PROCEEDINGS OF SCIENCE

Charged scalars in a Lopsided doublet model

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We consider a two-Higgs-doublet Model (2HDM) with a softly broken $\mathbb{Z}_2$ symmetry where only one of the doublets couples to fermions at tree-level. In addition, the other doublet does not acquire a vacuum expectation value. One can view this model as a generalization of the Inert Doublet Model (IDM), which has an exact $\mathbb{Z}_2$ symmetry. In this paper, the model is presented together with constraints from theory and the oblique parameters S and T. Some implications for collider phenomenology is outlined and in particular, we discuss the charged scalar in this model. At lowest order, the charged scalar decays into a pair of fermions proceed at one loop level. We also consider charged scalar decays into $W^{\pm}Z/\gamma$ which also occur at one loop level. We describe briefly how to calculate and renormalize those processes.

Third International Workshop on Prospects for Charged Higgs Discovery at Colliders - CHARGED2010, September 27-30, 2010
Uppsala Sweden

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1. Generic two-Higgs-doublet model

The two-Higgs-doublet model (2HDM) provides a simple extension of the Higgs sector of the Standard Model (SM) by the inclusion of one additional $SU(2)_L$, $Y=1$, complex scalar field $\Phi(x) = (\Phi^+(x), \Phi^0(x))^T$. The most general renormalizable potential that can be made out of two such doublets which respects the electroweak gauge symmetry is:

$$\begin{aligned} V_{2\mathrm{HDM}} &= m_{11}^2\,\Phi_1^\dagger\Phi_1 + m_{22}^2\,\Phi_2^\dagger\Phi_2 - \left[m_{12}^2\,\Phi_1^\dagger\Phi_2 + \mathrm{h.c.}\right] \\ &+ \frac{1}{2}\lambda_1(\Phi_1^\dagger\Phi_1)^2 + \frac{1}{2}\lambda_2(\Phi_2^\dagger\Phi_2)^2 + \lambda_3(\Phi_1^\dagger\Phi_1)(\Phi_2^\dagger\Phi_2) + \lambda_4(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1) \\ &+ \left\{\frac{1}{2}\lambda_5(\Phi_1^\dagger\Phi_2)^2 + \left[\lambda_6(\Phi_1^\dagger\Phi_1) + \lambda_7(\Phi_2^\dagger\Phi_2)\right](\Phi_1^\dagger\Phi_2) + \mathrm{h.c.}\right\}, \end{aligned} \tag{1.1}$$

where all parameters are real, except for $\lambda_{5,6,7}$ and m_{12}^2 which in general are complex-valued. We constrain ourselves to *CP*-conserving Higgs sectors and all parameters are taken to be real. The potential (1.1) is invariant under global $U(2)$-transformations among the Higgs-doublets. For $\lambda_{6,7}$ and $m_{12}^2 = 0$ the potential (1.1) is invariant under $\mathbb{Z}_2$-transformations of the Higgs-doublets: $\Phi_1 \to \Phi_1$, $\Phi_2 \to -\Phi_2$.

The $\mathbb{Z}_2$ breaking terms in (1.1) are m_{12}^2 which breaks this symmetry softly, *i.e.* the $\mathbb{Z}_2$-symmetry is restored in the UV-limit, and $\lambda_{6,7}$ which explicitly break this symmetry hard.

A generic basis for the vacuum expectation values (VEVs) which respects the unbroken $U(1)_{EM}$ symmetry is $\langle\Phi_1\rangle = (0, v\cos\beta)^T/\sqrt{2}$, $\langle\Phi_2\rangle = (0, v\sin\beta)^T/\sqrt{2}$, where we have introduced $\tan\beta \equiv v_2/v_1$ and $v^2 = v_1^2 + v_2^2$. The parameter $\tan\beta$ is not a physical parameter due to the $U(2)$ invariance of the Higgs Lagrangian. However it will become so when considering fermion-Higgs interactions which break the $U(2)$ invariance of the Higgs Lagrangian.

2. The Lopsided doublet model

We consider a 2HDM where the $\mathbb{Z}_2$ symmetry is softly broken and only Φ_1 have tree-level fermions couplings. In addition, only Φ_1 acquires a VEV, i.e. $v = v_1$, thus we have a physical realization of the so called *Higgs basis*. This model can be thought of as a generalization of the Inert Doublet Model (IDM) where the $\mathbb{Z}_2$ symmetry is exact [1]. The minimization conditions in the Higgs basis reads: $m_{11}^2 = -v^2\lambda_1/2$, $m_{12}^2 = v^2\lambda_6/2$, giving no constraint on m_{22}^2, which is therefore a free parameter in the model. Since the soft-$\mathbb{Z}_2$ breaking parameter m_{12}^2 is proportional to the hard breaking parameter λ_6, it naively seems that if soft breaking is introduced, the minimization condition implies that hard breaking also occurs. We shall see that this is not necessarily the case. Expanding the potential and using the minimization conditions, the mass matrix for the *CP*-even states is

$$M^2 = \begin{pmatrix} v^2\lambda_1 & v^2\lambda_6 \\ v^2\lambda_6 & m_A^2 + v^2\lambda_5 \end{pmatrix}, \tag{2.1}$$

where the masses for the *CP*-odd scalar A and the charged scalar $H^\pm$ are

$$m_{H^\pm}^2 = m_{22}^2 + v^2\lambda_3/2, \quad m_A^2 = m_{H^\pm}^2 - v^2(\lambda_5 - \lambda_4)/2\,. \tag{2.2}$$

For non-zero λ_6, the h and H will be states of indefinite $\mathbb{Z}_2$-parity due to mixing:

$$H = (\sqrt{2}\,\mathrm{Re}\,\Phi_1^0 - v)\cos\alpha + \sqrt{2}\,\mathrm{Re}\,\Phi_2^0\sin\alpha = \varphi_1\cos\alpha + \varphi_2\sin\alpha \tag{2.3}$$

$$h = -(\sqrt{2}\,\mathrm{Re}\,\Phi_1^0 - v)\sin\alpha + \sqrt{2}\,\mathrm{Re}\,\Phi_2^0\cos\alpha = -\varphi_1\sin\alpha + \varphi_2\cos\alpha, \tag{2.4}$$

where α is the angle that diagonalizes the mass matrix (2.1) by an orthogonal transformation. One notes that if $\mathbb{Z}_2$ is conserved and $\lambda_6 = 0$, as in the IDM, the mass matrix is diagonal and there will be no mixing between φ_1 and φ_2. Thus the doublets written in mass eigenstates are

$$\Phi_1 = \frac{1}{\sqrt{2}}\begin{pmatrix} \sqrt{2}G^+ \\ v - h\sin\alpha + H\cos\alpha + \mathrm{i}G^0 \end{pmatrix}, \quad \Phi_2 = \frac{1}{\sqrt{2}}\begin{pmatrix} \sqrt{2}H^+ \\ h\cos\alpha + H\sin\alpha + \mathrm{i}A \end{pmatrix}, \tag{2.5}$$

where G^+, G^0 are the Goldstone bosons. It should be noted that since only Φ_1 acquires a VEV, it is not approriate to call Φ_2 a *Higgs* doublet and the scalars *Higgs* bosons. We will from now on call all the mass eigenstates in this model *scalars*. One should note that the limit $\sin\alpha \to 0$ restores the $\mathbb{Z}_2$-symmetry and we recover the IDM. The *CP*-even scalars are given in terms of the λ_i's according to ($m_H > m_h$) :

$$m_h^2 = m_A^2\cos^2\alpha + v^2\lambda_1\sin^2\alpha + v^2\lambda_5\cos^2\alpha - v^2\lambda_6\sin 2\alpha\,, \tag{2.6}$$

$$m_H^2 = m_A^2\sin^2\alpha + v^2\lambda_1\cos^2\alpha + v^2\lambda_5\sin^2\alpha + v^2\lambda_6\sin 2\alpha\,. \tag{2.7}$$

Using these relations together with the ones for m_A^2 and $m_{H^\pm}^2$ one may solve for $\lambda_{1,3,4,5}$:

$$\lambda_1 = \left[m_H^2 + m_h^2 + \left(m_H^2 - m_h^2\right)/\cos 2\alpha - 2v^2\lambda_6\tan 2\alpha\right]/2v^2, \tag{2.8}$$

$$\lambda_3 = 2\left(m_{H^\pm}^2 - m_{22}^2\right)/v^2, \tag{2.9}$$

$$\lambda_4 = \left[m_H^2 + m_h^2 - \left(m_H^2 - m_h^2\right)/\cos 2\alpha + 2v^2\lambda_6\tan 2\alpha + 2m_A^2 - 4m_{H^\pm}^2\right]/2v^2, \tag{2.10}$$

$$\lambda_5 = \left[m_H^2 + m_h^2 - \left(m_H^2 - m_h^2\right)/\cos 2\alpha + 2v^2\lambda_6\tan 2\alpha - 2m_A^2\right]/2v^2. \tag{2.11}$$

The mixing angle can be written as $\sin 2\alpha = 2v^2\lambda_6/(m_H^2 - m_h^2)$. In the case of maximal mixing, $\sin 2\alpha = 1$, the above formulas for $\lambda_{1,4,5}$ are not valid. The correct formulas for the couplings for this case are found by simply omitting the factors with $\cos 2\alpha$ and $\tan 2\alpha$ in those expressions.

The free parameters of this model are chosen to be $m_h, m_H, m_A, m_{H^\pm}, \lambda_2, \lambda_7, \sin\alpha$ and m_{22}. Using the formalism of Davidson and Haber [2] one can find conditions when $\mathbb{Z}_2$ is only softly broken. It then turns out that it is possible to have soft breaking even with non-zero λ_6, provided that the following conditions are fullfilled:

$$(\lambda_1 - \lambda_2)\left[\lambda_{345}(\lambda_6 + \lambda_7) - \lambda_2\lambda_6 - \lambda_1\lambda_7\right] - 2(\lambda_6 - \lambda_7)(\lambda_6 + \lambda_7)^2 = 0, \tag{2.12}$$

$$(\lambda_1 - \lambda_2)m_{12}^2 + (\lambda_6 + \lambda_7)(m_{11}^2 - m_{22}^2) \neq 0, \tag{2.13}$$

where $\lambda_{345} = \lambda_3 + \lambda_4 + \lambda_5$. These conditions are valid if $\lambda_1 \neq \lambda_2$ in a basis where $\lambda_7 = -\lambda_6$.

The fermions will acquire mass through yukawa couplings with the Higgs doublet Φ_1. If one also assigns $\mathbb{Z}_2$ parities to the fermions in order to avoid flavor changing neutral currents at tree-level [3], one obtains for $\tan\beta = 0$

$$-\mathscr{L}_{\mathrm{mass}} = \frac{M_F}{v}\bar{F}\,F\,H\cos\alpha - \frac{M_F}{v}\bar{F}\,F\,h\sin\alpha \tag{2.14}$$

where $F = U, D, L$.

3. Constraints and Phenomenology

Since the charged scalar and the pseudoscalar do not couple to fermions at tree-level, basically all flavor constraints and some collider constraints do not apply. As a consequence they might have been produced already at LEP if light enough. At this stage we first examine theoretical constraints and constraints from the oblique parameters S,T. We enforce tree-level perturbativity and unitarity, *i.e.* the quartic scalar couplings λ_i should be smaller than 4π. The unitarity condition at tree level means that the S-matrix eigenvalues L_i should not be greater than 16π. Also the requirement that the potential (1.1) should be stable is checked. These constraints and the S,T parameters are evaluated using the software `2HDMC` [4]. The evaluated S,T parameters should fall within 1σ of figure 10.4 in [5] (reference value for SM m_h = 117 GeV). We find large regions in the parameter space which fulfill these constraints as long as the custodial symmetry is valid; $m_{H^\pm} \approx m_A$ or $m^2_{H^\pm} \approx m^2_H \sin^2\alpha + m^2_h \cos^2\alpha$. Figure 1 displays some examples of allowed regions in the $m_{H^\pm}, m_A$ plane with $\lambda_2 = \lambda_1$ and $\lambda_7 = \lambda_6$ in order to fulfill (2.12) and (2.13).

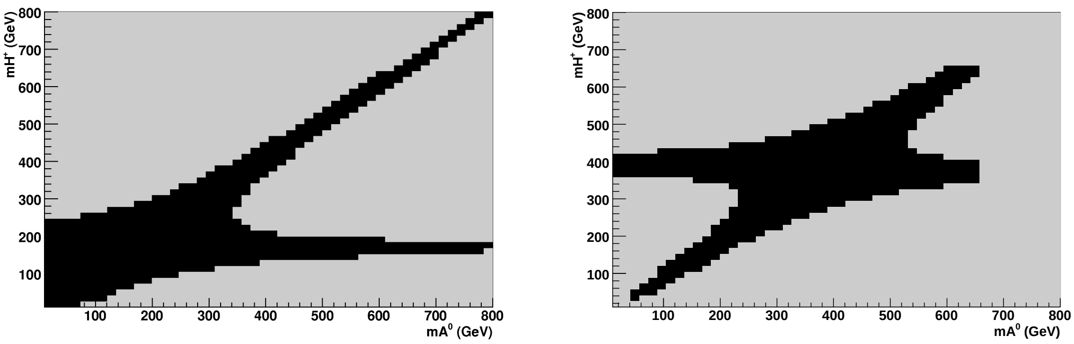

Figure 1: Examples of allowed regions in the $m_{H^\pm}, m_A$ plane from theory and S,T parameter constraints. The left (right) figure is for $m_h = 150(150)$ GeV, $m_H = 250(400)$ GeV, $\sin\alpha = 1/\sqrt{2}\,(0.3)$, $m_{22} = 50(100)$ GeV.

4. Decays of the charged scalar

If kinematically allowed, the decays $H^\pm \to W^\pm S$ ($S = h,H,A$), will dominate since they are the only occuring tree-level processes without internal propagators. The decay width is of the order a few to ten GeV and is shown in figure 2. Also shown in the same figure is the decay width for $H^\pm \to W^{\pm *}S$ which can be of the order MeV and are calculated with [4]. If $H^\pm \to W^{\pm *}S$ is off-shell enough, 3-body decays might compete with 4-body decays and the loop-induced 2-body decays $H^\pm \to f_i\overline{f_j}, W^\pm Z, W^\pm\gamma$.

The diagrams for the loop-induced 2-body decays can be divided into vertex-corrections (figure 3 a,d,e) and mixing-type ones (figure 3 b,c,f). Starting with $H^\pm \to W^\pm Z, W^\pm\gamma$, we follow [6] to obtain the relevant counterterms by expanding the doublets: $\Phi_i \to \sqrt{Z_i}\hat{\Phi}_i$ and their VEV's: $v_i \to \sqrt{Z_i}(\hat{v}_i - \delta_{v_i})$, $i = 1,2$ and ˆ denotes a renormalized quantity. This expansion is then performed in $\mathscr{L}^\Phi{}_{\text{kin}}$. At one-loop order, $Z_i = 1 + \delta_{z_i}$, where δ_{Z_i} is quadratic in the couplings g and g'.

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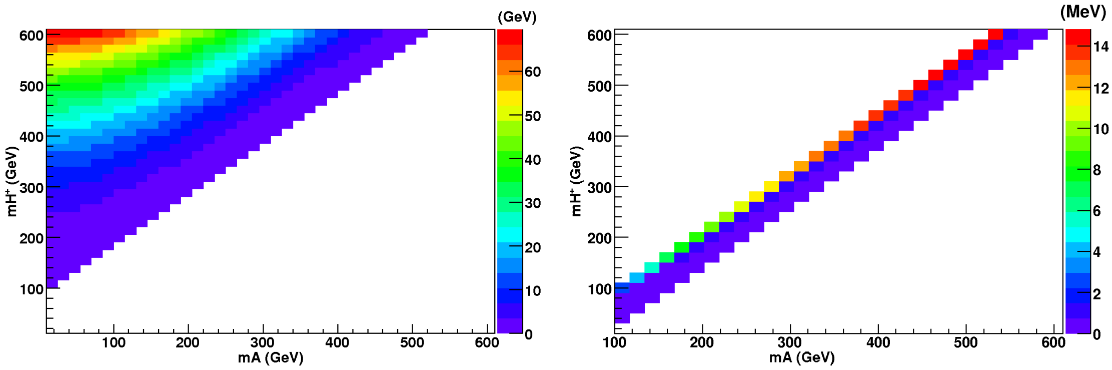

Figure 2: The left (right) figure shows $\Gamma(H^{\pm} \to W^{\pm}A)$, for on-shell (off-shell) $W^{\pm}$. For $H^{\pm} \to W^{\pm}h\,(H)$ the result should scale as $\sin^2\alpha$ $(\cos^2\alpha)$.

The result of the expansion is that the vertex counterterms for $H^{\pm}W^{\mp}Z$ and $H^{\pm}W^{\mp}\gamma$ are proportional to the one for $H^{\pm}W^{\pm}$-mixing, which is finite in dimensional regularization. Finally, we have to take into account $H^{\pm}G^{\pm}$-mixing, which is divergent. The renormalization proceeds by setting the renormalized h and H - tadpoles to zero and requiring that the real part of all renormalized self-energies vanishes on-shell, *e.g.*; $\mathrm{Re}\left[\hat{\Sigma}_{H^{\pm}W^{\mp}}(p^2 = m^2_{H^{\pm}})\right] = 0$. The self-energy for $H^{\pm}G^{\pm}$-mixing is determined on-shell by the self-energy $\hat{\Sigma}_{H^{\pm}W^{\mp}}$ by the Slavnov-Taylor identity.

For the process $H^{\pm} \to f_i\overline{f_j}$, only $H^{\pm}$ mixing with $W^{\pm}$ and $G^{\pm}$ requires renormalization, since the vertex diagrams (figure 3 a) have no counterterms. After renormalization, we require that $\hat{v}_2 = 0$, *i.e.* that we are still in the Higgs basis.

The summation of all diagrams and numerical evaluation of them is performed using `FormCalc` and related packages [7] using an on-shell renormalization scheme. For further discussion regarding renormalization conditions and results, we refer to [8].

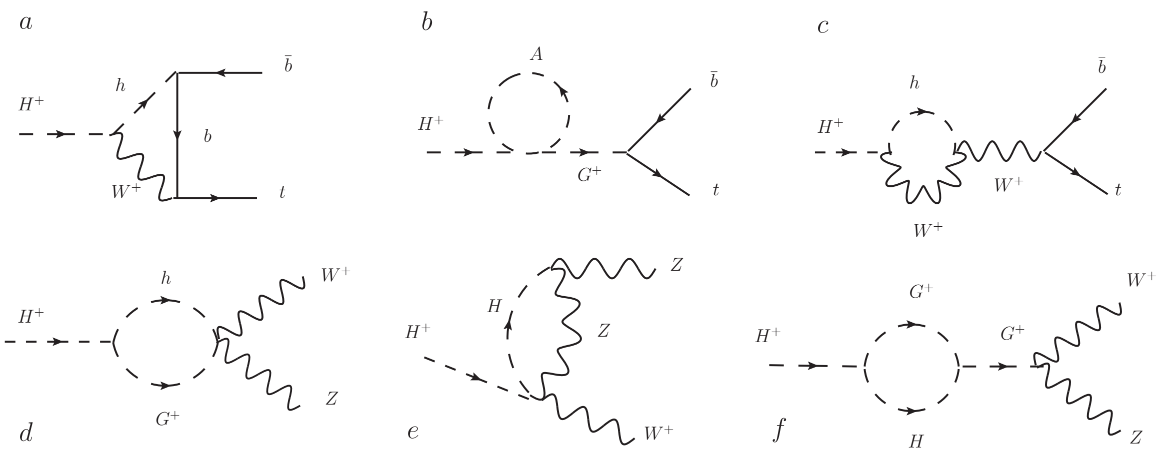

Figure 3: Some examples of Feynman diagrams which can contribute to $H^{\pm} \to f_i\overline{f_j}, W^{\pm}Z$.

5. Summary and Outlook

In this paper, we have presented a model where the charged scalar does not couple to fermions at tree-level but does so at higher orders in perturbation theory. This opens up for an interesting collider phenomenology and we outline how to calculate $H^{\pm}$ decays into two fermions. We show that there are potentially large regions in parameter space which are allowed from theoretical constraints and the S and T parameters. In [8] the model and constraints will be discussed in more detail. Also, the results of the loop-calculations accompanied with 3- and 4 body decays for the charged scalar will be presented.

Acknowledgements The author would like to thank the organizers of this workshop for the opportunity to present this interesting model. Many sincere thanks also to Thomas Hahn and Abdessslam Arhrib for help with FormCalc etc. Finally the author would like to thank his collaborators for helpful input and feedback.

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