Motivations for a stable scalar in an extended Higgs sector

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Physics Challenges in the face of LHC-14

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References

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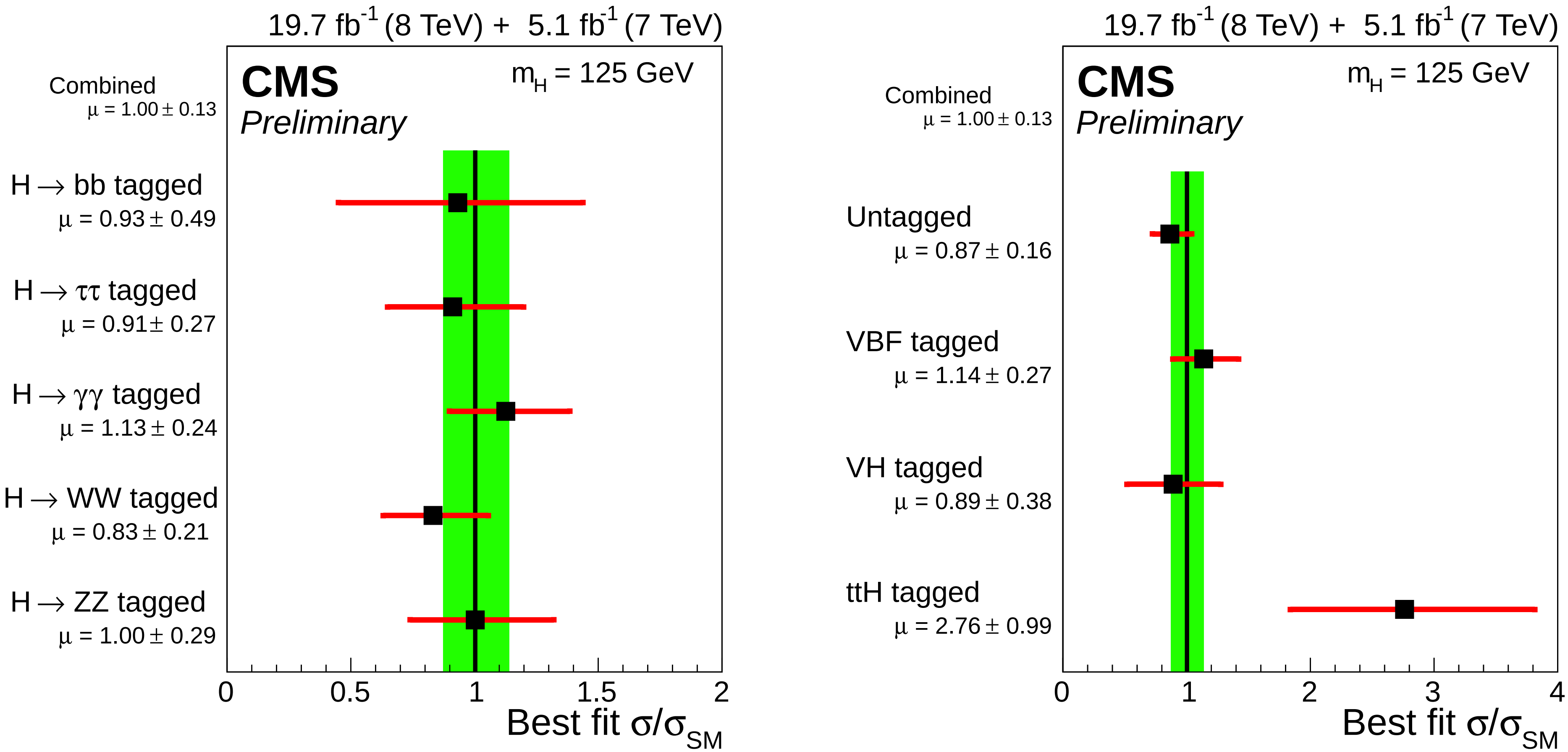
Outline

- A SM-like $h(125)$ and extended Higgs sectors
- Theoretical introduction to the 2HDM
  - The decoupling and alignment limits
- A symmetry origin for the alignment limit
  - The inert doublet model (IDM)
  - A stable scalar as the dark matter
  - benchmarks for collider searches
- Reducing the number of fine tunings in models of extended Higgs sectors
  - The IDM revisited

Evidence that the $h(125)$ is a SM-like Higgs boson



Values of the best-fit $\sigma/\sigma_{\rm SM}$ for the combination (solid vertical line) and for subcombinations by predominant decay mode [left pane], and for analysis tags targeting individual production mechanisms [right pane]. The vertical band shows the overall $\sigma/\sigma_{\rm SM}$ uncertainty. The $\sigma/\sigma_{\rm SM}$ ratio denotes the production cross section times the relevant branching fractions, relative to the Standard Model (SM) expectation. The horizontal bars indicate the 1 standard deviation uncertainties in the best-fit $\sigma/\sigma_{\rm SM}$ values for the individual modes; they include both statistical and systematic uncertainties. Taken from CMS-PAS-HIG-14-009 (July 2014).

Beyond the minimal Higgs sector

There is some motivation for going beyond the minimal Higgs sector.

- The flavor structure of the fermions in the Standard Model (SM) is non-minimal. Why not for the Higgs sector as well?
- The minimal Higgs sector of the MSSM is a two-Higgs doublet model. There is also some motivation for adding additional Higgs singlet fields (e.g. the NMSSM).
- Adding additional Higgs doublets and singlets does not spoil the tree-level relation of $\rho = m_W^2/m_Z^2 \cos^2\theta_W = 1$.
- The scalar sector could provide a stable particle, which would be a candidate for the dark matter.

Any viable extended Higgs sector must possess a SM-like Higgs boson. For simplicity, in this talk I shall focus on the two-Higgs doublet model (2HDM).

A theoretical Introduction to the 2HDM

The scalar fields of the 2HDM are complex SU(2) doublet, hypercharge-one fields, Φ_1 and Φ_2, where the corresponding vevs are $\langle\Phi_i\rangle = v_i/\sqrt{2}$, and $v^2 \equiv |v_1|^2 + |v_2|^2 = (246\ \mathrm{GeV})^2$. The most general renormalizable SU(2)×U(1) scalar potential is given by

$$\mathcal{V} = m_{11}^2\Phi_1^\dagger\Phi_1 + m_{22}^2\Phi_2^\dagger\Phi_2 - \left(m_{12}^2\Phi_1^\dagger\Phi_2 + \mathrm{h.c.}\right) + \tfrac{1}{2}\lambda_1\left(\Phi_1^\dagger\Phi_1\right)^2 + \tfrac{1}{2}\lambda_2\left(\Phi_2^\dagger\Phi_2\right)^2$$
$$+\lambda_3\Phi_1^\dagger\Phi_1\Phi_2^\dagger\Phi_2 + \lambda_4\Phi_1^\dagger\Phi_2\Phi_2^\dagger\Phi_1 + \left[\tfrac{1}{2}\lambda_5\left(\Phi_1^\dagger\Phi_2\right)^2 + [\lambda_6\Phi_1^\dagger\Phi_1 + \lambda_7\Phi_2^\dagger\Phi_2]\Phi_1^\dagger\Phi_2 + \mathrm{h.c.}\right],$$

where m_{12}^2 and λ_5, λ_6 and λ_7 are potentially complex; all other scalar potential parameters are real.

In the most general 2HDM, the fields Φ_1 and Φ_2 are indistinguishable. Thus, it is always possible to define two orthonormal linear combinations of the two doublet fields without modifying any prediction of the model. Performing such a redefinition of fields leads to a new scalar potential with the same form as above but with modified coefficients.

The Higgs basis

It is convenient to define new Higgs doublet fields:

$$H_1 = \begin{pmatrix} H_1^+ \\ H_1^0 \end{pmatrix} \equiv \frac{v_1^* \Phi_1 + v_2^* \Phi_2}{v}\,, \qquad H_2 = \begin{pmatrix} H_2^+ \\ H_2^0 \end{pmatrix} \equiv \frac{-v_2 \Phi_1 + v_1 \Phi_2}{v}\,.$$

It follows that $\langle H_1^0 \rangle = v/\sqrt{2}$ and $\langle H_2^0 \rangle = 0$. This is the *Higgs basis*, which is uniquely defined up to $H_2 \to e^{i\chi} H_2$. The scalar potential is:

$$\mathcal{V} = Y_1 H_1^\dagger H_1 + Y_2 H_2^\dagger H_2 + \left(Y_3 H_1^\dagger H_2 + \text{h.c.}\right) + \tfrac{1}{2}\, Z_1 \left(H_1^\dagger H_1\right)^2 + \tfrac{1}{2} Z_2 \left(H_2^\dagger H_2\right)^2$$

$$+ Z_3 H_1^\dagger H_1 H_2^\dagger H_2 + Z_4 H_1^\dagger H_2 H_2^\dagger H_1 + \left[\tfrac{1}{2} Z_5 \left(H_1^\dagger H_2\right)^2 + [Z_6 H_1^\dagger H_1 + Z_7 H_2^\dagger H_2] H_1^\dagger H_2 + \text{h.c.}\right],$$

where Y_1, Y_2 and $Z_1, \ldots, Z_4$ are real and uniquely defined, whereas Y_3, Z_5, Z_6 and Z_7 are potentially complex and transform under the rephasing of H_2,

$$[Y_3, Z_6, Z_7] \to e^{-i\chi} [Y_3, Z_6, Z_7] \quad \text{and} \quad Z_5 \to e^{-2i\chi} Z_5\,.$$

After minimizing the scalar potential, $Y_1 = -\frac{1}{2} Z_1 v^2$ and $Y_3 = -\frac{1}{2} Z_6 v^2$. This leaves 11 free parameters: 1 vev, 8 real parameters and two relative phases.

The Higgs mass eigenstates

The charged Higgs boson is the charged component of the Higgs-basis doublet H_2, and its mass is given by $m^2_{H^\pm} = Y_2 + \frac{1}{2}Z_3 v^2$. We identify the charged Goldstone bosons with $H_1^\pm$ and the neutral Goldstone boson with $\sqrt{2}\,\mathrm{Im}\,H_1^0$.

The neutral Higgs mass-eigenstates are linear combinations of $\{\sqrt{2}\,\mathrm{Re}\,H_1^0 - v\,,\ \mathrm{Re}\,H_2^0\,,\ \mathrm{Im}\,H_2^0\}$, which are determined by diagonalizing the following squared-mass matrix [cf. H.E. Haber and D. O'Neil, Phys. Rev. D **74**, 015018 (2006)]:

$$\mathcal{M}^2 = v^2 \begin{pmatrix} Z_1 & \mathrm{Re}(Z_6) & -\mathrm{Im}(Z_6) \\ \mathrm{Re}(Z_6) & \frac{1}{2}Z_{345} + Y_2/v^2 & -\frac{1}{2}\mathrm{Im}(Z_5) \\ -\mathrm{Im}(Z_6) & -\frac{1}{2}\mathrm{Im}(Z_5) & \frac{1}{2}Z_{345} - \mathrm{Re}(Z_5) + Y_2/v^2 \end{pmatrix},$$

where $Z_{345} \equiv Z_3 + Z_4 + \mathrm{Re}(Z_5)$. The diagonalizing matrix is a 3×3 real orthogonal matrix that depends on three angles, θ_{12}, θ_{13} and θ_{23}. The corresponding neutral Higgs mass eigenstates will be denoted by h_1, h_2 and h_3 with masses m_1, m_2 and m_3, respectively. Under the rephasing $H_2 \to e^{i\chi}H_2$,

$\theta_{12}\,,\ \theta_{13}$ are invariant, and $\quad \theta_{23} \to \theta_{23} - \chi\,.$

A SM-like Higgs boson

The couplings of $\sqrt{2}\,\mathrm{Re}\,H_1^0 - v$ are precisely those of the SM Higgs boson. However, in general this is not a mass eigenstate. In light of the structure of neutral Higgs squared-mass matrix,

$$\mathcal{M}^2 = v^2 \begin{pmatrix} Z_1 & \mathrm{Re}(Z_6) & -\mathrm{Im}(Z_6) \\ \mathrm{Re}(Z_6) & \frac{1}{2}Z_{345} + Y_2/v^2 & -\frac{1}{2}\mathrm{Im}(Z_5) \\ -\mathrm{Im}(Z_6) & -\frac{1}{2}\mathrm{Im}(Z_5) & \frac{1}{2}Z_{345} - \mathrm{Re}(Z_5) + Y_2/v^2 \end{pmatrix},$$

it follows that there will be one scalar state, h_1, that is SM-like if either:

- $|Z_6| \ll 1$. Indeed if we have $Z_6 = 0$, then the mass eigenstate h_1 is precisely "aligned" with $\sqrt{2}\,\mathrm{Re}\,H_1^0 - v$ and its squared mass is $m_1^2 = Z_1 v^2$. We will call this the **alignment limit**.

- $Y_2 \gg Z_i v^2$. In this case, $m_1^2 \simeq |Z_i| v^2 \ll m_2^2$, m_3^2, $m_{H^\pm}^2$. We call this the **decoupling limit**. (Indeed, in this case h_1 is approximately aligned with $\sqrt{2}\,\mathrm{Re}\,H_1^0 - v$.)

The decoupling/alignment limits in equations

To obtain the conditions in which h_1 is the SM-like Higgs boson, noting that:

$$\frac{g_{h_1 VV}}{g_{h_{\rm SM} VV}} = c_{12}c_{13}\,, \qquad \text{where } V = W \text{ or } Z\,,$$

where $h_{\rm SM}$ is the SM Higgs boson, we demand that

$$s_{12}\,,\, s_{13} \ll 1\,.$$

Here, $s_{12} \equiv \sin\theta_{12}$, $c_{12} \equiv \cos\theta_{12}$, etc. The following (exact) relations are noteworthy [H.E. Haber and D. O'Neil, op. cit.]:

$$\begin{aligned}
{\rm Re}(Z_6\, e^{-i\theta_{23}})\, v^2 &= c_{13}s_{12}c_{12}(m_2^2 - m_1^2)\,, \\
{\rm Im}(Z_6\, e^{-i\theta_{23}})\, v^2 &= s_{13}c_{13}(c_{12}^2 m_1^2 + s_{12}^2 m_2^2 - m_3^2)\,, \\
{\rm Im}(Z_5\, e^{-2i\theta_{23}})\, v^2 &= 2s_{12}c_{12}s_{13}(m_2^2 - m_1^2)\,.
\end{aligned}$$

For simplicity, we assume no mass degeneracies in the neutral scalar sector.

In the decoupling/alignment limit, $m_1^2 \simeq Z_1 v^2$ and

$$\boxed{\begin{aligned} s_{12} \equiv \sin\theta_{12} &\simeq \frac{\mathrm{Re}(Z_6 e^{-i\theta_{23}})v^2}{m_2^2 - m_1^2} \ll 1\,, \\ s_{13} \equiv \sin\theta_{13} &\simeq -\frac{\mathrm{Im}(Z_6 e^{-i\theta_{23}})v^2}{m_3^2 - m_1^2} \ll 1\,. \end{aligned}}$$

<u>The alignment limit</u> [Craig, Galloway and Thomas]

In the limit of $Z_6 \to 0$, s_{12} and s_{13} are small since the corresponding numerator factors are vanishingly small.

<u>The decoupling limit</u> [Haber and Nir]

In the limit of $Y_2 \gg v$ [with $|Z_i| \lesssim \mathcal{O}(1)$], we have m_2^2, $m_3^2 \gg m_1^2, |Z_6|v^2$. Hence s_{12} and s_{13} are small since the corresponding denominator factors are much larger than their respective numerator factors.

Alignment without decoupling

In the decoupling limit, $m_{h_1}^2 \simeq Z_1 v^2 \ll m_2, m_3, m_{H^\pm}$. However it may be difficult to discover the heavy h_2, h_3 and $H^\pm$ at the LHC. For example, in the case of the MSSM Higgs sector, there is no known way to observe the heavy Higgs states in the infamous "wedge region" of large $m_{H^\pm}$ and moderate $\tan\beta$.

In contrast, in the alignment limit with $|Z_6| \ll 1$ without decoupling, the masses of all Higgs bosons [h_1, h_2, h_3 and $H^\pm$] are of $\mathcal{O}(v)$. Thus the prospects for the discovery of additional Higgs scalars at the LHC may be significantly improved. Discovery at the ILC is assured if all scalar masses lie slightly below $\frac{1}{2}\sqrt{s}$.

Indeed if $Z_6 = 0$, then h_1 is exactly the SM-like Higgs boson, and the states of the second Higgs doublet may be accessible at the LHC and/or ILC.

A symmetry origin for the alignment limit

The exact alignment limit appears to required an unexplained tuning to achieve $Z_6 = 0$. However, there exists a symmetry that can impose the exact alignment limit—a discrete $\mathbb{Z}_2$ symmetry where the Higgs basis field H_1 is unchanged but $H_2 \to -H_2$. If we impose this symmetry on the scalar potential, then it follows that

$$Y_3 = Z_6 = Z_7 = 0\,.$$

Note that the minimum condition $Y_3 = -\frac{1}{2}Z_6 v^2$ requires that $Y_3 = 0$ if $Z_6 = 0$, so this $\mathbb{Z}_2$ symmetry *cannot* be softly broken.*

So far, we have not discussed the Higgs-fermion interactions. Having imposed the above $\mathbb{Z}_2$ symmetry in the bosonic sector of the theory, we can extend it to the Yukawa interactions. In general, there are a number a possible ways of doing this (Type-I, Type-II, ...). But to maintain the SM behavior of the H_1 couplings to fermions, the choice is unique.

*A paper on the arXiv claims otherwise, but the premise of this paper seems dubious.

To achieve SM couplings of $h_1 = \sqrt{2}\,\mathrm{Re}\ H_1^0 - v$ to the fermions, we must allow Yukawa couplings of the form

$$-\mathscr{L}_{\rm YUK} = \overline{Q}_L^{(0)} \widetilde{H}_1 \kappa_U^{(0)} U_R^0 + \overline{Q}_L^{(0)} H_1 \kappa_D^{(0)\,\dagger} D_R^{(0)} + \text{h.c.}\,,$$

where $\widetilde{H}_1 \equiv i\sigma_2 H_1^*$, $Q = (U, D)$ with $U = (u, c, t)$ and $D = (d, s, b)$. The superscript zeros indicate weak interaction eigenstates. The corresponding quark mass eigenstate fields are determined after diagonalizing the 3×3 quark mass matrices, $M_U^{(0)} = v\kappa_U^{(0)}/\sqrt{2}$ and $M_D^{(0)} = v\kappa_D^{(0)}/\sqrt{2}$. The Yukawa couplings above are allowed if all fermion fields are unchanged by the $\mathbb{Z}_2$ symmetry transformation. In this case, since H_2 is odd under the $\mathbb{Z}_2$ symmetry, it follows that the Yukawa couplings of H_2 to the fermions must be absent.

This is the inert doublet model (IDM). The lightest scalar of the H_2 doublet is therefore absolutely stable and a possible candidate for dark matter. That is, the inert 2HDM is a Type-I 2HDM in which there exists an unbroken $\mathbb{Z}_2$ symmetry in the Higgs basis.

Further details on the IDM

By imposing the discrete $\mathbb{Z}_2$ symmetry, the scalar potential is CP-conserving. The SM Higgs state is $h = \sqrt{2}\,\mathrm{Re}\; H_1^0 - v$. The inert doublet is

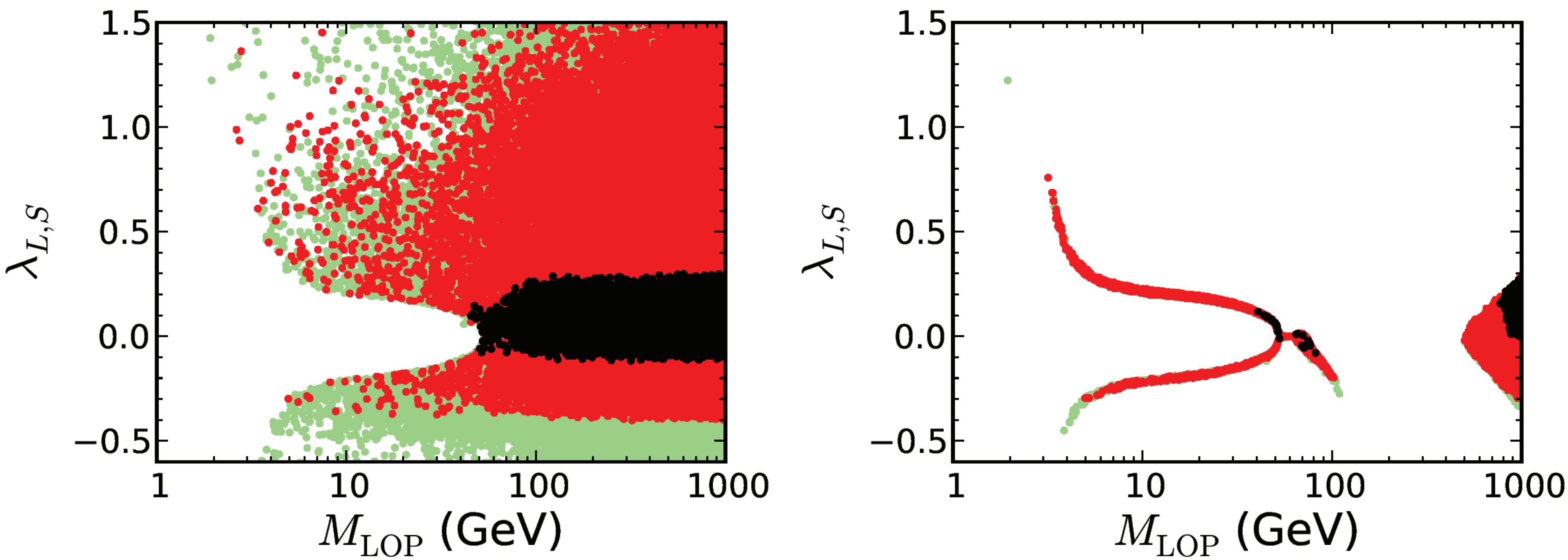
$$H_2 = \begin{pmatrix} H^+ \\ (H + iA)/\sqrt{2} \end{pmatrix} ,$$

where the mass eigenstates consist of two neutral scalars, H, A and a charged Higgs pair. One is free to rephase H_2 so that the model depends on Y_2, Z_1, Z_2, Z_3, Z_4 and $|Z_5|$. Vacuum stability implies that $Z_{1,2} > 0\,,\;\; Z_3 > -(Z_1 Z_2)^{1/2}$ and $Z_3 + Z_4 \pm |Z_5| > -(Z_1 Z_2)^{1/2}$. The physical Higgs masses are

$$m_h^2 = Z_1 v^2\,, \quad m_{H^\pm}^2 = Y_2 + \tfrac{1}{2} Z_3 v^2\,, \quad m_{H,A}^2 = m_{H^\pm}^2 + \tfrac{1}{2}(Z_4 \pm |Z_5|)v^2\,.$$

H and A have opposite CP-quantum numbers, but there is no interaction that can determine separate CP quantum number for these states. If one of these two neutral scalars, call it H_L, is to be a dark matter candidate, then we must demand that the lighter neutral is lighter than $H^\pm$, which yields $Z_4 < |Z_5|$.

Candidate for dark matter: the lightest $\mathbb{Z}_2$–odd particle (LOP)



The viable IDM parameter space projected on the $(M_{\rm LOP}\,,\,\lambda_{L,S})$ plane imposing only the upper limit (left) and the upper and lower limits (right) of the WMAP range, $0.1018 \leq M_{\rm LOP}h^2 \leq 0.1234$. The green points correspond to all valid points in the scan, while the red and black regions show the points which remain valid when the model satisfies stability and perturbativity up to a scale $\Lambda = 10^4$ GeV and the GUT scale $\Lambda = 10^{16}$ GeV, respectively. Taken from A. Goudelis, B. Herrmann and O. Stål, JHEP **1309** (2013) 106.

Assuming that $Z_4 < |Z_5|$, then we identify the squared-mass of the LOP as $M^2_{\rm LOP} = Y_2 + \frac{1}{2}(Z_3 + Z_4 - |Z_5|)v^2$. Above, $\lambda_{L,S} \equiv \frac{1}{2}(Z_3 + Z_4 - |Z_5|)$; when multiplied by v the latter corresponds to the hH_LH_L coupling.

Including the exclusion limits from the current dark matter direct detection experiments, a cosmologically relevant LOP is ruled out by Goudelis, Herrmann and Stål for all LOP masses below 500 GeV except for a narrow window around $\frac{1}{2}m_h$ (roughly 50 GeV $\lesssim M_{\rm LOP} \lesssim 80$ GeV).

LHC Benchmarks (work in progress by Haber and Stål)

We are formulating benchmark points for LHC studies of the IDM (assuming the LOP abundance is below the WMAP limits). One can search for the charged states $H^\pm$ and the heavier neutral state H_H. The production of H_L results in a missing energy signature.

- Regions of parameter space where $h \to H_L H_L$ is allowed by present data.

- Drell-Yan production of H^+H^-, followed by $H^\pm \to W^\pm + H_L$ and/or $H^\pm \to W^\pm + H_H$ followed by $H_H \to Z + H_L$.

- Drell-Yan production of $H_H H_L$, followed by $H_H \to Z + H_L$ ("mono-Zs") and/or $H_H \to W^\mp + H^\pm$ followed by $H^\pm \to W^\pm + H_L$.

Drell-Yan cross-sections at the LHC are small, so to see the relevant missing energy signals above background will require very large data samples. If the W and Z bosons in the decays are produced off-shell, these signals may be impossible to resolve.

In cases where the LHC cannot detect the inert doublet, the ILC may be useful. ILC benchmark points for such a scenario have been proposed by M. Aoki, S. Kanemura and H. Yokoya, Phys. Lett. B **725**, 302 (2013). The following results are given in the ILC Higgs White Paper for the case of $M_{\rm LOP} = 65$ GeV and a rather light inert doublet spectrum.

	Inert scalar masses [GeV]			ILC cross sections [$\sqrt{s} = 250$ GeV (500 GeV)]	
	m_{H_L}	m_{H_H}	$m_{H^\pm}$	$\sigma_{e^+e^-\to H_L H_H}$ [fb]	$\sigma_{e^+e^-\to H^+H^-}$ [fb]
(I)	65	73	120	152 (47)	11 (79)
(II)	65	120	120	74 (41)	11 (79)
(III)	65	73	160	152 (47)	0 (53)
(IV)	65	160	160	17 (35)	0 (53)

Reducing the number of fine-tunings of the 2HDM

Draper, Ruderman and I are exploring 2HDM scenarios in which the number of fine-tunings can be reduced from two to one. Introduce a new $\mathbb{Z}_2$ symmetry that exchanges the two Higgs doublet (in addition to the $\mathbb{Z}_2$ symmetry under which one of the two Higgs doublets changes sign). These two symmetries taken together can reduce the number of fine-tunings of the 2HDM to one. But, this $\mathbb{Z}_2 \times \mathbb{Z}_2$ symmetry cannot be maintained by the SM Yukawa couplings alone. To overcome this problem, we introduce mirror vector-like fermions. However, the mass terms for the vector-like fermions softly break the exchange $\mathbb{Z}_2$ symmetry (since the latter interchanges fermions with their mirrors).

Because the symmetry breaking is soft, the second fine-tuning can remain "under control" (when radiative corrections are considered) effectively leaving one fine-tuning to get the full 2HDM mass spectrum significantly below the ultraviolet cutoff scale. The IDM is one of the possible electroweak symmetry breaking phases that can arise in this framework. Details to follow ...