as done in the WBF $h \rightarrow b\bar{b}$ searches* [189]. The most natural approach to our signal would be request at least two jets in the event are b -tagged. Nevertheless according to the choice of the b -tag working point one of the jets could fail the b -tag selection.

In figure III.4 we show the number of b -tagged jets per event for the signal and some of the relevant backgrounds as a function of the working point after the full photon selection.

The signal extraction will be done by fitting the signal shape over a mass spectra. The categorization of signal on regions with different signal resolutions is a very powerful statistical strategy for shape analysis. In the SM $h \rightarrow \gamma\gamma$ analysis the barrel-endcap separation and $R9$ variables defines four purity categories to find di-photons with different di-photon mass resolution for signal.

The purity of the $h \rightarrow \gamma\gamma b\bar{b}$ signal however is driven by the presence of b -tagged jets on the event. We define a low and high purity based in the number of b -tags. The Medium working point is used. This choice for characterization is discussed in appendix (E). We consider the two categories summarized in table (III.3).

The low purity category suffers from a larger background however it helps to save efficiency for high mass resonance hypothesis, where the SM background is small. The low data statistics on the analyzed luminosity prevents us to make any further characterization. In figure (III.5) we show the efficiencies of the CSV tagger for a 300 GeV radion

*The jet probability tagger was also tested to show our results are not too much affected by the tagger choice.

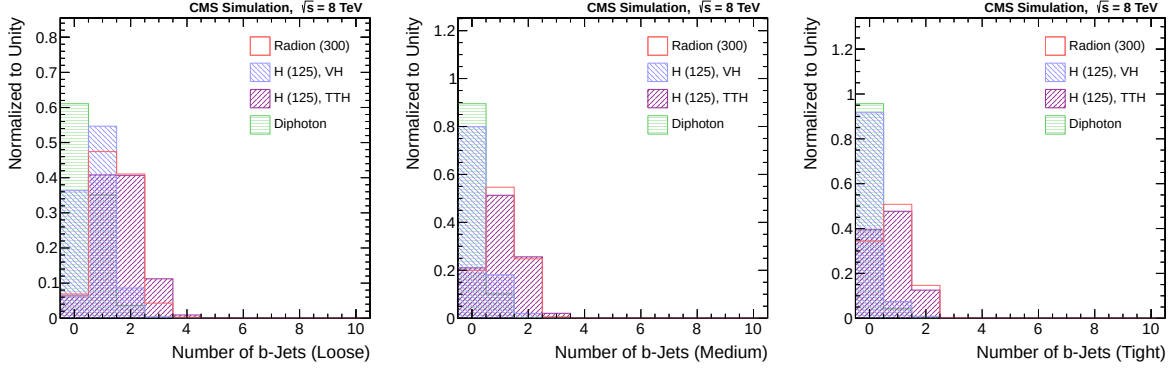


Figure III.4: Number of b-tagged jets as a function of the working point for the signal and some backgrounds.

High purity	Low purity
≥ 2 medium b -tagged jets	exactly 1 medium b -tagged jet

Table III.3: Events categorization.

signal as a function of the jet p_T , for different η -bins and flavors, after jets full selections.

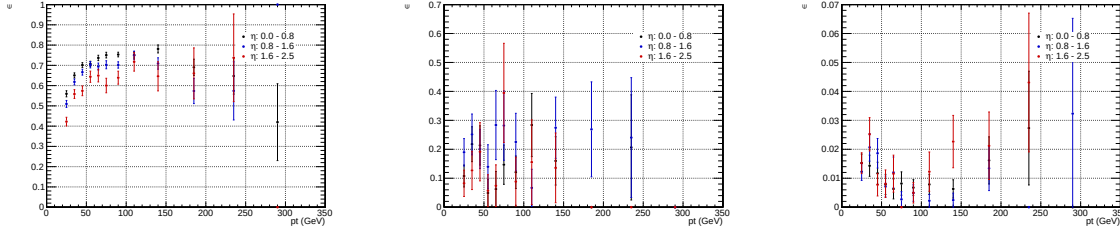


Figure III.5: b-tagging efficiencies as a function of the jet p_T for different η -bins and flavors (from the left to the right, for b, c and light flavors). The plots refer to jets from the 300 GeV radion sample that pass the jet selections described in table III.5. The vertical scale represents efficiency, it varies between 0 and 1 for b flavor and between 0 and 0.7 for the others. The horizontal scale is transverse momenta, it varies between 0 and 250 GeV. The red, blue and black points are respectively for a η range between 0 to 0.8, 0.8 to 1.6 and 1.6 to 2.5.

The detector simulation is not fully perfect therefore b -tag efficiencies for the jets are different between data and MC. To account for this fact the CMS BTV POG (b -Tag and Vertexing Physics Object Group) calculates Scaling Factors (SF) to corrects the MC efficiencies [190]. Those scale factors are typically between 0.8 and 1.6, depending on the flavor, p_T and η of the jet. These weights are applied per category, and therefore do not produce any migration-effect between the categories defined on table (III.3).

III.2.4.1 Jets combinatorics and MC truth

Even after the pileup rejection procedure some jet multiplicity remains on signal and background (see figure III.6). To choose the two jets which are expected to form the $h \rightarrow b\bar{b}$ candidate it is done a screening in the jets of each event based b -tag information:

- If in the jet collection there are exactly two medium b -tagged jets this pair will be the chosen. If there are more than two medium b -tagged jets the pair will be chosen by one of the pairing criteria listed below.
- If however there is only one medium b -tagged jet, this jet will be paired with a non b -tagged jet.
- If there are no b -tagged jet on the jet collection we use only the pairing criteria.

In a given jet collection define by the above screening we can define the $h \rightarrow b\bar{b}$ candidate using the different possible choices for jets pairing criteria:

1. The individually highest p_T jets.
2. The combination of jets with highest sum of p_T .
3. The combination of jets with highest ratio between sum of p_T and invariant mass ($p_T(jj)/m_{jj}$).
4. The combination of jets which have the invariant mass closest to $m_{\gamma\gamma}$.

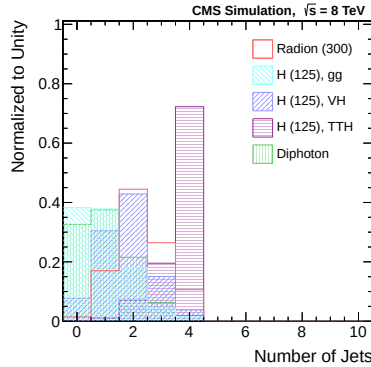


Figure III.6: Jet multiplicity after loose jet identification and pileup rejection.

The possible choices are compared using as the target the MC truth match in 300 GeV signal events*. We show the result of this study in table (III.4). To demonstrate the power of b -tag information on the jets pairing in this table we also show the same MC truth match when not using the b -tag screening. The b -tag information improves significantly the probability that the di-jet choice is correct, we also realize that it makes all the proposed methods to be equivalent. To avoid any danger of background shaping in the di-jet spectra due to the jet selection criteria we decide to use the highest di-jet p_T sum as pairing criteria.

*The MC truth is defined by the frequency that the chosen jet pairs are indeed the simulated ones.

Strategy	not using b -tag screening	using b -tag screening
1	0.67	0.87
2	0.64	0.87
3	0.66	0.87

Table III.4: Percentage of selected signal events ($M_X = 300$ GeV) in which the reconstructed pair of jets is closest to the MC truth for different di-jet combinatorics choices. We do not show the results to the combination of jets which have the invariant mass closest to $m_{\gamma\gamma}$ since it potentially shapes the background shape.

III.2.4.2 Kinematic fit

Once the di-jet pair is selected it is possible to improve the resolution in the 4-bodies spectrum exploiting the knowledge that the signal di-jet pair mass shall peak around the higgs mass (125 GeV). In order to optimally scale the di-jet quadri-momentum to the higgs boson mass it is used a kinematic fit to the two jets.

The kinematic fit consists on constraining the mass of the di-jet system to the closest value to the mass of di-photon system by modifying for each jet the values of E_T , η and ϕ in accordance to their expected resolutions. The method used to minimize the di-jet mass difference from 125 GeV based on the kinematic variables using a Lagrange Multiplier method. The official resolutions estimated on MC by [191] and [192] are used. Figure (III.7) shows the impact of the kinematic fit in the four body spectra for signal and the sum of all components of single SM higgs background.

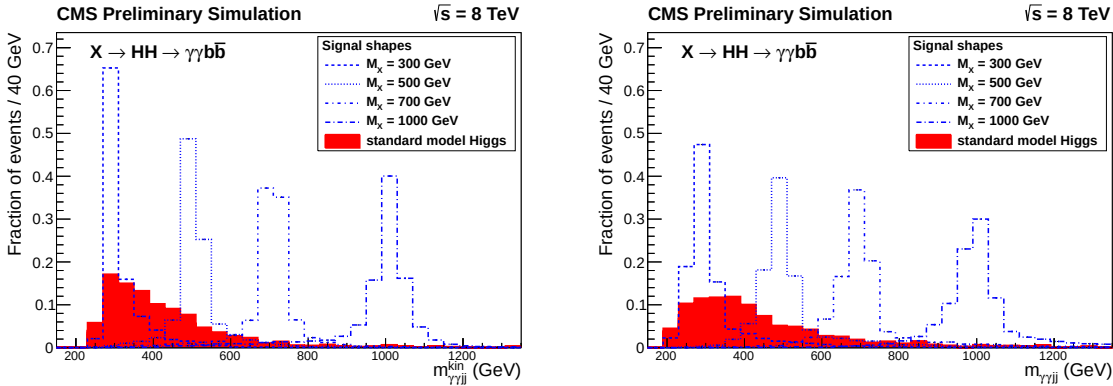


Figure III.7: The $M_{\gamma\gamma jj}$ spectrum of radion signal and single SM $h \rightarrow \gamma\gamma$ produced in association with a hadronic V boson. **Left:** With kinematic fit. **Right:** Without kinematic fit.

For the single SM higgs production the kinematical constrain does not exploit any physical process. Therefore the kinematic fit have the effect of only shuffling randomly the events in the final four-bodies invariant mass spectrum and therefore four bodies mass is not peaking even when the kinematic fit is applied. The same conclusion holds to the QCD background. For the resonant signal however the improvement in resolution is quite large, and more significant for the lower mass resonance hypothesis. The relative resolution improves with the jet p_T and therefore the di-jet higgs the smaller is the resolution gain by the kinematic fit is smaller with the increase the p_T .

III.2.5 Strategy for the extraction of a resonant signal

The obvious property of our signal is the presence of three mass spectra with a resonant fingerprint. The di-jet and di-photon mass possesses an excess on higgs mass while the 4-body mass shows resonance signal. Figure (III.8) shows the mentioned mass spectra. The two bottom plots in the figure are different scale versions of the same variable.

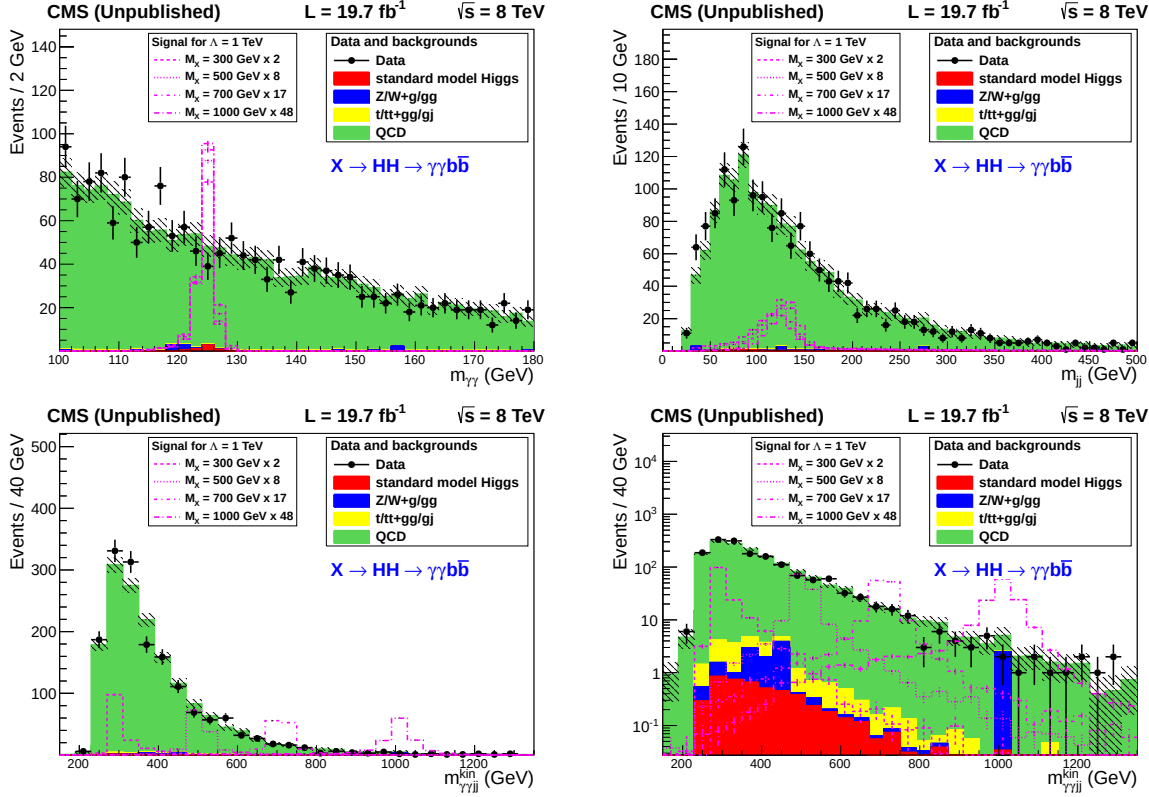


Figure III.8: Mass spectra after some basic selections on photons and jets, and requesting at least one loose b-tagged jet.

We would like to extract the signal shape from a fully data driven technique. Based in the masses spectra of the above figure we define three different kinematic regions:

- **Low mass region, m_X from 250 GeV to 400 GeV:** In this region we can explore the excellent resolution of signal in the di-photon invariant mass distribution. We perform a fit the $m_{\gamma\gamma}$ spectrum, after having imposed cuts on the di-jets and the 4-bodies mass. We assume the data to be background dominated and therefore the $m_{\gamma\gamma}$ fit of all selected data is considered for background estimation.

As a cross-check for the $m_{\gamma\gamma}$ fit method in the low mass region we also fit the 4-bodies mass spectrum. The 4-body spectrum exhibits a kinematic peak around 200-300 GeV. A resonant signal in this region would appear as a small peak on top of this kinematic peak. The background shape and normalization is estimated also from data but using sidebands in the $m_{\gamma\gamma}$ spectrum, as will be described.

- **Intermediate mass region, m_X from 400 GeV to 1 TeV:** In this mass region is smoothly falling. The signal is extracted with a fit in the four-bodies mass spectrum, after having imposed cuts on the di-jets and the di-photons mass. We also assume the data to be background dominated and use the fit of all data on signal region as estimation of background.

The $M_{\gamma\gamma}$ fit (after having imposed cuts on the di-jets and the 4-bodies mass) is used as a cross check method for this search region. In this method however the four body mass selection considerably reduces the background for high mass hypotheses, limiting the reliability of the results.

- **High mass region, $m_X > 1000$ GeV:** in the case of resonances with mass above 1 TeV the two b-quarks produced by the higgs boson start to merge in a single AkT05 jet. The study of this region therefore requires considering the possibility of a leading higgs jet with a substructure, rather than two separated b-jets.

This regime will not be explored in this thesis. A feasibility study in MC would help to understand up to which mass the isolation of the two photons would allow a efficient search.

For data driven specific studies of background shape and window selection optimizations a data control sample was constructed. The control background model is built selecting events where one of the photons is requested to be tightly selected and one to anti-selected. This sample is dominated by one $\gamma +$ three jets.

The main kinematic variables in this configuration have a slightly different shape than in the signal region and therefore one re-weigh the sidebands events to match the approximately the data spectra exploring the photons kinematics. A map of weights in 2 dimensions as function of the $p_{T,\gamma 1}$ and $p_{T,\gamma 2}$ is built. The photons re-weighting strategy is found to describe much better different kinematic variables than the jets based one. Figure (III.9) shows the main kinematic variables we use for signal extraction in data and control sample.

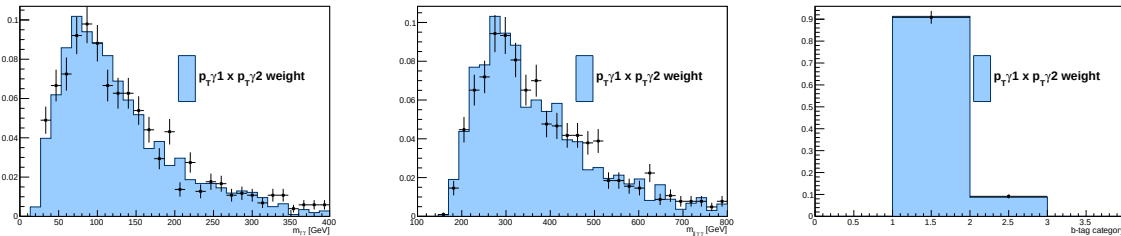


Figure III.9: Some of the main kinematic variables in data (dots) and in data control sample (filled histograms) after kinematic weighting normalized to shape after the baseline selection. **From left to right** m_{jj} spectra, $m_{\gamma\gamma jj}$ spectra and b-tag category population. Only events with $m_{\gamma\gamma} > 135$ GeV or $m_{\gamma\gamma} < 115$ GeV are used for the plots.

We summarize the baseline selection in table (III.5).

Photons	Jets	Events classification
tight photon identification	loose jet identification	
$p_{T\gamma_1}/m_{\gamma\gamma} > 1/3$	pileup rejection	
$p_{T\gamma_2}/m_{\gamma\gamma} > 1/4$	$p_{Tj} > 25$ GeV	
$ \eta_\gamma < 2.5$	$ \eta_j < 2.5$	≥ 2 b-tagged jets
$100 < m_{\gamma\gamma} < 180$ GeV	at least 1 b-tagged jet	exactly 1 b-tagged jet

Table III.5: Summary of the selections applied to photons and jets and of the events classification.

For both $M_{\gamma\gamma}$ and the $M_{\gamma jj}$ fit strategies we restrict the di-jets invariant mass spectrum. The optimization of the M_{jj} mass window was performed against the control sample data. A cut and count approach limit setting was used as figure of merit.

This procedure is done separately for the $m_{\gamma\gamma}$ and $m_{\gamma jj}$ -based extraction procedures, also for each category and mass hypothesis.

In the following we describe the signal extraction procedure in the considered resonance mass regions.

III.2.5.1 The low mass region, m_X from 250 GeV to 400 GeV

This region is still characterized by a large non-resonant background allowing an efficient data driven estimate. The limit extraction is done through a fit to the $M_{\gamma\gamma}$ spectrum. The $h \rightarrow b\bar{b}$ and X are tagged through the selection in the m_{jj} and $m_{\gamma jj}$ mass spectra.

To optimize the $m_{\gamma jj}$ window size we use a procedure identical to the one described to define the mass window selection for the m_{jj} window. The statistics in the high purity category is very low and therefore the cut optimization procedure is done only for the low purity category and applied to all the events. The mass selection based on these studies is given in table III.6.

$M_{\gamma jj}$ selection							M_{jj} selection
260	270	300	350	400	450	500	all hyp.
[225, 280]	[225, 295]	[255, 330]	[310, 395]	[370, 440]	[410, 495]	[445, 535]	[85, 155]

Table III.6: Mass cuts for the $M_{\gamma\gamma}$ fit method by mass hypothesis. All the numbers are in GeV units.

Despite its low cross compared with the QCD backgrounds the production of a single SM higgs with $h \rightarrow \gamma\gamma$ is peaked in the di-photon invariant mass spectrum at the same place as our signal. Given the signal extraction strategy used this background component is a potentially critical background. The Vh and WBF components are partially removed on requiring the presence of two b-tagged jets. The remaining contamination is mainly due to the overlap between the peaks of the vector going hadronically and the $h \rightarrow b\bar{b}$.

The presence of b -quarks and the mass resolution of the W bosons decaying hadronically introduce a contamination from $t\bar{t}h$ events. A small number of jets is a promising discriminant against this background.

The ggh channel is characterized by the largest cross section, despite the absence of hard jets on the hard event this component can be important in background estimation. The di-jet pair with m_{jj} in the region of 125 GeV is faked by the additional ISR jets. Radiation jets are usually produced with small p_T , the high invariant mass however is generated by the large angular distance. This implies that the fake hadronic higgs propagates usually quite forward and exhibits a large rapidity gap with the higgs in the di-photon final state. A selection on $\Delta\eta_{hh}$ could help against this background, however a cut on $\Delta\eta_{hh}$ also reduce the resonant signal significance against non-resonant backgrounds and we had decided not to use them for this round of the analysis.

The only discriminant that we use to reduce the single higgs production contribution are the di-jet window and the multiplicity of b -tagged jets.

Signal model

For the limits extraction we describe the signal with the sum of a crystal ball function and a Gaussian. The crystal ball function is constructed by a convolution between a Gaussian and a exponential tail. It is introduced to account possible non-symmetric effects on signal shape induced by detector smearings. In figure(III.11) we show the $M_{\gamma\gamma}$ spectrum for the 2 categories for the 300 GeV mass hypothesis.

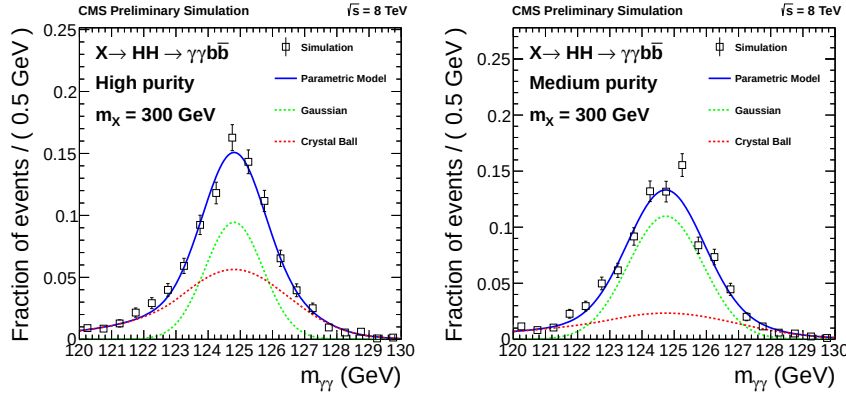


Figure III.10: Normalized signal mass shapes in the $M_{\gamma\gamma}$ spectrum for $M_X = 300$ GeV.

Background model and bias study

When estimating the background shape from data one might worry if the $M_{\gamma\gamma}$ spectrum may be shaped by the constraints in m_{jj} and $m_{\gamma\gamma jj}$. In order to fit for a signal peak in the $M_{\gamma\gamma}$ spectrum therefore one must ensure that we have a suitable background parameterization. On the other hand the use of a background model with a too large number of parameters may lead to a frequentist under-estimate. The background model can eat a potential signal by over-modelling. The desired background model is the one that gives an accurate signal strength with as few model parameters as possible.

Three functions were considered as truth models for the background parameterization:

$$\bullet \quad Ber(x) = \sum_{i=0}^N \beta_i b_{i,N}(x) \quad \bullet \quad Exp(x) = \sum_{i=1}^N \beta_i e^{-\alpha_i x} \quad \bullet \quad Pow(x) = \sum_{i=1}^N \beta_i x^{-\alpha_i}$$

The $Ber(x)$ is an expansion on Bernstein polynomials suitable in the case the spectra is not very steep. The other functions are expansions on strictly falling functions with different slopes. The order in which each of these functions is expanded is chosen using a likelihood-ratio test, treating each category separately*.

To choose a suitable fit function we test the goodness of the fit defining its bias. The truth models for this test are built by fitting the control sample data. They are then used as templates to generate 1000 toy MC datasets, following its shape. The candidate background models plus the signal model are fit to each toy dataset. We study the bias of the truth functions in the data control sample, including the mass windows in the m_{jj} and $m_{\gamma jj}$ spectra, with the latter window corresponding to a resonance mass of 300 GeV. The bias relative to how a given candidate background model performs with respect to a certain truth model is defined as:

$$\text{Bias} = \text{median} \left(\frac{\mu_{\text{gen}} - \mu_{\text{fit}}}{\sigma_{\mu_{\text{fit}}}} \right), \quad (\text{III.3})$$

where μ_{gen} is the number of signal events injected, μ_{fit} the number signal of events estimated from the truth model fit and the $\sigma_{\mu_{\text{fit}}}$ is the uncertainty on the signal strength from the signal plus background fit.

The candidate background function is considered unbiased for every truth model, if:

$$|\text{Bias}| < 0.14 \quad (\text{III.4})$$

holds, i.e. the deviation between the truth model and the candidate background function does not exceed 14% of the uncertainty of the candidate prediction.

There were more than one unbiased function with one parameter to choose in both purity categories. For consistency with the four-body spectrum discussed in the next section we chose the power law for both categories. To check for robustness in the background modelling, the bias was studied at three points on the $M_{\gamma\gamma}$ spectrum, corresponding to hypothesizediggs masses of 120, 125, and 130 GeV.

In figure (III.11) we show the unblinded fits on data on both categories.

*A the likelihood-ratio test is a statistical testing of the significance of discovery, quantified with the probability of do not find any signal on the presumption of signal plus background. This is calculated constructing the p -value variable with the likelihood-ratio between the test and truth model. The value of the p-value reveal how much the test function deviates from the false hypothesis of signal plus background to model background the defined truth model.

When the p-value is smaller than 0.05 it means the fit functions are very sensitive to signal injection, a p-value greater than 0.05 but smaller than 1.0 is taken as a sign that the fit functions have relative stability against signal hypothesis, while a p-value greater than 1.0 would mean that the signal hypothesis do not influence at all the likelihood function, and therefore the variable being tested is not suitable to study that signal, see for example [193]. The expansion is continued until the p-value of the test to be greater than 0.05.

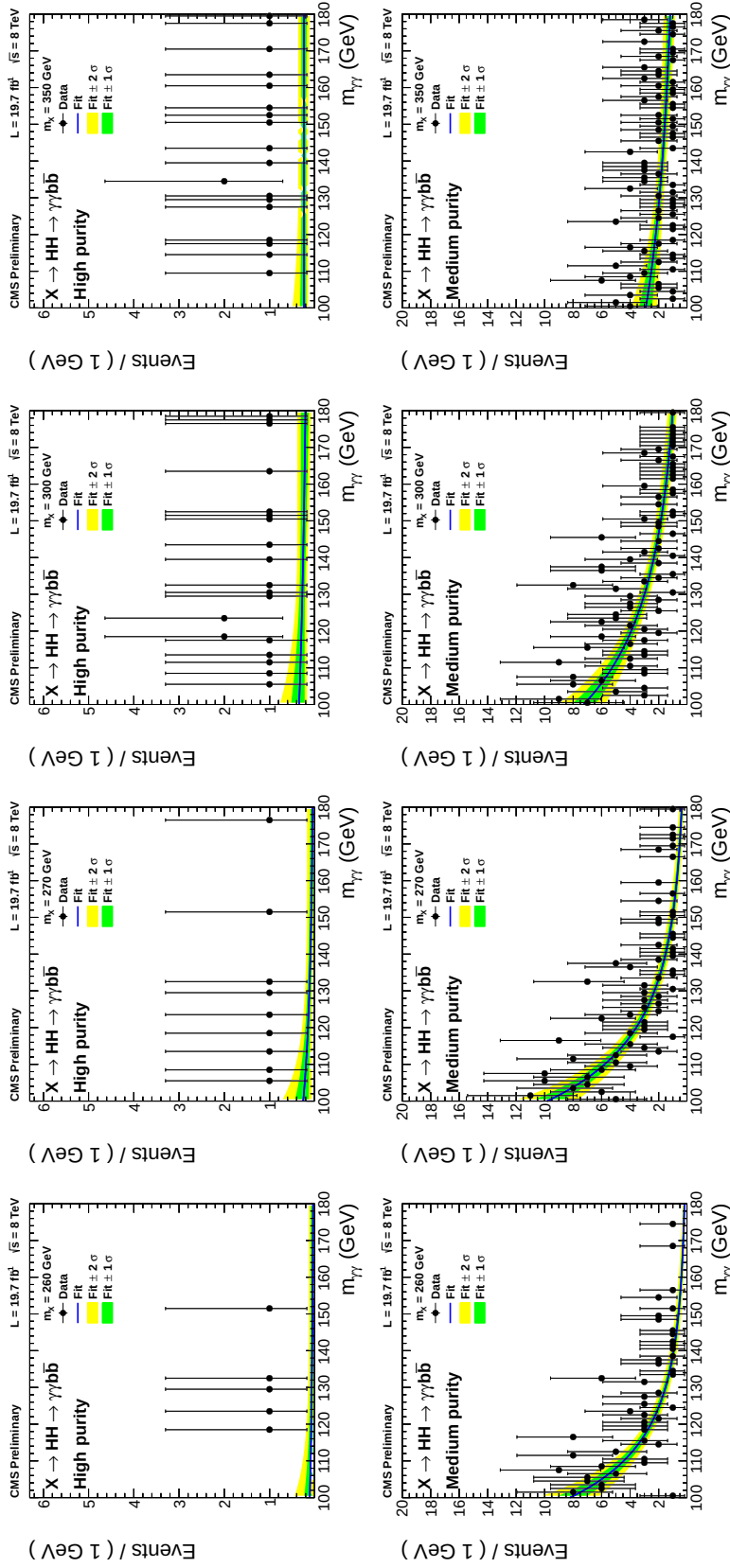


Figure III.11: Background mass shapes in the $m_{\gamma\gamma}$ spectrum on data. No kinematic fit is used. We also show the spectra of single SM higgs when possible to see it.

III.2.5.2 Sidebands signal extraction in low mass region

In this part we quickly comment on the method developed to extract signal from the four body spectra estimating the background shape from the $m_{\gamma\gamma}$ sidebands used to cross-check the results for a resonance of mass between 270 and 450 GeV.

As already stressed in this region the total background has a turn-on between 200 and 350 GeV plus a falling tail. We parametrize this behavior using a Landau distribution [194]. To assure a good background modelling we restrict the $m_{\gamma\gamma jj}$ window between 200 GeV and 700 GeV. The di-jets distribution is restricted between 90 GeV and 170 GeV.

A sideband method procedure in a given distribution spectra m_{YY} consists in taking data from a region aside the signal region (SR) to estimate the background shape in the distribution of interest m_{XX} . A transfer function $\alpha(m_{XX})$ is necessary to extrapolate data shape and normalization from the sideband to the signal region in the m_{XX} distribution. This can be derived from simulation or from an other data sideband region. The background shape on signal region is derived as:

$$F_{data}^{SR}(m_{XX}) = \alpha(m_{XX}) \times F_{data}^{sidebands}(m_{XX}) \quad (\text{III.5})$$

where the functions $F_B^A(M)$ is the result of the fit in the spectrum M in the A region for the sample B .

In the case of this analysis the sidebands and signal regions are defined in the di-photon mass distribution $m_{\gamma\gamma}$ as:

- Signal region (SR): $120 < m_{\gamma\gamma} < 130$ GeV
- Low-sideband region (LS): $100 < m_{\gamma\gamma} < 120$ GeV
- High-sideband region (HS): $130 < m_{\gamma\gamma} < 180$ GeV

The transfer function $\alpha(m_{\gamma\gamma jj})$ is determined from MC, to account for the correlations between $m_{\gamma\gamma jj}$ and $m_{\gamma\gamma}$. Only the di-photon Sherpa simulation is used. The $m_{\gamma\gamma jj}$ shape for the background is derived as equation (III.5), where in this case the weighting function is derived from the di-photon Sherpa MC as:

$$\alpha(m_{\gamma\gamma jj}) = \frac{F_{MC}^{SR}(m_{\gamma\gamma jj})}{F_{MC}^{LS}(m_{\gamma\gamma jj})} \quad (\text{III.6})$$

The upper sideband is not used mainly because the $m_{\gamma\gamma jj}$ distributions are quite different between the LS and the HS region. Despite the loss in statistics with respect to using both sidebands the α extraction is easier considering only the low sideband. To get a smooth shape for the α extraction from the fit on the SR and LS of the di-photon MC sample two million of toys events are generated following each obtained parametrization.

The $m_{\gamma\gamma jj}$ distributions for low purity events in SR and LS in the di-photon MC simulation and the same distribution in data LS are shown in figure III.12.

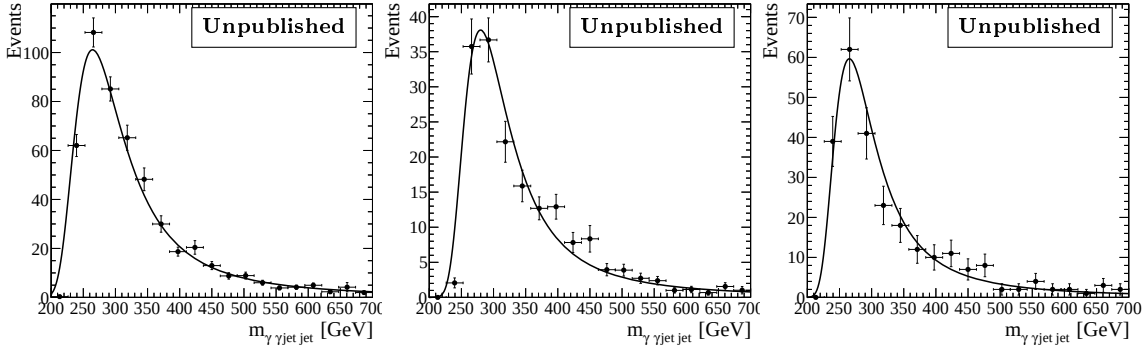


Figure III.12: Fit to the $m_{\gamma\gamma jj}$ distributions in the 1b-tag category with a Landau PDFs. Left: di-photons Sherpa MC events in the low $m_{\gamma\gamma}$ sideband. Middle: di-photons Sherpa MC events in the signal region. Right: data events in the low sideband.

The statistics available in data for the high purity category is not enough to parametrize its mass shapes. It was checked that the mass shapes of the low purity category is in reasonably agreement with the mass shapes of high purity category. Therefore the high purity category mass shapes are modelled from the templates estimated from the low purity events.

The $m_{\gamma\gamma}$ spectra is fitted with an exponential function. The normalization of the template background shape in the signal region is extracted from a fit to the data di-photon mass distribution in a sidebands region defined as $m_{\gamma\gamma} < 120$ GeV and $m_{\gamma\gamma} > 130$ GeV. The result of the fit in the SR for both high and low purity categories is shown in figure III.13 .

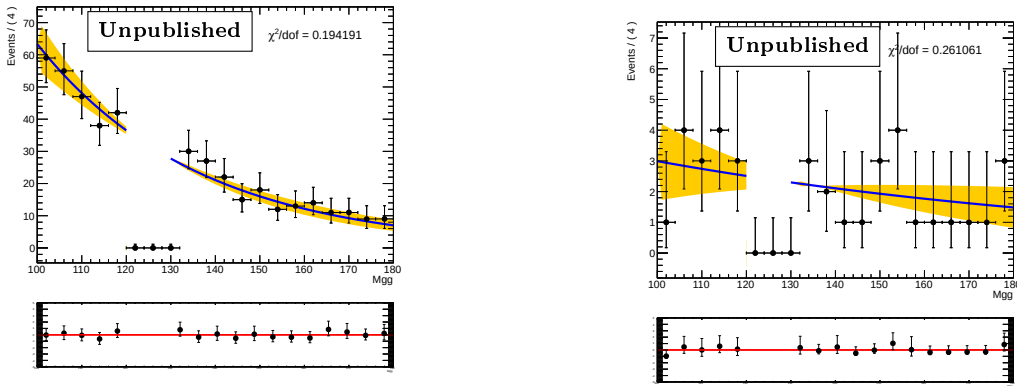


Figure III.13: Blinded fit in the $m_{\gamma\gamma}$ spectrum in the signal region. **Left:** Low purity category. **Right:** High purity category.

The sensitivity results obtained with this method were also cross-checked by extracting the α using the data control sample. The results are compatible with the ones based in the simulation.

III.2.6 The intermediate mass region, $M_R = 400 - 1000$ GeV

The analysis strategy for this mass region consists in a fit to the smoothly falling $m_{\gamma\gamma jj}$ spectrum. The number of events expected in a given $M_{\gamma\gamma jj}$ window around the resonance mass is very limited for high resonance masses and therefore a background model based on the fit to $M_{\gamma\gamma}$ does not benefit from enough events to be reliable. Together with the m_{jj} selection we also select the events with a selection in the $m_{\gamma\gamma}$ spectrum, exploring the excellent resolution in this variable. The mass cuts applied for the $m_{\gamma\gamma jj}$ fit method are:

M_{jj}	[90, 165]
$M_{\gamma\gamma}$	[120, 130]

Table III.7: Mass cuts applied for the $m_{\gamma\gamma jj}$ fit method.

Those selections were optimized against the data control sample in the low purity category. The m_{jj} mass windows used to select the resonance candidate events are optimized prior to the application of the kinematic fit.

As it was done in the low mass region the $m_{\gamma\gamma jj}$ spectrum for the signal is described by a sum of the crystal ball and a Gaussian functions. In figure (III.14) it is shown the signal fits on the 500 GeV and 1 TeV mass hypothesis for the different categories using kinematic fit.

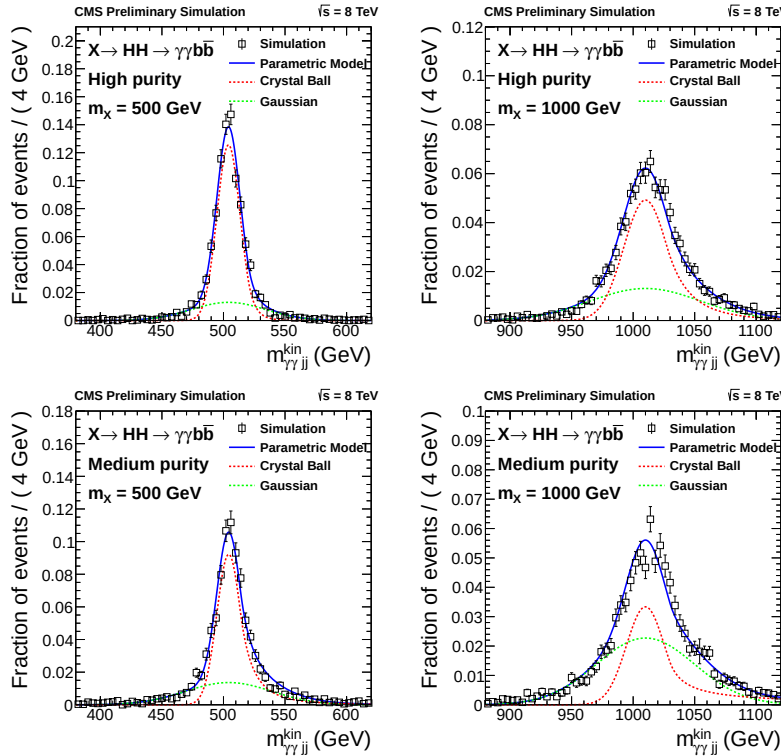


Figure III.14: Signal mass shapes in the $m_{\gamma\gamma jj}$ spectrum using kinematic fit.

Background estimation

The data is expected to be background dominated and therefore the background is estimated directly by a fit in the four body spectrum. To study the best function to fit this spectrum we repeat the bias studies procedure done for the low mass region.

For the bias study we include mass windows on the M_{jj} spectrum and kinematic fit. The four-body resonant mass is unknown therefore the bias was chosen to be evaluated injecting signal to 500, 700, and 1000 GeV mass hypotheses. Several candidate functions were tried against the data control sample. Only the unbiased functions for each category are presented. The low purity category is fitted with an extended power law:

$$1/(x^2 + b)^a \quad (\text{III.7})$$

The high purity category is fitted with the power function:

$$1/x^{2a} \quad (\text{III.8})$$

It was explicitly checked that the cuts applied on the $M_{\gamma\gamma}$ and M_{jj} spectra does not generate any local bump in background above in $m_{\gamma\gamma jj}$ spectrum when $m_{\gamma\gamma jj} > 320$ GeV [195]. The fits functions to the data are shown on the figure III.15.

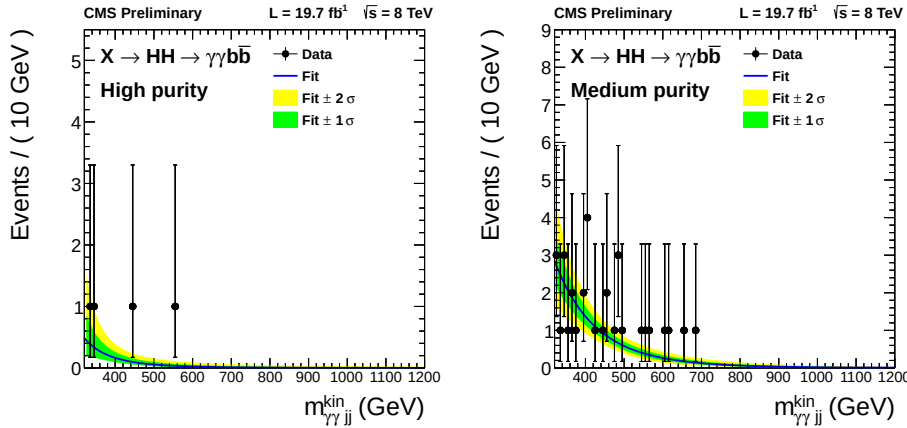


Figure III.15: Background mass shapes in the $M_{\gamma\gamma jj}$ mass spectrum on data. The kinematic fit is used. **Left:** 2-b-tag category; **Right:** 1-b-tag category.

III.2.6.1 Systematics uncertainties

The luminosity uncertainty is estimated to be 2.6% [196]. The experimental uncertainties on efficiencies, scales, and resolutions are applied to the reconstructed objects in simulated events by scaling and smearing the relevant observables and propagating the effects to the invariant masses used in the analysis. The systematic uncertainties are estimated only for the baseline signal extraction methods. They can be separated between the photon-related and jet-related ones:

For the jet related systematic uncertainties: The JES uncertainty is accounted for by varying the jet response by 1-2%, depending on the kinematics, while the JER uncertainty is accounted by varying the jet resolution by 10% [197]. An additional 1% uncertainty on the 4-body mass response accounts for potentially un-modelled effects in the high-mass region related to the partial overlap between the two b-jets coming from the higgs decay. The uncertainty on the b-tagging efficiency is obtained by varying the b-tagging efficiency and scale factor by one standard deviation in each category, the related systematics were shown to be anti-correlated between the two categories (this is the reason of the minus sign in table III.8) [198].

For the photon related systematic uncertainties: The photon-related uncertainties are taken from [199]. The offline photon selection efficiency and on the trigger efficiency normalization uncertainty is assumed to be 1%, the same as the $h \rightarrow \gamma\gamma$ analysis.

An uncertainty between 0.23% and 0.93% is considered on the photon energy resolution (PER) and between 0.12% and 0.88% on the photon energy scale (PES), depending on the photon pseudo-rapidity and the electromagnetic shower shape. They are estimated from the measurement of the SM processes $Z \rightarrow e^+e^-$ and $Z \rightarrow \mu^+\mu^-\gamma$. When $p_{T\gamma} > 100$ GeV there is no source from where to estimate the PER and PES uncertainties and the uncertainty of PES is conservatively increased up to 1%.

To account for the differences in the p_T spectrum of the signal photons and of the electrons from $Z \rightarrow e^+e^-$ used to estimate the PES and PER uncertainties an additional conservative normalization uncertainty of 5% is assumed for the high mass region, where there is no $Z \rightarrow e^+e^-$ statistics to check photon properties calibrations.

Finally an additional systematic uncertainty, taken to be 0.45 GeV, is assigned to the higgs mass, both for the signal and for the SM higgs background. It corresponds to the experimental uncertainty from the higgs mass measurement in the $h \rightarrow ZZ$ channel [200]. The quoted systematic uncertainties are summarized in the table III.8.

The SM single higgs contribution is used in the $M_{\gamma\gamma}$ fit method and therefore it is necessary to account the theoretical and experimental systematics related to it. For the PDF and QCD scale uncertainties we had used the CERN Yellow Report recommendations [108]. The other sources of uncertainty are associated to the description of the underlying event and to the contamination of the different higgs production modes to our final state. The impact of this description was studied in the context of the Vh with hadronic WBF signal contamination to the ggh component in $h \rightarrow \gamma\gamma$ search [187] and it was found to be 25% for ggh and 8% for WBF production. This study was not done for Vh and tth . The underlying event uncertainty to the Vh and tth components are taken to be respectively 8% and 25%.

The systematics uncertainties in the signal events selection related to the gluon splitting modelling is estimated from data and simulation observed difference in $t\bar{t}bb$ and $t\bar{t}qq$ [187]. We apply the same gluon splitting systematics found for ggh production to all the production modes. The theory systematics considered in the SM single signal are summarized in table (III.9).

Common normalization uncertainties	
Luminosity	2.6%
Di-photon trigger acceptance	1.0%
fit $m_{\gamma\gamma}$ - Background dominated	
Normalization uncertainties	
Photons selection acceptance	1.0%
"b-tag" eff. uncertainty 2 b-tag cat	4.6%
"b-tag" eff. uncertainty 1 b-tag cat	-1.2%
m_{jj} and $p_{T,j}$ cut acceptance (JES & JER)	1.5%
$m_{\gamma\gamma jj}$ cut acceptance (PES $\oplus$ JES & PER $\oplus$ JER)	2%
Shape uncertainties	
Parametric scale shift (PES $\oplus$ m_h uncertainty)	$\frac{\Delta m_{\gamma\gamma}}{m_{\gamma\gamma}} = 0.45 \oplus 0.35\%$
Parametric resolution shift (RES)	$\frac{\Delta\sigma}{m_{\gamma\gamma}} = 0.25\%$ $\frac{\Delta\sigma}{\sigma_{\gamma\gamma}} = 22\%$
fit $m_{\gamma\gamma jj}$ - Background dominated	
Normalization uncertainties	
Photons selection acceptance	1.0%
"b-tag" eff. uncertainty 2 b-tag cat	5.3%
"b-tag" eff. uncertainty 1 b-tag cat	-1.8%
m_{jj} and $p_{T,j}$ cut acceptance (JES & JER)	1.5%
$m_{\gamma\gamma}$ cut acceptance (PES & PER)	0.5%
Extra high pt norm. uncertainty	5.0%
Shape uncertainties	
Parametric abs. shift (PES $\oplus$ JES)	$\frac{\Delta m_{\gamma\gamma jj}}{m_{\gamma\gamma jj}} = 0.45 \oplus (0.8 \oplus 1.0) = 1.4\%$ $\frac{\Delta\sigma}{\sigma_{\gamma\gamma jj}} = 10\%$
Parametric shift (PER $\oplus$ JER)	

Table III.8: Systematic uncertainties by fit strategy.

production	PDF	QCD scale	underlying event	gluon splitting
$t\bar{t}h$	+8.1%,-8.1%	+3.8%,-9.3%	25%	40%
ggH	+7.5%,-6.9%	+7.2%,-7.8%	25%	40%
WBF	+2.6%,-2.8%	+0.2%,-0.2%	8%	40%
WH	+2.3%,-2.3%	+1%,-1%	8%	40%
ZH	+2.5%,-2.5%	+3.1%,-3.1%	8%	40%

Table III.9: Uncertainties for higgs Production. Both gluon fusion and $t\bar{t}$ associated production mechanisms are gluon initiated. WBF and V associated production are quark initiated.

III.2.7 Results

In this section we present here the analysis results based on the strategies discussed above. The high purity category drives the sensitivity results on the low mass region therefore the sensitivity obtained from the $m_{\gamma\gamma}$ fit using only the high purity category is presented as a second result.

As a reminder we summarize the strategy methods used for each mass hypothesis both for the nominal results and the cross-checks in the table bellow.

Mass (GeV)	260	270	300	350	400	450	500	550 to 1100
$M_{\gamma\gamma}$ fit (+ higgs)	B	B	B	B	S	C	C	-
$M_{\gamma\gamma} jj$ fit	-	-	-	-	B	B	B	B
$M_{\gamma\gamma}$ fit (2 b-tag only)	S	S	S	S	S	-	-	-
$M_{\gamma\gamma}$ Sidebands (statistical only)	-	C	C	C	C	C	-	-

Table III.10: Methods used for the analysis. B = baseline result. C = cross check. S = second result.

For the interpretation of the results we calculated the inclusive production cross section as described on section (I.5). We assume a branching ratio of 25 % for a radion decay to a di-higgs pair with $m_X > 300$ GeV. This value is taken from [85, 102, 103]. For the volume suppression of bulk KK-graviton coupling with gluons we had used $d_g = 0.0137$ (that is equivalent to assume $kr_c = 11$, see figure I.3). Along the approval process of the [201, 202] and [203] there was no time to be more precise in these two points, the numerical differences however are not significant (see figures I.12 and I.3) and will be corrected in future iterations.

The comparison of the alternative methods are shown in the left part of figure (III.16). All the methods are compatible within the one sigma band.

The right side of this figure we show the impact of the different elements of the analysis on the nominal results. The kinematic fit and the combination with the low purity category improves considerably the sensitivity of four body fit. The use of low purity category brings a modest improvement for the sensitivity of low mass region.

The left part of figure (III.17) shows the signal and SM single higgs background components cut flow evolution under the analysis selections. In the right part of this figure it is presented the final signal efficiencies for the benchmark signal separated by the b -tag purity categories.

The photon identification efficiency is better for boosted photons which breaks the mass scale invariance in the resonant signal efficiencies. The jet identification includes $p_{T_{jet}} > 25$ GeV cut, which kills most of the single higgs background. The impact of the mass cuts for the $m_{\gamma\gamma jj}$ fit are less stringent to signal than the mass cuts for the $m_{\gamma\gamma}$ fit. This is expected since the signal resolution is better in the $m_{\gamma\gamma}$ spectrum. The signal efficiencies are well split in the purity categories.

Apart from the $t\bar{t}H$ component the efficiencies of the single SM higgs drops considerably with the jet identification and even more with the b -tag requirement. The mass cuts in the 300 GeV hypothesis however makes the efficiencies for all the SM single signal to be below 4%. The full impact of including the single SM higgs as background in the sensitivity including all the relevant systematics was found to be negligible, see table below.

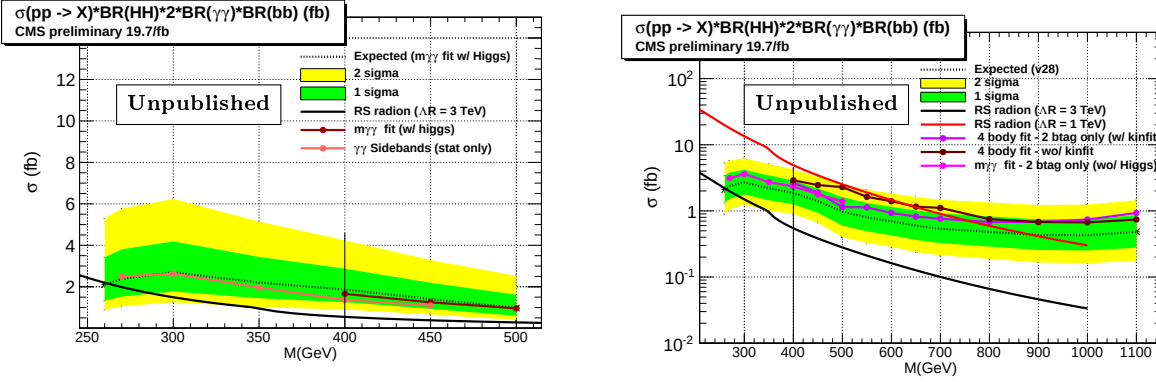


Figure III.16: **Left:** Comparison between the cross check methods. **Right:** Impact of the different ingredients of the signal extraction procedure on the expected limits. In this figure radion cross section is wrongly extended up to masses lower than 300 GeV.

Mass(GeV)	impact of SM single higgs
260	2.22%
270	2.61%
300	2.87%
350	2.81%

Table III.11: Impact of including the SM higgs as BKG on the results, including the relevant systematics.

The analysis was found to be statistics limited for all the inspected resonance mass range. The full set of systematic uncertainties worsen the expected limits by 1.7% at most.

The benchmark model used to derive the nominal results was the radion model simulated with **MadGraph5+Pythia6**, without any matching procedure for the hard jets. To test the results sensitivity to spin hypothesis and generators modelling the efficiencies obtained with the benchmark radion model was compared with the efficiencies obtained with a KK-graviton benchmark generated under same conditions and also with a heavy higgs generated and hadronized with **Pythia6**, using the default MSSM model implemented on this generator. In table (III.12) we compare the efficiency of the analysis in the different models quoting the statistical-only errors* and ignoring theory uncertainties.

From the above table it is possible to conclude that even without including theory uncertainties the performance of the analysis on all the tested models are similar. And therefore the analysis does not hold strong model-dependence on spin hypothesis or showering model. A more detailed study about the showering model and spin dependence is in progress.

*Calculated with the uncertainty formula from binomial model [204], the formula for the efficiency eff is $\delta\text{eff} = \sqrt{\frac{\text{eff}(1-\text{eff})}{N_{\text{processed}}}}$.

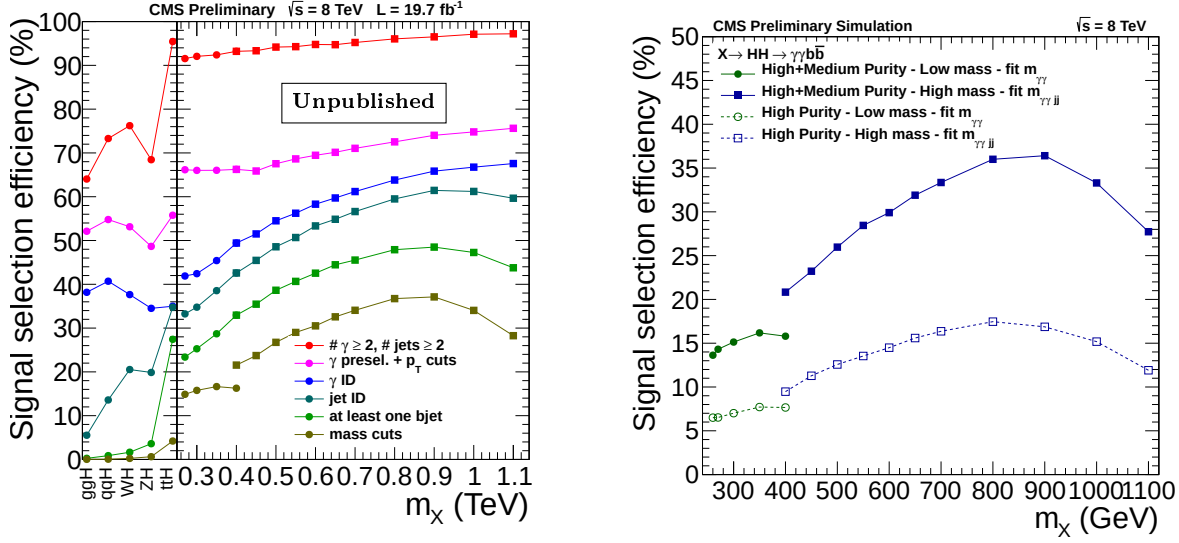


Figure III.17: **Left:** Signal selection efficiency flow under the analysis cuts. The mass cuts applied to the SM single higgs signal are the ones referent to the $M_{\gamma\gamma}$ fit in the 300 GeV signal hypothesis. **Right** Signal efficiency for baseline signal extraction method split in b -tag purity categories.

Mass (GeV)	MSSM		Radion		Graviton	
260	10.42	± 0.06	-	-	-	-
270	-	-	14.03	± 0.25	-	-
300	11.69	± 0.06	14.83	± 0.25	12.23	± 0.15
350	12.68	± 0.06	15.87	± 0.27	-	-
400	-	-	15.5	± 0.26	-	-
450	-	-	19.14	± 0.28	-	-
500	-	-	22.42	± 0.3	23.18	± 0.19

Table III.12: Efficiencies for different benchmark models and its statistical-only uncertainties. b -tag scale factors are included. The empty spaces means the point was not tested.

Figure (III.18) shows the expected and observed sensitivity results for the baseline methods. For the 400 GeV mass point the results of both low mass region and high mass methods are considered.

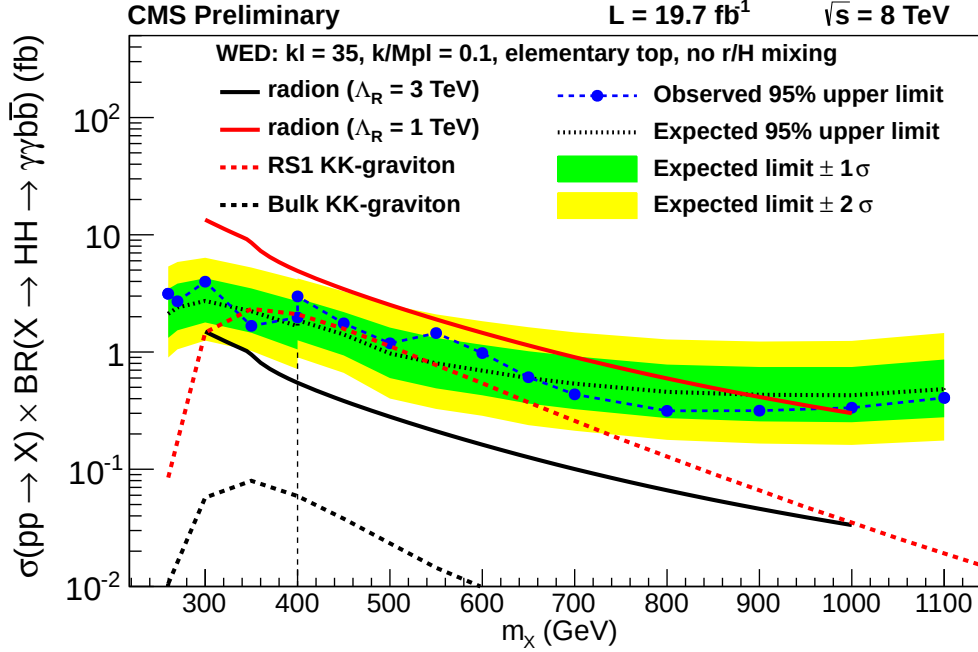


Figure III.18: Expected limits for baseline methods.

Figure (III.19) shows side by side a zoom of sensitivity results combining the purity categories and the case where only the high purity category is used.

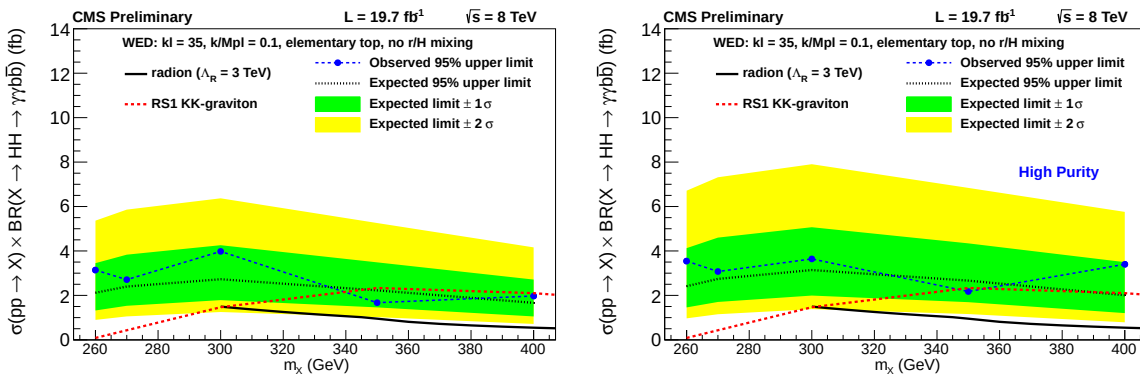


Figure III.19: Expected limits for baseline methods for low mass region

As it was already noted from figure (III.16) the low purity category does not change significantly the expected sensitivity. The low purity category however has more statistics than the high purity category and therefore its inclusion in the calculation of the exclusion limits decreases the statistical errors, narrowing the $\pm 1, 2\sigma$ bands.

III.3 $pp \rightarrow X \rightarrow h(b\bar{b})h(b\bar{b})$ at CMS

In this subsection we perform a non-extensive description of the search on each regime with the intent of feedback the studies presented in the section (II). The contribution made on these analyses was with the benchmark signal and theoretical interpretation.

The intrinsic experimental difficulty to apply the scale invariant tagger described in section (II.4) is related to the trigger strategy definition. Each experimental trigger is chosen and optimized accordingly with the physical goal of each search. And also the best jet and b -tag strategy for each regime will depend on the signal topology.

In subsection (III.3.2) we present the analysis designed to exploit exclusively the boosted regime [205]. In the subsection (III.3.1) we present the analysis restricted to the mass region between 250 GeV and 1100 GeV, when the fraction of the signal with four resolved jets is non negligible [206].

III.3.1 Boosted: $M_X > 1$ TeV

The search for fully hadronic di-higgs resonance in boosted regime ($r_M > 4$) is characterized by two jets with non-trivial substructure and the presence of b -quarks the final state. The main background is expected to be the QCD multi-jet production, the MC considered for the selection strategy studied was simulated using **MadGraph5** and **Pythia6**.

The search combines 19.6/fb of data collected combining two high level triggers, one selecting events with at least two AkT05 jets and other selecting multi-jets events with a minimum $H_T(\sum_{jets} p_T)$. A basic pre-selection is applied to the events. The part of selection acting in the stand-alone AkT05 triggered jets is:

$$p_{T_j} > 30 \text{ GeV}, \quad |\eta_j| < 2.5. \quad (\text{III.9})$$

The events are selected if $m_{jj} > 890$ GeV and $|\Delta\eta_{jj}| < 1.3$. This last angular cut is expected to remove t -channel QCD production while it keeps the s -channel $X \rightarrow hh$ production [207]. This cut is equivalent to the δy_{max}^{res} selection used in section (II.4). After the selection the AkT05 jets are matched to CA08 jets.

The N-subjettiness is used to tag the CA08 jets as possessing non-trivial substructure. The sub-jets are defined with the exclusive-kT [208] algorithm, a parameter τ_N is defined for each sub-jet as [209]:

$$\tau_N \equiv \frac{1}{d_0} \sum_k p_{T_k} \min(\Delta R_{1,k}, \dots, \Delta R_{N,k}) \quad (\text{III.10})$$

where (k) runs over jet constituents, $\Delta R_{N,k}$ is the distance between the constituent k and the sub-jet N , and $d_0 \equiv \sum_k p_{T_k} R_0$ where R_0 is the radius used on the original jet clustering.

The discriminating variable used when applying n-subjettiness is the ratio τ_N/τ_{N-1} . Differently from the mass-drop discriminator used on section (II) the τ_N/τ_{N-1} variable is sensitive to details of the radiation pattern and therefore is able to differentiate its core structure. The $N = 2$ discriminates a two core structure, while $N = 3$ is suitable for three core objects like a boosted top quark. The radius used for the sub-jets clustering is chosen as the one that produces a minimum value for the τ_N/τ_{N-1} when the re-clustering returns exactly N jets.

This analysis uses the ratio τ_2/τ_1 where the τ_1 and refers respectively to the hardest (second hardest) prong hadronic object. The mass-drop procedure was shown to be less discriminating to tag a boosted longitudinally polarized W [210]. The left part of figure (III.20) shows the distribution of the n-subjetiness parameter $\tau_{12} \equiv \tau_2/\tau_1$ in signal, background and data. Two purity categories are defined based in this variable. A high purity category is defined by $\tau_{12} < 0.5$ and a medium purity category is defined by $\tau_{12} \in [0.5, 0.75]$.

To clean the CA08 jets from soft radiation a pruning process is used [211]. In contrast with the filtering procedure used on section (II.4) the pruning is a process that clean soft and wide angle radiation radiation while re-clustering the jet. There are two parameters on pruning definition: the biggest distance allowed between two objects and how soft in p_T an object relatively with other. Pruning is used only for jet mass measurement. The right part of figure (III.20) shows the pruned jet mass distribution for signal and background. The pruning introduces a low mass "peak" on signal (around 20 GeV). This happens when the pruning algorithm rejects too many particles from the jet (over-pruning effect). This effect turns to be rather aggressive to reject a lot of QCD background even holding this small inefficiency. This low mass peak is also seen on tagging boosted weak bosons [202]

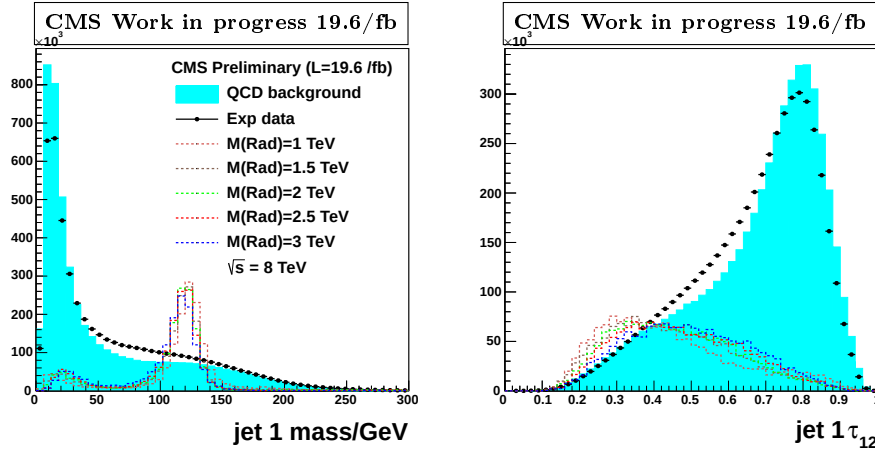


Figure III.20: Left: Pruned jet mass. Right: The $\tau_{12} \equiv \tau_2/\tau_1$ n-subjetiness variable. The background MC is normalized to the quantity of data events. The signal events are multiplied by an arbitrary number such its shape is visible in the plot.

The energy loss induced by the pruning process and also the neutrinos that are normally present on B hadron decays the higgs jet mass peak is at 120 GeV, lower than its generated value (125 GeV). A jet is tagged as coming from a boosted higgs if its pruned mass lies between 110 GeV and 130 GeV. This window selection was optimized against the MC background using as target the statistical-only signal sensitivity calculated in the di-jet mass spectrum.

The efficiency of higgs reconstruction after this selection is given in table (III.13). As effect of over-pruning very boosted particles and tracks the efficiency of the higgs-tag decreases with the resonance mass for $m_X > 1.5 \text{ TeV}$.

mass (TeV)	1 h-tag (%)	2 h-tag (%)
1.0	34.3	10.8
1.5	32.1	11.8
2.0	29.7	10.4
3.0	26	4.2

Table III.13: Efficiency of higgs tag with the resonance mass.

As it was done in the high mass region of the inclusive $X \rightarrow \gamma\gamma b\bar{b}$ the signal extraction is made by a shape analysis the in the di-jet mass spectra, assuming the selected data to be background dominated (see section III.2.5). The left part of figure (III.21) shows how the QCD di-jet background drops with the higgses selection. When two higgs mass tags are required the multi-jet background decreases by four orders of magnitude. The higgs mass cut induce a turn on around $m_j = 800$ GeV when one request a high mass cut to the pruned jet jet $M_{j,Prun} > 100$ GeV. We also see that the two higgs tag does not produce any unexpected bias in this spectrum and that from 900 GeV and therefore from this mass we can treat the di-jet shape with a strictly falling function. The right part of this figure shows the di-jet spectrum for signal after all selections.

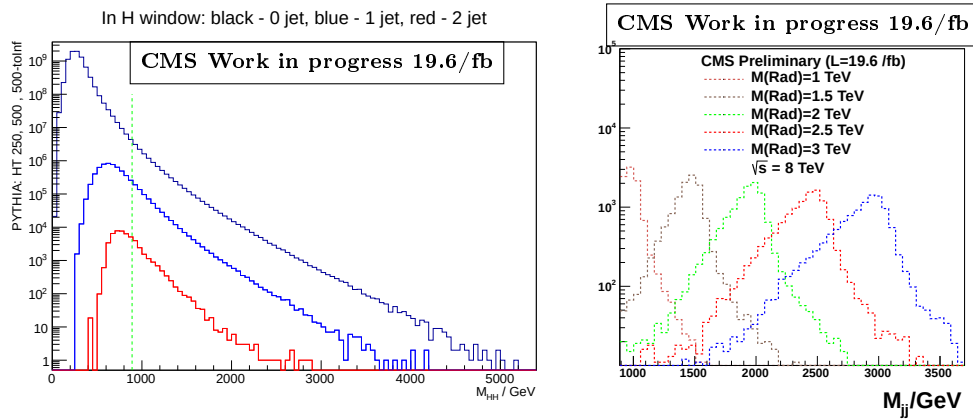


Figure III.21: **Left:** Shape normalized mass distribution of the reconstructed X resonance with mass 1.0 TeV, 1.5 TeV, and 2.0 TeV after full selection. **Right:** The M_{jj} distribution for QCD Pythia6 for 0 higgs tag, 1 higgs tag and 2 higgses tag.

To further reduce the QCD background and improve sensitivity of the analysis a dedicated b -tag algorithm designed to boosted topology is being tested [212]. The search for a heavy fully di-higgs resonance in 8 TeV data using sub-jet b -tag turns out to be statistics limited and therefore a more involved method for signal extraction and background estimation needs to be used.

III.3.2 Un-boosted: $m_X < 1$ TeV

A signal from the inclusive production of a di-higgs resonance in the $b\bar{b}b\bar{b}$ final state when this resonance has mass between 250 GeV and 1 TeV is characterized by four resolved b -jets. The four jet signature is particularly challenging to trigger due to the large QCD multi-jets contribution.

A multi-jet trigger with very low p_T threshold and any other requirement would record an overwhelming quantity of background events. A trigger exploring the b -tag information at HLT level was designed to store multi-jet events allowing low p_T thresholds while keeping acceptable the QCD background rate. Trigger performance was tested against SM processes that contain one or more b -quarks on final state, namely $t\bar{t}$ and one hadronic weak boson produced in association with jets (Vb , $Vb\bar{b}$ and V + light jets).

Each recorded event is required to have four AkT05 jets. The event is recorded if at least two calorimeter jets have $p_T > 80$ GeV or if four jets have $p_T > 30$ GeV, where four jets are required to be PF reconstructed. The triggered events are also required to have at least one of the jets tagged with a tight CSV b -tag and other jet tagged with at least medium CSV b -tag, and to have at least two tracks associated to a reconstructed primary vertex. This trigger is operating since May 2012 and had recorder 17.9/fb of data.

At pre-selection level the events are required to have at least two AkT05 jets with $p_T > 90$ GeV or the four jets with $p_T > 40$ GeV. The four jets must be b -tagged with medium CSV or at least two medium and two loose CSV tags. The four jets are paired as higgs candidates h_1 and h_2 minimizing the $|\Delta m_{h_1, h_2}|$. The selection cuts in the di-higgs system are:

$$|m_{h_1} - m_{h_2}| < 50 \text{ GeV}, \quad |m_{h_{1,2}} - 125 \text{ GeV}| < 35 \text{ GeV}. \quad (\text{III.11})$$

The selection in the individual higgses window and in the difference two higgs candidates is looser than the 15% of tolerance suggested in section (II.4). To maximize the significance of signal based on higgs tag performance two search regions are defined and different selection are used in each region:

Low mass region, $m_X \in [250, 450]$ GeV: The minimum p_T cut allowed by trigger is used ($p_T > 30$ GeV). All the four jets are required to be b -tagged, decreasing the higgs mis-tag rate due wrong jets pairing.

High mass region, $m_X > 450$ GeV: In this mass region for the candidate resonance the boost of each higgs is between $\gamma \sim 1.8$ and 3.5. Therefore the opening angle between the decay products of each higgs is expected to be lower than $\pi/2$. The jets are still paired minimizing the invariant mass difference of the pairs, but requiring first them to be separated by a angle $< \pi/2$. A harder p_T threshold is also required ($p_T > 40$ GeV). In this search region there is more restrictive cuts and signal selection is better

The background shape is not strictly falling in the four jet invariant mass in the mass window between 250 GeV and ~ 450 GeV, experiencing a kinematic peak around 400 GeV. Since the expected background peaks in the same region of the expected signal for most of the mass hypothesis probed it is not possible to estimate background from a fit from the full data.

This search extract signal form the four bodies mass spectra using the sidebands method in the invariant masses of the two higgses candidates (m_{h_1} and m_{h_2}) to estimate background (see section III.2.5.2). Similarly to the $pp \rightarrow X \rightarrow h(\gamma\gamma)h(b\bar{b})$ analysis a kinematic fit is applied to improve each di-jet higgs candidate signal resolution (see sub-section III.2.4.2).

The signal region is defined as a circle in the 2D plane centred in $(m_{H_1}, m_{H_2}) = (125, 125)$ GeV with 15 GeV of radius. The sideband region is defined as two quadrants of the region between the circle centered in $(m_{H_1}, m_{H_2}) = (125, 125)$ GeV with 35 GeV radius and the signal region. To avoid to pick the events where the mass of the two higgses is too asymmetric only the quadrants aside the SR are used for the sidebands background estimation.

The kinematic correlations to normalize the kinematics of sideband region to signal region were estimated from a data validation region (see section III.2.5.2). The validation signal region and its correspondent sideband is defined as a mirror of the true signal region, but centered in $(m_{H_1}, m_{H_2}) = (90, 90)$ GeV.

The right part of figure (III.22) shows the (m_{H_1}, m_{H_2}) scatter plot for the resonant signal hypothesis with $m_X = 600$ GeV superimposed with the above defined regions. The left part of this figure shows the blinded data distribution in the same plane.

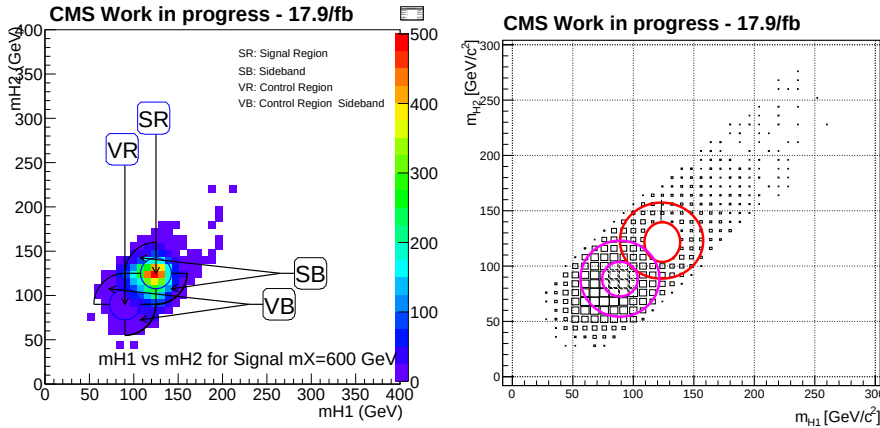


Figure III.22: Distributions in the scatter plane (m_{H_1}, m_{H_2}) after pre-selection. **Left:** The 600 GeV resonance signal, superimposed with the regions defined for background estimation. **Right:** Data, superimposed with the circumferences around the validation and signal regions as reference.

The $t\bar{t}$ background is found to be a non-negligible component. This component is simulated with *Pythia6*. In figure (III.23) we show the (m_{H_1}, m_{H_2}) scatter plot for the $t\bar{t}$ background. There is a part of the peak density that enters the signal region and it is not represented in the sidebands region. One might worry that this background component can fake a low mass signal, therefore this component is subtracted from data. This procedure is similar to the account the SM single higgs in the low mass region analysis for the $pp \rightarrow \gamma\gamma b\bar{b}$ channel. The right part of figure (III.23) shows the four body mass spectrum in the SB region before and after the $t\bar{t}$ subtraction. The final background shape is not modified in a significant way.

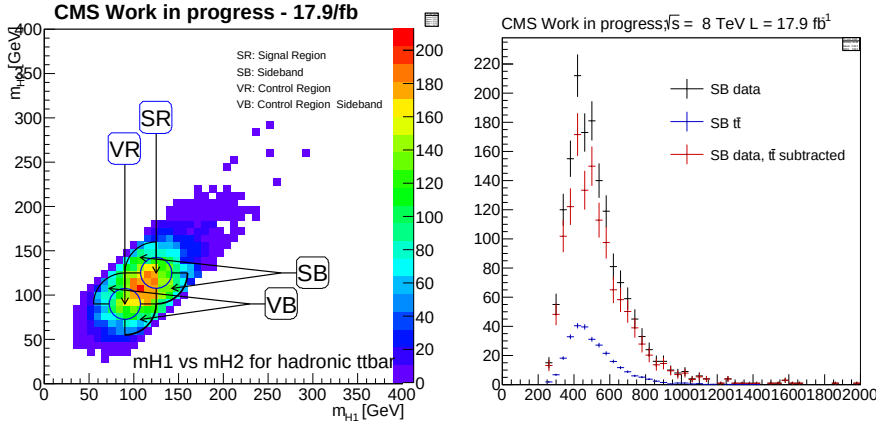


Figure III.23: **Left:** $t\bar{t}$ background in the scatter plane (m_{H_1}, m_{H_2}). **Right:** Data in SB region in four body spectra before and after $t\bar{t}$ background subtraction.

The specific sources of systematic uncertainties are related with the trigger acceptance, jet energy scale and resolution, b -tag scale factors, luminosity, mass selections and $t\bar{t}$ background subtraction method.

We thank Caterina Vernieri, Andrea Rizzi, Jacobo Konigsberg and Souvik Das for providing the figures necessary for this subsection discussion.

III.4 Interpretation of the di-higgs LHC results

In this section we compare the significance found by all the present LHC results in the di-higgs channel and interpret the best achieved sensitivities by mass hypothesis in the parameter space of radion and KK-graviton particles.

At the time of this thesis deliver three di-higgs analyses were public, all of them performed exclusively with 8 TeV data. Those are the CMS $hh \rightarrow \gamma\gamma b\bar{b}$ presented in section (III.2), the ATLAS $hh \rightarrow b\bar{b}b\bar{b}$ analysis [149] and the CMS analysis combining the di-higgs channels with multi-leptons and di-photon candidates [213].

To be able to compare the sensitivities of the analyses in the different di-higgs channels we assume the higgs branching fractions to be SM like (see table II.1). Figure (III.24) shows the expected cross section limits from the public di-higgs assuming the SM higgs branching fractions*.

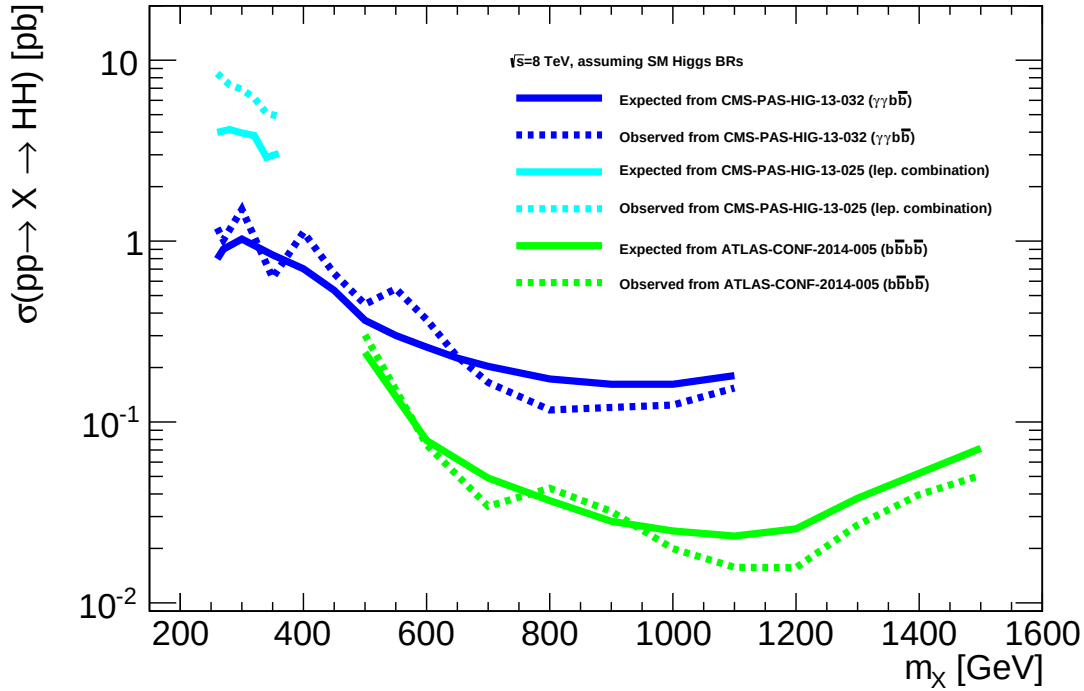


Figure III.24: Comparison of the limits of the public for the channel $X \rightarrow hh$ assuming SM higgs branching fractions.

The sensitivity reached by the di-higgs search in the multi-leptonic combination [213] is not competitive with the other two results. The expected and observed sensitivity of the ATLAS $hh \rightarrow b\bar{b}b\bar{b}$ analysis is better than the ones found by the CMS $\gamma\gamma b\bar{b}$ analysis in all the region for the resonance mass that was probed (from 500 GeV to 1.1 TeV). The CMS $\gamma\gamma b\bar{b}$ search was designed with the di-photon trigger. The lower p_T threshold of the di-photon trigger allows this channel to probe lower resonance masses.

If we extrapolate the ATLAS results in the $b\bar{b}b\bar{b}$ channel for mass hypotheses lower than 500 GeV we expect the sensitivity of the $\gamma\gamma b\bar{b}$ channel start to become better. This

*We thank Philip Hebda for this figure.

is a consequence of the large QCD background in the four jet final state and also from the fact of the combinatorics problems that increase the rate of wrong mass pairings in the fully hadronic case not be present in the $\gamma\gamma b\bar{b}$ channel.

No new physics was found with the current sensitivity, therefore the above results constrain the parameter space of the theories as presented in section (I). For the interpretation of the radion we assume a branching ratio of 25 % for a radion with $m_R > 300$ GeV. Differently from section (III.2.7) however we use the value for the volume suppression of bulk KK-graviton coupling with gluons from figure (I.3).

Based in the signal resolutions found by the $\gamma\gamma b\bar{b}$ analysis for the 1 TeV signal the negligible width hypothesis is reasonable for the radion case in no mixing scenario for a radion UV scale of $\Lambda_R \sim > 1$ TeV (see figure 7 of [83]). And also for a KK-graviton hypothesis for $\tilde{k}$ up to ~ 0.4 in RS1 scenario and up to ~ 1 in bulk scenario (see figure I.9).

However we draw the exclusion plots in a wider range of parameters. A wider resonance is not expected to change too much the results of the analyses. In the $\gamma\gamma b\bar{b}$ analysis in the low mass region the signal extraction is made in the $\gamma\gamma$ mass spectra. The detector effects in the reconstructed resonance mass are expected to be large in the $b\bar{b}b\bar{b}$ channel, which is confirmed by the ATLAS di-higgs results. A dedicated study of the width effects in the analysis in a given range of the gravity particles parameters space however deserves to be performed, and at least in the presented CMS analyses is in progress.

The gravity particles branching ratios are independent from the $\tilde{k}$ and Λ_R parameters. We summarize the parameter dependency in the KK-graviton and radion production cross section:

$$\sigma_{pp \rightarrow h_{\mu\nu}}(\tilde{k}, m_{Gr}) \propto \left(\frac{\tilde{k}}{m_{Gr}} \right)^2 \quad (\text{III.12})$$

$$\sigma(pp \rightarrow r)(\Lambda_R, m_R) \propto \left(\frac{1}{\Lambda_R} \right)^2 \quad (\text{III.13})$$

In figure (III.25) we show the impact of the best experimental sensitivities from the di-higgs channels in the KK-graviton parameter space in both RS1 scenario and the bulk scenario as proposed in [34, 89]. For the bulk scenario we consider the $\tilde{k} \equiv k/\overline{M}_{Pl}$ value up to 2, following the suggestion of [34].

It is not a surprise that the parameter $k/\overline{M}_{Pl}$ is more constrained when we assume the RS1 scenario. In this case the branching fraction of the KK-graviton to a pair of higgses is about two orders of magnitude lower than the considered bulk scenario, however its production cross section is about four orders of magnitude larger.

We compare the $\tilde{k}$ sensitivities in the bulk scenario with the results obtained by the boosted VV channels. The search for boosted pairs of weak bosons in the semi-leptonic [214] starts in $m_X > 600$ GeV. The parameter space constraint from this search is compatible with the one obtained from the ATLAS results in the $b\bar{b}b\bar{b}$ channel. For higher masses however the di-higgs channel is more stringent. In the RS1 scenario the boosted WW and ZZ searches [202] set more stringent limits to the $\tilde{k}$ parameter than the ATLAS resolved di-higgs analysis.

Given that three references [34, 86, 97] used by both ATLAS and CMS experiments to interpret results in KK-graviton had some mistakes we prefer not to comment about the comparison with results from other channels before to derive the parameter space limits ourselves. This compilation of results however was not ready up to the moment of deliver this thesis.

Figure III.25: The expected limits from the best limits from the public di-higgs searches in the KK-graviton parameter space. **Up:** RS1 scenario. **Down:** Bulk scenario under composite top hypothesis. We had assumed the effects of finite width to be negligible in the experimental results in all the probed parameter space. The higgs boson is SM like and have mass of 125 GeV. The yellow and green bands on the expected limit represent ± 1 and $\pm 2 \sigma$ variations. The lines correspond to the $pp \rightarrow X \rightarrow hh$ cross section isocurves. The dark gray region is excluded region. The limits for $M_{Gr} < 500$ GeV are from the CMS $\gamma\gamma b\bar{b}$ analysis and for $M_{Gr} \geq 500$ GeV are from the ATLAS $b\bar{b}b\bar{b}$ analysis.

In figure (III.26) we show the correspondent limits in the radion mass according with the UV scale Λ_R and ignoring the possibility of radion-higgs mixing. The presented di-higgs channels reaches to exclude $\Lambda_R \sim 2$ TeV in a mass range of $M_R = 300$ GeV to 600 GeV. Imposing more stringent constraint than the limits compiled by [102] (driven by the $ZZ \rightarrow 2l2\nu$ channel) in the mass region of $m_R > 450$ GeV, and similarly in the range $m_R \in [300, 450]$.

When the radion mass is in between 600 GeV and 1 TeV a radion UV scale $\Lambda_R > \sim 2.5$ TeV is excluded by the ATLAS $b\bar{b}b\bar{b}$ analysis. We note however that we had used conservative assumptions in the derivation of the radion cross section in this range of masses.

Updated bounds assuming the radiative corrections to the radion cross section in a more complete picture is in progress.

Figure III.26: The expected and observed 95% CL upper limits from the CMS $pp \rightarrow \gamma\gamma b\bar{b}$ channel in the radion parameter space, ignoring radion-higgs mixing. The color scheme and assumptions follows the caption of figure (III.25).

Conclusions

In this thesis we revised the physics of gravity particles that emerge as quantum excitations of the space-time metric in a world with a Warped Compact Extra Dimension. We discussed the approximations commonly applied in the theory parameter space. The feasibility of the direct detection of those particles in the di-higgs channel for different stages of the LHC program was investigated. The resonant di-higgs channel was studied in the context of CMS experiment in the 8 TeV LHC run using the model for gravity particles as benchmark.

The massless graviton particle is the manifestation of the metric perturbations in the 3+1 infinite dimension. The KK-graviton is its first massive excitation. The structure and strength of the KK-graviton couplings to fermions, gauge bosons and higgs boson depends of the behavior of these matter fields in the extra dimensional background. The KK-graviton mass scale for a reasonable range of the theory parameter is expected to be at the TeV scale and therefore we could expect this particle to be directly detected by the LHC experiments.

The radion is the zero mode of the metric perturbations in the direction of the compact extra dimension. A natural non-zero mass for the radion field is a by-product of a stabilization mechanism for the size of the compact dimension. This stabilization mechanism needs to modify the WED metric background with respect with to the free gravity solution. This conclusion is based on studying an explicit example where the stabilization mechanism is based on the presence of a scalar bulk field with a specific potential.

A characteristic fingerprint of the gravity particles at colliders would be the direct detection of a resonant signal in the di-higgs channel.

Based on the discovery of the higgs-like boson we investigated the LHC sensitivity for a resonance that decay on a higgs pair. At Monte Carlo level we studied the case where a very narrow resonance produced inclusively in a proton collider via gluon fusion in the di-higgs channel in the $b\bar{b}b\bar{b}$ final state. The experimental capabilities for identification of physics objects are estimated from the ATLAS and CMS literature.

With a semi-realistic study we confirmed that the higgses start to be boosted when the mass of the mother resonance is around 1 TeV. Beyond this resonance mass the products of each higgs merge in a single fat jet and therefore the techniques to identify jets substructure are a crucial element for the search of a heavy resonance search in the fully hadronic di-higgs channel.

A resonance tagging algorithm combining resolved and boosted analysis techniques allow the search to cover resonance masses where boosted and resolved events are present. The scale invariant tagger selects signal with high and almost flat efficiency and reject background coherently in all the resonance mass search range.

The cross section sensitivities predicted for a di-higgs resonance in a 8 TeV proton collider assuming a total luminosity of 25/fb are in the same order of magnitude of the signal sensitivities found by the ATLAS collaboration using 19.6/fb of data. This sensitivities drop by a factor half for the resolved jets regime if we assume the same experimental

capabilities in a proton-proton collider with center of mass energy of $\sqrt{s} = 14$ TeV and a total luminosity of 500/fb. We had found no considerable gain in the sensitivity for the boosted scenario. The increase on cross section with the collision center of mass energy increase is compensated by the increase of the di-jet background. Better techniques for the differentiation and b -tag of jets non-trivial substructure from new physics effects and ordinary QCD jets would change this picture.

The weak boson fusion production of a narrow resonance in the di-higgs channel in a proton collider was also investigated. Following the inclusive study we had presented the preliminary studies in the $b\bar{b}b\bar{b}$ final state. In this channel the recoil of the jets reminiscent of the weak boson fusion production mode makes $\sim 10\%$ of the parton level events to have one of the reconstructed higgs as one merged jet for all the resonance mass range allowed for the the di-higgs channel. If the mass of the resonance is larger than ~ 1 TeV both higgses become boosted (following what is was noted in the inclusive production mode). Jet substructure techniques play an important role from a resonance mass hypothesis as heavy as 350 GeV.

At the experimental level we had presented the CMS search for a resonance in the di-higgs channel in the $\gamma\gamma b\bar{b}$ final state and also commented about the searches for a di-higgs resonance in the $b\bar{b}b\bar{b}$ final state.

The search for the inclusive $X \rightarrow hh \rightarrow \gamma\gamma b\bar{b}$ was performed for a resonance mass between 260 GeV and 1.1 TeV. It was analyzed 19.7/fb of data collected from 8 TeV proton-proton collisions with a di-photon trigger. The search was based on a signal with two isolated photons and two separated jets and therefore it is expected to be efficient to look for a resonance with mass up to 1 TeV. Because of the background behavior the analysis strategy was different to search for a resonance with mass below and above 400 GeV.

The search strategy in the low mass region is similar to the search for a inclusive SM higgs in the $\gamma\gamma$ channel. The other resonant objects in the channel, namely the SM $h \rightarrow b\bar{b}$ and X resonance to hh are used to constraint background. In the high mass however the selection in the X resonance reduces efficiently the quantity of background events and the search is performed in the four bodies mass shape. The expected sensitivity of the two methods are compatible for the 400 GeV resonance mass hypothesis.

This search expect (observe) to exclude the bulk radion model for $\Lambda_R = 1$ TeV below the 900 GeV (1 TeV) mass hypothesis at 95% CL. The channel also expect (observe) to exclude at the RS1 KK-graviton between 340 GeV and 520 GeV (340 GeV and 400 GeV) for $k/\overline{M}_{Pl} = 0.1$.

Following the trigger set up the CMS search for a di-higgs resonance in the $b\bar{b}b\bar{b}$ final state is separated for resonances masses below and above ~ 1 TeV. In the search for a di-higgs resonance based in a signal with four well separated jets. The design of a multi-jet trigger with online b -tag information was a critical step to explore a resonant with mass in the region between 260 GeV and 400 GeV. The analysis strategy for the search for this resonance when the final state is characterized by four resolved jets differ consistently from what we had modelled on the naive feasibility study. The shape of QCD background makes the data driven method for signal extraction to be more sophisticated than on the fully boosted scenario.

The analysis strategy of the resonant search for a boosted di-higgs pair in the $b\bar{b}b\bar{b}$ final

state have very similar structure from what we had modelled on feasibility study: the pair of reconstructed higgses is selected while the signal is extracted from the resonance mass shape in the di-jet spectra. The substructure techniques found to be more performing for signal significance however are different from the ones used in our feasibility study. The addition of b -tag information for signal selection however constrain considerably the background. Alternative strategies for background estimation, following the low mass strategy are necessary.

When comparing the cross sections sensitivities reached with the di-higgs searches normalized to $X \rightarrow hh$ and assuming SM higgs branching fractions we find that at low masses ($m_X \sim < 400$ GeV) the di-higgs cross section sensitivities reached by the $\gamma\gamma b\bar{b}$ final state tend to be better than the sensitivities reached by the $b\bar{b}b\bar{b}$ final state due the large multi-jet background in this region. In the complementary region however ($m_X \sim > 400$ GeV) the final state with larger branching fraction ($b\bar{b}b\bar{b}$) is more limiting for the $X \rightarrow hh$ channel.

We interpret the present experimental results in the di-higgs channels in the gravity particles final states. We find the present di-higgs results to be more restrictive in the radion parameter space than the public di-boson results in the $m_X \subset [300, 600]$ GeV region. We also show how the public di-higgs analysis constrain the KK-graviton parameter space. The combinations with other KK-graviton channels would constrain even more the parameter space. However only the study of the structure of KK-graviton couplings with the weak bosons would reveal the nature of the resonance.

Appendices

A Signals of strongly coupled new physics in theories with a light higgs boson

In this appendix we consider the scenario where the Higgs doublet is realized as a bound state emerging from an underlying strong sector. A microscopic theory for a naturally light composite Higgs is not trivial to be constructed. However we draw the general idea of Higgs compositeness discussing a Topcolor theory as example.

In the 60s the lack of fundamental scalar degrees of freedom known in the nature, the success of BCS theory of superconductivity [215] and the application of Nambu-Jona-Lasinio mechanism to low energy QCD lead several authors to work on the possibility of the Higgs boson to be a top quark condensate [216] bounded by some new strong force mediated by a colored heavy vector boson with mass M . This idea was revisited in the 90s by [43] and also many other authors. When the heavy mediator is integrated out a four fermion interaction is induced the evolution of this interaction. The NP effects are parametrized by the dimension-full coupling constant of the four fermion operator $G \equiv g^2/M^2$, where g is the strength of the new force and M effectively act as the cutoff of the low energy theory. For a critical value of the coupling G an effective composite scalar field (the Higgs) is dynamically generated from the fermion loops. At low energies the mass of this scalar becomes tachyonic and the Higgs mechanism takes place.

The separation between the higgs vacuum expectation value and the mass of the mediator is controlled by the proximity of the bare coupling g is from the critical value. The compositeness conditions drives the Yukawa and quartic couplings to grow over the perturbative scale at the compositeness scale M . The top mass grows when the scale M decreases. To keep the dynamically generated higgs mass light and the compositeness scale at the TeV requires a fine-tuning between parameters is needed.

If the Higgs doublet emerges as a PGB of the global symmetry of the strong sector that incorporates the SM gauge group however its effective mass would be protected by this symmetry. When the global symmetry is explicitly broken by some underlying mechanism it allows a small mass to the scalar higgs, this mass be protected from receive large radiative corrections by the residual symmetry. Models constructed following this idea are known as *composite models*.

The microscopic theory describing the elementary sector however is unknown and since the elementary sector is strongly coupled it is non-calculable. Without the direct detection of strong resonances the low energy effects of a PGB higgs boson compositeness lie on the modifications of its couplings to matter. The allowed couplings of a PGB higgs with matter must follow the global symmetry that protects its mass (see for example [217]).

The Goldstone bosons of a global symmetry can be phenomenologically described by chiral Lagrangian [218, 219]. We consider a composite model where the strong sector obeys a global symmetry G spontaneously broken down to the coset G/H . A subgroup $P \subset G/H$ is gauged.

In this scenario the minimum ingredient to construct a Lagrangian is the matrix with the Goldstone bosons U . This object transforms as the adjoint under the unbroken group

H . If we denote $g \in G$ a general transformation, then:

$$gU(\phi) = U'(\phi' = \alpha\phi + \beta)h(\phi), \quad (\text{A.1a})$$

and the transformed matrix of Goldstones bosons can be written as:

$$U'(\phi' = \alpha\phi + \beta) = gU(\phi)h(\phi)^\dagger. \quad (\text{A.1b})$$

We denote the generators of the unbroken H group using the non-hatted indexes a and the coset generators using hatted indexes $\hat{a}$. The Goldstone fields are written as $\phi \equiv \pi^a T^a$, where the matrices T^a are the generator of infinitesimal transformations in the direction a . The U matrix reads:

$$U = e^{i\pi^a T^a}. \quad (\text{A.2})$$

A covariant ingredient $U^\dagger D_\mu U$ is constructed by means of the covariant derivative D_μ . The behavior of the covariant pieces of the chiral objects under a group transformation is:

$$U^\dagger D_\mu U \rightarrow h U^\dagger D_\mu U h^\dagger + ih \partial_\mu h^\dagger, \quad (\text{A.3})$$

We can decompose the covariant unit in the directions of the coset and the unbroken symmetry by projecting it in the direction of the generators:

$$d_\mu^a = i\text{Tr}[T^a U^\dagger D_\mu U], \quad E_\mu^a = i\text{Tr}[L^a U^\dagger D_\mu U], \quad (\text{A.4})$$

resulting that:

$$U^\dagger D_\mu U = d_\mu + E_\mu = d_\mu^a T^a + E_\mu^a L^a. \quad (\text{A.5})$$

The projected covariant components transform as:

$$d_\mu \rightarrow h d_\mu h^\dagger, \quad E_\mu \rightarrow h E_\mu h^\dagger + ih \partial_\mu h^\dagger. \quad (\text{A.6})$$

The EWSB Lagrangian can be written in terms only of the Goldstone components projected in the coset [8]

$$L_{EWSB} \equiv \frac{f^2}{4} \text{Tr}[D^\mu U^\dagger D_\mu U] = \frac{f^2}{4} \text{Tr}[d^\mu d_\mu]. \quad (\text{A.7})$$

We gauge a sub-group $P \subset G/H$ in the Lagrangian (A.7) by means of the covariant derivative.

The minimal global symmetry that would contain the EW SM gauge group and protect the mass relations between the weak bosons with global custodial symmetry is $G/H \equiv SO(5) \times U(1)/SO(4) \times U(1)_X$. The abelian symmetry $U(1)_X$ is necessary to set the correct fermions yipercharge [220].

The Goldstone bosons of the broken directions of the $SO(5) \times U(1)/SO(4) \times U(1)_X$ symmetry have the exact structure of a Higgs doublet. The couplings of the higgs boson (the Higgs neutral scalar) follows from an expansion around the vacuum $h(x) = 0$ of the

effective Lagrangian [221].

$$\frac{f^2}{4} \sin^2 \left(\theta + \frac{h(x)}{f} \right) \text{Tr}(d_\mu d^\mu), \quad (\text{A.8})$$

where $v = f \sin \theta$ and the Goldstone field is rescaled as $\pi^a \rightarrow \pi^a/v$.

This Lagrangian has a discrete symmetry under the parity transformation $h \rightarrow -h$ and $\pi \rightarrow -\pi$, although this is not obvious on this representation (see for example [221]).

The gauge group is taken as $P \equiv SO(4) \times U(1)_X = SU(2)_L \times SU(2)_R \times U(1)_X$. The covariant derivative is written as:

$$D_\mu = \partial_\mu - i(gW_{L,\mu} + g_1W_{R,\mu} + XB_\mu) \quad (\text{A.9})$$

The isospin triplet ‘‘pion’’ fields π^i ($i = 1, 2, 3$) will be identified with the longitudinally polarized weak bosons through the equivalence theorem [222]. The θ angle is determined from the weak bosons masses. At very high energies however one can use the Goldstone bosons to describe the interactions between the longitudinally polarized gauge bosons and self-couplings [223].

If we expand the Lagrangian (A.8) in powers of $h(x)/f$ the interactions of the higgs with the Goldstone bosons will be written as:

$$\mathcal{L}_{\text{eff}} = \frac{v^2}{4} \left(1 + 2a\frac{h}{v} + b\frac{h^2}{v^2} + b_3\frac{h^3}{v^3} + \dots \right) \text{Tr}[d_\mu d^\mu] + \dots, \quad (\text{A.10a})$$

We also introduce a parametrization to the higgs potential:

$$\frac{1}{2}(\partial_\mu h)^2 - \frac{1}{2}m_h^2 h^2 - d_3 \lambda v h^3 - d_4 \frac{\lambda}{4} h^4 + \dots. \quad (\text{A.10b})$$

if we define $\xi \equiv v/f$ the terms in the (A.10a) expansion are:

$$a = \sqrt{1 - \xi}, \quad b = 1 - 2\xi, \quad b_3 = -\frac{4}{3}\xi\sqrt{1 - \xi}, \quad (\text{A.11})$$

The format of the coefficients in the higgs potential will depend on the underlying theory. The SM is recovered when $a = b = d_3 = d_4 = 1$ and $b_3 = 0$. This parametrisation however is expect to be common to a large class of models, such as composite Higgs models and has been used to study anomalous Higgs couplings effects at the LHC [224].

In subsection (A.2) we comment the high energy behavior of $2 \rightarrow n$ processes involving only the Goldstone bosons. In the subsequent section we study the effect of anomalous couplings of Lagrangian (A.10) the partial unitarization of scattering processes. We then study non-unitarizing impact in partonic cross sections and also in the WBF production of three and four longitudinally polarized weak bosons or higgses at LHC. In section (A.4) we quickly comment about the LHC sensitivity to the presented anomalous couplings looking at the topology of the WBF di-higgs production at LHC.

A.1 About perturbative unitarity in a higgs-less scenario

The perturbative unitarity bound in the inelastic $2 \rightarrow n$ process assuming s -wave dominance for a given center-of-mass energy $\sqrt{s}$ is [225, 226]:

$$\sigma(2 \rightarrow n) < \frac{4\pi}{s}. \quad (\text{A.12})$$

The relativistic n -body phase space is proportional to s^{n-2} therefore the unitarity bound requires that the amplitude grows with energy no faster than:

$$A(2 \rightarrow n) \sim s^{1-n/2}. \quad (\text{A.13})$$

To assume the EWSB sector is described by a nonlinear sigma model is equivalent to assume the EWSB sector to be described by an equation like (A.7) where the d_μ contains only the $SU(2)_L$ Goldstone bosons and identify $f = v = 246$ GeV as the usual scale of electroweak. symmetry. If we assume the possible strong resonance states to be very heavy than by power-counting the scattering amplitude between the Goldstone bosons grows with energy as:

$$A_d(2 \rightarrow n) \sim \frac{s}{v^n}, \quad (\text{A.14a})$$

and hence naively:

$$\sigma(2 \rightarrow n) \sim \frac{1}{s} \left(\frac{s}{v^n} \right)^2 s^{n-2}. \quad (\text{A.14b})$$

Assuming that enough energy is available the growth of the cross section with energy in this model is faster for larger number of particles due to the kinematical factors in the phase space. Therefore there must be stronger cancellations in the scattering amplitudes due to new physics as the number of final state particles is increased.

For instance, unitarity requires that $A(2 \rightarrow 2) \sim \text{constant}$ and $A(2 \rightarrow 4) \sim 1/s$, whereas they both grow as $\sim s$ in the NL σ M. Given the fully relativistic n -body phase space given by [227]:

$$R_n(s) = \int \prod_{i=1}^n \frac{d^3 p_i}{(2\pi)^3 (2E_i)^3} (2\pi)^4 \delta^4(\sqrt{s} - \sum_{i=1}^n p_i) = \frac{(2\pi)^{4-3n} (\pi/2)^{n-1}}{(n-1)!(n-2)!} s^{n-2}, \quad (\text{A.15})$$

one can easily estimate the energy scale Λ_n at which perturbative unitarity is violated in $2 \rightarrow n$ processes in the non linear sigma model:

$$\Lambda_n = \left[\frac{2(n-1)!(n-2)!}{(2\pi)^{3-3n} (\pi/2)^{n-1}} \right]^{\frac{1}{2n}} v. \quad (\text{A.16})$$

Considering the fully relativistic phase space unitarity is violated in $2 \rightarrow 4$ processes at an energy which is almost 2.5 higher than that for the usual $2 \rightarrow 2$ processes.

A.2 Partial unitarization of $2 \rightarrow n$ ($n > 2$) processes

In case of the deviations of the couplings of the higgs sector with the gauge sector the bosons scattering, for example the ones as described by the Lagrangian (A.10), the scattering amplitudes among the elements of the theory do not fully unitarize [87, 222]. Therefore one could suspect that cross sections for $2 \rightarrow n$ processes among longitudinally polarized weak bosons could cross the unitarity bound before the $2 \rightarrow 2$ processes [228]. It is the purpose of this work to examine this impact in inelastic processes with 3 and 4 particles (either higgs bosons or V_L 's) in the final state.

In the familiar case of $2 \rightarrow 2$ amplitudes the only kinematical dependence is on the Mandelstam variables s and t . For instance, denoting the Goldstone bosons by their electric charge, the $00 \rightarrow +-$ amplitude, arising from only 4 diagrams, is given by:

$$\mathcal{M}_{00;+-} = \frac{s[(1-a^2)s - m_h^2]}{v^2(s - m_h^2)} \xrightarrow{s \gg m_h^2} (1-a^2) \frac{s}{v^2} \quad (\text{A.17})$$

There is a violation of unitarity even with the presence of the higgs boson if its coupling is not SM-like, *i.e.*, $a \neq 1$. In the SM case one obtains a constant amplitude at high energies, as expected.

In order to study the $2 \rightarrow 4$ scattering, one must expand each field U to order $\mathcal{O}(\pi^6)$:

$$\begin{aligned} \frac{v^2}{4} \text{Tr} [\partial_\mu U \partial^\mu U^\dagger] &= \frac{1}{2} (\partial^\mu \vec{\pi} \cdot \partial_\mu \vec{\pi}) \\ &+ \frac{1}{6v^2} \left(1 - \frac{2}{15v^2} \vec{\pi} \cdot \vec{\pi} \right) [(\vec{\pi} \cdot \partial_\mu \vec{\pi})^2 - (\vec{\pi} \cdot \vec{\pi}) (\partial^\mu \vec{\pi} \cdot \partial_\mu \vec{\pi})] + \mathcal{O}(\vec{\pi}^8) . \end{aligned} \quad (\text{A.18})$$

The number of Feynman diagrams to describe a $2 \rightarrow n$ scattering increases considerably with the number of final state particles making impractical to perform an analytical computation. We have implemented the Lagrangian given in equation (A.10) in **FormCalc** and **MadGraph** (using **FeynRules** UFO output format [229] to account the higher dimensional operators). We have also implemented this model in **CalcHep** using **LanHep** package[230]*.

The $2 \rightarrow 4$ amplitudes contain of the order of 100 diagrams involved that depends on several combinations of the scalar products of the different 4-momenta involved. However, some of their properties can be demonstrated with the following simple example for a given point in phase space, where all the particles lie in the same plane with an angle of $\pi/3$ between the nearest neighbors (we will keep the "pions" massless at the amplitude level since their masses are not relevant for issues of unitarity), in which case we obtain:

$$\begin{aligned} \mathcal{M}_{00;00+-} &\propto \frac{1}{v^4} [72s (13a^4 - a^2(7b+5) - 1) + 3m_h^2 (1580a^4 - 378a^3d_3 - 3a^2(245b+131) - 74) \\ &+ \frac{m_h^4}{s} (9774a^4 - 3087a^3d_3 - a^2(4494b+1289) + 52) + \dots], \end{aligned} \quad (\text{A.19})$$

The scattering amplitude grows with s , as expected. In the SM ($a = b = d_3 = 1$) one obtains in the limit $s \gg m_h^2$:

$$\mathcal{M}_{00;00+-} \propto \frac{1}{s} \frac{m_h^4}{v^4} \quad (\text{A.20})$$

*We thank Hong-Jian He for the insight of use auxiliary fields to be able to model higher dimensional interactions on **CalcHep** format.

The triple Higgs anomalous coupling parameterised by d_3 does not enter in the dominant contribution. In the following we will take $d_3 = 1$ for this channel. There is no contribution from the couplings d_4 and b_3 for the above processes. This result for the amplitude depends on the phase space configuration. The polynomials will be different but the features described above also happens at other phase space configurations and other channels.

For $2 \rightarrow 3$ processes a similar analysis can be performed. For a given configuration in phase space, where all particles lie in the same plane we find:

$$\mathcal{M}_{00;hhh} \propto \frac{1}{4v^3} [s(4ab - 4a^3 - 3b_3)] - m_h^2(3b_3 + 8ab - 8a^3) + \frac{4m_h^4}{s}(a^3 + ab - 6b_3 - 3a^2d_3) + \dots], \quad (\text{A.21})$$

for the angles between the nearest higgs bosons fixed at $2\pi/3$ and:

$$\mathcal{M}_{00;+-h} \propto \frac{a}{192v^3} \left[s(-1 + 2a^2 - b) + \frac{m_h^2}{4}(-164 + 386a^2 - 213b - 9ad_3) - \frac{3m_h^4}{2s}(-262 + 291a^2 - 93b + 81ad_3) + \dots \right], \quad (\text{A.22})$$

for the fixed configuration of two collinear pions, back-to-back with the Higgs boson and at right angles with the incident particles. The SM the first two terms in these amplitudes vanish, as it should. Notice also that the $\mathcal{M}_{00;hhh}$ amplitude is sensitive to b_3 , being the lowest multiplicity process in which this happens.

From the parity of the MCHM4 class of theories under which $\pi \rightarrow -\pi$ and $h \rightarrow -h$ the polynomial with largest growth in the $2 \rightarrow 3$ processes also vanish for the values of a , b and b_3 that obey the MCHM4 relations as in Eq.(A.11).

In summary, for $a \neq 1$ and $b \neq 1$ the squared amplitude grows as s^2 instead of decreasing as $1/s^2$. Therefore there is a sensitivity of $2 \rightarrow 3$ and $2 \rightarrow 4$ processes to non-SM higgs couplings compared to $2 \rightarrow 2$ processes whereas the SM amplitude goes to a constant for large s . For the $2 \rightarrow 3$ processes there is also a suppression if the values for the parameters predicted by the MCHM4 are used due to the symmetry of the coset. To quantify this sensitivity we study directly the cross section for these processes.

A.3 Cross section sensitivity

In this section we analyze the cross section for the $2 \rightarrow 3, 4$ processes at the parton level for couplings using the Lagrangian (A.10a) implemented in `CalcHep`. We had studied several channels but will report only on the most representative ones.

We show examples of the enhancements that result from the anomalous Higgs couplings we compute cross sections with a higgs mass $m_h = 125$ GeV. We implement a cut in the invariant mass of the final state pions $m_{+-} > 200$ GeV such that the higgs boson is never on-mass-shell when coupled to two pions, as is it the case in the SM. The left part of figure (A.1) shows the ratio of the cross section in several channels as a function of a (keeping the other parameters fixed in the SM case to the SM cross section at a fixed center-of-mass energy of $\sqrt{s} = 1$ TeV. The right part of this figure shows a similar plot in the MCHM4 case (changing the parameters according to equation A.11 and assuming $d_3 = \sqrt{1 - \xi}$). The notation $(00, + -)$ indicates that both 00 and $+ -$ initial states have been taken into account in the cross section.

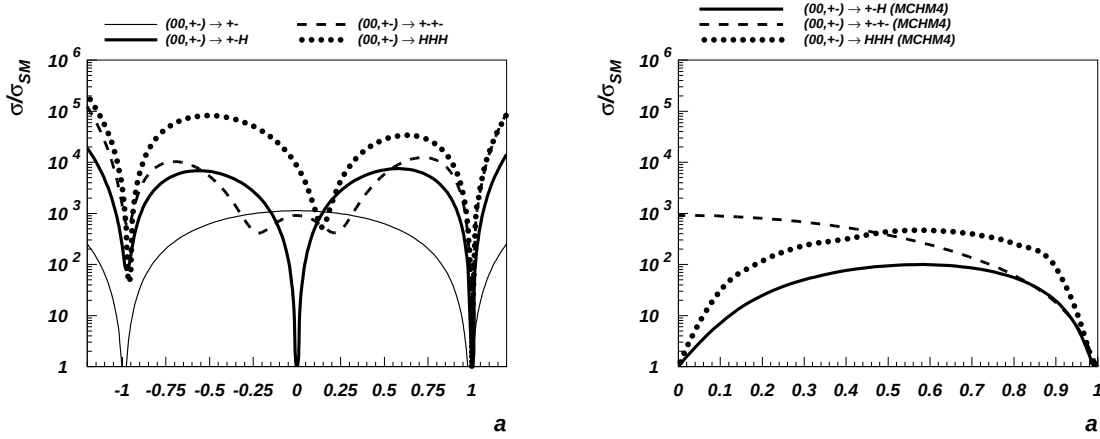


Figure A.1: Ratio of cross sections as a function of a parameter to the SM ones for two cases: the general effective Lagrangian with all parameters but a fixed at their SM values (left plot) and for MCHM4 model (right plot) at a fixed energy of $\sqrt{s} = 1$ TeV. The different channels are: $(00, + -) \rightarrow + - + -$ (dashed line), $(00, + -) \rightarrow + - h$ (thick solid line), $(00, + -) \rightarrow h h h$ (dotted line), and $(00, + -) \rightarrow + -$ (thin solid line) for comparison. The notation $(00, + -)$ indicates that both 00 and $+ -$ initial states were taken into account. Not all the values for the anomalous couplings shown in this plots are allowed by the present collider constraints [231] this figure is only meant for illustrative purposes.

Large enhancements of the order of 10^3 to 10^5 with respect to the SM value are obtained even for deviations as small as 10% of the couplings from their SM values. The largest case occurs for triple Higgs production. The 4, 3 and 2 dips in the cross section versus a for $2 \rightarrow 4$, $2 \rightarrow 3$ and $2 \rightarrow 2$ are due to the 4^{th} , 3^{rd} and 2^{nd} order polynomials in a in the amplitudes. The enhancements in $2 \rightarrow 2$ processes are modest compared to processes with higher multiplicities, at least at $\sqrt{s} = 1$ TeV. As expected from the parity symmetry of the coset the enhancements in the cross section $2 \rightarrow 3$ process is suppressed in the MCHM4 case with respect to the free variation of parameters.

We study the growth with center-of-mass energy of the cross section for different multiplicities for a few values of the anomalous coupling, namely $a = 0.9, 0.95$ and 1 (SM), keeping the other couplings at their SM values. We expect the same order-of-magnitude enhancements if the other couplings are also anomalous (but without obeying the MCHM4 relations).

In figure (A.2) we show a comparison of the cross section as a function of energy among representative processes with 2, 3 and 4 particles in the final state, for different values of the anomalous coupling. We also draw the unitarity limit (equation A.12).

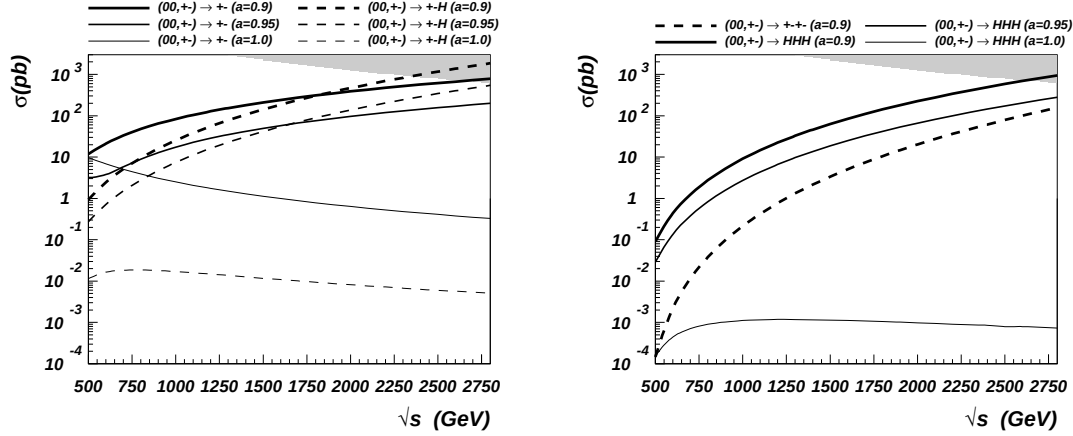


Figure A.2: Comparison among cross sections as a function of the center-of-mass energy for processes with 2, 3 and 4 particles in the final state. In the top plot, the solid lines are for $(00, +-) \rightarrow +-$ for $a = 0.9$ (thick), $a = 0.95$ (medium thick) and $a = 1$ (thin). Dashed lines are for $(00, +-) \rightarrow +-h$, with the same pattern for the thickness of the lines. In the bottom plot, the same pattern of lines show the results $(00, +-) \rightarrow hhh$ and the process $00 \rightarrow +-+-$ is shown as a dashed line for $a = 0.9$. In these plots only a deviates from the SM value. The unitarity bound is shown as a shaded area in the top right corner.

The SM cross section quickly stabilizes at a small value which depends on the specific process (of the order of 10^{-3} pb, 10^{-2} pb and 10^{-1} pb for $00 \rightarrow hhh$, $(00, +-) \rightarrow +-h$ and $(00, +-) \rightarrow +-$, respectively). The non-SM case the cross sections grow very fast with energy, reaching up to order of 100 pb and violating unitarity at center-of mass energies of the order of a few TeV.

Larger multiplicity processes, in the absence of a complete cancellation mechanism, grow faster due to phase space. In the examples shown the $2 \rightarrow 3$ start to overcome $2 \rightarrow 2$ at energies of $\mathcal{O}(1\text{TeV})$. The $2 \rightarrow 4$ process grows very rapidly for $a \neq 1$, but since it starts out very suppressed it surpasses the $2 \rightarrow 2$ only at very high energies, of the order of $\mathcal{O}(5\text{TeV})$.

So far we have only analyzed the scattering of the longitudinally polarized gauge bosons. Since it is difficult to separate out the contributions from these polarizations in an experimental setting, it is important to understand how the large enhancements found will affect the corresponding un-polarized cross section.

To study the complete weak boson scattering we use the Lagrangian (A.10) promoting the partial derivatives to full covariant derivatives and adopt the unitary gauge ($U = 1$). As an illustration we fix the partonic center-of-mass energy at 2 TeV, which is still below the unitarity bound since there is a partial unitarization due to the presence of a Higgs boson, and compare the cross sections for the longitudinally polarized gauge bosons with the full SM in table (A.1). We keep the notation 0, +, − to indicate longitudinally polarized gauge bosons and $Z, W^\pm$ to denote the un-polarized gauge bosons.

channel	$a = b = 1$ (SM)	$a = 0.9; b = 1$
$00 \rightarrow + -$	0.53 (0.13)	66.4 (295)
$ZZ \rightarrow W^+ W^-$	629 (610)	646 (655)
$00 \rightarrow + - h$	4.6×10^{-3} (2.0×10^{-3})	18.7 (350)
$ZZ \rightarrow W^+ W^- h$	5.49 (10.9)	6.17 (46.2)
$00 \rightarrow hh$	0.64 (0.18)	43.0 (158)
$ZZ \rightarrow hh$	7.18 (7.61)	4.31 (15.7)
$00 \rightarrow hhh$	5.6×10^{-4} (4.9×10^{-4})	4.5 (112)
$ZZ \rightarrow hhh$	1.7×10^{-2} (4.7×10^{-2})	0.61 (13.6)

Table A.1: Comparison of $2 \rightarrow 2$ and $2 \rightarrow 3$ cross sections (in picobarns) at $\sqrt{s} = 1$ TeV ($\sqrt{s} = 2$ TeV in parenthesis).

The processes with longitudinal polarizations are sub-dominant in the SM. However, as we discussed above, they are greatly enhanced with small deviations of the couplings and actually dominate the cross sections. Our results are consistent with the fact that $\sigma_{\text{all}} \approx \sigma_{LL}/9$ when the contribution from longitudinal polarizations is dominant.

When all polarizations are included, the contribution from transverse polarizations can mask the increase in the cross section for the longitudinally polarized gauge bosons. This can be seen in the case of $ZZ \rightarrow W^+ W^-$. For the un-polarized case the cross section for $ZZ \rightarrow W^+ W^-$ is still at least one order of magnitude larger than the typical $2 \rightarrow 3$ processes but $ZZ \rightarrow hh$ is of the same order as the $ZZ \rightarrow hhh$ cross section.

On the other hand, in cases where the contributions from the transverse polarizations are not large, as in the case of $ZZ \rightarrow hhh$, enhancements of $\mathcal{O}(10^3)$ can be obtained. The multiple higgs boson production offers the best channels to study anomalous couplings.

We estimate the impact of the enhancements found at the parton level arising from anomalous higgs couplings. We had performed a full calculation of $pp \rightarrow jj + X$, where $j = u, \bar{u}, d, \bar{d}, s, \bar{s}$ and $X = W^+ W^-, W^+ W^- h, hhh$ at the LHC (for $\sqrt{s} = 14$ and 33 TeV) using MadGraph5. We use CTEQ6L1 parton density function for the evaluation of the tree-level cross sections with the QCD scale equal to M_Z . The selection and acceptance cuts include the requirement of two jets with $P_{T_j} > 30$ GeV with $|\eta_j| < 5$ separated with $\Delta R = \sqrt{\Delta\phi_{jj}^2 + \Delta\eta_{jj}^2} > 0.4$. Additional cuts to select jets from the WBF process are also tested: each jet to be quite energetic with $E_j > 300$ GeV as well as a large rapidity gap between the two jets $|\Delta\eta_{jj}| > 4$ (see e.g. [232] for detailed motivation of this choice). In table (A.2) we present the results with and without the WBF cut.

When there are weak bosons in the final state the cross section actually decreases for most cases with $a = 1$ versus $a = 0.9$ ones. This is because we chose in our example

Process	14 TeV		33 TeV	
	with (without) WBF cuts		with (without) WBF cuts	
	a=1.0 b=1.0	a=0.9 b=1.0	a=1.0 b=1.0	a=0.9 b=1.0
$pp \rightarrow jjW^+W^-$	95.2 (1820)	99.3 (1700)	512 (5120)	540 (5790)
$pp \rightarrow jjW^+W^-h$	0.011 (0.206)	0.0088 (0.172)	0.0765 (0.914)	0.0626 (0.758)
$pp \rightarrow jjh hh$	1.16×10^{-4} (3.01×10^{-4})	0.0566 (0.0613)	0.00115 (0.00165)	1.85 (1.46)

Table A.2: Cross section (in fb) for $pp \rightarrow jjW^+W^-$, $pp \rightarrow jjW^+W^-h$ and $pp \rightarrow jjh hh$ with WBF cuts. The number in parenthesis is the cross section without WBF cuts.

$a < a_{SM} = 1$ and since the transverse polarizations dominate the cross section, reducing the coupling a results in a smaller cross section. In the case of triple higgs production, the enhancements are substantial: roughly a factor of 500 for $\sqrt{s} = 14$ TeV (LHC14) and 1600 for $\sqrt{s} = 33$ TeV (LHC33), with WBF cuts. In figure (A.3) we show the results for the triple higgs production cross section for both LHC14 and LHC33 for the large range of anomalous coupling a . The enhancements with respect to the SM case $a = 1$ are large and don't change significantly once $|\Delta a/a| > 0.1$.

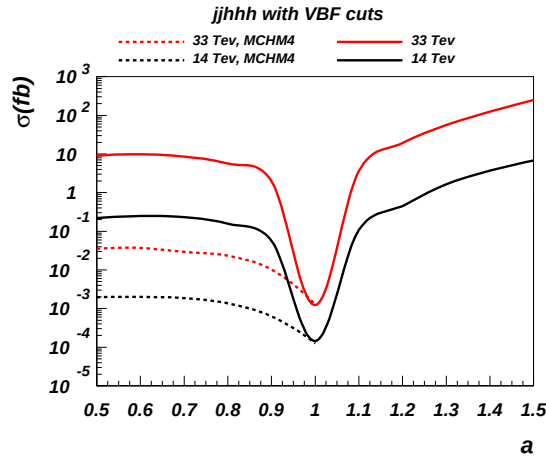


Figure A.3: Cross section for triple higgs production $pp \rightarrow jjh hh$ with WBF cuts as a function of the anomalous coupling a for LHC14 (dark lines) and LHC33 (light lines). Dashed lines are for other parameters fixed to SM values and solid lines are for parameters given by MCHM4 relations.

Though the enhancement can be as large as 10^5 for $a = 1.5$, the absolute value of the respective cross sections are 10 fb for $\sqrt{s} = 14$ TeV with WBF cuts. The study of these processes challenging at the LHC. A dedicated analysis is necessary to understand the LHC or LHC33 sensitivity to the processes above.

A.4 Probing anomalous couplings at LHC in WBF di-higgs production

The kinematics of the WBF jets are sensitive to the structure of the VV initiated internal process. In this study we want to describe the couplings of the higgs scalar with the massive weak bosons in a complete calculation description of the kinematics of the WBF di-higgs production at hadron colliders. The effective operators relevant to WBF di-higgs production are the parameters a , b and d_3 of the Lagrangian (A.10), see figure (A.4).

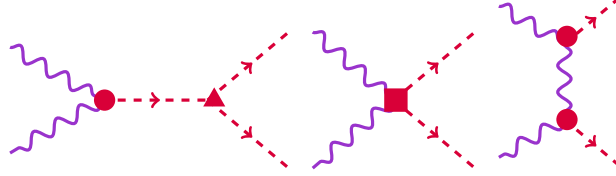


Figure A.4: Feynman diagrams for di-higgs WBF production. Up to dimension six operators. The vertex marked by the triangle receives corrections dependent of d_3 parameter, while the vertexes marked by a the box and circle are respectively corrected by b and a parameters.

The parameters that potentially modify the behavior of the WBF jets are a and b . This study was primary done by [233], this reference however did not studied the impact of substructure in the study. In a preliminary parton level study following the study presented in section II.5. We had verified that in the case a and b variates a portion of the higgses in the $pp \rightarrow hhjj$ process is boosted and therefore the use of the scale invariant tagger is an important tool in the measurement of the SM anomalous couplings in the electroweak sector. The study of this process at LHC and HL-LHC environment is in progress. To visually demonstrate we show in figure the p_T of the leading reconstructed higgs varying one of the a , b and d_3 anomalous couplings from the SM value separately.

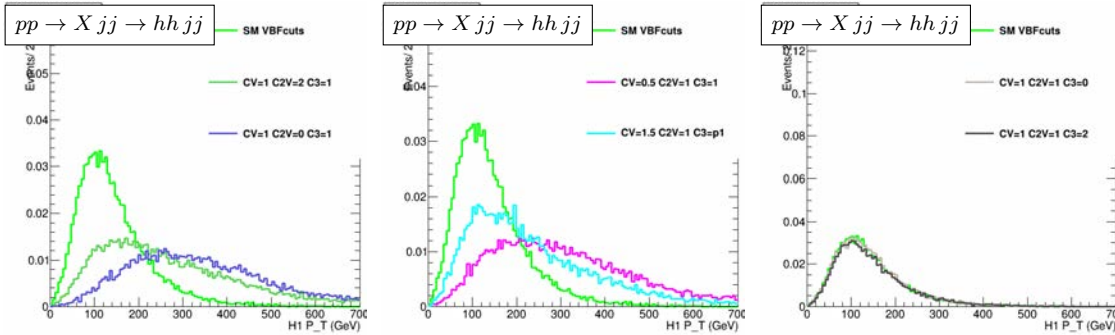


Figure A.5: Transverse momentum of the leading reconstructed higgs under variation of the anomalous couplings. The plots are shape normalized. **Left:** Variation of a (c_V). **Center:** Variation of b (c_{2V}). **Right:** Variation of d_3 . In the case of variation of d_3 (c_3) coupling there is no variation in the process kinematics. The new physics effect of variation of the higgs triple coupling is in the signal cross sections.

B Bulk Standard Model fields

On this appendix we develop the Kaluza-Klein expansions of bulk matter fields in a more detailed way that made for the graviton field on section (I.3). Apart from details related with gauge invariance and field spin hypothesis, the methodology for KK decomposition of Graviton, gauge bosons and fermions five dimensional fields are very similar and we derive only the equations for the profiles of the gauge fields. The dedicated calculation of the four dimensional projection of a five dimensional fermion field can be find for example in [234, 235, 33]. We ignore the fields back-reaction on the metric and also possible boundary terms.

B.1 Bulk gauge bosons

We derive the wave functions for an abelian five dimensional gauge field. The extension of this formalism to non-abelian gauge groups is straightforward.

The most suitable gauge for the full derivation is the R_ξ gauge, this gauge avoids mixing between the space-time (A_μ) and the extra dimensional (A_5) components of the gauge fields on the Lagrangian, allowing a clear separation of the equations of motion*.

The gauge fixing term, on the R_ξ gauge reads:

$$S_{GF} = \int d^5x \frac{-1}{2\xi} (\partial_\mu A^\mu - \xi \partial_5 (e^{-2\sigma(\phi)} A_5))^2, \quad (\text{B.1})$$

In this gauge the full 5D action for the abelian gauge field S_A is:

$$S_A = \int d^5x \left(-\frac{1}{4} F_{\mu\nu} F^{\mu\nu} + \frac{1}{2} e^{-2\sigma(\phi)} \partial_\mu A_5 \partial^\mu A_5 + \frac{e^{-2A}}{2} (\partial_5 A_\mu)^2 - \frac{1}{2\xi} (\partial_\mu A^\mu)^2 - \frac{\xi}{2} (\partial_5 (e^{-2\sigma(\phi)} A_5))^2 \right) - \int d^4x (e^{-2\sigma(\phi)} A_5 \partial_\mu A^\mu)|_{0,\pi}. \quad (\text{B.2})$$

From the minimization of the above action we arrive to the equations of motion:

$$-\partial_\mu^2 A_5 + \xi \partial_5^2 (e^{-2\sigma(\phi)} A_5) = 0, \quad (\text{B.3a})$$

$$[\eta^{\mu\nu} \partial^2 - \eta^{\mu\nu} \partial_5 (e^{-2\sigma(\phi)} \partial_5) - (1 - \frac{1}{\xi}) \partial^\mu \partial^\nu] A_\nu = 0, \quad (\text{B.3b})$$

and the boundary terms:

$$S_{Boundaries} = \int d^4x e^{-2\sigma(\phi)} \left(-A_5 \partial_\mu A^\mu - \frac{\xi}{2} A_5 \partial_5 (e^{-2\sigma(\phi)} A_5) + \frac{1}{2} A_\mu \partial_5 A^\mu \right) |_{0,\pi}. \quad (\text{B.4})$$

To vanish all the boundary terms on each brane with 2 conditions we can choose two

*Do not confuse the gauge parameter ξ , used only on this appendix, with the parameter controlling higgs-radion mixing.

combinations, denoted by the signs:

$$\begin{array}{l|l} (+) & A_5|_{0,\pi} = \partial_5 A^\mu|_{0,\pi} = 0 \\ (-) & A^\mu|_{0,\pi} = \partial_5(e^{-2\sigma(\phi)} A_5)|_{0,\pi} = 0 \\ (0) & A_5|_{0,\pi} = A_\mu|_{0,\pi} = 0. \end{array} \quad (\text{B.5})$$

The equation (B.3) is solved performing a Kaluza-Klein (kk) expansion, separating the behaviour of each mode along the extra dimension on the - so called - profile functions $f_A^{(n)}(\phi)$ and $f_5^{(n)}(\phi)$:

$$A_\mu(x, y) = \frac{1}{\sqrt{L}} \sum_n A_\mu^{(n)}(x) f_A^{(n)}(\phi), \quad (\text{B.6a})$$

$$A_5(x, y) = \frac{1}{\sqrt{L}} \sum_n A_5^{(n)}(x) f_5^{(n)}(\phi). \quad (\text{B.6b})$$

If we substitute the above KK expansion in the five dimensional equations of motion (B.3) we find that each mode obeys:

$$\partial_\nu^2 A_n^\nu(x) - (1 - \frac{1}{\xi}) \partial^\nu (\partial^\mu A_\mu^{(n)}(x)) = -m_n^2 A_n^\nu(x) \quad (\text{B.7a})$$

$$\partial_\mu^2 A_5^{(n)} = -\xi \frac{m_n'}{2} A_5^{(n)} \quad (\text{B.7b})$$

where the arbitrary constants m_n and m_n' are the separation variables. The $A_{\nu,n}$ behave like a massive gauge field with mass m_n , while the the fifth component $A_{5,n}$ behaves like a scalar field with mass $\xi m_n'$. Their profiles follow the equations:

$$\partial_5(e^{-2\sigma(\phi)} \partial_5 f_A^{(n)}) = -m_n^2 f_A^{(n)}, \quad (\text{B.8a})$$

$$\partial_5^2(e^{-2\sigma(\phi)} f_5^{(n)}) = -m_n'^2 f_5^{(n)}. \quad (\text{B.8b})$$

The orthonormality relations reads:

$$\int_0^\pi d\phi f_A^{(n)} f_A^{(m)} = \frac{\delta_{mn}}{r_c}. \quad (\text{B.9})$$

Assuming the five dimensional gauge parameter $(\alpha(x, \phi))$ also have a KK decomposition $(\alpha(x, \phi) = \sum \alpha^{(n)}(x) f_\alpha^{(n)}(\phi))$ the gauge invariance of to each $A_\mu^{(n)}$ modes holds only if we use the boundary conditions conditions (+) or (-) of (B.5). These boundaries conditions translate on the profile functions as:

$$\begin{array}{l|l} (+) & f_5^{(n)}|_{0,\pi} = \partial_5 f_A^{(n)}|_{0,\pi} = 0, \\ (-) & f_A^{(n)}|_{0,\pi} = \partial_5(e^{-2\sigma(\phi)} f_5^{(n)})|_{0,\pi} = 0. \end{array} \quad (\text{B.10})$$

B.1.1 Massless zero modes

There is no natural value for the zero mode, we take the zero modes to be massless ($m_0 = m'_0 = 0$). In this case equations (B.3) can be explicitly solved, the profiles solutions reads:

$$f_5^{(0)} = e^{2\sigma(\phi)}(c\phi + d) \quad (\text{B.11a})$$

$$f_A^{(0)} = ae^{2\sigma(\phi)} + b \quad (\text{B.11b})$$

If we chose the $(++)$ boundaries conditions the fifth gauge field component to vanish while the $A_{\mu,n}$ components holds a constant non-zero value, like the SM field pattern:

$$f_5^{(0)}(\phi) = 0, \quad (\text{B.12a})$$

$$f_A^{(0)}(\phi) = \frac{1}{\sqrt{\pi r_c}}. \quad (\text{B.12b})$$

C Effects of higgs-radion mixing in LHC production

In this section we investigate the effects of higgs-radion mixing on the scalar LHC production cross section. Reference [236] exploit the amount of higgs-radion mixing allowed by LHC8 data, they consider both the precision measurements of the 125 GeV boson and exclusion limits on the other scalar. They reach the conclusion that if we consider the 125 GeV scalar boson as the lightest scalar (h_ξ), the parameter $\Lambda_R = 3$ TeV and also the exclusion limits on the mass of the radion m_R the reasonable values for mixing are already very constrained to $\xi < 0.1$ at 98% CL when the radion mass varies from 125 GeV (degenerated case) to 600 GeV.

We follow [73] and define the parameters a, b, c and d relating the gauge and mass scalar eigenstates like:

$$h = d h_\xi + c r_\xi, \quad (\text{C.1a})$$

$$r = a r_\xi + b h_\xi. \quad (\text{C.1b})$$

The explicit format of the parameters with the fundamental parameters of the theory can be seen on equation [73] however these expressions are not important to our discussion. The modification the couplings of radion and higgs scalars to weak bosons and quarks are proportional to the factor:

$$\text{radion : } (c + \gamma a), \quad (\text{C.2a})$$

$$\text{higgs : } (d + \gamma b), \quad (\text{C.2b})$$

where $\gamma \equiv \frac{v_0}{\Lambda_R}$ and the value of v_0 is defined such as the mass of the higgs gauge eigenstate is $m_h^2 = 2\lambda v_0^2$.

The Λ_R scale is supposed to be at TeV scale, while for small mixing we expect the value of $v_0 \sim v \sim 100$ GeV. For small mixing the value of γ is expected to be small, suppressing the dependences on a and b on the couplings. For a near zero mixing the d and b values to approach zero, resulting the biggest mixing effects on the non-gluon fusion couplings to be on the radion production.

In the figure (C.1) we show the leading order effect of mixing for the non-gluon fusion production modes. The effects of mixing can reach the 50% level for a low mass radion (~ 300 GeV), not modifying too much the cross section production of a 125 GeV SM higgs by these modes.

The modification of radion (higgs) couplings to massless gauge bosons are proportional to the loop function and point-like interaction (generically denoted by *loop* and *point*) as:

$$\text{radion : } (c + \gamma a) \times \text{loop} + \gamma a \times \text{point}, \quad (\text{C.3a})$$

$$\text{higgs : } (d + \gamma b) \times \text{loop} + \gamma b \times \text{point}. \quad (\text{C.3b})$$

Figure C.1: Effects of mixing on leading order production on non-gluon fusion production mode. The red (green) curves represent a positive (negative) mixing parameter. The continuous (dashed) curve represents $|\xi| = 0.1$ ($|\xi| = 0.01$). The black curve stands for non mixing scenario ($\xi = 0$), for reference. **Right:** radion. **Left:** higgs.

In the figure (C.1) we show the effects of two specific values for the mixing parameter $|\xi| = 0.1$ and $|\xi| = 0.01$ on the ratio between the gluon fusion production cross sections given the ξ parameter with the same cross section on non-mixing scenario.

Figure C.2: Effects of mixing on leading order production on gluon fusion production mode. The red (green) curves represent a positive (negative) mixing parameter. The continuous (dashed) curve represents $|\xi| = 0.1$ ($|\xi| = 0.01$). The black curve stands for non mixing scenario ($\xi = 0$), for reference. **Right:** radion. **Left:** higgs.

The impact of mixing on higgs gluon fusion production is bigger than the impact on radion gluon fusion production for all mass range. What is a clear effect of the dominance of the point-like interaction upon the loop function on scalar-gluon interactions. For phenomenological reasonable range for the mixing parameter ξ the maximum mixing effect we find on gluon fusion radion production is 8%.

D Counting experiment

For each candidate resonance mass, m_X , we compute the number of background events in a mass window of width $f_m = 15\%$ around M_X . The local p -value for each mass point M based on the expected number of signal and background events, N_s and N_b respectively, in the mass window considered if $N_b \ll 1$ and $N_s \ll 1$ is given by (see for example [193]):

$$p = \frac{1}{2} \left(1 - \text{Erf} \left[\frac{N_s}{\sqrt{2N_b}} \right] \right), \quad (\text{D.1})$$

where Erf is the error function, and one assumes that the number of background events in each mass bin is a Poisson distribution with mean N_b . If $N_b \sim < 1$ and $N_s \sim < 1$ one has to use the discrete Poisson formula for the p -value:

$$p(M) = 1 - \frac{\Gamma(N_s + N_b, N_b)}{\Gamma(N_s + N_b)}, \quad (\text{D.2})$$

which involves the incomplete Gamma function.

To require $p = 0.05$ determines the number of signal events N_s that would allow an exclusion of the background-only hypothesis at the 95% confidence level, namely:

$$N_s = \sqrt{2N_b} \cdot \text{Erf}^{-1}(1 - 2 \cdot 0.05). \quad (\text{D.3})$$

We can obtain the model independent bound on the combination $\sigma(pp \rightarrow X)\text{BR}(X \rightarrow hh)$ by correcting the number of events for the signal tagging efficiency, with and without b -tagging, the $h \rightarrow b\bar{b}$ branching fraction and the assumed total integrated luminosity L , the total excluded cross section reads:

$$[\sigma(pp \rightarrow X)\text{BR}(X \rightarrow hh)]_{\text{excl. 95\% CL}}(M) = \frac{N_s}{(\text{BR}(h \rightarrow b\bar{b}))^2 \text{SignalEff}(M)L}, \quad (\text{D.4})$$

where the signal efficiency $\text{SignalEff}(M)$ is derived from the **MadGraph** radion and graviton samples.

E Classification for the limits

In this appendix we study the best categorization scheme, the comparison target are the statistical-only expected limits for $M_{\gamma\gamma}$ method. The main classification variable for the events is the number of b -tags. We define two b -tag with the presence of at least two b -tagged jets (high purity) or exactly 1 b -tagged jet (low purity) – we built working points requiring one or two jets to pass a given b -tag selection. The b -tag WPs we can construct with listed in table E.1.

	2 b -tag category	1 b -tag category
WP1	loose-loose	loose
WP2	medium-loose	loose
WP3	tight-loose	loose
WP4	medium-medium	medium
WP5	tight-medium	medium
WP6	tight-tight	tight

Table E.1: The six event working points built combining different b -tag requests on the two jets.

On top of b -tag classification an additional possible events classification is based on the goodness of the photons. The shower shape variable $R9$ is strongly related to the probability that the photon has converted, so we considered splitting the events having at least one *low* $R9$ photon (high probability for a conversion) from those with both *high* $R9$ photons, this classification is inherited from $H\gamma\gamma$ cut based analysis [187].

Taking into account both $R9$ and b -tag could result on the construction of four categories, however the statistics on the 2 b -tag category is very limited resulting on the definition of three categories (see table E.2). We compare these categorization definitions with the performance of use only two b -tag based categories.

	cat0	cat1	cat2	cat3
2 categories	2 btag	1btag	-	-
3 categories	2 btag	1btag high R9	1 btag low R9	- -
4 categories	2 btag high R9	2 btag low R9	1btag high R9	1 btag low R9

Table E.2: Definition of the categories under analysis.

In figure E.1 we show the statistical only expected limits for the described classifications for the 300 GeV and 500 GeV mass hypothesis. At 500 GeV hypothesis the data events yield at higher masses is too low to allow too many classes.

The separation in 2 b -tag categories improves the limits with respect to one b -tag category only. This is related to the fact that the limits are dominated by the high purity category, while the low purity one mainly allows to save a fraction of the signal in case it is very small. The classification in photons categories brings a marginal improvement.

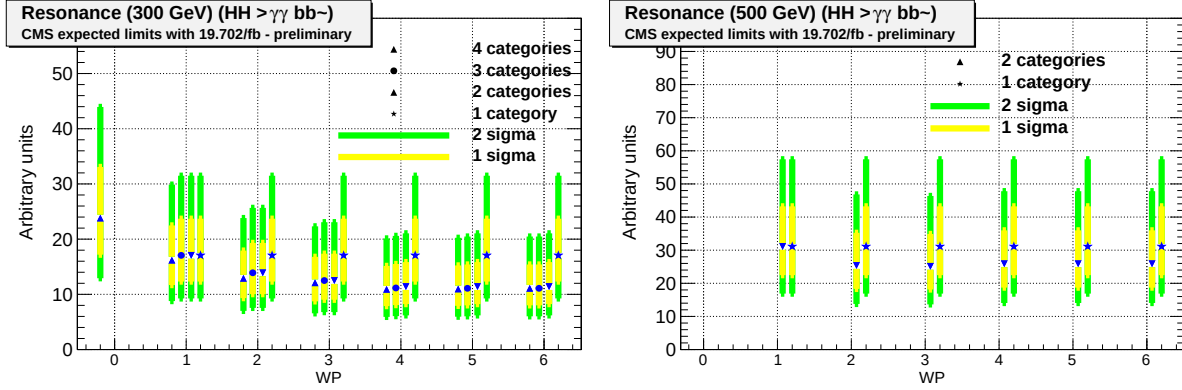


Figure E.1: Preliminary limits as a function of the CSVM b -tagging and classification. WP0 means photon-only categories, for the other WPs see table E.1. The selection and the samples are not the final ones: non optimized dijet mass window ($98 < m_{jj} < 151$ GeV) and no kinematic fit are applied. A third order Bernstein polynomial is used as a background model. Only the statistical uncertainties are considered. **disclaimer: The figure is not obtained with the exact final selection and samples, but the conclusions are not expected to change.**

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