Two-Higgs doublet model and the LHC

Chuene Mosomane

National Institute for Theoretical Physics: School of Physics and Mandelstam Institute for theorecical Physics, University of Witwatersrand, Wits 2050, South Africa

E-mail: cjmosomane@gmail.com

Abstract. The discovery of a 125 GeV Higgs boson at the Large Hadron Collider (LHC) in 2012, which has properties as expected in the Standard Model (SM), was a major milestone for particle physics. More recently the ATLAS and CMS experiments have reported several excesses in the search for di-Higgs boson resonances. In these proceedings, we shall review the study of the general two-Higgs doublet model (2HDM) together with the Inert doublet model (IDM). The SM Higgs boson can also be realised within the Inert Doublet Model (IDM); a version of the two Higgs doublet model (2HDM), with an unbroken Z_2 symmetry under which one of the $SU(2)$ is the SM Higgs doublet (with one being the SM Higgs boson), and the second $SU(2)$ doublet transforms non-trivially with no vacuum expectation value, and does not interact with fermions. The 2HDM is a model which goes beyond the SM; it has a richer particle spectrum. The most general potential of a 2HDM is Lorentz invariant and renormalizable, containing 14 free parameters, with the most general Yukawa Lagrangian.

1. Introduction

The discovery of the Higgs boson is a milestone in particle physics and an extraordinary success of the LHC machine and the ATLAS and CMS collaborations [1]. The existence of the Higgs boson ensures unitary in the scattering of longitudinal polarised W bosons, which allows the introduction of the gauge bosons ($W^{\pm}, Z^0$) and the quarks masses without breaking the gauge symmetry of the theory through the Higgs mechanism [2]. The underlying idea of this mechanism is the introduction of a complex scalar field into the theory, whose ground state acquires a non-zero vacuum expectation value (vev). The potential of the Higgs field does not share the same symmetry as the full Lagrangian of the theory, and hence the symmetry is spontaneously broken by the vev of the Higgs field. Due to the Goldstone theorem [3], which implies that a spontaneously broken local symmetry leads to massless Goldstone bosons whose degrees of freedom are then eaten up by the gauge bosons, only one degree of freedom of the Higgs field remains as a physical particle in the theory. This is the Higgs boson whose mass is a parameter in the theory[4]. Despite its great success in explaining all available data the SM has several inadequacies for the fermions and mass spectrum. It does not contain any dark matter particles [5], the SM also does not explain the strong CP problem [6]. These deficiencies led particle physicists to explore theory Beyond the SM (BSM). The two Higgs doublet model (2HDM) predicts more physical scalars such as a charged scalar, two neutral scalars and a pseudoscalar. We shall also discuss the Inert doublet model (IDM), which is a 2HDM with an unbroken Z_2 symmetry under which one of the doublets transform non-trivially, and all other SM fields are invariant.

2. 2HDM

The SM (with only one scalar field) has been a triumphant achievement in particle physics in terms of explaining the laws of nature, it has been regarded as having inadequacies and incomplete because there are phenomena that it doesn't explain (such as the darkmatter candidate and the baryonic-asymmetry of the univere). This gives us the impetus to investigate more complete theories that better address these questions, these are what we consider Beyond the SM (BSM). The 2HDM is a BSM theoty which answers some of these questions, such as CP-violation etc. The motivation behind the 2HDM is that when breaking the electroweak symmetry does nature only allow for one scalar field, where by adding another scalar doublet to the SM we observe some interesting properties, such as the production of additional Higgs particles, sources of CP-violation and Flavour Changing Neutral Currents (FCNCs). We shall write the 2HDM Lagrangian which requires it to be renormalisable and Lorentz invariant.

$$\mathcal{L}_{2HDM} = (D_\mu \Phi_1)^\dagger (D^\mu \Phi_1) + (D_\mu \Phi_2)^\dagger (D^\mu \Phi_2) - V(\Phi_1 \Phi_2), \tag{1}$$

where the covariant derivative in the standard notation is given by:

$$D_\mu = \partial_\mu + ig\frac{\tau^i}{2} \cdot W^i_\mu + i\frac{g'}{2} Y B_\mu, \tag{2}$$

and the τ^i are the Pauli matrices and the Y is the hypercharge. The potential is:

$$\begin{aligned} V &= m_{11}^2 \Phi_1^\dagger \Phi_1 + m_{22}^2 \Phi_2^\dagger \Phi_2 - (m_{12}^2 \Phi_1^\dagger \Phi_2 + h.c) \\ &+ \frac{\lambda_1}{2}(\Phi_1^\dagger \Phi_1)^2 + \frac{\lambda_2}{2}(\Phi_2^\dagger \Phi_2)^2 + \lambda_3 (\Phi_1^\dagger \Phi_1)(\Phi_2^\dagger \Phi_2) + \lambda_4 (\Phi_1^\dagger \Phi_2)(\Phi_2^\dagger \Phi_1) \\ &+ \left[\frac{\lambda_5}{2}(\Phi_1^\dagger \Phi_2)^2 + \lambda_6 (\Phi_1^\dagger \Phi_1)(\Phi_1^\dagger \Phi_2) + \lambda_7 (\Phi_2^\dagger \Phi_2)(\Phi_1^\dagger \Phi_2)\right]. \end{aligned} \tag{3}$$

The 2HDM potential exhibits two CP-conserving neutral minima $\langle \Phi_1 \rangle$ and $\langle \Phi_2 \rangle$ of the form:

$$\langle \Phi_1 \rangle = \begin{pmatrix} 0 \\ \frac{v_1}{\sqrt{2}} \end{pmatrix}, \qquad \langle \Phi_2 \rangle = \begin{pmatrix} 0 \\ \frac{v_2}{\sqrt{2}} \end{pmatrix},$$

with v_1 and v_2 being the vev's of the scalar doublet fields. In order for the doublets Φ_a to have their minima at $\langle \Phi_a \rangle$, the minimisation conditions apply, the end result gives us the following expressions

$$m_{11}^2 = m_{12}^2 \frac{v_2}{v_1} - \frac{\lambda_1}{2} v_1^2 - (\lambda_3 + \lambda_4 + \lambda_5)\frac{v_2^2}{2}, \tag{4}$$

$$m_{22}^2 = m_{12}^2 \frac{v_1}{v_2} - \frac{\lambda_1}{2} v_1^2 - (\lambda_3 + \lambda_4 + \lambda_5)\frac{v_1^2}{2}. \tag{5}$$

Introducing eight fields $\phi_a^\pm$, ρ_a, and η_a ($a = 1, 2$) the doublets may be expanded around the minima, taking the form:

$$\Phi_a = \begin{pmatrix} \Phi_a^\pm \\ \frac{1}{\sqrt{2}}(v_a + \rho_a + i\eta_a) \end{pmatrix}, \qquad a = 1, 2.$$

Inserting the scalar fields into the 2HDM potential generates terms that are linear in the two fields ρ_a. The expansion of the scalar fields into the 2HDM potential generates terms that are bilinear in the fields $\phi_a^\pm, \rho_a$ and η_a. Since these bilinear terms contribute to the propagators of the eight fields, they give rise to the mass terms. All bilinear terms in the 2HDM potential can be transformed into the explicit form of non-diagonal mass-squared matrices given by,

$$\mathcal{M}_\rho^2 =- \begin{pmatrix}\rho_1, & \rho_2\end{pmatrix} \begin{pmatrix} m_{12}^2\frac{v_2}{v_1}+\lambda_1 v_1^2 & -m_{12}^2+\lambda_{345}v_1v_2 \\ -m_{12}^2+\lambda_{345}v_1v_2 & m_{12}^2\frac{v_2}{v_1}+\lambda_2 v_2^2 \end{pmatrix} \begin{pmatrix}\rho_1 \\ \rho_2\end{pmatrix}, \tag{6}$$

$$\mathcal{M}_\eta^2 =\frac{m_A^2}{v_1^2+v_2^2} \begin{pmatrix}\eta_1, & \eta_2\end{pmatrix} \begin{pmatrix} v_2^2 & -v_1v_2 \\ -v_1v_2 & v_1^2 \end{pmatrix} \begin{pmatrix}\eta_1 \\ \eta_2\end{pmatrix}, \tag{7}$$

$$\mathcal{M}_{\phi^\pm}^2 =[m_{12}^2-(\lambda_4+\lambda_5)v_1v_2] \begin{pmatrix}\phi_1^-, & \phi_2^-\end{pmatrix} \begin{pmatrix} \frac{v_2}{v_1} & -1 \\ -1 & \frac{v_1}{v_2} \end{pmatrix} \begin{pmatrix}\phi_1^+ \\ \phi_2^+\end{pmatrix}. \tag{8}$$

To have physical propagating particles in the 2HDM, we shall consider the eigenstates with specific masses. This can be achieved by diagonalising the mass-squared matrices in Eq (6-8). The fields $\phi_i^\pm, \rho_i$ and η_i in the gauge basis are transformed into physical fields.The massless particles G^0 and $G^\pm$ are the three Goldstone bosons of the 2HDM and are "eaten"up to give the $W^\pm$, Z^0 gauge bosons. The remaining five physical particles $H^0(heavy)$ and $h^0(lighter)$ are the CP-even, A^0 is a CP-odd and the $H^\pm$ are the charged Higgs:

$$h^0 =\rho_1 \sin\alpha - \rho_2\cos\alpha, \tag{9}$$

$$H^0 =-\rho_1\cos\alpha - \rho_2\sin\alpha, \tag{10}$$

$$H^\pm =\phi_1^\pm \sin\beta + \phi_2^\pm\cos\beta, \tag{11}$$

$$A^0 =\eta_1\sin\beta - \eta_2\cos\beta, \tag{12}$$

where the parameter α is a mixing angle between the CP-even scalars (h^0, H^0) and the parameter is the β is a rotational angle that diagonalises the mass-squared matrices of the charged Higgs and the pseudoscalars abd is defined as; $\tan\beta = v_2/v_1$. The mass matrices Eq. (5-7) are diagonal in this basis; therefore these corresponding fields are referred to as the mass basis of the 2HDM potential. The tree-level masses of the particles in the mass basis are given by

$$m_{H^0}^2 =\frac{1}{2}\left[\mathcal{M}_{11}^2+\mathcal{M}_{12}^2+\sqrt{(\mathcal{M}_{11}^2-\mathcal{M}_{22}^2)^2+4(\mathcal{M}_{12}^2)^2}\right], \tag{13}$$

$$m_{h^0}^2 =\frac{1}{2}\left[\mathcal{M}_{11}^2+\mathcal{M}_{12}^2-\sqrt{(\mathcal{M}_{11}^2-\mathcal{M}_{22}^2)^2+4(\mathcal{M}_{12}^2)^2}\right], \tag{14}$$

$$m_{H^0}^2 =v^2\left(\frac{m_{12}^2}{v_1v2}-\lambda_5\right), \qquad m_{G^0}^2=0, \tag{15}$$

$$m_{H^\pm}^2 =v^2\left(\frac{m_{12}^2}{v_1v_2}-\frac{\lambda_4+\lambda_5}{2}\right), \qquad m_{G^\pm}^2=0, \tag{16}$$

where $\mathcal{M}_{11}, \mathcal{M}_{12}, \mathcal{M}_{22}$ are the mass-squared matrices and the λ_4, λ_5 are the quartic couplings.

3. Yukawa Lagrangian

In the SM, diagonalising the mass matrix automatically diagonalises the Yukawa interactions, therefore, there are no tree-level FCNC. A general 2HDM introduces the possibility of having FCNCs at tree-level. However, in general Yukawa interactions (Y_1 and Y_2) will not be simultaneously diagonalised, and thus the Yukawa couplings will not be flavour diagonal. The

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interaction of the fermions and scalar bosons of the 2HDM is determined by the form of the Yukawa Lagrangian. The transformation from a general to a flavour-conserving 2HDM can be achieved naturally by imposing a discrete or continous symmetries on the Higgs doublets Φ_1 and Φ_2 [7]. Imposing these discrete symmetries will suppress FCNCs at tree-level with the same quantum numbers, i.e those that potentially mix, and couple only to the same Higgs doublet. The 2HDM will have four different model types and can be summarised in the table 1.

Table 1. Natural flavour conservation models. The superscript i is a generation index.

Models	Type I	Type II	(Leptonic-specific)	(Flipped)
u_R^i	Φ_2	Φ_2	Φ_2	Φ_2
d_R^i	Φ_2	Φ_1	Φ_2	Φ_1
e_R^i	Φ_2	Φ_2	Φ_1	Φ_2

We can write the Yukawa Lagrangian that conserves the FCNCs as follows:

$$\begin{aligned}\mathcal{L}_Y =&\bar{Q}_{L.i}(Y_{u,1}^{ij}\tilde{\Phi}_1 + Y_{u,2}^{ij}\Phi_2)u_{R,j} + \bar{Q}_{L.i}(Y_{d,1}^{ij}\Phi_1 + Y_{u,2}^{ij}\tilde{\Phi}_2)d_{R,j} \\ &+ \bar{L}_{L.i}(Y_{l,1}^{ij}\Phi_1 + Y_{l,2}^{ij}\tilde{\Phi}_2)l_{R,j} + h.c,\end{aligned} \tag{17}$$

with $\tilde{\Phi}_i = -i[\Phi_i^\dagger \tau_2]^T, i = 1,2$ and $Y_{u,1}^{ij}$ being the 3×3 Yukawa matrices. The (u_R, d_R) represent the right handed up-quark and down-quark, the l_R represent the right handed lepton, and the $(\bar{Q}_{L.i}, \bar{L}_{L.i})$ are the left handed quarks and left handed leptons. The coupling constants of the Yukawa interaction can be summarised as follows[8]:

4. Inert doublet model

The IDM is introduced to allow for the possibility of several mirror families of fermions. In the IDM the Higgs doublet Φ_2 does not couple to matter and the vev is zero, leaving the Z_2 symmetry unbroken. The scalar spectrum consists of the SM-like Higgs obtained from Φ_1 and one charged and two neutral states from Φ_2 (inert scalars). The scalar potential is the same as in Eq. (2) with $m_{12}^2 = 0$. The asymmetry phase $\langle\phi_1^0\rangle = v/\sqrt{2}$ and $\langle\phi_2^0\rangle = 0$ corresponds to a sizeable region of parameter space, and the scalar masses are as follows:

$$m_h^2 = \lambda_1 v^2, \qquad m_S^2 = m_{22} + \frac{(\lambda_3 + \lambda_4 + \lambda_5)v^2}{2}, \tag{18}$$

$$m_+^2 = m_{22}^2\frac{\lambda_3 v^2}{2}, \qquad m_A^2 = m_{22} + \frac{(\lambda_3 + \lambda_4 + \lambda_5)v^2}{2}. \tag{19}$$

The $m_+^2 = m_-^2$ is the mass-squared matrix for the charged Higgs boson and it is the same, the Φ_2 can be produced at colliders through their couplings to the electroweak gauge bosons, subjected to the Z_2 symmetry. In addition they also participate in the cubic and quartic Higgs couplings:

$$\begin{aligned}V_{int} =& \frac{\lambda_2}{2}\left(H^+H^- + \frac{S^2 + A^2}{2}\right)^2 + \lambda_3\left(vh + \frac{h^2}{2}\right)\left(H^+H^- + \frac{S^2 + A^2}{2}\right) \\ &+ \frac{\lambda_4 + \lambda_5}{2}\left(vh + \frac{h^2}{2}\right)S^2 + \frac{\lambda_4 - \lambda_5}{2}\left(vh + \frac{h^2}{2}\right)A^2.\end{aligned} \tag{20}$$

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Assuming the mass hierarcy $m_+^2 > m_A^2 > m_S^2$, the dominant decay of $A \to Sf\bar{f}$, where A is a CP-odd scalar and f is fermion and S (singlet scalar) appears as missing energy. The decay pattern and the ratios of the neutral scalars can be examined at the LHC. We would observe S with mass ≈ 50 GeV. In this case S could be observed as a resonance through $pp \to S \to VV$ modes, with $V = ZZ, W^+W^-$; it also alters the coupling strength of the known interaction in the theory[8].

5. Phenomenology of the 2HDM at the LHC

We shall also discuss the phenomenology of the 2HDM particle spectrum. We are interested in studying the production and the decay of CP-even scalars (and the couplings), the Heavy Higgs and lighter Higgs as summarised in Table 2. In this proceedings we shall concern ourselves with the CP-even scalars. The phenomenology of h is similar to that of the SM Higgs boson coupling. The vector boson coupling constants for the lighter Higgs boson, hZZ and hW^+W^-, are given by the SM Higgs boson times $\sin(\beta - \alpha)$, similarly HZZ and HW^+W^- are proportional to $\cos(\beta - \alpha$; these vertices are the most important for our phenomenology. The scalars h and H thus share the Higgs field vev and strength of the coupling of the W^+W^- and ZZ to scalar fields. Furthermore the decay of the Higgs boson in the 2HDM depends on the model for the Yukawa interaction and also when $\sin(\beta - \alpha) = 1$ the decay pattern of h is almost the same as that of the SM. Therefore the production of H can vary over a large range. In Figure 1, we have conducted a general search for an additional heavy Higgs boson, $H \to \gamma\gamma$, in a mass range betwwen $2m_h < m_H < 2m_t$[9]. The analyses are performed in a combination of 50000 events category; we showed the energy scales in GeV against the number of events which is the normalized differential cross-sections.

Table 2. Couplings of of h^0 and H^0 to gauge bosons pairs:

$\cos(\beta - \alpha)$	$\sin(\beta - \alpha)$
HW^+W^-	hW^+W^-
HZZ	hZZ
ZAh	ZAH
$W^\pm H^\mp h$	$W^\pm H^\mp H$
$ZW^\pm H^\mp h$	$ZW^\pm H^\mp h$
$\gamma W^\pm H^\mp h$	$\gamma W^\pm H^\mp h$

6. Conclusion

In this proceedings we discussed the particle spectrum that arises because of electroweak symmetry breaking in the 2HDM. There are many possibilities for the decay branching ratios of these particles. We briefly reviewed the 2HDM, which is an extension of the SM with an additional scalar doublet. In particular we have studied the IDM, which is obtained by setting one of vevs to zero. The ongoing searches at the LHC rely on specific production and decay mechanisms that occupy only a part of the complete parameter space[**?**]. The study of the IDM is done by imposing several constraints on the exploration of the Higgs sector, where the main constraint comes from the discovery of the resonace 125 GeV by ATLAS and CMS. In the context of 2HDM, this resonance might be interpreted as the h an H. In the $H \to \gamma\gamma$, the final state particles can be very precisely measured and the reconstructed mass m_H is excellent

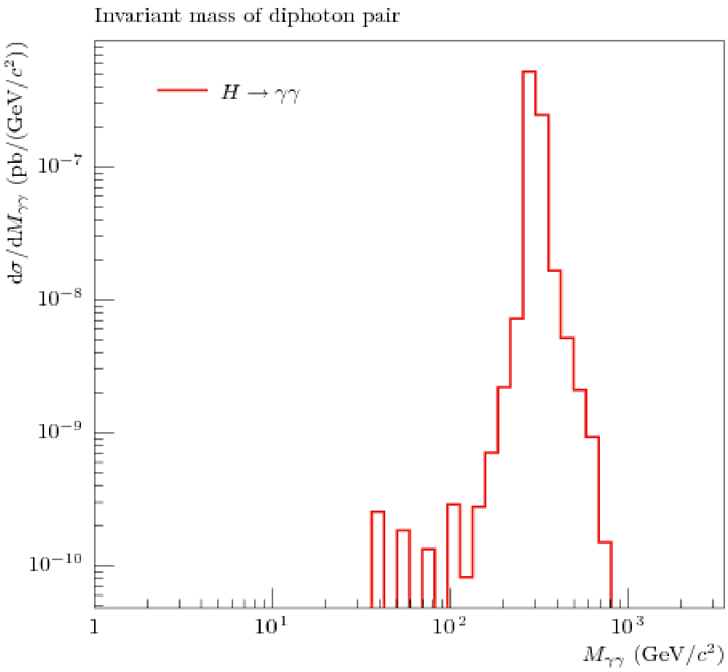

Figure 1. Left:Invariant mass of the kinematic fit in the process $H \to \gamma\gamma$ in 2HDM with $m_H \approx 300$GeV

as shown in Fig.1. Currently, we can make a demonstration of how the phenomenology of the 2HDM particle spectrum could be observed at the LHC using some Monte-Carlo simulators.

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