

Investigation of mass reconstruction techniques for resonances in the scattering of $W^\pm W^\pm \rightarrow W^\pm W^\pm$ at the LHC

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Abstract

Physics scenarios of electroweak symmetry breaking beyond the Standard Model introduce new resonances in the scattering of massive, weak gauge bosons. Due to the best signal to background ratio, the like-sign $W^\pm W^\pm jj$ channel is the most favourable final state for a first glance at resonances in vector boson scattering (VBS) at a hadron collider such as the LHC. Resonances in this scattering process can be modelled with the approach of an effective field theory and the K-matrix unitarisation method.

This thesis presents a study of mass reconstruction of resonances in the fully leptonic decay channel of like-sign $W^\pm W^\pm$ scattering. Special emphasis lies on the technique of constrained minimisation leading to mass bound variables. For different resonance types, variables providing the best discovery potential in proton-proton collisions at a centre-of-mass energy of 8 TeV and an integrated luminosity of 20 fb^{-1} are determined and characterised.

Kurzdarstellung

Durch Szenarien neuer Physik im Bereich der elektroschwachen Symmetriebrechung jenseits des Standard Models werden neue Resonanzen in der Streuung von massiven, schwachen Vektorbosonen eingeführt. Aufgrund des besten Verhältnisses von Signal zu Untergrund an Hadronbeschleunigern wie dem LHC, erweist sich der Kanal der Streuung von gleichgeladenen $W^\pm$ Bosonen im Endzustand $W^\pm W^\pm jj$ als günstig, um erste Studien zu Resonanzen in Vektorbosonstreuung (VBS) durchzuführen. Resonanzen in diesem Streuprozess können mittels einer effektiven Feldtheorie und der K-Matrix-Unitarisierungsmethode simuliert werden.

In dieser Diplomarbeit wird die Rekonstruktion der Massen von Resonanzen im voll-leptonischen Zerfallskanal der gleichgeladenenen $W^\pm W^\pm$ -Streuung untersucht. Der Schwerpunkt liegt dabei auf der Massenrekonstruktion mittels der Technik der zwangsbedingten Minimierung, welche zu Variablen mit Schrankencharakter führt. Für ausgewählte Spezies von Resonanzen werden jene Massenrekonstruktionsvariablen ermittelt und charakterisiert, welche das beste Entdeckungspotential in Proton-Proton-Kollisionen bei einer Schwerpunktsenergie von 8 TeV und einer integrierten Luminosität von 20 fb^{-1} aufzeigen.

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1 Introduction

In the last half century, the field of particle physics has developed immensely. From the early theoretical concept of electroweak interactions [1, 2], a theory evolved which is able to describe the whole spectrum of observed particles with their properties and interactions. Various precision measurements [3] and the recent discovery of the Higgs boson, the last missing particle in this framework, at the Large Hadron Collider (LHC) [4, 5] established this theory as the fundamental theory of particle physics. It is called the Standard Model (SM).

Nonetheless, the SM is not able to give answers to some striking observations such as dark matter [6], the baryon asymmetry [7], the hierarchy problem [8] or the origin of the masses of the elementary fermions, not to mention the inability to incorporate gravity in the SM. Some of these questions are inevitably related to electroweak symmetry breaking (EWSB), the mechanism which is responsible for generating the masses of the fundamental particles and for which the Higgs boson is direct evidence. Thus, it is “now widely believed that [the EWSB as assumed in the SM] cannot be the whole answer.” [9]. Various new approaches draw upon more elaborated structures of Higgs fields such as in supersymmetry or composite particle theories for which evidence could evolve in the form of new resonances at reasonably high energies above 1 TeV. However, restrictions on the experimental energy range accessible to date inhibit the exploration of these theories. Thus, these new resonances need to be embedded in an effective theory which parametrises model-independently the unidentified underlying physics of EWSB at these high energies and converges to the SM where its validity has been well established. Therefore, new, hypothetical degrees of freedom that are allowed by the laws of quantum number conservation are added to this effective approach. These degrees of freedom can then be related to new particles in the different explicit models of EWSB beyond the SM.

The scattering of massive, weak vector bosons offers an excellent starting point to explore new resonances in this field as it is directly related to EWSB. These scattering processes are barely investigated in the context of the SM. With the measurement of electroweak production of like-charge $W^\pm W^\pm$ in association with two jets at the ATLAS detector, a first evidence of the existence of such scattering processes was recently published [10]. Due to its superior background rejection at hadron colliders such as the LHC, this channel is chosen to serve as the pioneer process to investigate resonances in the scope of EWSB.

Aiming at detecting and investigating new resonances, their discovery potential needs to be evaluated in a set of physics observables which are most sensitive to the influence of these resonances. The study of an appropriate measure of the resonance’s mass, such as the invariant composite mass of all final state decay products of the resonance, offers a constructive starting point. At hadron colliders, leptonic final states constitute the cleanest detector signatures. Translating this into the channel of

like-charge $W^\pm W^\pm$ scattering, fully leptonic decays of the $W^\pm$ bosons into a charged lepton and an (anti-)neutrino each are the most favourable topologies to investigate. This, however, bears the problem of two (anti-)neutrinos in the final state which are undetectable by detectors. Solely the sum of the energy carried away by both neutrinos in the transverse plane can be detected. As a consequence, the full invariant mass of the resonance cannot be reconstructed in these decay channels. Therefore, strategies need to be found which allow for an identification of the resonance and, in a later stage, a most precise mass measurement despite the lack of important information.

The objective of this thesis is to investigate mass reconstruction variables in the context of resonances in like-charge $W^\pm W^\pm$ scattering. Numerous mass reconstruction variables have been designed in recent years for the analysis of complex decay topologies at particle physics experiments, especially in terms of the search for the Higgs boson [11–13] and supersymmetric particles [14–16]. In 2011, a comprehensive prescription for obtaining well defined mass reconstruction variables was presented [17], which employs the method of constrained minimisation in order to develop variables that emerge as a lower bound on the probed resonance mass. These variables are referred to as “mass bound variables” and shall be investigated along with variables that have been introduced without the descriptive approach of mass bounds, amongst these the visible mass, the effective mass, the vectorial mass, and the collinear mass.

In this thesis, the mass reconstruction variables are constructed for the distinct event topology of resonance production in like-charge $W^\pm W^\pm$ in association with two jets in the fully leptonic final state. Further, their reconstruction performance is compared on the level of event distributions. Characteristic quantities of these distributions are studied such as the mean, the bias and the width. Finally, their discovery potential in proton-proton collisions with a centre-of-mass energy of 8 TeV and an integrated luminosity of 20 fb^{-1} is determined. Contributions from non-resonant production of the same final state in the SM are considered to be the major background in the determination of the discovery potential.

This thesis is organised as follows: Chapter 2 introduces the underlying theoretical concept with emphasis on electroweak interactions and possible new resonances. The experimental conditions for the study of these resonances at the LHC and the ATLAS detector are given in Chapter 3. The process of vector boson scattering and the channel of like-sign $W^\pm W^\pm$ and its event topology is described in Chapter 4. Chapter 5 outlines the procedure of the simulation of Monte Carlo event generation for this scattering process. The steps taken for preparing the events for the analysis are specified in Chapter 6. In Chapter 7, the variables used for the mass reconstruction of the resonances are explained. Finally, the results of the mass reconstruction of resonances in $W^\pm W^\pm$ scattering are presented in Chapter 8. This chapter comprises the study and comparison of the mass reconstruction variables as well as the investigation of the detection potential of these variables for different resonance parameters.

2 Theoretical foundation

In order to verify experimental observation in high energy particle physics a theory describing the fundamental microscopic interactions between elementary particles is required. An extraordinarily well tested candidate theory is the Standard Model of Particle Physics (SM). It presents a combination of quantum mechanics and relativity theory to form a relativistic quantum field theory of the particle content and the forces of our nature. The SM emerged from a conglomerate of many theories which have been formulated in the last century. Known from results of experiments covering the to date accessible energy range, nature can be described by four fundamental forces between the various elementary particles.

A brief introduction to the theoretical concept of the SM is given in the first section of this chapter. A more detailed account can be found in various textbooks [18–20]. The emphasis, however, is placed on electroweak interactions, described in the second section. An approach to describe physics beyond the SM (BSM) in the currently experimentally accessible energy range is an effective field theory (EFT), where new resonances come into play to account for the several shortcomings of the SM.

2.1 Phenomenology of the Standard Model of Particle Physics

Theoretical concept

The SM is a renormalisable relativistic quantum field theory and formulated as a non-Abelian gauge theory based on the approach of Yang and Mills [2]. It describes the dynamics of the elementary constituents of matter and their fundamental interactions, the electromagnetic force, the weak and the strong force. Each interaction is connected to a continuous symmetry, therefore the SM Lagrangian features the invariance under the group of local gauge transformations of the group $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$. The $SU(3)_C$ is the basis for Quantum Chromodynamics (QCD), the theory of the strong force, whereas the unified electroweak force is based on the $SU(2)_L \otimes U(1)_Y$ symmetry group. More details on QCD can be found in [21, 22]. Electroweak interactions are discussed in more detail in the following section. According to Noether's theorem [23], each continuous symmetry is associated to a conserved operator acting on the fermionic spin one-half matter fields. The eigenvalues of these operators represent conserved quantum numbers, the charges which corresponds to each interaction. In addition, vector fields have to be introduced to ensure gauge invariance. In terms of a gauge theory they are also called gauge fields. In quantised theories, such as the SM, the excitations of the gauge fields result in the existence of spin-one gauge bosons, the mediators of the forces between the fermions. The charges corresponding to the interactions are the three colour charges (red, blue, green), the third component of the

weak isospin I_3^W , and the electric charge Q for the strong, weak, and electromagnetic force, respectively. The fourth fundamental force in nature, gravitation, is not part of the SM.

Particle content

Table 2.1 summarises the masses m and electric charges Q of all gauge bosons and their correspondence to the fundamental interactions. The photon γ is the mediator of the electromagnetic interaction, the $W^\pm$ and Z^0 bosons correspond to the weak interaction and eight gluons g act on strongly interacting particles. Photons and gluons are massless in theory due to the exact conservation of the corresponding symmetries. However, weak bosons have a large mass hinting at broken corresponding symmetries. Spontaneous symmetry breaking (SSB) induces masses for the electroweak gauge bosons via the Higgs mechanism. This mechanism introduces the existence of a spin-zero particle, the Higgs boson, which was recently discovered at the LHC¹ with a mass of 125 GeV [4, 5]. The SSB and the Higgs mechanism are briefly outlined in Section 2.2.

bosons		mass m	electric charge Q	interaction	symmetry group
photon	γ	$< 10^{-18}$ eV	0	electromagnetic	} $SU(2)_L \otimes U(1)_Y$
W bosons	$W^\pm$	80.385 GeV	± 1	weak	
Z boson	Z^0	91.188 GeV	0	weak	
gluons	g	0	0	strong	$SU(3)_c$

Table 2.1: The Standard Model spin-one gauge bosons with their masses, electric charges [24] and the interactions they mediate. Masses are given in units of eV which arises from the convention in particle physics to use “natural units” ($c = \hbar = 1$). The mass of gluons is the theoretical value. A mass as large as a few MeV may not be precluded [24].

Matter particles in this theory are represented by spin one-half fermion fields which can be classified into quarks and leptons according to their interactions. Leptons carry electroweak charges and are therefore participating in the weak and electromagnetic force. Quarks additionally have colour charges and hence they are the only particles interacting via the strong force. Both, quarks and leptons, can be arranged in three generations. The only difference between fermions of different generations are their masses and flavours. Other quantum numbers are identical for corresponding fermions of different generations and therefore they identically couple to all spin-one particles. Table 2.2 summarises the masses m and electric charges Q of all fermions of the SM grouped together in the three generations. Each fermion has a partner particle with opposite charges and equal mass, its anti-fermion. The first generation of these fermions form the stable matter while the second and third generation particles are heavier and therefore unstable copies of the first.

¹For an explanation of the experimental set-up refer to Section 3.1.

generation	1 st		2 nd		3 rd		electric charge Q
	flavour	mass	flavour	mass	flavour	mass	
quarks	u	2.3 MeV	c	1275 MeV	t	173.21 GeV	$+\frac{2}{3}$
	d	4.8 MeV	s	95 MeV	b	4.18 GeV	$-\frac{1}{3}$
leptons	ν_e	$< 2 \text{ eV}$	ν_μ	$< 2 \text{ eV}$	ν_τ	$< 2 \text{ eV}$	0
	e^-	511 keV	μ^-	106 MeV	τ^-	1.777 GeV	-1

Table 2.2: The Standard Model fermions with their experimental masses and electric charges [24]. The quark masses are approximate values since they depend on the renormalisation scheme used. For the neutrinos only upper limits on the mass have been confirmed experimentally. The given upper limits correspond to the linear parametrisation of the flavour eigenstates ν_α ($\alpha \in \{e, \mu, \tau\}$) with respect to the mass eigenstates ν_i ($i \in \{1, 2, 3\}$), e.g. $\nu_\alpha = \sum_i U_{\alpha i} \nu_i$ and therefore $m_{\nu_\alpha} = \sum_i |U_{\alpha i}|^2 m_i$. Notice the wide spectrum of the fermion masses covering a range of at least 11 orders of magnitude from the electron-neutrino to the top-quark.

Leptons have an integer electric charge of -1 or 0 for charged leptons and neutrinos, respectively. The charged leptons are the electron e^- , the muon μ^- , and the tau lepton τ^- with their associated neutrinos, the electron-neutrino ν_e , the muon-neutrino ν_μ and the tau-neutrino ν_τ . The corresponding anti-leptons are denoted e^+ , μ^+ , τ^+ , $\bar{\nu}_e$, $\bar{\nu}_\mu$, and $\bar{\nu}_\tau$, which are neutral².

Quarks can be separated by electric charge into up-type quarks ($Q = +2/3$) and down-type quarks ($Q = -1/3$). The up-type quark flavours are up u , charm c and top t and the down-type quark flavours down d , strange s and bottom b . Together with their anti-particles ($\bar{u}$, $\bar{d}$, $\bar{s}$, $\bar{c}$, $\bar{b}$ and $\bar{t}$) they count 12 quark types². As an experimental fact no single quarks have ever been observed³. Only colourless bound states of three quarks (with different colour each) or a quark and an anti-quark (with colour and corresponding anti-colour) are experimentally accessible. Those bound states are called hadrons, amongst them the constituents of the atomic nuclei, the protons and neutrons.

Open questions

The SM, although consolidating an incredible success in modern particle physics, leaves various questions unanswered. For example, it has 61 elementary particles. Why are there so many of them? In addition, it has over 20 arbitrary parameters. What is the origin of these parameters? Can the number of free parameters be reduced to form a more general theory (“Grand Unification” [25, 26])? Another question arises when regarding the masses of the fermions. Why do they cover such a huge range

² In general however, the corresponding term for particles will always refer to the particle and the anti-particle hereafter, e.g. the notation “lepton” comprises leptons and anti-leptons.

³The top quarks might be an exception here, since it decays before forming composite particles.

(“hierarchy problem” [8])? Moreover, why are there three generations of fermions?

Experiments exploring the observable nature of our universe have never observed regions that exist of anti-matter (“baryon asymmetry” [7]). However the SM predicts an almost equal production of matter and anti-matter in the beginning of our universe. As well, only about 4 % of the total energy of the universe consists of baryonic matter as described by the SM. What is the remaining part, the dark matter and dark energy [6], made of? Most obviously is as well another question: How can gravitation be included in the SM?

Many theories are currently being developed as extensions of the SM to seek for answers to those questions. One extension in the electroweak sector of the SM is the approach of effective field theory, which is explained in Section 2.3.

2.2 Electroweak interactions of the Standard Model

The first theory of weak interactions was formulated by E. Fermi in 1933 as an effective four point interaction [27]. This theory, however, only holds for energies up to approximately 100 GeV. Above this range, the heavy vector bosons $W^\pm$ and Z^0 need to be considered as independent dynamical degrees of freedom. In the 1960s S. Glashow could establish a first gauge theory of the weak force based on the Yang-Mills theory [1]. This has only been possible by including the electromagnetic force and thus forming a unified electroweak (EW) theory with the simplest possible gauge symmetry $SU(2)_L \otimes U(1)_Y$. However, Glashow’s theory only contained massless gauge bosons. This problem was solved by A. Salam and S. Weinberg with the introduction of spontaneous symmetry breaking of the underlying gauge symmetry via the Higgs mechanism [28, 29]. This theory was able to predict the existence and masses of heavy gauge bosons. Great effort was put into experiments to find evidence for this theory and a first success was achieved in 1973 with the Gargamelle experiment that discovered the “neutral current” interactions mediated by the Z^0 [30]. Four years later the $W^\pm$ and Z^0 bosons were discovered at the SPS⁴ with masses in the energy range predicted by the electroweak theory [31–34].

The Lagrangian for the electroweak theory can be described by two major parts:

$$\mathcal{L}_{\text{EW}} = \mathcal{L}_{\text{symm}} + \mathcal{L}_{\text{SSB}}. \quad (2.1)$$

$\mathcal{L}_{\text{symm}}$ describes the Yang-Mills theory with the gauge symmetry group $SU(2)_L \otimes U(1)_Y$ and their dynamical degrees of freedom, and $\mathcal{L}_{\text{SSB}}$ contains all necessary terms describing the spontaneous symmetry breaking.

The $SU(2)_L \otimes U(1)_Y$ symmetry of the EW theory

$\mathcal{L}_{\text{symm}}$ can further be split into one part that contains the dynamics and interactions of the fermions fields and a second part for describing the dynamics of the spin-one gauge bosons:

$$\mathcal{L}_{\text{symm}} = \mathcal{L}_{\text{fermion}} + \mathcal{L}_{\text{gauge}}. \quad (2.2)$$

⁴The Super Proton Synchrotron at CERN.

fermion fields			I_3^W	Y	Q
$\begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L$	$\begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix}_L$	$\begin{pmatrix} \nu_\tau \\ \tau^- \end{pmatrix}_L$	$+\frac{1}{2}$	$-\frac{1}{2}$	0
e_R^-	μ_R^-	τ_R^-	$-\frac{1}{2}$	$-\frac{1}{2}$	-1
e_R^-	μ_R^-	τ_R^-	0	-1	-1
$\begin{pmatrix} u \\ d \end{pmatrix}_L$	$\begin{pmatrix} c \\ s \end{pmatrix}_L$	$\begin{pmatrix} t \\ b \end{pmatrix}_L$	$+\frac{1}{2}$	$+\frac{1}{6}$	$+\frac{2}{3}$
u_R	c_R	t_R	$-\frac{1}{2}$	$+\frac{1}{6}$	$-\frac{1}{3}$
d_R	s_R	b_R	0	$+\frac{2}{3}$	$+\frac{2}{3}$
d_R	s_R	b_R	0	$-\frac{1}{3}$	$-\frac{1}{3}$

Table 2.3: Electroweakly interacting fermions and their quantum numbers of the $SU(2)_L \otimes U(1)_Y$ symmetry group. Left-handed fermions form $SU(2)_L$ doublets whereas right-handed fermions are $SU(2)_L$ singlets. I_3^W denotes the third component of the weak isospin, the charge of $SU(2)_L$, Y is the hypercharge corresponding to $U(1)_Y$ and Q is the electric charge. The corresponding anti-fermions have opposite charges. Notice that this table does not comprise right-handed neutrinos, since they have not been observed in any experiment so far and therefore are considered to not interact electroweakly, i.e. all corresponding quantum numbers are zero. However, most scenarios for neutrino masses require the existence of right-handed neutrinos.

As observed in experiments, the electroweak theory is a theory, where left- and right-handed components of the fermion fields⁵ ψ transform differently under the gauge group. This is denoted by the “ L ” of $SU(2)_L$. Left-handed fermions form a $SU(2)_L$ doublet whereas right-handed fermions are $SU(2)_L$ singlets. Table 2.3 contains all observed $SU(2)_L$ fermions with their electroweak charges. Only left-handed fermions and right-handed anti-fermions participate in the charged current, mediated by the $W^\pm$ bosons. This observations result in the final form of the fermionic matter Lagrangian

$$\begin{aligned} \mathcal{L}_{\text{fermion}} = & \bar{L}_L^j i \gamma^\mu D_\mu L_L^j + \bar{l}_R^j i \gamma^\mu D_\mu l_R^j \\ & + \bar{Q}_L^j i \gamma^\mu D_\mu Q_L^j + \bar{u}_R^j i \gamma^\mu D_\mu u_R^j + \bar{d}_R^j i \gamma^\mu D_\mu d_R^j. \end{aligned} \quad (2.3)$$

$L_L^j(Q_L^j)$ denote left-handed $SU(2)_L$ lepton(quark) doublets whereas l_R^j, u_R^j, d_R^j are the right-handed $SU(2)_L$ singlets for leptons, up-type and down-type quarks, respectively. The index j indicates the sum over all fermion families (compare Table 2.3). $\bar{\psi}$ is the adjoint spinor

$$\bar{\psi} \equiv \psi^\dagger \gamma^0 \quad (2.4)$$

with $\psi^\dagger$ the hermitian conjugate and the 4×4 Dirac matrix $\gamma^0 = \begin{pmatrix} \mathbb{1}_{2 \times 2} & 0 \\ 0 & -\mathbb{1}_{2 \times 2} \end{pmatrix}$. Left- and right-handed components of the fermions can be obtained by applying the

⁵Fermion fields are represented by Dirac-spinors, the solutions to the Dirac equation for fermionic matter particles. They have four components and special behaviour under Lorentz transformation.

projection operators P_L and P_R on the spinor

$$\begin{aligned}\psi &= P_L\psi + P_R\psi \\ &= \psi_L + \psi_R\end{aligned}\tag{2.5}$$

where P_L, P_R are defined as⁶

$$P_{L/R} = \frac{1}{2}(1 \mp \gamma_5).\tag{2.6}$$

The projection operators satisfy the relations $P_LP_R = 0$, $P_{L,R}P_{L,R} = P_{L,R}$ and $P_L + P_R = 1$.

Furthermore in Equation (2.3) a sum over μ is implied. γ^μ are the Dirac matrices and D_μ is the covariant derivative with respect to the group $SU(2)_L \otimes U(1)_Y$

$$D_\mu = \partial_\mu + i g_W T^a W_\mu^a + i g_Y Y B_\mu.\tag{2.7}$$

W_μ^a ($a = 1, 2, 3$) and B_μ denote the gauge fields with their corresponding gauge couplings g_W and g_Y , which are introduced to ensure local gauge invariance under $SU(2)_L \otimes U(1)_Y$

$$\begin{aligned}\psi_L &\longrightarrow U_L \quad U_Y \quad \psi_L \\ &\longrightarrow e^{i\alpha_a(x)T^a} e^{i\beta(x)Y} \psi_L, \\ \psi_R &\longrightarrow \quad U_Y \quad \psi_R \\ &\longrightarrow e^{i\beta(x)Y} \psi_R.\end{aligned}\tag{2.8}$$

$T^a = \frac{1}{2}\tau^a$ are the three generators of $SU(2)_L$ with the Pauli matrices τ^a and Y , the hypercharge⁷, is the generator of $U(1)_Y$. The corresponding vector fields are B_μ and W_μ^a , respectively. The four gauge parameters $\alpha_a(x)$ and $\beta(x)$ depend on the space-time referring to a local transformation. The required local gauge invariance of $\mathcal{L}_{EW}$ induces the transformation properties of the gauge fields

$$\begin{aligned}W_\mu^a &\longrightarrow W_\mu^a - \frac{1}{g_W} \partial_\mu \alpha^a(x) - \epsilon^{abc} \alpha^b(x) W_\mu^c \\ B_\mu &\longrightarrow B_\mu - \frac{1}{g_Y} \partial_\mu \beta(x).\end{aligned}\tag{2.9}$$

The second term in Equation (2.2) describes the dynamics of the gauge bosons

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4} W_{\mu\nu}^a W^{a\mu\nu} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu}\tag{2.10}$$

with the antisymmetric field strength tensors

$$B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu,\tag{2.11a}$$

$$W_{\mu\nu}^a = \partial_\mu W_\nu^a - \partial_\nu W_\mu^a + g_W \epsilon^{abc} W_{b,\mu} W_{c,\nu},\tag{2.11b}$$

⁶with $\gamma^5 = \begin{pmatrix} 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & 1 \\ 1 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \end{pmatrix}$

⁷The denotation was chosen on the basis of the electric charge since the underlying symmetry for electromagnetic interactions is as well a different $U(1)$ symmetry.

where ϵ^{abc} are the group structure constants, i.e. the totally antisymmetric Levi-Civita tensors.

The observed interactions determine the dynamical degrees of freedom and the couplings of the fermions to the gauge bosons. The observed charged current leads to the physical degrees of freedom of the charged $W^\pm$ bosons with the relation

$$W_\mu^\pm = \frac{1}{\sqrt{2}} (W_\mu^1 \mp i W_\mu^2) \quad (2.12)$$

to the gauge fields. The charged current couplings are g_W which follows from Equations (2.7) and (2.12). The couplings to the photon need to fulfil the requirements of equal couplings to left- and right-handed fermions with a strength of the electric charge of the positron e as well as no couplings to neutrinos. This leads to the rotation

$$\begin{pmatrix} A_\mu \\ Z_\mu \end{pmatrix} = \begin{pmatrix} \cos \theta_W & \sin \theta_W \\ -\sin \theta_W & \cos \theta_W \end{pmatrix} \begin{pmatrix} B_\mu \\ W_\mu^3 \end{pmatrix} \quad (2.13)$$

of B_μ and W_μ^3 to form the fields of the photon and Z^0 boson, A_μ and Z_μ , respectively. These are the mediators of electroweak neutral currents. As well the fermion couplings to gauge bosons are fixed and have the relation:

$$g_W \sin \theta_W = g_Y \cos \theta_W = e, \quad (2.14)$$

which determines the electric charge generator Q

$$Q = I_3^W + Y. \quad (2.15)$$

With this, the coupling to the Z^0 boson is fixed as well

$$g_Z = \frac{e}{\sin \theta_W \cos \theta_W}. \quad (2.16)$$

Since the $SU(2)_L \otimes U(1)_Y$ is a non-Abelian symmetry group, Equation (2.10) contains, in addition to the dynamics of the gauge fields, cubic and quartic self-interactions amongst those fields. Using Equations (2.12) and (2.13) the possible cubic vertices W^+W^-V and quartic vertices W^+W^-VV , $W^\pm W^\pm W^\pm W^\pm$ and $W^+W^-W^+W^-$ with $V = Z^0, \gamma$ arise (see Figure 2.1).

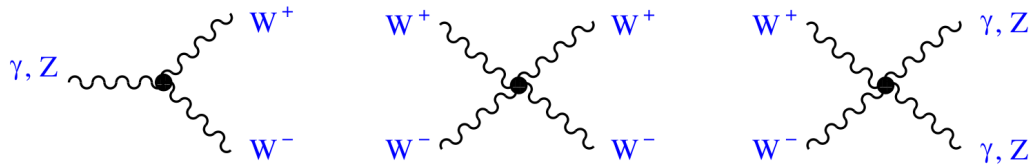


Figure 2.1: Examples for Feynman diagrams for cubic and quartic gauge boson self-interactions [35].

Spontaneous symmetry breaking in the EW theory

With the Lagrangian (2.2) and its explicit parts (2.3) and (2.10) the electroweak interactions of the SM have been introduced. However, this Lagrangian has little to do with nature since all described particles do not have any mass yet. A mass term for spin-one degrees of freedom

$$\frac{1}{2}M_W^2 W_\mu^a W^{a\mu} + \frac{1}{2}M_B^2 B_\mu B^\mu \quad (2.17)$$

violates local gauge invariance of the group $SU(2)_L \otimes U(1)_Y$ because of their inhomogeneous transformation laws (2.9) and fermion mass terms

$$m\psi\bar{\psi} \quad (2.18)$$

lead to mixed terms for left- and right-handed fields which behave differently under the gauge transformation (2.8). This means, in order to generate masses, the local gauge symmetry $SU(2)_L \otimes U(1)_Y$ needs to be spontaneously broken, i.e. the electroweak Lagrangian remains fully symmetric while its vacuum ground state is non-symmetric. This electroweak symmetry breaking (EWSB) is achieved by the Higgs mechanism which was proposed by several physicists⁸ in 1964 [36–38]. Besides the achievement of giving mass to the gauge bosons, as well as to the charged fermions, this mechanism is able to restore unitarity in the process of vector boson scattering.

The Higgs mechanism introduces a new scalar, complex $SU(2)_L$ doublet field

$$\Phi = \begin{pmatrix} \theta^2(x) + i\theta^1(x) \\ \frac{1}{\sqrt{2}}(v + H(x) + i\theta^3(x)) \end{pmatrix} \quad (2.19)$$

with four degrees of freedom, θ^i , the Goldstone modes, and $H(x)$, being real fields. The Lagrangian

$$\mathcal{L}_\Phi = (D^\mu \Phi)^\dagger (D_\mu \Phi) - V(\Phi) \quad (2.20)$$

includes a kinetic term for the scalar field in addition to the potential

$$V(\Phi) = \mu^2 |\Phi|^2 + \lambda |\Phi|^4. \quad (2.21)$$

$V(\Phi)$ is invariant under $SU(2)_L \otimes U(1)_Y$. If $\mu^2 < 0$ and $\lambda > 0$ are chosen the potential features a set of degenerate, non-vanishing ground states

$$\langle 0|\Phi|0\rangle \neq 0 \quad (2.22)$$

where one of these is selected to form the vacuum state

$$\Phi_0 = \begin{pmatrix} 0 \\ \frac{1}{\sqrt{2}}v \end{pmatrix} \quad (2.23)$$

with $v = \sqrt{-\mu^2/\lambda}$. Notice the particular choice of the electrically charged field $\Phi_0^+ = 0$ which ensures the electromagnetic $U(1)_{QED}$ remaining an unbroken symmetry⁹ and

⁸Three parties were working independently on this model: P. Higgs, F. Englert and R. Brout, T. Kibble et al.

⁹Applying the electric charge generator $Q = I_3^W + Y$ on the vacuum state Φ_0 leaves it invariant.

the photon being massless. This leads to the collapse of $SU(2)_L \otimes U(1)_Y$ to the subgroup $U(1)_{QED}$ and the generation of massive weak gauge bosons.

In general one might consider the parametrised form of Φ as in Equation (2.19). However, with local and continuous $SU(2)_L$ invariance of the Lagrangian the θ^i can be eliminated. In fact, this is due to the infinite number of ground states which exhibit the rotational invariance

$$\Phi_0 = e^{i \frac{\tau^a}{2} \theta^a(x)} \begin{pmatrix} 0 \\ \frac{1}{\sqrt{2}} v \end{pmatrix}. \quad (2.24)$$

In the unitary gauge $\theta^a(x) = 0$, Φ becomes

$$\Phi = \begin{pmatrix} 0 \\ \frac{1}{\sqrt{2}} (v + H) \end{pmatrix}. \quad (2.25)$$

Thus, in the unitary gauge, the Goldstone modes decouple completely. With the covariant derivative (2.7) and respecting the mixing of the gauge bosons (Equations (2.12), (2.13)), the kinetic term of the Lagrangian (2.20)

$$(D^\mu \Phi)^\dagger (D_\mu \Phi) = \frac{1}{2} \partial_\mu H \partial^\mu H + (v + H)^2 \left\{ \frac{g_W^2}{4} W_\mu^\dagger W^\mu + \frac{g_W^2}{8 \cos^2 \theta_W} Z_\mu^\dagger Z^\mu \right\} \quad (2.26)$$

reveals mass terms for the weak vector bosons $W^\pm$ and Z^0 at leading order

$$M_W = v \cdot \frac{g_W}{2} = M_Z \cos \theta_W \quad (2.27)$$

as well as the interactions between the neutral spin-zero Higgs particle H and the gauge bosons. Notice that, by construction, the photon remains massless. The Higgs potential

$$V(\Phi) = -\frac{v^2 \lambda}{2} (v + H)^2 + \frac{\lambda}{4} (v + H)^4 \quad (2.28)$$

then exhibits a mass term for the newly introduced Higgs boson

$$M_H = v \sqrt{2\lambda}. \quad (2.29)$$

Since λ is not fixed by theory, the SM cannot make any direct predictions on the Higgs mass. In addition, the Higgs potential $V(\Phi)$ contains the self-interactions of the Higgs boson.

Now, it is worth considering once more the number of degrees of freedom of the Higgs doublet (Equation (2.19)) [35]. The three θ^i can be eliminated, however, the corresponding degrees of freedom are not lost. Since the formerly massless gauge bosons $W^\pm$ and Z^0 gain masses through the SSB, they each obtain a longitudinal polarisation state. The number of degrees of freedom is conserved and the introduction of the θ^i was necessary for obtaining massive gauge bosons. These θ^i are massless spin-zero particles which are referred to as Goldstone bosons in the Goldstone theorem [39, 40]. They correspond to the “flat directions” in the Higgs potential $V(\Phi)$ which connect the degenerate states of minimal energy.

Masses for fermions can be introduced manually via a Yukawa Lagrangian

$$\mathcal{L}_{\text{Yukawa}} = y_{Lj} \bar{L}_L^j \Phi_L^j + y_{uj} \bar{Q}_L^j \Phi^c u_R^j + y_{dj} \bar{Q}_L^j \Phi d_R^j + \text{h.c.} \quad (2.30)$$

and Yukawa couplings y_f with $f = \{L^j, u^j, d'^j\}$. The Higgs doublet Φ connects the left- and right-handed fermion fields, with $\Phi^c = i\sigma_2\Phi^*$ denoting the charge-conjugate scalar field. Applying the unitary gauge, fermion-Higgs interactions and mass terms for fermions evolve with

$$m_f = v \cdot \frac{y_f}{\sqrt{2}}. \quad (2.31)$$

These masses are again not predictable since y_f is a free parameter.

Electroweak symmetry breaking adds additional terms to the $SU(2)_L \otimes U(1)_Y$ Lagrangian (2.2)

$$\mathcal{L}_{\text{SSB}} = \mathcal{L}_\Phi + \mathcal{L}_{\text{Yukawa}} \quad (2.32)$$

adding particle masses as well as a massive, spin-zero particle, the Higgs boson. Its interactions with other particles is always proportional to the particle's mass. Via a non-zero vacuum expectation value v , the Higgs mechanism introduces a weak scale

$$v = \frac{1}{\sqrt{\sqrt{2}G_F}} = 246 \text{ GeV} \quad (2.33)$$

with G_F the Fermi coupling constant, which can be determined by applying the four-fermion Fermi theory to the decay of the muon [27, 41]. The mechanism of EWSB introduced here represents the simplest model including exactly one Higgs doublet. Scenarios of BSM physics in the EWSB sector such as two Higgs doublets as in the MSSM¹⁰ [42] or a more elaborate Higgs multiplet structure would as well solve the mass problem in a $SU(2)_L \otimes U(1)_Y$ theory and are not excluded with the discovery of the Higgs boson with a mass of 125 GeV. These and many other theories include a series of new particles which would evolve as resonances in interactions relevant for EWSB. The following section composes a short theoretical introduction to these phenomena.

2.3 Resonances in the electroweak sector

As discussed in Section 2.1, the SM features some unsolved questions which physics beyond the SM has to account for. The aim of this work is to further explore the EWSB mechanism after the discovery of the Higgs boson. Scenarios of BSM physics in this context shall be approached as model-independently as possible by parametrising the hidden theory while making just a few generic assumptions. The fact that this new theory is not yet known introduces an infinite number of new parameters as there are infinite possibilities to be still tested. This leads to the choice of a non-renormalisable effective field theory (EFT) approach which can be considered as a perturbative expansion in the characteristic energy scale for the underlying physics beyond the SM. Any fixed order of this expansion then contains only a finite subset of parameters which can be tested in experiments. The model discussed in the next sections introduces an effective chiral Lagrangian formalism for the scattering of weak bosons and the further study of EWSB. This Lagrangian needs to reproduce the SM gauge symmetry $SU(2)_L \otimes U(1)_Y$ and its breaking by fermion and boson masses in the low-energy regime where it is already confirmed. At high energies and with the

¹⁰Minimal Supersymmetric Standard Model.

assumption of a conserved custodial isospin symmetry new degrees of freedom can be added to the Lagrangian in form of distinct resonances. However, as the result of an effective theory approach at fixed order, the inclusion of new degrees of freedom leads to the violation of the unitarity of the scattering amplitudes in most of the cases. As long as no complete theory is available, the approach must be manually prevented from violating unitarity, i.e. a formalism has to be introduced in order to prohibit the resonant scattering amplitudes to raise with energy.

In brief, the electroweak chiral EFT Lagrangian, additional resonances and unitarity restoration are the theoretical ingredients for this work. They have been established in [9, 43] and are summarised briefly in the subsequent sections.

Effective field theory approach

In order to parametrise new contributions in the sector of EWSB, an additional 2×2 isospin matrix field

$$\Sigma = e^{-\frac{i}{v}\theta^a T^a} \quad (2.34)$$

with the covariant derivative $D_\mu \Sigma = \partial_\mu \Sigma + i g_W T^a W_\mu^a \Sigma + i g_Y Y \Sigma B_\mu$ is introduced which is invariant under the local $SU(2)_L$ group. θ^a are the scalar massless Goldstone bosons resulting from the spontaneous EWSB. Equation (2.34) denotes only one possible parametrisation of the Σ field. Explicitly introducing the Higgs field as excitation of Σ leads to

$$\Sigma \longrightarrow \left(1 + \frac{1}{v}H\right) \Sigma. \quad (2.35)$$

Now, with the custodial symmetry $SU(2)_c \cong SU(2)_L \otimes SU(2)_R$ doublet notation for both, the left- and right-handed fermion fields

$$Q_{L/R}^j = \begin{pmatrix} u^j \\ d^j \end{pmatrix}_{L/R}, \quad L_{L/R}^j = \begin{pmatrix} \nu^j \\ l^j \end{pmatrix}_{L/R}. \quad (2.36)$$

an EW effective chiral Lagrangian can be constructed as expansion¹¹ in E/Λ . In the leading order of this expansion it can be written as [9]:

$$\begin{aligned} \mathcal{L}_{EFT} = & - \left(\frac{v}{2}\right)^2 \text{tr} [V_\mu V^\mu] - \frac{1}{4} B_{\mu\nu} B^{\mu\nu} - \frac{1}{4} W_{\mu\nu}^a W^{a,\mu\nu} \\ & - (\bar{Q}_L^i \Sigma M_Q^{ij} Q_R^j + \bar{L}_L^i \Sigma M_L^{ij} L_R^j + \text{h.c.}) \\ & - \bar{L}_L^{ic} \Sigma^* M_{N_L}^{ij} \frac{1 + \tau^3}{2} \Sigma L_L^{jc} - \bar{L}_R^{ic} M_{N_R}^{ij} \frac{1 + \tau^3}{2} L_R^j \\ & + \bar{Q}_L^j i \gamma^\mu D_\mu Q_L^j + \bar{Q}_R^j i \gamma^\mu D_\mu Q_R^j + \bar{L}_L^j i \gamma^\mu D_\mu L_L^j + \bar{L}_R^j i \gamma^\mu D_\mu L_R^j \\ & - \mu^2 \left(\frac{v}{2}\right)^2 \text{tr} [\Sigma^\dagger \Sigma] + \lambda \left(\frac{v}{2}\right)^4 \text{tr} [\Sigma^\dagger \Sigma]^2. \end{aligned} \quad (2.37)$$

The M_f^{ij} denote the possibly non-diagonal mass matrix of the corresponding fermions $f \equiv \{Q, L, N_L, N_R\}$, which includes the general case of right-handed neutrinos¹², and

$$V_\mu = \Sigma (D_\mu \Sigma)^\dagger. \quad (2.38)$$

¹¹ $\Lambda = 4\pi v$ denotes the energy scale where physics beyond the SM dominates, i.e. perturbation theory breaks down.

¹²The subscript “c” for the fermion fields in Equation (2.37) denotes the charge conjugate, e.g. $L^c = -i\tau^2 L^*$

The first term contains the mass terms of the weak gauge bosons and the generalised Higgs boson kinematics. The two following terms correspond to the SM terms of the transversal degrees of freedom of the weak gauge bosons in Equation (2.2) with the field strength tensors defined according to Equation (2.11). The second and third line introduces the fermion masses as well as couplings to longitudinal weak gauge bosons. The fourth line again corresponds to the SM kinetic terms of the fermions in (2.2). The covariant derivative is as in (2.7). The last two terms denote the general form of the EWSB potential with the special case being the Higgs potential (2.21).

Resonances

Additional degrees of freedom, which represent new resonances in EWSB, can be added to this leading order in E/Λ chiral Lagrangian. These resonances can either represent elementary particles of the new theory or they are itself bound states of more fundamental objects of a yet more fundamental theory. If assumed that the custodial symmetry $SU(2)_c$ is conserved at high energies only five possible spin-isospin combinations for resonances which can occur in quasi-elastic weak boson scattering $VV \rightarrow VV$ remain. They are listed in Table 2.4. These resonances transform as matter fields under local $SU(2)_L$ according to their isospin representation.

type	isospin I	spin J	electric charge Q	width $\Gamma_{\text{res}_i} = \frac{\Gamma_{\text{res}}}{\Gamma_0}$
σ	0	0	0	6
f	0	2	0	$\frac{1}{5}$
ρ	1	1	$-1, 0, +1$	$\frac{4}{3} \frac{v^2}{M^2}$
ϕ	2	0	$-2, -1, 0, +1, +2$	1
t	2	2	$-2, -1, 0, +1, +2$	$\frac{1}{30}$

Table 2.4: Properties of all resonances that can couple to the electroweak sector of the SM without violation of isospin, i.e. requiring conservation of the custodial symmetry $SU(2)_c$: isospin I , spin J , possible electric charge Q configurations and the scaled width Γ_{res_i} with $\Gamma_0 = \frac{g^2 M^2}{64\pi v^2}$ [43].

These resonances can be related to different explicit models of BSM physics. For example, the scalar isotensor resonance ϕ corresponds to Little Higgs models [44–48]. The vector isovector resonance ρ occurs in Technicolor [49–51] as well as extra-dimension models [52–55]. The tensor isoscalar resonance f can represent a graviton resonance [56]. The tensor isotensor resonance t is rather exotic and plays a role only in extended composite models. Following the scope of this work, only resonances that can occur in the s-channel scattering $W^\pm W^\pm \rightarrow W^\pm W^\pm$ are further discussed in detail. These are the doubly electrically charged components of the isotensor resonances ϕ and t .

These two resonances can be added explicitly to the effective Lagrangian (2.37) [43]

$$\mathcal{L}_\phi = \mathcal{L}_{\text{kin}}^\phi - \frac{1}{2} \text{tr} [\phi j_\phi], \quad (2.39a)$$

$$\mathcal{L}_t = \mathcal{L}_{\text{kin}}^t - \frac{M_t^2}{4} \text{tr} [t_{\mu\nu} t^{\mu\nu}] + \frac{1}{2} \text{tr} [t_{\mu\nu} j_t^{\mu\nu}], \quad (2.39b)$$

with kinetic and mass terms $\mathcal{L}_{\text{kin}}$, which are not necessary to introduce for describing the interactions of the resonances. In addition, a term including the linear coupling of each resonance to longitudinal vector boson currents is included with

$$j_\phi = -\frac{g_\phi v}{2} \left(V_\mu \otimes V^\mu - \frac{\tau^{aa}}{6} \text{tr} [V_\mu V^\mu] \right), \quad (2.40a)$$

$$j_t^{\mu\nu} = -\frac{g_t v}{2} \left[\frac{1}{2} (V^\mu \otimes V^\nu + V^\nu \otimes V^\mu) - \frac{g^{\mu\nu}}{4} V_\rho \otimes V^\rho \right] - \frac{g_t v}{2} \left[-\frac{\tau^{aa}}{6} \text{tr} [V^\mu V^\nu] + \frac{g^{\mu\nu} \tau^{aa}}{24} \text{tr} [V_\rho V^\rho] \right] \quad (2.40b)$$

expressed by the field V_μ (Equation 2.38) and $\tau^{aa} = \tau^a \otimes \tau^a = \tau^3 \otimes \tau^3 + 2\tau^+ \otimes \tau^- + 2\tau^- \otimes \tau^+$, the tensor product of Pauli matrices τ^a ($a = 1, 2, 3$) with $\tau^\pm = \tau^1 \pm i\tau^2$.

Equations (2.39) constitute the minimal Lagrangian for resonances coupling to the EWSB sector. Couplings to fermions and transversal gauge bosons including the photon are omitted since those phenomena do not contribute to EWSB. The latter are in any case expressed via the couplings g_W and g_Y which are suppressed at high energies where the terms E/Λ are dominant. Then, the partial decay width for the decay into longitudinal vector bosons becomes the total width and can be calculated in leading order using the Goldstone-Boson equivalence theorem (GBET) [57]. This approach therefore is able to set a lower bound on the total width of the resonances

$$\Gamma_{\text{res}} = \Gamma_{\text{res}_i} \cdot \frac{g_{\text{res}}^2}{64\pi} \cdot \frac{m_{\text{res}}^3}{v^2}. \quad (2.41)$$

where for each resonance Γ_{res_i} is listed in Table (2.4) and g_{res} is, together with the mass of the resonance, a free parameter of the model. For the two resonances considered here, ϕ features a considerably broader width than t , assuming identical masses.

Naturally one would suppose that nature has realised a superposition of several of those resonances. However, in this thesis only one resonance is assumed to be present at a time.

Unitarity

The experimental approach for searching for resonances is studying their production cross sections

$$\sigma = \int \frac{1}{64\pi^2 s} |\mathcal{M}|^2 d\Omega \quad (2.42)$$

at a certain centre-of-mass energy $\sqrt{s}$ and with $d\Omega = d\cos\theta d\varphi$ being the area element. Directly related are their scattering amplitudes

$$A = \frac{1}{64\pi^2 s} |\mathcal{M}|^2. \quad (2.43)$$

$\mathcal{M}$ can be calculated as transition matrix element of the scattering matrix S between an initial ψ_i and a final ψ_f state

$$\mathcal{M} = \langle \psi_f | S | \psi_i \rangle. \quad (2.44)$$

The scattering matrix inherits the theoretical concept of the underlying process, here the resonant scattering of two vector bosons. Since an effective theory in the form of a truncated perturbative expansion is employed for the case at hand the completeness of the theory is missing and therefore unitarity is not guaranteed any more. In fact the next-to-leading order (NLO) in the E/Λ expansion of the resonant scattering amplitudes¹³ asymptotically rises with energy [43]

$$A^\phi(s, t, u) = -\frac{g_\phi^2}{4v^2} \left(\frac{t^2}{t - M_\phi^2} + \frac{u^2}{u - M_\phi^2} - \frac{2}{3} \frac{s^2}{s - M_\phi^2} \right), \quad (2.45a)$$

$$\begin{aligned} A^t(s, t, u) = & -\frac{g_t^2}{24v^2} \left(\frac{t^2}{t - M_t^2} P_2(s, t, u) + \frac{u^2}{u - M_t^2} P_2(s, t, u) \right) \\ & + \frac{g_t^2}{24v^2} \left(\left(\frac{2}{3} \frac{s^2}{s - M_t^2} - \frac{s^2}{6M_t^2} \right) P_2(s, t, u) \right) \end{aligned} \quad (2.45b)$$

with $P_2(s, t, u) = [3(t^2 + u^2) - 2s^2]/s^2$ and t and u being the Mandelstam variables. This behaviour is forbidden since the scattering amplitudes in this case violate unitarity at some energy, i.e. the scattering probability exceeds the value one. To prevent this a unitarisation formalism needs to be introduced manually to the theory. Therefore, the scattering amplitudes need to be expanded in partial waves, i.e. spin-isospin eigenamplitudes [9]

$$\mathcal{M} = 32\pi \sum_l a_l(s) (2l + 1) P_l(\cos \theta). \quad (2.46)$$

Following the optical theorem

$$\sigma = \frac{\Im(\mathcal{M}(\theta = 0))}{s}, \quad (2.47)$$

which is directly connected to the unitarity of the S-matrix, a condition for each partial scattering wave $a_l(s)$ can be derived

$$\left| a_l(s) - \frac{i}{2} \right| = \frac{1}{2} \quad (2.48)$$

if one assumes only elastic scattering processes. This means that the spin-isospin eigenamplitudes for elastic scattering have to stay on the Argand circle [9], which is a circle centred at $i/2$ and with a radius of $1/2$. The WHIZARD [58] event generator used in these studies uses the method of K-matrix unitarisation. This unitarisation formalism projects the eigenamplitude onto the Argand circle. The projected eigenamplitude $a_K(s)$ takes the form

$$a_K(s) = \frac{1}{\text{Re}(1/a_l(s)) - i}. \quad (2.49)$$

¹³Notice that these amplitudes do not include widths of the resonances.

Geometrically, $a_K(s)$ is the intersection point between the Argand circle and the straight line connecting the points $a_l(s)$ on the real axis and i on the imaginary axis (see Figure (2.2)). The unitarised amplitude as a result asymptotically approaches saturation, which can formally be seen as a resonance with infinite mass and width. Resonant amplitudes which can be derived from Lagrangian (2.39) and are shown in Equations (2.45) are transformed into Breit-Wigner amplitudes through the unitarisation process. The resulting K-matrix unitarised cross section for $VV \rightarrow VV$ scattering with the presence of resonances is illustrated exemplarily for one resonance parameter combination in Figure 2.3.

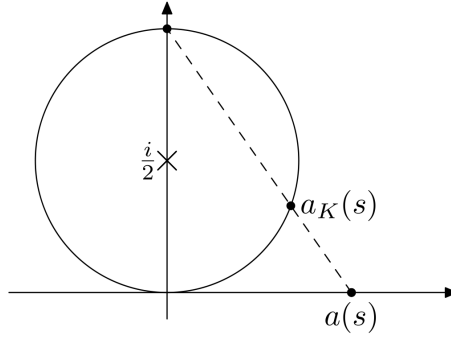


Figure 2.2: Illustration of the K-matrix unitarisation procedure: The real spin-isospin eigenamplitude $a(s)$ is projected onto the Argand-circle (solid) along the straight (dashed) line which connects the points $a(s)$ and i . $a_K(s)$ is the new unitarised scattering eigenamplitude [43].

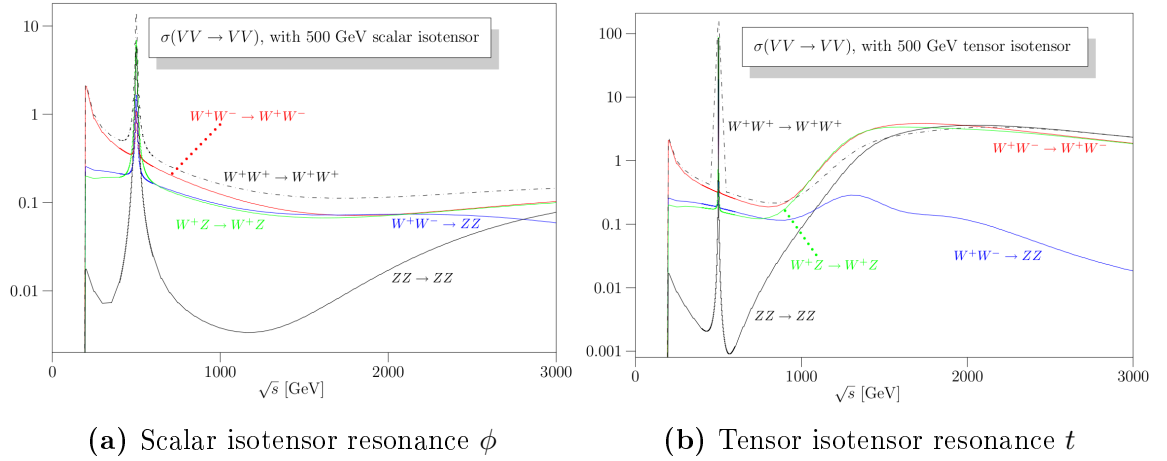


Figure 2.3: K-matrix unitarised cross sections for $VV \rightarrow VV$ scattering with the presence of resonances of mass 500 GeV and $g = 1$ [43].

3 Experimental set-up

This work presents a feasibility study on resonances in EWSB following the recently published analysis on electroweak production of $W^\pm W^\pm jj$ [10]. Since this analysis was conducted with the ATLAS experiment at the LARGE HADRON COLLIDER (LHC), this chapter provides a brief introduction to the experimental set-up for which this study is designed. More information about the design, construction and operation of the LHC can be found in Ref. [59].

The experiment is hosted by the European Organisation for Nuclear Research CERN¹ on the border between Switzerland and France near Geneva. Founded in 1954, CERN now counts 21 mostly European member states stating the organisation as the leading international research facility in particle and high energy physics. With the plan of building the LHC international collaborations such as ATLAS have been established at CERN. The ATLAS collaboration involves about 3 000 scientists from over 38 countries. Their aim is to test predictions of theories of particle physics and search for evidence of yet unknown phenomena.

3.1 The Large Hadron Collider

The Large Hadron Collider [59] is a circular particle accelerator and collider. It provides the worlds largest experimental tool to study particle collisions at an exceptional energy and luminosity range. This chapter summarises the technical details of the LHC which are stated in [59].

Built in the tunnel of the former LEP² accelerator the LHC started operation in September 2008. The first phase of physics data-taking (called Run 1) began in 2010. Since early 2013 the LHC is shut down for upgrades and renovations of the accelerator complex as well as the detectors. No data-taking takes place in this “Long Shutdown 1” which endures until beginning of 2015.

In Run 1, a centre-of-mass energy of up to 8 TeV and a peak instantaneous luminosity of $7.7 \times 10^{33} \text{ cm}^{-2} \text{ s}^{-1}$ was reached. A number of 1 380 bunches of 50 ns spacing formed the beam for this running period.

Technical details

The LHC is a proton-proton³ accelerator designed for collisions at a centre-of-mass energy of 14 TeV and a peak luminosity of $10^{34} \text{ cm}^{-2} \text{ s}^{-1}$. As can be seen in Figure 3.1, four major experiments (ALICE [60], ATLAS [61], CMS [62] and LHCb [63]) are arranged along the accelerator ring, which has a circumference of circa 26.7 km and is

¹CERN originally French for Conseil Européen pour la Recherche Nucléaire

²Large Electron-Positron Collider, 1989 - 2000

³The LHC also performs heavy ion (Pb) collisions.

located between 45 m and 170 m under ground. Two particle beams travel in opposite directions in separate beam pipes being intersected at the locations of the four experiments. The beam pipe is operated in ultra-high vacuum. Before reaching the LHC the protons need to be pre-accelerated by a system of smaller accelerators, which built an entire complex of accelerators of different size and energy range.

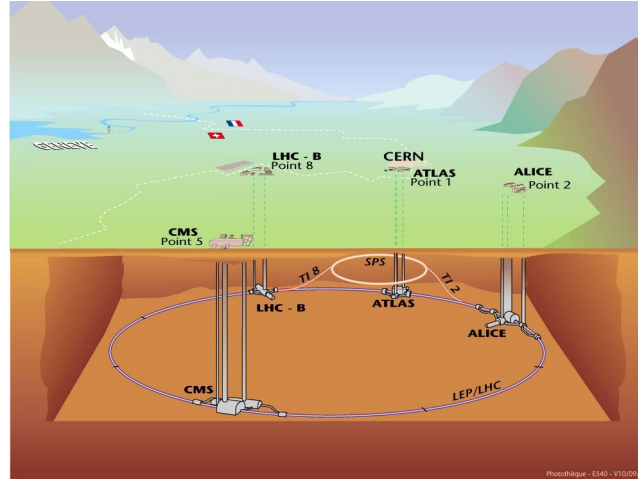


Figure 3.1: Scheme of the LHC showing the location of the four experiments ALICE [60], ATLAS [61], CMS [62] and LHCb [63]. Picture taken from [64].

At design operation each beam is composed of 2 808 bunches with a 25 ns spacing and an intensity of circa 1.6×10^{11} protons per bunch. A range of 1 232 superconducting dipole magnets with a magnetic field strength of up to 8.3 T bend the beam along the circular trajectory in the ring. In order to achieve such high magnetic field strengths, the dipole magnets along with 392 quadrupole magnets for beam focusing are operated at a temperature of 1.9 K, where superconductivity is ensured. This temperature is reached through a cryogenic system running with super-fluid helium-4.

The event rate of a given process

$$\frac{dN_{\text{process}}}{dt} = \mathcal{L} \cdot \sigma_{\text{process}} \quad (3.1)$$

is proportional to the cross section σ of this process which is energy dependent. The factor of proportionality is the luminosity

$$\mathcal{L} = \frac{N_b^2 f n_b}{4\pi\sigma_x\sigma_y} \cdot F \quad (3.2)$$

which gives a measure of the number of collisions happening with a certain accelerator set-up. It depends on machine parameters only. N_b is the number of protons per bunch, f the revolution frequency, n_b the number of bunches and $\sigma_{x,y}$ the beam width in x,y direction for Gaussian beam profiles. F denotes the various geometric corrections that need to be considered for real shaped beams [65]. These considerations motivate a high centre-of-mass energy and high beam intensities for the LHC.

The design beam conditions give an average crossing rate of 32 MHz producing several petabytes of data every year taken by the detectors. In order to handle these huge amounts of data, the LHC has its own computing grid infrastructure, the largest worldwide.

3.2 The ATLAS detector

The ATLAS⁴ detector [61] is a general purpose detector designed, constructed and operated by the ATLAS collaboration. The cylindrical installation of concentric layers around the beam pipe with a total dimension of 44 m in height and 25 m in diameter is displayed in Figure 3.2. Furthermore, noticeable is the forward-backward symmetry with respect to the interaction point. Related to this symmetry, the detector is sectioned into the central barrel region with one end-cap region perpendicular to the beam pipe on either side.

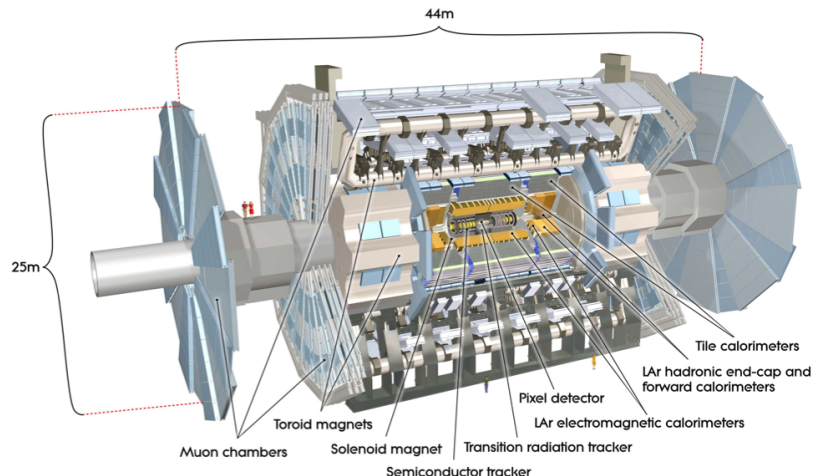


Figure 3.2: A sketch of the ATLAS detector at the LHC [61].

Four major components constitute the detector's layout: the Inner Detector, the calorimeters, the Muon Spectrometer and the magnet system, which follow the conceptual requirements for the detector. Identifying outgoing particles and reconstructing the underlying kinematics of a collision is the core purpose for this detector. Hence, a high precision measurement of particle energies, momenta and charges must be ensured resulting in a highly hermetic detector with reasonable choices for materials and detection mechanisms. Furthermore, the detector has to withstand large collision rates with approximately 1 000 particles per collision. It, therefore, has to provide a fast readout and trigger system as well as radiation hardness.

In 2012, at a centre-of-mass-energy of 8 TeV, the total integrated luminosity collected by the ATLAS detector was 20.3 fb^{-1} .

In order to handle the large instantaneous event rates of 40 MHz, a trigger system is implemented in the data taking process. The object is to reject elastic scattering events while keeping interesting inelastic scattering events. The trigger efficiency is an important characteristics for physics analysis.

Following in brief is an introduction of the ATLAS coordinate system, since it will be referenced throughout the analysis. After that, a description of the different parts of the detector is given in which the emphasis is placed on the calorimeters.

⁴A Toroidal LHC ApparatuS

3.2.1 The ATLAS coordinate system

The interaction point of interest (nominal vertex) is chosen as the origin of a right-handed coordinate system. The beam pipe defines the z -axis and the x - y plane forms the transverse plane which lies perpendicular to the beam pipe providing the azimuthal angle ϕ . θ defines the angle between the measured particle and the positive z -axis. The rapidity is defined as

$$y = \frac{1}{2} \ln \left[\frac{E + p_z}{E - p_z} \right], \quad (3.3)$$

which has the essential property of being invariant under Lorentz boosts in the longitudinal direction. In the limit of massless particles, the rapidity is equal to the geometrically defined pseudorapidity

$$\eta = -\ln \tan \left(\frac{\theta}{2} \right). \quad (3.4)$$

Since mainly leptons can be considered massless in respect to the nominal centre-of-mass energy the pseudorapidity is a good approximation of y for these particles. An angular distance between two objects in the detector can then be defined through the distance parameter

$$\Delta R = \sqrt{(\Delta\eta)^2 + (\Delta\phi)^2}. \quad (3.5)$$

The transverse momentum p_T , the transverse energy E_T and the missing transverse energy E_T^{miss} are defined in the x - y plane, here exemplary for p_T :

$$p_T^2 = p_x^2 + p_y^2. \quad (3.6)$$

3.2.2 The tracking systems

Measurements of tracks of particles are obtained in the tracking detectors of the ATLAS detector. Two tracking detectors are incorporated in this experiment, the Inner Detector and the Muon Spectrometer which built the inner- and outermost parts of the ATLAS detector, respectively. Tracking detectors are composed of several layers of material, which trace the trajectory of charged particles through it. When penetrating one of the layers the particles interact via ionisation, thus only the tracking of charged particles is possible. Equipping the tracking volume with an appropriate magnetic system parallel to the beam axis allows to measure the transverse momentum and the charge of charged particles using the Lorentz force. The particles will be bend in the transverse plane with a curvature of radius r_T whose magnitude depends on the magnetic field strength, and the transverse momentum p_T and absolute charge $|q|$ of the particle. The direction of the curvature depends on the sign of the charge of the particle. With the known magnetic field strength the magnitude of the transverse momentum can be determined

$$p_T = |q| B r_T. \quad (3.7)$$

The Inner Detector Figure 3.3 shows a scheme of the Inner Detector which contains three types of tracking detectors. The magnetic field for momentum measurement is supplied by a superconducting solenoid of 2 T which surrounds the Inner Detector.

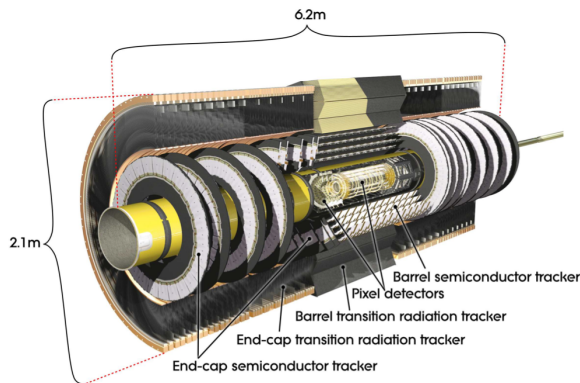


Figure 3.3: The ATLAS Inner Detector with the innermost semiconductor pixel detector, the Semiconductor Tracker (SCT) and the Transition Radiation Tracker (TRT) in the barrel and forward region [61].

Innermost located is a semiconductor pixel detector providing a high momentum and vertex resolution in the region with the highest track density very close to the beam pipe. It covers a region of $|\eta| < 2.5$. It consists of three pixel layers which have an accuracy of $10\,\mu\text{m}$ in $R - \phi$ and $115\,\mu\text{m}$ in z or R for the barrel or end-cap region, respectively.

The middle detector, the Semi-Conductor Tracker (SCT), is a silicon strip detector. As the pixel detector, the SCT covers the region $|\eta| < 2.5$ with four layers being crossed by each particle. Each layer consists of two sensor panels, slightly twisted against each other making it possible to measure both, the $R - \phi$ and z direction simultaneously in each layer. Each space point can be measured with an accuracy of $17\,\mu\text{m}$ in $R - \phi$ and $580\,\mu\text{m}$ in z in the barrel and R in the end-cap region, respectively.

The Transition Radiation Tracker (TRT) constitutes the outermost part of the Inner Detector. It is a straw tube detector filled with a xenon-based gas mixture in 4mm straws. It provides a large number of hits, typically 36, per particle and a coverage of up to $|\eta| < 2.0$. Per construction, the TRT only gives information in the $R - \phi$ plane with an accuracy of approximately $130\,\mu\text{m}$ per straw. The TRT features the capability to distinguish between electrons and other, heavier particles via the generation and detection of transition radiation⁵.

The large number of measurements in the TRT and the high precision measurements in the pixel detector and SCT play together to perform very robust track reconstruction, precise momentum and vertex measurement and an enhanced electron identification. The overall transverse momentum resolution achieved in the Inner Detector is

$$\frac{\sigma_{p_T}}{p_T} = 0.05\% \, p_T(\text{GeV}) \oplus 1\%, \quad (3.8)$$

where the first term arises from the momentum measurement via the curvature in a magnetic field and the second term includes multiple scattering.

⁵A relativistic, charged particle produces transition radiation when it passes the surface of division of two media with different dielectric constants. The generation of transition radiation depends on the mass of the particle.

The Muon Spectrometer The Muon Spectrometer is the outermost part of the ATLAS detector. Besides muons all particles that can be observed by the ATLAS detector are absorbed in the layers between the Inner Detector and the Muon Spectrometer, the calorimeters (see the following section). Muons are the only measurable particles that can penetrate all layers of the detector, thus escaping the detector without being stopped. This observation is a result of the underlying interaction mechanism of muons with the detector material. In principle, muons are leptons, as electrons, but 200 times heavier and therefore not interacting predominately by Bremsstrahlung, but by ionisation leaving only little energy in the inner layers of the detector.

Consisting of three layers of high precision tracking chambers and large toroid magnets the Muon Spectrometer supplies a complementary measurement of the muon momentum and charge compared to the Inner Detector. Different types of ionisation chambers are used for detection: Monitored Drift Tubes (MDT), Cathode Strip Chambers (CSC), Resistive Plate Chambers (RPC) and Thin Gap Chambers (TGC). The acceptance range in pseudorapidity is $|\eta| < 2.7$. The relative resolution in transverse momentum in the Muon Spectrometer is 10 % for 1 TeV muons.

3.2.3 The calorimeters

In high energy particle physics experiments, calorimeters are commonly used to measure the energy of particles. In the ATLAS detector the calorimeters are situated between the Inner Detector and the Muon Spectrometer which is an inevitable requirement because the energy of a particle is measured through its total absorption. Besides muons all electromagnetically or hadronically interacting particles up to a certain energy range leave their total energy in the layers of the calorimeters i.e. being completely stopped. As a result, unlike in the tracking detectors, also neutral particles such as photons and neutral hadrons are detected. The construction of highly hermetic calorimeters is an important feature which allows to measure the energy of jets as well as the missing transverse energy $\vec{E}_T^{\text{miss}}$ in the event.

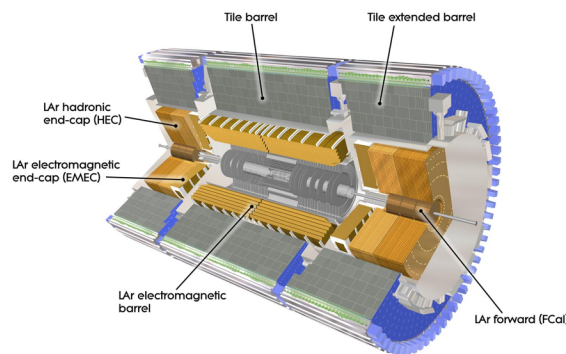


Figure 3.4: The ATLAS Electromagnetic and Hadronic Calorimeter [61].

The ATLAS calorimeters consist of an inner Electromagnetic and outer Hadronic Calorimeter as shown in Figure 3.4. Together with the Forward Calorimeter it consists

of three distinct parts. Electrons and photons can be identified in the Electromagnetic Calorimeter leaving their energy by forming electromagnetic particle showers. Hadrons, which are significantly heavier and mainly interact via the strong force, propagate further causing more wide-spread hadronic showers mainly in the denser Hadronic Calorimeter.

Both, the Electromagnetic and Hadronic Calorimeter are sampling calorimeters which feature an inhomogeneous body of two alternating materials. One medium serves as the detection medium whereas the other medium, a dense metal, enhances the shower propagation. With this construction only the energy deposited in the sensitive medium is detected enforcing the estimation of the total deposited energy. This reduces the resolution but also the costs. Besides in the barrel region of the hadronic calorimeter, which consists of scintillation tiles, liquid argon (LAr) is used as the sensitive medium. The whole calorimeter covers a region of $|\eta| < 4.9$. Since the total absorption of high-energy particles must be guaranteed the thickness plays an important role in the design of a calorimeter. This is specified in radiation or interaction lengths, X_0 and λ , respectively⁶. The total thickness of the whole active calorimeter is $\lambda = 9.7$ at $\eta = 0$. This sufficiently prevents high-energetic particles to penetrate to the Muon Spectrometer while adequately letting pass muons.

Energy resolution Since, the measurement of $\vec{E}_T^{\text{miss}}$ plays a crucial role in this analysis, an emphasis is placed on the resolution of the calorimeters. The energy resolution is an important characteristics of calorimeters which denotes the extend to which a detector can distinguish two close-by energies. It is mainly statistically dominated resulting in an energy measurement representing a roughly Gaussian distribution with a non-zero width. The resolution is then given as the full width at half maximum FWHM σ_E . The relative resolution of a calorimeter can be parametrised as follows

$$\frac{\sigma_E}{E} = \frac{a}{\sqrt{E(\text{GeV})}} \oplus b \oplus \frac{c}{E(\text{GeV})}. \quad (3.9)$$

The first term is the stochastic term with contributions from quantum as well as sampling fluctuations. Quantum fluctuations include shower fluctuations which determine the intrinsic resolution of the physical shower process. The number of particles produced in the shower is dependent on the minimal energy needed for ionisation W ,

$$N = \frac{E}{W} \quad (3.10)$$

with E the energy of the incoming particle. Thus, the higher the incoming particle's energy, the higher the number of particles produced in the shower leading to an increasing resolution with the energy of the incoming particle

$$\frac{\sigma_E}{E} \propto \frac{\sqrt{N}}{N} = \frac{1}{\sqrt{N}} \longrightarrow \frac{\sigma_E}{E} \propto \frac{1}{\sqrt{E}}. \quad (3.11)$$

⁶The term radiation length X_0 is related to the energy loss of high-energy, electromagnetically interacting particles in matter. It denotes the mean distance over which the energy of an electron/positron is diminished by $1/e$, or $7/9$ of the mean free path for high-energy photons making pair production. The interaction length λ is used for hadronic calorimeters and denotes the mean free path of a strongly interacting particle between two inelastic scatterings.

In addition to shower fluctuations, other statistical fluctuations such as photo-electron statistics for scintillating calorimeters, sampling fluctuations and landau fluctuations contribute to the first term in Equation (3.9). The second term of the parametrisation comes from shower leakage due to the finite size of the calorimeter with contributions especially from the longitudinal leakage. The third term is the constant term caused by instrumental effects such as electronic noise. The target in energy resolution for the different calorimeters used in the ATLAS detector are summarised in Table 3.1. As can be seen from Equations (3.8) and (3.9) for the p_T resolution in the Inner Detector and the energy resolution in the calorimeters, respectively, calorimeters constitute a complementary measurement to the Inner Detector. They show better resolution for high energies while the p_T resolution diminishes with increasing transverse momentum. In addition, the calorimeters are designed in thickness so that the shower leakage is well below the statistical limitations.

calorimeter component	required resolution
Electromagnetic Calorimeter	$\sigma_E/E = 10\%/\sqrt{E(\text{GeV})} \oplus 0.7\%$
Hadronic Calorimeter	$\sigma_E/E = 50\%/\sqrt{E(\text{GeV})} \oplus 3\%$
Forward Calorimeter	$\sigma_E/E = 100\%/\sqrt{E(\text{GeV})} \oplus 10\%$

Table 3.1: Resolution requirements for the different calorimeter components of the ATLAS detector [61]

4 Electroweak vector boson scattering at the LHC

As discussed in Section 2.2, the SM comprises the self-interactions of electroweak gauge bosons, amongst these vertices the triple and quartic self-couplings. Together with the Higgs production and exchange these processes are referred to as “Vector Boson Scattering” (VBS). While triple gauge boson couplings have already been observed at LEP [66], the observation of quartic self-coupling process remained one of the last unseen processes after the discovery of the Higgs boson [4]. A first evidence for VBS including this process in the channel $W^\pm W^\pm \rightarrow W^\pm W^\pm$, however, has recently been published [10].

Since the scattering of the longitudinal components of massive weak bosons is directly connected to the mechanism of EWSB, the study of VBS in this channel promises to reveal a deeper understanding of these phenomena. This process can not only provide a complementary study on EWSB to direct Higgs searches, but as well deliver a first insight into physics beyond the SM in this sector. Open questions can be addressed such as whether the found Higgs boson is the SM Higgs boson or if EWSB is realised in a more complex theory with new heavy resonances to be discovered.

At the LHC the process of $W^\pm W^\pm \rightarrow W^\pm W^\pm$ scattering with leptonic decay of both $W^\pm$ bosons is especially interesting due to its preeminent signal-to-background ratio [10].

4.1 The vector boson scattering process

The tree level process of VBS at the LHC can only be realised through the radiation of vector bosons from the initial quarks of the colliding protons. In the framework of the SM, the fundamental, quasi-elastic interaction of the radiated vector bosons comprises five Feynman diagrams, which are shown in Figure 4.1. These are the triple and quartic gauge boson self-couplings, as well as SM Higgs exchange and production, which provides unitarity conservation for the whole process of VBS. The possible outgoing massive weak vector boson pairs are then same-sign $W^\pm W^\pm$, opposite-sign $W^\pm W^\mp$, ZZ , and $W^\pm Z$. The final states of these VBS processes contain six particles, two quarks and the decay products of both vector bosons. Considering the case of only leptonic vector boson decays, which is reasonable for a hadron collider as the LHC, these can be charged leptons and neutrinos. The same final state can also be produced via processes that do not contain VBS. All processes which have this same final state can be grouped together into two sets¹ [10]:

¹Interference between those processes are possible as well [10].

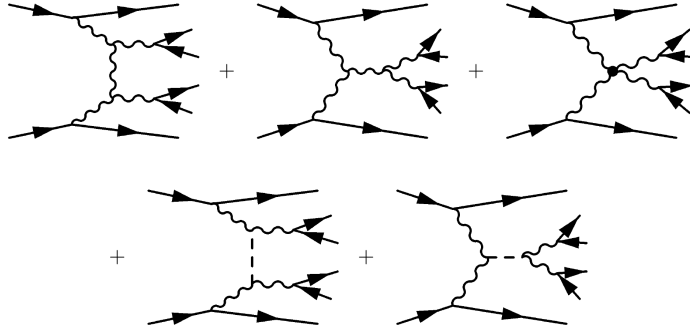


Figure 4.1: Feynman diagrams which are comprised in VBS at the LHC. The upper row are the triple and quartic self-interaction processes of the vector bosons, while the lower processes depict the interaction of vector bosons with the Higgs boson [10].

- The set of VVjj-EW contains VBS as well as all possible SM processes of order $\mathcal{O}(\alpha_{EW}^6)$ in the electroweak coupling parameter. They are purely electroweak. Example Feynman diagrams of EW non-VBS processes are shown in Figure 4.2, amongst them in particular diagrams which show different kinematic behaviour with respect to VBS such as the triple vector boson production with one of them decaying hadronically.
- The set of VVjj-QCD contains processes with strong contributions in the only possible remaining order of $\mathcal{O}(\alpha_{EW}^4) \otimes \mathcal{O}(\alpha_s^2)$ at tree level, with α_s being the strong coupling parameter. Example processes are shown in Figure 4.3, amongst them in particular gluon-gluon or quark-anti-quark induced processes. As shown in [10] these processes can be suppressed by topological selection requirements.

VVjj-QCD processes are not included in the study at hand, but their intrinsic production properties lead the motivation for studying the distinct process of like-sign $W^\pm W^\pm$ scattering.

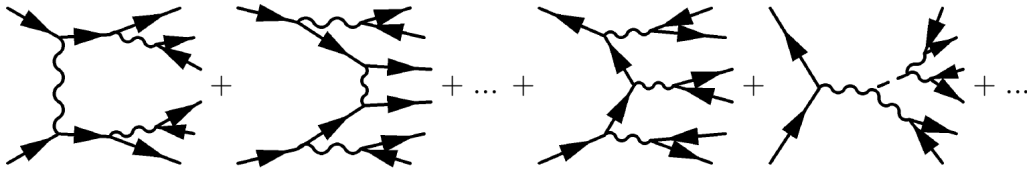


Figure 4.2: Example Feynman diagrams for the order $\mathcal{O}(\alpha_{EW}^6)$ VVjj-EW background processes to VBS at the LHC [10].

4.2 The channel of like-sign $W^\pm W^\pm$ scattering

The scattering of like-sign $W^\pm W^\pm$ features a higher signal-to-background ratio than other VBS processes, which makes this channel well suited for a first look into physics beyond the SM in EWSB. For this particular process, where two $W^\pm$ bosons of same

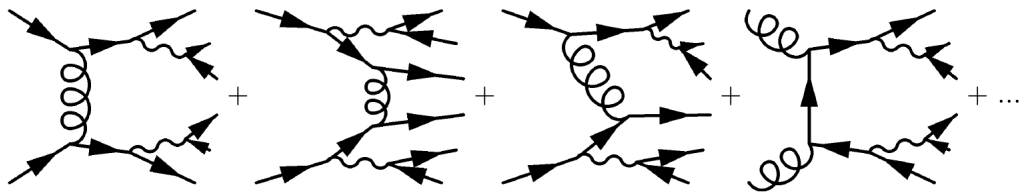


Figure 4.3: Example Feynman diagrams for the order $\mathcal{O}(\alpha_{EW}^4) \otimes \mathcal{O}(\alpha_s^2)$ VVjj-QCD background processes to VBS at the LHC [10].

electric charge are produced, no processes at leading order exist that are induced by two gluons or a quark-anti-quark pair. This is due to electric charge conservation between the initial and the final state. Hence, the QCD contributions to this scattering channel are smaller than in any other scattering channel. At a centre-of-mass energy of 8 TeV the cross sections for $W^\pm W^\pm$ jj-EW and $W^\pm W^\pm$ jj-QCD are of the same order [10] and there is much lower background from $t\bar{t}$ or Z +jets. In addition, no processes with resonant production of any SM particle is possible in the scattering of two like-charge $W^\pm$ bosons.

For like-sign $W^\pm W^\pm$ scattering in the fully leptonic decay channel of the two $W^\pm$ bosons, the process to be studied at the LHC is the following:

$$qq \longrightarrow qql^\pm \nu l^\pm \nu, \quad (4.1)$$

where ν denotes either the neutrino for the W^+W^+ case or the anti-neutrino for W^-W^- scattering. This general notation for the neutrinos in this process will be used throughout this work. The typical VBS event topology is illustrated in Figure 4.4. Following the studied process (4.1), it features two energetic quarks in the forward direction, the initial quarks which radiated the vector bosons. They will propagate as jets² into the detector and are referred to as “tagging jets”. The tagging jets have large separation in pseudorapidity $\Delta\eta$ and large invariant mass m_{jj} . For the double leptonic decay channel, there are additionally two central like-sign leptons and two neutrinos. The two neutrinos almost certainly leave the detector without interaction and therefore cause an imbalance in the transverse momentum, which can be measured.

4.3 Resonances in like-sign $W^\pm W^\pm$ scattering

A most reasonable beginning for the study of resonances in any channel is to investigate observables which are well suited for a discovery. The most suited property for that is the resonance’s mass which is only accessible in s-channel production of the resonance. In the channel of same-sign $W^\pm W^\pm$ only doubly charged resonances can be exchanged in the s-channel and are therefore especially interesting in this study. As can be seen in Table 2.4, the scalar and tensor isospin quintets ϕ and t , respectively, have a doubly charged component. These are the resonances which are discussed further in this thesis.

²The term “jet” is explained in Section 6.1.2.

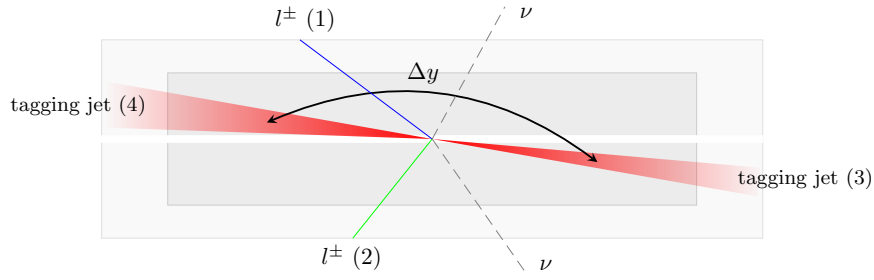


Figure 4.4: Event topology of the VBS process in the channel of $W^\pm W^\pm$ scattering with leptonically decaying vector bosons at the LHC. The scheme shows the two forward jets and the decay products of the $W^\pm$ bosons, two charged leptons and two neutrinos [10].

The topology of resonant VBS is assumed to be very similar to the SM VBS topology. With higher masses of the resonance, however, the outgoing vector bosons tend to be boosted and by that transfer their angular correlation on to their decay products, the leptons and neutrinos.

The presence of two neutrinos makes it impossible to fully reconstruct the invariant mass of the diboson pair m_{WW} . This issue will be addressed with the following study.

5 Monte Carlo event generation

This work presents a study based on Monte Carlo generated events. The BSM scenario with electroweak resonances in like-sign $W^\pm W^\pm$ scattering is the model under investigation. Processes occurring in this model need to be simulated in order to make predictions for the behaviour of such new resonances.

The event simulation needs to be undertaken in the scope of the underlying experiment, the LHC, which provides the hard collisions. Since the LHC is a hadron collider, the event simulation is dominated by QCD effects producing typically hundreds of particles in the final state of a hard scattering process. Those have a momentum range of many orders of magnitude. Perturbation theory provides the mathematical tool to calculate matrix elements and cross sections of a given hard process. However, non-perturbative, soft QCD phenomena, which introduce the issue of colour confinement, are as well involved in the simulation of events. These phenomena need to be modelled due to the absence of a theoretical description.

In this chapter, firstly, the general structure of event generation for the LHC, following the description in [67], is introduced. Afterwards, the simulation of the scattering process of interest, $qq \rightarrow qql^\pm \nu l^\pm \nu$ at parton level with the WHIZARD [58] event generator and the procedure of parton showering with PYTHIA8 [68] are discussed.

5.1 The structure of event generation

The LHC collides protons, which are composite particles built of quarks and gluons, called “partons”. At a centre-of-mass energy of 8 TeV, the collisions studied at hand do not occur at the level of observable particles but instead at much smaller distances at the level of elementary particles, which are the partons inside the proton. This needs to be incorporated in the simulation process. However, although the initial states of the colliding protons are well known (4 TeV of energy each), those of the partons are unknown up to the fact that they carry an arbitrary fraction of the proton’s energy each. Here, parton distribution functions (PDFs) need to be introduced, which give the parton’s longitudinal momentum distribution in the proton as a function of the momentum transfer in the hard process Q and the factorisation scale μ_F at which they are calculated. They are obtained e.g. from Deep Inelastic Scattering experiments. When generating the hard process, the event generator has to pick randomly initial state partons from these distributions [69].

Figure 5.1 sketches the process of event generation. Event generators are built of several components, which describe the hard scattering process at very high scales and continue down to very low scales where hadrons are formed. Due to these huge scale differences the components can be factorised by introducing a factorisation scale, i.e. treated separately in different regimes defined by their scale of momentum transfer.

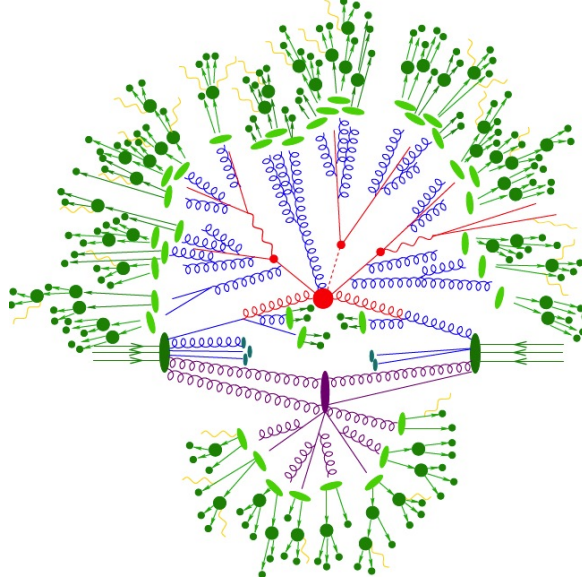


Figure 5.1: Illustration of the structure of Monte Carlo event generation for a pp-collision. In red is the hard scattering process, in blue the parton shower contributions, in green the hadronisation and decays of unstable particles, in violet the underlying event and in orange QED final state radiation [70].

The individual processes for event generation can therefore be treated independently. They are in brief outlined hereafter according to [67].

The hard scattering process The heart of the simulation is the hard scattering process at the level of partons. Matrix elements for all relevant processes involved can be calculated with perturbation theory, mostly at leading order. The fixed order perturbative calculations introduce remaining dependence on the renormalisation and factorisation scale¹. In order to obtain the process' cross section, one needs to integrate over the phase space of all final state particles' momenta in the hard process. This integration has variable dimension, $3n - 4$ for an n -particle final state². Here, the Monte Carlo method comes into play, which is very suitable for such high dimensional integration, since its accuracy is independent of the dimension. With these tools events at “parton level” can be generated.

The QCD parton shower The regime of the hard scattering and the soft QCD phenomena are connected by an evolutionary process from the scale defined by the hard process down to a cut off scale where perturbation theory is no longer applicable due to QCD confinement. The purpose is to simulate additional parton radiation which is not associated to the hard scattering process but a natural consequence of the non-Abelian gauge structure of QCD. Partons, the initial and also final state ones,

¹In a full expansions the dependence on these scales would vanish completely. However, since event generators always perform fixed order calculations the dependence remains and the scales have to be set manually. The choice of these scales is not fixed and relies on experience.

²Three momentum components for each produced particle, diminished by 4 overall energy momentum conservation constraints.

are coloured and emit gluons which themselves emit gluons which can again produce quark-anti-quark pairs. This leads to a “shower of partons”. This phenomenon can be formulated in a sequential Markov process which employs probabilistic decision about the emission of one or more partons to the final state.

Hadronisation In order to get back to a level where observable particles occupy the final state, partons need to be transferred into colourless hadrons. The formation of hadrons arises at the cut off scale of the parton shower, where the scale of momentum transfer is low. Therefore, it cannot be calculated perturbatively and needs to rely on QCD-inspired models with tunable parameters.

Hadron and τ lepton decay During the hadronisation process unstable hadrons can be produced which need to decay further into, at the time-scale of the collider, stable hadrons. This is implemented on the basis of matrix elements for different decay modes. In addition, τ leptons decay in this step of the event simulation, according to their branching ratios circa 65% into hadrons plus τ neutrino, and 35% into leptons plus neutrinos.

Soft QCD contributions and underlying event Since most of the collisions at the LHC are of soft nature, a model for these processes need to be included in the event generation. Related to this is the “underlying event”, the remaining components of the protons which participated in the hard interaction. These are also colour charged partons which need to be connected to the colour flow of the parton shower and the final state particles from the hard scattering. It can also occur that these partons have an additional interaction to the hard process. This phenomenon is called “multiple parton interaction”. Here as well only models are available.

QED radiation Finally the simulation of final state photon radiation of electrically charged particles is included in the full particle-level event simulations. This photon radiation is referred to as “QED radiation”.

All these phenomena together form the tools of various particle-level event generators. A more detailed description of these tools can be found in [67]. In the presented studies, the WHIZARD event generator [58] provides the “parton-level” events for all used samples and is then interfaced to the PYTHIA8 event generator [68] in order to simulate all remaining phenomena. The result are “particle-level” events.

In order to compare Monte Carlo generated events to real ATLAS data, a detector simulation needs to follow. It simulates the interaction of the produced particles with the detector matter and the electronic response of the detector. In the scope of this work, however, a different path is taken with the manual implementation of a first step towards a full detector simulation, so called “missing transverse momentum smearing” (refer to Section 6.3).

5.2 Parton-level event generation

5.2.1 The event generator WHIZARD

WHIZARD [58] is currently the only generator that is able to generate heavy, K-matrix unitarised resonances in VBS. Therefore, it is the event generator of choice to simulate the process under investigation $qq \rightarrow qql^\pm \nu l^\pm \nu$. It contains the full matrix element generator O'MEGA [71], which is able to compute complete tree-level matrix elements for collider processes including up to 8 final state particles in the hard scattering. This is a particular feature of WHIZARD.

5.2.2 Simulation of the process $qq \rightarrow qql^\pm \nu l^\pm \nu$

For the simulation and event generation of resonant scattering of vector bosons, the SM extending effective field theory approach, introduced in Section 2.3, is implemented in WHIZARD. This model (“SM_km”³) comprises K-matrix unitarisation and the new resonances in the EW sector discussed in Section 2.3.

At leading order matrix element level, the process $qq \rightarrow qql^\pm \nu l^\pm \nu$ is a two-to-six process. This can directly be calculated with WHIZARD by considering all possible tree level processes which can be built from any connection of the initial and the final state particles. This procedure, in particular, has the advantage that the W bosons are not strictly required to be produced on their mass shell and that it includes all correlations between the $W^\pm$, which is especially important for any diboson invariant mass reconstruction. In addition, the generated processes also include irreducible background to the resonant VBS.

Standard Model electroweak $qq \rightarrow qql^\pm \nu l^\pm \nu$

As all samples, the SM-EW $qq \rightarrow qql^\pm \nu l^\pm \nu$ process is generated at leading order for proton-proton collisions with a centre-of-mass energy of 8 TeV. Only electroweak contributing processes are included in these studies by setting α_s to zero. The quark flavour summation in the initial and final state only contains the light quarks $q \in \{u, d, s, c\}$, whose masses are set to zero. The heavy bottom- and top-quark have a small contribution due to their small occurrence within the proton and are therefore neglected. Besides the τ lepton, all leptons are considered to be massless. All possible leptonic decays for the W bosons are included, leading to a total of six final state configurations at parton level, $e^\pm e^\pm$, $e^\pm \mu^\pm$, $\mu^\pm \mu^\pm$, $e^\pm \tau^\pm$, $\mu^\pm \tau^\pm$, $\tau^\pm \tau^\pm$. However, τ leptons will decay in the parton shower and are therefore not considered in the analysis. Hence, the notation “lepton” will hereafter correspond to electrons and muons as well as their anti-particles positrons and anti-muons only. The recently discovered Higgs boson with a mass of 125 GeV is implemented as the SM Higgs boson with the corresponding width and couplings. All additional parameters for the event generation are summarised in Table 5.1.

In order to limit the event generation towards the final event selection requirements, loose selection criteria on the final state particles were already imposed at generator level. These selection criteria are listed in Table 5.2.

³This model name corresponds to the WHIZARD version 2.1.1

	[GeV]		[GeV]		
m_W	80.399	m_H	126.0	α_s	0.0
Γ_W	2.085	Γ_H	0.00418	$\sqrt{s}$	8 TeV
m_Z	91.1876	m_τ	1.77705		
Γ_Z	2.4952	m_{top}	172.5		

Table 5.1: List of additional parameters for the event generation with WHIZARD. Masses of the light quarks u, d, s, c and all leptons, besides the τ lepton mass, are set to zero.

$\Delta R(q, q)$	> 0.4
$\eta(q)$	$< 5.0 $
$p_T(q)$	$> 15 \text{ GeV}$
$\eta(l)$	$< 5.0 $
$p_T(l)$	$> 8 \text{ GeV}$

Table 5.2: List of generation phase space selection cuts for the event generation with WHIZARD. “q” denotes all occurring quark flavours $\{u, d, s, c\}$, and “l” all leptons $\{e^\pm, \mu^\pm\}$ besides the $\tau^\pm$.

Parton density functions (PDF) are used to describe the probability to find a parton with a certain proton’s energy fraction inside the proton. In these studies the PDF set “cteq6ll.LHpdf” from the LHAPDF library [72] is used. These PDF depend on the momentum transfer of the hard interaction and the renormalisation and factorisation scale. Both scales are set simultaneously and dynamically to the invariant mass of all leptons in the event, $m_{l\nu l\nu}$. K-matrix is switched on already for the SM event simulation, although this would not have been necessary and has no influence on the generated events [73].

Resonant $qq \rightarrow qql^\pm \nu l^\pm \nu$

For the generation of the resonant $qq \rightarrow qql^\pm \nu l^\pm \nu$ all previous parameters and settings are maintained. Additionally, one resonance at a time is switched on by setting its mass m , width Γ and coupling g to longitudinal vector bosons. Including the electroweak resonances in the simulation changes the $W^\pm W^\pm \rightarrow W^\pm W^\pm$ vertex. Here, the K-matrix unitarisation becomes important as it keeps the $W^\pm W^\pm \rightarrow W^\pm W^\pm$ amplitude physically reasonable.

For each possible resonance type in $W^\pm W^\pm$ scattering, ϕ and t , several mass points and couplings are chosen. One set of samples varies the mass of the resonance $m \in \{500 \text{ GeV}, 600 \text{ GeV}, 750 \text{ GeV}, 800 \text{ GeV}, 1000 \text{ GeV}, 1200 \text{ GeV}\}$ while fixing the coupling to $g = 2.5$. A second set fixes the mass to $m = 500 \text{ GeV}$ and varies the couplings $g \in \{0.5, 1.0, 1.5, 2.0\}$. This gives a total of 10 samples per resonance type plus the SM case without any resonance switched on. All samples used in the analysis are listed in Table 5.3 specifying the properties of the resonances as well as the cross

section for a centre-of-mass energy $\sqrt{s} = 8 \text{ TeV}$. This cross section corresponds to the leading order WHIZARD phase space integration with the consideration of only the four lightest quark flavours and the implemented generator level selection criteria in Table 5.2. Statistical uncertainties calculated by WHIZARD are well below 1 % and are, therefore, not listed separately.

resonance	mass m [GeV]	coupling g	width Γ [GeV]	cross section [fb] (at $\sqrt{s} = 8 \text{ TeV}$)
ϕ	500	2.5	64.21	34.02
	600	2.5	110.95	26.72
	750	2.5	216.70	20.48
	800	2.5	263.00	19.02
	1000	2.5	513.66	16.78
	1200	2.5	887.61	15.69
ϕ	500	0.5	2.57	15.27
	500	1.0	10.27	17.72
	500	1.5	23.11	20.77
	500	2.0	41.09	27.88
t	500	2.5	2.14	21.78
	600	2.5	3.70	17.87
	750	2.5	7.22	15.99
	800	2.5	8.77	15.96
	1000	2.5	17.12	15.28
	1200	2.5	29.59	15.02
t	500	0.5	0.09	14.89
	500	1.0	0.34	15.35
	500	1.5	0.77	16.84
	500	2.0	1.37	18.18
SM				14.78

Table 5.3: List of Monte Carlo samples generated with the WHIZARD event generator with K-matrix unitarisation. The mass and the coupling of the resonances are free parameters, whereas the width is calculated using Equation (2.41). Cross sections are given at generator level before parton shower for a centre-of-mass energy of 8 TeV and the generator level selection requirements in Table 5.2. “SM” refers to the SM-EW $l^\pm \nu l^\pm \nu jj$ sample without any resonance.

6 Event preparation

This chapter comprises all reconstruction and selection steps which are required in order to use the generated events for further analysis. At first, physics objects need to be reconstructed. Then, suitable selection requirements for them are imposed. The topology of VBS events includes leptons, jets and neutrinos. The reconstruction and selection procedure of these objects is outlined in the first section of this chapter. Afterwards, the entire event topology needs to be reconstructed where selection criteria are introduced which promote the VBS event structure and especially permit the reconstruction of the invariant mass of the $W^\pm W^\pm$ system for resonant scattering processes. Finally, the method of “missing transverse momentum smearing” is introduced which is used to simulate a first step of the full ATLAS detector simulation.

6.1 Object reconstruction and selection

In this section the reconstruction of the different objects used in the analysis is outlined. These reconstruction instructions are used independently of the level of event simulation and reconstruction introduced at the end of Section 5.1. The selection criteria for transverse momentum and rapidity ranges are based on those at detector level for a full analysis of ATLAS data which are motivated by the detector layout and energy and momentum measurement limitations outlined in Section 3.2.

According to the event topology and the chosen final state $qq \rightarrow qql^\pm \nu l^\pm \nu$ (refer to Section 4.2), the reconstruction of jets, charged leptons and neutrinos is required.

6.1.1 Dressed leptons

In all generated MC samples, $\tau^\pm$ leptons are included in the generation of the final state $qq \rightarrow qql^\pm \nu l^\pm \nu$. In the parton shower simulation with PYTHIA8, $\tau^\pm$ leptons decay either leptonically or hadronically. For this analysis only electrons and muons are considered as final state leptons. This means, at parton level τ leptons are not included, while from particle level on, they are included in the analysis if they decay leptonically and leave an observable electron or muon in the event.

For the reconstruction of leptons, the method of “photon dressing” is used. Dressed leptons are obtained by collecting all final state photons in a cone of $\Delta R = 0.1$ around the direction of the undressed lepton. The four-momenta of these photons are added to the lepton’s four momentum finally building the dressed lepton’s momentum. This is necessary in order to make the reconstruction less dependent on the implementation of collinear final state QED radiation which can differ between the various MC event generators. Of course, this does not change anything for parton-level leptons since there QED radiation is not performed yet.

Requirements are placed on the transverse momentum and the rapidity of the dressed leptons in order to be considered for the analysis. These are $p_T > 30 \text{ GeV}$ and $|\eta| < 2.5$, respectively.

6.1.2 Jets

Fundamental particle reactions in the quark sector are treated in terms of quarks and gluons. However, QCD phenomena make it impossible to observe single partons in nature due to color confinement. Immediately after their production, high-energy quarks and gluons fragment and cluster to form hadrons which results in a collimated spray of hadrons, referred to as “jet”. Hence, not the initial final state partons are observable in the experiment but the jets resulting from them. The idea of parton reconstruction is therefore, to find a robust definition which relates the theoretical predictions of partons to the in experiments measurable quantities of jets. Different jet definitions were developed for this purpose, constituting a jet algorithm, a set of rules for grouping particles into jets, and their parameters [74].

A widely used jet algorithm in LHC physics and used in this analysis is the anti- k_t algorithm [75]. It is a sequential recombination jet algorithm, which performs an iterative four-momentum recombination in the $y - \phi$ space. It takes all objects in an event (either partons, particles or at detector level calorimeter clusters) and associates two distance measures to each of them

$$\rho_{ij} = \min(p_{Ti}^{2p}, p_{Tj}^{2p}) \frac{(\Delta R_{ij})}{R}, \quad (6.1)$$

$$\rho_{iB} = p_{Ti}^{2p}. \quad (6.2)$$

p_{Ti} is the transverse momentum of object i , ΔR_{ij} is the distance in $y - \phi$ space between object i and j . R is the distance parameter of the jet, i.e. the spread in $y - \phi$ space. The first quantity, ρ_{ij} , is a p_T -dependent distance measure between two objects, the second, ρ_{iB} , is the distance between object i and the beam line. If $\rho_{iB} > \rho_{ij}$, then a new object with the momentum $p_{Ti} + p_{Tj}$ replaces the objects i and j . If $\rho_{iB} < \rho_{ij}$, then object i is added to the list of final jets. For the anti- k_t algorithm, the exponent p equals -1 which means that objects with largest transverse momentum are clustered first. This provides stable and robust jets which are primarily defined by the highest- p_T component [76].

The jets used in this analysis are reconstructed with the FASTJET [77] implementation and have a distance parameter $R = 0.4$. They have to fulfil the transverse momentum and pseudorapidity requirements $p_T > 30 \text{ GeV}$ and $|\eta| < 4.5$, respectively.

Due to the demand of universality of the jet algorithms with respect to the event generation level, the composition of jets in the event should be stable through all levels considered. This has to hold, although, the input for the jet reconstruction is quite different at parton and particle level, due to the effects of the parton shower.

6.1.3 Neutrinos

The ATLAS detector is unable to measure any quantity of neutrinos directly. The reason is that neutrinos solely interact weakly whereas most detection mechanism rely

on electromagnetic or strong interactions due to their considerably larger cross sections in the energy range of detection. Nevertheless, it is possible to trace the existence of neutrinos in an event through the consideration of momentum conservation. The initial colliding partons in an event are supposed to have exclusively longitudinal momentum. Hence, momentum conservation in the transverse plane imposes zero momentum in this plane after the collision. A measurement of an imbalance of the transverse momentum is, therefore, considered as energy which is carried away by undetected neutrinos. This quantity of transverse momentum imbalance is referred to as missing transverse energy¹ $\vec{E}_T^{\text{miss}}$. The missing transverse energy in an event is calculated as the negative sum of all object's momenta in x and y direction separately in the range of $|\eta| < 4.5$ in the event, besides, of course, those of the neutrinos

$$p_{x(y)}^{\text{miss}} = - \left(\sum_i p_{x(y)}^{\text{jet}_i} + \sum_i p_{x(y)}^{\text{lep}_i} \right). \quad (6.3)$$

The transverse component is considered only, since the incoming parton's energy fractions in beam direction are unknown and, therefore, no longitudinal momentum conservation constraint can be imposed.

Only in an ideal and hermetic detector neutrinos are the only sources of missing transverse energy. However, it is impossible to build a detector which covers the full $y - \phi$ space around the collision point. The consequence is mis-measurement of the amount of transverse momentum. In the scope of these studies no detector simulation was conducted. Nevertheless, the impact of detector resolution on the measurement of $\vec{E}_T^{\text{miss}}$ is studied. This is done via the smearing of the missing transverse energy through a simulation of the energy resolution of the ATLAS detector (Section 6.3).

6.2 Event selection

The event selection is motivated by supporting the VBS topology over any possible background and the guarantee that the reconstruction of the invariant mass of the $W^\pm W^\pm$ system is made possible. The demand of leaving a maximum number of events for the mass reconstruction studies results in only the few necessary selection requirements which are imposed on every event:

- exactly two charged dressed leptons,
- a minimum of two jets,
- both leptons have the same electric charge,
- both leptons have to have a minimum distance of $\Delta R \geq 0.3$ to any reconstructed jet in the event and to each other.

The two jets with the highest transverse momentum are referred to as “tagging jets”. All combinations of leptonic decays, besides $\tau^\pm$ leptons, of the two $W^\pm$ bosons are treated equally.

¹Despite $\vec{E}_T^{\text{miss}}$ is actually a momentum, the notation of an “energy” has been established due to the measurement strategy of $\vec{E}_T^{\text{miss}}$ in ATLAS.

6.3 Missing transverse momentum smearing

A first step of the event generation is the simulation of the hard scattering process at parton level with WHIZARD. Further on all relevant effects, to obtain observable physics objects at the particle level, are added by the PYTHIA8 event generator. On top of these simulations, a first step towards a real measurement at detector level is taken with the simulation of detector resolution effects. In the context of this work, the kinematical information of the relevant final state particles, the decay products of both W bosons, plays a crucial role. In a real physics experiment at the ATLAS detector, the kinematic properties of charged leptons are obtained in the tracking detectors. Neutrinos are traceable by the existence of missing transverse energy. The reconstruction of $\vec{E}_T^{\text{miss}}$ at detector level is performed using energy deposits in the calorimeters and reconstructed tracks of muons in the Muon Spectrometer [78].

Referring to the energy and momentum resolution performance of the ATLAS detector, in Section 3.2.3, Equation (3.9) and Section 3.2.2, Equation (3.8), respectively, the energy resolution of the calorimeters is the dominant performance limitation. Hence, as a beginning, the simulation of the resolution of missing transverse momentum is conducted. This is referred to in the following as “ $\vec{E}_T^{\text{miss}}$ smearing”.

The smearing of $\vec{E}_T^{\text{miss}}$ is based on the ATLAS performance of missing energy reconstruction [78]. It is implemented separately for the magnitude E_T^{miss} and the azimuthal direction ϕ^{miss} of the missing transverse momentum. In [78], a study of the magnitude of E_T^{miss} resolution is conducted using the ratio

$$R = \frac{\text{RMS} \left(\frac{E_T^{\text{miss}}(\text{reco})}{E_T^{\text{miss}}(\text{true})} \right)}{\left\langle \frac{E_T^{\text{miss}}(\text{reco})}{E_T^{\text{miss}}(\text{true})} \right\rangle}, \quad (6.4)$$

where “RMS” is short for the root mean squared of the event distribution in n bins in the given variable:

$$\text{RMS} \left(\frac{E_T^{\text{miss}}(\text{reco})}{E_T^{\text{miss}}(\text{true})} \right) = \sqrt{\frac{\sum_j^n \left(\frac{E_T^{\text{miss}}(\text{reco})}{E_T^{\text{miss}}(\text{true})} \right)_j^2}{n}} \quad (6.5)$$

and $\left\langle \frac{E_T^{\text{miss}}(\text{reco})}{E_T^{\text{miss}}(\text{true})} \right\rangle$ denotes the mean value of the $\frac{E_T^{\text{miss}}(\text{reco})}{E_T^{\text{miss}}(\text{true})}$ distribution. Figure 6.1a shows the ratio R in dependency of the magnitude of the true missing transverse momentum $E_T^{\text{miss}}(\text{true})$ as obtained in ATLAS [78]. The performance of the E_T^{miss} measurement depends on the event topology. It is shown for $t\bar{t}$ events in this figure. The $t\bar{t}$ event topology is chosen from the available studies in [78] due to the most similar event topology to $W^\pm W^\pm$ scattering with two $W^\pm$ bosons. To extract the relative resolution of the E_T^{miss} measurement from this ratio, some assumptions need to be made. Taking the histogram in Figure 6.1a, the $E_T^{\text{miss}}(\text{true})$ is approximately constant within each bin $i \in \{1, n\}$ of the $E_T^{\text{miss}}(\text{true})$ distribution and $\text{RMS}(E_T^{\text{miss}}(\text{true})) = 0$. This leads to

$$\text{RMS} \left(\frac{E_T^{\text{miss}}(\text{reco})}{E_T^{\text{miss}}(\text{true})} \right)_i = \frac{\text{RMS} (E_T^{\text{miss}}(\text{reco}))_i}{E_T^{\text{miss}}(\text{true})_i} \quad (6.6)$$

and

$$\left\langle \frac{E_T^{\text{miss}}(\text{reco})}{E_T^{\text{miss}}(\text{true})} \right\rangle_i = \frac{\langle E_T^{\text{miss}}(\text{reco}) \rangle_i}{E_T^{\text{miss}}(\text{true})_i}. \quad (6.7)$$

Assuming that the mean of the E_T^{miss} distribution in each bin should stay unaffected by the $\vec{E}_T^{\text{miss}}$ smearing

$$\langle E_T^{\text{miss}}(\text{reco}) \rangle_i = E_T^{\text{miss}}(\text{true})_i \quad (6.8)$$

leads to the result

$$R_i = \frac{\text{RMS}(E_T^{\text{miss}}(\text{reco}))_i}{\langle E_T^{\text{miss}}(\text{true}) \rangle_i}. \quad (6.9)$$

This quantity is the relative resolution in each $E_T^{\text{miss}}(\text{true})$ -bin of the magnitude of the missing transverse momentum. From Figure 6.1a the ratio R is approximated with two linear and a constant function for different intervals of $E_T^{\text{miss}}(\text{true})$

$$\begin{aligned} E_T^{\text{miss}}(\text{true}) < 95 \text{ GeV} : \quad R &= -3.56 \cdot 10^{-3} \frac{1}{\text{GeV}} \cdot E_T^{\text{miss}}(\text{true}) + 0.59, \\ E_T^{\text{miss}}(\text{true}) \in (95, 200) \text{ GeV} : \quad R &= -0.74 \cdot 10^{-3} \frac{1}{\text{GeV}} \cdot E_T^{\text{miss}}(\text{true}) + 0.29, \\ E_T^{\text{miss}}(\text{true}) \geq 200 \text{ GeV} : \quad R &= 0.15. \end{aligned} \quad (6.10)$$

These equations are used as the input for the smearing of the true $\vec{E}_T^{\text{miss}}$, which is imitated by a Gaussian probability density function of mean $E_T^{\text{miss}}(\text{true})$ and width $R \cdot E_T^{\text{miss}}(\text{true})$. If the magnitude drops below zero after smearing, the direction of $\vec{E}_T^{\text{miss}}$ is turned by 180° , i.e. π is added to the azimuthal angle ϕ^{miss} , and the

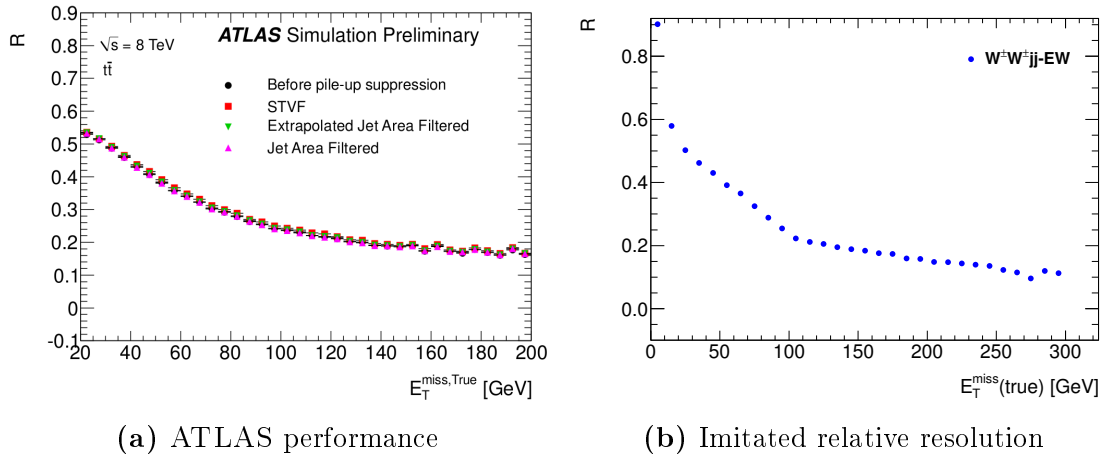


Figure 6.1: Distributions of relative resolution ratio R as defined in Equation (6.4) as a function of the true E_T^{miss} for $t\bar{t}$ events in ATLAS ([78], Figure 14(d)) (left) and for the reconstruction in this study as relative resolution input for the smearing of the magnitude of the missing transverse momentum (right). As a reference for the right plot, the “Pile-up suppression Extrapolated Jet Area Filtered” graph (green) from the left plot is taken. The notations $E_T^{\text{miss, true}}$ and $E_T^{\text{miss}}(\text{true})$ refer to the same quantity.

positive value of the smeared E_T^{miss} is taken. The resulting reconstructed distribution of the ratio R as a function of $E_T^{\text{miss}}(\text{true})$ is displayed in Figure 6.1b which shows a good agreement with the input used (Figure 6.1a). In fact these graphs recover the behaviour of the energy resolution in the ATLAS detector, Equation (3.9). For small energies the $1/\sqrt{E}$ term is dominant, whereas for higher energies a constant behaviour is observable. The resulting reconstructed distribution of the ratio $(E_T^{\text{miss}}(\text{reco}) - E_T^{\text{miss}}(\text{true}))/E_T^{\text{miss}}(\text{true})$ is shown in Figure 6.2. The Gaussian distribution is clearly visible as well as the observation of no entries below -1 since these values would result in negative reconstructed E_T^{miss} values.

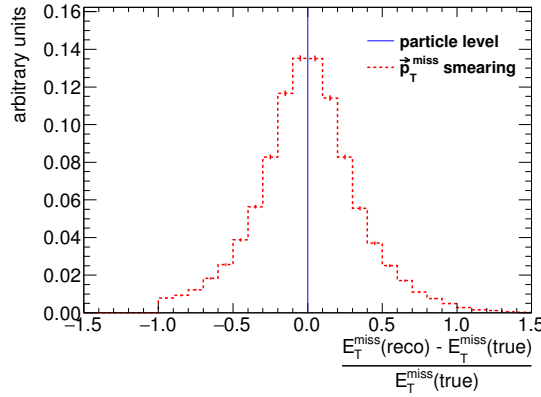


Figure 6.2: Reconstruction of the relative resolution $(E_T^{\text{miss}}(\text{reco}) - E_T^{\text{miss}}(\text{true}))/E_T^{\text{miss}}(\text{true})$ of the magnitude of $\vec{E}_T^{\text{miss}}$ according to Figure 6.1 and Equations (6.10).

The smearing of the azimuthal angle ϕ^{miss} is handled in a similar way as for the magnitude. Figure 6.3a shows the width of the difference of the reconstructed and true azimuthal angle of $\vec{E}_T^{\text{miss}}$ as a function of the $E_T^{\text{miss}}(\text{true})$ for $W^\pm \rightarrow e^\pm \nu$ events² in ATLAS [78]. This figure is used to obtain the absolute resolution for $\Delta\phi^{\text{miss}} = \phi^{\text{miss}}(\text{reco}) - \phi^{\text{miss}}(\text{true})$ in dependence of $E_T^{\text{miss}}(\text{true})$. Again, the histogram is approximated by linear and constant functions to reconstruct the resolution for different intervals of $E_T^{\text{miss}}(\text{true})$

$$\begin{aligned}
 E_T^{\text{miss}}(\text{true}) < 20 \text{ GeV} : & \quad \sigma(\Delta\phi^{\text{miss}}) = 0.8 \text{ rad}, \\
 E_T^{\text{miss}}(\text{true}) \in [20, 30) \text{ GeV} : & \quad \sigma(\Delta\phi^{\text{miss}}) = 0.55 \text{ rad}, \\
 E_T^{\text{miss}}(\text{true}) \in [30, 75) \text{ GeV} : & \quad \sigma(\Delta\phi^{\text{miss}}) = -5 \cdot 10^{-3} \frac{\text{rad}}{\text{GeV}} \cdot E_T^{\text{miss}}(\text{true}) + 0.6 \text{ rad}, \\
 E_T^{\text{miss}}(\text{true}) \in [75, 200) \text{ GeV} : & \quad \sigma(\Delta\phi^{\text{miss}}) = -0.57 \cdot 10^{-3} \frac{\text{rad}}{\text{GeV}} \cdot E_T^{\text{miss}}(\text{true}) + 0.2 \text{ rad}, \\
 E_T^{\text{miss}}(\text{true}) \geq 200 \text{ GeV} : & \quad \sigma(\Delta\phi^{\text{miss}}) = 0.126 \text{ rad}.
 \end{aligned} \tag{6.11}$$

The Gaussian resolution function for the smearing of the azimuthal angle is constructed with the mean of zero and the width $\Delta\phi^{\text{miss}}$, where the assumption was made that the smearing preserves the symmetry in the azimuthal angle and, therefore, leaves

²This figure was not available for $t\bar{t}$ the event topology in [78].

its mean unaffected. The resulting reconstructed distribution of the absolute resolution of $\Delta\phi^{\text{miss}}$ as a function of $E_T^{\text{miss}}(\text{true})$ is depicted in Figure 6.3b. Here as well the validity of the procedure can be concluded. The result of the smearing of $\Delta\phi^{\text{miss}}$ is illustrated in Figure 6.4 which shows the expected Gaussian behaviour around a mean of zero.

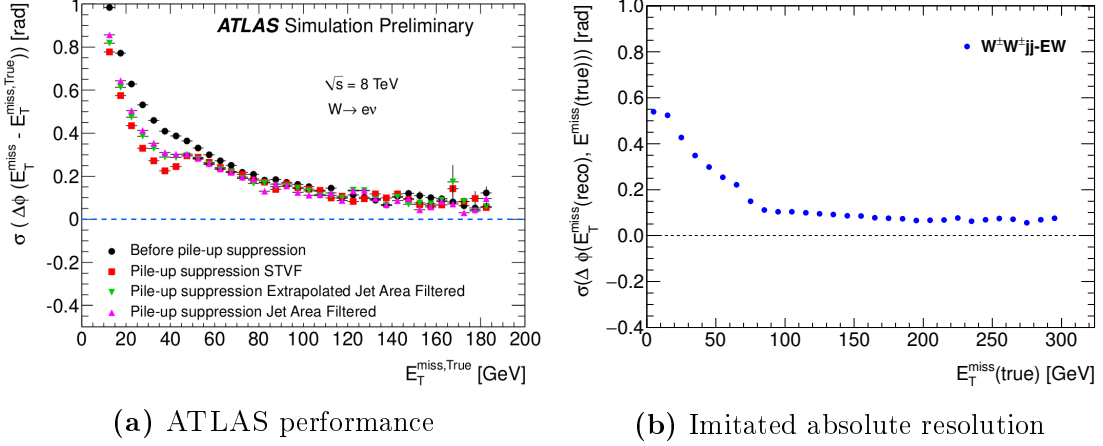


Figure 6.3: Distributions of absolute resolution of the difference between reconstructed and true transversal direction $\Delta\phi^{\text{miss}}$ as a function of the true E_T^{miss} for $W^\pm \rightarrow e^\pm \nu$ events in ATLAS ([78], Figure 26(a)) (left) and the reconstruction of this study as absolute resolution input for the smearing of the transversal direction of the missing transverse momentum (right).

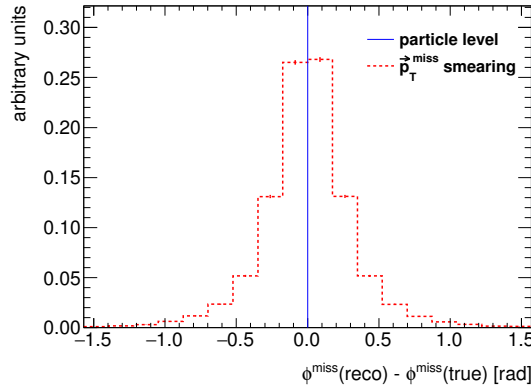


Figure 6.4: Reconstruction of the absolute resolution of $\Delta\phi^{\text{miss}}$ of the transversal direction of $\vec{E}_T^{\text{miss}}$ according to Figure 6.3 and Equations (6.11).

The result of the smearing on the magnitude and the azimuthal angle of $\vec{E}_T^{\text{miss}}$ is depicted in Figure 6.5. As expected, the distribution of the magnitude E_T^{miss} becomes broader with respect to the distribution before smearing at particle level. The mean is slightly shifted towards smaller values. This effect is due to events where the smearing of E_T^{miss} results in small but negative magnitudes which are then flipped back to small, positive values with opposite direction. This leads to a slightly asymmetric smearing result for low magnitudes of $\vec{E}_T^{\text{miss}}$. The distribution of the azimuthal angle of $\vec{E}_T^{\text{miss}}$ remains flat which is expected for a random smearing.

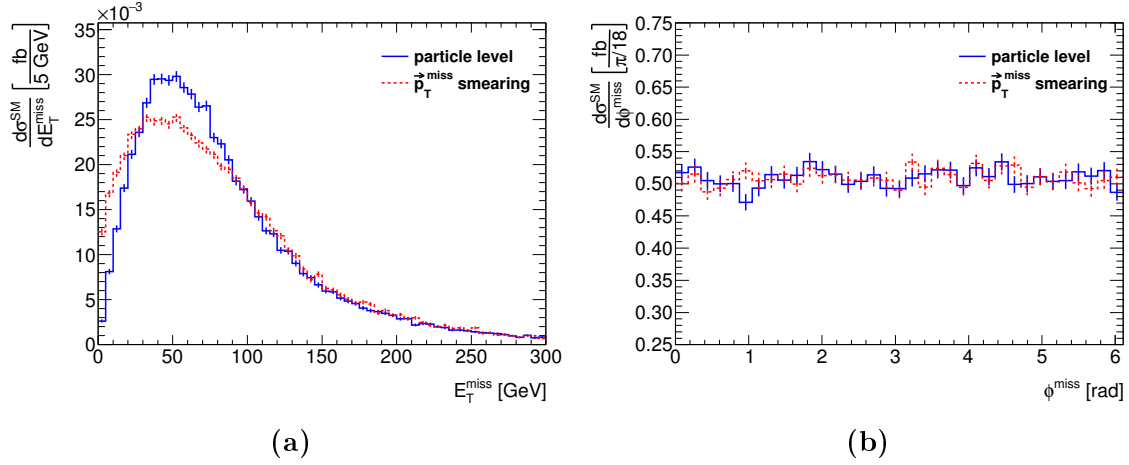


Figure 6.5: Comparison of smeared (red, dashed) and unsmeared (blue, solid) magnitude E_T^{miss} (left) and transversal direction ϕ^{miss} (right) of $\vec{E}_T^{\text{miss}}$.

6.4 Event reconstruction levels

In summary, three event reconstruction levels are used in this analysis which are hereafter referred to as

- “parton level” — pure hard scattering events obtained from WHIZARD, in literature sometimes referred to as “generator truth level”,
- “particle level” — parton-level events interfaced to PYTHIA8 for simulation of all remaining physics processes after the hard scattering simulation, which are not related to detector simulation,
- “ $\vec{E}_T^{\text{miss}}$ smearing” — “particle-level” events with simulation of “missing transverse momentum smearing”.

Notice that all three event reconstruction levels cannot be “detected” within an experiment. A particle detector has always an intrinsic detection resolution which in addition to the pure particle physics process enters in the result of a measurement. The missing transverse energy is smeared with detector resolution effects in the $\vec{E}_T^{\text{miss}}$ smearing event reconstruction level, however, this includes only a small part of the resolution effects. Hence, a comparison with real ATLAS data cannot be provided. However, the underlying hypothesised physics can be studied which is the main goal of this work.

7 Mass reconstruction of resonances in $W^\pm W^\pm$ scattering

The intention of this thesis is the reconstruction of the mass of heavy resonances in $W^\pm W^\pm$ scattering and the study of its properties and performance. Two resonance types are possible for this process. Yet, the general technique of “re-constructing” the resonance’s mass is independent on the type of resonance. The reason for this is that these resonances have rather identical event topologies, which is the crucial point for mass reconstruction. The event topology for the process at hand has been discussed briefly in Section 4.3 and shall, at the beginning of this chapter, be recapitulated in the context of reconstructing the resonance’s mass from the final decay products, which can – only partly – be accessed in measurement.

Taking this as a basis, mass reconstruction variables are constructed in the subsequent sections. These variables have been divided into two classes. One selection is developed by following a concrete prescription, the “technique of constrained minimisation”, leading to “mass bound variables” [17]. The second selection agglomerates some subset of traditionally used variables, which in most cases have been introduced in a rather “ad-hoc” manner to serve a concrete case of mass reconstruction.

7.1 The problem statement

The reconstruction of the mass of a resonance, ϕ or t , denoted universally $X^{\pm\pm}$, of any mass or coupling parameter shall be investigated commonly. The fully leptonic decay via two $W^\pm$ bosons

$$X^{\pm\pm} \longrightarrow W^\pm W^\pm \longrightarrow l^\pm l^\pm \nu \nu \quad (7.1)$$

has four final state particles, two equally charged leptons and two (anti-)neutrinos. The resonance mass can be obtained, from the relativistic relation of mass and energy, by forming the invariant mass of all final decay products

$$m_X^2 = m_{WW}^2 = P_{WW}^\mu P_{\mu WW} = E^2 - |\vec{p}|^2, \quad (7.2)$$

where P_{WW}^μ denotes the composite four-momentum of the $W^\pm W^\pm$ system

$$P_{WW}^\mu = P_{l_1}^\mu + P_{l_2}^\mu + P_{\nu_1}^\mu + P_{\nu_2}^\mu. \quad (7.3)$$

The momenta of the charged leptons, any combination of electrons and muons, are well measurable in the detector. However, neutrinos leave without any interaction. The only chance to retain information about the neutrinos’ momenta is to consider the measured missing transverse energy E_T^{miss} . Assuming that the measured missing

transverse momentum totally stems from the two decay neutrinos is the only reasonable way to identify them. Nevertheless, the separate momenta of both neutrinos $\vec{q}_{\nu_1}$, $\vec{q}_{\nu_2}$ remain unknown. With the assumption of massless leptons, both neutrinos together have 6 degrees of freedom, while the measurement of $\vec{E}_T^{\text{miss}}$ imposes only two constraints on them

$$\vec{E}_T^{\text{miss}} = \vec{q}_{\nu_1 T} + \vec{q}_{\nu_2 T}. \quad (7.4)$$

Hence, the event information obtained in measurements is incomplete and reconstruction of the invariant mass of the resonance suffers from ambiguities. Of course, in a real experiment there are always other sources of missing transverse momentum as discussed in Section 6.1.3. The impact on miss-measurement of $\vec{E}_T^{\text{miss}}$ is simulated in Section 6.3 and studied in a later stage of this work.

Obviously, the problem of insufficient event information for the reconstruction of heavy particles' masses did not just appear recently. The transverse mass variable m_T [79, 80] can be traced back to the discovery of the $W^\pm$ boson [31, 33] and was commonly embedded in the ATLAS and CMS design reports [81, 82]. It is now an inherent part of analyses with the issue of W +jets and $t\bar{t}$ background suppression [83, 84]. In recent years, the reconstruction of more complex topologies with more than one invisible particle became eminent. Numerous generalised approaches for transverse mass reconstructions were developed, especially in the context of searches for the Higgs boson in the W^+W^- [11, 85–88] and $\tau^+\tau^-$ [12, 13, 89] decay channels, as well as for supersymmetric particles and Dark matter candidates [14–16, 90–93]. In 2011, most of the effort of previous studies was collected, summarised and arranged in the descriptive pattern of “mass bound variables” in [17].

In the frame of resonances in VBS the generic transverse mass was investigated [69]. This former study seeks to be expanded to examine the properties and discovery potential of mass bound variables as well as “traditional” variables such as the “visible mass” m_{vis} , the “effective mass” m_{eff} [94], the “vectorial mass” m_{vec} and the “collinear mass” m_{col} [95].

7.2 Mass bound variables

Mass bound variables are a result of a well-motivated kinematical approach. It combines the formation of transverse projected mass variables with applying the technique of constrained minimisation in decay topologies with invisible particles. This class of mass reconstruction variables was introduced in [17]. They, by construction, inherit the feature of constraining the resonance mass from below, if applied to events with the correct event topology. Therefore, they are most useful when applied to cases where the mass regime of BSM physics is higher than the energy scale of the most important backgrounds.

These variables can be obtained via a concrete prescription. For the specified case of reconstructing *one* resonance in the event, the steps that need to be taken are the following:

1. Concretise the decay topology under investigation,
2. Form the invariant mass $\mathcal{M}_1$ corresponding to this topology by choosing
 - the transverse projection
 - the order of operations $\in \{\text{formation of composite particles, transverse projection}\},$
3. Concretise constraints on mass reconstruction,
4. Minimise over remaining free parameters¹.

The first step has already been accomplished by formulating the problem statement. The second step is the one with the most ambiguities and choices, leading to a number of different transverse mass bounds, each appropriate by itself but with different properties. The choices of transverse projection and the order of operations are discussed in the two following sections. The third step is fully adjustable to the case at hand. Different constraints which are appropriate for the decay topology $X^{\pm\pm} \rightarrow W^{\pm}W^{\pm} \rightarrow l^{\pm}l^{\pm}\nu\nu$ are discussed. Finally, the constraining character of the resulting mass bound variable is obtained by minimising over the space of free parameters.

7.2.1 Transverse projections

Forming transverse masses is reasonable with the application at hadron colliders. This makes the resulting variables less dependent on the unknown longitudinal momenta of the incoming partons which is due to the limited detector acceptance in the direction of the beam and the PDFs of the proton. In fact, mass bound variables need to be intrinsically independent under longitudinal boost since the bound should only depend on the resonance's mass.

For the calculation of transverse masses, a four-momentum has to be “projected” into the plane transverse to the beam direction. The result is a transverse projected (1+2)-momentum

$$p_{T\alpha} = (e_T, p_x, p_y) \quad (7.5)$$

analogous to the known Lorentz four-momentum

$$P_{\mu} = (E, p_x, p_y, p_z). \quad (7.6)$$

The notation will be as in [17]: capital letters and indices $\mu, \nu, \dots$ for (1+3) dimensional quantities, lower case letters for quantities of any dimension less than (1+3) and indices $\alpha, \beta, \dots$ for (1+2)-momenta. Then, the energy component E is the full (1+3) dimensional energy $E^2 = M^2 + |\vec{p}|^2$ with the (1+3) dimensional mass M , whereas e_T is the (1+2) dimensional transverse projected energy. However, the latter is not uniquely defined. While the transverse projection of the momentum components of a four-vector is unambiguous (see Figure 7.1), the projection of the energy component is not trivial.

Three different ways of transverse projections exist. Only two of them are of interest in this work, the “mass-preserving $\top$ ” and the “massless o” projection.

¹ $\mathcal{M}_i$ with any index expressing the way of projection denotes the invariant mass of the $W^{\pm}W^{\pm}$ system *before* the minimisation in the following.

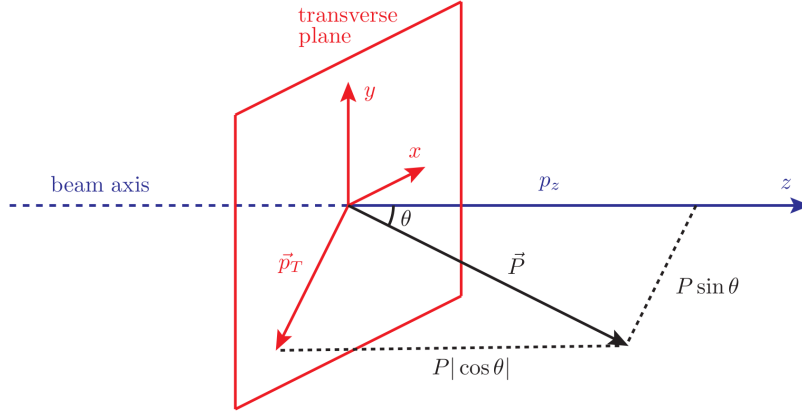


Figure 7.1: Illustration of the transverse projection at a collider experiment according to the ATLAS coordinate system introduced in Section 3.2.1. The beam forms the oriented z -axis. The transverse plane is spanned by the x - and y -axis. A 3-dimensional momentum vector $\vec{P}$ can uniquely be decomposed into a longitudinal component p_z and a transverse component $\vec{p}_T = (p_x, p_y)$ [17].

The mass-preserving “ $\top$ ” projection

$$p_{\top\alpha} \equiv (e_{\top}, \vec{p}_{\top}) \quad (7.7)$$

with

$$e_{\top} \equiv \sqrt{M^2 + |\vec{p}_T|^2} = \sqrt{E^2 - p_z^2}, \quad (7.8)$$

$$\vec{p}_{\top} \equiv \vec{p}_T \quad (7.9)$$

is defined in terms of the (1+3) dimensional mass M in order that the mass is preserved

$$p_{\top}^{\alpha} p_{\top\alpha} = m_{\top}^2 = M^2. \quad (7.10)$$

The massless “ o ” projection

$$p_{o\alpha} \equiv (e_o, \vec{p}_o) \quad (7.11)$$

with

$$e_o \equiv |\vec{p}_T|, \quad (7.12)$$

$$\vec{p}_o \equiv \vec{p}_T \quad (7.13)$$

is defined such that the invariant scalar product is always

$$p_o^{\alpha} p_{o\alpha} = m_o = 0. \quad (7.14)$$

The result are two different kind of transverse projected (1+2) dimensional momenta where the space-like momentum components $\vec{p}_T$ are completely identical, i.e. $\vec{p}_T \equiv \vec{p}_{\top} \equiv \vec{p}_o$, while the energy components differ. Both projections are equivalent in terms of usage but lead to different results for transverse mass reconstruction variables. The symbol “ T ” as an index is used for transverse projections in general, while “ $\top$ ” and “ o ” are used when referring to the special cases only.

7.2.2 The order of summing and projecting

Due to energy and momentum conservation, the energy and momentum of the resonance to be reconstructed, are obtained by summing the measured energies and momenta of all final decay particles. When forming the transverse mass there is a difference however, whether to form the composite momenta before or after the projection into the transverse plane. While again for the space-like components of the four-momentum this makes no difference

$$\sum_i \vec{p}_{iT} = \left(\sum_i \vec{P}_i \right)_T, \quad (7.15)$$

it does for the energy component

$$\sum_i e_{iT} \neq \left(\sum_i E_i \right)_T, \quad (7.16)$$

$$\sum_i e_{io} \neq \left(\sum_i E_i \right)_o. \quad (7.17)$$

The index i stands for the individual decay products. This behaviour for the energy components leads at first view to four different mass bound variables m_{1T} , m_{T1} , m_{o1} , m_{1o} . The index “1” stands for the number of reconstructed parent particles, in this case only one resonance. The order of the indices indicates the order of the operations from left to write, i.e. m_{1T} first sums all components and projects afterwards, whereas m_{T1} is early-projected and late-summed, equally for the “o” projected masses. The notation “ m_{WW}^{reco} ” is used whenever it is referred to any of the different mass reconstruction variables but no special one.

7.2.3 Minimisation

Once the event topology is ascertained and the choice for the projection and the order of operations for forming the invariant mass $\mathcal{M}_1$ was made, the constrained minimisation can be done. Equation (7.4) is taken as a general constraint. Others will be added later. The minimisation is conducted over all free parameters, i.e. all possible values of the unknown invisible momenta consistent with the measurement of $\vec{E}_T^{\text{miss}}$. This operation is always well-defined and unambiguous [17] with a numerical result

$$m_{1T} = \min_{\sum \vec{q}_{iT} = \vec{E}_T^{\text{miss}}} [\mathcal{M}_{1T}], \quad (7.18)$$

$$m_{T1} = \min_{\sum \vec{q}_{iT} = \vec{E}_T^{\text{miss}}} [\mathcal{M}_{T1}], \quad (7.19)$$

with $T = \{T, o\}$. The minimisation ensures that the result is the smallest invariant mass of all constituents consistent with the event’s kinematics, which automatically includes that it remains always below the probed resonance mass. This allows for a

concrete event-by-event statement for the reconstructed mass. It produces an invariant mass distribution of the $W^\pm W^\pm$ system with a distinct endpoint at or at least below the resonance mass

$$m_{WW}^{\text{reco}} \equiv m_{WW}^{\text{event}} \leq m_X, \quad (7.20)$$

i.e. a *lower bound* on the resonance mass. This is, of course, only true for events where the true event topology is the one assumed in Equation (7.1). As well, the assumption of $\vec{E}_T^{\text{miss}}$ coming only from the neutrinos of the decay has to be true. The minimisation fixes the separate invisible momenta q_{iT} and q_{iz} and makes the mass bounds independent of them. Minimising over the q_{iz} balances the rapidity of the visible and invisible subsystems to $\Delta y(\vec{P}_1, \vec{Q}_1) = 0$ which is in accordance to the mass bound variables being invariant under longitudinal boosts.

7.2.4 The resulting mass bound variables

The set of possible transverse mass bound variables is now introduced and constructed. The summation over the final decay products is performed separately for the visible and invisible particles. P^μ , p^α denote the composite momenta of all visible decay products, here the two charged leptons, in the respective dimension. Q^μ , q^α are the combined hypothesised invisible momenta of the neutrinos. E , e_T and $\tilde{E}$, $\tilde{e}_T$ denote the visible and invisible energies in the different dimensions, respectively.

The late-projected invariant mass $\mathbf{m}_{1\top}$

This variable is already in use for the ATLAS Higgs analysis in the W^+W^- decay channel [96–98]. The respective four-momenta are first summed and then projected into the transverse plane via the mass-preserving “ $\top$ ” projection

$$\begin{aligned} \mathcal{M}_{1\top}^2 &= (P^\mu + Q^\mu)_\top (P_\mu + Q_\mu)_\top \\ &= (e_{1\top} + \tilde{e}_{1\top})^2 - (\vec{p}_{1T} + \vec{q}_{1T})^2. \end{aligned} \quad (7.21)$$

with the momentum components

$$\vec{p}_{1T} \equiv \vec{p}_{l_1T} + \vec{p}_{l_2T}, \quad (7.22)$$

$$\vec{q}_{1T} \equiv \vec{q}_{\nu_1T} + \vec{q}_{\nu_2T} \quad (7.23)$$

and the energy components

$$e_{1\top} \equiv \sqrt{M_{ll}^2 + \vec{p}_{1T}^2}, \quad (7.24)$$

$$\tilde{e}_{1\top} \equiv \sqrt{\tilde{E}_1^2 - q_{1z}^2}. \quad (7.25)$$

M_{ll} denotes the invariant mass of the two charged leptons. The summed energy and longitudinal momentum of the invisible system are $\tilde{E}_1 = \sum_i \sqrt{M_{\nu_i}^2 + \vec{q}_{\nu_iT}^2 + q_{\nu_i z}^2}$ and $q_{1z} = q_{\nu_1 z} + q_{\nu_2 z}$, respectively, with M_{ν_i} being the masses of the two neutrinos². The

²Notice that the masses of the leptons are neglected once the minimisation is performed.

resulting invariant mass

$$\mathcal{M}_{1\top}^2 = \left\{ \sqrt{M_{ll}^2 + \vec{p}_{1T}^2} + \sqrt{\left(\sum_i \sqrt{M_{\nu_i}^2 + \vec{q}_{\nu_i T}^2 + q_{\nu_i z}^2} \right)^2 - q_{1z}^2} \right\}^2 - (\vec{p}_{1T} + \vec{q}_{1T})^2 \quad (7.26)$$

can then be minimised over the 6 invisible degrees of freedom $\vec{q}_{\nu_i T}$ and $q_{\nu_i z}$ conform to the $\vec{E}_T^{\text{miss}}$ constraint, leading to

$$\vec{q}_{\nu_i T} = \vec{E}_T^{\text{miss}} \frac{M_{\nu_i}}{M_{\nu_1} + M_{\nu_2}}, \quad (7.27)$$

$$q_{\nu_i z} = q_{1z} \frac{M_{\nu_i}}{M_{\nu_1} + M_{\nu_2}}, \quad (7.28)$$

with an arbitrary choice for q_{1z} . The resulting late-projected “mass-preserving” mass bound with the assumption of massless leptons is

$$m_{1\top}^2 = \left(\sqrt{M_{ll}^2 + \vec{p}_{1T}^2} + E_T^{\text{miss}} \right)^2 - \left(\vec{p}_{1T} + \vec{E}_T^{\text{miss}} \right)^2. \quad (7.29)$$

It can be noticed that, although this mass is call “transverse”, it still incorporates the dependence on visible longitudinal momenta in M_{ll} introduced through the energy components. This is due to the formation of composite particles before projecting their quantities.

The early-projected invariant masses $\mathbf{m}_{\top 1}$ and $\mathbf{m}_{o1}$

This is the case of inverse order for the summation and projection. Considering first the mass-preserving “ $\top$ ” projection. First, all quantities are projected into the transverse plane before adding them together to represent composite quantities

$$\begin{aligned} \mathcal{M}_{\top 1}^2 &= (p_{\top}^\alpha + q_{\top}^\alpha)(p_{\alpha\top} + q_{\alpha\top}) \\ &= (e_{\top 1} + \tilde{e}_{\top 1})^2 - (\vec{p}_{1T} + \vec{q}_{1T})^2. \end{aligned} \quad (7.30)$$

The momentum components are as in 7.22, while the energy components are

$$e_{\top 1} \equiv \sum_i \sqrt{M_{l_i}^2 + \vec{p}_{l_i T}^2}, \quad (7.31)$$

$$\tilde{e}_{\top 1} \equiv \sum_i \sqrt{M_{\nu_i}^2 + \vec{q}_{\nu_i T}^2}, \quad (7.32)$$

with M_{l_i} being the masses of the two charged leptons and M_{ν_i} as before. These expressions for the energy components do not contain any longitudinal momentum component which is the difference to the late-projected case. The invariant mass is then

$$\mathcal{M}_{\top 1}^2 = \left\{ \sum_i \sqrt{M_{l_i}^2 + \vec{p}_{l_i T}^2} + \sum_i \sqrt{M_{\nu_i}^2 + \vec{q}_{\nu_i T}^2} \right\}^2 - (\vec{p}_{1T} + \vec{q}_{1T})^2. \quad (7.33)$$

The minimisation is done equivalently to the previous case. The result is

$$m_{\top 1}^2 = \left(\sum_i \sqrt{M_{l_i}^2 + \vec{p}_{l_i T}^2} + \sqrt{(M_{\nu_2} + M_{\nu_2})^2 + E_T^{\text{miss}2}} \right)^2 - \left(\vec{p}_{1T} + \vec{E}_T^{\text{miss}} \right)^2. \quad (7.34)$$

Considering now again massless leptons this expression reduces to

$$\begin{aligned} m_{\top 1}^2 &= (|\vec{p}_{l_1 T}| + |\vec{p}_{l_2 T}| + E_T^{\text{miss}})^2 - \left(\vec{p}_{1T} + \vec{E}_T^{\text{miss}} \right)^2 \\ &= m_{o1}^2 \end{aligned} \quad (7.35)$$

In fact, this is exactly the same result as one would get with the early-projected massless “ o ” transverse invariant mass m_{o1} since there the masses of the individual particles are not considered right from the beginning. The energy components immediately are

$$e_{o1} \equiv \sum_i |\vec{p}_{l_i T}|, \quad (7.36)$$

$$\tilde{e}_{o1} \equiv \sum_i |\vec{q}_{\nu_i T}|. \quad (7.37)$$

The result of the minimisation already is Equation (7.35). Since, in this case of projection and the order of the operations, the masses of the leptons are neglected from the beginning, this is the more appropriate procedure in the case of massless leptons. Thus, the further investigation will use m_{o1} only.

The late-projected invariant mass m_{1o}

The procedure is the same as for the $m_{1\top}$ mass bound besides projecting the quantities with the massless “ o ” method

$$\begin{aligned} \mathcal{M}_{1o}^2 &= (P^\mu + Q^\mu)_o (P_\mu + Q_\mu)_o \\ &= (e_{1o} + \tilde{e}_{1o})^2 - (\vec{p}_{1T} + \vec{q}_{1T})^2. \end{aligned} \quad (7.38)$$

The energy components are

$$e_{1o} \equiv \left| \sum_i \vec{p}_{l_i T} \right|, \quad (7.39)$$

$$\tilde{e}_{1o} \equiv \left| \sum_i \vec{q}_{\nu_i T} \right|. \quad (7.40)$$

leading to

$$\mathcal{M}_{1o}^2 = \left\{ \left| \sum_i \vec{p}_{l_i T} \right| + \left| \sum_i \vec{q}_{\nu_i T} \right| \right\}^2 - (\vec{p}_{1T} + \vec{q}_{1T})^2 \quad (7.41)$$

which is minimised with the result

$$m_{1o}^2 = \{ |\vec{p}_{l_1 T} + \vec{p}_{l_2 T}| + E_T^{\text{miss}} \} - \left(\vec{p}_{1T} + \vec{E}_T^{\text{miss}} \right)^2. \quad (7.42)$$

7.2.5 Properties and relations of the mass bound variables

Saturation

The constructed mass bounds are calculated event by event resulting in a lower bound on the true resonance mass. The question now arises in which circumstances the constructed mass bound variables are saturated, i.e. when do they provide an endpoint at the actual resonance mass

$$m_{WW}^{\text{reco}} \equiv m_{WW}^{\text{event}} = m_X, \quad \text{with } m_{WW}^{\text{reco}} \in \{m_{1\top}, m_{o1}, m_{1o}\} \quad (7.43)$$

This is the case for events where the minimisation recovers the true kinematics of the event. For the mass-preserving “ $\top$ ” projection this is theoretically possible while for the massless “ o ” projection this will never occur if the true event topology comprises massive final decay particles. However, in the case under consideration, i.e. Monte Carlo simulated samples with only massless electrons and muons in the final state, the massless “ o ” projection will as well provide an endpoint at the resonance mass. Although these leptons are not massless in nature, this approximation is reasonable at the high energies probed at the LHC. However, the influence of the detector resolution, the background processes and the intrinsic width of the resonance will result in a smeared out endpoint.

Hierarchy

An ideal mass bound variable would cluster all resonant events just below the resonance mass leading to a sharp peak structure below the endpoint. However, the minimisation is only in few cases able to regain the real event kinematics. In most of the cases the minimisation leads to reconstructed kinematics which have an invariant mass which is below the true invariant mass of the final decay products of the resonance. Yet, the individually constructed mass bound variables do behave differently under minimisation due to their contained event information. In fact, it is possible to retrieve a hierarchy amongst these mass bounds which is found to be

$$m_{1\top} \geq m_{\top1} = m_{o1} \geq m_{1o}. \quad (7.44)$$

The equality between $m_{\top1}$ and m_{o1} was shown in Section 7.2.4 for the case of massless decay products. In addition, the process of projection discards information. The early projection weakens the bound as stated in the first relation between $m_{1\top}$ and $m_{\top1}$. While the late-projected mass bound $m_{1\top}$ retains some of the longitudinal momentum components, the $m_{\top1}$ drops all of them. Even weaker is the bound the case of m_{1o} where the visible invariant mass M_{ll} is discarded.

7.2.6 Constraints for on-shell $W^\pm$ bosons

The previously constructed mass bound variables use information obtained from the final decay kinematics. Thereby, the variable $m_{1\top}$ uses most information available. However, none of the previous variables uses constraints on the intermediate level of the decay topology. As determined by the event topology (7.1), the probed resonances decay via $W^\pm$ bosons. In addition, those resonances are considered to have masses

well above twice the W boson mass m_W (refer to Table 5.3). Therefore, it is reasonable to assume that – either one or even both – $W^\pm$ bosons are produced close to their mass shells. This introduces the mass-shell constraints for one or both $W^\pm$ bosons

$$\begin{aligned} m_W^2 &= (P_{l_i}^\mu + Q_{\nu_i}^\mu)(P_{\mu l_i} + Q_{\mu \nu_i}) \\ &= 2 \left(|\vec{p}_{l_i}| \cdot \sqrt{\vec{q}_{\nu_i T}^2 + q_{\nu_i z}^2} - \vec{p}_{l_i T} \vec{q}_{\nu_i T} - p_{l_i z} q_{\nu_i z} \right) \end{aligned} \quad (7.45)$$

in the limit of narrow $W^\pm$ boson width. Hence, based on the variable $m_{1\top}$, two new mass bound variables can be formed, $m_{1\top\text{star}}$ [11] for requiring one intermediate on-shell $W^\pm$ boson, and $m_{1\top\text{bound}}$ [12, 88] for both $W^\pm$ bosons being produced on their mass shell. However, in contrast to the previous “simple” mass bound variables, no analytical solution exists for these mass bounds. They alternatively have to be computed numerically. The relevant code is available on-line [99] and is used in these studies. This numerical approach seeks at splitting the measured missing transverse momentum amongst the two neutrinos. From a randomly chosen start configuration the smallest invariant mass of the $W^\pm W^\pm$ system m_{WW} is evaluated by scanning the possible splitting parameter space. The procedure is repeated up to a minimal scanning step size.

The result of the additional constraints is a new global minimum [17] for the invariant mass of the $W^\pm W^\pm$ system which is larger compared to the unconstrained case. In addition, these mass bound variables impose a new kinematical threshold for the invariant mass

$$m_{1\top\text{star}} \geq m_W, \quad (7.46)$$

$$m_{1\top\text{bound}} \geq 2m_W. \quad (7.47)$$

That means that values of m_{WW} for each event are pushed above the corresponding kinematical threshold and closer to the endpoint of the distribution. This effect arises for all events, for the resonant signal as well as for the background, whose topology might not even exhibit any on-shell $W^\pm$ boson. Again, for these variables a hierarchy can be found

$$m_{1\top\text{bound}} \geq m_{1\top\text{star}} \geq m_{1\top}. \quad (7.48)$$

However, while $m_{1\top\text{star}}$ is always bounded from above by the probed resonance mass, this does not need to be the case for $m_{1\top\text{bound}}$ [12]. Additionally, the first is always well-defined for any momentum configuration probed, while the existence of the latter is not always guaranteed [88].

7.3 Traditional mass reconstruction variables and their relation to mass bound variables

In addition to mass bound variables, traditionally in LHC experiments used variables shall be studied and their relationship to mass bound variables demonstrated. The construction of the mass variables is arranged from the most simple “visible mass” m_{vis} over the “effective mass” m_{eff} and the “vectorial mass” m_{vec} culminating in the collinear mass m_{col} .

The visible mass m_{vis}

The visible mass was used before in studies of $Z^0 \rightarrow \tau^+\tau^-$ [100], as well as in $H \rightarrow \tau^+\tau^-$ analyses of both LHC experiments, CMS [101] and ATLAS [102]. With ignoring the missing transverse momentum $\vec{E}_T^{\text{miss}}$ completely, it is the mass variable which uses the least event information. It is simply computed as the invariant mass of the visible lepton pair from the decay

$$\begin{aligned} m_{\text{vis}}^2 &\equiv M_{ll}^2 \\ &= (P_{l_1}^\mu + P_{l_2}^\mu)(P_{\mu l_1} + P_{\mu l_2}) \\ &= 2 \{ |\vec{p}_{l_1}| |\vec{p}_{l_2}| (1 - \cos \Delta\theta(l_1, l_2)) \}, \end{aligned} \quad (7.49)$$

where massless leptons are assumed and $\Delta\theta(l_1, l_2)$ is the angle between both leptons.

The effective mass m_{eff}

The effective mass is taken to be the scalar sum of the transverse momenta of both visible decay leptons plus missing transverse momentum $\vec{E}_T^{\text{miss}}$ as in [17]

$$m_{\text{eff}}^2 = (|\vec{p}_{l_1 T}| + |\vec{p}_{l_2 T}| + E_T^{\text{miss}})^2. \quad (7.50)$$

The relation with m_{o1} can be quantised with Equation (7.35)

$$m_{\text{eff}}^2 = m_{o1}^2 + \left(\vec{p}_{1T} + \vec{E}_T^{\text{miss}} \right)^2. \quad (7.51)$$

It is worthwhile considering the case where m_{eff} and m_{o1} become equal. This is the case when the resonance is produced without any other accompanying particle or radiation. Then, its decay products are the only particles that carry transverse momentum in the final state. Hence, with momentum conservation the last term in (7.51) vanishes and the both mass variables become equal. Regarding the VBS event topology (refer to Section 4.1) this cancellation of transverse momenta will be very rare due to the two tagging jets which always need to be present.

The vectorial mass m_{vec}

The vectorial mass is not a transverse mass variable in the classical meaning as it includes the longitudinal momenta of the visible leptons. It is basically the sum of four-momenta of the two charged leptons and missing transverse momentum

$$\begin{aligned} m_{\text{vec}}^2 &= (P_{l_1}^\mu + P_{l_2}^\mu + P^{\mu \text{miss}})(P_{\mu l_1} + P_{\mu l_2} + P_{\mu}^{\text{miss}}) \\ &= (|\vec{p}_{l_1}| + |\vec{p}_{l_2}| + E_T^{\text{miss}})^2 - \left(\vec{p}_{l_1 T} + \vec{p}_{l_2 T} + \vec{E}_T^{\text{miss}} \right)^2 - (p_{l_1 z} + p_{l_2 z})^2, \end{aligned} \quad (7.52)$$

where $P_\mu^{\text{miss}} = (E_T^{\text{miss}}, \vec{E}_T^{\text{miss}}, 0)$. Therefore, the relation

$$m_{\text{vec}}^2 = m_{o1}^2 + \left(\vec{p}_{1T} + \vec{E}_T^{\text{miss}} \right)^2 + (p_{l_1 z} + p_{l_2 z})^2 \quad (7.53)$$

arises.

The collinear mass m_{col}

The collinear mass [95] was used in the context of τ decays of the Z^0 boson as well as for Higgs $\rightarrow \tau^+ \tau^-$ decays [84, 103] and in the context of VBS resonances [69]. The approach assumes that the $W^\pm$ bosons from the decay of a heavy resonance receive a large Lorentz boost due to the large difference between the $W^\pm$ boson's and the resonance's mass. The collinear mass is then constructed by taking the decay products as being collinear to their corresponding parent $W^\pm$ boson. Hence, all three particles, $W^\pm$ boson, charged lepton and neutrino, have the same momentum direction and in particular the measurable momentum of the visible lepton can be expressed as

$$\vec{p}_{l_i T} = x_i \cdot \vec{p}_{W_i T}. \quad (7.54)$$

That means the charged lepton carries a momentum fraction x_i of the corresponding $W^\pm$ boson. With momentum conservation in the transverse plane

$$\vec{p}_{W_1 T} + \vec{p}_{W_2 T} = \vec{p}_{l_1 T} + \vec{p}_{l_2 T} + \vec{E}_T^{\text{miss}}, \quad (7.55)$$

the momentum fractions x_i can be determined

$$x_1 = \frac{p_{l_1 x} p_{l_2 y} - p_{l_1 y} p_{l_2 x}}{p_{l_2 y} p_x^{\text{miss}} - p_{l_2 x} p_y^{\text{miss}} + p_{l_1 x} p_{l_2 y} - p_{l_1 y} p_{l_2 x}}, \quad (7.56)$$

$$x_2 = \frac{p_{l_1 x} p_{l_2 y} - p_{l_1 y} p_{l_2 x}}{p_{l_1 x} p_y^{\text{miss}} - p_{l_1 y} p_x^{\text{miss}} + p_{l_1 x} p_{l_2 y} - p_{l_1 y} p_{l_2 x}}. \quad (7.57)$$

Only momentum fractions $0 < x_i < 1$ are physical. The invariant mass of the $W^\pm W^\pm$ system

$$m_{\text{col}}^2 = (P_{W_1}^\mu + P_{W_2}^\mu) (P_{\mu W_1} + P_{\mu W_2}) \quad (7.58)$$

can then be formed.

However, a problem arises when the resonance is decaying at rest, since then the two $W^\pm$ bosons are emitted in opposite direction. Thus, $\vec{E}_T^{\text{miss}}$ is equal to zero and the momentum fractions tend to one, which is not a physical solution since the neutrino momenta would be zero in that case.

7.4 Summary of mass reconstruction variables

The following mass bound variables are used in the following studies:

- $m_{1\top}$ — late-projected mass-preserving “ $\top$ ”,
- m_{o1} — early-projected massless “ o ”,
- m_{1o} — late-projected massless “ o ”,
- $m_{1\top\text{star}}$ — $W^\pm W^{\pm*}$ decay topologies,
- $m_{1\top\text{bound}}$ — $W^\pm W^\pm$ decay topologies.

The following “traditional” mass variables are used:

- m_{vis} — invariant mass of lepton pair (i.e. M_{ll}),
- m_{eff} — scalar sum of transverse momenta of both leptons and $\vec{E}_T^{\text{miss}}$,
- m_{vec} — invariant mass of lepton pair + $\vec{E}_T^{\text{miss}}$ (with p_z of leptons),
- m_{col} — collinear approach.

8 Study on mass reconstruction for resonances in $W^\pm W^\pm$ Scattering

Scenarios of BSM physics such as resonances in the scattering of two like-sign $W^\pm$ bosons have never been observed to date. In order to approach a possible discovery, an estimate on the discovery potential of a resonance needs to be evaluated. In addition, by facing the problem of lacking event information for the reconstruction of the full invariant mass of any possible resonance in this channel with the ATLAS detector, ways of measuring the mass of such resonances need to be found. In the following study, different mass reconstruction variables are employed in order to accomplish an identification of the resonance parameters. Before making best use of these variables, their properties need to be investigated and their performance compared. This is realised on the basis of Monte Carlo generated events of $W^\pm W^\pm$ scattering in associated production with two jets and with additional resonances, as explained in Chapter 5. The approach followed solely fixates the attention on the kinematics of the VBS in the channel $W^\pm W^\pm jj$ as explained in Chapter 4.

In this chapter, first the parton level quantity of the true invariant diboson mass is investigated in order to subsequently construct a possible signal definition. Further, the mass reconstruction variables, explained in the previous chapter, are compared and examined with respect to their generic properties. Finally, the discovery potential of the resonances using these mass reconstruction variables is explored in the context of the SM $W^\pm W^\pm jj$ -EW background at the LHC.

8.1 Reconstruction of the true invariant mass m_{WW}^{true}

At first, to get a handle on the resonances' intrinsic properties, the true invariant mass m_{WW}^{true} of the $W^\pm W^\pm$ system at parton level¹ is reconstructed. At parton level, i.e. the direct result from the hard scattering process $qq \longrightarrow qql^\pm \nu l^\pm \nu$, the $W^\pm$ bosons may not be accessed directly. However, since the six final state particles of the hard process are the only particles present at this stage of event reconstruction, a selected event (see event selection 6.2) comprises the particle content of exactly two jets, two charged leptons and two neutrinos². Hence, there is no ambiguity for the selection of the decay particles of the $W^\pm W^\pm$ system and the true invariant $W^\pm W^\pm$ diboson mass can be reconstructed by simply combining the four-momenta of the charged leptons and neutrinos as in Equation (7.2). Thus, in fact $m_{WW}^{\text{true}} = m_{l\nu l\nu}$.

In Figures 8.1 and 8.2, the m_{WW}^{true} distribution for the scalar isotensor resonance ϕ and the tensor isotensor resonance t are shown for different mass and coupling pa-

¹Refer to Section 6.4 for the event reconstruction levels.

²Neutrinos do not have any selection criteria, neither in p_T nor in $|\eta|$.

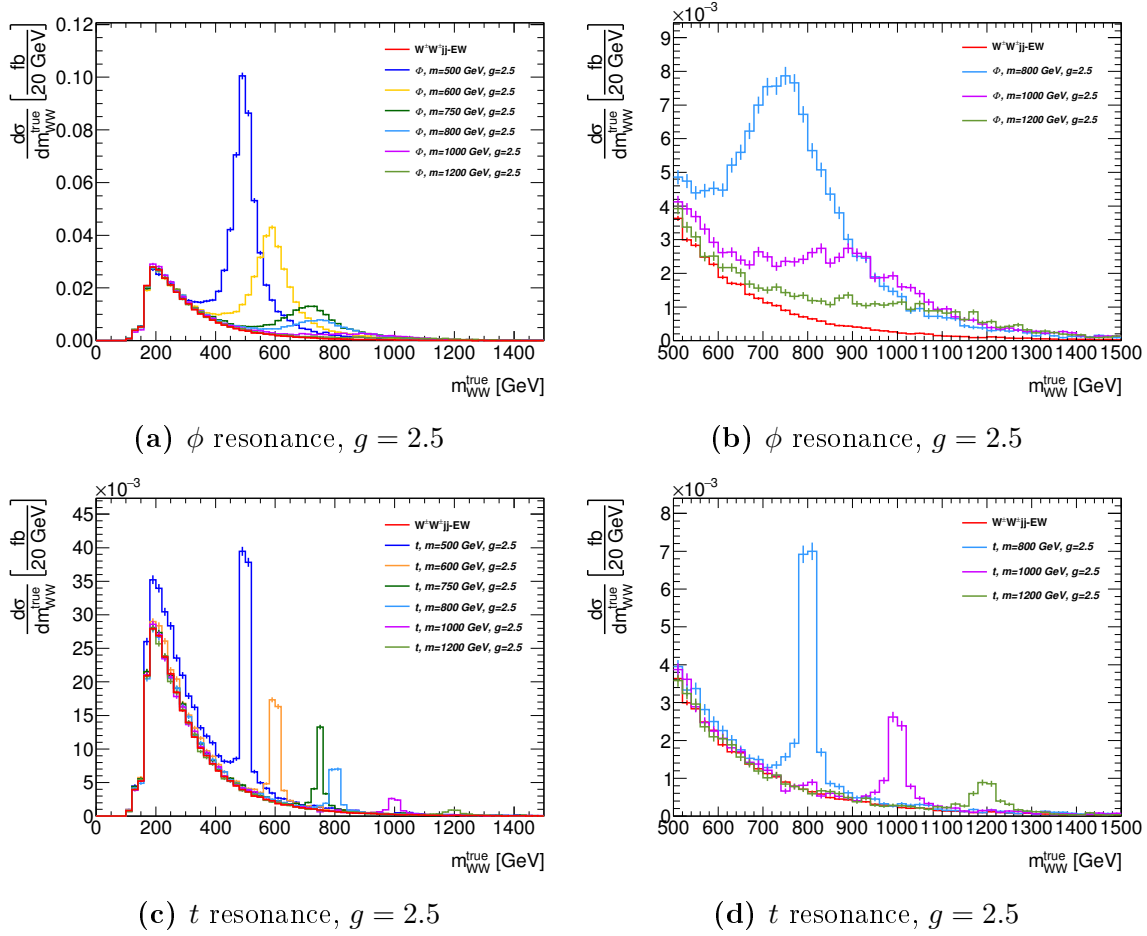


Figure 8.1: Distribution of the true invariant mass m_{WW}^{true} of the $W^\pm W^\pm$ system at parton level for the scalar isotensor resonance ϕ (a, b) and the tensor isotensor resonance t (c, d) for variable resonance masses $m \in \{500 \text{ GeV}, 600 \text{ GeV}, 750 \text{ GeV}, 800 \text{ GeV}, 1000 \text{ GeV}, 1200 \text{ GeV}\}$ and fixed coupling $g = 2.5$. Figures on the left show the full spectrum of the variation of the resonance masses while on the right the region of high m_{WW}^{true} is re-shown enlarged. All histograms are normalised to the corresponding parton level cross sections (refer to Table 5.3).

rameters, respectively. For comparison, the SM-EW $l^\pm \nu l^\pm \nu jj$ contribution is as well shown. It is referred to as $W^\pm W^\pm jj$ -EW as it will be throughout the entire analysis. Only events which pass the event selection, described in Section 6.2, are included in these and all following figures. The resonance masses and couplings are varied separately, as summarised in Table 5.3. The different intrinsic widths of the two resonance types, listed in Table 2.4, are clearly visible. While the ϕ resonance is broad and almost evolves to a flat continuum over the SM $W^\pm W^\pm jj$ -EW contribution for high masses, the t resonance's contribution is rather sharp. The position of the resonant peak occurs at the generated resonance mass for all t resonance samples. This is also observable for the ϕ resonance with not too high masses. However, for masses of the ϕ resonance above 750 GeV the peak is noticeably shifted towards lower m_{WW}^{true} . This can be due to the steeply falling continuum SM $W^\pm W^\pm jj$ -EW background, but,

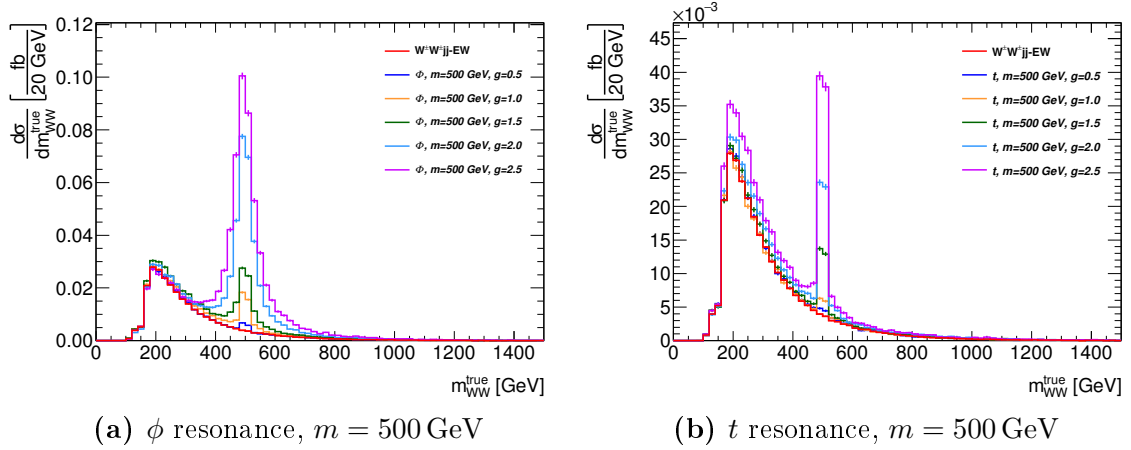


Figure 8.2: Distribution of the true invariant mass m_{WW}^{true} of the $W^{\pm}W^{\pm}$ system at parton level for the scalar isotensor resonance ϕ (left) and the tensor isotensor resonance t (right) for a fixed resonance mass $m = 500$ GeV and variable coupling $g \in \{0.5, 1.0, 1.5, 2.0\}$. All histograms are normalised to the corresponding parton level cross sections (refer to Table 5.3).

as shown in the next section, this shift remains also after subtraction of the SM $W^{\pm}W^{\pm}\text{jj-EW}$ contribution. Another consideration is the effect of unitarisation of the resonant amplitudes, which suppresses the resonant contribution before reaching the maximum value at the actual resonance mass [104]. As expected, the peak becomes more pronounced for larger coupling parameters for both resonance types. For the ϕ resonance the broadening of the resonant peak with increased coupling is visible.

Additionally, it can be noticed that, for all resonance samples, the SM $W^{\pm}W^{\pm}\text{jj-EW}$, besides some statistical fluctuations, is completely embedded in each resonant sample, i.e. it fully overlaps with the resonance samples. Contributions of new resonances are then on top of this contribution. For the t resonance, an excess can as well be observed for values of m_{WW}^{true} which are far below the actual resonance mass. An explanation would be non-resonant exchange of uncharged components of the t resonance in the t-channel of the $W^{\pm}W^{\pm}$ scattering.

8.2 Signal extraction – Construction of a mass window in m_{WW}^{true}

For the investigation of the performance and characteristics of the mass reconstruction variables it is reasonable to only select those events which have a resonant contribution. Therefore, a method has to be developed to efficiently select resonant events while rejecting non-resonant contributions. Here, an approach is taken which extracts the “signal” contribution with the application of a window cut on the true parton-level invariant mass m_{WW}^{true} .

As has been observed in the previous section, the SM $W^{\pm}W^{\pm}\text{jj-EW}$ sample is fully embedded in each resonance sample. This sample represents an estimate of the SM $W^{\pm}W^{\pm}\text{jj-EW}$ background contribution in this process. Thus, a reasonable straightforward first step to retrieve the resonant (“signal”) contribution, is the “subtraction”

of this estimate on the SM-EW background contribution from the resonance sample. This, however, cannot be done on an event basis, but has to be implemented via a mathematical description for histograms. Taking the distribution of m_{WW}^{true} with a total number of n bins, the remaining events with resonant contribution in bin $i \in [1, n]$ can be expressed as

$$N_i^{\text{signal, true}}(m, g) = N_i^{\text{res, true}}(m, g) - N_i^{\text{SM, true}}, \quad (8.1)$$

where $N_i^{\text{SM, true}}$ and $N_i^{\text{res, true}}$ are the numbers of events in bin i in the m_{WW}^{true} distribution in the SM $W^\pm W^\pm \text{jj}$ -EW and resonance sample, respectively. Clearly, as $N_i^{\text{res, true}}$ depends on the chosen mass and coupling parameter of the resonance, this is equally true for $N_i^{\text{signal, true}}$. The second step selects only a subset of bins to contribute to the total signal events $N^{\text{signal, true}}(m, g)$. This is done by finding the shortest window in the m_{WW}^{true} distribution which comprises the bin with the maximal number of entries and at least 68 % of the sum of all $N_i^{\text{signal, true}}$ in all bins i . Hence, the total number of signal events is

$$N^{\text{signal, true}}(m, g) = \sum_{i \in [i_{\min}^{68\%}, i_{\max}^{68\%}]} N_i^{\text{signal, true}}(m, g) \quad (8.2)$$

in a certain m_{WW}^{true} window including the bins from $i_{\min}^{68\%}$ to $i_{\max}^{68\%}$, which are determined by the algorithm. The distribution of $N^{\text{signal, true}}$ in m_{WW}^{true} and the resulting 68 % window bounds can be seen in Figure 8.3 for both resonance types, ϕ and t , and a selection of resonance parameters, $m \in \{500 \text{ GeV}, 750 \text{ GeV}, 1000 \text{ GeV}\}$ and fixed coupling $g = 2.5$. Distributions for the remaining probed resonance parameters can be found in Appendix A. The depicted distributions are normalised to the corresponding signal cross sections

$$\sigma^{\text{signal}}(m, g) = \sigma^{\text{res}}(m, g) - \sigma^{\text{SM}}, \quad (8.3)$$

with $\sigma^{\text{res}}(m, g)$ being the parton-level cross section for the resonant samples and similarly σ^{SM} for the SM $W^\pm W^\pm \text{jj}$ -EW sample, as listed in Table 5.3. The resulting window bounds in m_{WW}^{true} as well as the mean of the true invariant mass in this window and the position of the maximum of the m_{WW}^{true} distribution, $m_{WW}^{\text{true max}}$, are summarised in Table 8.1 for all probed points in the resonance parameter space.

The motivation for the application of a window requirement in m_{WW}^{true} in addition to the subtraction of the SM $W^\pm W^\pm \text{jj}$ -EW background contribution can be estimated from Figure 8.3b and the observation of an excess of events in m_{WW}^{true} below the generated resonance mass for the t resonance. Since those contribution do not inherit an s-channel resonance exchange, they would not contribute favourably to the reconstruction of the resonance mass. It may be interesting to study the impact of a variation of the fraction of events which is included in the mass window, especially for the t resonance and smaller percentage values.

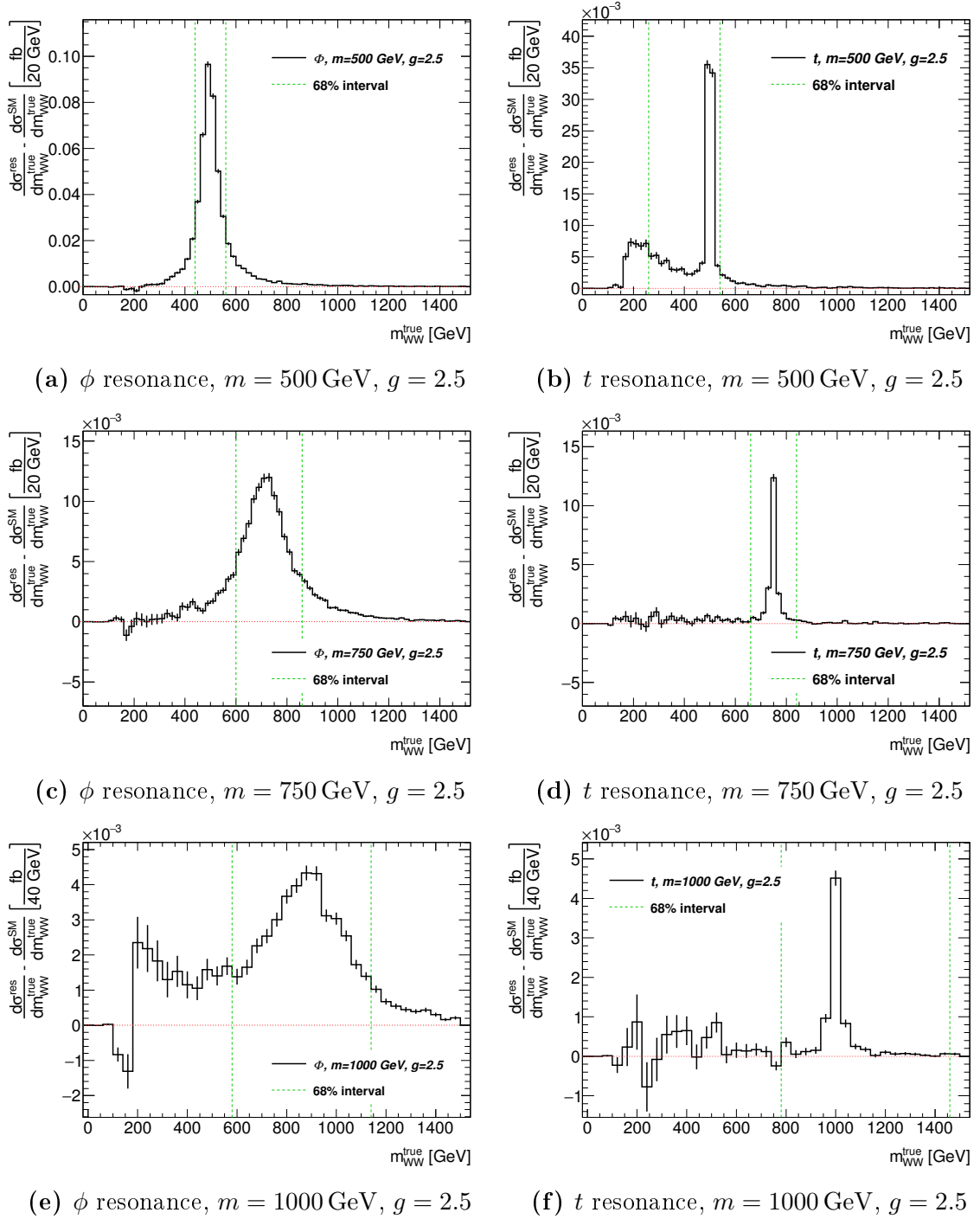


Figure 8.3: Signal extraction in the true invariant diboson mass at parton level m_{WW}^{true} for the ϕ (left) and t (right) resonance for masses $m \in \{500 \text{ GeV}, 750 \text{ GeV}, 1000 \text{ GeV}\}$ and coupling $g = 2.5$. The distribution in m_{WW}^{true} of the resonance samples was subtracted by the m_{WW}^{true} SM $W^\pm W^\pm jj$ -EW distribution. The window in m_{WW}^{true} (green, dashed) comprises 68% of the remaining subtracted contribution. The distributions are normalised to the corresponding signal cross section $\sigma^{\text{signal}}(m, g)$ from Equation (8.3). The figures for the remaining resonance parameter combinations in Table 8.1 can be found in Appendix A.

	resonance mass m [GeV]	coupling g	68% mass window in m_{WW}^{true} [GeV]	$\langle m_{WW}^{\text{true}} \rangle$ [GeV]	$m_{WW}^{\text{true max}}$ [GeV]
ϕ	500	2.5	440 – 560	497	490
	600		500 – 660	583	590
	750		600 – 860	721	730
	800		580 – 900	744	750
	1000		580 – 1140	865	880
	1200		700 – 1260	976	1000
t	500	2.5	260 – 540	451	490
	600		380 – 680	580	590
	750		660 – 840	749	750
	800		660 – 940	798	810
	1000		780 – 1460	1006	1000
	1200		1180 – 1220	1200	1200
ϕ	500	0.5	360 – 540	489	490
		1.0	360 – 560	484	490
		1.5	340 – 580	481	490
		2.0	440 – 540	495	490
t	500	0.5	220 – 620	440	490
		1.0	480 – 520	500	490
		1.5	300 – 540	471	490
		2.0	300 – 560	469	490

Table 8.1: Bounds of the 68% mass window construction in the true invariant diboson mass at parton level m_{WW}^{true} as well as the true mean $\langle m_{WW}^{\text{true}} \rangle$ and the location of the maximum of the m_{WW}^{true} distribution for the ϕ and t resonance signal with different mass and coupling parameters according to Figure 8.3 and the figures in Appendix A.

8.3 Distributions of mass reconstruction variables

In this section, different constellations and comparisons of distributions of mass reconstruction variables are investigated. At first, the mass bound variables are compared to the mass reconstruction variables without the feature of a mass bound. Further, the distributions for the different resonance types ϕ and t are examined as well as the differences between the mass reconstruction at parton level and after $\vec{E}_T^{\text{miss}}$ smearing.

The considered mass variables are reconstructed as explained in Section 7 for every event which passes the event selection requirements (Section 6.2). These variables are summarised in Section 7.4. The two selected charged leptons form the visible sector. The measured $\vec{E}_T^{\text{miss}}$ is included, if used, to account for the invisible neutrinos. The distributions comprise only signal events in the m_{WW}^{true} mass window for the corresponding resonance parameters as defined in the previous section. The distributions are shown at parton level in order to emphasise the generic properties of the mass reconstruction variables, unless stated otherwise. Furthermore, they are normalised to the signal cross section as defined in Equation (8.3), unless stated otherwise.

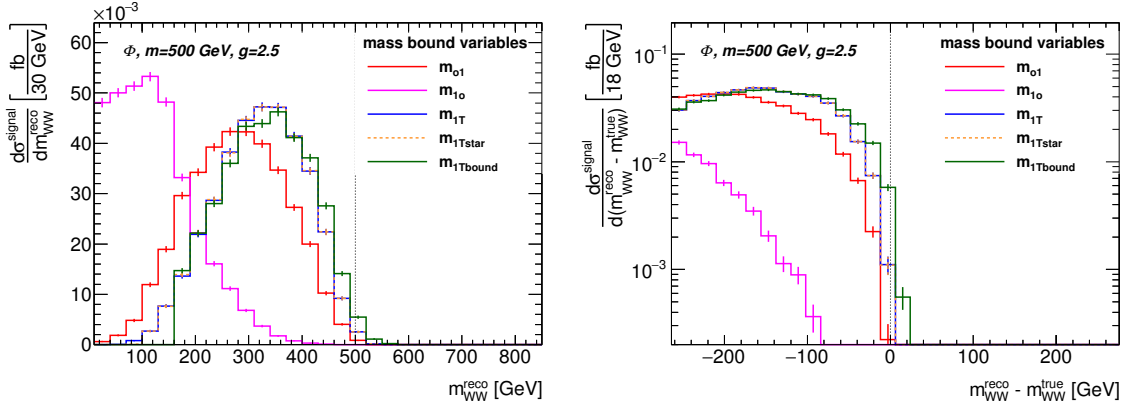
8.3.1 Comparison of different mass reconstruction variables

The focus of this investigation lies, firstly, in the comparison of the distributions of different mass bound variables and, secondly, in the comparison of these distributions to those of the traditionally used mass reconstruction variables.

Figure 8.4a depicts the different mass bound variables in the distributions of the invariant reconstructed mass³ m_{WW}^{reco} at parton level exemplary for the scalar isotensor resonance ϕ with parameters $m = 500 \text{ GeV}$ and $g = 2.5$. The difference between m_{WW}^{reco} and m_{WW}^{true} at parton level for the mass bound variables is shown in Figure 8.4b for the same resonance parameters. The reconstructed diboson mass for all these variables show the expected behaviour of building a lower mass bound for the resonance mass. Only few events lie beyond the generated resonance mass, which can be explained by respecting the generic width of the resonance.

Additionally, the predicted hierarchy (7.44) can be recovered. The variable m_{1o} clearly provides the weakest bound with a maximum far below the resonance mass and a gentle decline towards higher m_{WW}^{reco} . The endpoint of the distribution is located below the resonance mass due to the omission of the visible invariant mass of the dilepton pair. As emphasised in Figure 8.4b, the difference between the reconstructed and true invariant diboson mass is always below zero, i.e. the reconstructed mass is in no case able to recover the true invariant mass, but instead underestimates it systematically. Hence, the m_{1o} variable cannot provide a saturated mass bound in the probed topology. In contrary, the late-projected m_{o1} approaches the resonance mass, providing a weak, but saturated mass bound due to the approximation of massless leptons in the event simulation. Evidently, the class of mass-preserving variables $m_{1\top}$, $m_{1\top\text{star}}$ and $m_{1\top\text{bound}}$ exhibit the most pronounced mass bounds. The variable constrained by two on-shell $W^\pm$ bosons, $m_{1\top\text{bound}}$, accumulates the most events close to the resonance mass by shifting them closer to the endpoint with respect to $m_{1\top}$

³The notation “ m_{WW}^{reco} ” is used whenever it is referred to any of the different mass reconstruction variables but no special one.



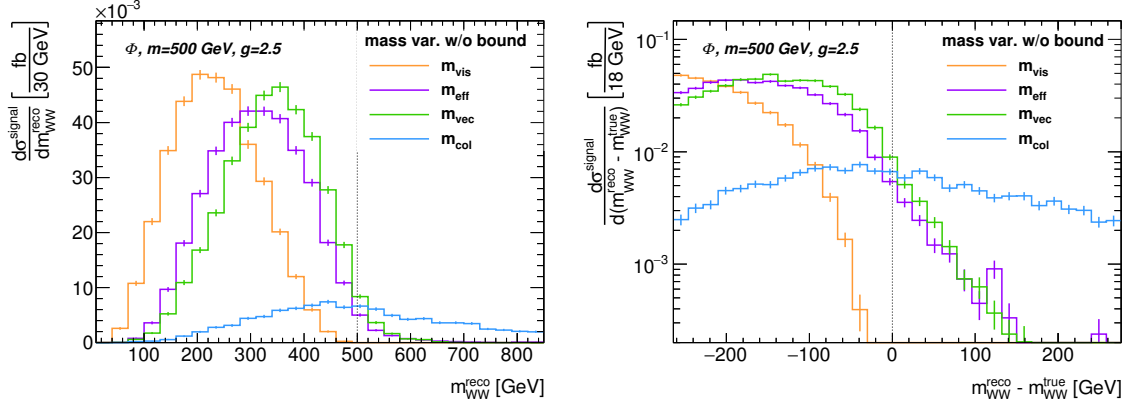
(a) m_{WW}^{reco} : mass bound variables, ϕ resonance, $m = 500$ GeV, $g = 2.5$

(b) $m_{WW}^{\text{reco}} - m_{WW}^{\text{true}}$: mass bound variables, ϕ resonance, $m = 500$ GeV, $g = 2.5$

Figure 8.4: Reconstructed invariant diboson mass m_{WW}^{reco} (left) and the difference between m_{WW}^{reco} and m_{WW}^{true} (right) at parton level for signal events in the true invariant mass window for the ϕ resonance with mass $m = 500$ GeV and coupling $g = 2.5$ for mass bound variables. The distributions are normalised to the corresponding signal cross section $\sigma^{\text{signal}}(m, g)$ from Equation (8.3).

and $m_{1T\text{star}}$. However it simultaneously overestimates the true invariant diboson mass for the highest fraction of events amongst the “T”-variables, although this fraction is small. Additionally, this variable obviously features the kinematic constraint $m_{WW}^{\text{reco}} \geq 2m_W$. As mentioned in Section 7.2.6, $m_{1T\text{bound}}$ can fail to give a physical solution, which is observed to be the case in less than one percent of the selected signal events for all probed resonance types and parameter combinations. This reconstruction efficiency is comparable to what was found for the process $H \rightarrow W^+W^-$ [88]. The overall picture shows only small differences between the “T”-variables. This is even more evident when comparing m_{1T} and $m_{1T\text{star}}$, since they show minor differences only for low m_{WW}^{reco} , a region which is of minor importance in these studies. This small difference is due to m_W being small compared to the resonance mass and the constraint on on-shell $W^\pm$ bosons becomes insignificant in the limit $m_{\text{resonance}}/m_W \rightarrow 0$ where $m_{1T\text{star}}$ becomes equal to m_{1T} . In conclusion, the variable $m_{1T\text{star}}$ is chosen to be dismissed for further studies due to the much greater computational effort for obtaining this variable compared to m_{1T} .

The distributions at parton level of the reconstructed invariant mass m_{WW}^{reco} and as well the difference between m_{WW}^{reco} and m_{WW}^{true} for the traditional mass reconstruction variables are illustrated in Figure 8.5. The effective mass m_{eff} , the vectorial mass m_{vec} , and the collinear mass m_{col} fail to give a lower mass bound. While the m_{eff} and m_{vec} variables overestimate the true invariant mass for only a few percent of the selected signal events for all probed resonance parameter configuration, this is the case in about 33 % to 70 % of the events in m_{col} depending on the resonance type and parameters. In addition, the collinear mass is not able to give a physical solution for the momentum fractions x_i of the leptons for circa 70 % of the selected signal events throughout the parameter space of probed resonances. It can be concluded that the assumption of collinear decay products of each $W^\pm$ boson cannot hold for this event topology. As this approach performs well in the case for Z^0 decays into $\tau^\pm$ lepton



(a) m_{WW}^{reco} : traditional mass variables, ϕ resonance, $m = 500$ GeV, $g = 2.5$

(b) $m_{WW}^{\text{reco}} - m_{WW}^{\text{true}}$: traditional mass variables, ϕ resonance, $m = 500$ GeV, $g = 2.5$

Figure 8.5: Reconstructed invariant diboson mass m_{WW}^{reco} (left) and the difference between m_{WW}^{reco} and m_{WW}^{true} (right) at parton level for signal events in the true invariant mass window for the ϕ resonance with mass $m = 500$ GeV and coupling $g = 2.5$ for traditional mass reconstruction variables. The distributions are normalised to the corresponding signal cross section $\sigma^{\text{signal}}(m, g)$ from Equation (8.3).

pairs, this might be due to the greater mass ratio $m_W/m_{\text{resonance}} > m_\tau/m_Z$ [69]. Due to this small reconstruction efficiency, m_{col} is discarded for further investigations. The visible mass underestimates the true invariant diboson mass and peaks at lower m_{WW}^{reco} values compared to the other traditional variables, which was expected due to the disuse of $\vec{E}_T^{\text{miss}}$.

The observations discussed for the scalar isotensor resonance ϕ can likewise be made for the tensor isotensor resonance t . The distributions at parton level of the reconstructed invariant mass m_{WW}^{reco} and as well the difference between m_{WW}^{reco} and m_{WW}^{true} for all examined mass reconstruction variables are shown in Figure 8.6.

In addition, no major changes in the quantitative result can be observed for higher mass parameters of the resonance or diminished resonance coupling. Example distributions for varied resonance parameters can be found in Appendix B. However, for masses greater than 1000 GeV and couplings less than 1.5 event statistics become very limited due to the small signal cross sections.

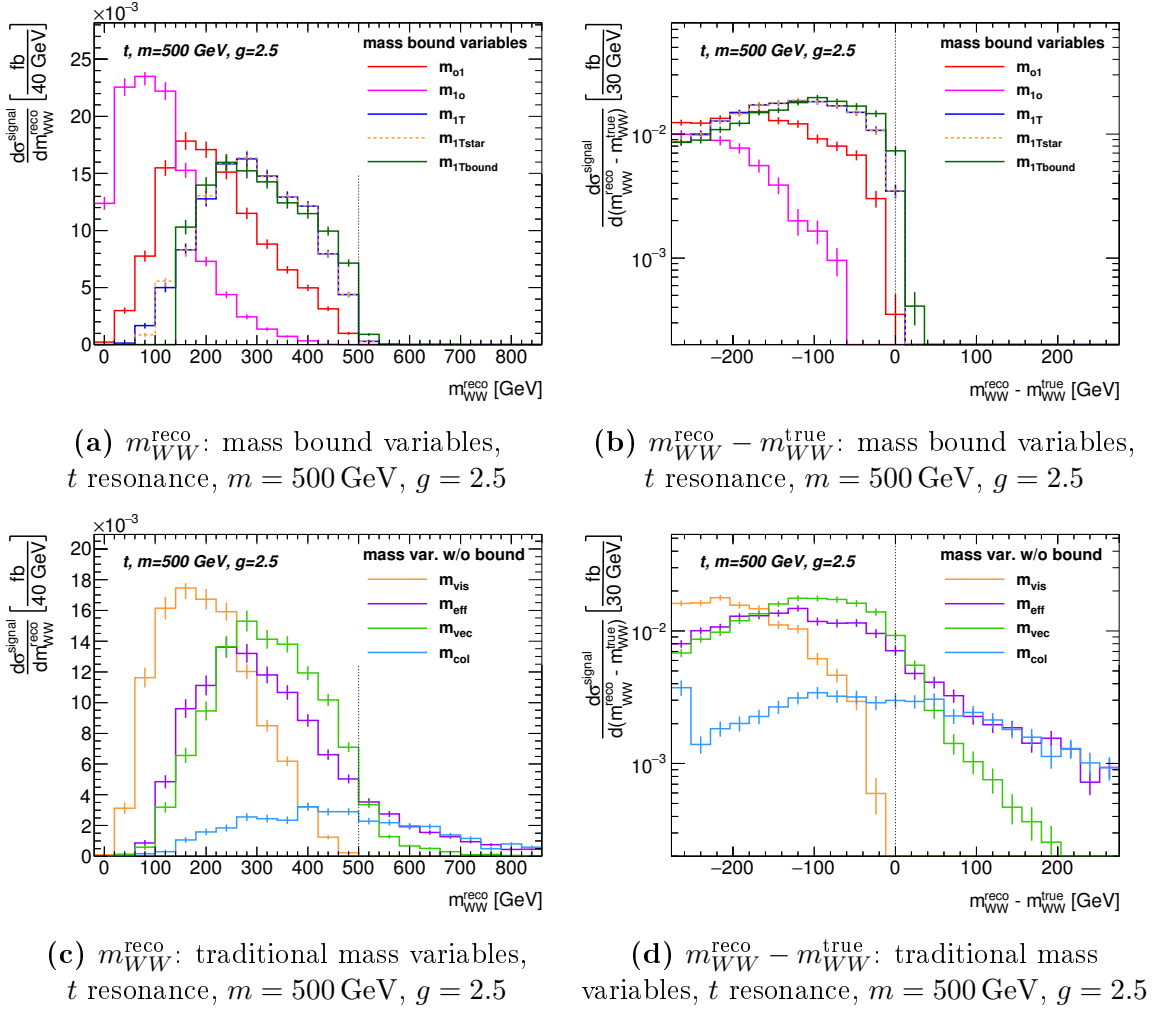


Figure 8.6: Reconstructed invariant diboson mass m_{WW}^{reco} (left) and the difference between m_{WW}^{reco} and m_{WW}^{true} (right) at parton level for signal events in the true invariant mass window for the t resonance with mass $m = 500$ GeV and coupling $g = 2.5$ for mass bound variables (a, b) as well as traditional mass reconstruction variables (c, d). The distributions are normalised to the corresponding signal cross section $\sigma^{\text{signal}}(m, g)$ from Equation (8.3).

8.3.2 Comparison of mass reconstruction variables for different resonance types

A direct comparison between the scalar isotensor resonance ϕ and the tensor isotensor resonance t at parton level is shown in Figure 8.7. The reconstructed invariant diboson mass m_{WW}^{reco} is illustrated for the mass variables m_{1T} , m_{o1} , and m_{vec} and different resonance mass parameters and coupling $g = 2.5$. For all mass variables a difference in shape between both resonance types can be observed. For low masses the maximum of the distribution of the t resonance is shifted to lower values of m_{WW}^{reco} with respect to the ϕ resonance. Hence, the distribution for ϕ features a sharper drop around the resonance mass for mass bound variables. For higher masses the shift of the maximum remains for m_{o1} , while it is reversed for m_{1T} and m_{vec} . These shape differences leave space for further investigation such as the spin dependence of the mass reconstruction.

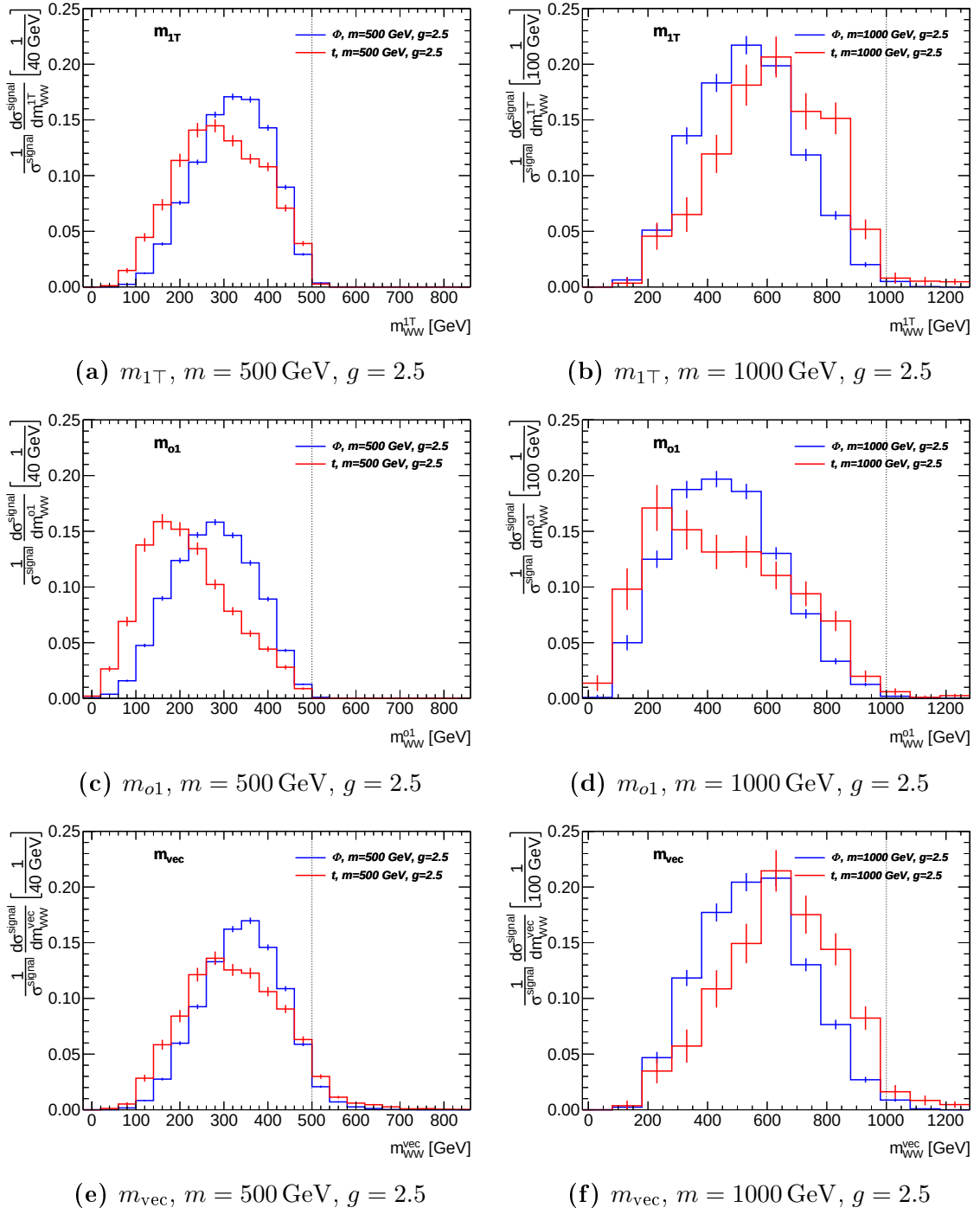


Figure 8.7: Comparison of the reconstructed invariant diboson mass m_{WW}^{reco} for the ϕ and t resonance signal at parton level in the true invariant mass window with resonance mass $m = 500$ GeV (left) and $m = 1000$ GeV (right) and coupling $g = 2.5$. Distributions for the mass reconstruction variables m_{1T} (a, b), m_{o1} (c, d) and m_{vec} (e, f) are shown. The distributions are normalised to unity in order to emphasise shape differences between the resonance types.

8.3.3 Comparison of mass reconstruction variables at parton level and after $\vec{E}_T^{\text{miss}}$ smearing

The previously shown studies were conducted at parton level in order to elaborate the intrinsic properties of the mass reconstruction variables. Following, however, is a study on the impact of particle level simulation and $\vec{E}_T^{\text{miss}}$ smearing. In Figure 8.8 the reconstructed invariant diboson mass m_{WW}^{reco} as well as the difference $m_{WW}^{\text{reco}} - m_{WW}^{\text{true}}$ for the ϕ resonance is compared between the parton level and the event reconstruction level after with $\vec{E}_T^{\text{miss}}$ smearing. The parameters $m \in \{500 \text{ GeV}, 1000 \text{ GeV}\}$ and $g = 2.5$ and the mass bound variable, m_{o1} , are chosen as an example. After $\vec{E}_T^{\text{miss}}$ smearing, the event selection efficiency increases due to more jets in the final state at particle level. In addition, the particle level comprises events that inherit $W^\pm$ bosons which decayed into $\tau^\pm + \nu_\tau$ and where these $\tau^\pm$ decayed leptonically. These events

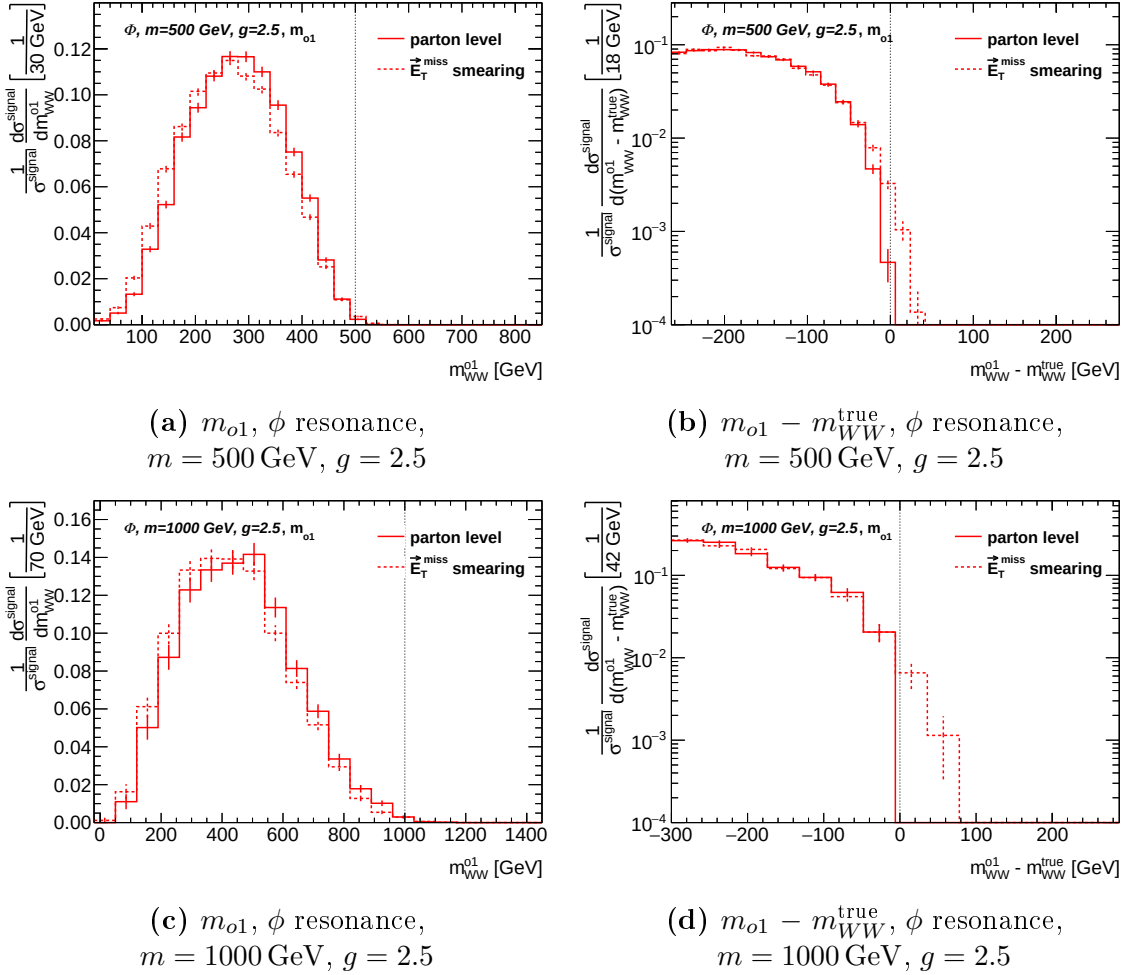


Figure 8.8: Comparison of the reconstructed invariant mass m_{WW}^{reco} (left) and the difference of m_{WW}^{reco} and m_{WW}^{true} (right) at parton level and after $\vec{E}_T^{\text{miss}}$ smearing for signal events in the true invariant mass window for the ϕ resonance with mass $m \in \{500 \text{ GeV}, 1000 \text{ GeV}\}$ and coupling $g = 2.5$. The mass reconstruction variable m_{o1} is depicted. The distributions are normalised to unity in order to emphasis shape differences.

might be selected at particle level due to additional $e^\pm$ or $\mu^\pm$ in the final state, whereas they are not selected at parton level. In order to focus on the shape difference of the distributions between the different event reconstruction levels, the distributions are normalised to unity. Throughout the tested resonance mass range, equal observations are made. The maximum of the reconstructed invariant mass is shifted to slightly lower values after $\vec{E}_T^{\text{miss}}$ smearing. As expected, the fraction of events overestimating the true diboson mass increases due to the selection of leptons which come from the parton shower rather than the actual leptons from the $W^\pm$ decay. However, no impact on the width of the distribution can be seen, which would be expected due to the broadening of the E_T^{miss} distribution after smearing as shown in Section 6.3. Figure 8.9 shows the comparison of the reconstructed invariant mass for the mass bound variable $m_{1\text{T}}$ and the t resonance at parton level and after $\vec{E}_T^{\text{miss}}$ smearing.

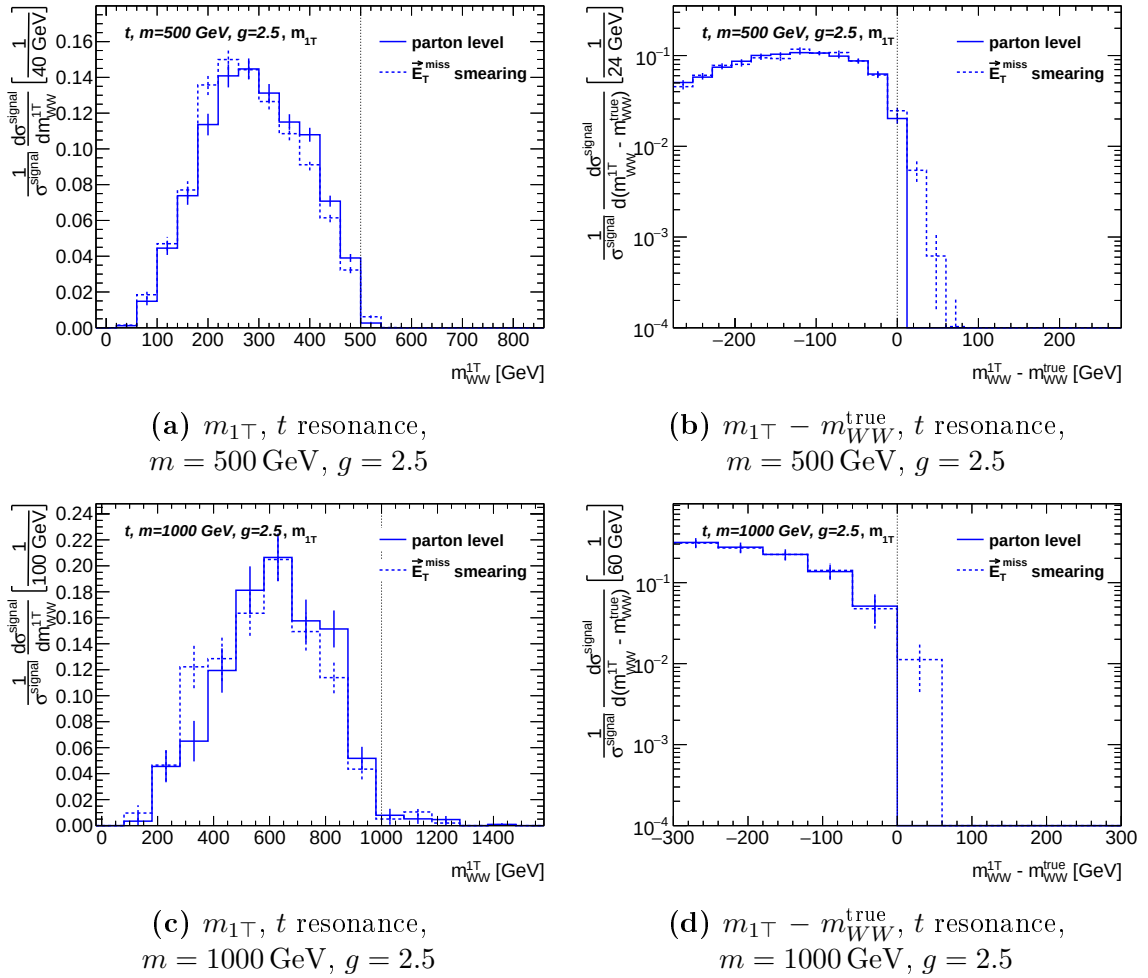


Figure 8.9: Comparison of the reconstructed invariant mass m_{WW}^{reco} (left) and the difference $m_{1\text{T}}^{\text{reco}} - m_{WW}^{\text{true}}$ (right) at parton level and after $\vec{E}_T^{\text{miss}}$ smearing for signal events in the true invariant mass window for the t resonance with mass $m \in \{500 \text{ GeV}, 1000 \text{ GeV}\}$ and coupling $g = 2.5$. The mass bound variable $m_{1\text{T}}$ is chosen. The distributions are normalised to unity in order to emphasis shape differences.

8.4 Characteristic parameters of the mass reconstruction variables

This section investigates different characteristics of the mass reconstruction variables for both resonance types with variable mass and coupling parameters. These studies are conducted at the event reconstruction level after $\vec{E}_T^{\text{miss}}$ smearing and with the signal events in the true mass window obtained as explained in Section 8.2, unless stated otherwise. The characteristic quantities are

- the maximum position $m_{WW}^{\text{reco}}{}^{\text{max}}$ of the m_{WW}^{reco} spectrum,
- the mean of the distribution $\langle m_{WW}^{\text{reco}} \rangle$,
- the width of the distribution $\sigma_{m_{WW}^{\text{reco}}}$,
- the resulting relative width $\frac{\sigma_{m_{WW}^{\text{reco}}}}{\langle m_{WW}^{\text{reco}} \rangle}$,
- and the corrected bias $\langle m_{WW}^{\text{reco}} \rangle - \langle m_{WW}^{\text{true}} \rangle$.

The corrected bias is investigated instead of the traditional bias $\langle m_{WW}^{\text{reco}} \rangle - m_{\text{resonance}}$ due to the shift of the resonance peak for the ϕ resonance with high mass parameters as explained in Section 8.1. The complete list of resulting characteristic quantities for the different resonance types and resonance mass and coupling parameters can be found in the tables in Appendix C. Subsequently, the most important quantities are discussed in detail, which are the mean $\langle m_{WW}^{\text{reco}} \rangle$, the corrected bias of the mean and the width $\sigma_{m_{WW}^{\text{reco}}}$.

8.4.1 Mean of the reconstructed invariant mass $\langle m_{WW}^{\text{reco}} \rangle$

In Figure 8.10 the dependency of the reconstructed mean $\langle m_{WW}^{\text{reco}} \rangle$ on the true mean $\langle m_{WW}^{\text{true}} \rangle$ is shown for all mass reconstruction variables and the ϕ and t resonance with $g = 2.5$. The mean of the true invariant mass $\langle m_{WW}^{\text{true}} \rangle$ is chosen instead of the generated resonance mass again because of the shift of the resonance peak in the true invariant mass spectrum shown in Section 8.1. The statistical uncertainties on the reconstructed mean m_{WW}^{reco} lie below one percent for all mass reconstruction variables and all probed resonance parameters, except for the t resonance with $m = 1200$ GeV and $g = 2.5$, where they are up to three percent due to limited event statistics. For all mass reconstruction variables the reconstructed mean shows a linear dependency on the true mean. The graphs are fitted with the linear function

$$m_{WW}^{\text{reco}} = a \cdot m_{WW}^{\text{true}} + b. \quad (8.4)$$

The resulting parameters of the fit are listed in Table 8.2 for both resonance types and all mass reconstruction variables. The slope for each mass reconstruction variable is comparable between the ϕ and t resonance, except for the variables m_{o1} and m_{eff} , where it is smaller in the case of the t resonance with respect to the ϕ resonance.

Further it can be observed that the mean of the spectrum of all mass reconstruction variables shows no correlation with the coupling parameter of the ϕ resonance. For

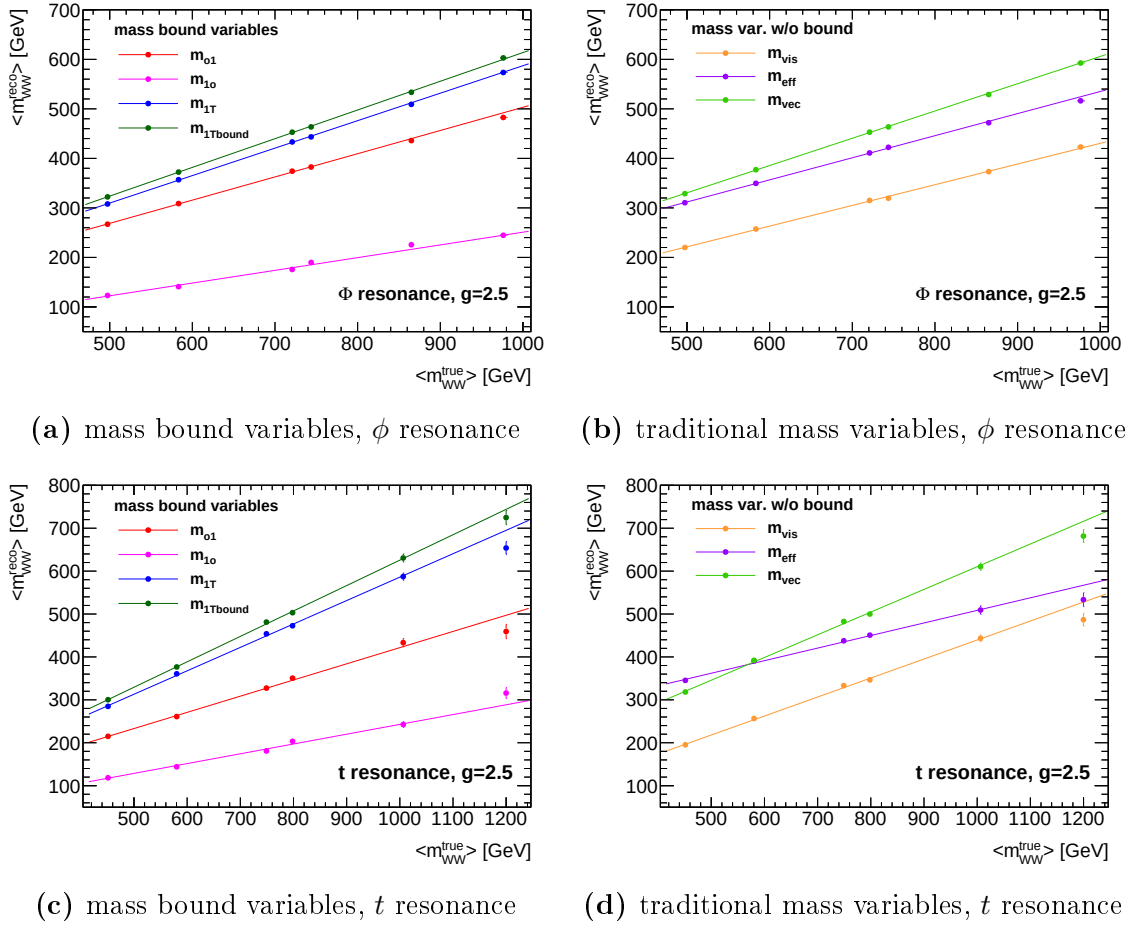


Figure 8.10: Mean of the distribution of the reconstructed invariant diboson mass $\langle m_{WW}^{\text{reco}} \rangle$ as a function of the mean of the true invariant diboson mass $\langle m_{WW}^{\text{true}} \rangle$ for signal events in the true invariant mass window and after $\vec{E}_T^{\text{miss}}$ smearing. The ϕ (a, b) and t (c, d) resonance with coupling $g = 2.5$ are depicted for mass bound variables (left) and traditional mass reconstruction variables (right). The distributions are each fitted with a linear function with parameters listed in Table 8.2.

the t resonance, a slight decrease with the coupling parameter can be observed. Figure C.1, which illustrates this behaviour, and the corresponding fit parameters listed in Table C.1 can be found in Appendix C.

An interesting issue is the possibility to extract the true position of the resonance peak from the measured reconstructed mean $\langle m_{WW}^{\text{reco}} \rangle$ in an experiment in order to determine a possible resonance's mass. However, the reconstructed mean must not depend on the chosen signal definition in the m_{WW}^{true} window. Therefore, for the given purpose, the full subtracted true signal spectrum is chosen without any requirement on the true invariant mass m_{WW}^{true} , i.e. without the application of a mass window. The reconstructed mean $\langle m_{WW}^{\text{reco}} \rangle$ as a function of the position of the maximum $m_{WW}^{\text{true max}}$ is illustrated in Figure 8.11 for all mass variables and the ϕ and t resonance. For lower resonance masses a linear behaviour can be observed, while the slope slightly reduces towards higher masses. An interpolation between the probed resonance masses might

be possible when narrowing the step size in between the mass parameters. In general a stronger dependency, i.e. a larger slope, is advantageous in order to extract the true peak position due to limited resolution.

resonance	m_{WW}^{reco} variable	linear fit parameters		
		a	b [GeV]	χ^2/NDF
$\phi, g = 2.5$	m_{o1}	0.469 ± 0.004	34.3 ± 2.4	9.1/4
	m_{1o}	0.257 ± 0.003	-6.0 ± 2.0	64.7/4
	$m_{1\top}$	0.555 ± 0.004	32.2 ± 2.4	4.7/4
	$m_{1\top\text{bound}}$	0.580 ± 0.004	33.8 ± 2.5	2.2/4
	m_{vis}	0.418 ± 0.004	12.4 ± 2.2	11.6/4
	m_{eff}	0.446 ± 0.004	88.8 ± 2.5	5.4/4
	m_{vec}	0.551 ± 0.004	55.0 ± 2.5	2.7/4
$t, g = 2.5$	m_{o1}	0.377 ± 0.008	44.5 ± 4.6	8.2/4
	m_{1o}	0.228 ± 0.006	14.4 ± 3.5	14.1/4
	$m_{1\top}$	0.546 ± 0.008	40.1 ± 4.4	14.3/4
	$m_{1\top\text{bound}}$	0.592 ± 0.008	33.2 ± 4.7	3.4/4
	m_{vis}	0.443 ± 0.007	-3.3 ± 4.1	13.8/4
	m_{eff}	0.293 ± 0.010	215.5 ± 6.3	7.3/4
	m_{vec}	0.529 ± 0.008	81.3 ± 4.8	11.5/4

Table 8.2: Parameter of the linear fit of the dependency of $\langle m_{WW}^{\text{reco}} \rangle$ on $\langle m_{WW}^{\text{true}} \rangle$ according to Figure 8.10 for resonances of both types with coupling $g = 2.5$. “NDF” denotes the number of degrees of freedom for the fit.

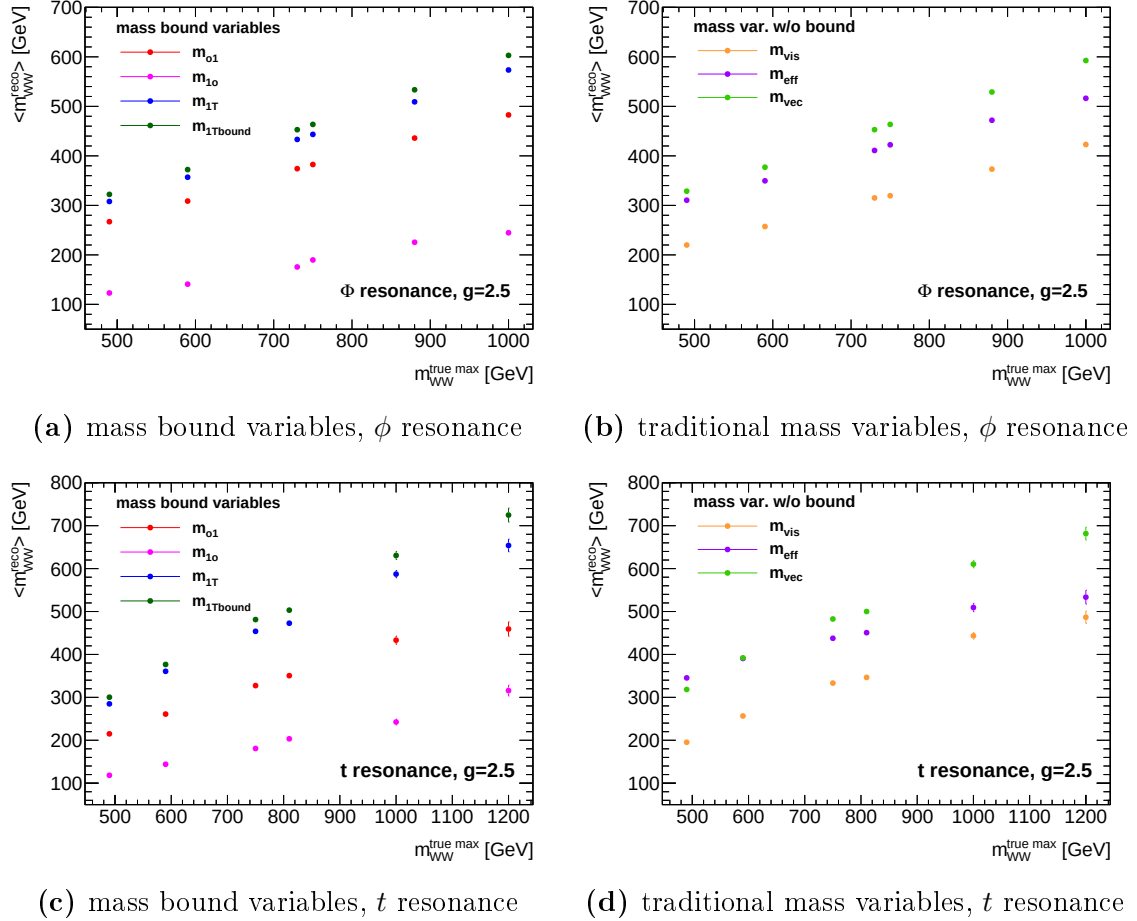


Figure 8.11: Mean of the distribution of the reconstructed invariant diboson mass $\langle m_{WW}^{\text{reco}} \rangle$ as a function of the location of the maximum in the true invariant diboson mass $m_{WW}^{\text{true max}}$ after $\vec{E}_T^{\text{miss}}$ smearing. $\langle m_{WW}^{\text{reco}} \rangle$ is determined for signal events without the requirement of a mass window on m_{WW}^{true} . The ϕ (a, b) and t (c, d) resonance with coupling $g = 2.5$ are depicted for mass bound variables (left) and traditional mass reconstruction variables (right).

8.4.2 Bias of the reconstructed mean $\langle m_{WW}^{\text{reco}} \rangle$

From Figure 8.10 in the previous section, it can already be observed that all mass reconstruction variables underestimate the mean of the true invariant mass m_{WW}^{true} . The corrected bias, $\langle m_{WW}^{\text{reco}} \rangle - \langle m_{WW}^{\text{true}} \rangle$, quantifies this behaviour and is shown for all mass variables and both resonance types in Figure 8.12 as a function of the mean of the true invariant mass $\langle m_{WW}^{\text{true}} \rangle$. Throughout the mass variables and resonance types the bias is negative and its absolute value increases with increasing resonance mass. For the mass bound variables, the order of the bias values reflects the hierarchy among these variables. Stronger mass bounds have smaller bias. The variable m_{1o} has the largest bias in accordance with the observations in Section 8.3.1. The bias of the traditional mass reconstruction variables is comparable with the values for the mass bounds. The visible mass, m_{vis} , inherits the largest bias amongst this class of variables, due to the least information input. Again, it can be found that the bias is independent of the coupling parameter of the resonances, except for small couplings of the t resonance, where statistics is very low. (see Figure C.2 in Appendix C).

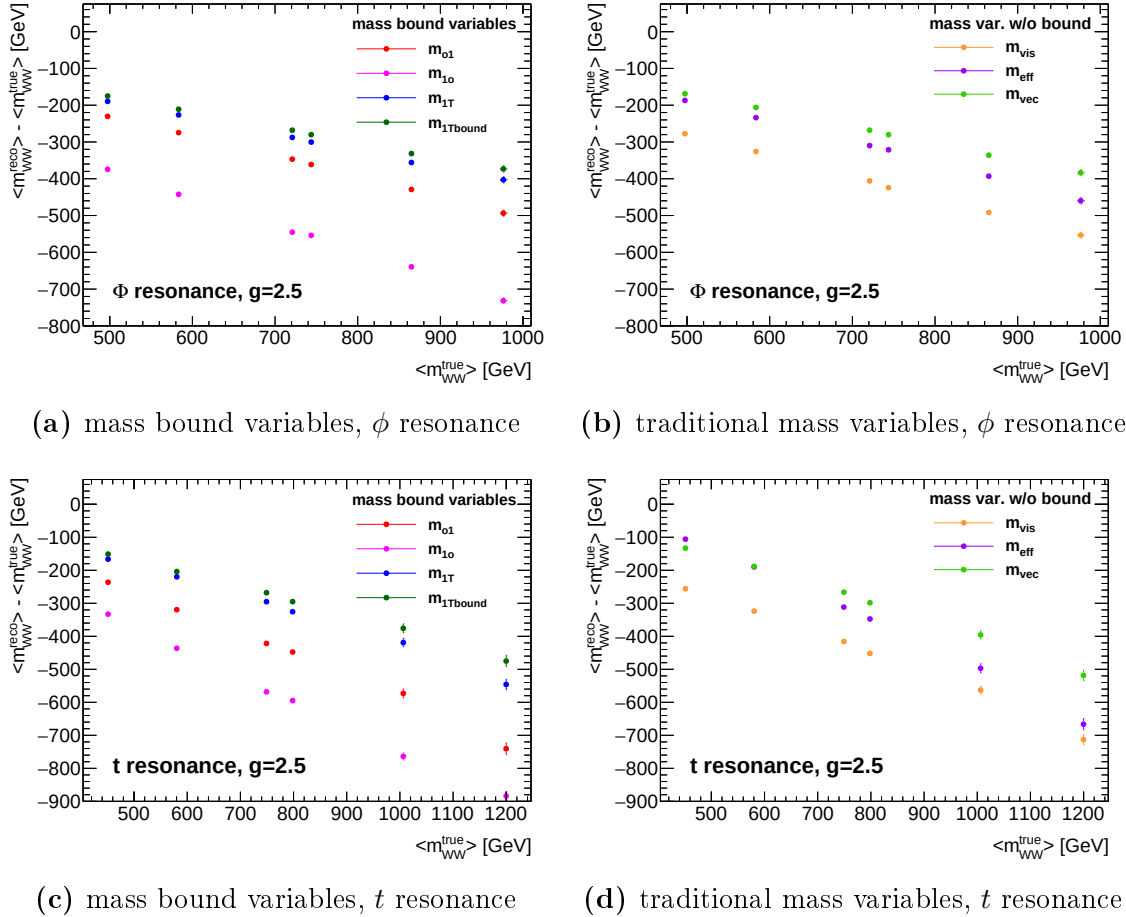
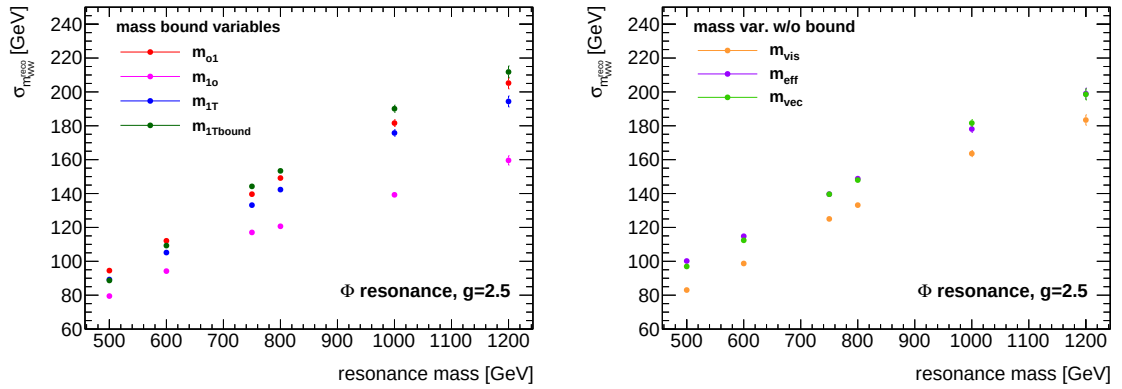
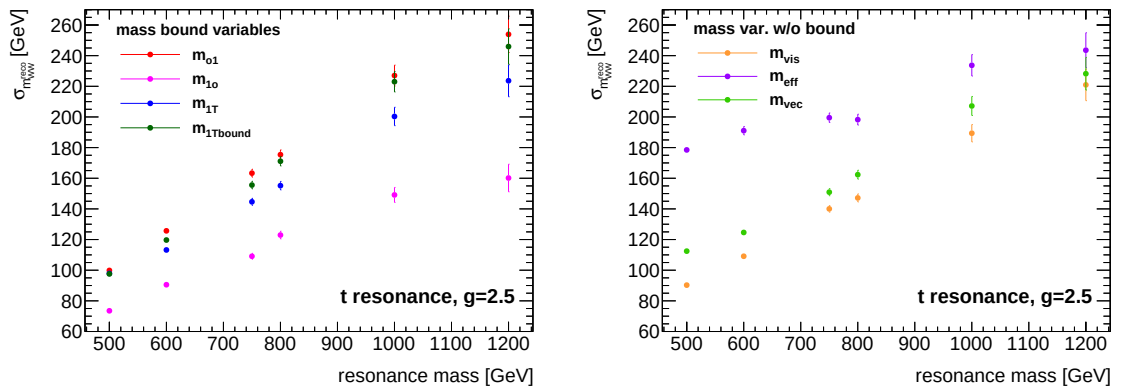


Figure 8.12: Corrected bias $\langle m_{WW}^{\text{reco}} \rangle - \langle m_{WW}^{\text{true}} \rangle$ as a function of the mean of the true invariant mass $\langle m_{WW}^{\text{true}} \rangle$ for signal events in the true invariant mass window and after $\vec{E}_T^{\text{miss}}$ smearing. The ϕ (a, b) and t (c, d) resonance with coupling $g = 2.5$ are depicted for mass bound variables (left) and traditional mass reconstruction variables (right).

8.4.3 Width $\sigma_{m_{WW}^{\text{reco}}}$ of the reconstructed invariant mass

The width of the reconstructed invariant mass is an important characteristics of mass reconstruction as it determines the resolution with which a possible resonance's mass can be reconstructed. It is illustrated in Figure 8.13 for both resonance types as a functions of the resonance mass. The width of m_{WW}^{reco} increases with the resonance mass for all investigated mass variables and both resonance types. The statistical uncertainties do not exceed one percent, except for resonance mass parameters above $m = 1000 \text{ GeV}$, where they rise up to five percent. The variable m_{1o} has the smallest width since it accumulates many events at low m_{WW}^{reco} . Due to the minimal information input, the visible mass m_{vis} has as well a smaller width than other variables. In addition, the effective mass m_{eff} performs worst for the t resonance while it is comparable with m_{vec} for the ϕ resonance. Furthermore, the variable m_{1T} shows a smaller width than its constrained version $m_{1T\text{bound}}$. Summing the overall picture, the width reflects the amount of information which is used to reconstruct the invariant diboson mass, separately in each class of reconstruction, the “ o ”-projected, the “ T ”-projected and the


 (a) mass bound variables, ϕ resonance

 (b) traditional mass variables, ϕ resonance

 (c) mass bound variables, t resonance

 (d) traditional mass variables, t resonance

Figure 8.13: Width of the reconstructed diboson mass $\sigma_{m_{WW}^{\text{reco}}}$ as a function of the resonance mass for signal events in the true invariant mass window and after $\vec{E}_T^{\text{miss}}$ smearing. The ϕ (a, b) and t (c, d) resonance with coupling $g = 2.5$ are depicted for mass bound variables (left) and traditional mass reconstruction variables (right).

traditional variables. Comparing the resonance types, the width for the reconstructed invariant mass for the tensor isotensor t is larger than the width of the reconstructed mass for the scalar isotensor ϕ , despite the narrower resonance width of the first. In addition, only a small decrease of the width of m_{WW}^{reco} with the resonance coupling can be observed (see Figure C.3 in Appendix C).

The relative width $\sigma_{m_{WW}^{\text{reco}}}/\langle m_{WW}^{\text{reco}} \rangle$ as a function of the resonance mass is depicted in Figure 8.14 for all mass variables. The variation of the relative width of each mass variable is relatively small throughout the full range of resonance mass. However, comparing the mass bound variables the relative width is quite different. The relative width amongst the different traditional mass variables does not spread largely. The best performing mass variables, for both resonance types, are the two “ $\top$ ”-projected mass variables and m_{vec} with an approximately constant value around 28 – 35 % over the probed resonance mass range. The overall same picture evolves for the dependency on the resonance coupling parameter (see Figure C.4 in Appendix C).

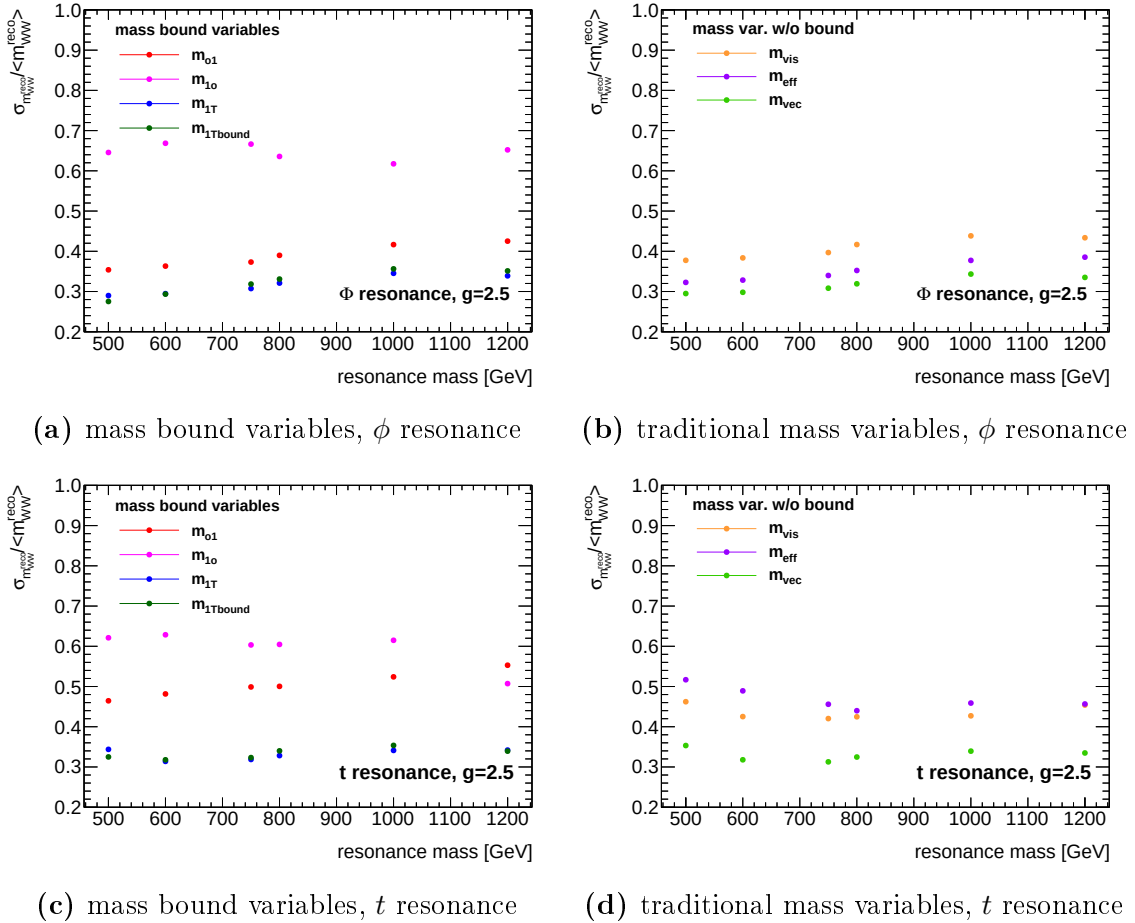


Figure 8.14: Relative width $\sigma_{m_{WW}^{\text{reco}}}/\langle m_{WW}^{\text{reco}} \rangle$ of the reconstructed diboson mass as a function of the resonance mass for signal events in the true invariant mass window and after $\vec{E}_T^{\text{miss}}$ smearing. The ϕ and t resonance with coupling $g = 2.5$ are depicted for mass bound variables (left) and traditional mass reconstruction variables (right).

It is worthwhile comparing the scales of the width of m_{WW}^{reco} , the generic width of the resonance (see Table 5.3) and the range of the m_{WW}^{true} mass window (see Table 8.1). As previously mentioned, $\sigma_{m_{WW}^{\text{reco}}}$ is slightly smaller for the broader ϕ resonance. However, for both resonance types, the width of the reconstructed mass distributions is of the same order of magnitude, while the width between these both resonance types differ by at least one order of magnitude. The range of the mass window is, except for the case of the t resonance with $m = 1200 \text{ GeV}$ and $g = 2.5$, larger than the width of m_{WW}^{reco} . Especially for the t resonance, where the range of the mass window shows no conform trend, no direct correlation can be observed. Therefore, the conclusion can be made that $\sigma_{m_{WW}^{\text{reco}}}$ represents the intrinsic resolution of the mass reconstruction variables, which is dependent on the mass parameter of the resonance, when taking resonant signal events.

8.5 Discovery potential of the mass reconstruction variables

The discovery potential is evaluated in terms of the signal, which is the resonant contribution of the $W^\pm W^\pm \text{jj}$ resonance distributions obtained after $\vec{E}_T^{\text{miss}}$ smearing. The background in this study comprises the dominant contribution of SM $W^\pm W^\pm \text{jj}$ -EW production. At first, the discovery potential is defined in the context of the signal and background event contribution. Finally, the discovery potential of studied mass reconstruction variables is compared and the best performing variables are determined. With these variables an estimation of the resonance parameters mass and coupling is conducted for a discovery potential of five standard deviations, which is defined to be the lower threshold for discovery. Further, differences for the different resonance types and the event reconstruction levels are investigated.

8.5.1 Technical approach

The discovery potential is investigated for signal events in the resonance samples over the SM $W^\pm W^\pm \text{jj}$ -EW background at the event reconstruction level after $\vec{E}_T^{\text{miss}}$ smearing. In order to obtain the expected number of events, both distributions are normalised to the product $\sigma \cdot \int \mathcal{L} dt$ with an integrated luminosity of 20 fb^{-1} corresponding to the luminosity collected with the ATLAS detector in 2012 with a centre-of-mass energy of 8 TeV . Signal events are then obtained similarly to the procedure explained in Section 8.2, except that, in the following case, the reconstructed distributions of each mass reconstruction variable m_{WW}^{reco} are taken instead of the true invariant mass distribution. For a histogram of n bins in any m_{WW}^{reco} variable, the signal events per bin $i \in [1, n]$ are obtained by subtracting the SM $W^\pm W^\pm \text{jj}$ -EW contribution $N_i^{\text{res, reco}}$ from the resonance events $N_i^{\text{SM, reco}}$ with parameters mass m and coupling g in each bin i :

$$S_i \equiv N_i^{\text{signal, reco}}(m, g) = N_i^{\text{res, reco}}(m, g) - N_i^{\text{SM, reco}} \quad (8.5)$$

with $N_i^{\text{SM, reco}} \equiv B_i$ hereafter. Further, for every mass reconstruction variable

$$r \in \{m_{\text{vis}}, m_{\text{eff}}, m_{\text{vec}}, m_{o1}, m_{1o}, m_{1\top}, m_{1\top\text{bound}}\}$$

the window $i \in [i_{r_{\min}}, i_{r_{\max}}]$ with

$$S = \sum_{i=i_{r_{\min}}}^{i_{r_{\max}}} S_i, \quad B = \sum_{i=i_{r_{\min}}}^{i_{r_{\max}}} B_i \quad (8.6)$$

is determined, such that $S/\sqrt{B}$ is maximised. The procedure for obtaining this window starts from the bin i which bears the maximal $S_i/\sqrt{B_i}$ and includes neighbour bins as long as $S/\sqrt{B}$ increases. The resulting value for $S/\sqrt{B}$ is referred to as discovery potential in the following.

Example distribution of the determination of the discovery potential are depicted in Figures 8.15 and 8.16 for varied resonance couplings and masses, respectively. For illustration purpose only the distribution for those mass reconstruction variables are shown, which give the best discovery potential for the given resonance parameter combination. The SM $W^\pm W^\pm \text{jj}$ -EW background contribution, the $W^\pm W^\pm \text{jj}$ contribution with resonance as well as the window for maximal $S/\sqrt{B}$ can be seen. In the tables in Appendix D, a summary of the discovery potential for each mass reconstruction variable and the corresponding m_{WW}^{reco} window is given for the scalar isotensor resonance ϕ and the tensor isotensor resonance t with variable mass and coupling parameters.

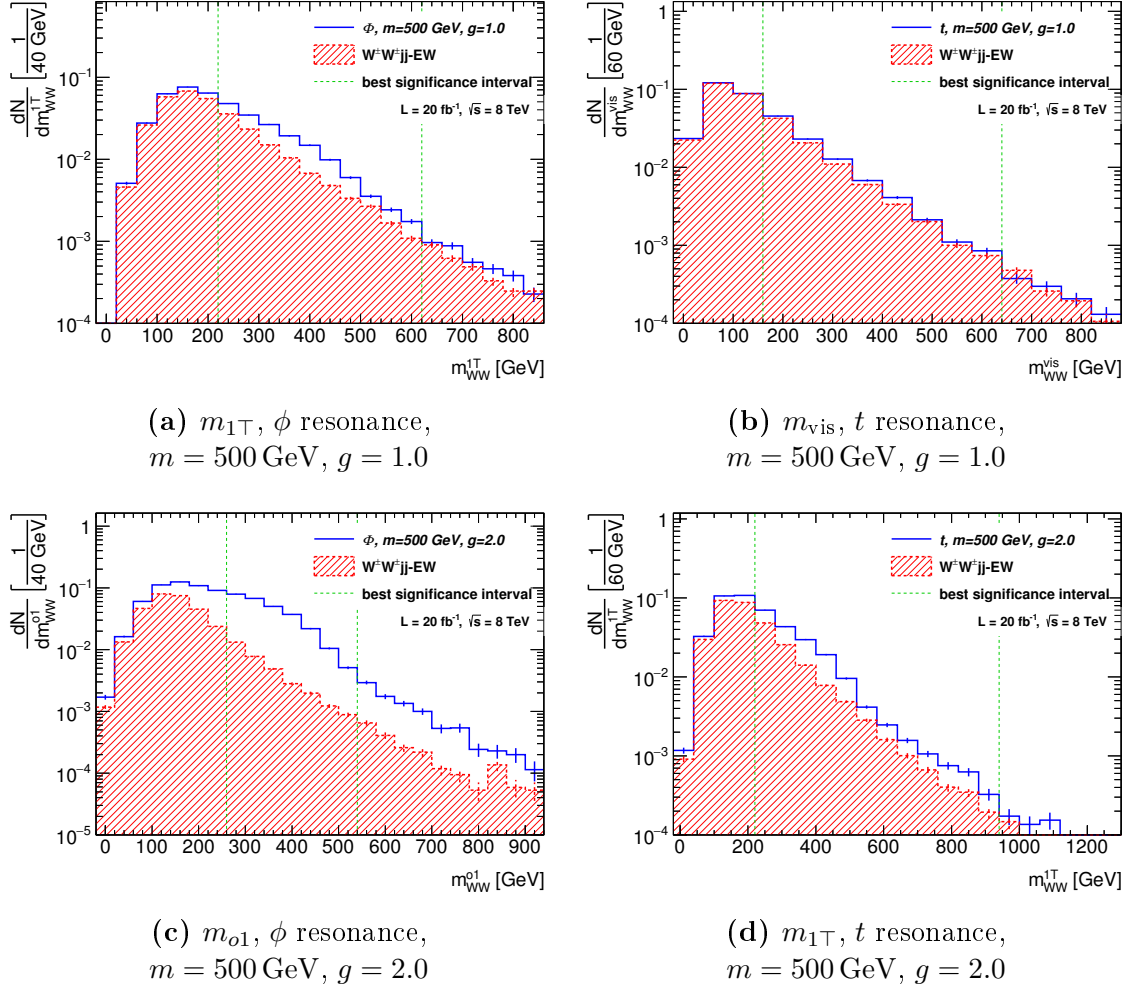


Figure 8.15: Reconstructed invariant diboson mass m_{WW}^{reco} for the determination of the discovery potential $S/\sqrt{B}$ for the ϕ (left) and t (right) resonance for the mass $m = 500 \text{ GeV}$ and couplings $g \in \{1.0, 2.0\}$. For demonstration purposes, for each resonance parameter combination, the distribution of the mass variable is shown which gives the best discovery potential (compare Table D.1 and D.2 in the appendix). The SM $W^\pm W^\pm \text{jj-EW}$ distribution (red, filled) represents an estimate on the expected background. The $W^\pm W^\pm \text{jj}$ contribution with resonances is shown in blue. The window comprising the bins with best discovery potential is marked (green, dashed). All distributions are normalised to the product of the corresponding parton level cross section (Table 5.3) and a luminosity of 20 fb^{-1} .

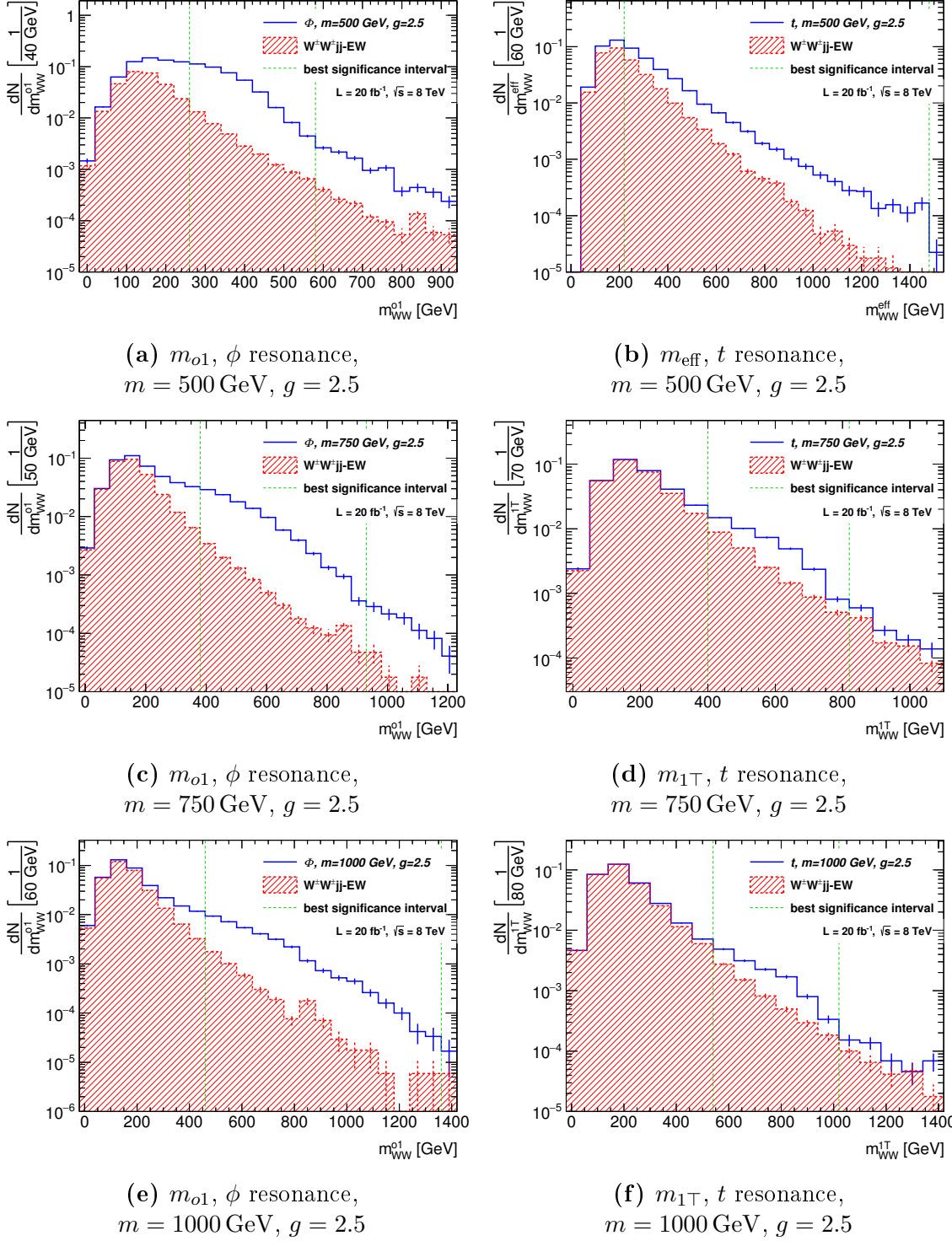


Figure 8.16: Reconstructed invariant diboson mass m_{WW}^{reco} for the determination of the discovery potential $S/\sqrt{B}$ for the ϕ (left) and t (right) resonance for masses $m \in \{500 \text{ GeV}, 750 \text{ GeV}, 1000 \text{ GeV}\}$ and coupling $g = 2.5$. For demonstration purposes, for each resonance parameter combination, the distribution of the mass variable is shown which gives the best discovery potential (compare Tables D.3 and D.4 in the appendix). The SM $W^\pm W^\pm \text{jj-EW}$ distribution (red, filled) represents an estimate on the expected background. The $W^\pm W^\pm \text{jj}$ contribution with resonances is shown in blue. The window comprising the bins with best discovery potential is marked (green, dashed). All distributions are normalised to the product of the corresponding parton level cross section (Table 5.3) and a luminosity of 20 fb^{-1} .

8.5.2 Comparison of the discovery potential for different mass reconstruction variables

The performance of the mass reconstruction variables with respect to the discovery potential shall be discussed in detail. Table 8.3 summarises the best performing mass reconstruction variables for both resonance types and all probed resonance parameter combinations. In Figure 8.17 the best performing mass reconstruction variables for each resonance type and points in the resonance parameter space are compared. The figure shows the discovery potential as a function of the resonance mass and coupling. As expected, the discovery potential decreases with increasing mass, while it increases with increasing coupling due to the corresponding cross sections.

resonance		variable with		$S/\sqrt{B}$
mass m [GeV]	coupling g	best $S/\sqrt{B}$		
ϕ	500	2.5	m_{o1}	284
	600			231
	750			149
	800			114
	1000			66
	1200			40
t	500	2.5	m_{eff}	54
	600		$m_{1\top(\text{bound})}$	31
	750		$m_{1\top(\text{bound})}$	22
	800		$m_{\text{vec}}, m_{1\top(\text{bound})}, m_{o1}$	18
	1000		$m_{1\top}$	13
	1200		$m_{1\top\text{bound}}$	7.9
ϕ	500	0.5	$m_{1\top}$	4.65
		1.0	$m_{1\top(\text{bound})}, m_{\text{vis}}$	27
		1.5	$m_{1\top(\text{bound})}$	56
		2.0	m_{o1}	186
t	500	0.5	m_{vec}	2.36
		1.0	m_{vis}	4.33
		1.5	$m_{\text{eff}}, m_{1\top}$	15
		2.0	$m_{1\top(\text{bound})}$	32

Table 8.3: Best mass reconstruction variables with respect to the discovery potential $S/\sqrt{B}$ together with their resulting performance for both resonance types, ϕ and t , and all probed resonance parameter combinations.

For the scalar isotensor resonance ϕ with a coupling of $g = 2.5$ and variable mass a clear picture arises. For all probed resonance masses the variable m_{o1} performs best with a fair enhancement over all other variables. This can be explained with the favourable background suppression of this variable with respect to the other variables. For smaller couplings of the ϕ resonance and a fixed mass of $m = 500$ GeV, the mass-

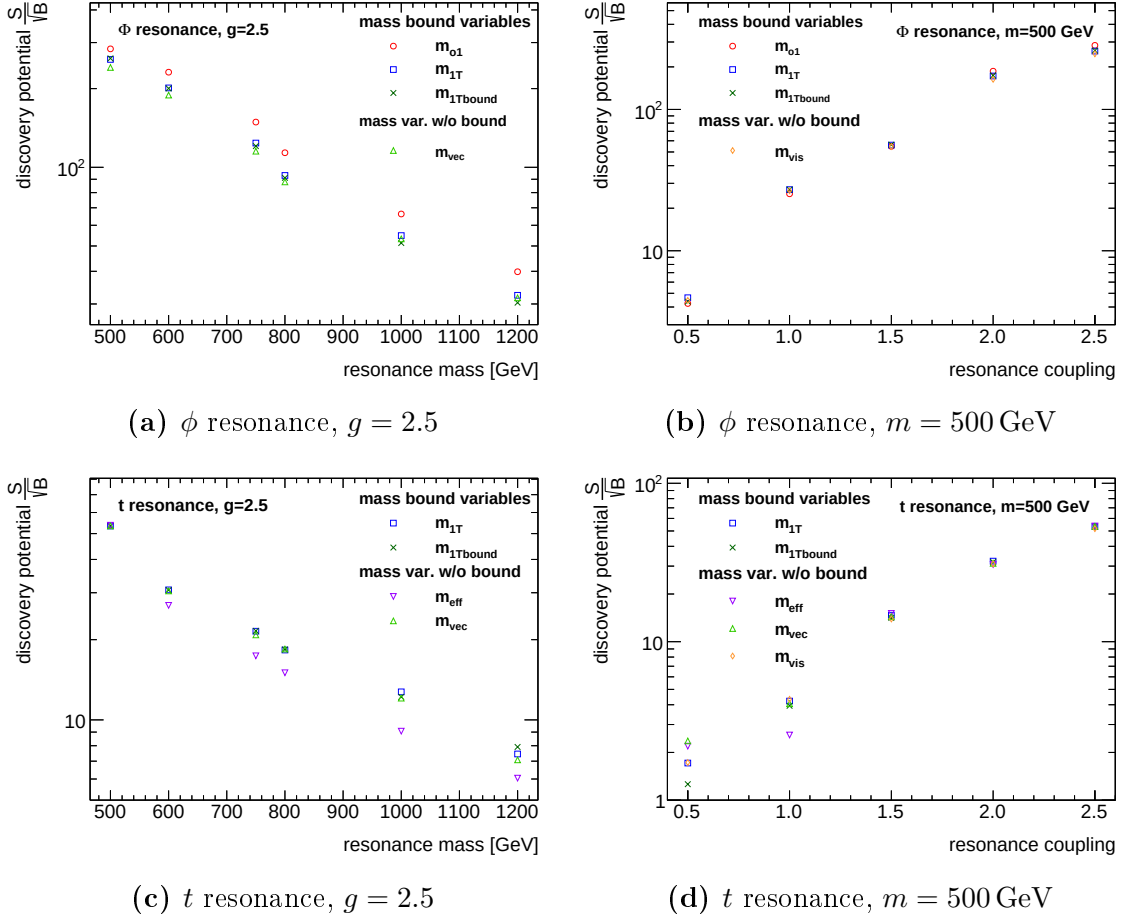


Figure 8.17: Discovery potential for the ϕ and t resonance as a function of the resonance mass (left) and the resonance coupling (right). A selection of the best performing mass reconstruction variables for each case with respect to the discovery potential is chosen (compare Table D.3 and Table D.1 in Appendix D). The distributions contain signal events after $\vec{E}_T^{\text{miss}}$ smearing.

preserving “T”-family variables m_{1T} and $m_{1Tbound}$ become best performing. For a coupling of $g = 1.0$, m_{vis} is amongst the variables with best discovery potential. Notwithstanding, the “T”-variables, m_{o1} and m_{vis} lie close together for any coupling parameter.

For the tensor isotensor resonance t with a high coupling and variable mass the “T”-family variables m_{1T} and $m_{1Tbound}$ show a slight advantage over the other variables. However, m_{eff} and m_{vec} are as well amongst the best performing variables for some cases. For lower coupling parameters the traditional mass reconstruction variables gain advantage over the mass bound variables. The vectorial mass m_{vec} , the visible mass m_{vis} and the effective mass m_{eff} are the best performing mass variables in this range of resonance parameters. Yet, the “T”-mass bound variables lie always close behind.

The overall picture shows that different types of the mass bound variables are always amongst the best performing mass reconstruction variables with respect to the discovery potential. Only for the t resonance with mass $m = 500$ GeV and coupling

$g = 0.5$ the mass bound variables show a considerably smaller $S/\sqrt{B}$ than traditional variables. Furthermore, the two different resonance types favour different mass bound variables, due to different characteristics such as spin and width. Moreover, $m_{1\text{Tbound}}$ differs only slightly from its unconstrained version $m_{1\text{T}}$ for all considered resonance parameter combinations. The computational effort for the $m_{1\text{Tbound}}$ reconstruction variable with respect to $m_{1\text{T}}$ exceeds considerably the gain in discovery potential for the probed event topology.

	resonance mass m [GeV]	coupling g	range for linear extrapolation [GeV]	m_5 [GeV]	fit parameters $a \cdot x + b$ $a \text{ GeV}^{-1}$ b	
ϕ	variable	2.5	{1000, 1200}	1464	-0.13	198.43
t	variable	2.5	{1000, 1200}	1320	-0.02	36.89
	Resonance mass m [GeV]	coupling g	range for linear interpolation	g_5	fit parameters $a \cdot x + b$ a b	
ϕ	500	variable	{0.5, 1.0}	0.51	44.92	-17.81
t	500	variable	{1.0, 1.5}	1.03	21.50	-17.17

Table 8.4: Summary of the resonance parameter extraction for a discovery potential $S/\sqrt{B} = 5$ for both resonance types. Listed are the range which is considered for obtaining the linear fit, the resulting values for the fit parameters of the linear fit function $a \cdot x + b$ as well as the resulting values for m_5 and g_5

It is now worthwhile to extract those resonance parameters (mass and coupling) which induce a discovery potential of $S/\sqrt{B} = 5$. This can be achieved separately for the different resonance types ϕ and t . For the resonances with mass $m = 500 \text{ GeV}$ and variable coupling, a coupling parameter $g = g(S/\sqrt{B} = 5) \equiv g_5$ can be obtained by linearly interpolating between the points of examined resonance parameters. For the resonances with fixed coupling $g = 2.5$ and variable mass, a mass parameter $m = m(S/\sqrt{B} = 5) \equiv m_5$ is extracted by linearly extrapolating towards masses larger than 1200 GeV since all investigated resonance mass point exhibit a discovery potential larger than five standard deviations. Table 8.4 summarises the range in which the linear fit parameters a and b are computed and their resulting values as well as the final coupling g_5 or mass m_5 for the respective case. The resonance mass at which a discovery potential of five standard deviations is obtained are 1464 GeV and 1320 GeV for the scalar isotensor resonance ϕ and tensor isotensor resonance t , respectively. Larger resonance masses have a smaller discovery potential. The resonance coupling parameter g_5 is 0.51 and 1.03 for ϕ and t , respectively. Smaller couplings exhibit smaller $S/\sqrt{B}$.

8.5.3 Comparison of the discovery potential for the different resonance types

The discovery potential as a function of the resonance mass and coupling, for both resonance types, scalar isotensor ϕ and tensor isotensor t , is compared in Figure 8.18. Due to the overall larger cross section for the ϕ resonance (compare Table 5.3), the discovery potential is larger for this type than for the t resonance. The Figure shows the comparison for the variable $m_{1\top}$, however results for other mass reconstruction variables show the same quantitative behaviour (compare the tables in Appendix D).

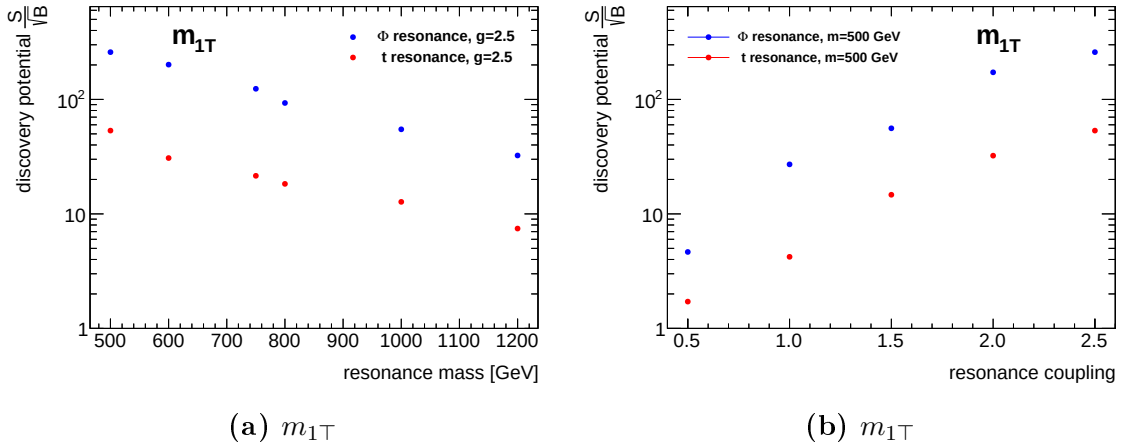


Figure 8.18: Discovery potential for the ϕ and t resonance as a function of the resonance mass (left) and the resonance coupling (right) for the mass bound variable $m_{1\top}$. Results for other mass reconstruction variables are comparable. The distributions contain signal events after $\vec{E}_T^{\text{miss}}$ smearing.

8.5.4 Comparison of the discovery potential for different event reconstruction levels

Three different event reconstruction levels are investigated, which are explained in Section 6.4. The reconstruction of selected events at parton level, particle level and after $\vec{E}_T^{\text{miss}}$ smearing is considered which differs between the event reconstruction levels. Figure 8.19 illustrates the comparison of the discovery potential at these three different stages of event reconstruction as a function of the resonance mass and coupling. For the ϕ and t resonance, one of the best performing mass variables for the resonance type at hand are chosen which are m_{o1} and $m_{1\top}$, respectively. Except for resonances with mass $m = 500$ GeV and coupling $g = 0.5$ and for the ϕ resonance with mass $m = 1200$ GeV and coupling $g = 2.5$, the parton level discovery potential is lower than the value at particle level or after $\vec{E}_T^{\text{miss}}$ smearing. This is due to the better event selection efficiency after the particle level simulation as explained in Section 8.3.3. The $\vec{E}_T^{\text{miss}}$ smearing has almost no impact with respect to the discovery potential at particle level. This is in accordance with the observations of Section 8.3.3 where no impact on the width of the distribution of the mass reconstruction variables at this level was seen.

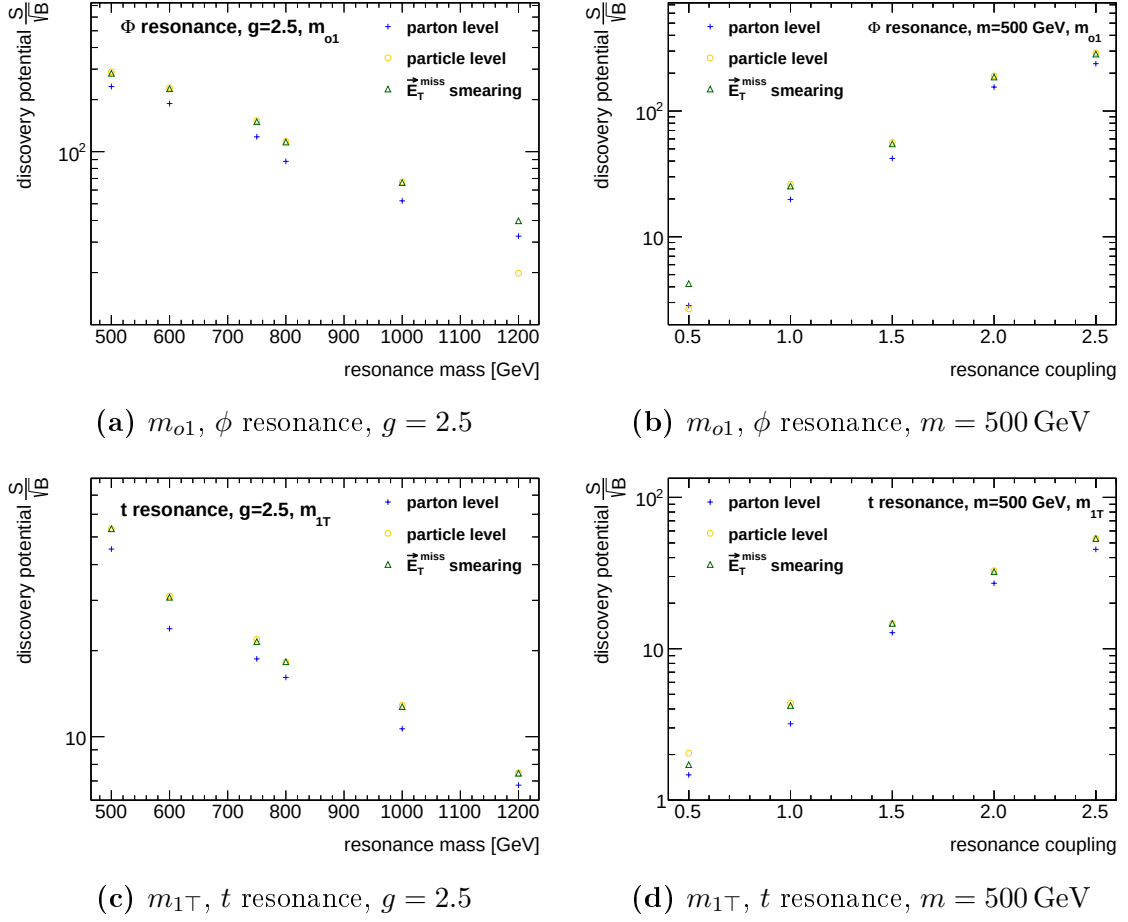


Figure 8.19: Discovery potential for the ϕ and t resonance as a function of the resonance mass (left) and the resonance coupling (right) for the three different event reconstruction levels: parton level (blue), particle level (yellow) and after $\vec{E}_T^{\text{miss}}$ smearing (green), for the best performing mass bound variables, m_{o1} for the ϕ resonance and m_{1T} for the t resonance. Results for other mass reconstruction variables are comparable.

8.6 Discussion and conclusion

The study of mass reconstruction variables for resonances in $W^\pm W^\pm$ scattering is conducted in the context of two different resonance types, the scalar isoscalar resonance ϕ with spin zero and isospin two, and the tensor isotensor resonance t with spin and isospin of two. Different mass reconstruction variables are investigated, the mass bound variables originating from the technique of constrained minimisation

$$m_{1o}, m_{o1}, m_{1T}, m_{1T\text{star}} \text{ and } m_{1T\text{bound}},$$

and traditional mass reconstruction variables formerly used in particle physics experiments

$$m_{\text{vis}}, m_{\text{eff}}, m_{\text{vec}}, \text{ and } m_{\text{col}}.$$

The investigation of the true invariant mass of the $W^\pm W^\pm$ system demonstrates the intrinsic properties of both resonance types for different parameter setups for the

resonance mass and coupling. While the t resonance shows the expected behaviour of a narrow resonance, the ϕ resonance exhibits a shift of the broad resonance peak towards smaller m_{WW}^{true} for resonance mass larger than 750 GeV.

The presented method for the signal extraction in a mass window of m_{WW}^{true} offers a precise instruction in order to select events with resonant contributions, which are suitable for an investigation of mass reconstruction variables. The content of the mass window is a free parameter which is set to 68 % in this work. Nevertheless, the impact on the variation of this parameter can be studied further. Especially for the narrow t resonance, a smaller m_{WW}^{true} mass window might be more appropriate for the signal event selection and might enhance the performance.

The distributions of the mass bound variables show the expected behaviour of representing a lower mass bound on the probed resonance mass due to the technique of constrained minimisation. Further constraints on the production of on-shell $W^\pm$ bosons alter the distributions and performance of the variable “ $\top$ ”-projected variables only slightly due to the very large probed resonance masses in comparison to the mass of the $W^\pm$ boson.

The traditional mass reconstruction variables are not designed to give a lower mass bound on the resonance mass. Hence, except the visible mass, each of these mass variables fail to give a lower mass bound. The collinear mass has a reconstruction efficiency of approximately 30 % in the full space of probed resonance parameters. It can be concluded that the assumption of collinear decay products does not hold for the topology of resonant $W^\pm W^\pm$ scattering. Therefore, m_{col} is discarded for all further studies.

Qualitatively, all probed resonance masses and couplings, and both resonance types reveal the same overall observations with respect to the performance of the mass reconstruction variables. However, a clear difference in the shape of the distributions of the mass variables can be discovered. A more pronounced investigation would show if these differences are solely explainable due to the influence of spin.

The comparison of the distributions of different mass reconstruction variables with respect to the event reconstruction levels exhibit only minor changes between the parton level and the reconstruction after $\vec{E}_T^{\text{miss}}$ smearing. The broadening of the distributions was expected for the event reconstruction level after $\vec{E}_T^{\text{miss}}$ smearing due to the deterioration of the energy resolution on the basis of the performance of the ATLAS detector. However, this is not observed. Nevertheless, clearly observable is the greater amount of events which overestimates the resonance mass for mass bound variables.

In addition, a linear dependency of the mean of the reconstructed invariant mass $\langle m_{WW}^{\text{reco}} \rangle$ on the mean of the true invariant mass $\langle m_{WW}^{\text{true}} \rangle$ is demonstrated for all explored mass variables. The dependency of $\langle m_{WW}^{\text{reco}} \rangle$ on the position of the maximum in the true invariant mass distribution, $m_{WW}^{\text{true max}}$, might be able to be used for specifying the mass of a resonance in a possible measurement of an excess of signal events in real data. The bias of the reconstructed mean is negative throughout all mass variables and probed resonance parameters and its magnitude rises with the mass of the resonance. The absolute width of the distributions of the mass reconstruction variables increases with the resonance mass. The variable m_{10} features the smallest width, however, as well is the weakest bound of all mass bound variables

and accumulates most events at low m_{WW}^{reco} . Due to the disuse of $\vec{E}_T^{\text{miss}}$, m_{vis} has as well a small width compared to other variables. Comparing the absolute width with respect to the resonance types shows a slightly larger width for the narrow t resonance than for the broad ϕ resonance. The relative width for all resonance masses ranges between 28 – 35 % for the best performing mass variables, the “ $\top$ ”-projected variables and m_{vec} .

The discovery potential for both resonance types and different mass and coupling parameters is investigated against the SM $W^\pm W^\pm \text{jj}$ -EW production at a centre-of-mass energy $\sqrt{s} = 8 \text{ TeV}$ and an integrated luminosity $\int \mathcal{L} dt = 20 \text{ fb}^{-1}$. As expected, the discovery potential decreases with the resonance mass and grows with the resonance coupling due to the behaviour of the corresponding cross section. The best performing mass reconstruction variables in terms of the discovery potential differ for the various points in the resonance parameter space. The variable m_{o1} has a superior discovery potential for ϕ resonances with large coupling due to the best background rejection in this case. For smaller couplings of the ϕ resonance, the “ $\top$ ”-family variables $m_{1\top}$ and $m_{1\top\text{bound}}$ and m_{vis} become best performing. For the t resonance picture is more complex. For reasonably high coupling the “ $\top$ ”-projected variables $m_{1\top}$ and $m_{1\top\text{bound}}$ are among the variables which are best performing in most of the cases. m_{eff} and m_{vec} give good results for the cases of resonance coupling $g = 2.5$ and masses $m = 500 \text{ GeV}$ and $m = 800 \text{ GeV}$, respectively. For smaller couplings the latter two and m_{vis} are best performing for one distinct coupling each.

It can be noticed that, except for the t resonance with small couplings, the mass bound variables are always amongst the best performing variables. Yet, different types of mass bounds are preferable for different types of resonances. Only for small coupling the mass-projected “ $\top$ ” variables are the best performing mass bound variables for both resonance types.

The overall discovery potential is better for the scalar isotensor resonance ϕ than for the tensor isotensor resonance t due to the larger cross section of the first. Comparing the discovery potential with respect to the event reconstruction levels shows again no great impact for the $\vec{E}_T^{\text{miss}}$ smearing.

The limits for the resonance mass and coupling for a discovery potential $S/\sqrt{B} = 5$ for the event reconstruction level after $\vec{E}_T^{\text{miss}}$ smearing and background contributions from only SM $W^\pm W^\pm \text{jj}$ -EW production are extracted. For a resonance coupling parameter of $g = 2.5$ the limits on the resonance mass for $S/\sqrt{B} = 5$ are 1464 GeV and 1320 GeV for the ϕ and t resonance, respectively. The limits on the resonance coupling for a resonance mass $m = 500 \text{ GeV}$ for ϕ and t are 0.51 and 1.03, respectively. Due to the overall superior discovery potential for the ϕ resonance, the limits for this resonance type are improved in comparison to the t resonance.

9 Conclusion

The objective of this thesis is to investigate mass reconstruction techniques for resonances in the scattering of like-charge $W^\pm$ bosons. Many physics scenarios, especially in the context of electroweak symmetry breaking, predict new resonances in massive, weak vector boson scattering which enhance the electroweak gauge boson scattering cross section. Recently, a first measurement of the cross section of the production of like-charge $W^\pm W^\pm$ in association with two jets in pp collisions was conducted [10]. The cross section for purely electroweak contributions to this process was found to be in agreement with the Standard Model prediction. However, the experimental uncertainty of more than 46 % still leaves room for physics beyond the Standard Model. Especially, the region of higher energy which is not yet in reach of the LHC could exhibit contributions from these new models.

This thesis presents the modelling of two types of heavy resonances in $W^\pm W^\pm$ scattering with the WHIZARD event generator, the scalar isotensor resonance ϕ and the tensor isotensor resonance t . The aim is the investigation of the reconstruction of the mass of such resonances. In this special context, the performance of mass reconstruction variables is studied. In addition to traditional mass reconstruction variables obtained from observable quantities in a particle detector, special emphasis is put on the study of mass bound variables originating from the mass reconstruction with the technique of constrained minimisation. These mass variables are investigated in the context of their mass reconstruction performance in the channel of $W^\pm W^\pm jj$ production. In particular their characteristic parameters in event distributions are investigated at the parton level event reconstruction level. Furthermore, the discovery potential for contributions exceeding the Standard model electroweak $W^\pm W^\pm jj$ production are discussed for pp collisions at a centre-of-mass energy $\sqrt{s} = 8 \text{ TeV}$ and an integrated luminosity $\int \mathcal{L} dt = 20 \text{ fb}^{-1}$. The discovery potential is investigated at the particle event reconstruction level after missing energy smearing.

The discovery potential falls with the mass of the resonance and grows with the resonance coupling parameter. The best performing mass reconstruction variables in terms of the discovery potential $S/\sqrt{B}$ are summarised in Table 9.1.

Different mass reconstruction variables are best performing for the various points in the resonance parameter space. Yet, it can be noticed that, except for the t resonance with small couplings, the mass bound variables are always amongst the best performing variables. However, different types of mass bounds are optimal for different types of resonances. The variable m_{o1} has a superior discovery potential for ϕ resonances with large coupling. For smaller couplings of the ϕ resonance, the “ $\top$ ”-family variables $m_{1\top}$ and $m_{1\top\text{bound}}$ become best performing. These are as well the variables which are among those which perform best for the case of a t resonance with reasonably high coupling. Only for smaller couplings the effective mass, m_{eff} , the vectorial mass, m_{vec} , and the visible mass, m_{vis} are best performing.

	resonance mass m [GeV]	coupling g	variables with best $S/\sqrt{B}$
ϕ	[500, 1200]	2.5	m_{o1}
	500	2.0	m_{o1}
	500	[0.5, 1.5]	$m_{1\top\text{bound}}, m_{1\top}, m_{\text{vis}}$
t	500	[600, 1200]	$m_{1\top\text{bound}}, m_{1\top}, m_{\text{vec}}$
		2.5	m_{eff}
		2.0	$m_{1\top\text{bound}}, m_{1\top}$
		1.5	$m_{\text{eff}}, m_{1\top}$
		1.0	m_{vis}
		0.5	m_{vec}

Table 9.1: Best performing mass reconstruction variables with respect to the discovery potential.

The overall discovery potential is shown to be superior for the ϕ resonance in comparison to the t resonance. In the context of these studies, limits in the best performing mass variable for the resonance mass and coupling parameters for a discovery potential of five are extracted. These limits are in the context of studies on Monte Carlo events at the particle event reconstruction level after the smearing of the missing transverse energy. In addition, only the background contribution from non-resonant Standard model electroweak $W^\pm W^\pm \text{jj}$ production is considered. The results are summarised in Table 9.2.

resonance	$m_5 = m(S/\sqrt{B} = 5)$ (for $g = 2.5$)	$g_5 = g(S/\sqrt{B} = 5)$ (for $m = 500$ GeV)
ϕ	1464 GeV	0.51
t	1320 GeV	1.03

Table 9.2: Summary of the resonance parameter extraction for a discovery potential $S/\sqrt{B} = 5$ for both resonance types.

A linear dependency of the mean of the reconstructed invariant mass $\langle m_{WW}^{\text{reco}} \rangle$ of the $W^\pm W^\pm$ system on the true invariant diboson mass for all mass variables is demonstrated. In addition, $\langle m_{WW}^{\text{reco}} \rangle$ reveals a dependency on the location of the maximum in the true invariant diboson mass spectrum which can be used to specify the mass of a resonance in a possible measurement of an excess of signal events over the expected background contributions. The relative width of the mass reconstruction is at 28 – 35 % for the variables with the smallest relative width, the “ $\top$ ”-projected variables and m_{vec} .

10 Outlook

In [10], limits on the anomalous quartic gauge coupling (aQGC) parameters α_4 and α_5 with the application of K-matrix unitarisation are derived. An interpretation of these limits in terms of the studied resonances shall be given as an outlook for further investigations. These aQGC parametrise the low energy effects of physics beyond the Standard Model in the context of an effective field theory approach. The relation between the aQGC parameters α_4 and α_5 and the ϕ and t resonance parameters is the following¹ [105]:

$$\alpha_{(4,5)} = \Delta\alpha_{(4,5)} \frac{\Gamma_{\text{res}_i} g_{\text{res}}^2}{M_{\text{res}}^2} \frac{v^2}{4}. \quad (10.1)$$

The scaled width of the resonance, Γ_{res_i} , and the contribution of the two resonances to the aQGC, $\Delta\alpha_{(4,5)}$, are listed in Table 10.1.

resonance	Γ_{res_i}	$\Delta\alpha_4$	$\Delta\alpha_5$
ϕ	1	0.25	-0.083
t	1/30	-0.625	4.375

Table 10.1: Scaled width and contributions to the anomalous quartic gauge coupling parameters α_4 and α_5 of the scalar isotensor resonance ϕ and the tensor isotensor resonance t [105].

The expected and observed two-dimensional confidence intervals in the (α_4, α_5) plane for the 95 % confidence level are shown in Figure 10.1. In addition, the contributions of the two resonances are marked with green lines which follow the values in Table 10.1. The intersection of these lines with the contours of the expected and observed 95 % confidence intervals on the aQGC parameters α_4 and α_5 give the two-dimensional limits in the aQGC parameter space with respect to the contributions of the studied resonances. Constraining one of the two resonance parameters, i.e. mass or coupling, limits on the respective other parameter can be obtained. Limits on the mass of the resonances can, therefore, be obtained by fixing the resonance coupling. The results with two example coupling parameters, $g = 1$ and $g = 2.5$, are listed in Table 10.2. Due to the narrow width of the t resonance, the corresponding limits on the resonance mass are weaker than for the ϕ resonance.

¹This relation is derived without the application of any unitarisation model. However, this shall suffice in this context of an outlook.

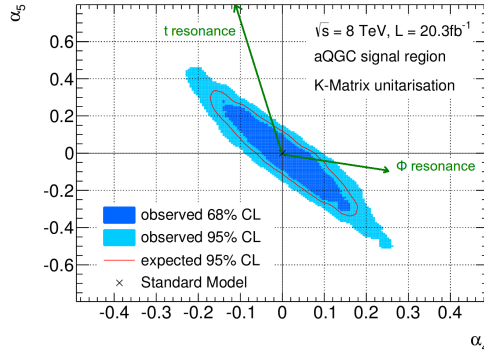


Figure 10.1: Two-dimensional confidence intervals on the aQGC parameters α_4 and α_5 . The area outside the given curve/filled region is excluded at the respective confidence level. The expected exclusion contour at the 95 % confidence level is marked by the red curve [106]. The contributions of the scalar isotensor resonance ϕ and the tensor isotensor resonance t to the aQGC are marked as green lines.

resonance	limits on (α_4, α_5)		limits on mass m [GeV]	
	at 95 % confidence level		coupling $g = 1$	coupling $g = 2.5$
ϕ	expected	$(0.09, -0.03)$	205	513
	observed	$(0.12, -0.035)$	189	473
t	expected	$(-0.021, 0.15)$	122	305
	observed	$(-0.027, 0.19)$	118	295

Table 10.2: Expected and observed two-dimensional limits on the aQGC parameters α_4 and α_5 with respect to contributions from the two resonance types ϕ and t . These limits are retrieved from Figure 10.1. They are interpreted as limits on the resonance mass for two fixed coupling parameter points, $g = 1$ and $g = 2.5$.

Further studies in the context of mass reconstruction of resonances in like-charge $W^\pm W^\pm$ scattering may include an examination of the performance of the mass reconstruction variables with the full ATLAS detector simulation. This allows for a direct comparison with real collision data. Furthermore, additional major background contributions from like-charge $W^\pm W^\pm jj$ -QCD and $W^\pm Z$ scattering could be included in the determination of the discovery potential. Moreover, the influence of the spin of the resonance on the mass reconstruction variables can be investigated with a fit to the spectrum of the reconstructed invariant diboson mass. Furthermore, these studies can be conducted with an increased centre-of-mass energy of $\sqrt{s} = 13$ TeV of the LHC from 2015 on.

A Truth mass window m_{WW}^{true} construction

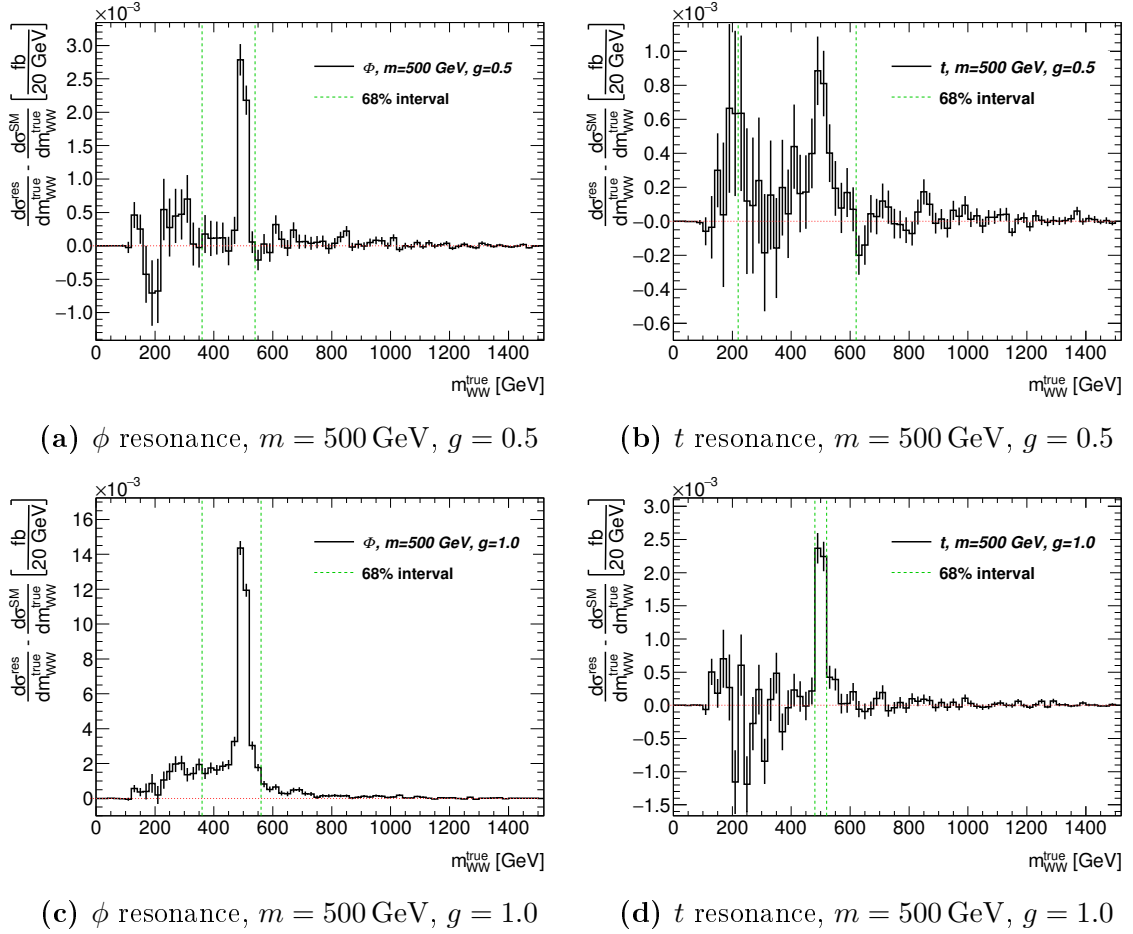
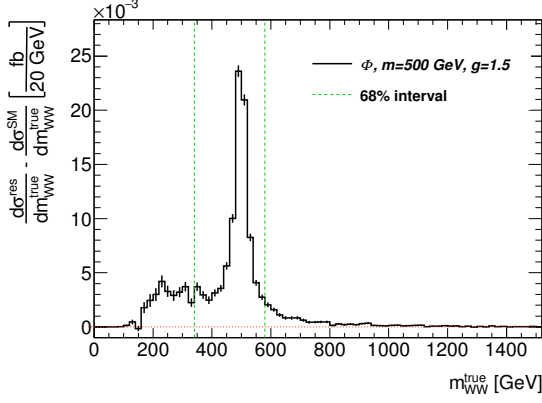
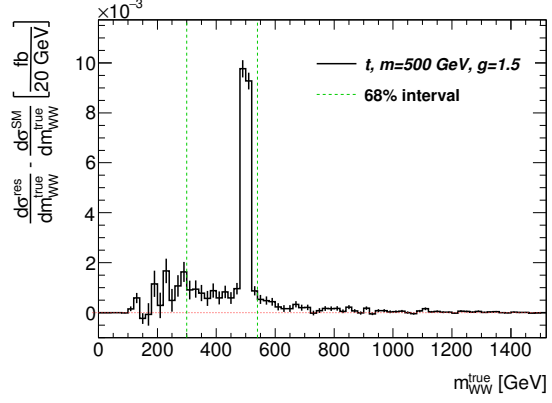


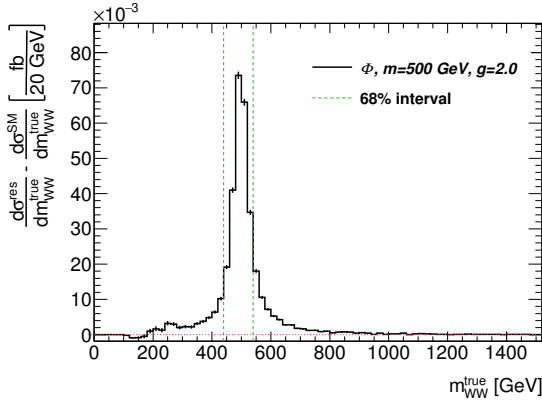
Figure A.1: Signal extraction in the true invariant diboson mass at parton level m_{WW}^{true} for the ϕ (left) and t (right) resonance of mass $m = 500$ GeV and coupling $g \in \{0.5, 1.0\}$. The distribution in m_{WW}^{true} of the resonance samples was subtracted by the m_{WW}^{true} SM $W^\pm W^\pm$ jj-EW distribution. The window in m_{WW}^{true} (green, dashed) comprises 68% of the remaining subtracted contribution. The distributions are normalised to the corresponding signal cross section $\sigma^{\text{signal}}(m, g)$ from Equation (8.3).



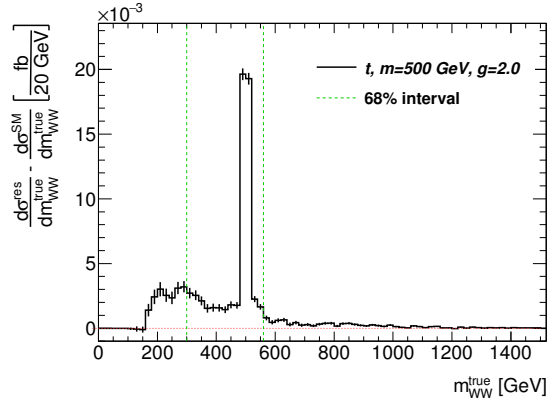
(a) ϕ resonance, $m = 500$ GeV, $g = 1.5$



(b) t resonance, $m = 500$ GeV, $g = 1.5$

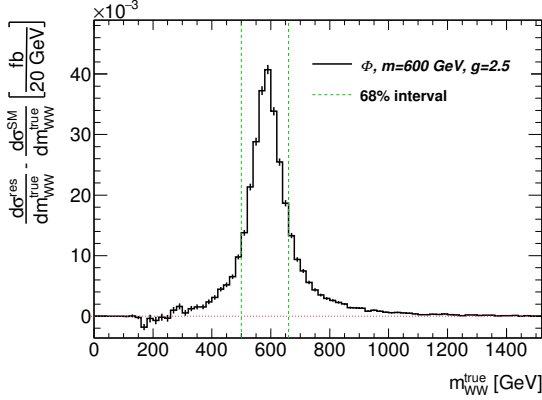


(c) ϕ resonance, $m = 500$ GeV, $g = 2.0$

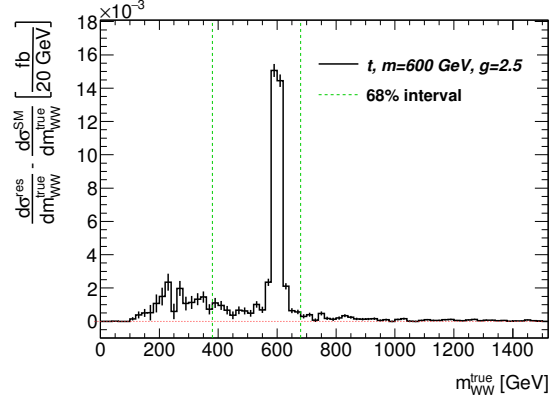


(d) t resonance, $m = 500$ GeV, $g = 2.0$

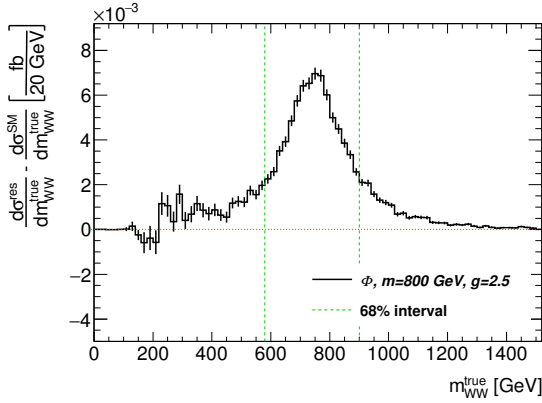
Figure A.2: Signal extraction in the true invariant diboson mass at parton level m_{WW}^{true} for the ϕ (left) and t (right) resonance of mass $m = 500$ GeV and coupling $g \in \{1.5, 2.0\}$. The distribution in m_{WW}^{true} of the resonance samples was subtracted by the m_{WW}^{true} SM $W^\pm W^\pm \text{jj}$ -EW distribution. The window in m_{WW}^{true} (green, dashed) comprises 68% of the remaining subtracted contribution. The distributions are normalised to the corresponding signal cross section $\sigma^{\text{signal}}(m, g)$ from Equation (8.3).



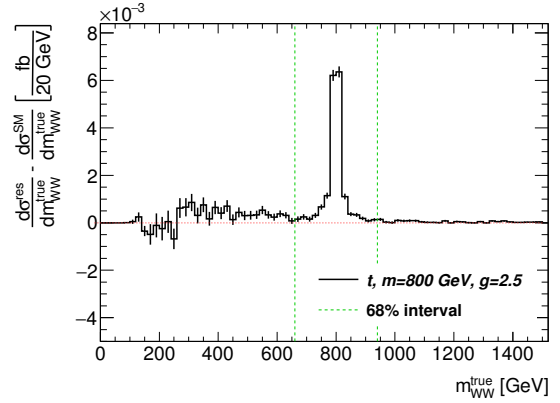
(a) ϕ resonance, $m = 600$ GeV, $g = 2.5$



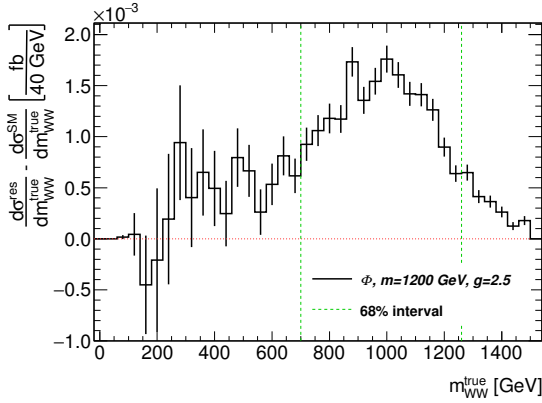
(b) t resonance, $m = 600$ GeV, $g = 2.5$



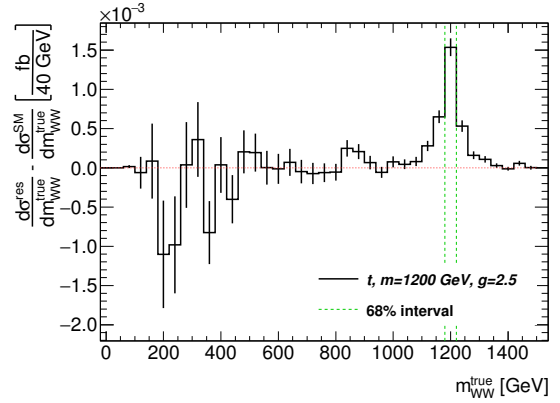
(c) ϕ resonance, $m = 800$ GeV, $g = 2.5$



(d) t resonance, $m = 800$ GeV, $g = 2.5$



(e) ϕ resonance, $m = 1200$ GeV, $g = 2.5$



(f) t resonance, $m = 1200$ GeV, $g = 2.5$

Figure A.3: Signal extraction in the true invariant diboson mass at parton level m_{WW}^{true} for the ϕ (left) and t (right) resonance for masses $m \in \{600 \text{ GeV}, 800 \text{ GeV}, 1200 \text{ GeV}\}$ and coupling $g = 2.5$. The distribution in m_{WW}^{true} of the resonance samples was subtracted by the m_{WW}^{true} SM $W^\pm W^\pm \text{jj-EW}$ distribution. The window in m_{WW}^{true} (green, dashed) comprises 68 % of the remaining subtracted contribution. The distributions are normalised to the corresponding signal cross section $\sigma^{\text{signal}}(m, g)$ from Equation (8.3).

B Distributions of mass reconstruction variables

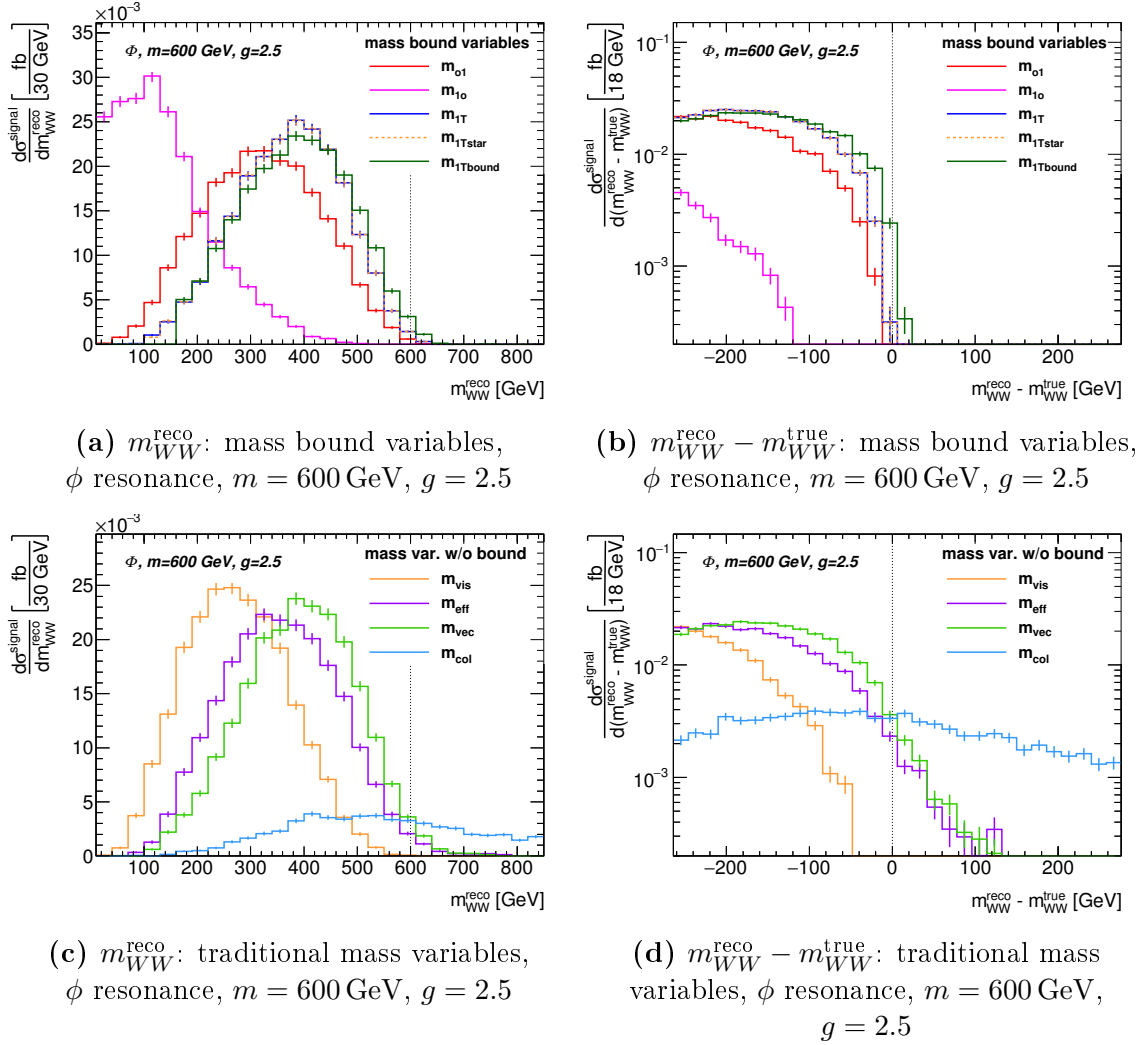
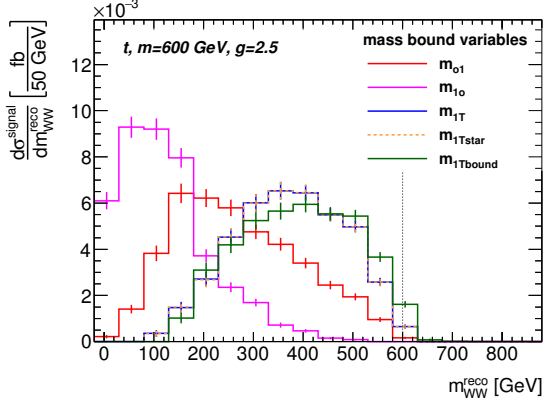
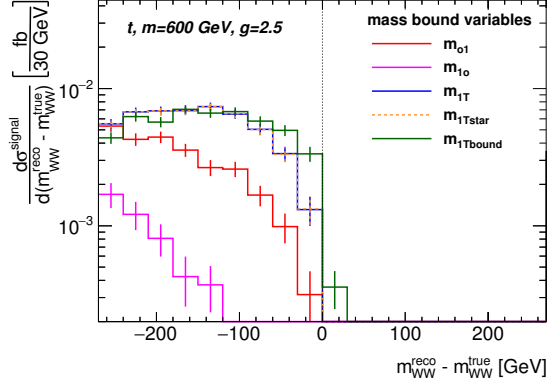


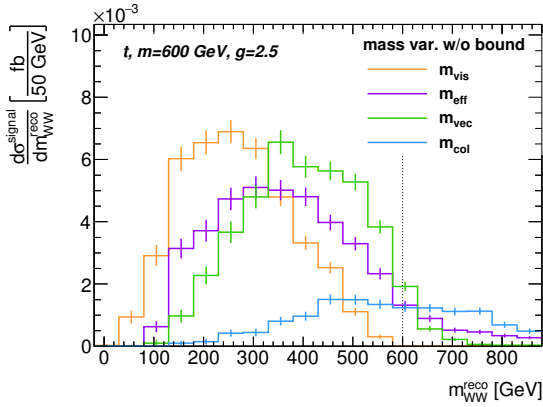
Figure B.1: Reconstructed invariant diboson mass m_{WW}^{reco} (left) and the difference between m_{WW}^{reco} and m_{WW}^{true} (right) at parton level for signal events in the true invariant mass window for the ϕ resonance with mass $m = 600$ GeV and coupling $g = 2.5$ for mass bound variables. The distributions are normalised to the corresponding signal cross section $\sigma^{\text{signal}}(m, g)$ from Equation (8.3).



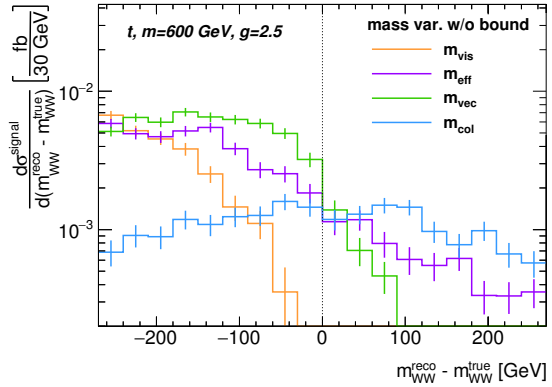
(a) m_{WW}^{reco} : mass bound variables, t resonance, $m = 600$ GeV, $g = 2.5$



(b) $m_{WW}^{\text{reco}} - m_{WW}^{\text{true}}$: mass bound variables, t resonance, $m = 600$ GeV, $g = 2.5$



(c) m_{WW}^{reco} : traditional mass variables, t resonance, $m = 600$ GeV, $g = 2.5$



(d) $m_{WW}^{\text{reco}} - m_{WW}^{\text{true}}$: traditional mass variables, t resonance, $m = 600$ GeV, $g = 2.5$

Figure B.2: Reconstructed invariant diboson mass m_{WW}^{reco} (left) and the difference between m_{WW}^{reco} and m_{WW}^{true} (right) at parton level for signal events in the true invariant mass window for the t resonance with mass $m = 600$ GeV and coupling $g = 2.5$ for mass bound variables. The distributions are normalised to the corresponding signal cross section $\sigma^{\text{signal}}(m, g)$ from Equation (8.3).

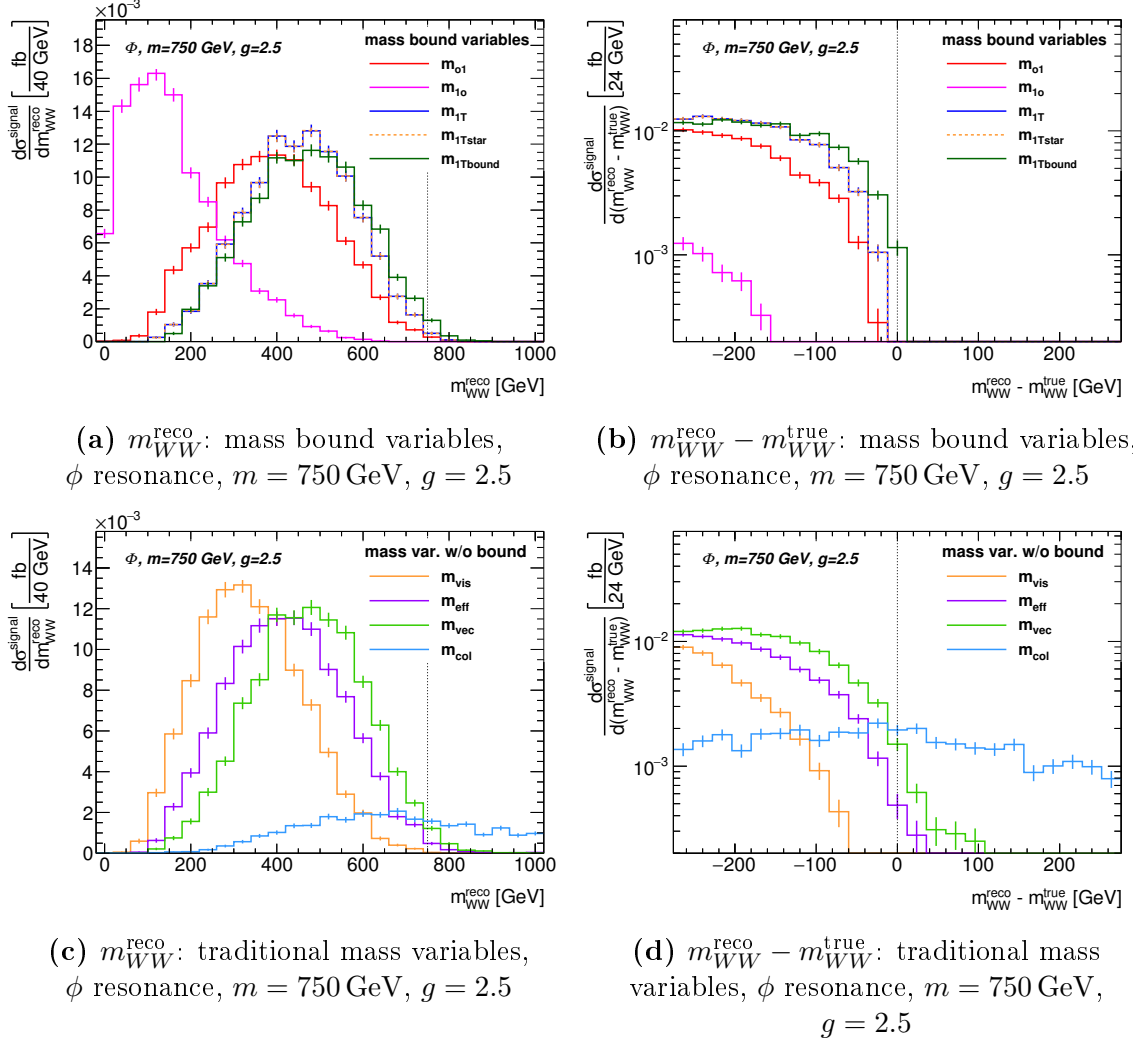
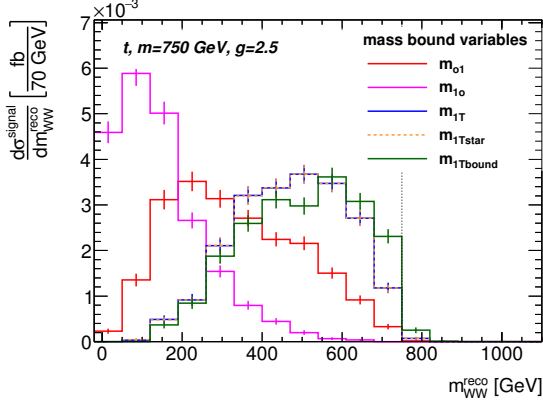
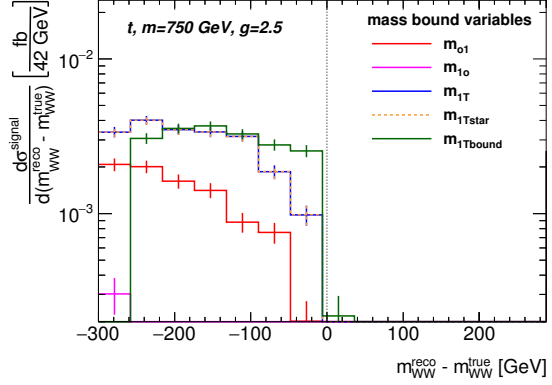


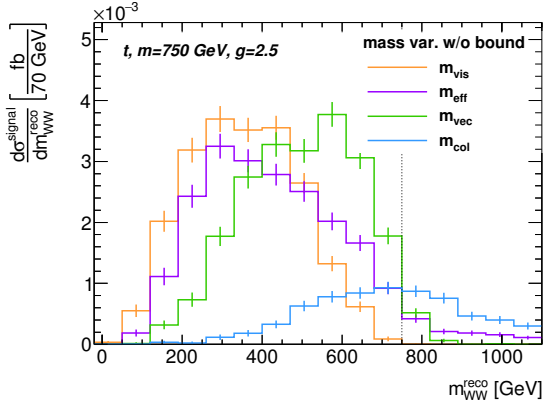
Figure B.3: Reconstructed invariant diboson mass m_{WW}^{reco} (left) and the difference between m_{WW}^{reco} and m_{WW}^{true} (right) at parton level for signal events in the true invariant mass window for the ϕ resonance with mass $m = 750$ GeV and coupling $g = 2.5$ for mass bound variables. The distributions are normalised to the corresponding signal cross section $\sigma^{\text{signal}}(m, g)$ from Equation (8.3).



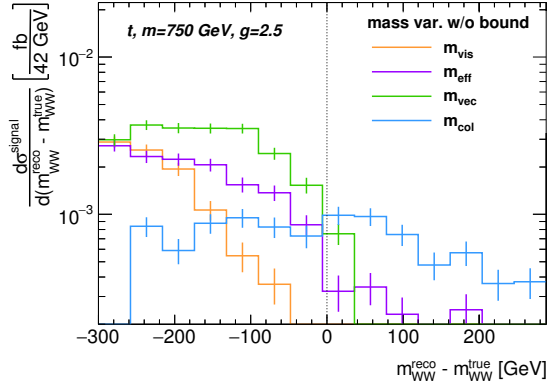
(a) m_{WW}^{reco} : mass bound variables, t resonance, $m = 750$ GeV, $g = 2.5$



(b) $m_{WW}^{\text{reco}} - m_{WW}^{\text{true}}$: mass bound variables, t resonance, $m = 750$ GeV, $g = 2.5$



(c) m_{WW}^{reco} : traditional mass variables, t resonance, $m = 750$ GeV, $g = 2.5$



(d) $m_{WW}^{\text{reco}} - m_{WW}^{\text{true}}$: traditional mass variables, t resonance, $m = 750$ GeV, $g = 2.5$

Figure B.4: Reconstructed invariant diboson mass m_{WW}^{reco} (left) and the difference between m_{WW}^{reco} and m_{WW}^{true} (right) at parton level for signal events in the true invariant mass window for the t resonance with mass $m = 750$ GeV and coupling $g = 2.5$ for mass bound variables. The distributions are normalised to the corresponding signal cross section $\sigma^{\text{signal}}(m, g)$ from Equation (8.3).

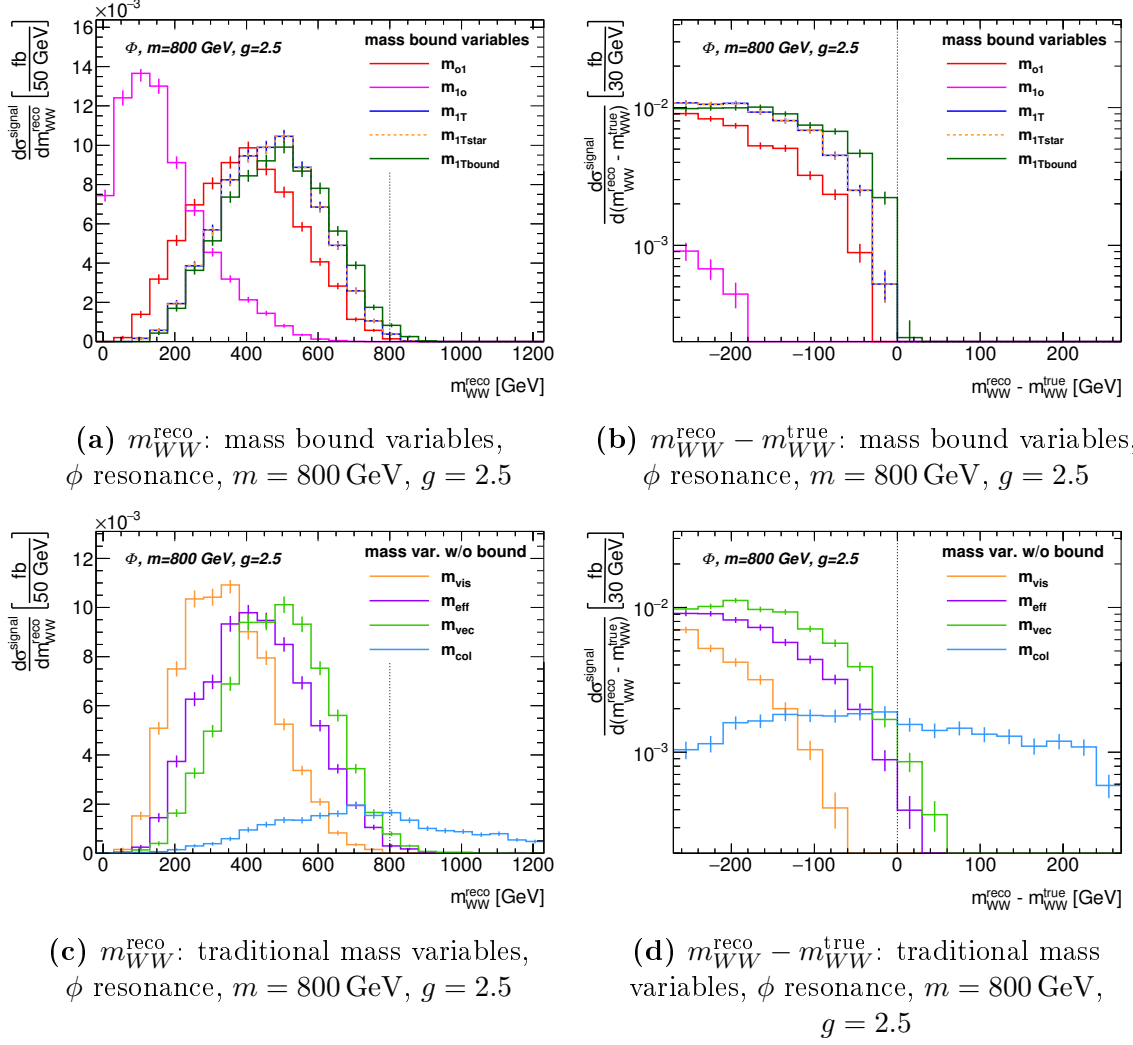
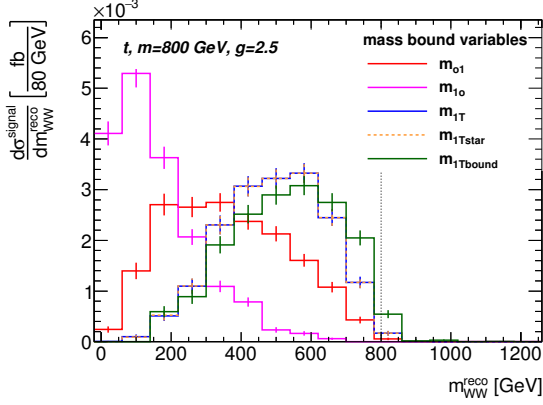
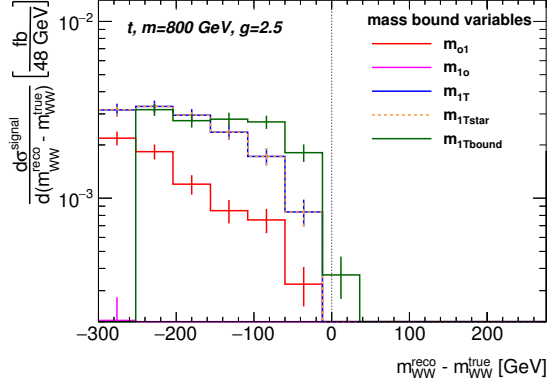


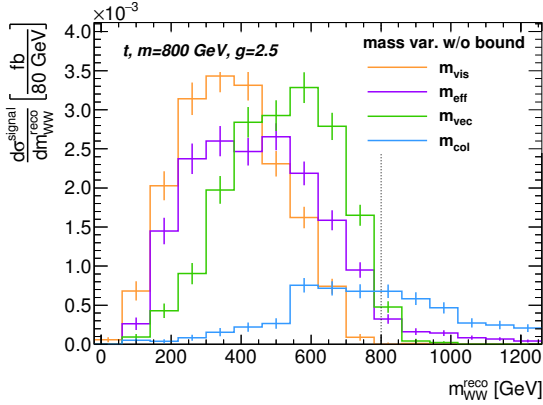
Figure B.5: Reconstructed invariant diboson mass m_{WW}^{reco} (left) and the difference between m_{WW}^{reco} and m_{WW}^{true} (right) at parton level for signal events in the true invariant mass window for the ϕ resonance with mass $m = 800$ GeV and coupling $g = 2.5$ for mass bound variables. The distributions are normalised to the corresponding signal cross section $\sigma^{\text{signal}}(m, g)$ from Equation (8.3).



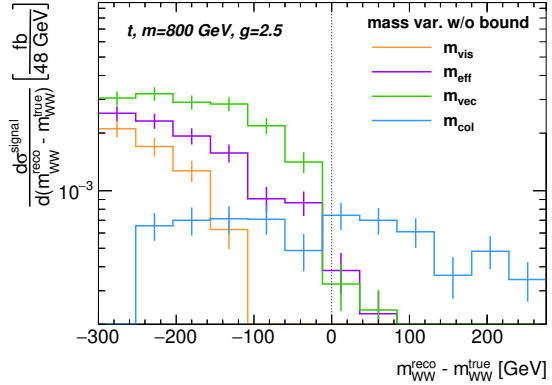
(a) m_{WW}^{reco} : mass bound variables, t resonance, $m = 800$ GeV, $g = 2.5$



(b) $m_{WW}^{\text{reco}} - m_{WW}^{\text{true}}$: mass bound variables, t resonance, $m = 800$ GeV, $g = 2.5$

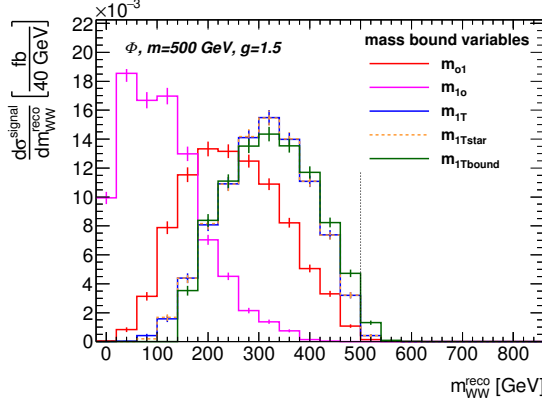


(c) m_{WW}^{reco} : traditional mass variables, t resonance, $m = 800$ GeV, $g = 2.5$

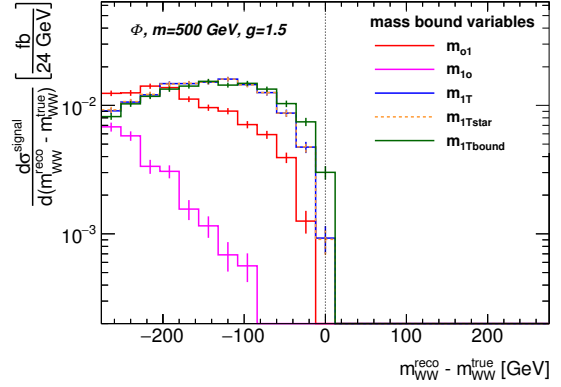


(d) $m_{WW}^{\text{reco}} - m_{WW}^{\text{true}}$: traditional mass variables, t resonance, $m = 800$ GeV, $g = 2.5$

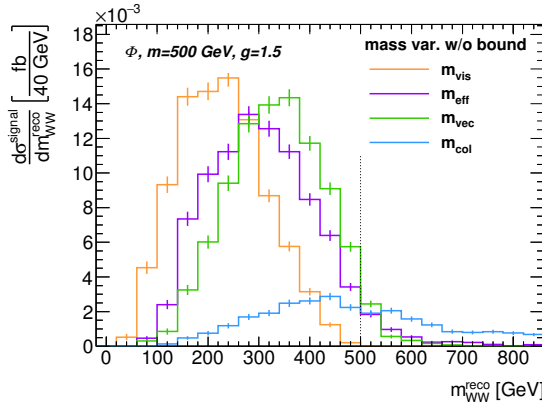
Figure B.6: Reconstructed invariant diboson mass m_{WW}^{reco} (left) and the difference between m_{WW}^{reco} and m_{WW}^{true} (right) at parton level for signal events in the true invariant mass window for the t resonance with mass $m = 800$ GeV and coupling $g = 2.5$ for mass bound variables. The distributions are normalised to the corresponding signal cross section $\sigma^{\text{signal}}(m, g)$ from Equation (8.3).



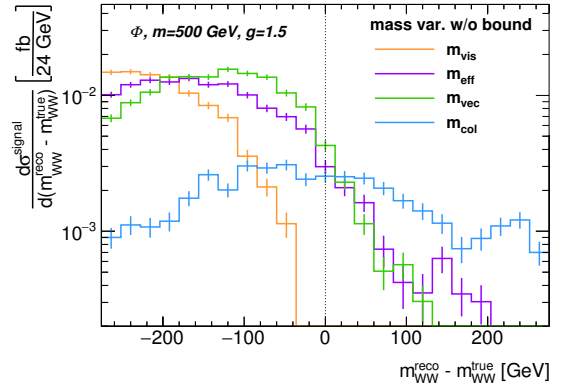
(a) m_{WW}^{reco} : mass bound variables, ϕ resonance, $m = 500$ GeV, $g = 1.5$



(b) $m_{WW}^{\text{reco}} - m_{WW}^{\text{true}}$: mass bound variables, ϕ resonance, $m = 500$ GeV, $g = 1.5$

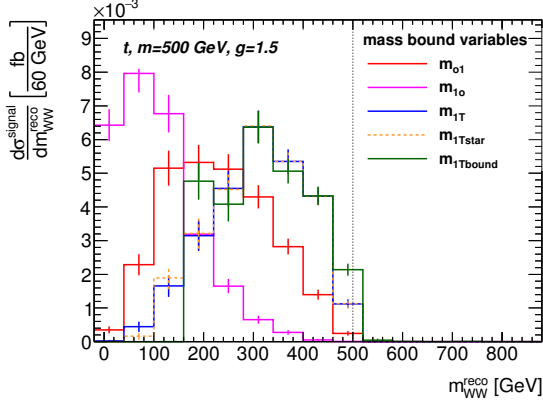


(c) m_{WW}^{reco} : traditional mass variables, ϕ resonance, $m = 500$ GeV, $g = 1.5$

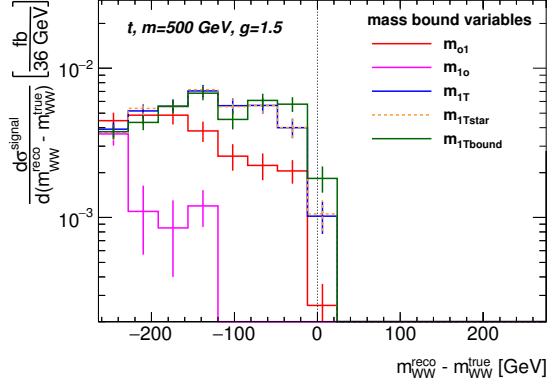


(d) $m_{WW}^{\text{reco}} - m_{WW}^{\text{true}}$: traditional mass variables, ϕ resonance, $m = 500$ GeV, $g = 1.5$

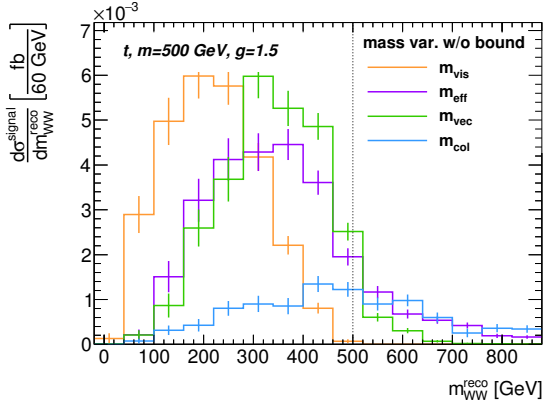
Figure B.7: Reconstructed invariant diboson mass m_{WW}^{reco} (left) and the difference between m_{WW}^{reco} and m_{WW}^{true} (right) at parton level for signal events in the true invariant mass window for the ϕ resonance with mass $m = 500$ GeV and coupling $g = 1.5$ for mass bound variables. The distributions are normalised to the corresponding signal cross section $\sigma^{\text{signal}}(m, g)$ from Equation (8.3).



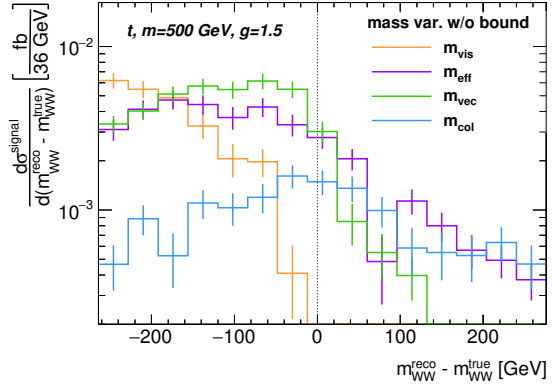
(a) m_{WW}^{reco} : mass bound variables, t resonance, $m = 500$ GeV, $g = 1.5$



(b) $m_{WW}^{\text{reco}} - m_{WW}^{\text{true}}$: mass bound variables, t resonance, $m = 500$ GeV, $g = 1.5$

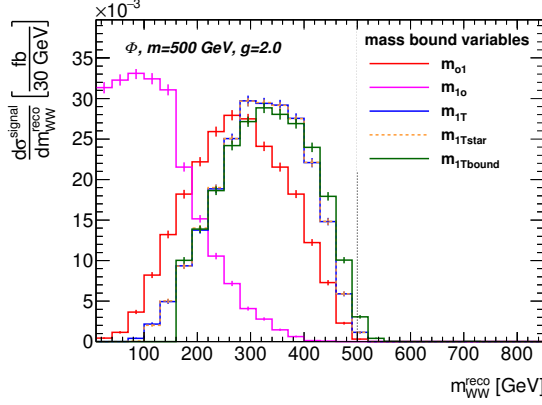


(c) m_{WW}^{reco} : traditional mass variables, t resonance, $m = 500$ GeV, $g = 1.5$

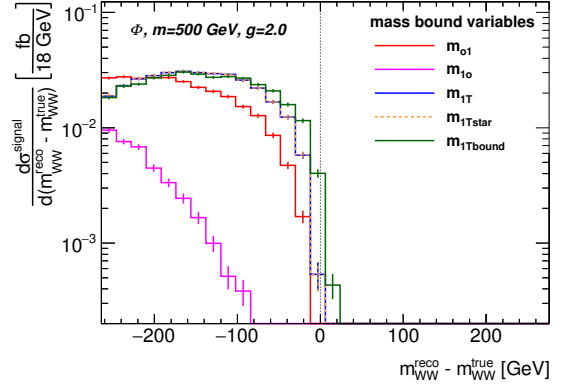


(d) $m_{WW}^{\text{reco}} - m_{WW}^{\text{true}}$: traditional mass variables, t resonance, $m = 500$ GeV, $g = 1.5$

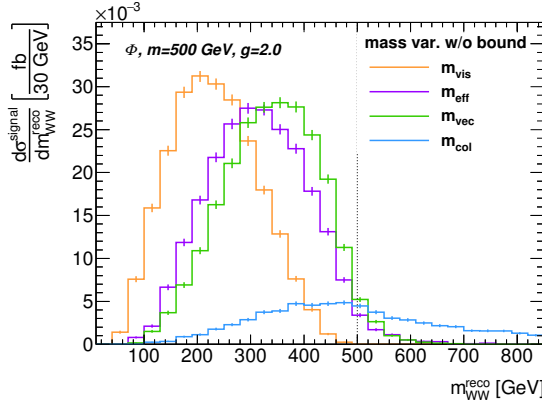
Figure B.8: Reconstructed invariant diboson mass m_{WW}^{reco} (left) and the difference between m_{WW}^{reco} and m_{WW}^{true} (right) at parton level for signal events in the true invariant mass window for the t resonance with mass $m = 500$ GeV and coupling $g = 1.5$ for mass bound variables. The distributions are normalised to the corresponding signal cross section $\sigma^{\text{signal}}(m, g)$ from Equation (8.3).



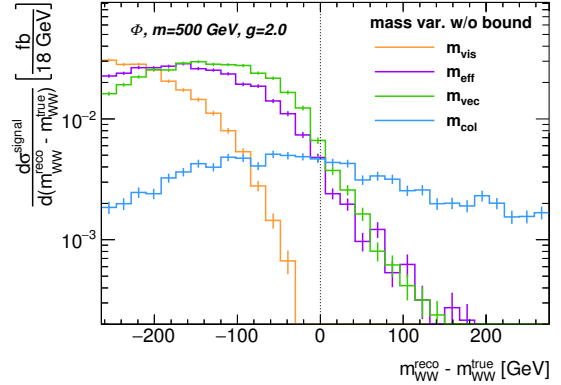
(a) m_{WW}^{reco} : mass bound variables, ϕ resonance, $m = 500$ GeV, $g = 2.0$



(b) $m_{WW}^{\text{reco}} - m_{WW}^{\text{true}}$: mass bound variables, ϕ resonance, $m = 500$ GeV, $g = 2.0$

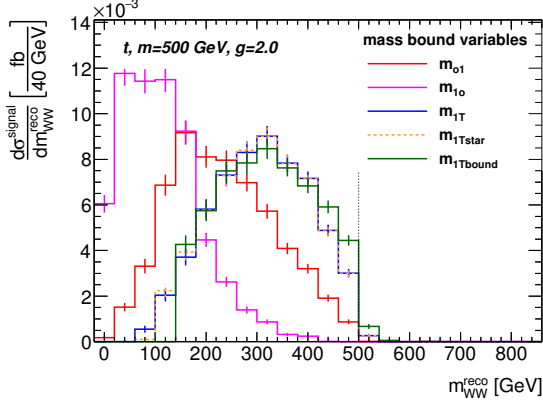


(c) m_{WW}^{reco} : traditional mass variables, ϕ resonance, $m = 500$ GeV, $g = 2.0$

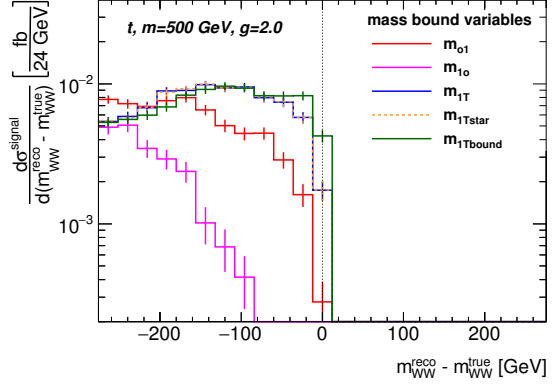


(d) $m_{WW}^{\text{reco}} - m_{WW}^{\text{true}}$: traditional mass variables, ϕ resonance, $m = 500$ GeV, $g = 2.0$

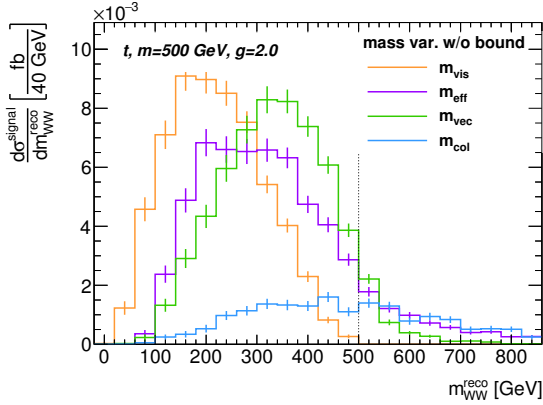
Figure B.9: Reconstructed invariant diboson mass m_{WW}^{reco} (left) and the difference between m_{WW}^{reco} and m_{WW}^{true} (right) at parton level for signal events in the true invariant mass window for the ϕ resonance with mass $m = 500$ GeV and coupling $g = 2.0$ for mass bound variables. The distributions are normalised to the corresponding signal cross section $\sigma^{\text{signal}}(m, g)$ from Equation (8.3).



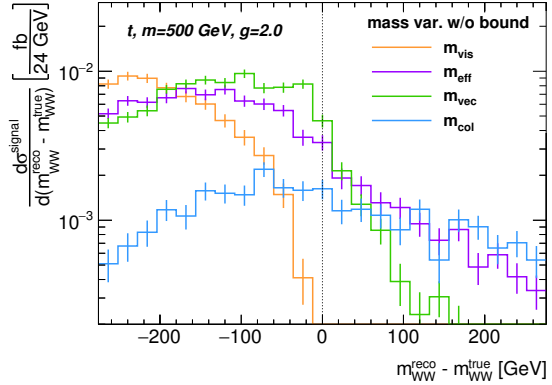
(a) m_{WW}^{reco} : mass bound variables, t resonance, $m = 500$ GeV, $g = 2.0$



(b) $m_{WW}^{\text{reco}} - m_{WW}^{\text{true}}$: mass bound variables, t resonance, $m = 500$ GeV, $g = 2.0$



(c) m_{WW}^{reco} : traditional mass variables, t resonance, $m = 500$ GeV, $g = 2.0$



(d) $m_{WW}^{\text{reco}} - m_{WW}^{\text{true}}$: traditional mass variables, t resonance, $m = 500$ GeV, $g = 2.0$

Figure B.10: Reconstructed invariant diboson mass m_{WW}^{reco} (left) and the difference between m_{WW}^{reco} and m_{WW}^{true} (right) at parton level for signal events in the true invariant mass window for the t resonance with mass $m = 500$ GeV and coupling $g = 2.0$ for mass bound variables. The distributions are normalised to the corresponding signal cross section $\sigma^{\text{signal}}(m, g)$ from Equation (8.3).

C Characteristic parameters of the mass reconstruction variables

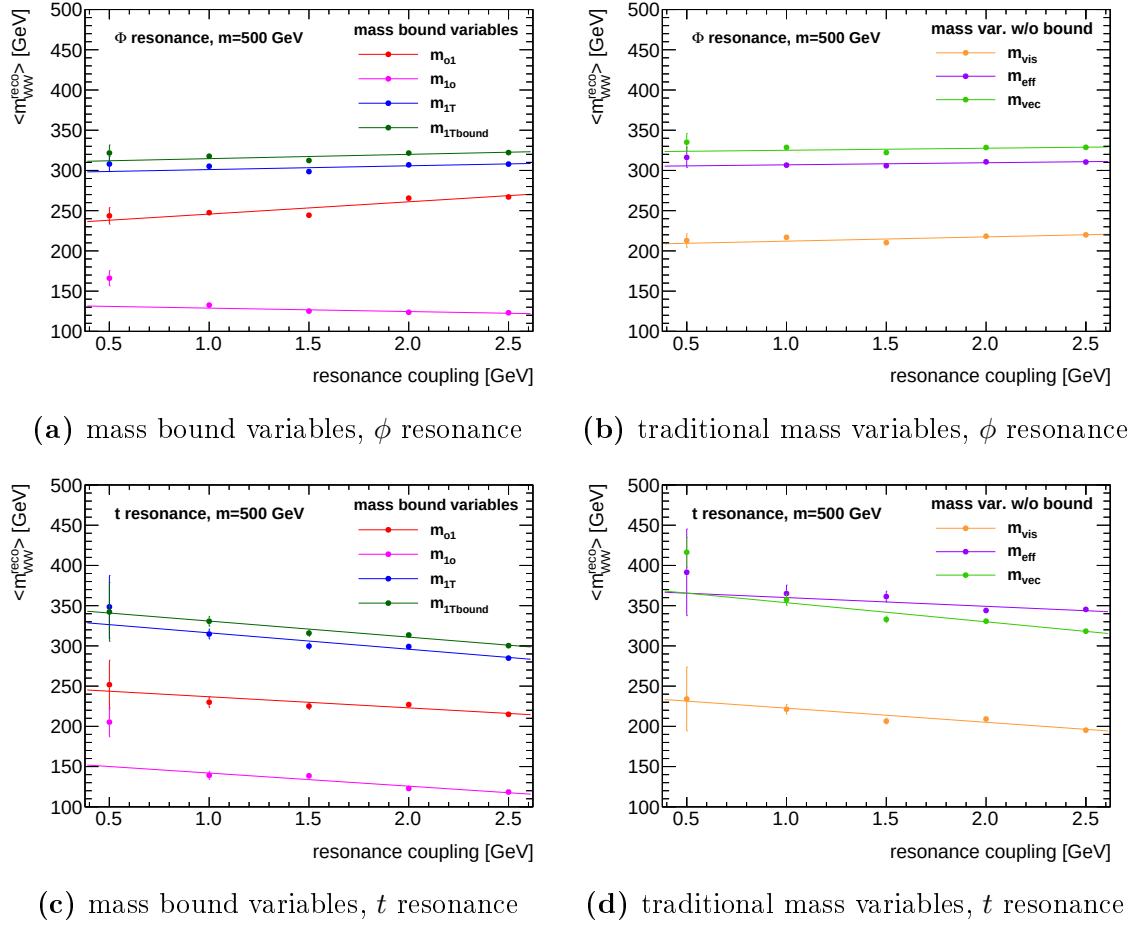
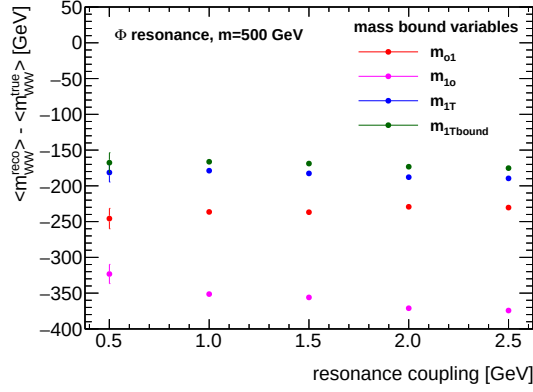


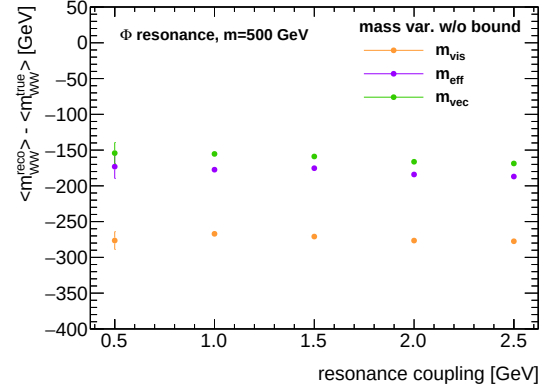
Figure C.1: Mean of the distribution of the reconstructed invariant diboson mass $\langle m_{WW}^{\text{reco}} \rangle$ as a function of the resonance coupling parameter for signal events in the true invariant mass window and after $\vec{E}_T^{\text{miss}}$ smearing. The ϕ and t resonance with mass $m = 500$ GeV are depicted for mass bound variables (left) and traditional mass reconstruction variables (right). The distributions are each fitted with a linear function with parameters listed in Table C.1.

Resonance	m_{WW}^{reco} variable	linear fit parameters		
		a	b [GeV]	χ^2/NDF
$\phi, g = 2.5$	m_{o1}	15.3 ± 1.0	230 ± 2	111.4/3
	m_{1o}	-4.3 ± 0.8	133 ± 2	25.1/3
	$m_{1\top}$	4.7 ± 0.9	296 ± 2	27.5/3
	$m_{1\top\text{bound}}$	5.3 ± 0.9	309 ± 2	28.8/3
	m_{vis}	5.3 ± 0.9	207 ± 2	27.3/3
	m_{eff}	2.6 ± 1.1	304 ± 2	6.8/3
	m_{vec}	2.5 ± 1.0	323 ± 2	16.5/3
$t, g = 2.5$	m_{o1}	-13.8 ± 2.9	251 ± 6	7.2/3
	m_{1o}	-16.3 ± 2.2	158 ± 5	15.8/3
	$m_{1\top}$	-20.4 ± 2.7	337 ± 6	6.7/3
	$m_{1\top\text{bound}}$	-20.0 ± 2.7	351 ± 6	3.8/3
	m_{vis}	-17.5 ± 2.6	240 ± 6	10.0/3
	m_{eff}	-10.9 ± 4.8	371 ± 11	4.3/3
	m_{vec}	-23.8 ± 3.0	377 ± 7	12.5/3

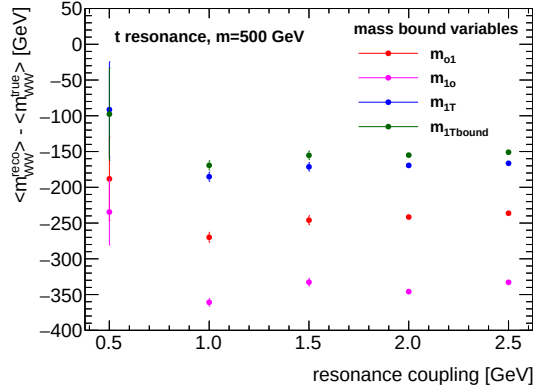
Table C.1: Parameter of the linear fit of the dependency of $\langle m_{WW}^{\text{reco}} \rangle$ on resonance coupling parameter according to Figure C.1 for resonances of both types with coupling $m = 500 \text{ GeV}$. “NDF” denotes the number of degrees of freedom for the fit.



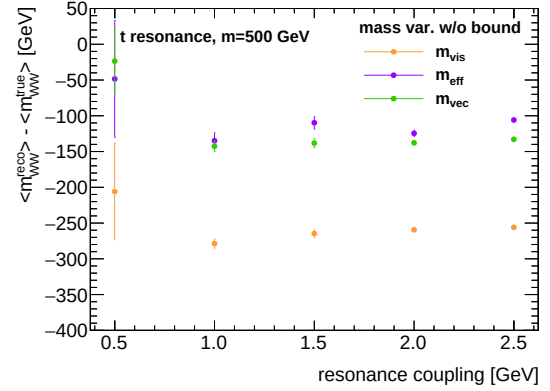
(a) mass bound variables, ϕ resonance



(b) traditional mass variables, ϕ resonance

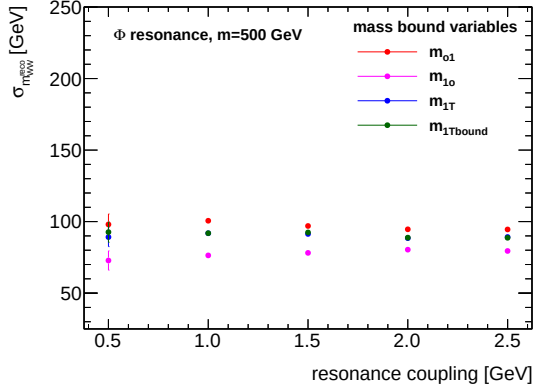


(c) mass bound variables, t resonance

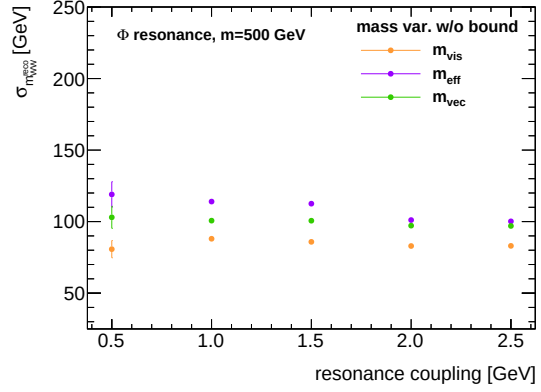


(d) traditional mass variables, t resonance

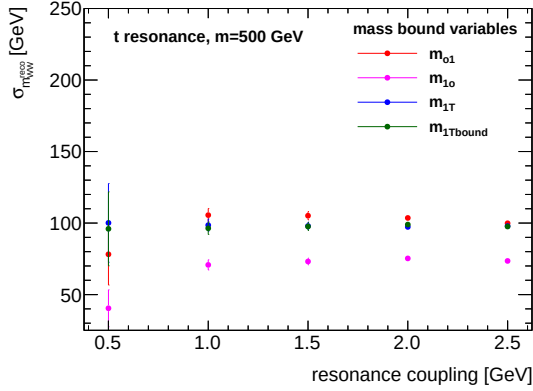
Figure C.2: Corrected bias $\langle m_{WW}^{\text{reco}} \rangle - \langle m_{WW}^{\text{true}} \rangle$ as a function of the resonance coupling parameter for signal events in the true invariant mass window and after $\vec{E}_T^{\text{miss}}$ smearing. The ϕ (a, b) and t (c, d) resonance with mass $m = 500 \text{ GeV}$ are depicted for mass bound variables (left) and traditional mass reconstruction variables (right).



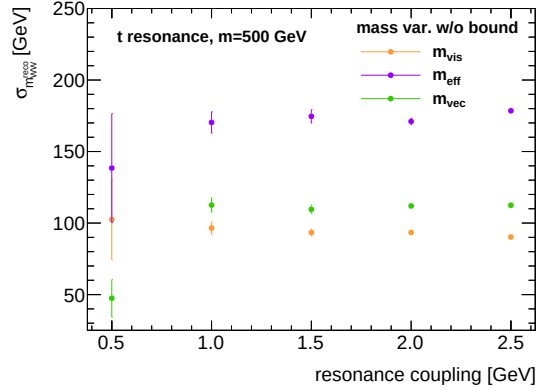
(a) mass bound variables, ϕ resonance



(b) traditional mass variables, ϕ resonance

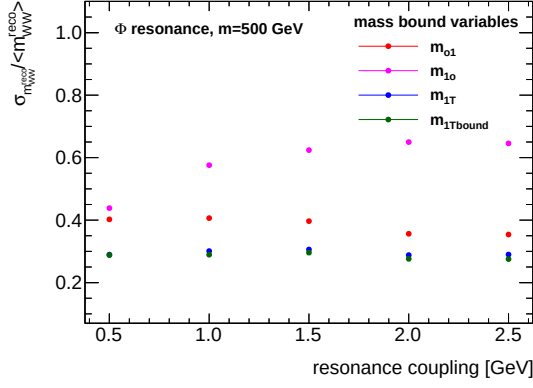


(c) mass bound variables, t resonance

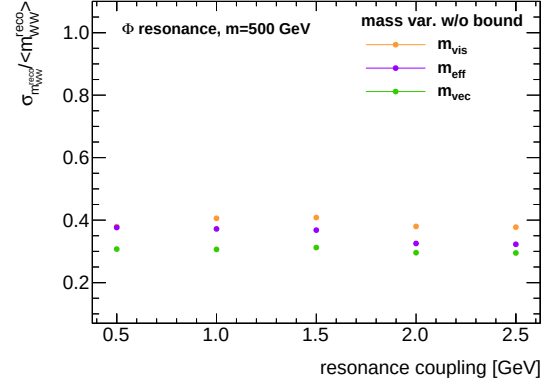


(d) traditional mass variables, t resonance

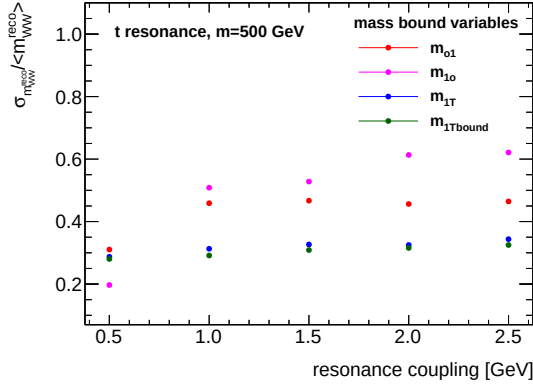
Figure C.3: Width of the reconstructed diboson mass $\sigma_{m_{WW}^{reco}}$ as a function of the resonance coupling parameter for signal events in the true invariant mass window and after $\vec{E}_T^{\text{miss}}$ smearing. The ϕ (a, b) and t (c, d) resonance with mass $m = 500$ GeV are depicted for mass bound variables (left) and traditional mass reconstruction variables (right).



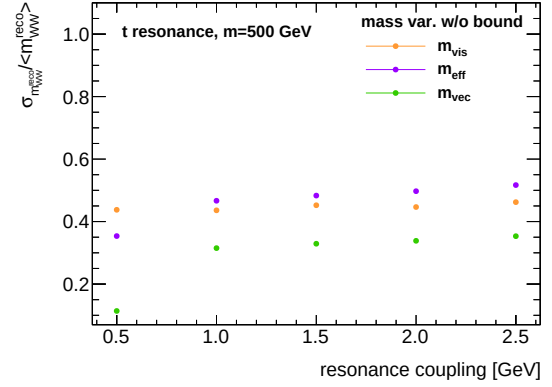
(a) mass bound variables, ϕ resonance



(b) traditional mass variables, ϕ resonance



(c) mass bound variables, t resonance



(d) traditional mass variables, t resonance

Figure C.4: Relative width of the reconstructed diboson mass $\sigma_{m_{WW}^{\text{reco}}}$ as a function of the resonance coupling parameter for mass bound variables (left) and traditional mass reconstruction variables (right) for signal events in the true invariant mass window m_{WW}^{true} for the ϕ (a, b) and t (c, d) resonance with mass $m = 500$ GeV after $\vec{E}_T^{\text{miss}}$ smearing.

Resonance	m_{WW}^{reco} variable	$m_{WW}^{\text{reco max}}$ [GeV]	$\langle m_{WW}^{\text{reco}} \rangle$ [GeV]	$\sigma_{m_{WW}^{\text{reco}}}$ [GeV]	$\frac{\sigma_{m_{WW}^{\text{reco}}}}{\langle m_{WW}^{\text{reco}} \rangle}$	bias [GeV]	bias corr. [GeV]
ϕ , $m = 500 \text{ GeV}$, $g = 2.5$	m_{vis}	205	219.9	83.0	0.38	-280.1	-277.5
	m_{eff}	295	310.5	100.2	0.32	-189.5	-187.0
	m_{vec}	325	328.7	97.0	0.29	-171.3	-168.7
	m_{o1}	265	267.1	94.5	0.35	-233.0	-230.4
	m_{1o}	55	123.1	79.5	0.65	-376.9	-374.4
	$m_{1\top}$	325	307.9	89.2	0.29	-192.1	-189.6
	$m_{1\top\text{bound}}$	325	322.2	88.7	0.28	-177.8	-175.2
ϕ , $m = 600 \text{ GeV}$, $g = 2.5$	m_{vis}	235	257.3	98.7	0.38	-342.7	-326.0
	m_{eff}	355	349.6	114.8	0.33	-250.3	-233.7
	m_{vec}	385	377.1	112.4	0.30	-222.9	-206.2
	m_{o1}	295	308.7	112.1	0.36	-291.3	-274.6
	m_{1o}	85	140.9	94.2	0.67	-459.1	-442.4
	$m_{1\top}$	385	356.9	105.2	0.29	-243.1	-226.5
	$m_{1\top\text{bound}}$	355	372.3	109.3	0.29	-227.7	-211.0
ϕ , $m = 750 \text{ GeV}$, $g = 2.5$	m_{vis}	240	315.1	125.0	0.40	-435.0	-405.8
	m_{eff}	360	410.9	139.6	0.34	-339.1	-309.9
	m_{vec}	440	452.9	139.6	0.31	-297.1	-267.9
	m_{o1}	400	374.1	139.6	0.37	-375.9	-346.7
	m_{1o}	80	175.6	117.0	0.67	-574.4	-545.2
	$m_{1\top}$	440	433.2	133.2	0.31	-316.8	-287.6
	$m_{1\top\text{bound}}$	480	452.9	144.2	0.32	-297.1	-268.0
ϕ , $m = 800 \text{ GeV}$, $g = 2.5$	m_{vis}	255	319.5	133.2	0.42	-480.5	-424.2
	m_{eff}	405	422.3	148.8	0.35	-377.7	-321.4
	m_{vec}	455	463.6	148.0	0.32	-336.4	-280.2
	m_{o1}	405	382.6	149.2	0.39	-417.5	-361.2
	m_{1o}	105	189.8	120.7	0.64	-610.2	-553.9
	$m_{1\top}$	455	443.4	142.3	0.32	-356.6	-300.3
	$m_{1\top\text{bound}}$	505	463.5	153.4	0.33	-336.5	-280.2
ϕ , $m = 1000 \text{ GeV}$, $g = 2.5$	m_{vis}	225	373.2	163.6	0.44	-626.8	-491.7
	m_{eff}	365	472.0	178.1	0.38	-528.1	-393.0
	m_{vec}	575	528.9	181.6	0.34	-471.1	-336.0
	m_{o1}	365	435.9	181.6	0.42	-564.1	-429.1
	m_{1o}	85	225.5	139.3	0.62	-774.5	-639.4
	$m_{1\top}$	435	509.1	175.8	0.35	-490.9	-355.8
	$m_{1\top\text{bound}}$	435	533.5	190.1	0.36	-466.5	-331.5
ϕ , $m = 1200 \text{ GeV}$, $g = 2.5$	m_{vis}	365	423.0	183.4	0.43	-777.0	-553.3
	m_{eff}	505	516.3	199.0	0.39	-683.7	-459.9
	m_{vec}	645	592.6	198.6	0.34	-607.4	-383.7
	m_{o1}	365	482.7	205.2	0.43	-717.3	-493.5
	m_{1o}	85	244.7	159.6	0.65	-955.3	-731.5
	$m_{1\top}$	645	573.5	194.4	0.34	-626.5	-402.8
	$m_{1\top\text{bound}}$	645	602.9	211.8	0.35	-597.1	-373.3

Table C.2: Characteristic parameters for the ϕ resonance with variable mass $m \in \{500 \text{ GeV}, 600 \text{ GeV}, 750 \text{ GeV}, 800 \text{ GeV}, 1000 \text{ GeV}, 1200 \text{ GeV}\}$ and coupling $g = 2.5$: Position of maximum $m_{WW}^{\text{reco max}}$, reconstructed mean $\langle m_{WW}^{\text{reco}} \rangle$, width $\sigma_{m_{WW}^{\text{reco}}}$, relative width $\frac{\sigma_{m_{WW}^{\text{reco}}}}{\langle m_{WW}^{\text{reco}} \rangle}$, bias $\langle m_{WW}^{\text{reco}} \rangle - m_{\text{resonance}}$, and corrected bias $\langle m_{WW}^{\text{reco}} \rangle - \langle m_{WW}^{\text{true}} \rangle$.

Resonance	m_{WW}^{reco} variable	$m_{WW}^{\text{reco max}}$ [GeV]	$\langle m_{WW}^{\text{reco}} \rangle$ [GeV]	$\sigma_{m_{WW}^{\text{reco}}}$ [GeV]	$\frac{\sigma_{m_{WW}^{\text{reco}}}}{\langle m_{WW}^{\text{reco}} \rangle}$	bias [GeV]	bias corr. [GeV]
t , $m = 500 \text{ GeV}$, $g = 2.5$	m_{vis}	160	195.3	90.3	0.46	-304.7	-255.9
	m_{eff}	280	345.4	178.5	0.52	-154.7	-105.9
	m_{vec}	280	318.3	112.4	0.35	-181.7	-133.0
	m_{o1}	200	215.0	99.9	0.46	-285.0	-236.3
	m_{1o}	40	118.4	73.5	0.62	-381.7	-332.9
	$m_{1\top}$	240	284.8	97.9	0.34	-215.2	-166.4
	$m_{1\top\text{bound}}$	240	300.3	97.6	0.32	-199.7	-151.0
t , $m = 600 \text{ GeV}$, $g = 2.5$	m_{vis}	205	256.6	109.1	0.43	-343.4	-323.8
	m_{eff}	305	390.5	191.1	0.49	-209.5	-189.9
	m_{vec}	355	392.1	124.6	0.32	-207.9	-188.3
	m_{o1}	155	261.0	125.7	0.48	-339.0	-319.4
	m_{1o}	55	144.0	90.5	0.63	-456.0	-436.4
	$m_{1\top}$	355	360.7	113.2	0.31	-239.3	-219.8
	$m_{1\top\text{bound}}$	355	376.6	119.7	0.32	-223.4	-203.8
t , $m = 750 \text{ GeV}$, $g = 2.5$	m_{vis}	295	333.2	140.0	0.42	-416.8	-415.9
	m_{eff}	295	437.6	199.5	0.46	-312.4	-311.5
	m_{vec}	435	482.6	150.9	0.31	-267.4	-266.5
	m_{o1}	225	327.4	163.3	0.50	-422.6	-421.8
	m_{1o}	85	180.8	109.1	0.60	-569.2	-568.3
	$m_{1\top}$	505	453.9	144.6	0.32	-296.1	-295.2
	$m_{1\top\text{bound}}$	575	481.2	155.6	0.32	-268.8	-267.9
t , $m = 800 \text{ GeV}$, $g = 2.5$	m_{vis}	260	346.5	147.2	0.42	-453.5	-451.7
	m_{eff}	260	450.7	198.3	0.44	-349.3	-347.5
	m_{vec}	420	500.0	162.3	0.32	-300.0	-298.2
	m_{o1}	260	350.5	175.4	0.50	-449.5	-447.7
	m_{1o}	100	203.4	123.0	0.60	-596.7	-594.9
	$m_{1\top}$	420	472.7	155.2	0.33	-327.3	-325.5
	$m_{1\top\text{bound}}$	500	503.2	171.1	0.34	-296.8	-295.1
t , $m = 1000 \text{ GeV}$, $g = 2.5$	m_{vis}	430	443.4	189.4	0.43	-556.6	-562.9
	m_{eff}	430	509.3	233.7	0.46	-490.7	-497.0
	m_{vec}	630	610.5	207.2	0.34	-389.5	-395.9
	m_{o1}	230	433.2	227.0	0.52	-566.8	-573.1
	m_{1o}	130	242.5	149.1	0.61	-757.5	-763.8
	$m_{1\top}$	630	587.2	200.3	0.34	-412.8	-419.1
	$m_{1\top\text{bound}}$	630	630.6	223.0	0.35	-369.4	-375.7
t , $m = 1200 \text{ GeV}$, $g = 2.5$	m_{vis}	470	486.8	220.9	0.45	-713.3	-713.1
	m_{eff}	470	533.5	243.6	0.46	-666.5	-666.4
	m_{vec}	750	681.6	228.2	0.33	-518.4	-518.2
	m_{o1}	470	459.2	253.9	0.55	-740.8	-740.6
	m_{1o}	190	315.8	160.2	0.51	-884.2	-884.0
	$m_{1\top}$	610	653.9	223.6	0.34	-546.1	-545.9
	$m_{1\top\text{bound}}$	750	724.8	245.9	0.34	-475.2	-475.0

Table C.3: Characteristic parameters for the t resonance with variable mass $m \in \{500 \text{ GeV}, 600 \text{ GeV}, 750 \text{ GeV}, 800 \text{ GeV}, 1000 \text{ GeV}, 1200 \text{ GeV}\}$ and coupling $g = 2.5$: Position of maximum $m_{WW}^{\text{reco max}}$, reconstructed mean $\langle m_{WW}^{\text{reco}} \rangle$, width $\sigma_{m_{WW}^{\text{reco}}}$, relative width $\frac{\sigma_{m_{WW}^{\text{reco}}}}{\langle m_{WW}^{\text{reco}} \rangle}$, bias $\langle m_{WW}^{\text{reco}} \rangle - m_{\text{resonance}}$, and corrected bias $\langle m_{WW}^{\text{reco}} \rangle - \langle m_{WW}^{\text{true}} \rangle$.

Resonance	m_{WW}^{reco} variable	$m_{WW}^{\text{reco max}}$ [GeV]	$\langle m_{WW}^{\text{reco}} \rangle$ [GeV]	$\sigma_{m_{WW}^{\text{reco}}}$ [GeV]	$\frac{\sigma_{m_{WW}^{\text{reco}}}}{\langle m_{WW}^{\text{reco}} \rangle}$	bias [GeV]	bias corr. [GeV]
ϕ , $m = 500$ GeV, $g = 0.5$	m_{vis}	180	212.9	80.7	0.38	-287.1	-276.5
	m_{eff}	260	316.2	119.0	0.38	-183.8	-173.1
	m_{vec}	340	335.1	103.0	0.31	-164.9	-154.2
	m_{o1}	260	243.6	98.0	0.40	-256.4	-245.8
	m_{1o}	100	166.1	72.8	0.44	-333.9	-323.3
	$m_{1\top}$	340	308.0	89.2	0.29	-192.0	-181.4
	$m_{1\top\text{bound}}$	340	321.7	92.7	0.29	-178.3	-167.6
ϕ , $m = 500$ GeV, $g = 1.0$	m_{vis}	205	216.9	88.0	0.41	-283.1	-267.1
	m_{eff}	255	306.6	114.0	0.37	-193.4	-177.4
	m_{vec}	305	328.6	100.7	0.31	-171.4	-155.4
	m_{o1}	205	247.5	100.6	0.41	-252.5	-236.5
	m_{1o}	55	132.6	76.4	0.58	-367.4	-351.4
	$m_{1\top}$	305	305.2	91.9	0.30	-194.8	-178.8
	$m_{1\top\text{bound}}$	305	317.7	91.9	0.29	-182.3	-166.3
ϕ , $m = 500$ GeV, $g = 1.5$	m_{vis}	160	210.23	85.8	0.41	-289.7	-270.9
	m_{eff}	280	305.9	112.5	0.37	-194.2	-175.3
	m_{vec}	320	322.2	100.6	0.31	-177.8	-159.0
	m_{o1}	200	244.3	96.9	0.40	-255.7	-236.9
	m_{1o}	40	125.2	78.1	0.62	-374.8	-356.0
	$m_{1\top}$	320	298.6	91.4	0.31	-201.4	-182.6
	$m_{1\top\text{bound}}$	320	312.4	92.4	0.30	-187.7	-168.9
ϕ , $m = 500$ GeV, $g = 2.0$	m_{vis}	175	218.3	82.9	0.38	-281.7	-276.5
	m_{eff}	295	310.7	101.1	0.33	-189.3	-184.1
	m_{vec}	325	328.52	97.2	0.30	-171.5	-166.3
	m_{o1}	265	265.6	94.6	0.36	-234.5	-229.3
	m_{1o}	55	123.7	80.4	0.65	-376.3	-371.1
	$m_{1\top}$	295	306.9	88.4	0.29	-193.1	-187.9
	$m_{1\top\text{bound}}$	295	321.6	88.8	0.28	-178.4	-173.3

Table C.4: Characteristic parameters for the ϕ resonance with mass $m = 500$ GeV and variable coupling $g \in \{0.5, 1.0, 1.5, 2.0\}$: Position of maximum $m_{WW}^{\text{reco max}}$, reconstructed mean $\langle m_{WW}^{\text{reco}} \rangle$, width $\sigma_{m_{WW}^{\text{reco}}}$, relative width $\frac{\sigma_{m_{WW}^{\text{reco}}}}{\langle m_{WW}^{\text{reco}} \rangle}$, bias $\langle m_{WW}^{\text{reco}} \rangle - m_{\text{resonance}}$, and corrected bias $\langle m_{WW}^{\text{reco}} \rangle - \langle m_{WW}^{\text{true}} \rangle$.

Resonance	m_{WW}^{reco} variable	$m_{WW}^{\text{reco max}}$ [GeV]	$\langle m_{WW}^{\text{reco}} \rangle$ [GeV]	$\sigma_{m_{WW}^{\text{reco}}}$ [GeV]	$\frac{\sigma_{m_{WW}^{\text{reco}}}}{\langle m_{WW}^{\text{reco}} \rangle}$	bias [GeV]	bias corr. [GeV]
t , $m = 500 \text{ GeV}$, $g = 0.5$	m_{vis}	190	234.1	102.5	0.44	-265.9	-205.8
	m_{eff}	330	391.5	138.4	0.35	-108.5	-48.4
	m_{vec}	330	416.4	47.4	0.11	-83.6	-23.5
	m_{o1}	190	251.8	78.2	0.31	-248.3	-188.2
	m_{1o}	190	205.3	40.4	0.20	-294.7	-234.6
	$m_{1\top}$	330	348.6	100.1	0.29	-151.5	-91.4
	$m_{1\top\text{bound}}$	330	342.3	96.0	0.28	-157.8	-97.7
t , $m = 500 \text{ GeV}$, $g = 1.0$	m_{vis}	250	221.3	96.5	0.44	-278.7	-278.7
	m_{eff}	250	365.1	170.4	0.47	-134.9	-134.9
	m_{vec}	370	357.4	112.6	0.32	-142.7	-142.6
	m_{o1}	250	230.0	105.6	0.46	-270.0	-269.9
	m_{1o}	70	139.2	70.8	0.51	-360.8	-360.8
	$m_{1\top}$	310	314.8	98.5	0.31	-185.2	-185.2
	$m_{1\top\text{bound}}$	310	330.6	96.3	0.29	-169.4	-169.3
t , $m = 500 \text{ GeV}$, $g = 1.5$	m_{vis}	190	206.4	93.4	0.45	-293.6	-264.7
	m_{eff}	250	361.4	174.6	0.48	-138.6	-109.8
	m_{vec}	310	332.9	109.6	0.33	-167.1	-138.2
	m_{o1}	190	225.2	105.1	0.47	-274.8	-246.0
	m_{1o}	70	138.5	73.1	0.53	-361.5	-332.7
	$m_{1\top}$	310	299.8	98.0	0.33	-200.2	-171.4
	$m_{1\top\text{bound}}$	310	315.9	97.6	0.31	-184.1	-155.3
t , $m = 500 \text{ GeV}$, $g = 2.0$	m_{vis}	160	209.1	93.4	0.45	-290.9	-259.5
	m_{eff}	240	344.0	171.1	0.50	-156.0	-124.6
	m_{vec}	280	330.8	112.0	0.34	-169.2	-137.8
	m_{o1}	160	227.0	103.6	0.46	-273.0	-241.6
	m_{1o}	40	122.8	75.3	0.61	-377.3	-345.9
	$m_{1\top}$	280	299.2	97.3	0.33	-200.8	-169.4
	$m_{1\top\text{bound}}$	280	313.6	99.0	0.32	-186.4	-155.0

Table C.5: Characteristic parameters for the t resonance with variable mass $m = 500 \text{ GeV}$ and variable coupling $g \in \{0.5, 1.0, 1.5, 2.0\}$: Position of maximum $m_{WW}^{\text{reco max}}$, reconstructed mean $\langle m_{WW}^{\text{reco}} \rangle$, width $\sigma_{m_{WW}^{\text{reco}}}$, relative width $\frac{\sigma_{m_{WW}^{\text{reco}}}}{\langle m_{WW}^{\text{reco}} \rangle}$, bias $\langle m_{WW}^{\text{reco}} \rangle - m_{\text{resonance}}$, and corrected bias $\langle m_{WW}^{\text{reco}} \rangle - \langle m_{WW}^{\text{true}} \rangle$.

D Discovery Potential of the mass reconstruction variables

Resonance	m_{WW}^{reco} variable	S	B	discovery potential	best window in m_{WW}^{reco} [GeV]
ϕ , $m = 500$ GeV, $g = 0.5$	m_{vis}	204	2096	4.45	(140,420)
	m_{eff}	62	373	3.24	(380,500)
	m_{vec}	14	26	2.78	(580,620)
	m_{o1}	144	1151	4.24	(220,620)
	m_{1o}	202	5053	2.85	(20,180)
	$m_{1\top}$	207	1984	4.65	(220,500)
	$m_{1\top\text{bound}}$	199	2065	4.37	(220,500)
ϕ , $m = 500$ GeV, $g = 1.0$	m_{vis}	1237	2166	27	(140,500)
	m_{eff}	1232	2581	24	(220,660)
	m_{vec}	1107	1848	26	(260,620)
	m_{o1}	850	1129	25	(220,540)
	m_{1o}	1544	5585	21	(20,500)
	$m_{1\top}$	1239	2093	27	(220,620)
	$m_{1\top\text{bound}}$	1282	2236	27	(220,740)
ϕ , $m = 500$ GeV, $g = 1.5$	m_{vis}	2559	2197	55	(140,580)
	m_{eff}	2592	2617	51	(220,820)
	m_{vec}	2363	1919	54	(260,860)
	m_{o1}	1871	1165	55	(220,820)
	m_{1o}	3225	5590	43	(20,580)
	$m_{1\top}$	1994	1271	56	(260,500)
	$m_{1\top\text{bound}}$	2130	1431	56	(260,660)
ϕ , $m = 500$ GeV, $g = 2.0$	m_{vis}	5931	1321	163	(180,460)
	m_{eff}	6292	1733	151	(260,660)
	m_{vec}	6916	1870	160	(260,660)
	m_{o1}	4773	655	186	(260,540)
	m_{1o}	8558	5594	114	(20,700)
	$m_{1\top}$	6273	1324	172	(260,540)
	$m_{1\top\text{bound}}$	6528	1411	174	(260,620)

Table D.1: Discovery potential and corresponding range in m_{WW}^{reco} for the ϕ resonance with mass $m = 500$ GeV and variable coupling $g \in \{0.5, 1.0, 1.5, 2.0\}$ at smeared $\vec{E}_T^{\text{miss}}$ event reconstruction level for an integrated luminosity of 20 fb^{-1} .

Resonance	m_{WW}^{reco} variable	S	B	discovery potential	best window in m_{WW}^{reco} [GeV]
t , $m = 500$ GeV, $g = 0.5$	m_{vis}	8	20	1.73	(520,580)
	m_{eff}	82	1409	2.19	(280,640)
	m_{vec}	19	65	2.36	(580,700)
	m_{o1}	22	269	1.37	(280,340)
	m_{1o}	63	2543	1.26	(100,340)
	$m_{1\top}$	39	512	1.71	(280,340)
	$m_{1\top\text{bound}}$	29	516	1.26	(280,340)
t , $m = 500$ GeV, $g = 1.0$	m_{vis}	181	1746	4.33	(160,640)
	m_{eff}	9	12	2.58	(700,760)
	m_{vec}	160	1556	4.06	(280,700)
	m_{o1}	108	1142	3.18	(220,580)
	m_{1o}	162	3944	2.58	(40,160)
	$m_{1\top}$	137	1048	4.22	(280,520)
	$m_{1\top\text{bound}}$	92	541	3.94	(340,520)
t , $m = 500$ GeV, $g = 1.5$	m_{vis}	585	1772	14	(160,1120)
	m_{eff}	579	1472	15	(280,1300)
	m_{vec}	576	1596	14	(280,940)
	m_{o1}	407	1142	12	(220,580)
	m_{1o}	768	4808	11	(40,580)
	$m_{1\top}$	684	2164	15	(220,1300)
	$m_{1\top\text{bound}}$	646	2092	14	(220,520)
t , $m = 500$ GeV, $g = 2.0$	m_{vis}	1840	3527	31	(100,700)
	m_{eff}	1576	2635	31	(220,1480)
	m_{vec}	1248	1596	31	(280,940)
	m_{o1}	957	1165	28	(220,820)
	m_{1o}	1722	4808	25	(40,580)
	$m_{1\top}$	1498	2156	32	(220,940)
	$m_{1\top\text{bound}}$	1519	2280	32	(220,1240)

Table D.2: Discovery potential and corresponding range in m_{WW}^{reco} for the t resonance with mass $m = 500$ GeV and variable coupling $g \in \{0.5, 1.0, 1.5, 2.0\}$ at smeared $\vec{E}_T^{\text{miss}}$ event reconstruction level for an integrated luminosity of 20 fb^{-1} .

Resonance	m_{WW}^{reco} variable	S	B	discovery potential	best window in m_{WW}^{reco} [GeV]
ϕ , $m = 500$ GeV, $g = 2.5$	m_{vis}	8985	1321	247	(180,460)
	m_{eff}	9521	1770	226	(260,820)
	m_{vec}	8554	1260	241	(300,660)
	m_{o1}	7339	668	284	(260,580)
	m_{1o}	12705	5596	170	(20,980)
	$m_{1\top}$	9606	1379	259	(260,620)
	$m_{1\top\text{bound}}$	9703	1382	261	(260,580)
ϕ , $m = 600$ GeV, $g = 2.5$	m_{vis}	5626	911	186	(220,780)
	m_{eff}	5866	1200	169	(300,820)
	m_{vec}	5665	899	189	(340,900)
	m_{o1}	4775	427	231	(300,820)
	m_{1o}	1789	242	115	(220,700)
	$m_{1\top}$	6255	966	201	(300,820)
	$m_{1\top\text{bound}}$	5307	702	200	(340,820)
ϕ , $m = 750$ GeV, $g = 2.5$	m_{vis}	2406	498	108	(280,780)
	m_{eff}	2415	572	101	(380,1280)
	m_{vec}	2339	412	115	(430,1130)
	m_{o1}	1993	179	149	(380,930)
	m_{1o}	1095	207	76	(230,780)
	$m_{1\top}$	2697	475	124	(380,1130)
	$m_{1\top\text{bound}}$	2755	522	121	(380,1230)
ϕ , $m = 800$ GeV, $g = 2.5$	m_{vis}	1789	504	80	(280,880)
	m_{eff}	1668	478	76	(400,1480)
	m_{vec}	1562	315	88	(460,1120)
	m_{o1}	1392	150	114	(400,1300)
	m_{1o}	959	243	62	(220,820)
	$m_{1\top}$	1863	401	93	(400,1120)
	$m_{1\top\text{bound}}$	1918	443	91	(400,1120)
ϕ , $m = 1000$ GeV, $g = 2.5$	m_{vis}	568	163	45	(400,880)
	m_{eff}	720	279	43	(460,1420)
	m_{vec}	569	114	53	(580,1240)
	m_{o1}	611	85	66	(460,1360)
	m_{1o}	358	102	36	(280,1480)
	$m_{1\top}$	670	149	55	(520,1300)
	$m_{1\top\text{bound}}$	706	189	51	(520,1300)
ϕ , $m = 1200$ GeV, $g = 2.5$	m_{vis}	265	100	27	(460,1060)
	m_{eff}	245	101	24	(580,1360)
	m_{vec}	338	114	32	(580,1240)
	m_{o1}	217	30	40	(580,1360)
	m_{1o}	202	102	20	(280,1480)
	$m_{1\top}$	312	93	32	(580,1300)
	$m_{1\top\text{bound}}$	337	124	30	(580,1360)

Table D.3: Discovery potential and corresponding range in m_{WW}^{reco} for the ϕ resonance with variable mass $m \in \{500 \text{ GeV}, 600 \text{ GeV}, 750 \text{ GeV}, 800 \text{ GeV}, 1000 \text{ GeV}, 1200 \text{ GeV}\}$ and coupling $g = 2.5$ at smeared $\vec{E}_T^{\text{miss}}$ event reconstruction level for an integrated luminosity of 20 fb^{-1} .

Resonance	m_{WW}^{reco} variable	S	B	discovery potential	best window in m_{WW}^{reco} [GeV]
t , $m = 500$ GeV, $g = 2.5$	m_{vis}	3078	3545	52	(100,1480)
	m_{eff}	2759	2635	54	(220,1480)
	m_{vec}	2825	2793	53	(220,820)
	m_{o1}	3964	6277	50	(40,820)
	m_{1o}	3141	4814	45	(40,940)
	$m_{1\top}$	2473	2145	53	(220,820)
	$m_{1\top\text{bound}}$	2548	2283	53	(220,1480)
t , $m = 600$ GeV, $g = 2.5$	m_{vis}	835	880	28	(220,580)
	m_{eff}	777	834	27	(340,1480)
	m_{vec}	909	890	30	(340,820)
	m_{o1}	587	541	25	(280,820)
	m_{1o}	1341	4813	19	(40,760)
	$m_{1\top}$	794	666	31	(340,820)
	$m_{1\top\text{bound}}$	821	717	31	(340,940)
t , $m = 750$ GeV, $g = 2.5$	m_{vis}	327	296	19	(330,680)
	m_{eff}	279	257	17	(470,1450)
	m_{vec}	343	271	20	(470,820)
	m_{o1}	224	143	19	(400,820)
	m_{1o}	654	4434	9.8	(50,1030)
	$m_{1\top}$	422	384	22	(400,820)
	$m_{1\top\text{bound}}$	441	419	22	(400,820)
t , $m = 800$ GeV, $g = 2.5$	m_{vis}	222	189	16	(380,780)
	m_{eff}	251	279	15	(460,1340)
	m_{vec}	221	143	18	(540,860)
	m_{o1}	155	77	18	(460,780)
	m_{1o}	140	242	9.0	(220,780)
	$m_{1\top}$	185	102	18	(540,780)
	$m_{1\top\text{bound}}$	216	138	18	(540,860)
t , $m = 1000$ GeV, $g = 2.5$	m_{vis}	103	95	11	(460,860)
	m_{eff}	78	74	9.1	(620,1500)
	m_{vec}	153	160	12	(540,1260)
	m_{o1}	75	43	12	(540,1340)
	m_{1o}	45	79	5.0	(300,1500)
	$m_{1\top}$	140	122	13	(540,1020)
	$m_{1\top\text{bound}}$	118	93	12	(620,1340)
t , $m = 1200$ GeV, $g = 2.5$	m_{vis}	4	1	5.2	(940,1020)
	m_{eff}	38	39	6.1	(700,1500)
	m_{vec}	50	50	7.1	(700,1260)
	m_{o1}	6	1	6.8	(1020,1260)
	m_{1o}	14	13	3.91	(460,780)
	$m_{1\top}$	47	40	7.5	(700,1260)
	$m_{1\top\text{bound}}$	59	56	7.9	(700,1260)

Table D.4: Discovery potential and corresponding range in m_{WW}^{reco} for the t resonance with variable mass $m \in \{500 \text{ GeV}, 600 \text{ GeV}, 750 \text{ GeV}, 800 \text{ GeV}, 1000 \text{ GeV}, 1200 \text{ GeV}\}$ and coupling $g = 2.5$ at smeared $\vec{E}_T^{\text{miss}}$ event reconstruction level for an integrated luminosity of 20 fb^{-1} .

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