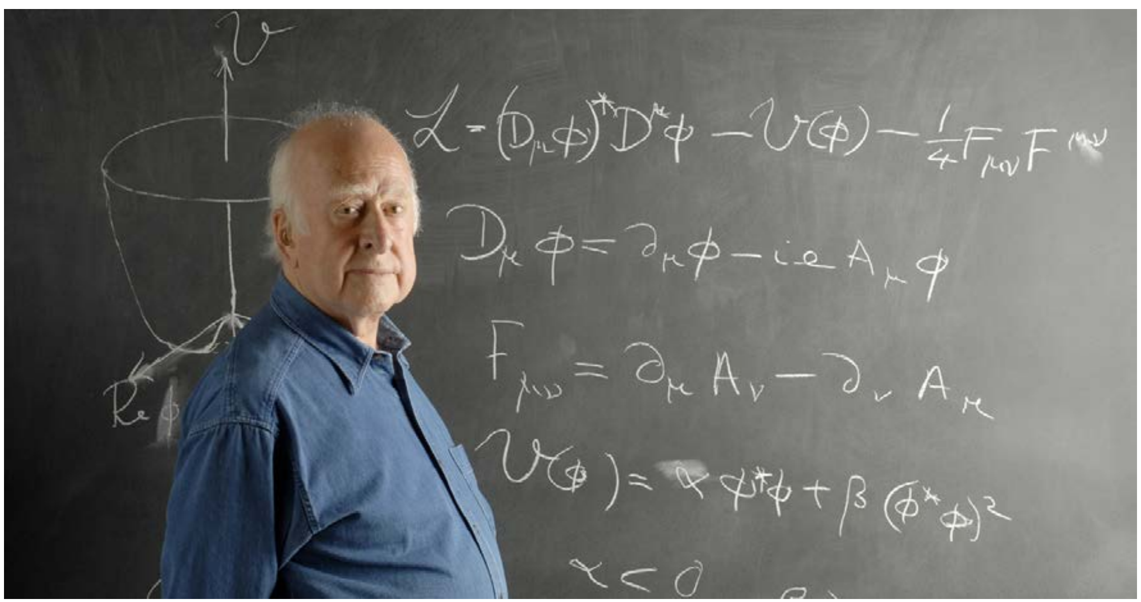
A Higgs Hunter's Perspective

The Higgs Centre
for Theoretical Physics
University of Edinburgh

Howard E. Haber
The Higgs Symposium
9—11 January 2013



The LHC Discovery of 4 July 2012

The CERN update of the search for the Higgs boson, simulcast at ICHEP-2012 in Melbourne, Australia

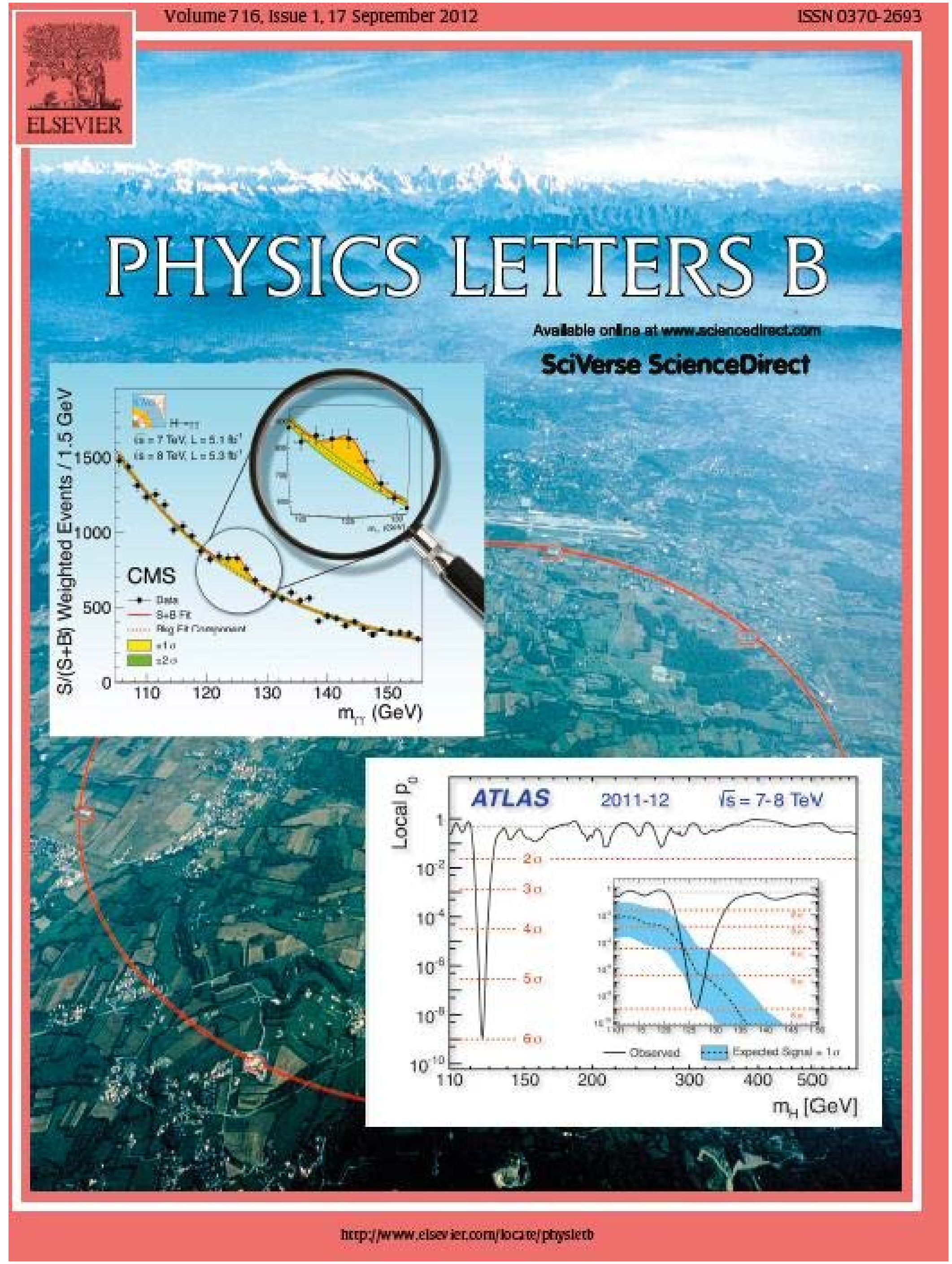
The discovery of the new boson is published in Physics Letters B.

ATLAS Collaboration:

Physics Letters B716 (2012) 1—29

CMS Collaboration:

Physics Letters B716 (2012) 30—61




REACTIONS TO THE LATEST HIGGS BOSON ANNOUNCEMENT...
MAYBE WE WILL BEGIN TO UNDERSTAND HOW MATTER HOLDS TOGETHER!
MAYBE WE CAN DEVELOP A NEW GENERATION OF WEAPONS!
GOSH, I WONDER WHAT KIM KARDASHIAN IS DOING RIGHT NOW.
SCIENTISTS
MILITARY
THE PUBLIC

Outline

1. Higgs Hunting—the early years
2. Beyond the Standard Model Higgs boson
3. The Two-Higgs Doublet Model (2HDM)
4. The Higgs sector of the MSSM
5. The Decoupling Limit
6. A cautionary tale: the wrong Higgs couplings
7. Basis-independent studies of the 2HDM

My first paper on the experimental implications of the Higgs boson was published nearly 35 years ago, and served as a basis for my Ph.D. thesis. In this paper, we considered the possibility of detecting observable consequences of a very massive Higgs boson (near the unitarity limit).

Nuclear Physics B144 (1978) 525–535
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WILL LARGE WEAK INTERACTION EFFECTS BE OBSERVABLE AT VERY HIGH ENERGIES? *

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Received 2 June 1978

We examine the possibility that some weak interaction cross sections become large at very high energies. Such effects could lead to observable weak interaction phenomena in very high energy pp and $p\overline{p}$ colliding beams. Apart from direct production of W's and Z's, the possibility of observable weak effects requires the Higgs mass to be much larger than the mass of the vector bosons. However, for feasible experiments such effects are suppressed in many popular gauge theories, and we expect that observable weak cross sections will not get large at high energies. Theories with more than one Higgs boson may allow large observable effects.

At the time, the experimental challenges presented by the Higgs boson of the Standard Model seemed daunting. These had been spelled out in a ground breaking paper by John Ellis, Mary K. Gaillard and Dmitri Nanopoulos a few years earlier, assuming that the collider was energetic enough to produce Higgs bosons which would be subsequently observed via their decay products.

Veltman argued that below the Higgs boson threshold, the indirect effects of Higgs bosons were especially difficult to observe due to the logarithmic sensitivity to the Higgs mass in gauge boson interactions. We extended this analysis to processes involving fermions, although we did not anticipate the existence of a top quark with a mass larger than the W and Z bosons.

Nuclear Physics B106 (1976) 292–340

A PHENOMENOLOGICAL PROFILE OF THE HIGGS BOSON

John ELLIS, Mary K. GAILLARD * and D.V. NANOPOULOS **
CERN, Geneva

Received 7 November 1975

A discussion is given of the production, decay and observability of the scalar Higgs boson H expected in gauge theories of the weak and electromagnetic interactions such as the Weinberg-Salam model. After reviewing previous experimental limits on the mass of

We should perhaps finish with an apology and a caution. We apologize to experimentalists for having no idea what is the mass of the Higgs boson, unlike the case with charm [3,4] and for not being sure of its couplings to other particles, except that they are probably all very small. For these reasons we do not want to encourage big experimental searches for the Higgs boson, but we do feel that people performing experiments vulnerable to the Higgs boson should know how it may turn up.

We were then motivated to examine extended Higgs sectors.

- New experimental Higgs observables and signals
- Enhanced Higgs couplings to fermions (additional sources for fermion mass terms)
- New sources of CP-violation
- Needed in the MSSM to avoid higgsino anomalies

This led us to the development of the two-Higgs doublet model with Type-I Higgs-fermion interactions.

Nuclear Physics B161 (1979) 493–532

THE FERMION MASS SCALE AND POSSIBLE EFFECTS OF HIGGS BOSONS ON EXPERIMENTAL OBSERVABLES

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(Revised 5 July 1979)

We consider a conventional SU(2) ⊗ U(1) gauge theory with two (or more) Higgs doublets, but with the fermion mass scale determined by the vacuum expectation values of the Higgs particles rather than determined only by widely differing Higgs couplings. Such an alternative to the standard theory cannot be excluded by current data; the Higgs-fermion coupling is allowed to be at least 70 times that of the standard Weinberg-Salam theory. In such a model, one has the possibility of observing large and interesting effects due to the Higgs particles in the theory. These include decays of heavy quarkonium states and Drell-Yan production of $\ell^+\ell^-$ ($\ell = e, \mu, \tau$). Restrictions due to charged Higgs scalar currents (in π, μ and β decay and in neutrino production) take unexpected forms and are not too stringent. The best place to search for a Higgs is in $K^+N \to \mu^+\mu^-X$ (subsect. 3.5).

Which extended Higgs sectors should we consider?

For an arbitrary Higgs sector, the tree-level ρ-parameter is given by

$$\rho_0 \equiv \frac{m_W^2}{m_Z^2 \cos^2\theta_W} = 1 \quad \Longleftrightarrow \quad (2T+1)^2 - 3Y^2 = 1\,,$$

independently of the Higgs vevs, where T and Y specify the weak-isospin and the hypercharge of the Higgs representation to which it belongs. Y is normalized such that the electric charge of the scalar field is $Q = T_3 + Y/2$. The simplest solutions are Higgs singlets $(T,Y) = (0,0)$ and hypercharge-one complex Higgs doublets $(T,Y) = (\frac{1}{2}, 1)$.

Thus, I shall focus on multi-Higgs models with additional Higgs doublets. The simplest of such model extensions is the two-Higgs doublet model (2HDM).

Extended Higgs sectors present an immediate challenge when coupled to fermions. Namely, one must avoid potential flavor-changing neutral currents (FCNCs) due to tree-level neutral Higgs exchange, which would be in conflict with experimental data.

Example: the Higgs-fermion Yukawa couplings in the 2HDM

In a generic basis, the 2HDM Higgs-fermion Yukawa Lagrangian is:

$$-\mathscr{L}_{\rm Y} = \overline{U}_L \Phi_a^{0\,*} h_a^U U_R - \overline{D}_L K^\dagger \Phi_a^- h_a^U U_R + \overline{U}_L K \Phi_a^+ h_a^{D\,\dagger} D_R + \overline{D}_L \Phi_a^0 h_a^{D\,\dagger} D_R + \text{h.c.}\,,$$

where K is the CKM mixing matrix, and there is an implicit sum over $a = 1, 2$. The $h^{U,D}$ are 3×3 Yukawa coupling matrices and

$$\langle \Phi_a^0 \rangle \equiv \frac{v_a}{\sqrt{2}}\,, \quad (a = 1, 2)\,, \qquad v^2 \equiv v_1^2 + v_2^2 = (246~\text{GeV})^2\,.$$

In the Standard Model, the diagonalization of the quark mass matrix automatically diagonalizes the neutral Higgs—fermion interactions. In the 2HDM, this is not possible for a general set of Yukawa coupling matrices.

Glashow and Weinberg proved that the diagonalization of the quark mass matrix automatically diagonalizes the neutral Higgs—fermion interactions if at most one neutral Higgs field couples to fermions of a given electric charge.

The Glashow—Weinberg conditions are satisfied for:

- Type-I Yukawa couplings: $h_2^U = h_2^D = 0$,
- Type-II Yukawa couplings: $h_1^U = h_2^D = 0$,

which can be enforced by a discrete symmetry [or supersymmetry].

The 2HDM received a huge boost when the minimal supersymmetric extension of the Standard Model (MSSM) became a leading candidate for providing a *natural* explanation for the electroweak symmetry breaking scale relative to that of the Planck scale.

- A supersymmetric Higgs-fermion interaction is necessarily holomorphic, implying the need for two Higgs doublets to provide masses for both up-type and down-type fermions.

- The fermionic higgsino partners of the $Y = \pm 1$ Higgs doublets is vector-like and hence anomaly free.

- The superymmetric constraints on the Higgs interactions yields important relations among Higgs couplings and Higgs masses.

The tree-level MSSM Higgs sector

The Higgs sector of the MSSM is a 2HDM, whose Yukawa couplings and Higgs potential are constrained by SUSY. Instead of employing to hypercharge-one scalar doublets $\Phi_{1,2}$, it is more convenient to introduce a $Y=-1$ doublet $H_d \equiv i\sigma_2\Phi_1^*$ and a $Y=+1$ doublet $H_u \equiv \Phi_2$:

$$H_d = \begin{pmatrix} H_d^1 \\ H_d^2 \end{pmatrix} = \begin{pmatrix} \Phi_1^{0\,*} \\ -\Phi_1^- \end{pmatrix}, \qquad H_u = \begin{pmatrix} H_u^1 \\ H_u^2 \end{pmatrix} = \begin{pmatrix} \Phi_2^+ \\ \Phi_2^0 \end{pmatrix}.$$

Yukawa couplings (Type-II):

$$\mathcal{L}_{\rm Yukawa} = -h_u^{ij}(\bar{u}_R^i u_L^j H_u^2 - \bar{u}_R^i d_L^j H_u^1) - h_d^{ij}(\bar{d}_R^i d_L^j H_d^1 - \bar{d}_R^i u_L^j H_d^2) + {\rm h.c.}$$

The Higgs potential of the MSSM is:

$$V = \left(m_d^2 + |\mu|^2\right) H_d^{i*} H_d^i + \left(m_u^2 + |\mu|^2\right) H_u^{i*} H_u^i - m_{ud}^2 \left(\epsilon^{ij} H_d^i H_u^j + \text{h.c.}\right)$$
$$+\tfrac{1}{8}\left(g^2 + g'^2\right)\left[H_d^{i*} H_d^i - H_u^{j*} H_u^j\right]^2 + \tfrac{1}{2} g^2 |H_d^{i*} H_u^i|^2 \,,$$

where $\epsilon^{12} = -\epsilon^{21} = 1$ and $\epsilon^{11} = \epsilon^{22} = 0$, and the sum over repeated indices is implicit. Above, μ is a supersymmetric Higgsino mass parameter and m_d^2, m_u^2, m_{ud}^2 are soft-supersymmetry-breaking masses. The quartic Higgs couplings are related to the SU(2) and $\mathrm{U(1)_Y}$ gauge couplings as a consequence of SUSY.

Minimizing the Higgs potential, the neutral components of the Higgs fields acquire vevs:‖

$$\langle H_d \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} v_d \\ 0 \end{pmatrix}, \qquad \langle H_u \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_u \end{pmatrix},$$

where $v^2 \equiv v_d^2 + v_u^2 = 4m_W^2/g^2 = (246~\text{GeV})^2$. The ratio of the two vevs is an important parameter of the model:

$$\tan\beta \equiv \frac{v_u}{v_d}\,, \qquad 0 \leq \beta \leq \tfrac{1}{2}\pi\,.$$

‖The phases of the Higgs fields can be chosen such that the vacuum expectation values are real and positive. That is, the tree-level MSSM Higgs sector conserves CP, which implies that the neutral Higgs mass eigenstates possess definite CP quantum numbers.

The five physical Higgs particles consist of a charged Higgs pair

$$H^{\pm} = H_d^{\pm} \sin\beta + H_u^{\pm} \cos\beta\,,$$

one CP-odd scalar

$$A^0 = \sqrt{2}\left(\mathrm{Im}\,H_d^0 \sin\beta + \mathrm{Im}\,H_u^0 \cos\beta\right)\,,$$

and two CP-even scalars

$$h^0 = -(\sqrt{2}\,\mathrm{Re}\,H_d^0 - v_d)\sin\alpha + (\sqrt{2}\,\mathrm{Re}\,H_u^0 - v_u)\cos\alpha\,,$$
$$H^0 = (\sqrt{2}\,\mathrm{Re}\,H_d^0 - v_d)\cos\alpha + (\sqrt{2}\,\mathrm{Re}\,H_u^0 - v_u)\sin\alpha\,,$$

where we have now labeled the Higgs fields according to their electric charge. The angle α arises when the CP-even Higgs squared-mass matrix (in the H_d^0—H_u^0 basis) is diagonalized to obtain the physical CP-even Higgs states.

All Higgs masses and couplings can be expressed in terms of two parameters usually chosen to be m_A and $\tan\beta$.

The charged Higgs mass is given by

$$m_{H^\pm}^2 = m_A^2 + m_W^2 \,,$$

and the CP-even Higgs bosons h^0 and H^0 are eigenstates of the squared-mass matrix

$$\mathcal{M}_0^2 = \begin{pmatrix} m_A^2 \sin^2\beta + m_Z^2 \cos^2\beta & -(m_A^2 + m_Z^2)\sin\beta\cos\beta \\ -(m_A^2 + m_Z^2)\sin\beta\cos\beta & m_A^2 \cos^2\beta + m_Z^2 \sin^2\beta \end{pmatrix} .$$

The eigenvalues of $\mathcal{M}_0^2$ are the squared-masses of the two CP-even Higgs scalars

$$m_{H,h}^2 = \tfrac{1}{2}\left(m_A^2 + m_Z^2 \pm \sqrt{(m_A^2 + m_Z^2)^2 - 4m_Z^2 m_A^2 \cos^2 2\beta} \right) ,$$

and α is the angle that diagonalizes the CP-even Higgs squared-mass matrix. It follows that

$$\boxed{m_h \leq m_Z|\cos 2\beta| \leq m_Z \,.}$$

Note the contrast with the SM where the Higgs mass is a free parameter, $m_h^2 = \frac{1}{2}\lambda v^2$.

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North-Holland, Amsterdam

HIGGS BOSONS IN SUPERSYMMETRIC MODELS (I)*

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Received 3 June 1985
(Revised 6 August 1985)

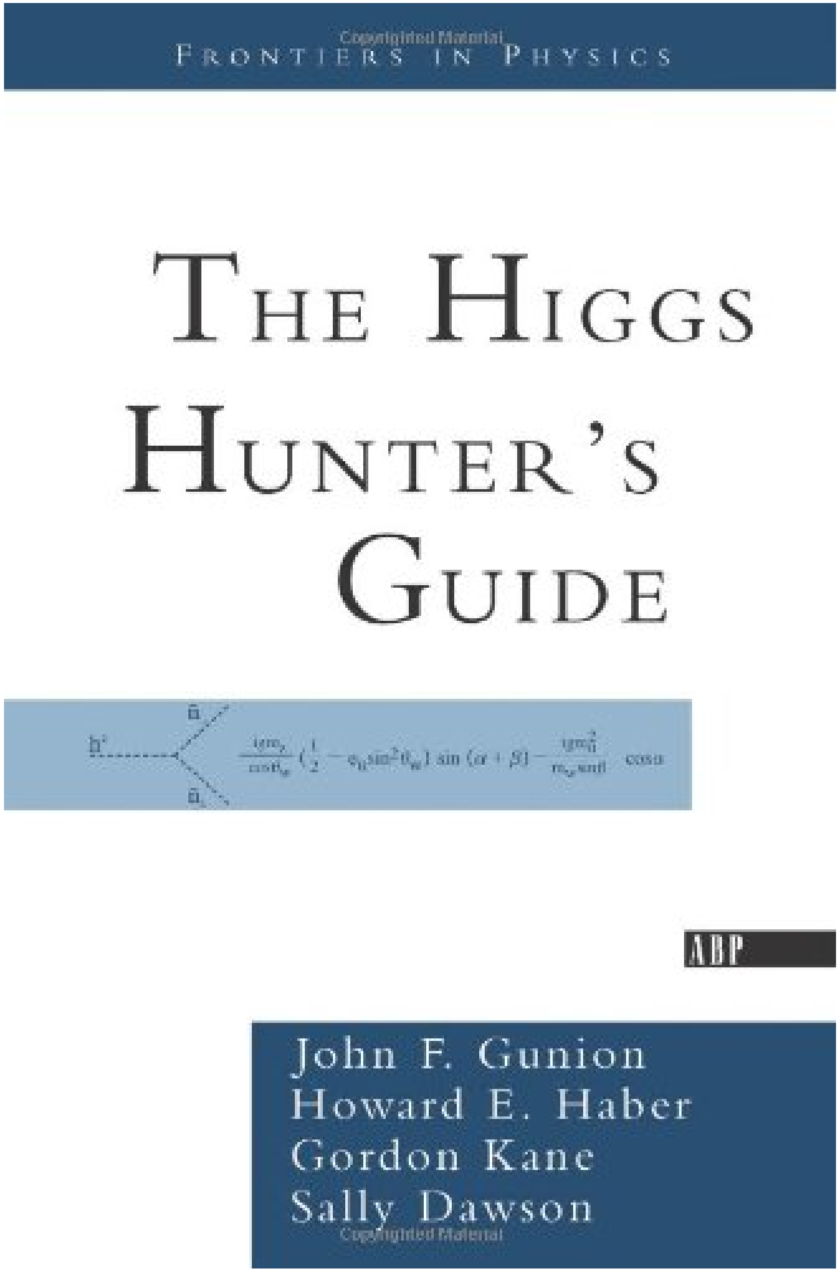
We describe the properties of Higgs bosons in a class of supersymmetric theories. We consider models in which the low-energy sector contains two weak complex doublets and perhaps one complex gauge-singlet Higgs field. Supersymmetry is assumed to be either softly or spontaneously broken, thereby imposing a number of restrictions on the Higgs boson parameters. We elucidate the Higgs boson masses and present Feynman rules for their couplings to the gauge bosons, fermions and scalars of the theory. We also present Feynman rules for vertices which are related by supersymmetry to the above couplings. Exact analytic expressions are given in two useful limits – one corresponding to the absence of the gauge-singlet Higgs field and the other corresponding to the absence of a supersymmetric Higgs mass term.

By 1990, there was a huge literature on the phenomenology of the Higgs boson of the Standard Model and of Higgs bosons of extended Higgs sectors.

The LEP collider was ready to extend the Higgs search to masses of order the Z mass and beyond.

On the horizon was the SSC (soon to be cancelled) and the LHC.

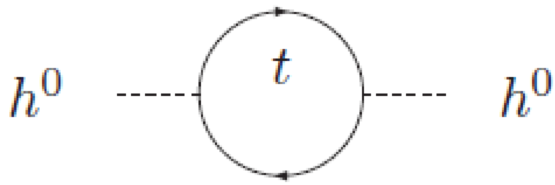
The hunt for the Higgs boson was on!

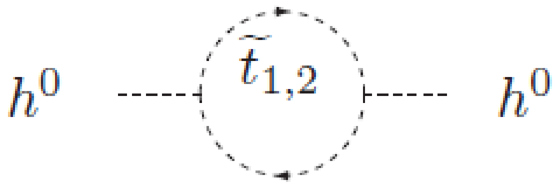


An aside: saving the MSSM Higgs sector

The absence of a Higgs boson discovery at LEP appeared to be in conflict with the upper bound on the mass of the h^0 ($m_h \leq m_Z|\cos 2\beta| \leq m_Z$) quoted earlier. But, radiative corrections save the day.

The Higgs mass can be shifted due to loops of particles and their superpartners (an incomplete cancelation, which would have been exact if supersymmetry were unbroken):





$$m_h^2 \lesssim m_Z^2 + \frac{3g^2 m_t^4}{8\pi^2 m_W^2}\left[\ln\left(\frac{M_S^2}{m_t^2}\right) + \frac{X_t^2}{M_S^2}\left(1 - \frac{X_t^2}{12M_S^2}\right)\right],$$

where $X_t \equiv A_t - \mu\cot\beta$ governs stop mixing and M_S^2 is the average squared-mass of the top-squarks $\widetilde{t}_1$ and $\widetilde{t}_2$ (which are the mass-eigenstate combinations of the interaction eigenstates, $\widetilde{t}_L$ and $\widetilde{t}_R$).

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Can the Mass of the Lightest Higgs Boson of the Minimal Supersymmetric Model be Larger than m_Z?

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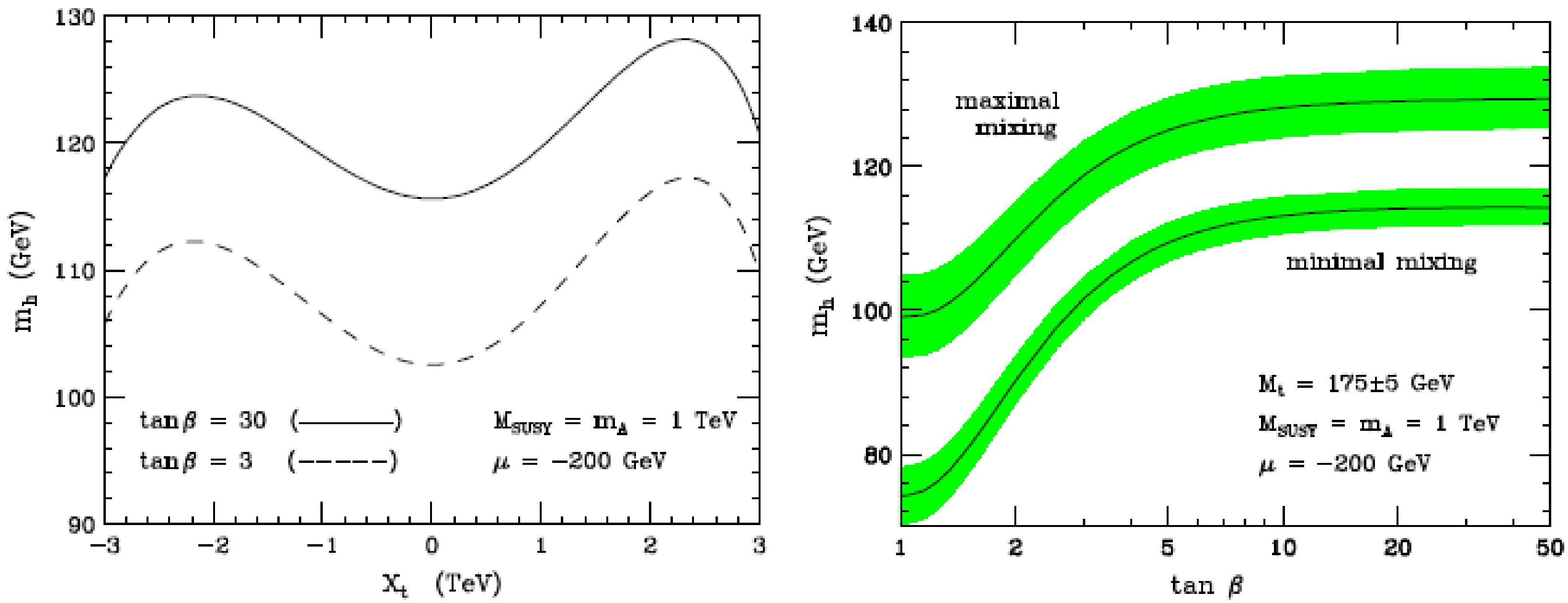
(Received 3 January 1991)

In the minimal supersymmetric model (MSSM), the *tree-level* mass of the lightest Higgs scalar h^0 cannot be larger than the mass of the Z boson. We have computed the one-loop radiative correction to the upper bound on m_{h^0} as a function of the free parameters of the MSSM. We find that the dominant correction to $m_{h^0} - m_Z$ is large and positive and grows like m_t^4, where m_t is the top-quark mass. As a result, the MSSM cannot be ruled out if the CERN e^+e^- collider LEP-200 fails to discover the Higgs boson.

Papers by our group and by two other groups (J. Ellis, G. Ridolfi and F. Zwirner and Y. Okada, M. Yamaguchi and T. Yanagida) published in 1991 demonstrated that radiative corrections could yield a significant enhancement to the Higgs mass upper bound in the MSSM, pushing the mass bound above the eventual LEP limit.

The state-of-the-art computation includes the full one-loop result, all the significant two-loop contributions, some of the leading three-loop terms, and renormalization-group improvements. The final conclusion is that $m_h \lesssim 130$ GeV [assuming that the top-squark mass is no heavier than about 2 TeV].



Maximal mixing corresponds to choosing the MSSM Higgs parameters in such a way that m_h is maximized (for a fixed $\tan\beta$). This occurs for $X_t/M_S \sim 2$. As $\tan\beta$ varies, m_h reaches is maximal value, $(m_h)_{\max} \simeq 130$ GeV, for $\tan\beta \gg 1$ and $m_A \gg m_Z$.

The fate of extended Higgs sectors

- Extended Higgs sectors provide new phenomenological signatures which, if observable, would reveal a richer structure for the electroweak symmetry breaking dynamics.
  - Charged Higgs bosons
  - CP-odd neutral Higgs bosons
  - Neutral Higgs bosons of indefinite CP (and new sources of CP violation)
- But, the mass scale that characterizes the scalar states beyond the Standard Model (SM) Higgs boson could be somewhat separated from the electroweak symmetry breaking (EWSB) scale.
- In such a scenario, integrating out physics above the new mass scale would yield an effective theory of a single Higgs scalar whose properties would be very close to those of the SM Higgs boson. The true nature of the EWSB dynamics would only be revealed by detecting the deviations from SM-like Higgs behavior.

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North-Holland

MULTI-SCALAR MODELS WITH A HIGH-ENERGY SCALE*

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Received 6 November 1989

We study multi-Higgs models under the assumption that new physics exists at some high-energy scale (Λ_{NP}). If we perform the minimally required fine-tuning in order to set the electroweak scale (Λ_{EW}), we find that the low-energy scalar spectrum is identical to that of the Standard Model with minimal Higgs content, up to corrections of order $\Lambda^2_{EW}/\Lambda^2_{NP}$. If, in

This was the introduction of the Decoupling limit scenario.

Example: decoupling of the non-minimal Higgs bosons of the MSSM Higgs sector (tree-level analysis)

In the limit of $m_A \gg m_Z$, the expressions for the Higgs masses and mixing angle simplify and one finds

$$\begin{aligned} m_h^2 &\simeq m_Z^2 \cos^2 2\beta\,, \\ m_H^2 &\simeq m_A^2 + m_Z^2 \sin^2 2\beta\,, \\ m_{H^\pm}^2 &= m_A^2 + m_W^2\,, \\ \cos^2(\beta-\alpha) &\simeq \frac{m_Z^4 \sin^2 4\beta}{4m_A^4}\,. \end{aligned}$$

Two consequences are immediately apparent. First, $m_A \simeq m_H \simeq m_{H^\pm}$, up to corrections of $\mathcal{O}(m_Z^2/m_A)$. Second, $\cos(\beta-\alpha) = 0$ up to corrections of $\mathcal{O}(m_Z^2/m_A^2)$. This is the decoupling limit, since at energy scales below approximately common mass of the heavy Higgs bosons $H^\pm$ H^0, A^0, the effective Higgs theory is precisely that of the SM.

A survey of tree-level Higgs couplings in the MSSM

1. Higgs couplings to gauge boson pairs ($V = W$ or Z)

$$g_{h^0VV} = g_V m_V \sin(\beta - \alpha)\,, \qquad\qquad g_{H^0VV} = g_V m_V \cos(\beta - \alpha)\,,$$

where $g_V \equiv 2m_V/v$. There are no tree-level couplings of A^0 or $H^\pm$ to VV.

2. Higgs couplings to a single gauge boson

The couplings of V to two neutral Higgs bosons (which must have opposite CP-quantum numbers) is denoted by $g_{\phi A^0 Z}(p_\phi - p_A^0)$, where $\phi = h^0$ or H^0 and the momenta p_ϕ and p_A^0 point into the vertex, and

$$g_{h^0A^0Z} = \frac{g\cos(\beta - \alpha)}{2\cos\theta_W}\,, \qquad\qquad g_{H^0A^0Z} = \frac{-g\sin(\beta - \alpha)}{2\cos\theta_W}\,.$$

3. Summary of Higgs boson–vector boson couplings

The properties of the three-point and four-point Higgs boson-vector boson couplings are conveniently summarized by listing the couplings that are proportional to either $\sin(\beta-\alpha)$ or $\cos(\beta-\alpha)$ or are angle-independent. As a reminder, $\cos(\beta-\alpha)\to 0$ in the decoupling limit.

$\cos(\beta-\alpha)$	$\sin(\beta-\alpha)$	angle-independent
$H^0W^+W^-$	$h^0W^+W^-$	—
H^0ZZ	h^0ZZ	—
ZA^0h^0	ZA^0H^0	ZH^+H^-, γH^+H^-
$W^\pm H^\mp h^0$	$W^\pm H^\mp H^0$	$W^\pm H^\mp A^0$
$ZW^\pm H^\mp h^0$	$ZW^\pm H^\mp H^0$	$ZW^\pm H^\mp A^0$
$\gamma W^\pm H^\mp h^0$	$\gamma W^\pm H^\mp H^0$	$\gamma W^\pm H^\mp A^0$
—	—	$VV\phi\phi$, VVA^0A^0, VVH^+H^-

where $\phi=h^0$ or H^0 and $VV=W^+W^-$, ZZ, $Z\gamma$ or $\gamma\gamma$.

4. Higgs-fermion couplings

Supersymmetry imposes a Type-II structure for the Higgs-fermion Yukawa couplings. Since the neutral Higgs couplings to fermions are flavor-diagonal, we list only the Higgs coupling to 3rd generation fermions. The couplings of the neutral Higgs bosons to $f\bar{f}$ relative to the Standard Model value, $gm_f/2m_W$, are given by

$$
\begin{aligned}
h^0 b\bar{b} \quad (\text{or } h^0\tau^+\tau^-): &\qquad -\frac{\sin\alpha}{\cos\beta} = \sin(\beta-\alpha) - \tan\beta\cos(\beta-\alpha)\,, \\
h^0 t\bar{t}: &\qquad \frac{\cos\alpha}{\sin\beta} = \sin(\beta-\alpha) + \cot\beta\cos(\beta-\alpha)\,, \\
H^0 b\bar{b} \quad (\text{or } H^0\tau^+\tau^-): &\qquad \frac{\cos\alpha}{\cos\beta} = \cos(\beta-\alpha) + \tan\beta\sin(\beta-\alpha)\,, \\
H^0 t\bar{t}: &\qquad \frac{\sin\alpha}{\sin\beta} = \cos(\beta-\alpha) - \cot\beta\sin(\beta-\alpha)\,, \\
A^0 b\bar{b} \quad (\text{or } A^0\tau^+\tau^-): &\qquad \gamma_5\tan\beta\,, \\
A^0 t\bar{t}: &\qquad \gamma_5\cot\beta\,,
\end{aligned}
$$

where the γ_5 indicates a pseudoscalar coupling. Note that the $h^0 f\bar{f}$ couplings approach their SM values in the decoupling limit, where $\cos(\beta-\alpha)\to 0$.

Similarly, the charged Higgs boson couplings to fermion pairs, with all particles pointing into the vertex, are given by**

$$g_{H^-t\bar{b}} = \frac{g}{\sqrt{2}m_W}\left[m_t \cot\beta\, P_R + m_b \tan\beta\, P_L\right],$$
$$g_{H^-\tau^+\nu} = \frac{g}{\sqrt{2}m_W}\left[m_\tau \tan\beta\, P_L\right].$$

Especially noteworthy is the possible $\tan\beta$-enhancement of certain Higgs-fermion couplings. The general expectation in MSSM models is that $\tan\beta$ lies in a range:

$$1 \lesssim \tan\beta \lesssim \frac{m_t}{m_b}.$$

Near the upper limit of $\tan\beta$, we have roughly identical values for the top and bottom Yukawa couplings, $h_t \sim h_b$, since

$$h_b = \frac{\sqrt{2}\,m_b}{v_d} = \frac{\sqrt{2}\,m_b}{v\cos\beta}, \qquad h_t = \frac{\sqrt{2}\,m_t}{v_u} = \frac{\sqrt{2}\,m_t}{v\sin\beta}.$$

**Including the full flavor structure, the CKM matrix appears in the charged Higgs couplings in the standard way for a charged-current interaction.

In general, in the limit of $\cos(\beta-\alpha)\to 0$, all the h^0 couplings to SM particles approach their SM limits. In particular, if λ_V is a Higgs coupling to vector bosons and λ_f is a Higgs couplings to fermions, then

$$\frac{\lambda_V}{[\lambda_V]_{\rm SM}} = \sin(\beta-\alpha) = 1 + \mathcal{O}\left(m_Z^4/m_A^4\right) \,, \qquad\qquad \frac{\lambda_f}{[\lambda_f]_{\rm SM}} = 1 + \mathcal{O}\left(m_Z^2/m_A^2\right) \,.$$

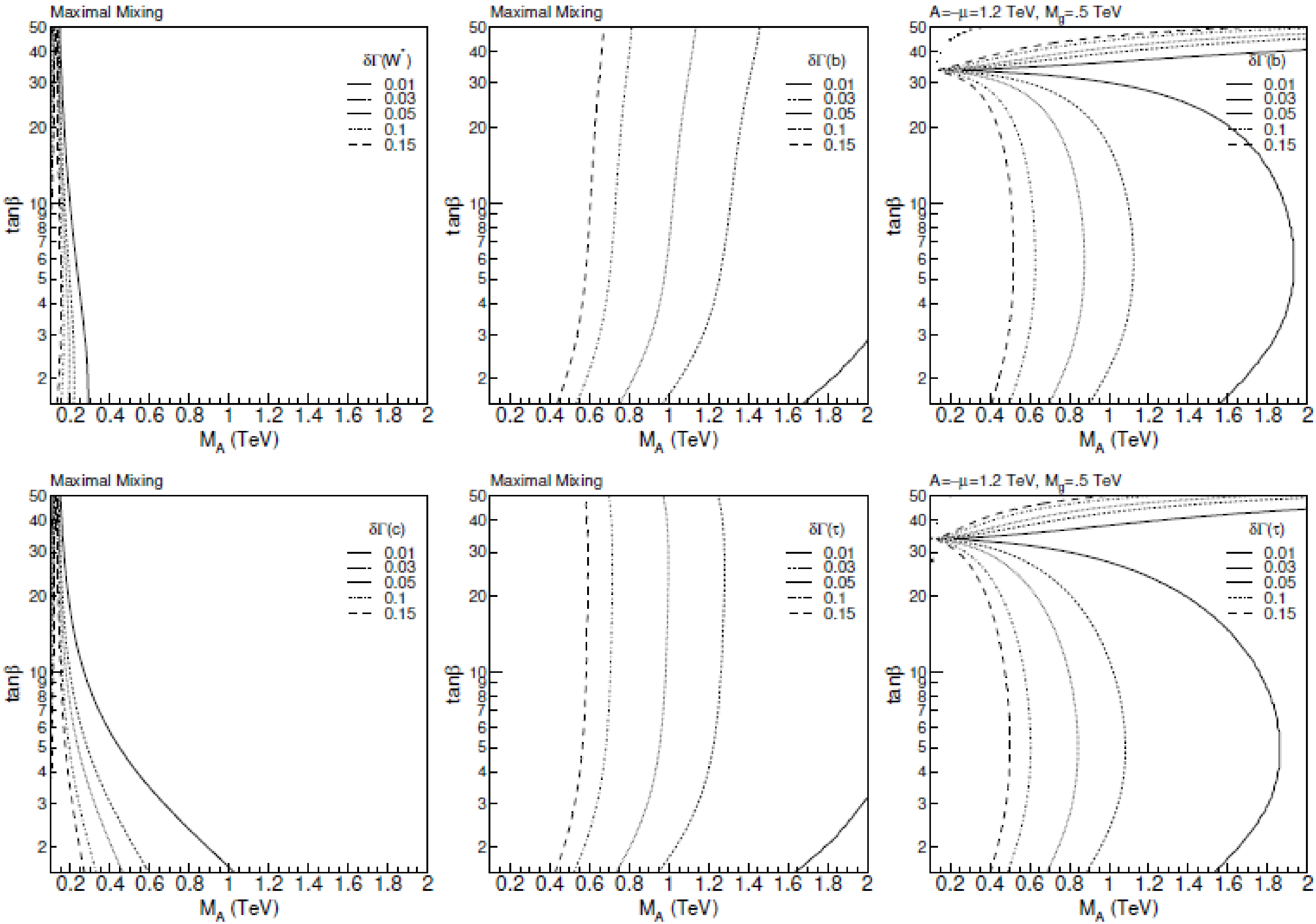
The behavior of the $h^0 ff$ coupling is repeated below:

$$h^0 b\bar{b} \quad (\text{or } h^0\tau^+\tau^-): \qquad -\frac{\sin\alpha}{\cos\beta} = \sin(\beta-\alpha) - \tan\beta\cos(\beta-\alpha)\,,$$

$$h^0 t\bar{t}: \qquad \frac{\cos\alpha}{\sin\beta} = \sin(\beta-\alpha) + \cot\beta\cos(\beta-\alpha)\,.$$

Note the extra $\tan\beta$ enhancement in the deviation of λ_{hbb} from $[\lambda_{hbb}]_{\rm SM}$.

Thus, the approach to decoupling is fastest for the h^0VV couplings, and slowest for the couplings of h^0 to down-type quarks and leptons (if $\tan\beta$ is large).



Deviations of Higgs partial widths from their SM values in two different MSSM scenarios (Carena, Haber, Logan and Mrenna).

The extension of the decoupling limit to the most general 2HDM is straightforward. For the CP-conserving 2HDM, the decoupling limit is controlled by $\cos(\beta - \alpha)$, and the first-order corrections due to the non-minimal Higgs sector to various observables are correlated.

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CP-conserving two-Higgs-doublet model: The approach to the decoupling limit

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A *CP*-even neutral Higgs boson with standard-model-like couplings may be the lightest scalar of a two-Higgs-doublet model. We study the decoupling limit of the most general *CP*-conserving two-Higgs-doublet model, where the mass of the lightest Higgs scalar is significantly smaller than the masses of the other Higgs bosons of the model. In this case, the properties of the lightest Higgs boson are nearly indistinguishable from those of the standard model Higgs boson. The first nontrivial corrections to Higgs boson couplings in the approach to the decoupling limit are also evaluated. The importance of detecting such deviations in precision Higgs boson measurements at future colliders is emphasized. We also clarify the case in which a neutral Higgs

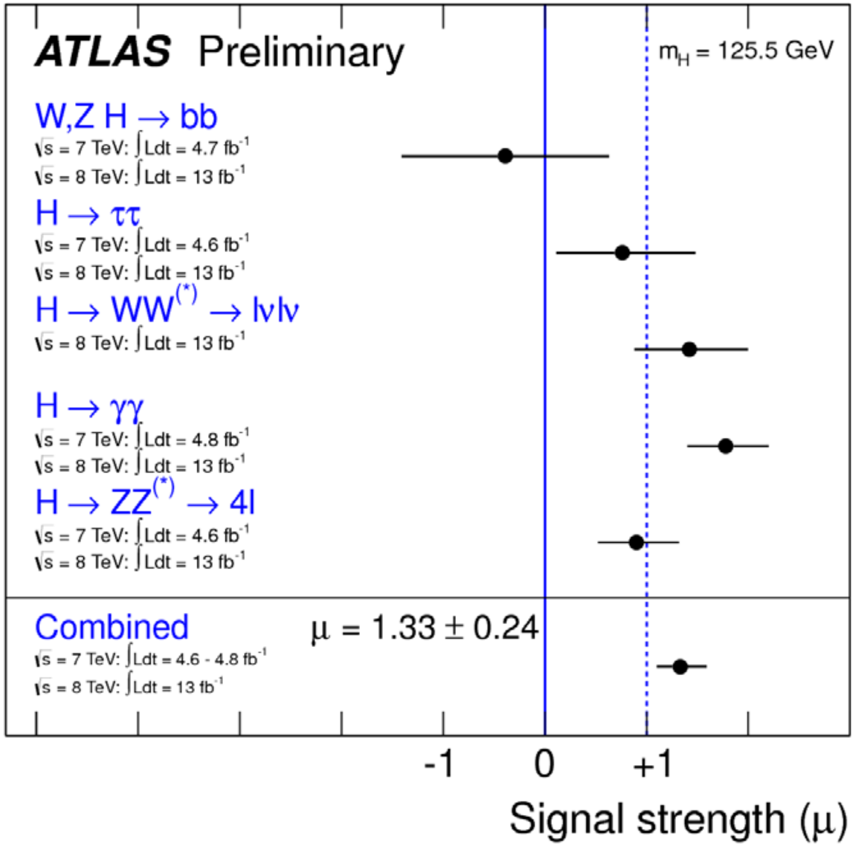
Interpreting the LHC Higgs data and the decoupling limit

- It is important to distinguish two energy scales:
  - Λ_H : the scale of the heavy non-minimal Higgs bosons.
  - Λ_{NP} : the scale of new physics beyond the Higgs-extended SM.

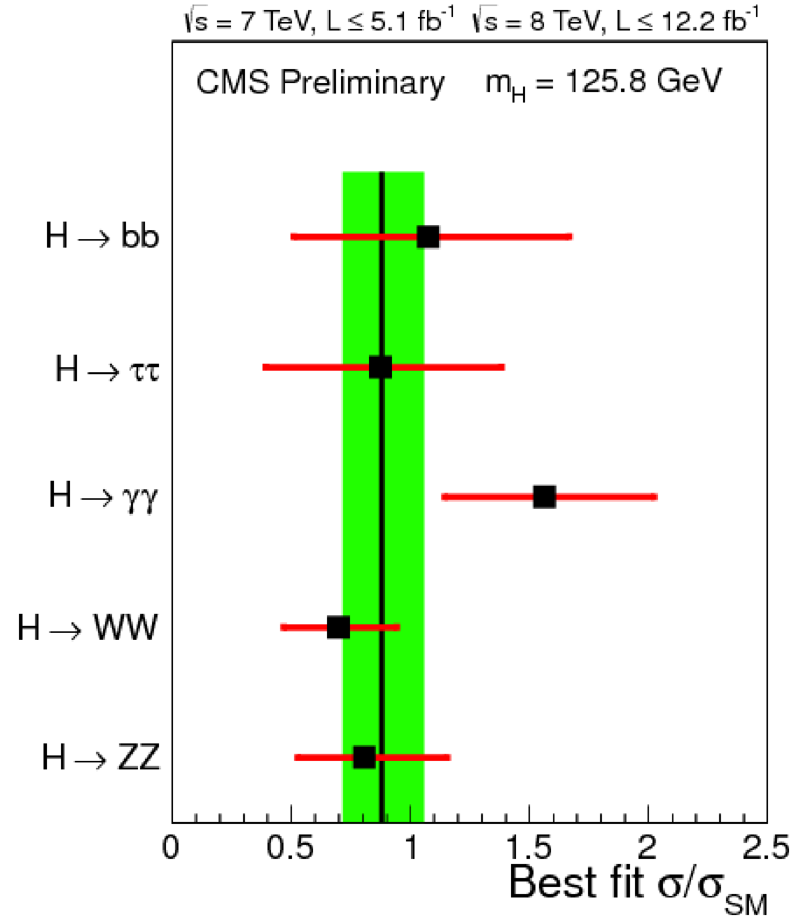
- The departure from the decoupling limit can receive contributions from both the heavy Higgs states via tree-level mixing and from new physics via one-loop radiative correction effects.

  - Separating out these two effects if deviations from SM Higgs couplings are confirmed will be important (and challenging).

What is the LHC Higgs data telling us?



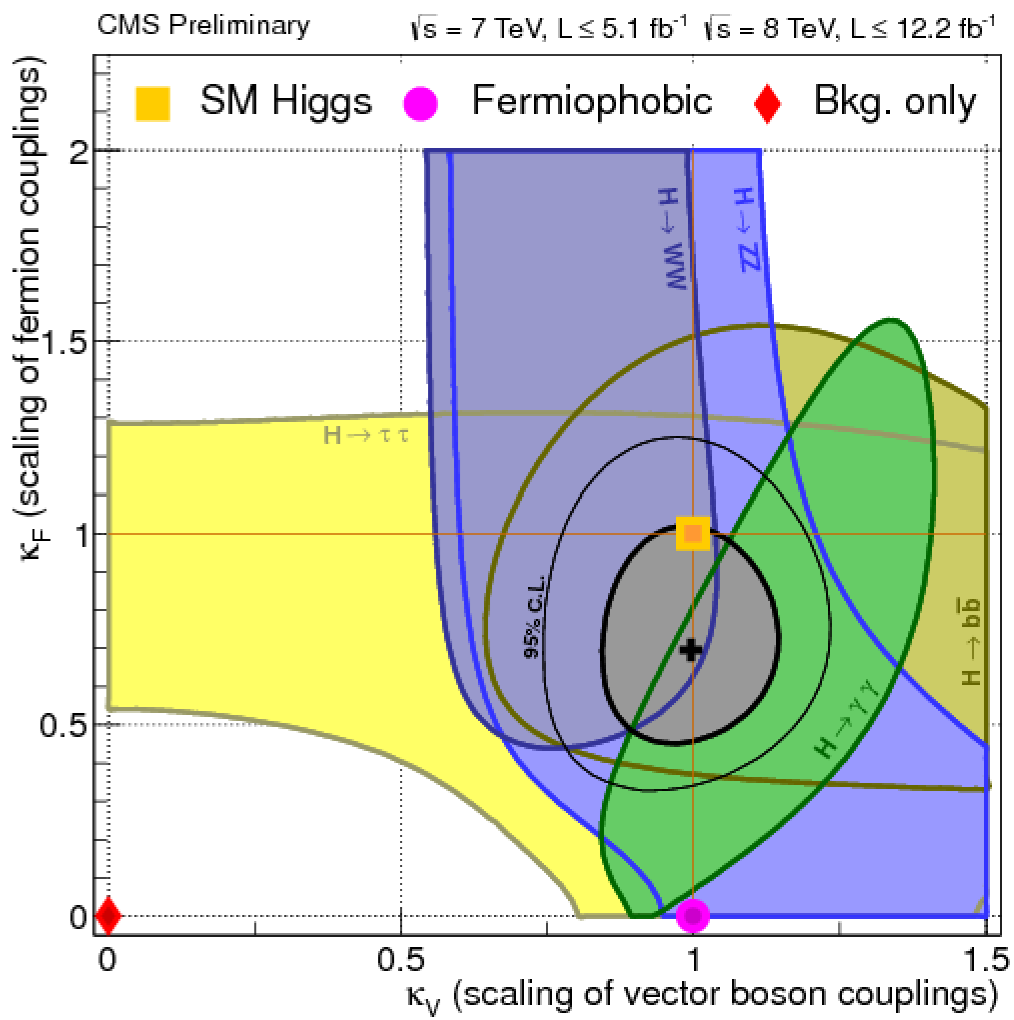
Summary of the individual and combined best-fit values of the strength parameter for a Higgs boson mass hypothesis of 125.5 GeV. Taken from ATLAS-CONF-2012-170, 13 December 2012.



Values of $\hat{\mu} = \sigma/\sigma_{SM}$ for the combination (solid vertical line) and for sub-combinations grouped by decay mode (points). The vertical band shows the overall $\hat{\mu}$ value 0.88 ± 0.21. The horizontal bars indicate the ±1σ uncertainties (both statistical and systematic) on the $\hat{\mu}$ values for individual channels. Taken from CMS-PAS-HIG-12-045, 16 November 2012.

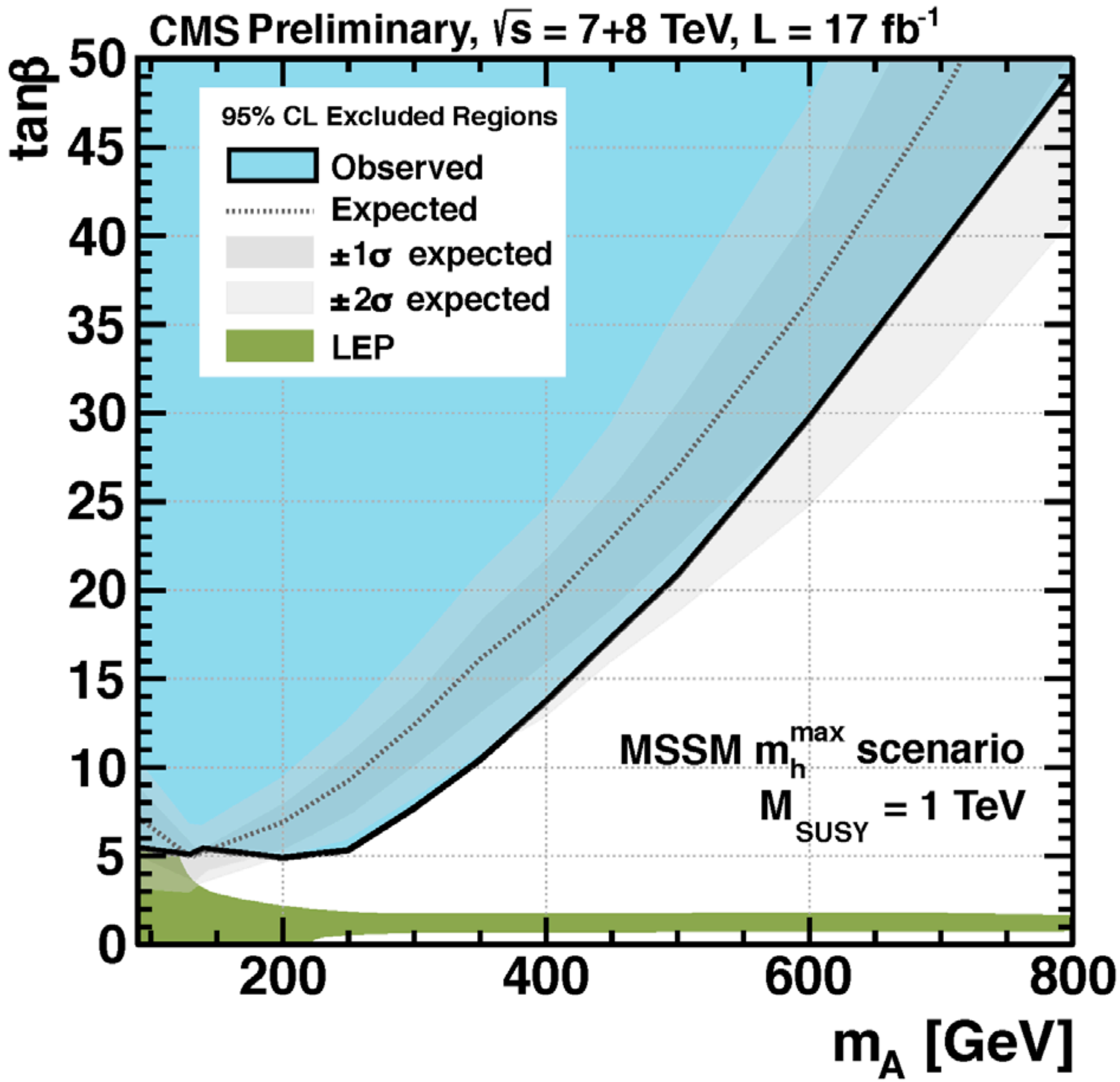
The γγ excess has received the most attention. ATLAS currently quotes a 2.4 σ discrepancy from the SM prediction.

Nevertheless, global fits suggest no statistically significant deviation from SM predictions.



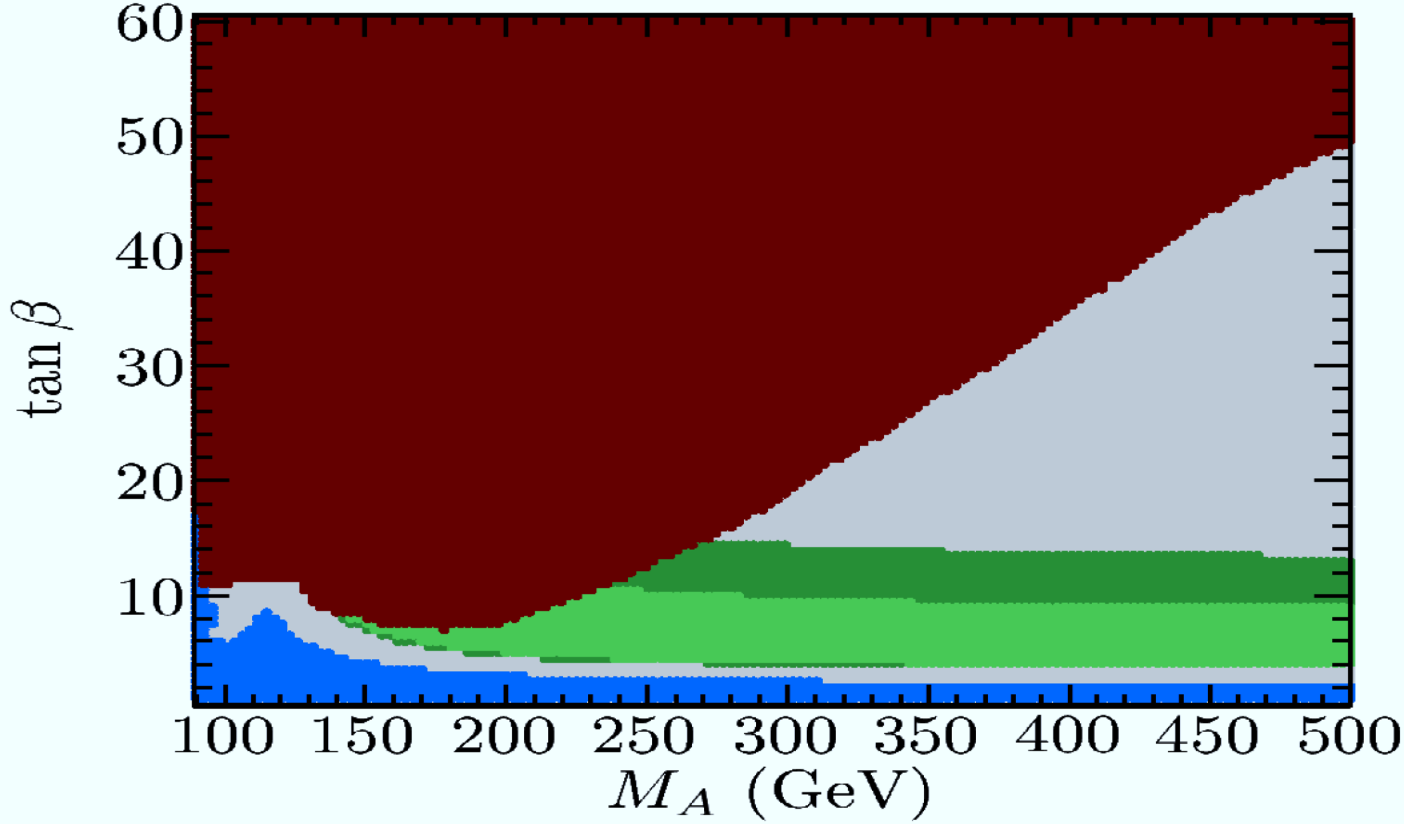
2D test statistics q(κ_V, κ_F) scan, including individual channels, assuming κ_V and κ_F have the same sign. CMS finds a slightly better fit when the relative sign of κ_V and κ_F is negative. Taken from CMS-PAS-HIG-12-045, 16 November 2012.

More evidence for the decoupling limit?



Lower bounds on M_A *and* $\tan\beta$ *from interpreting signal at* ~ 126 GeV *as light MSSM Higgs boson* h

Red: LHC limits from $H, A \to \tau^+\tau^-$ search; Blue: LEP limits
Green: compatible with interpreting signal at 126 GeV as light MSSM Higgs h ($+\ m_t$ variation) [*S. Heinemeyer, O. Stål, G. W. '11, '12*]



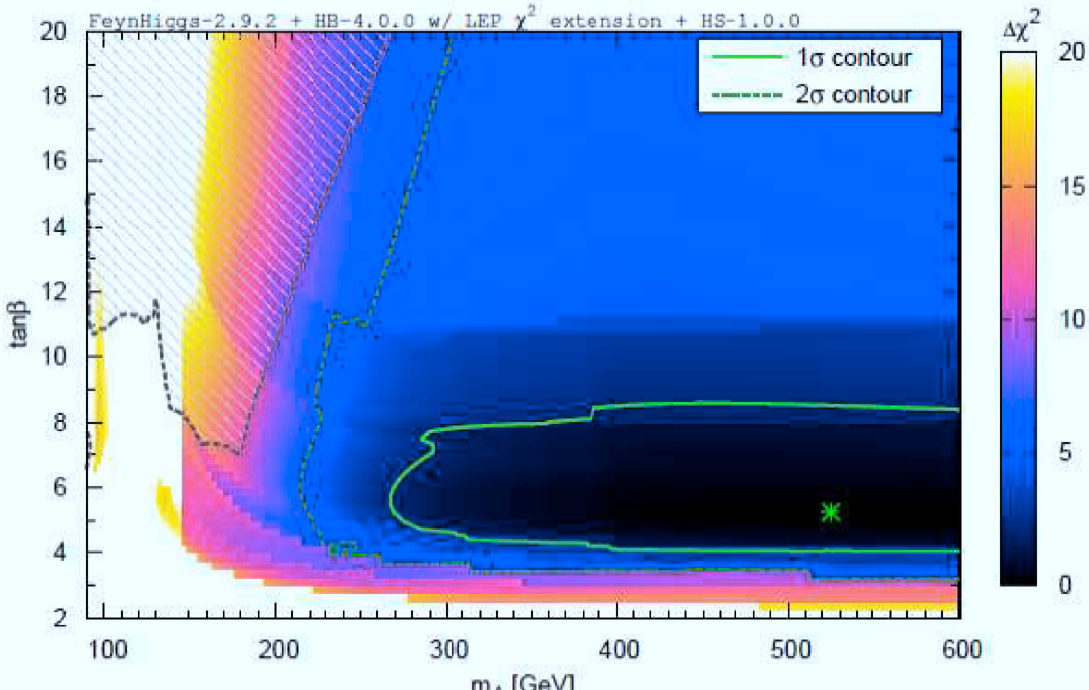
$\Rightarrow \tan\beta \gtrsim 4,\ M_A \gtrsim 140$ GeV, $M_{H^\pm} \gtrsim 160$ GeV

Analysis in $m_h^{\max}$ benchmark scenario

The $m_h^{\max}$ scenario

$(M_{SUSY} = 1$ TeV, $|X_t| = 2$ TeV, $\mu = 200$ GeV, $M_1 = 100$ GeV, $M_2 = 200$ GeV, $M_3 = 1200$ GeV$)$

5. Take into account $\Delta m_h^{th} = \Delta m_H^{th} = 2$ GeV in HiggsBounds.



- Exclusion in large m_A region vanishes (now, $m_h \lesssim 130$ GeV is allowed).

A cautionary tale: the wrong-Higgs couplings of the MSSM

Although radiative corrections to couplings tend to be at the few-percent level, there is some potential for significant effects:

- large radiative corrections due to a $\tan\beta$-enhancement (assuming $\tan\beta \gg 1$)
- CP-violating effects induced by complex SUSY-breaking parameters that enter in loops

In the MSSM, the tree-level Higgs–quark Yukawa Lagrangian is supersymmetry-conserving and is given by Type-II structure,

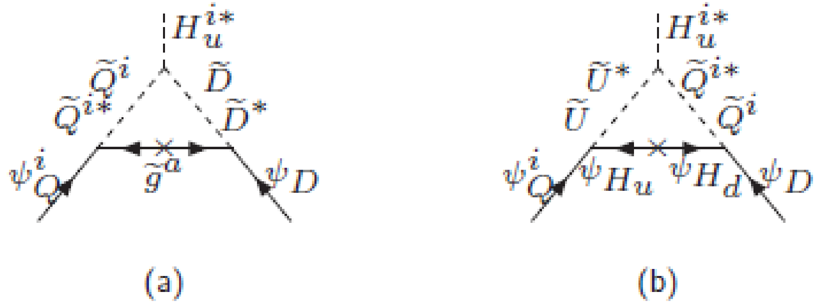
$$\mathcal{L}_{\rm yuk}^{\rm tree} = -\epsilon_{ij} h_b H_d^i \psi_Q^j \psi_D + \epsilon_{ij} h_t H_u^i \psi_Q^j \psi_U + {\rm h.c.}$$

Two other possible dimension-four gauge-invariant non-holomorphic Higgs-quark interactions terms, the so-called wrong-Higgs interactions,

$$H_u^{k*} \psi_D \psi_Q^k \qquad {\rm and} \qquad H_d^{k*} \psi_U \psi_Q^k \,,$$

are not supersymmetric (since the dimension-four supersymmetric Yukawa interactions must be holomorphic), and hence are absent from the tree-level Yukawa Lagrangian.

Nevertheless, the wrong-Higgs interactions can be generated in the effective low-energy theory below the scale of SUSY-breaking. In particular, one-loop radiative corrections, in which supersymmetric particles (squarks, higgsinos and gauginos) propagate inside the loop can generate the wrong-Higgs interactions. (See review by M. Carena and H.E. Haber Prog. Part. Nucl. Phys. **50**, 63 (2003) and references therein.)



One-loop diagrams contributing to the wrong-Higgs Yukawa effective operators. In (a), the cross ($\times$) corresponds to a factor of the gluino mass M_3. In (b), the cross corresponds to a factor of the higgsino Majorana mass parameter μ. Field labels correspond to annihilation of the corresponding particle at each vertex of the triangle.

If the superpartners are heavy, then one can derive an effective field theory description of the Higgs-quark Yukawa couplings below the scale of SUSY-breaking ($M_{\rm SUSY}$), where one has integrated out the heavy SUSY particles propagating in the loops.

The resulting effective Lagrangian is:

$$\mathcal{L}^{\text{eff}}_{\text{yuk}} = -\epsilon_{ij}(h_b + \delta h_b)\psi_b H^i_d \psi^j_Q + \Delta h_b \psi_b H^{k*}_u \psi^k_Q$$
$$+\epsilon_{ij}(h_t + \delta h_t)\psi_t H^i_u \psi^j_Q + \Delta h_t \psi_t H^{k*}_d \psi^k_Q + \text{h.c.}$$

In the limit of $M_{\text{SUSY}} \gg m_Z$,

$$\Delta h_b = h_b \left[\frac{2\alpha_s}{3\pi} \mu M_3 \mathcal{I}(M_{\tilde{b}_1}, M_{\tilde{b}_2}, M_g) + \frac{h_t^2}{16\pi^2} \mu A_t \mathcal{I}(M_{\tilde{t}_1}, M_{\tilde{t}_2}, \mu) \right] ,$$

where, M_3 is the Majorana gluino mass, μ is the supersymmetric Higgs-mass parameter, and $\widetilde{b}_{1,2}$ and $\widetilde{t}_{1,2}$ are the mass-eigenstate bottom squarks and top squarks, respectively. The loop integral is given by

$$\mathcal{I}(a,b,c) = \frac{a^2 b^2 \ln(a^2/b^2) + b^2 c^2 \ln(b^2/c^2) + c^2 a^2 \ln(c^2/a^2)}{(a^2 - b^2)(b^2 - c^2)(a^2 - c^2)} .$$

In the limit where at least one of the arguments of $\mathcal{I}(a,b,c)$ is large,

$$\mathcal{I}(a,b,c) \sim 1/\max(a^2, b^2, c^2) .$$

Thus, in the limit where $M_3 \sim \mu \sim A_t \sim M_{\tilde{b}} \sim M_{\tilde{t}} \sim M_{\text{SUSY}} \gg m_Z$, the one-loop contributions to Δh_b do *not* decouple.

Phenomenological consequences of the wrong-Higgs Yukawas

The effects of the wrong-Higgs couplings are $\tan\beta$-enhanced modifications of some physical observables. To see this, rewrite the Higgs fields in terms of the physical mass-eigenstates (and the Goldstone bosons):

$$H_d^1 = \frac{1}{\sqrt{2}}(v\cos\beta + H^0\cos\alpha - h^0\sin\alpha + iA^0\sin\beta - iG^0\cos\beta)\,,$$
$$H_u^2 = \frac{1}{\sqrt{2}}(v\sin\beta + H^0\sin\alpha + h^0\cos\alpha + iA^0\cos\beta + iG^0\sin\beta)\,,$$
$$H_d^2 = H^-\sin\beta - G^-\cos\beta\,,$$
$$H_u^1 = H^+\cos\beta + G^+\sin\beta\,,$$

with $v^2 \equiv v_u^2 + v_d^2 = (246~\mathrm{GeV})^2$ and $\tan\beta \equiv v_u/v_d$. For simplicity, we neglect below possible CP-violating effects due to complex couplings. Then, the b-quark mass is:

$$m_b = \frac{h_b v}{\sqrt{2}}\cos\beta\left(1 + \frac{\delta h_b}{h_b} + \frac{\Delta h_b \tan\beta}{h_b}\right) \equiv \frac{h_b v}{\sqrt{2}}\cos\beta(1+\Delta_b)\,,$$

which defines the quantity Δ_b.

In the limit of large $\tan\beta$ the term proportional to δh_b can be neglected, in which case,

$$\Delta_b \simeq (\Delta h_b/h_b)\tan\beta\,.$$

Thus, Δ_b is $\tan\beta$–enhanced if $\tan\beta \gg 1$. As previously noted, Δ_b survives in the limit of large $M_{\rm SUSY}$; this effect does not decouple.

From the effective Yukawa Lagrangian, we can obtain the couplings of the physical Higgs bosons to third generation fermions. Neglecting possible CP-violating effects,

$$\mathcal{L}_{\rm int} = -\sum_{q=t,b,\tau}\left[g_{h^0 q\bar{q}}h^0 q\bar{q} + g_{H^0 q\bar{q}}H^0 q\bar{q} - i g_{A^0 q\bar{q}}A^0 \bar{q}\gamma_5 q\right] + \left[\bar{b}g_{H^- t\bar{b}}tH^- + {\rm h.c.}\right]\,.$$

The one-loop corrections can generate measurable shifts in the decay rate for $h^0 \to b\bar{b}$:

$$g_{h^0 b\bar{b}} = -\frac{m_b}{v}\frac{\sin\alpha}{\cos\beta}\left[1 + \frac{1}{1+\Delta_b}\left(\frac{\delta h_b}{h_b} - \Delta_b\right)(1+\cot\alpha\cot\beta)\right]\,.$$

At large $\tan\beta \sim 20$—50, Δ_b can be as large as 0.5 in magnitude and of either sign, leading to a significant enhancement or suppression of the Higgs decay rate to $b\bar{b}$.

Non-decoupling effects in $h^0 \to b\bar{b}$: a closer look

The origin of the non-decoupling effects can be understood by noting that below the scale $M_{\rm SUSY}$, the effective low-energy Higgs theory is a completely general 2HDM. Thus, it is not surprising that the wrong-Higgs couplings do not decouple in the limit of $M_{\rm SUSY} \to \infty$.

However, suppose that $m_A \sim \mathcal{O}(M_{\rm SUSY})$. Then, the low-energy effective Higgs theory is a one-Higgs doublet model, and thus $g_{h^0 b\bar{b}}$ must approach its SM value. Indeed in this limit,

$$\cos(\beta-\alpha) = \frac{m_Z^2 \sin 4\beta}{2m_A^2} + \mathcal{O}\left(\frac{m_Z^4}{m_A^4}\right) ,$$

$$1 + \cot\alpha\cot\beta = -\frac{2m_Z^2}{m_A^2}\cos 2\beta + \mathcal{O}\left(\frac{m_Z^4}{m_A^4}\right) .$$

Thus the non-decoupling SUSY radiative corrections previously obtained now decouple as expected. [Haber, Herrero, Logan, Penaranda, Rigolin and Temes]

In the 2HDM, symmetries are imposed to restrict the structure of the Higgs-fermion interactions in order to avoid tree-level Higgs-mediated FCNCs. Two choices for this symmetry are:

- **Supersymmetry**
- **A subgroup of the U(2) Higgs-flavor symmetry. Most commonly, a $\mathbb{Z}_2$ discrete symmetry under which one of the two Higgs fields changes sign.**

However, these symmetries are typically broken at some scale, $\Lambda_{\textbf{sym}}$. If $\Lambda_{\textbf{sym}}$ lies somewhat above the EWSB scale but below the mass scale of the non-minimal Higgs bosons, then the effective theory below $\Lambda_{\textbf{sym}}$ is a completely general 2HDM.

Thus, I was motivated to study the most general 2HDM, where the two Higgs doublet fields are indistinguishable.

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Basis-independent methods for the two-Higgs-doublet model. II. The significance of $\tan\beta$

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In the most general two-Higgs-doublet model (2HDM), there is no distinction between the two complex hypercharge-one $SU(2)_L$ doublet scalar fields, Φ_a ($a = 1, 2$). Thus, any two orthonormal linear combinations of these two fields can serve as a basis for the Lagrangian. All physical observables of the model must be basis-independent. For example, $\tan\beta \equiv \langle\Phi_2^0\rangle/\langle\Phi_1^0\rangle$ is basis-dependent and thus cannot be a physical parameter of the model. In this paper, we provide a basis-independent treatment of the Higgs sector with particular attention to the neutral Higgs boson mass-eigenstates, which generically are *not* eigenstates of CP. We then demonstrate that all physical Higgs couplings are indeed independent of $\tan\beta$. In specialized versions of the 2HDM, $\tan\beta$ can be promoted to a physical parameter of the Higgs-fermion interactions. In the most general 2HDM, the Higgs-fermion couplings can be expressed in terms of a number of physical "$\tan\beta$-like" parameters that are manifestly basis-independent. The minimal supersymmetric extension of the standard model provides a simple framework for exhibiting such effects.

Since the two Higgs fields are indistinguishable, one can redefine the basis of the two Higgs fields by applying a general U(2) transformation. True observables must be basis-independent. In particular, the parameter tan β is no longer a physical quantity.

Start with the 2HDM fields, Φ_1 and Φ_2, in a generic basis, where $\langle\Phi_i\rangle = v_i$, and $v^2 \equiv |v_1|^2 + |v_2|^2 = (246~\text{GeV})^2$. 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 an overall rephasing, $H_2 \to e^{i\chi}H_2$. In the Higgs basis, the scalar potential is given by:

$$\begin{aligned}\mathcal{V} = {} & Y_1 H_1^\dagger H_1 + Y_2 H_2^\dagger H_2 + [Y_3 H_1^\dagger H_2 + \text{h.c.}] + \tfrac{1}{2}Z_1(H_1^\dagger H_1)^2 \\ & + \tfrac{1}{2}Z_2(H_2^\dagger H_2)^2 + Z_3(H_1^\dagger H_1)(H_2^\dagger H_2) + Z_4(H_1^\dagger H_2)(H_2^\dagger H_1) \\ & + \left\{\tfrac{1}{2}Z_5(H_1^\dagger H_2)^2 + \left[Z_6(H_1^\dagger H_1) + Z_7(H_2^\dagger H_2)\right]H_1^\dagger H_2 + \text{h.c.}\right\},\end{aligned}$$

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 complex and transform under the rephasing of H_2,

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

The Higgs mass-eigenstate basis

The physical charged Higgs boson is the charged component of the Higgs-basis doublet H_2, and its mass is given by $m^2_{H^\pm} = Y_2 + \frac{1}{2} Z_3 v^2$.

The three physical neutral Higgs boson mass-eigenstates are determined by diagonalizing a 3×3 real symmetric squared-mass matrix that is defined in the Higgs basis.†

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

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}. Under the rephasing $H_2 \to e^{i\chi} H_2$,

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

It is convenient to define the $q_{k\ell}$ which are defined in terms of the invariant angles θ_{12} and θ_{13}, where $c_{ij} \equiv \cos\theta_{ij}$ and $s_{ij} \equiv \sin\theta_{ij}$.

k	q_{k1}	q_{k2}
0	i	0
1	$c_{12}c_{13}$	$-s_{12} - ic_{12}s_{13}$
2	$s_{12}c_{13}$	$c_{12} - is_{12}s_{13}$
3	s_{13}	ic_{13}

The neutral Goldstone boson (h_0) and the physical neutral Higgs states ($h_{1,2,3}$) are given by:

$$h_k = \frac{1}{\sqrt{2}}\left\{q_{k1}^*\left(H_1^0 - \frac{v}{\sqrt{2}}\right) + q_{k2}^* H_2^0 e^{i\theta_{23}} + \text{h.c.}\right\} .$$

If we also define the physical charged Higgs state by $H^\pm = e^{\pm i\theta_{23}} H_2^\pm$, then all the mass eigenstate fields are invariant under the rephasing $H_2 \to e^{i\chi} H_2$.

The gauge boson–Higgs boson interactions

$$\mathscr{L}_{VVH} = \left(gm_W W_\mu^+ W^{\mu -} + \frac{g}{2c_W} m_Z Z_\mu Z^\mu\right) \mathrm{Re}(q_{k1}) h_k + em_W A^\mu (W_\mu^+ G^- + W_\mu^- G^+)$$

$$- gm_Z s_W^2 Z^\mu (W_\mu^+ G^- + W_\mu^- G^+)\,,$$

$$\mathscr{L}_{VVHH} = \left[\tfrac{1}{4} g^2 W_\mu^+ W^{\mu -} + \frac{g^2}{8c_W^2} Z_\mu Z^\mu\right] \mathrm{Re}(q_{j1}^* q_{k1} + q_{j2}^* q_{k2})\, h_j h_k$$

$$+ \left[\tfrac{1}{2} g^2 W_\mu^+ W^{\mu -} + e^2 A_\mu A^\mu + \frac{g^2}{c_W^2}\left(\tfrac{1}{2} - s_W^2\right)^2 Z_\mu Z^\mu + \frac{2ge}{c_W}\left(\tfrac{1}{2} - s_W^2\right) A_\mu Z^\mu\right] (G^+ G^- + H^+ H^-)$$

$$+ \left\{ \left(\tfrac{1}{2} eg A^\mu W_\mu^+ - \frac{g^2 s_W^2}{2c_W} Z^\mu W_\mu^+\right) (q_{k1} G^- + q_{k2} H^-) h_k + \mathrm{h.c.} \right\},$$

$$\mathscr{L}_{VHH} = \frac{g}{4c_W} \mathrm{Im}(q_{j1} q_{k1}^* + q_{j2} q_{k2}^*) Z^\mu h_j \overleftrightarrow{\partial}_\mu h_k - \tfrac{1}{2} g \left\{ i W_\mu^+ \left[q_{k1} G^- \overleftrightarrow{\partial}^\mu h_k + q_{k2} H^- \overleftrightarrow{\partial}^\mu h_k \right] + \mathrm{h.c.} \right\}$$

$$+ \left[ie A^\mu + \frac{ig}{c_W}\left(\tfrac{1}{2} - s_W^2\right) Z^\mu \right] (G^+ \overleftrightarrow{\partial}_\mu G^- + H^+ \overleftrightarrow{\partial}_\mu H^-)\,,$$

where $s_W \equiv \sin\theta_W$ and $c_W \equiv \cos\theta_W$.

The cubic and quartic Higgs couplings

$$\mathscr{L}_{3h} = -\tfrac{1}{2}v\, h_j h_k h_\ell \Big[q_{j1}q^*_{k1}\mathrm{Re}(q_{\ell 1})Z_1 + q_{j2}q^*_{k2}\,\mathrm{Re}(q_{\ell 1})(Z_3+Z_4) + \mathrm{Re}(q^*_{j1}q_{k2}q_{\ell 2}Z_5\, e^{-2i\theta_{23}})$$

$$+\mathrm{Re}\Big([2q_{j1}+q^*_{j1}]q^*_{k1}q_{\ell 2}Z_6\, e^{-i\theta_{23}}\Big) + \mathrm{Re}(q^*_{j2}q_{k2}q_{\ell 2}Z_7\, e^{-i\theta_{23}})\Big]$$

$$-v\, h_k G^+G^-\Big[\mathrm{Re}(q_{k1})Z_1 + \mathrm{Re}(q_{k2}\, e^{-i\theta_{23}}Z_6)\Big] + v\, h_k H^+H^-\Big[\mathrm{Re}(q_{k1})Z_3 + \mathrm{Re}(q_{k2}\, e^{-i\theta_{23}}Z_7)\Big]$$

$$-\tfrac{1}{2}v\, h_k\Big\{G^-H^+\Big[q^*_{k2}Z_4 + q_{k2}\, e^{-2i\theta_{23}}Z_5 + 2\mathrm{Re}(q_{k1})Z_6\, e^{-i\theta_{23}}\Big] + \mathrm{h.c.}\Big\}\,,$$

$$\mathscr{L}_{4h} = -\tfrac{1}{8}h_j h_k h_l h_m\Big[q_{j1}q_{k1}q^*_{\ell 1}q^*_{m1}Z_1 + q_{j2}q_{k2}q^*_{\ell 2}q^*_{m2}Z_2 + 2q_{j1}q^*_{k1}q_{\ell 2}q^*_{m2}(Z_3+Z_4)$$

$$+2\mathrm{Re}(q^*_{j1}q^*_{k1}q_{\ell 2}q_{m2}Z_5\, e^{-2i\theta_{23}}) + 4\mathrm{Re}(q_{j1}q^*_{k1}q^*_{\ell 1}q_{m2}Z_6\, e^{-i\theta_{23}}) + 4\mathrm{Re}(q^*_{j1}q_{k2}q_{\ell 2}q^*_{m2}Z_7\, e^{-i\theta_{23}})\Big]$$

$$-\tfrac{1}{2}h_j h_k G^+G^-\Big[q_{j1}q^*_{k1}Z_1 + q_{j2}q^*_{k2}Z_3 + 2\mathrm{Re}(q_{j1}q_{k2}Z_6\, e^{-i\theta_{23}})\Big]$$

$$-\tfrac{1}{2}h_j h_k H^+H^-\Big[q_{j2}q^*_{k2}Z_2 + q_{j1}q^*_{k1}Z_3 + 2\mathrm{Re}(q_{j1}q_{k2}Z_7\, e^{-i\theta_{23}})\Big]$$

$$-\tfrac{1}{2}h_j h_k\Big\{G^-H^+\Big[q_{j1}q^*_{k2}Z_4 + q^*_{j1}q_{k2}Z_5\, e^{-2i\theta_{23}} + q_{j1}q^*_{k1}Z_6\, e^{-i\theta_{23}} + q_{j2}q^*_{k2}Z_7\, e^{-i\theta_{23}}\Big] + \mathrm{h.c.}\Big\}$$

$$-\tfrac{1}{2}Z_1G^+G^-G^+G^- - \tfrac{1}{2}Z_2H^+H^-H^+H^- - (Z_3+Z_4)G^+G^-H^+H^-$$

$$-\tfrac{1}{2}(Z_5\, e^{-2i\theta_{23}}H^+H^+G^-G^- + \mathrm{h.c.}) - G^+G^-(Z_6\, e^{-i\theta_{23}}H^+G^- + \mathrm{h.c.}) - H^+H^-(Z_7\, e^{-i\theta_{23}}H^+G^- + \mathrm{h.c.})\,.$$

Higgs-fermion Yukawa couplings in the 2HDM

In the Higgs basis, $\kappa^{U,D}$ and $\rho^{U,D}$, are the 3×3 Yukawa coupling matrices,

$$-\mathscr{L}_{\rm Y} = \overline{U}_L(\kappa^U H_1^{0\,\dagger} + \rho^U H_2^{0\,\dagger})U_R - \overline{D}_L K^\dagger(\kappa^U H_1^- + \rho^U H_2^-)U_R$$
$$+\overline{U}_L K(\kappa^{D\,\dagger}H_1^+ + \rho^{D\,\dagger}H_2^+)D_R + \overline{D}_L(\kappa^{D\,\dagger}H_1^0 + \rho^{D\,\dagger}H_2^0)D_R + {\rm h.c.},$$

where $U = (u, c, t)$ and $D = (d, s, b)$ are the physical quark fields and K is the CKM mixing matrix. (Repeat for the leptons.)

By setting $H_1^0 = v/\sqrt{2}$ and $H_2^0 = 0$, one obtains the quark mass terms. Hence, κ^U and κ^D are proportional to the diagonal quark mass matrices M_U and M_D, respectively,

$$M_U = \frac{v}{\sqrt{2}}\kappa^U = {\rm diag}(m_u\,,\, m_c\,,\, m_t)\,, \qquad M_D = \frac{v}{\sqrt{2}}\kappa^{D\,\dagger} = {\rm diag}(m_d\,,\, m_s\,,\, m_b)\,.$$

Note that $\rho^Q \to e^{-i\chi}\rho^Q$ under the rephasing $H_2 \to e^{i\chi}H_2$, (for $Q = U, D$).

It follows that the Yukawa couplings of the mass-eigenstate Higgs bosons and the Goldstone bosons to the quarks are:

$$-\mathscr{L}_Y = \frac{1}{v}\overline{D}\sum_k \Big\{ M_D(q_{k1}P_R + q_{k1}^* P_L) + \frac{v}{\sqrt{2}}\Big[q_{k2}\,[e^{i\theta_{23}}\rho^D]^\dagger P_R + q_{k2}^*\,e^{i\theta_{23}}\rho^D P_L\Big]\Big\} D h_k$$

$$+\frac{1}{v}\overline{U}\sum_k \Big\{ M_U(q_{k1}P_L + q_{k1}^* P_R) + \frac{v}{\sqrt{2}}\Big[q_{k2}^*\,e^{i\theta_{23}}\rho^U P_R + q_{k2}\,[e^{i\theta_{23}}\rho^U]^\dagger P_L\Big]\Big\} U h_k$$

$$+\Big\{\overline{U}\Big[K[e^{i\theta_{23}}\rho^D]^\dagger P_R - [e^{i\theta_{23}}\rho^U]^\dagger K P_L\Big] D H^+ + \frac{\sqrt{2}}{v}\overline{U}\,[K M_D P_R - M_U K P_L]\, D G^+ + \text{h.c.}\Big\},$$

where $P_{L,R} = \frac{1}{2}(1 \mp \gamma_5)$ are left and right-handed projection operators.

- The combinations $e^{i\theta_{23}}\rho^U$ and $e^{i\theta_{23}}\rho^U$ that appear in the interactions above are invariant under the rephasing of H_2.
- Note that no $\tan\beta$ parameter appears above! This is because $\tan\beta$ is an unphysical parameter in the general 2HDM.
- If ρ^U and ρ^D are complex non-diagonal 3×3 matrices, then the 2HDM exhibits (tree-level) flavor changing neutral currents (FCNCs) mediated by neutral Higgs exchange and new sources of CP-violation.

How to avoid tree-level Higgs-mediated FCNCs

- Arbitrarily declare ρ^U and ρ^D to be diagonal matrices. This is an unnaturally fine-tuned solution.

- Impose a discrete symmetry or supersymmetry (e.g."Type-I" or "Type-II" Higgs-fermion interactions), which selects out a special basis of the 2HDM scalar fields. In this case, ρ^Q is automatically proportional to M_Q (for $Q = U, D, L$), and is hence diagonal.

- Impose alignment without a symmetry: $\rho^Q = \alpha^Q \kappa^Q \,, (Q = U, D, L) \,,$ where the α^Q are complex scalar parameters [e.g. see Pich and Tuzon (2009)].

- Impose the decoupling limit. Tree-level Higgs-mediated FCNCs will be suppressed by factors of squared-masses of heavy Higgs states. (How heavy is sufficient?)

Is $\tan\beta$ a physical observable?

In a generic 2HDM, $\tan\beta$ is meaningless, unless there is some additional symmetry, which picks out a special basis for the scalar fields. But, if you do not know the symmetry *a priori*, how should you proceed?

- You can test for specific models, such as a Type-I or II model.
- You can measure observables that are basis-independent, and determine whether additional symmetries are present.

Example: For simplicity, ignore the first two generations of quarks and assume CP conservation. Then, you should measure ρ^D and ρ^U.

$$\begin{aligned} \rho^D/\rho^U &= m_b/m_t\,, && \text{for Type-I couplings}\,, \\ \rho^D\rho^U &= -2m_t m_b/v^2\,, && \text{for Type-II couplings}\,. \end{aligned}$$

Thus, measuring ρ^D and ρ^U would help determine the underlying structure of the Higgs-fermion interaction.

The Decoupling Limit of the 2HDM revisited

In the decoupling limit, one of the two Higgs doublets of the 2HDM receives a very large mass which then decouples from the theory. This is achieved when $Y_2 \gg v^2$ and $|Z_i| \lesssim \mathcal{O}(1)$ [for all i]. The effective low energy theory is a one-Higgs-doublet model, which yields the SM Higgs boson.

We order the neutral scalar masses according to $m_1 < m_{2,3}$ and define the Higgs mixing angles accordingly. The conditions for the decoupling limit are:

$$|\sin\theta_{12}| \lesssim \mathcal{O}\left(\frac{v^2}{m_2^2}\right) \ll 1\,, \qquad |\sin\theta_{13}| \lesssim \mathcal{O}\left(\frac{v^2}{m_3^2}\right) \ll 1\,,$$

$$\text{Im}(Z_5\, e^{-2i\theta_{23}}) \lesssim \mathcal{O}\left(\frac{v^2}{m_3^2}\right) \ll 1\,.$$

In the decoupling limit, $m_1 \ll m_2, m_3, m_{H^\pm}$. In particular, the properties of h_1 coincide with the SM Higgs boson with $m_1^2 = Z_1 v^2$ up to corrections of $\mathcal{O}(v^4/m_{2,3}^2)$, and $m_2 \simeq m_3 \simeq m_{H^\pm}$ with squared mass splittings of $\mathcal{O}(v^2)$.

In the exact decoupling limit, where $s_{12} = s_{13} = \text{Im}(Z_5\, e^{-2i\theta_{23}}) = 0$, the interactions of h_1 are precisely those of the SM Higgs boson. In particular, the interactions of the h_1 in the decoupling limit are CP-conserving and diagonal in quark flavor space.

In the most general 2HDM, CP-violating and neutral Higgs-mediated FCNCs are suppressed by factors of $\mathcal{O}(v^2/m_{2,3}^2)$ in the decoupling limit. In contrast, the interactions of the heavy neutral Higgs bosons (h_2 and h_3) and the charged Higgs bosons ($H^\pm$) in the decoupling limit can exhibit both CP-violating and quark flavor non-diagonal couplings (proportional to the ρ^Q).

The decoupling limit is a generic feature of extended Higgs sectors.

- The observation of a SM-like Higgs boson does not rule out the possibility of an extended Higgs sector in the decoupling regime.
- Deviations from SM Higgs behavior would provide clues to the structure of the extended Higgs sector and/or the structure of new physics beyond the SM.

We are at the dawn of an exciting era for Higgs hunters.

- We eagerly await the updates and clarifications from ATLAS and CMS on the full 2012 Higgs data set.
- Will the newly discovered Higgs boson signal the presence of new physics beyond the Standard Model?
- Meanwhile, I would like to congratulate Peter Higgs for receiving the Companion of Honour.
- I would also like to express my gratitude to Peter for his instrumental role in creating an exciting and stimulating career path for me and my fellow Higgs hunters.