High-scale validity of a two-Higgs doublet scenario: a study including LHC data

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Abstract

We consider the conditions for the validity of a two-Higgs doublet model at high energy scales, together with all other low-and high-energy constraints. The constraints on the parameter space at low energy, including the measured value of the Higgs mass and the signal strengths in channels are juxtaposed with the conditions of vacuum stability, perturbativity and unitarity at various scales. We find that a scenario with an exact $Z_2$ symmetry in the potential cannot be valid beyond about 10 TeV without the intervention of additional physics. On the other hand, when the $Z_2$ symmetry is broken, the theory can be valid even up to the Planck scale without any new physics coming in. The interesting feature we point out is that such high-scale validity is irrespective of the top quark mass uncertainty, in contrast with the standard model with a single Higgs doublet. It is also shown that the presence of a CP-violating phase which is allowed when the $Z_2$ symmetry is relaxed. The allowed regions in the parameter space are presented for each case. Though our results are illustrated in the context of a Type-II model, their qualitative features are preserved for a general two-Higgs doublet scenario.
1 Introduction

The Higgs sector of the standard electroweak model (SM) continues to appear enigmatic from several angles. The existence of such a sector, comprising at least one scalar doublet, and driving the spontaneous symmetry breakdown $SU(2)_L \times U(1)_Y \rightarrow U(1)_{EM}$ is almost impossible to deny now. It is also widely agreed that the Large Hadron Collider (LHC) has found [1, 2] a neutral boson with mass around 125 GeV, which is almost certainly of spin zero [3] and dominantly a CP-even field [4–7]. However, despite the properties of the boson being consistent with that of the SM Higgs, rather persistent enquiries are on, to find out whether the electroweak symmetry breaking sector also contains some signature of physics beyond the standard model. The LHC data till date leaves room for such new physics.

Two sets of standpoints are noticed in such enquiries. First of all, with spin-1/2 fermions showing family replication, it is not obvious why the part of the matter sector containing spin-zero particles should also not have similar repetition. With this in view, multi-doublet scenarios are under regular scrutiny, the most widely investigated being models with two Higgs doublets. An extended electroweak symmetry breaking sector entails a rich phenomenology, including additional sources of CP violation [8]. Of course, scalars belonging to higher representations of $SU(2)$ have also attracted attention, especially triplets which can play a role in the so-called Type-II mechanism of neutrino mass generation [9]. Secondly, even with just one doublet (leading to a single physical scalar), the Higgs mass is not stable under quadratically divergent radiative corrections, and it is somewhat artificial (or ‘fine-tuned’) to have a 125-GeV Higgs if the cut-off for the SM is much higher than a TeV or so. Furthermore, it is also not clear that the SM scalar potential retains a finite and stable minimum at high scales. But for the yet uncertain measurement of the top quark mass, which is crucial in governing the evolution of the Higgs self-coupling via Yukawa interactions, we may be doomed to live in an unstable or metastable vacuum new physics intervenes at a scale no greater than $10^{8–10}$ GeV [10–13]. Therefore, the ultraviolet incompleteness of the current scheme of electroweak symmetry breaking looms up as a distinct possibility, even if one disregards the somewhat philosophical issue of naturalness.

In this paper, we follow these two standpoints in tandem. We take up a two-Higgs doublet scenario as the minimal extension of the standard electroweak theory, assessing its viability as well sufficiency modulo all available constraints. The intuitive motivations for the study is that the proportionality constant between the top quark mass and its coupling to the 125 GeV scalar is different from its SM value when more doublets are around. Consequently, the dependence of the vacuum stability limit on the top quark mass is expected to be different.
However, one can make precise and quantitative statements on the matter only when one takes cognizance of the exact scenario, and includes the complete set of renormalisation group equations appropriate for it. This is precisely what we aim to do here, using a two-Higgs doublet scenario at various levels of generality.

The desired suppression of flavour-changing Yukawa interactions is best implemented by imposing a discrete symmetry on such models, thus preventing both the doublets from coupling with $T_3 = +1/2$ and $-1/2$ fermions simultaneously. It is possible to go beyond such imposition and examine two Higgs doublets in a ‘basis-independent’ formulation [14–16]. However, we feel that our central issue, namely, the evolution of the Higgs self-interaction(s), is amenable to a more transparent study if one adheres to a specific Yukawa scheme. With this in view, we adopt the so-called Type-II scenario for our study [17], though our broad conclusions do not depend on this choice.

We begin by examining the situation when the discrete symmetry is exact, and derive the constraints on the low-energy values of the parameters of this scenario. The lighter neutral scalar mass being around 125 GeV is of course the prime requirement here, and constraints from rare processes such as $b \to s\gamma$ are also included. In addition, the constraints from perturbativity of all scalar quartic couplings are considered, together with those from vacuum stability. The parameter space thus validated is further examined in the light of the perturbativity and vacuum stability conditions at high scales. Thus we identify the parameter regions that keep a two-Higgs doublet scenario valid upto different levels of high scales— an exercise that reveals rather severe limits. The same investigation is carried out for cases where the discrete symmetry is broken by soft (dimension-2) and hard (dimension-4) terms in turn, with the Yukawa coupling assignment remaining (for simplicity) the same as in the case with unbroken symmetry. The effect of a CP-violating phase is also demonstrated. Finally, the regions found to be allowed from all the above considerations, at both low-and high-scales, are pitted against the existing data from the LHC in different channels. Thus we identify parameter regions that are consistent with the measured signal strengths in different channels. This entire study is aimed at indicating how far a two-Higgs doublet model can remain valid, not only at the LHC energy but also upto various high scales without further intervention of new physics.

Although a number of recent studies have addressed some similar questions, the present study has gone beyond them on the following points:

- Our study reveals that the precise value of the top quark mass is rather unimportant in deciding the high-scale validity of the theory. Regions in the parameter space are
identified, for which the theory has no cut-off till the Planck scale, even though the
top quark mass can be at the upper edge of the allowed band.

- We find that it is rather difficult to retain the validity of a two-Higgs doublet scenario
  well above a TeV with the discrete symmetry intact. Also, large values of $\tan\beta$, the
  ratio of the vacuum expectation values (vev) of the two doublets, are mostly disfavoured
  in this case.

- With the discrete symmetry broken, there is a correlation between allowed $\tan\beta$ and
  the extent of symmetry breaking, when it comes to validity upto the Planck scale.

- We examine the constraints on the model including a CP-violating phase [18–21]. In
  fact, since the existence of a phase is a natural consequence of relaxing the discrete
  symmetry, the high-scale validity of a two Higgs doublet model may be argued to be
  contingent on the possibility of CP-violation in the scalar potential.

- We have performed a detailed examination of the validity of the scenario at both
  low and high scales, including dimension-4 discrete symmetry breaking terms in our
  analysis. The LHC constraints are also imposed in this situation.

We remind the reader of the broad features of a two-Higgs doublet scenario in section 2.
In section 3, we list and explain all the constraints that the scenario is subjected to, at both
the low and high scales. Sections 4, 5 and 6 contain, in turn, the results of our analysis,
with the discrete symmetry intact, softly broken and broken by hard terms, respectively. We
summarise and conclude in section 7.

2 The two-Higgs-doublet scenario and the scalar potential: basic features

In the present work, we consider the most general renormalizable scalar potential for two
doublets $\Phi_1$ and $\Phi_2$, each having hypercharge (+1).

\[
V(\Phi_1, \Phi_2) = m_{11}^2 \Phi_1^\dagger \Phi_1 + m_{22}^2 \Phi_2^\dagger \Phi_2 - m_{12}^2 (\Phi_1^\dagger \Phi_2 + \Phi_2^\dagger \Phi_1) + \lambda_1 \frac{1}{2} (\Phi_1^\dagger \Phi_1)^2 + \lambda_2 \frac{1}{2} (\Phi_2^\dagger \Phi_2)^2 + \lambda_3 \Phi_1^\dagger \Phi_1 \Phi_2^\dagger \Phi_2 + \lambda_4 \Phi_1^\dagger \Phi_2 \Phi_2^\dagger \Phi_1 + \lambda_5 \frac{1}{2} \left[ (\Phi_1^\dagger \Phi_1)^2 + (\Phi_2^\dagger \Phi_2)^2 \right] + \lambda_6 \Phi_1^\dagger \Phi_2 \left( \Phi_1^\dagger \Phi_2 + \Phi_2^\dagger \Phi_1 \right) + \lambda_7 \Phi_2^\dagger \Phi_2 \left( \Phi_1^\dagger \Phi_2 + \Phi_2^\dagger \Phi_1 \right)
\] (2.1)
The parameters $m_{12}$, $\lambda_5$, $\lambda_6$ and $\lambda_7$ could be complex in general, although the phase in one of them can be removed by redefinition of of the relative phase between $\Phi_1$ and $\Phi_2$. Thus this scenario in general has the possibility of CP-violation in the scalar sector.

In a general two-Higgs-doublet model (2HDM), a particular fermion can couple to both $\Phi_1$ and $\Phi_2$. However this would lead to the flavor changing neutral currents (FCNC) at the tree level [22–25]. One way to avoid such FCNC is to impose a $Z_2$ symmetry, such as one that demands invariance under $\Phi_1 \to -\Phi_1$ and $\Phi_2 \to \Phi_2$. This type of symmetry puts restrictions on the scalar potential. The $Z_2$ symmetry is exact as long as $m_{12}$, $\lambda_6$ and $\lambda_7$ vanish, when the scalar sector also becomes CP-conserving. The symmetry is said to be broken softly if it is violated in the quadratic terms only, i.e., in the limit where $\lambda_6$ and $\lambda_7$ vanish but $m_{12}$ does not. Finally, a hard breaking of the $Z_2$ symmetry is realized when it is broken in the quartic terms as well. Thus in this case, $m_{12}$, $\lambda_6$ and $\lambda_7$ all are non-vanishing in general.

As mentioned in the introduction, we focus on a specific scheme of coupling fermions to the doublets. This scheme is referred to in the literature as the Type-II 2HDM, where the down type quarks and the charged leptons couple to $\Phi_1$ and the up type quarks, to $\Phi_2$ [26]. This can be ensured through the discrete symmetry $\Phi_1 \to -\Phi_1$ and $\psi_R^i \to -\psi_R^i$, where $\psi$ is charged leptons or down type quarks and $i$ represents the generation index. It has been already mentioned that this choice is purely for illustrative purpose; our general results are largely independent of it. Although we start by analysing the high-scale validity of the model with $m_{12} = \lambda_6 = \lambda_7 = 0$, we subsequently include the effects of both soft and hard breaking of $Z_2$ in turn, which bring back these parameters. The two simplifications that we still make are as follows: (a) the phases of $\lambda_6$ and $\lambda_7$ are neglected though that of $m_{12}$ is considered, and (b) the Yukawa coupling assignments of $\Phi_1$ and $\Phi_2$ are left unchanged.

Minimization of the scalar potential in Eq. 2.1 yields

$$
\langle \Phi_1 \rangle = \left( \begin{array}{c}
0 \\
\frac{v_1}{\sqrt{2}}
\end{array} \right), \quad \langle \Phi_2 \rangle = \left( \begin{array}{c}
0 \\
\frac{v_2}{\sqrt{2}}
\end{array} \right),
$$

(2.2)

where the vacuum expectation values (vev) are often expressed in terms of the $m_Z$ and the ratio

$$
\tan \beta = \frac{v_2}{v_1}.
$$

(2.3)

We parametrize the doublets in the following fashion,

$$
\Phi_i = \frac{1}{\sqrt{2}} \left( \begin{array}{c}
\sqrt{2}w_i^+ \\
vi + hi + iz_i
\end{array} \right) \text{ for } i = 1, 2.
$$

(2.4)
Since the basis used in $V(\Phi_1, \Phi_2)$ allows mixing between the two doublets, one diagonalises the charged and neutral scalar mass matrices to obtain the physical states. There are altogether eight mass eigenstates, three of which become the longitudinal components of the $W^\pm$ and $Z$ gauge bosons. Of the remaining five, there is a mutually conjugate pair of charged scalars ($H^\pm$), two neutral scalars ($H, h$) and a neutral pseudoscalar ($A$), when there is no CP-violation. Otherwise, a further mixing occurs between ($H, h$) and $A$. The mass eigenstates $H$ and $h$ are decided by a mixing angle $\alpha$.

In the absence of CP-violation, the squared masses of these physical scalars and the mixing angle $\alpha$ can be expressed as [27],

\begin{align}
  m_A^2 &= m_{12}^2 \frac{s_\beta c_\beta}{2} - \frac{1}{2} v^2 \left(2 \lambda_5 + \frac{\lambda_6}{t_\beta} + \lambda_7 t_\beta\right), \quad (2.5a) \\
  m_{H^\pm}^2 &= m_A^2 + \frac{1}{2} v^2 (\lambda_5 - \lambda_4), \quad (2.5b) \\
  m_h^2 &= \frac{1}{2} \left[(A + B) - \sqrt{(A - B)^2 + 4C^2}\right], \quad (2.5c) \\
  m_H^2 &= \frac{1}{2} \left[(A + B) + \sqrt{(A - B)^2 + 4C^2}\right], \quad (2.5d) \\
  \tan 2\alpha &= \frac{2C}{A - B}, \quad (2.5e)
\end{align}

where we have defined,

\begin{align}
  A &= m_A^2 s_\beta^2 + v^2 (\lambda_1 c_\beta^2 + \lambda_5 s_\beta^2 + 2\lambda_6 s_\beta c_\beta), \quad (2.6a) \\
  B &= m_A^2 c_\beta^2 + v^2 (\lambda_2 s_\beta^2 + \lambda_5 c_\beta^2 + 2\lambda_7 s_\beta c_\beta), \quad (2.6b) \\
  C &= -m_A^2 s_\beta c_\beta + v^2 [(\lambda_3 + \lambda_4) s_\beta c_\beta + \lambda_6 c_\beta^2 + \lambda_7 s_\beta^2]. \quad (2.6c)
\end{align}

Furthermore, the interactions of the various charged and neutral scalars to the up-and down-type fermions are functions of $\alpha$ and $\beta$. Their detailed forms in different 2HDM scenarios, including the Type-II model adopted here for illustration, can be found in the literature [28].

### 3 Theoretical and experimental constraints

Next, we subject the Type-II 2HDM using various theoretical and experimental constraints (though the most binding ones are often irrespective of the specific type of 2HDM). It should be remembered at the outset that the most general $Z_2$ violating 2HDM has seven quartic couplings, namely, $\lambda_i$ ($i = 1, \ldots, 7$), in addition to $\tan \beta$ and $m_{12}$, totalling to nine free parameters. Though such a nine-dimensional parameter is prima facie large enough
to accommodate any phenomenology, the set of constraints under consideration below can ultimately become quite restrictive.

We discuss the theoretical constraints in subsection 3.1, and take up the experimental/phenomenological ones in the subsequent subsections. It should be noted that the parameter space is being constrained in two distinct ways. Subsections 3.2 - 3.4 list essentially low-energy constraints which apply at the energy scale of the subprocesses leading to Higgs production. The various masses and couplings get restricted by the requirement of satisfying them. However, while such a strategy is valid for the discussion of subsection 3.1 as well, we additionally require the conditions laid down there to hold at various high scales, too. This not only restricts the low-energy parameters more severely, but also answers the main question asked in this paper, namely, to what extent the 2HDM can be deemed ‘ultraviolet complete’.

3.1 Perturbativity, unitarity and vacuum stability

For the 2HDM to behave as a perturbative quantum field theory at any given scale, one must impose the conditions $|\lambda_i| \leq 4\pi$ ($i = 1, \ldots, 7$) and $|y_i| \leq \sqrt{4\pi}$ ($i = t, b, \tau$) at that scale\(^1\). On applying such conditions, one implies upper bounds on the values of the couplings at low as well as high scales.

Next, we impose the more stringent condition of unitarity on the tree-level scattering amplitudes involving the scalar degrees of freedom. In a model with an extended scalar sector, the scattering amplitudes are taken between various two-particle states constituted out of the fields $w_{i}^{\pm}$, $h_i$ and $z_i$ corresponding to the parametrization of Eq. 2.4. Maintaining this, there will be neutral two-particle states (e.g., $w_{i}^{+}w_{j}^{-}$, $h_ih_j$, $z_iz_j$, $h_iz_j$) as well as singly charged two-particle states (e.g., $w_{i}^{\pm}h_j$, $w_{i}^{\pm}z_j$). The various two particle initial and final states give rise to a $2 \rightarrow 2$ scattering matrix whose elements are the lowest order partial wave expansion coefficients in the corresponding amplitudes. The method used by Lee, Quigg and Thacker (LQT) \([29]\) prompts us to consider the eigenvalues of this two-particle scattering matrix \([30–32]\). These eigenvalues, labelled as $a_i$, should satisfy the condition $\text{Re}[a_i] < 1/2$. Again, these conditions apply to high scales as well, if we expect perturbativity to hold.

When the quartic part of the scalar potential preserves CP \([33,34]\) and $\mathbb{Z}_2$ symmetries, the LQT eigenvalues are discussed in \([35–37]\). For $\lambda_6, \lambda_7 = 0$, we follow the procedure described

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\(^1\)The conditions are slightly different for the two types of couplings. The reason becomes clear if we note that the perturbative expansion parameter for $2 \rightarrow 2$ processes driven by the quartic couplings is $\lambda_i$. The corresponding parameter for Yukawa-driven scattering processes is $|y_i|^2$
The general formulas including $\lambda_6, \lambda_7$, is given in Appendix B.

The condition to be taken up next is that of vacuum stability. For the scalar potential of a theory to be stable, it must be bounded from below in all directions. This condition is threatened if the quartic part of the scalar potential, which is responsible for its behaviour at large field values, turns negative. Avoiding such a possibility up to any given scale ensures vacuum stability up to that scale. The issue of vacuum stability in context of a 2HDM has been discussed in detail in [38–41].

The 2HDM potential has eight real scalar fields. By studying the behaviour of the quartic part of its scalar potential along different field directions, one arrives at the following conditions [28, 42],

\begin{align*}
\text{vsc}_1 & : \lambda_1 > 0 \quad \text{(3.1a)} \\
\text{vsc}_2 & : \lambda_2 > 0 \quad \text{(3.1b)} \\
\text{vsc}_3 & : \lambda_3 + \sqrt{\lambda_1 \lambda_2} > 0 \quad \text{(3.1c)} \\
\text{vsc}_4 & : \lambda_3 + \lambda_4 - |\lambda_5| + \sqrt{\lambda_1 \lambda_2} > 0 \quad \text{(3.1d)} \\
\text{vsc}_5 & : \frac{1}{2} (\lambda_1 + \lambda_2) + \lambda_3 + \lambda_4 + \lambda_5 - 2|\lambda_6 + \lambda_7| > 0 \quad \text{(3.1e)}
\end{align*}

The couplings in the general $Z_2$ violating Type-II 2HDM evolve from a low scale to a high scale according to a set of renormalization group (RG) equations listed in the Appendix A. If one proposes the UV cut-off scale of the model to be some $\Lambda_{UV}$, it might so happen that the couplings grow with the energy scale and hit the Landau pole before $\Lambda_{UV}$. A second, still unacceptable, possibility is that of the LQT eigenvalues crossing their perturbativity/unitarity limits. The RG evolution of the 2HDM couplings has been recently studied in [43, 44]. Finally, the stability conditions can get violated below $\Lambda_{UV}$, making the scalar potential unbounded from below. All these problems are avoided if one postulates that all of the conditions laid down above are valid up to $\Lambda_{UV}$, which marks the maximum energy up to which the 2HDM can be valid without the intervention of any additional physics.

### 3.2 Higgs mass constraints

The spectrum of a generic 2HDM consists of a charged scalar, a CP-odd neutral scalar and two CP-even neutral scalars. Since the LHC has observed a CP-even neutral boson around 125 GeV, we allow only those regions in the parameter space for which $h$, the lighter neutral scalar, lies in the mass range 123–127 GeV (keeping an error bar). In addition, the charged
scalar is required to have a mass greater than 315 GeV due to low energy constraints, coming mainly from $b \to s\gamma$ \cite{45,46}.

### 3.3 Oblique parameter constraints

The presence of an additional $SU(2)$ doublet having a hypercharge $Y = 1$ modifies the electroweak oblique parameters \cite{47}. It is to be noted that since the couplings of the fermions to gauge bosons remain unaltered even after the introduction of the second doublet, all the additional contributions come from the scalar sector of the 2HDM. The oblique parameters can be decomposed as,

\begin{align}
S &= S_{SM} + \Delta S, \\
T &= T_{SM} + \Delta T,
\end{align}

where $S_{SM}$ and $T_{SM}$ denote the Standard Model (SM) contributions and $\Delta S$ and $\Delta T$ denote any new physics effect. The central value is the contribution coming from the standard model with the reference values $m_{h,\text{ref}} = 126$ GeV and $M_{t,\text{ref}} = 173$ GeV where $M_t$ denotes the pole mass of the top quark. The expressions for $\Delta S$ and $\Delta T$ for a general 2HDM can be found in \cite{16,48–50}. The corresponding bounds we have used are $|\Delta S| < 0.10$ and $|\Delta T| < 0.12$ \cite{51}. While $\Delta S$ hardly imposes any constraint on this scenario, the splitting amongst the scalar masses affects the $T$ parameter. Typically for $m_{12} = 0$, $T$ prevents large mutual splitting among states other than the lightest neutral scalar. For $m_{12} \neq 0$, the scalars other than the light neutral one have masses $\sim m_{12}$. As $m_{12}$ is increased, the masses approach the decoupling limit, and in that case, the $T$ parameter constraint ceases to play a significant role. The oblique parameters constraints turn out to be redundant in our analysis in such a case. The consistency with these parameters has nevertheless been explicitly ensured at each allowed point of the parameter space.

### 3.4 Collider constraints

Apart from the theoretical constraints discussed above, we also strive to find the region of parameter space of a 2HDM allowed by the recent Higgs data. The ATLAS \cite{52–54} and CMS \cite{55} collaborations have measured the production cross section for a $\sim 125$ GeV Higgs multiplied by its branching ratios to various possible channels. In our case, since the underlying theory is a 2HDM, all the cross sections and decay widths get modified compared to the corresponding SM values. For example, the production cross section of
the light neutral Higgs through gluon fusion will get rescaled in the case of a 2HDM due to the fact that the fermionic couplings of the 125 GeV Higgs are now changed with respect to the SM values by appropriate multiplicative factors. Similarly, the loop induced decay $h \to \gamma\gamma$ will now draw an additional contribution from the charged scalars. Some recent investigations in this area, can be found in [56–66]. Also, model-independent analysis of the data, which impose constraints on non-SM couplings of the scalar discovered, have to allow such contributions [67–72]. In order to check the consistency of a 2HDM with the measured rates in various channels, We theoretically compute the signal strength $\mu^i$ for the $i^{th}$ channel using the relation:

$$\mu^i = \frac{R_{\text{prod}} \times R^i_{\text{decay}}}{R_{\text{width}}}$$

(3.3)

Here $R_{\text{prod}}, R^i_{\text{decay}}$ and $R_{\text{width}}$ denote respectively the ratios of the theoretically calculated production cross section, the decay rate to the $i^{th}$ channel and the total decay width for a $\sim 125$ GeV Higgs to their corresponding SM counterparts. Thus, our analysis strategy is to generate a region in parameter space allowed by the constraints coming from vacuum stability, perturbative unitarity and electroweak precision data. We subsequently compute $\mu^i$ for each point in that allowed region and compare them to the experimentally measured signal strengths, $\hat{\mu}^i$, supplied by the LHC. This exercise carves out a sub-region, which is allowed by the recent Higgs data, from the previously obtained parameter space.

For our numerical analysis, we have taken gluon fusion to be the dominant production mode for the SM-like Higgs. As for the subsequent decays of $h$, we have considered all the decay channels mentioned in Table 1. Unless otherwise stated we use $1\sigma$ allowed ranges of $\hat{\mu}^i$.

4 Results with exact discrete symmetry

In this section, we set out to obtain the allowed parameter space of a Type-II 2HDM having an exact $\mathbb{Z}_2$ symmetry consistent with the various theoretical and collider constraints described above. In this particular case, one naturally has $m_{12} = 0$, $\lambda_6, \lambda_7 = 0$. Thus, we scan over the quartic couplings $\lambda_i$ ($i = 1, \ldots, 5$) within their perturbative limits ($\lambda_{1,2} \in [0,4\pi]$ and $\lambda_{3,4,5} \in [-4\pi,4\pi]$) and allow them to evolve from a low scale to a higher scale, designated by $\Lambda_{UV}$. The RG equations for the evolution of all the 2HDM couplings are listed in Appendix A. In our analysis, the scale from which the evolution starts, has been chosen to be the top quark pole mass $M_t = 173.1$ GeV. This pins down the values of the Yukawa couplings at that scale through the relations $y_t(M_t) = \sqrt{2}m_t(M_t)/v_2$ and
\[
y_i(M_t) = \sqrt{2m_i(M_t)/v_1} \quad \text{for } i = b \text{ and } \tau. \quad \text{Here } m_j(M_t) \text{ refers to the running mass of the } j\text{-th fermion at the scale } M_t. \quad \text{We choose } m_t(M_t), m_b(M_t) \text{ and } m_\tau(M_t) \text{ to be 163.3, 4.2 and 1.77 GeV respectively [73, 74].}
\]

We obtain the allowed values of \( \lambda_i(M_t) \) (\( i = 1, \ldots, 5 \)) which, in course of evolution towards \( \Lambda_{UV} \), satisfy all the constraints of perturbativity, unitarity and vacuum stability at all intermediate scales. Choosing \( \Lambda_{UV} = 1 \) TeV and \( \tan \beta = 2 \) yields the following bounds on the quartic couplings,

\[
\lambda_1 \in [1.30, 4.00], \quad \lambda_2 \in [0.35, 1.95], \quad \lambda_3 \in [1.18, 5.70], \quad \lambda_4 \in [-3.50, -0.10], \quad \lambda_5 \in [-3.32, -0.10]
\]

(4.1)

Recent data indicate that \( M_t \) the top quark pole mass is \( [173.07 \pm 0.52 \pm 0.72] \) GeV [75]. Different values for \( M_t \) (within the allowed band) necessarily alter the running masses as well. However, choosing different values of the top quark mass does not cause any noticeable change to the allowed region in the parameter space of the quartic couplings to change significantly, as it is evident from Fig. 1.

We thus see that the uncertainity in the top quark mass measurement has no bearing on the allowed region of the parameter space. This result is a precursor to a much stronger one.

| Channel | Experiment | \( \hat{\mu} \) | Energy in TeV (Luminosity in fb\(^{-1}\)) |
|---------|------------|-----------------|-----------------------------------------------|
| \( h \rightarrow \gamma \gamma \) | ATLAS | 1.55\(^{+0.33}_{-0.28}\) | 7 (4.8) + 8 (20.7) |
| | CMS | 0.78\(^{+0.28}_{-0.26}\) | 7 (5.1) + 8 (19.6) |
| \( h \rightarrow 4l \) | ATLAS | 1.43\(^{+0.40}_{-0.35}\) | 7 (4.6) + 8 (20.7) |
| | CMS | 0.93\(^{+0.29}_{-0.25}\) | 7 (5.1) + 8 (19.7) |
| \( h \rightarrow 2l2\nu \) | ATLAS | 0.99\(^{+0.31}_{-0.28}\) | 7 (4.6) + 8 (20.7) |
| | CMS | 0.72\(^{+0.20}_{-0.18}\) | 7 (4.9) + 8 (19.4) |
| \( h \rightarrow b\bar{b} \) | ATLAS | 0.20\(^{+0.70}_{-0.60}\) | 7 (4.7) + 8 (20.3) |
| | CMS | 1.00\(^{+0.50}_{-0.50}\) | 7 (5.1) + 8 (18.9) |
| \( h \rightarrow \tau\bar{\tau} \) | ATLAS | 1.4\(^{+0.50}_{-0.40}\) | 8 (20.3) |
| | CMS | 0.78\(^{+0.27}_{-0.27}\) | 7 (4.9) + 8 (19.7) |
| \( h \rightarrow 2l2\nu \) | ATLAS | 1.4\(^{+0.70}_{-0.60}\) | 7 (4.6) + 8 (20.7) |
| | CMS | 0.62\(^{+0.58}_{-0.47}\) | 7 (4.9) + 8 (19.4) |

Table 1: The signal strengths in various channels with their 1\( \sigma \) uncertainties.
Figure 1: A comparison of the allowed parameter spaces at $\Lambda_{UV} = 1$ TeV, $\tan \beta = 2$ and $m_{12} = 0$ GeV for two values of $M_t$.

that we obtain in the next sections, namely, the high scale validity of the 2HDM irrespective of the measured value of the top quark mass.

The allowed parameter space in the $\lambda_2 - \lambda_3$ as well as $\lambda_4 - \lambda_5$ plane, using $m_t = 163.3$ GeV has been shown in Fig. 2, where we also show the parameter space allowed by the Higgs data at the $2\sigma$ level. Since we have considered simultaneous variation of all the quartic couplings, these plots are just the two dimensional projections of an allowed region in a five dimensional space. An examination of the plots reveals that while $\lambda_5 < -3.3$ get disallowed on grounds of vacuum stability, $\lambda_5 > -0.2$ is unable to satisfy the charged Higgs mass bounds. The lower bounds on $\lambda_1$ and $\lambda_2$ are due to the constraint on the light neutral Higgs mass, whereas the corresponding upper bounds arise from the violation of unitarity. The various scalar masses (in GeV) are observed to lie in the following range.

$$m_H \in [130, 357], \ m_{H\pm} \in [315, 416], \ m_A \in [78, 448]$$

The above range of the scalar masses are in accordance with the $T$ parameter constraints discussed before. Note that the $m_{12} = 0$ limit of the model comes up with a small region in the parameter space when $m_h$ and $m_H$ are nearly mass degenerate, without violating any other constraint.

To illustrate the RG running of the various couplings, the vacuum stability conditions and the LQT eigenvalues, we choose the following initial conditions.

$$\lambda_1(M_t) = 1.3, \ \lambda_2(M_t) = 0.35, \ \lambda_3(M_t) = 3.00, \ \lambda_4(M_t) = -1.80, \ \text{and} \ \lambda_5(M_t) = -1.20$$

$$m_H \in [130, 357], \ m_{H\pm} \in [315, 416], \ m_A \in [78, 448]$$
Figure 2: The allowed region for an exact $\mathbb{Z}_2$ symmetric Type-II 2HDM when $\Lambda_{UV} = 1$ TeV

The above initial condition corresponds to an allowed region in the $\lambda_i(M_t)$ ($i = 1, \ldots, 5$) parameter space with $\lambda_1(M_t)$ and $\lambda_2(M_t)$ fixed at their lower bounds. Fig. 3 describes the

Figure 3: RG running of $\lambda_i$, the LQT eigenvalues and the stability conditions with the energy scale for $\tan \beta = 2$ and $m_{12} = 0$

RG running of the aforementioned quantities. Since $\lambda_1$ starts evolving from a very large value, it rises steeply with the energy scale, inspite of the attempts of $y_b$ and $y_\tau$ to slow it down. On the other hand, $y_t$, being the dominant Yukawa coupling in this case, prevents
\( \lambda_2 \) to rise as sharply as \( \lambda_1 \). The LQT eigenvalues (see Eq. B.3) \( a_+, b_+ \) and \( c_+ \) evolve in a manner as shown in Fig 3b. Also, the running of the stability conditions is shown in Fig 3c.

This leads to the observation that although \( \lambda_1 \) remains within the perturbative limit upto 10 TeV, the LQT eigenvalue \( |a_+| \) crosses its unitarity limit of \( 8\pi \) beyond that scale. Thus, this example illustrates the interplay among perturbativity and unitarity in determining the UV fate of this scenario and it appears that unitarity often proves stronger as a constraint than perturbativity. Further, it follows from Eq. 2.5 that relatively larger values of \( \tan \beta \) are difficult to accommodate in this case, since one has to comply with the measured value of \( m_h \). Also, one can generally conclude that in order to push the UV limit of 2HDM to higher scales, one must look beyond the exact \( Z_2 \) symmetric case.

5 Results with softly broken discrete symmetry

This section illustrates the effects of the various constraints imposed on the model with non-zero \( m_{12} \), i.e., in presence of a soft \( Z_2 \) symmetry violating term. The RG runnings of the various couplings in the model are just like the ones in exact \( Z_2 \) symmetric case, the only differences being in the expressions for the scalar masses as evident from Eq. 2.5. After performing a scan over the space of the quartic couplings within the perturbative limits, we project the allowed region of the parameter space in the \( \lambda_3 - \lambda_5 \) as well as the \( \lambda_4 - \lambda_5 \) plane which satisfy all the constraints listed in Sec. 3 upto \( \Lambda_{UV} = 10^3, 10^{11}, 10^{16}, 10^{19} \) GeV in Figs. 4, and 5. The benchmarks chosen are \( \tan \beta = 2, 10, 20 \) and \( m_{12} = 200, 1000 \) GeV. Having \( \tan \beta \) higher than in the previous section is possible in this case, so long as \( m_{12} \) is correspondingly large, thus generating an acceptable \( m_h \).

Some recent studies have put constraints on the 2HDM parameters using the current Higgs data. We find that for the benchmark point \( \tan \beta = 2 \) and \( m_{12} = 200 \) GeV, a substantial region of the parameter space allowed by the theoretical constraints also gets allowed by the Higgs data even at the 1\( \sigma \) confidence limit. For other benchmark points, the Higgs data allow the entire region allowed by the theoretical constraints. Hence in those cases, the regions allowed by the Higgs data are not shown separately.

We note here that though \( m_{12} \) does not appear in the RG equations themselves, it indirectly puts constraints on \( \lambda_i \) through the mass constraints. In particular, higher values of \( m_{12} \) push the theory towards the decoupling limit, where the couplings of \( h \) approach the corresponding SM values. Hence, in this case, the 2HDM signal strengths are expected to be closer to the SM values. In Figs. 4, and 5, we also indicate the regions allowed by the
Figure 4: The allowed parameter space for $\Lambda_{UV} = 10^3$ GeV. The Higgs data allowed region has also been shown. The $\tan \beta$ and $m_{12}$ values are shown in the plots.
Figure 5: The allowed parameter spaces for $\Lambda_{UV} = 10^{11}$ (green), $10^{16}$ (blue) and $10^{19}$ GeV (red). The $\tan \beta$ and $m_{12}$ values are shown in the plots.
recent Higgs data. An inspection of the results so obtained shows that as $\Lambda_{UV}$ is pushed towards higher scales, the allowed parameter shrinks, and finally at the Planck scale, it is most constrained. Though the bound obtained on one quartic coupling is correlated to the bounds on the other quartic couplings, one can still make a few remarks on the obtained results. As Fig. 5 shows, for the benchmark point $\tan\beta = 2$, $m_{12} = 200\text{GeV}$, the theory is not UV complete beyond $10^{16}$ GeV. However, higher $\tan\beta$ allow for an UV completion upto the Planck scale.

It is worth commenting further on the bounds on the $\lambda_i$ that the plots in Fig. 5 show. While high $|\lambda_3|$ get disfavoured on grounds of perturbative unitarity, vacuum stability in general disallows high $|\lambda_5|$. Similar is the case with $|\lambda_4|$. Firstly, in a Type-II 2HDM, since the top quark does not couple to $\Phi_1$, $\lambda_1$ always increases with energy and unitarity tends to get violated. Thus, higher the value of $\Lambda_{UV}$ is, tighter is the upper bound on $\lambda_1$. We summarise the bounds on $\lambda_1$ and $\lambda_2$ in table 2.

| $\Lambda_{UV}$   | $m_{12}$ | $\tan\beta$ | $\lambda_1$ | $\lambda_2$ |
|------------------|----------|--------------|--------------|--------------|
| $10^{11}$ GeV    | 200 GeV  | 2            | 0.63         | 0.16-0.70    |
|                  | 10       | 0.72         | 0.24-0.28    |              |
|                  | 20       | 0.65         | 0.21-0.26    |              |
|                  | 1000 GeV | 2            | 0.27         | 0.16-0.48    |
|                  | 10       | 0.72         | 0.24-0.26    |              |
|                  | 20       | 0.65         | $\sim 0.25$ |              |
| $10^{19}$ GeV    | 200 GeV  | 10           | 0.40         | 0.25-0.27    |
|                  | 20       | 0.35         | $\sim 0.25$ |              |
|                  | 1000 GeV | 2            | 0.27         | 0.36-0.48    |

Table 2: $\lambda_1$ and $\lambda_2$ bounds at $\Lambda_{UV} = 10^{11}$ and $10^{19}$ GeV.

The variation observed in the upper bound is due to two different values of $m_{12}$ chosen. Note that although $\tan\beta$ determines the initial conditions for the Yukawa couplings, it does not alter the RG running of $\lambda_1$, and hence the bounds on it, because $y_b(M_t)$ and $y_\tau(M_t)$ still remain very small owing to the smallness of $m_b(M_t)$ and $m_\tau(M_t)$. On the other hand, the top quark affects the running of $\lambda_i$ significantly, and hence also plays around a bit with the allowed parameter space. This fact is illustrated in Fig. 6 where we choose $M_t = 171, 175$ GeV and highlight the difference in the parameter spaces so obtained. However, the role
played by $M_t$ is not as crucial as it is in the SM. Hence, in the subsequent sections, we keep
$M_t = 175$ GeV and $m_t(M_t) = 163.3$ GeV. A small enough initial value of $\lambda_2$ causes $\lambda_2(Q)$
to turn negative at some scale destroying the vacuum stability of the theory in the process.
To illustrate the point better, we choose an initial condition for the quartic couplings at
$\tan \beta = 2$ and $m_{12} = 1000$ GeV, out of the allowed set of couplings which obey all the
imposed constraints up to the $\Lambda_{UV} = 10^{19}$ GeV, and display the RG running of the $\lambda_i$, the
stability conditions and the LQT eigenvalues in Fig. 7.

$$
\begin{align*}
\lambda_1(M_t) &= 0.03, \\
\lambda_2(M_t) &= 0.39, \\
\lambda_3(M_t) &= 0.49, \\
\lambda_4(M_t) &= -0.50, \text{ and } \\
\lambda_5(M_t) &= 0.03 \\
\end{align*}
$$

(5.1)

As indicated in Fig. 7, $a_+(Q)$ grows most sharply amongst the other LQT eigenvalues and
hence violates unitarity just after crossing the Planck scale in this case. Thus it turns out
that $|a_+(Q)| \leq 8\pi$ proves to be the strongest constraint in determining an upper bound on
$|\lambda_i|$.

The most important observation that emerges from this part of the study is that the $2HDM$ can be valid all the way up to the GUT scale or even the Planck scale without the intervention of any new physics. This is true even if the top quark mass is at the upper end of the currently allowed range. The additional quartic couplings can counterbalance the effect of
the Yukawa coupling threatening vacuum stability, while still remaining acceptable from the
angle of perturbativity. It is seen that we get allowed parameter space for $\Lambda_{UV} = 10^{19}$ GeV

Figure 6: A comparison of the allowed parameter spaces at $\Lambda_{UV} = 10^{19}$ GeV, $\tan \beta = 2$ and
$m_{12} = 1000$ GeV for two values of $M_t$. 

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Figure 7: RG running of $\lambda_i$, the LQT eigenvalues and the stability conditions with the energy scale for $\tan \beta = 2$ and $m_{12} = 1000$ GeV.

corresponding to several values of $\tan \beta$ and $m_{12}$. There is, however, a noticeable correlation - large $m_{12}$ tends to favour small values of $\tan \beta$.

The splitting amongst the scalar masses becomes the narrowest at the Planck scale, albeit being dependent on the values of $\tan \beta$ and $m_{12}$. For example, for $\tan \beta = 10$, $m_{12} = 200$ GeV and $\Lambda_{UV} = 10^{19}$ GeV, the masses (in GeV) are observed to lie in the following range,

$$m_H \in [635, 636], \ m_{H^\pm} \in [619, 652], \ m_A \in [618, 653]$$

(5.2)

Current measurements yield a value $0.1184 \pm 0.0007$ for $\alpha_s(M_z)$. In our analysis, we have used $\alpha_s(M_z) = 0.1184$ throughout. However, we demonstrate the effect of a $3\sigma$ variation of $\alpha_s(M_z)$ on the running of $\lambda_2$, the quartic coupling where the effect is expected to be more pronounced compared to the other ones.

We took $\lambda_2(M_t) = 0.39$ in Fig. 8. It is seen that the RG running is not significantly altered even by a $3\sigma$ variation of $\alpha_s(M_z)$. Hence, for any value of $\alpha_s(M_z)$ within this band, the parameter spaces will not change in a major fashion, and whatever constraints apply to $\lambda_2(M_t)$ will continue to be valid rather insensitively to $\alpha_s(M_z)$.

The implications of having a complex $m_{12}$ in the scalar potential [33, 34] is also investigated here. We rewrite the quadratic part of the scalar potential as,

$$V_{quad}(\Phi_1, \Phi_2) = m_{11}^2 \Phi_1^\dagger \Phi_1 + m_{22}^2 \Phi_2^\dagger \Phi_2 - |m_{12}| \left( e^{i\alpha} \Phi_1^\dagger \Phi_2 + e^{-i\alpha} \Phi_2^\dagger \Phi_1 \right).$$

(5.3)

The quartic couplings are kept real as in the previous case. The presence of an arbitrary phase $\alpha$ in $m_{12}^2$, leads to a charged scalar $H^+$, three neutral scalars $H_1$, $H_2$ and $H_3$ which are not eigenstates of CP, and of course the charged and neutral Goldstone bosons. The
masses of the neutral scalars can not be obtained in closed form in this case, rather, the corresponding mass matrix has to be diagonalized numerically. In the process, we choose the lightest neutral scalar, say $H_3$ to be around 125 GeV and the charged scalar to have a mass higher than 300 GeV. The quartic couplings satisfying these conditions are selected and are further constrained by the imposition of the theoretical constraints under RG.

Figure 8: Running of $\lambda_2$ for three different values for $\alpha_s(M_Z)$.

Our conclusion therefore is that the regions in the parameter space of a 2HDM, consistent

Fig. 9 shows the bounds imposed on $\lambda_i$ under RG evolution. We had chosen the values $\alpha = \frac{\pi}{4}$, $|m_{12}| = 200$ GeV and $\tan \beta = 2$ as a benchmark.

Scatter plots in mass planes are also presented Fig. 10. For higher $\Lambda_{UV}$, the bounds on the scalar masses become tighter. To make the effect of the added phase in changing the scalar masses, we also show the mass bounds in the situation with $\alpha = 0$.

Our conclusion therefore is that the regions in the parameter space of a 2HDM, consistent
Figure 10: The allowed regions in mass plane as a function of $\Lambda_{UV}$. The upper and lower two plots correspond to the $\alpha = \frac{\pi}{4}$ and $\alpha = 0$ respectively.
with UV completion at the GUT/Planck scale, are dependent on the phase of the complex parameter(s) of the scalar potential. Together with the less crucial role played by the top mass uncertainty, this is the other important lesson to take home from this section.

6 When the discrete symmetry is broken by quartic terms

We now come to the last part of our study where the $Z_2$ symmetry is broken both at the soft and hard level (i.e., $m_{12}, \lambda_6, \lambda_7 \neq 0$). In this case however, the RG running of the various couplings in the theory is different with respect to the soft breaking case, owing to the introduction of $\lambda_6$ and $\lambda_7$ (see Appendix A for the complete set of RG equations). While scanning the $\lambda_i$ parameter space, we try to reduce the number of free parameters so that the analysis does not become unwieldy. We therefore fix some parameters studied earlier within their allowed ranges. In this spirit, we choose $\lambda_1(M_t) = 0.02$ and $\lambda_6(M_t) = \lambda_7(M_t)$ for computational convenience. $\lambda_1(M_t)$ has been deliberately chosen to be small on purpose so that it respects perturbative unitarity even upto the Planck scale. We, in this case too, project the allowed regions in the $\lambda_3 - \lambda_5$ and the $\lambda_4 - \lambda_5$ planes which satisfy all the conditions upto $\Lambda_{UV} = 10^3, 10^{11}, 10^{16}, 10^{19}$ GeV. The benchmarks are $\tan \beta = 2, 10, 20$ and $m_{12} = 200, 1000$ GeV. The results of the scans are shown in Fig. 12.

In the case where $\Lambda_{UV} = 10^3$ GeV, we show the subregions in the parameter spaces which are also allowed by the recent Higgs data. The major constraint, however, comes from the signal strength corresponding to $h \rightarrow \gamma \gamma$. It is clearly seen in Fig. 11 that $m_{12} = 1000$ GeV allows for a bigger region in the parameter space that is allowed by the Higgs data, compared to what $m_{12} = 200$ GeV does. This is obviously expected, given the fact that a high value of $m_{12}$ takes the theory towards the decoupling limit and hence, the 125 GeV Higgs becomes more SM like in the process.

We demonstrate the UV completion of the hard $Z_2$ violating case by showing the RG evolution of the various quartic couplings and stability conditions upto $\Lambda_{UV} = 10^{19}$ GeV. We choose the following initial condition for the quartic couplings at $\tan \beta = 2$ and $m_{12} = 1000$ GeV.

$$\lambda_1(M_t) = 0.02, \lambda_2(M_t) = 0.48, \lambda_3(M_t) = 0.40, \lambda_4(M_t) = -0.30,$$
$$\lambda_5(M_t) = -0.01, \lambda_6(M_t) = -0.05, \text{and } \lambda_7(M_t) = -0.05$$

(6.1)
Figure 11: The blue region gives the parameter space for $\Lambda_{UV} = 10^3$ satisfying all the constraints coming from the theoretical constraints and the green region shows the parameter space that is allowed by the Higgs data. The $\tan \beta$ and $m_{12}$ values are shown in the plots.
Figure 12: The allowed parameter spaces for $\Lambda_{UV} = 10^{11}$ (grey), $10^{16}$ (green) and $10^{19}$ GeV (blue). The $\tan \beta$ and $m_{12}$ values are shown in the plots.
As shown in Fig. 13, $\lambda_3$ increases most sharply whereas $\lambda_2$ first plunges down due to the top quark effect and then starts increasing. Choosing same initial conditions for $\lambda_6$ and $\lambda_7$
causes their evolutions to become fairly similar. In this section, it should be noted that the allowed parameter spaces found are not expected to be exhaustive as we have not scanned over all $\lambda_i(M_t)$ independently, rather, have put $\lambda_1(M_t) = 0.02$ and $\lambda_6(M_t) = \lambda_7(M_t)$ while doing so. However, given the similar structure of the 1-loop beta functions of $\lambda_6$ and $\lambda_7$, the bounds obtained on them would have not substantially changed even if an independent scanning would have been allowed.

7 Summary and Conclusions

We set out to investigate the high-scale behaviour of a 2HDM. The results are illustrated in the context of a Type-II scenario. We have subjected the parameter space of this model to the theoretical constraints based on perturbativity, unitarity and vacuum stability. The relatively less stringent constraints from oblique parameters, and also the LHC constraints on the signal strength in each decay channel of a Higgs around 125 GeV have also been taken into account.

We find that a 2HDM with a discrete $Z_2$ symmetry (thereby forbidding some cross-terms in the two doublets in the potential) cannot be valid beyond 10 TeV, since otherwise the requirement of keeping one neutral scalar mass around 125 GeV cannot be met. With the discrete symmetry broken, on the other hand, it is possible to fulfill all the constraints over a much larger region of the parameter space. Thus the theory with a 2HDM can distinctly be valid up to energies as high as $10^{16}$ GeV or even the Planck scale, without the intervention of any additional physics. More strikingly, this feature holds irrespectively of the uncertainty in the measured value of the top quark mass, which is in contrast to what is expected in the standard model with a single Higgs doublet. In addition, the uncertainty in the strong coupling $\alpha_s(M_Z)$ cannot destroy such high-scale validity of this scenario. The effect of a CP-violating phase in the potential is also considered, and the claim about high-scale validity is robust against variation of this phase. The allowed regions of the parameter space, in terms of the various quartic couplings as well as the scalar mass eigenvalues, are presented by us in detail, in the light of theoretical as well as collider bounds. The inclusion of $Z_2$-breaking quartic couplings, too, is found to retain the high-scale validity of the theory over a large region.

The results obtained here hold qualitatively for a more general 2HDM as well. The only situation where some departure can take place is where the Yukawa coupling of the bottom quark becomes comparable to, or more than, that of the top quark. Even then the large
number of quartic couplings are mostly effective in rescuing the scenario from an unstable vacuum.

The results presented here are based on one-loop RG equations, in consonance with most similar studies in the context of 2HDM. The main conclusions, however, are not expected to change qualitatively on inclusion of higher-loop effects.

On the whole, our conclusion is that it is possible to live with a 2HDM till scales as high as the Planck mass without any additional physics. While the issue of naturalness remains unaddressed in this statement, it is interesting to see that no current experimental measurement or theoretical restriction can affect high-scale validity, which is not the case for the single-doublet scenario.

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Appendix

A Renormalization group (RG) equations

The RG equations for the gauge couplings, for this model, are given by [28],

\[
16\pi^2 \frac{dg_s}{dt} = -7g_s^3, \quad (A.1a)
\]

\[
16\pi^2 \frac{dg}{dt} = -3g^3, \quad (A.1b)
\]

\[
16\pi^2 \frac{dg'}{dt} = 7g'^3. \quad (A.1c)
\]

Since we want to avoid CP violation coming from the quartic sector of the Higgs potential, we choose to keep \( \lambda_i \) \((i = 1, \ldots, 7)\) real. In that case, the quartic couplings evolve according
to,

\[ 16\pi^2 \frac{d\lambda_1}{dt} = 12\lambda_1^2 + 4\lambda_2^2 + 4\lambda_4 \lambda_4 + 2\lambda_4^2 + 2\lambda_5^2 + 24\lambda_6^2 + \frac{3}{4} (3g^4 + g'^4 + 2g^2 g'^2) \]

\[ -\lambda_1 (9g^2 + 3g'^2 - 12y^2 - 4y'^2) - 12y^4 - 4y'^4, \]  
(A.2a)

\[ 16\pi^2 \frac{d\lambda_2}{dt} = 12\lambda_2^2 + 4\lambda_3^2 + 4\lambda_3 \lambda_4 + 2\lambda_4^2 + 2\lambda_5^2 + 24\lambda_7^2 \]

\[ + \frac{3}{4} (3g^4 + g'^4 + 2g^2 g'^2) - 3\lambda_2 (3g^2 + g'^2 - 4y^2) - 12y^4, \]  
(A.2b)

\[ 16\pi^2 \frac{d\lambda_3}{dt} = (\lambda_1 + \lambda_2) (6\lambda_3 + 2\lambda_4) + 4\lambda_3^2 + 2\lambda_4^2 + 2\lambda_5^2 + 4 (\lambda_6^2 + \lambda_7^2) \]

\[ + 16 (\lambda_6 \lambda_7) + \frac{3}{4} (3g^4 + g'^4 - 2g^2 g'^2) \]

\[ - \lambda_3 (9g^2 + 3g'^2 - 6y^2 - 6y'^2 - 2y^2) - 12y^2 y'^2, \]  
(A.2c)

\[ 16\pi^2 \frac{d\lambda_4}{dt} = 2 (\lambda_1 + \lambda_2) \lambda_4 + 8\lambda_3 \lambda_4 + 4\lambda_4^2 + 8\lambda_5^2 + 10 (\lambda_6^2 + \lambda_7^2) + 4 (\lambda_6 \lambda_7) \]

\[ + 3g^2 g'^2 - \lambda_4 (9g^2 + 3g'^2 - 6y^2 - 6y'^2 - 2y^2) + 12y^2 y'^2, \]  
(A.2d)

\[ 16\pi^2 \frac{d\lambda_5}{dt} = (2\lambda_1 + 2\lambda_2 + 8\lambda_3 + 12\lambda_4) \lambda_5 + 10 (\lambda_6^2 + \lambda_7^2) + 4\lambda_6 \lambda_7 \]

\[ - \lambda_5 (9g^2 + 3g'^2 - 6y^2 - 6y'^2 - 2y^2), \]  
(A.2e)

\[ 16\pi^2 \frac{d\lambda_6}{dt} = (12\lambda_1 + 6\lambda_3 + 8\lambda_4) \lambda_6 + (6\lambda_3 + 4\lambda_4) \lambda_7 + 10\lambda_5 \lambda_6 + 2\lambda_5 \lambda_7 \]

\[ - \lambda_6 (9g^2 + 3g'^2 - 6y^2 - 3y'^2), \]  
(A.2f)

\[ 16\pi^2 \frac{d\lambda_7}{dt} = (12\lambda_2 + 6\lambda_3 + 8\lambda_4) \lambda_7 + (6\lambda_3 + 4\lambda_4) \lambda_6 + 10\lambda_5 \lambda_7 + 2\lambda_5 \lambda_6 \]

\[ - \lambda_7 (9g^2 + 3g'^2 - 9y^2 - 3y'^2 - y^2). \]  
(A.2g)

For the Yukawa couplings, the corresponding set of RG equations are,

\[ 16\pi^2 \frac{dy_b}{dt} = y_b \left( -8g^2 - g'^2 - \frac{9}{4} g^2 - \frac{12}{5} g'^2 - 2y^2 + y'^2 + \frac{1}{2} y'^2 \right), \]  
(A.3a)

\[ 16\pi^2 \frac{dy_t}{dt} = y_t \left( -8g^2 - g'^2 - \frac{17}{12} g^2 - \frac{9}{2} g'^2 + \frac{1}{2} y^2 \right), \]  
(A.3b)

\[ 16\pi^2 \frac{dy_r}{dt} = y_r \left( -\frac{9}{4} g^2 - \frac{15}{4} g^2 + 3y^2 + \frac{5}{2} y'^2 \right), \]  
(A.3c)

### B Lee-Quigg-Thacker unitarity bounds

We perform a coupled channel analysis of $2 \to 2$ scattering involving fields in the scalar sector, to the leading order. The basis of neutral two-particle states is given by,

\[
\left\{ w_1^+ w_2^-, w_2^+ w_1^-, h_1 z_2, h_2 z_1, z_1 z_2, h_1 h_2, h_1 z_1, h_2 z_2, w_1^+ w_1^-, w_2^+ w_2^-, \frac{z_1 z_1}{\sqrt{2}}, \frac{z_2 z_2}{\sqrt{2}}, h_1 h_1, h_2 h_2 \right\}
\]  
(B.1)
For the general \( \lambda_6, \lambda_7 \neq 0 \) case, the \((14 \times 14)\) two-particle scattering matrix is given as follows:

\[
\mathcal{M}_{NC} = \begin{pmatrix}
\mathcal{A}_{7 \times 7} & \mathcal{B}_{7 \times 7} \\
\mathcal{B}_{7 \times 7}^\dagger & \mathcal{C}_{7 \times 7}
\end{pmatrix},
\]

where \( \mathcal{A}, \mathcal{B} \) and \( \mathcal{C} \) are given by,

\[
\mathcal{A}_{7 \times 7} = \begin{pmatrix}
\lambda_3 + \lambda_4 & 2\lambda_5 & \frac{1}{2}(\lambda_4 - \lambda_3) & \frac{1}{2}(\lambda_4 + \lambda_5) & \frac{1}{2}(\lambda_4 + \lambda_3) & 0 \\
2\lambda_5 & \lambda_3 + \lambda_4 & \frac{1}{2}(\lambda_4 - \lambda_5) & \frac{1}{2}(\lambda_4 + \lambda_5) & \frac{1}{2}(\lambda_4 + \lambda_3) & 0 \\
\frac{1}{2}(\lambda_4 - \lambda_5) & \frac{1}{2}(\lambda_4 - \lambda_3) & \lambda_5 & (\lambda_3 + \lambda_4 - \lambda_5) & 0 & 0 \\
\frac{1}{2}(\lambda_4 + \lambda_5) & \frac{1}{2}(\lambda_4 + \lambda_3) & 0 & 0 & (\lambda_3 + \lambda_4 + \lambda_5) & \lambda_5 \\
0 & 0 & \lambda_6 & \lambda_6 & 0 & 0 \\
0 & 0 & 0 & 0 & \lambda_1 & 0
\end{pmatrix},
\]

\[
\mathcal{B}_{7 \times 7} = \begin{pmatrix}
0 & 2\lambda_6 & 2\lambda_7 & \frac{\lambda_6}{\sqrt{2}} & \frac{\lambda_7}{\sqrt{2}} & \frac{\lambda_6}{\sqrt{2}} & \frac{\lambda_7}{\sqrt{2}} \\
0 & 2\lambda_6 & 2\lambda_7 & \frac{\lambda_6}{\sqrt{2}} & \frac{\lambda_7}{\sqrt{2}} & \frac{\lambda_6}{\sqrt{2}} & \frac{\lambda_7}{\sqrt{2}} \\
\lambda_7 & 0 & 0 & 0 & 0 & 0 & 0 \\
\lambda_7 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & \lambda_6 & \lambda_7 & \frac{3\lambda_6}{\sqrt{2}} & \frac{3\lambda_7}{\sqrt{2}} & \frac{\lambda_6}{\sqrt{2}} & \frac{\lambda_7}{\sqrt{2}} \\
0 & \lambda_6 & \lambda_7 & \frac{3\lambda_6}{\sqrt{2}} & \frac{3\lambda_7}{\sqrt{2}} & \frac{3\lambda_6}{\sqrt{2}} & \frac{3\lambda_7}{\sqrt{2}} \\
\lambda_5 & 0 & 0 & 0 & 0 & 0 & 0
\end{pmatrix},
\]

\[
\mathcal{C}_{7 \times 7} = \begin{pmatrix}
\lambda_2 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 2\lambda_1 & (\lambda_3 + \lambda_4) & \frac{\lambda_1}{\sqrt{2}} & \frac{\lambda_2}{\sqrt{2}} & \frac{\lambda_1}{\sqrt{2}} & \frac{\lambda_2}{\sqrt{2}} \\
0 & (\lambda_3 + \lambda_4) & 2\lambda_2 & \frac{\lambda_2}{\sqrt{2}} & \frac{\lambda_2}{\sqrt{2}} & \frac{\lambda_2}{\sqrt{2}} & \frac{\lambda_2}{\sqrt{2}} \\
0 & \frac{\lambda_1}{\sqrt{2}} & \frac{\lambda_2}{\sqrt{2}} & \frac{\lambda_1}{\sqrt{2}} & \frac{3\lambda_1}{2} & \frac{\lambda_1}{2} & \frac{3\lambda_1}{2} \\
0 & \frac{\lambda_2}{\sqrt{2}} & \frac{\lambda_1}{\sqrt{2}} & \frac{\lambda_2}{\sqrt{2}} & \frac{1}{2}(\lambda_3 + \lambda_4 + \lambda_5) & \frac{3\lambda_2}{2} & \frac{1}{2}(\lambda_3 + \lambda_4 - \lambda_5) \\
0 & \frac{\lambda_2}{\sqrt{2}} & \frac{\lambda_1}{\sqrt{2}} & \frac{\lambda_2}{\sqrt{2}} & \frac{1}{2}(\lambda_3 + \lambda_4 - \lambda_5) & \frac{3\lambda_1}{2} & \frac{1}{2}(\lambda_3 + \lambda_4 + \lambda_5) \\
0 & \frac{\lambda_2}{\sqrt{2}} & \frac{\lambda_1}{\sqrt{2}} & \frac{\lambda_2}{\sqrt{2}} & \frac{1}{2}(\lambda_3 + \lambda_4 - \lambda_5) & \frac{3\lambda_1}{2} & \frac{1}{2}(\lambda_3 + \lambda_4 + \lambda_5)
\end{pmatrix}
\]

The constraint imposed by unitarity is then given by \(|a_i| \leq 8\pi\), where \( a_i \) (i = 1, ..., 14) are eigenvalues of the matrix \( \mathcal{M} \). The eigenvalues of \( \mathcal{M} \) are evaluated numerically in the present study. However, in the absence of hard \( \mathbb{Z}_2 \) breaking, i.e., when \( \lambda_6, \lambda_7 = 0 \), the matrix decomposes into blocks and analytical expressions for its eigenvalues can be obtained in simple forms which are listed below.

\[
a_{\pm} = \frac{3}{2}(\lambda_1 + \lambda_2) \pm \sqrt{\frac{9}{4}(\lambda_1 - \lambda_2)^2 + (2\lambda_3 + \lambda_4)^2},
\]

(B.3a)
\[ b_\pm = \frac{1}{2} (\lambda_1 + \lambda_2) \pm \sqrt{\frac{1}{4} (\lambda_1 - \lambda_2)^2 + \lambda_4^2}, \]  
(B.3b)

\[ c_\pm = \frac{1}{2} (\lambda_1 + \lambda_2) \pm \sqrt{\frac{1}{4} (\lambda_1 - \lambda_2)^2 + \lambda_5^2}, \]  
(B.3c)

\[ e_1 = (\lambda_3 + 2\lambda_4 - 3\lambda_5), \]  
(B.3d)

\[ e_2 = (\lambda_3 - \lambda_5), \]  
(B.3e)

\[ f_1 = f_2 = (\lambda_3 + \lambda_4), \]  
(B.3f)

\[ f_+ = (\lambda_3 + 2\lambda_4 + 3\lambda_5), \]  
(B.3g)

\[ f_- = (\lambda_3 + \lambda_5). \]  
(B.3h)

The matrix corresponding to the overall singly charged states,

\[ \{ h_1 w_1^+, h_2 w_1^+, z_1 w_1^+, z_2 w_1^+, h_1 w_2^+, h_2 w_2^+, z_1 w_2^+, z_2 w_2^+ \} \]  
(B.4)

is given by,

\[
M_{CC} = \begin{pmatrix}
\lambda_1 & \lambda_6 & 0 & 0 & \lambda_6 & \frac{1}{2} (\lambda_4 + \lambda_5) & 0 & \frac{i(\lambda_4 - \lambda_5)}{2} \\
\lambda_6 & \lambda_3 & 0 & 0 & \frac{\lambda_4 + \lambda_5}{2} & -i(\lambda_4 - \lambda_5) & \lambda_7 & 0 \\
0 & 0 & \lambda_1 & \lambda_6 & 0 & -i(\lambda_4 - \lambda_5) & \lambda_6 & \frac{\lambda_4 + \lambda_5}{2} \\
0 & 0 & \lambda_6 & \lambda_3 & 0 & i(\lambda_4 - \lambda_5) & \lambda_7 & 0 \\
\frac{\lambda_4 + \lambda_5}{2} & \frac{i(\lambda_4 - \lambda_5)}{2} & \lambda_3 & \lambda_7 & 0 & 0 & 0 \\
\frac{i(\lambda_4 - \lambda_5)}{2} & \frac{\lambda_4 + \lambda_5}{2} & \lambda_7 & 0 & 0 & 0 & 0 \\
\frac{i(\lambda_4 - \lambda_5)}{2} & 0 & \frac{\lambda_4 + \lambda_5}{2} & \lambda_7 & 0 & 0 & 0 \\
0 & 0 & \lambda_6 & \lambda_3 & \lambda_7 & 0 & 0 & \lambda_2
\end{pmatrix}.
\]

Again for the case \( \lambda_6, \lambda_7 = 0 \), the eigenvalues of \( M_{CC} \) are, \( b_\pm, c_\pm, e_2, f_1, f_- \) and \( p = (\lambda_3 - \lambda_4) \).

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