Supersymmetric Heavy Higgses at the LHC

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LPTHE Paris, 7 Juin 2013

Contents

The hierarchy problem :

When calculating the radiative corrections to the SM Higgs boson mass, one encounters quadratic divergences in the cut–off scale Λ at which the theory stops to be valid and New Physics should appear :

$$\Delta^F M_H^2 \quad = \quad N_f \frac{\lambda_f^2}{8\pi^2}\left[-\Lambda^2 + 6m_f^2 \log\frac{\Lambda}{m_f} - 2m_f^2\right] + \mathcal{O}(1/\Lambda^2) \quad ; \quad \Delta M_H^2 \propto \Lambda^2$$

If $\Lambda \sim M_{\rm GUT} \sim 10^{16}$ GeV, the Higgs boson mass would be huge ($\gg$ 126 GeV).

For the SM Higgs boson to stay relatively light ($M_H \lesssim 1$ TeV from unitarity and perturbativity) we need to add a counterterm to M_H^2 and adjust it with a precision of $\mathcal{O}(10^{-30})$, which seems highly unnatural $\Rightarrow$ the naturalness or fine–tuning problem.

A related question, why $\Lambda \gg M_Z$? : the hierarchy problem.

The problem can be seen as a lack of a symmetry which protects M_H against very high scales.

A new scalar particles would contribute to M_H as :

$$\Delta^S M_H^2 = \frac{\lambda_S N_S}{16\pi^2}\left[-\Lambda^2 + 2m_S^2 \log\left(\frac{\Lambda}{m_S}\right)\right] - \frac{\lambda_S^2 N_S}{16\pi^2} v^2\left[-1 + 2\log\left(\frac{\Lambda}{m_S}\right)\right] + \mathcal{O}\left(\frac{1}{\Lambda^2}\right)$$

If $\lambda_f^2 = 2m_f^2/v^2 = -\lambda_S$ and $N_S = 2N_f$ then

$$\Delta^{F+2S} M_H^2 = \frac{\lambda_f^2 N_f}{4\pi^2}\left[(m_f^2 - m_S^2)\log\left(\frac{\Lambda}{m_S}\right) + 3m_f^2 \log\left(\frac{m_S}{m_f}\right)\right] + \mathcal{O}\left(\frac{1}{\Lambda^2}\right)$$

If there are scalar particles $\Rightarrow$ no quadratic divergence : the hierarchy and naturalness problems solved.

If $m_f = m_s$: no divergences at all $\Rightarrow$ M_H is protected by this 'supersymmetry'.

If this symmetry is broken, $m_s \gg m_f$ $\Rightarrow$ the hierarchy and naturalness problems would be reintroduced again in the theory...

Basics of Supersymmetry :

The SUSY generators : $\mathcal{Q}|\mathrm{Fermion}\rangle = |\mathrm{Boson}\rangle \quad \mathcal{Q}|\mathrm{Boson}\rangle = |\mathrm{Fermion}\rangle$,
The particles are combined into superfields :

$$\mathcal{L}_{\mathrm{kin}} = \sum_i \left\{ (D_\mu S_i^*)(D^\mu S_i) + i\overline{\psi}_i D_\mu \gamma^\mu \psi_i \right\} + \sum_a \left\{ -\frac{1}{4} F^a_{\mu\nu} F^{\mu\nu a} + \frac{i}{2} \overline{\lambda}_a \sigma^\mu D_\mu \lambda_a \right\}$$

The interactions among the fields are specified by SUSY and gauge invariance :

$$\mathcal{L}_{\mathrm{int.\ scal-fer.-gauginos}} = -\sqrt{2} \sum_{i,a} g_a \left[S_i^* T^a \overline{\psi}_{iL} \lambda_a + \mathrm{h.c.} \right]$$

$$\mathcal{L}_{\mathrm{int.\ quartic\,scal.}} = -\frac{1}{2} \sum_a \left(\sum_i g_a S_i^* T^a S_i \right)^2$$

Interactions are given in terms of the gauge coupling constants.

$\Rightarrow$ when SUSY is exact, everything is completely specified and there is no adjustable parameter.

The only freedom : the choice of the superpotential W which gives the form of the scalar potential and the Yukawa interactions between fermion and scalar fields.

The interaction Lagrangian :

$$\mathcal{L}_W = -\sum_i \left| \frac{\partial W}{\partial z_i} \right|^2 - \frac{1}{2} \sum_{ij} \left[\overline{\psi}_{iL} \frac{\partial^2 W}{\partial z_i \partial z_j} \psi_j + \mathrm{h.c.} \right]$$

$V_{\mathrm{tree}}^{SUSY} = V_F + V_D$:

- $V_F = \sum_i |W^i|^2$ with $W^i = \partial W / \partial S_i$
- $V_D = \frac{1}{2} \sum_{a=1}^{3} \left(\sum_i g_a S_i^* T^a S_i \right)^2$ with $\mathrm{U(1)_Y, SU(2)_L, SU(3)_C}$

No fundamental scalar particles have the same mass as the known fermions
$\Rightarrow$ SUSY must be broken.

One possibility : introduce by hand terms that break SUSY explicitly and parametrize our ignorance of the fundamental SUSY–breaking mechanism (low energy effective SUSY theory).
$\Rightarrow$ the Minimal Supersymmetric Standard Model (MSSM).

The Minimal Supersymmetric Standard Model :

The unconstrained MSSM, defined by 4 assumptions :

(a) Minimal gauge group: The MSSM is based on the group $\mathrm{SU(3)_C \times SU(2)_L \times U(1)_Y}$, i.e. the SM gauge symmetry.

(b) Minimal particle content:

quarks and leptons + squarks and sleptons: $\hat{Q}, \hat{U}_R, \hat{D}_R, \hat{L}, \hat{E}_R$. (3 gen.)
2 chiral superfields $\hat{H}_1, \hat{H}_2$.

(c) Minimal Yukawa interactions and R–parity conservation: a discrete symmetry called R–parity is imposed (enforce lepton and baryon number conservation)

$$R_p = (-1)^{2s+3B+L}; \quad \mathrm{R}_p = +1 \text{ SM part.}, \; R_p = -1 \text{ SUSY part.}$$

The most general superpotential, compatible with gauge invariance, renormalizability and R–parity conservation :

$$W = \sum_{i,j=gen} -Y^u_{ij}\, \hat{u}_{Ri} \hat{H}_2 \cdot \hat{Q}_j + Y^d_{ij}\, \hat{d}_{Ri} \hat{H}_1 \cdot \hat{Q}_j + Y^\ell_{ij}\, \hat{\ell}_{Ri} \hat{H}_1 \cdot \hat{L}_j + \mu \hat{H}_2 \cdot \hat{H}_1$$

(d) Minimal set of soft SUSY–breaking terms: To break SUSY while preventing the reappearance of the quadratic divergences (soft SUSY–breaking).

One adds to the Lagrangian terms which explicitly break SUSY :

- Mass terms for the gluinos, winos and binos:

$$-\mathcal{L}_{\rm gaugino} = \frac{1}{2}\left[M_1 \tilde{B}\tilde{B} + M_2 \sum_{a=1}^{3} \tilde{W}^a \tilde{W}_a + M_3 \sum_{a=1}^{8} \tilde{G}^a \tilde{G}_a \ + \ \text{h.c.}\right]$$

- Mass terms for the scalar fermions:

$$-\mathcal{L}_{\rm sfermions} = \sum_{i=gen} m^2_{\tilde{Q}_i} \tilde{Q}_i^\dagger \tilde{Q}_i + m^2_{\tilde{L}_i} \tilde{L}_i^\dagger \tilde{L}_i + m^2_{\tilde{u}_i} |\tilde{u}_{R_i}|^2 + m^2_{\tilde{d}_i} |\tilde{d}_{R_i}|^2 + m^2_{\tilde{\ell}_i} |\tilde{\ell}_{R_i}|^2$$

- Mass and bilinear terms for the Higgs bosons:

$$-\mathcal{L}_{\rm Higgs} = m^2_{H_2} H_2^\dagger H_2 + m^2_{H_1} H_1^\dagger H_1 + B\mu(H_2 \cdot H_1 + \text{h.c.})$$

- Trilinear couplings between sfermions and Higgs bosons

$$-\mathcal{L}_{\rm tril.} = \sum_{i,j=gen} \left[A^u_{ij} Y^u_{ij} \tilde{u}^*_{R_i} H_2 \cdot \tilde{Q}_j + A^d_{ij} Y^d_{ij} \tilde{d}^*_{R_i} H_1 \cdot \tilde{Q}_j + A^l_{ij} Y^\ell_{ij} \tilde{\ell}^*_{R_i} H_1 \cdot \tilde{L}_j \ + \ \text{h.c.}\right]$$

$$\Rightarrow \quad V_{\rm soft} = -\mathcal{L}_{\rm sfermions} - \mathcal{L}_{\rm Higgs} - \mathcal{L}_{\rm tril.}$$

The MSSM defined by the four hypotheses (a)–(d) is the unconstrained MSSM.

unconstrained MSSM : the soft SUSY–breaking terms introduce 105 parameters (+19 for SM) $\Rightarrow$ phenomenological analysis complicated.

A more viable MSSM with more assumptions :
(i) all the soft SUSY–breaking parameters are real (no new source of CP–violation);
(ii) the matrices for the sfermion masses and for the trilinear couplings are all diagonal (no FCNCs at the tree–level) ;
(iii) the soft SUSY–breaking masses and trilinear couplings of the first and second sfermion generations are the same at low energy.

$\Rightarrow$ 22 input parameters : the phenomenological MSSM (pMSSM)

- $\tan\beta$: the ratio of the vevs of the two–Higgs doublet fields;
- $m^2_{H_1}, m^2_{H_2}$: the Higgs mass parameters squared;
- M_1, M_2, M_3: the bino, wino and gluino mass parameters;
- $m_{\tilde{q}}, m_{\tilde{u}_R}, m_{\tilde{d}_R}, m_{\tilde{l}}, m_{\tilde{e}_R}$: the 1/2 gen. sfermion mass param;
- A_u, A_d, A_e: the first/second generation trilinear couplings;
- $m_{\tilde{Q}}, m_{\tilde{t}_R}, m_{\tilde{b}_R}, m_{\tilde{L}}, m_{\tilde{\tau}_R}$: 3^{rd} gen. sfermion mass param;
- A_t, A_b, A_τ: 3^{rd} gen. trilinear couplings.

In the MSSM to break the electroweak symmetry one need two doublets of complex scalar fields :

$$H_1 = \begin{pmatrix} H_1^0 \\ H_1^- \end{pmatrix} \text{ with } Y_{H_1} = -1 \;\; , \;\; H_2 = \begin{pmatrix} H_2^+ \\ H_2^0 \end{pmatrix} \text{ with } Y_{H_2} = +1$$

the terms contributing to the scalar Higgs potential V_H come from 3 different sources :

- The D terms containing the quartic Higgs interactions

$$\mathrm{U(1)_Y} \quad : \quad V_D^1 = \frac{1}{2}\left[\frac{g_1}{2}(|H_2|^2 - |H_1|^2)\right]^2$$

$$\mathrm{SU(2)_L} \quad : \quad V_D^2 = \frac{1}{2}\left[\frac{g_2}{2}(H_1^{i*}\tau_{ij}^a H_1^j + H_2^{i*}\tau_{ij}^a H_2^j)\right]^2$$

- The F term of the Superpotential which can be written as $V_F = \sum_i |\partial W(\phi_j)/\partial\phi_i|^2$. From the term $W \sim \mu \hat{H}_1 \cdot \hat{H}_2$, one obtains the component $V_F = \mu^2(|H_1|^2 + |H_2|^2)$
- Soft SUSY–breaking scalar Higgs mass terms and the bilinear term

$$V_{\rm soft} = m_{H_1}^2 H_1^\dagger H_1 + m_{H_2}^2 H_2^\dagger H_2 + B\mu(H_2 \cdot H_1 + \mathrm{h.c.})$$

The full scalar potential involving the Higgs fields :

$$V_H = (|\mu|^2 + m_{H_1}^2)|H_1|^2 + (|\mu|^2 + m_{H_2}^2)|H_2|^2 - \mu B \epsilon_{ij}(H_1^i H_2^j + \mathrm{h.c.}) + \frac{g_2^2 + g_1^2}{8}(|H_1|^2 - |H_2|^2)^2 + \frac{1}{2} g_2^2 |H_1^\dagger H_2|^2$$

Expanding the Higgs fields in terms of their charged and neutral components and defining the mass squared terms :

$$\overline{m}_1^2 = |\mu|^2 + m_{H_1}^2 \,, \quad \overline{m}_2^2 = |\mu|^2 + m_{H_2}^2 \,, \quad \overline{m}_3^2 = B\mu$$

one obtains :

$$V_H = \overline{m}_1^2(|H_1^0|^2 + |H_1^-|^2) + \overline{m}_2^2(|H_2^0|^2 + |H_2^+|^2) - \overline{m}_3^2(H_1^- H_2^+ - H_1^0 H_2^0 + \mathrm{h.c.}) + \frac{g_2^2 + g_1^2}{8}(|H_1^0|^2 + |H_1^-|^2 - |H_2^0|^2 - |H_2^+|^2)^2 + \frac{g_2^2}{2}|H_1^{-*} H_1^0 + H_2^{0*} H_2^+|^2$$

We require then that the minimum of the potential V_H breaks the $\mathrm{SU(2)_L \times U(1)_Y}$ group while preserving the electromagnetic symmetry $\mathrm{U(1)_Q}$.

Some important remarks :

• To have EWSB, one needs a combination of the H_1^0 and H_2^0 fields to have a negative squared mass term. This occurs if

$$\overline{m}_3^2 > \overline{m}_2^2 \overline{m}_2^2$$

if not, $\langle H_1^0 \rangle = \langle H_2^0 \rangle$ is a stable minimum of the potential.
• In the direction $|H_1^0|=|H_2^0|$, there is no quartic term. V_H is bounded from below for large values of the field H_i only if the following condition is satisfied:

$$\overline{m}_1^2 + \overline{m}_2^2 > 2|\overline{m}_3^2|$$

• To have explicit EWSB the potential at the minimum should have a saddle point :

$$\mathrm{Det}\left(\frac{\partial^2 V_H}{\partial H_i^0 \partial H_j^0}\right) < 0 \Rightarrow \overline{m}_1^2\,\overline{m}_2^2 < \overline{m}_3^4$$

• The 2 above conditions on the masses $\bar{m}_i$ are not satisfied if $\overline{m}_1^2 = \overline{m}_2^2$ and, thus, we must have non–vanishing soft SUSY–breaking scalar masses m_{H_1} and m_{H_2}

$$\overline{m}_1^2 \neq \overline{m}_2^2 \Rightarrow m_{H_1}^2 \neq m_{H_2}^2$$

- To break EW symmetry, we need also to break SUSY $\Rightarrow$ close connection between gauge symmetry breaking and SUSY–breaking.

- In constrained models (ex. mSUGRA), $m_{H_1} = m_{H_2}$, but the running to lower energies via the contributions of top/bottom quarks and their SUSY partners in the RGEs makes that this degeneracy is lifted at the weak scale.

- In the running one obtains $m^2_{H_2} < 0$ or $m^2_{H_2} \ll m^2_{H_1}$ which thus triggers EWSB: this is the radiative breaking of the symmetry.

- EWSB is more natural and elegant in the MSSM than in the SM since, in the latter case, we needed to make the ad hoc choice $\mu^2 < 0$ while in the MSSM this comes simply from radiative corrections.

Minimizing the Higgs potential :

neutral component of the Higgs fields acquire vacuum expectetion values :

$$\langle H_1^0 \rangle = \frac{v_1}{\sqrt{2}} \quad , \quad \langle H_2^0 \rangle = \frac{v_2}{\sqrt{2}}$$

where : $(v_1^2 + v_2)^2 = v^2 = \frac{4M_Z^2}{g_2^2+g_1^2} = (246 \text{ GeV})^2$
defining the ratio of the two vevs : $\tan\beta = \frac{v_2}{v_1} = \frac{(v \sin\beta)}{(v \cos\beta)}$

one obtains two minimization conditions :

$$B\mu = \frac{(m_{H_1}^2 - m_{H_2}^2)\tan 2\beta + M_Z^2 \sin 2\beta}{2}$$
$$\mu^2 = \frac{m_{H_2}^2 \sin^2\beta - m_{H_1}^2 \cos^2\beta}{\cos 2\beta} - \frac{M_Z^2}{2}$$

$(m_{H_1}, m_{H_2}, \tan\beta) \approx (B, \mu^2)$ (sign(μ) is undetermined).

To obtain the Higgs physical fields and their masses :

- develop the 2 doublet complex scalar fields H_1 and H_2 around the vacuum, into real and imaginary parts :

$$H_1 = (H_1^0, H_1^-) = \frac{1}{\sqrt{2}}\left(v_1 + H_1^0 + iP_1^0\ ,\ H_1^-\right)$$
$$H_2 = (H_2^+, H_2^0) = \frac{1}{\sqrt{2}}\left(H_2^+\ ,\ v_2 + H_2^0 + iP_2^0\right)$$

  The real parts correspond to the CP–even Higgs bosons.
  The imaginary parts corresponds to the CP–odd Higgs and the Goldstone bosons.

- diagonalize the mass matrices evaluated at the vacuum

$$\mathcal{M}_{ij}^2 = \frac{1}{2}\frac{\partial^2 V_H}{\partial H_i \partial H_j}\bigg|_{\langle H_1^0\rangle = v_1/\sqrt{2}, \langle H_2^0\rangle = v_2/\sqrt{2}, \langle H_{1,2}^{\pm}\rangle = 0}$$

For the neutral Goldstone and CP–odd Higgs bosons :

$$\mathcal{M}_I^2 = \begin{bmatrix} -\bar{m}_3^2 \tan\beta & \bar{m}_3^2 \\ \bar{m}_3^2 & -\bar{m}_3^2 \cot\beta \end{bmatrix}$$

$\mathrm{Det}(\mathcal{M}_I^2) = 0, \Rightarrow$ 1 eigenvalue is zero and corresponds to the Goldstone boson mass, while the other corresponds to the pseudoscalar Higgs mass:

$$M_A^2 = -\bar{m}_3^2(\tan\beta + \cot\beta) = -\frac{2\bar{m}_3^2}{\sin 2\beta}$$

The mixing angle θ which gives the physical fields is simply the angle β :

$$\begin{pmatrix} G^0 \\ A \end{pmatrix} = \begin{pmatrix} \cos\beta & \sin\beta \\ -\sin\beta & \cos\beta \end{pmatrix} \begin{pmatrix} P_1^0 \\ P_2^0 \end{pmatrix}$$

Same thing in the case of the charged Higgs boson :

$$\begin{pmatrix} G^\pm \\ H^\pm \end{pmatrix} = \begin{pmatrix} \cos\beta & \sin\beta \\ -\sin\beta & \cos\beta \end{pmatrix} \begin{pmatrix} H_1^\pm \\ H_2^\pm \end{pmatrix}$$

with a massless charged Goldstone and a charged Higgs boson with a mass

$$M_{H^\pm}^2 = M_A^2 + M_W^2$$

In the case of the CP–even Higgs bosons :

$$\mathcal{M}_R^2 = \begin{bmatrix} -\bar{m}_3^2 \tan\beta + M_Z^2 \cos^2\beta & \bar{m}_3^2 - M_Z^2 \sin\beta\cos\beta \\ \bar{m}_3^2 - M_Z^2 \sin\beta\cos\beta & -\bar{m}_3^2 \cot\beta + M_Z^2 \sin^2\beta \end{bmatrix}$$

injecting the expression of M_A^2 into $\mathcal{M}_R^2$:

$$M_R^2 = M_Z^2 \begin{pmatrix} \cos^2\beta & -\cos\beta\sin\beta \\ -\cos\beta\sin\beta & \sin^2\beta \end{pmatrix} + M_A^2 \begin{pmatrix} \sin^2\beta & -\cos\beta\sin\beta \\ -\cos\beta\sin\beta & \cos^2\beta \end{pmatrix}$$

one obtains for the CP–even Higgs boson masses :

$$M_{h,H}^2 = \frac{1}{2}\left[M_A^2 + M_Z^2 \mp \sqrt{(M_A^2 + M_Z^2)^2 - 4M_A^2 M_Z^2 \cos^2 2\beta}\right]$$

The physical CP–even Higgs bosons are obtained from the rotation of angle α :

$$\begin{pmatrix} H \\ h \end{pmatrix} = \begin{pmatrix} \cos\alpha & \sin\alpha \\ -\sin\alpha & \cos\alpha \end{pmatrix} \begin{pmatrix} H_1^0 \\ H_2^0 \end{pmatrix}$$

where the mixing angle α is given by :

$$\alpha = \frac{1}{2}\arctan\left(\tan 2\beta \, \frac{M_A^2 + M_Z^2}{M_A^2 - M_Z^2}\right) , \quad -\frac{\pi}{2} \leq \alpha \leq 0$$

- Out of the 6 parameters which describe the MSSM Higgs sector, $M_h, M_H, M_A, M_{H^\pm}, \beta$ and α, only 2 parameters, which can be taken as $\tan\beta$ and M_A, are free parameters at the tree–level.

- A strong hierarchy is imposed on the mass spectrum :

$$\mathrm{M}_H > \max(M_A, M_Z)$$

$$\mathrm{M}_{H^\pm} > M_W$$

very important constraint at the tree–level :

$$M_h \quad \leq \quad \min(M_A, M_Z) \cdot |\cos 2\beta| \leq M_Z$$

The Higgs couplings to gauge bosons:

Obtained from the kinetic terms of the fields H_1 and H_2 in the Lagrangian

$$\mathcal{L}_{\text{kin.}} = (D^\mu H_1)^\dagger (D_\mu H_1) + (D^\mu H_2)^\dagger (D_\mu H_2)$$

Expanding the covariant derivative D_μ and performing the transformations on the gauge and scalar fields to obtain the physical fields, one can identify the couplings $V_\mu V_\nu H_i$, $V_\mu H_i H_j$ (and $V_\mu V_\nu H_i H_j$) :

$$Z_\mu Z_\nu h \ : \ i g_Z M_Z \sin(\beta-\alpha) g_{\mu\nu} \quad , \quad Z_\mu Z_\nu H \ : \ i g_Z M_Z \cos(\beta-\alpha) g_{\mu\nu}$$
$$W_\mu^+ W_\nu^+ h \ : \ i g_W M_W \sin(\beta-\alpha) g_{\mu\nu} \quad , \quad W_\mu^+ W_\nu^- H \ : \ i g_W M_W \cos(\beta-\alpha) g_{\mu\nu}$$

$$Z_\mu h A : +\frac{g_Z}{2}\cos(\beta-\alpha)(p+p')_\mu \quad , \quad Z_\mu H A : -\frac{g_Z}{2}\sin(\beta-\alpha)(p+p')_\mu$$
$$Z_\mu H^+ H^- : -\frac{g_Z}{2}\cos 2\theta_W (p+p')_\mu \quad , \quad \gamma_\mu H^+ H^- : -ie(p+p')_\mu$$
$$W_\mu^\pm H^\pm h : \mp i\frac{g_W}{2}\cos(\beta-\alpha)(p+p')_\mu \quad , \quad W_\mu^\pm H^\pm H : \pm i\frac{g_2}{2}\sin(\beta-\alpha)(p+p')_\mu$$
$$W_\mu^\pm H^\pm A : \frac{g_W}{2}(p+p')_\mu \quad , \quad W_\mu^\pm G^\pm G^0 : \frac{g_W}{2}(p+p')_\mu$$

– CP–invariance forbids WWA, ZZA and $WZH^{\pm}$ couplings.
– The couplings of the neutral CP–even Higgs bosons : $G_{hVV} \propto \sin(\beta-\alpha)$ and $G_{HVV} \propto \cos(\beta-\alpha)$,

$$\cos^2(\beta-\alpha) = \frac{M_h^2(M_Z^2 - M_h^2)}{M_A^2(M_H^2 - M_h^2)}$$

The couplings G_{hVV} and G_{HVV} are complementary :

$$G_{hVV}^2 + G_{HVV}^2 = g_{H_{\rm SM}VV}^2$$

– From CP–invariance there are no Zhh, ZHh, ZHH and ZAA couplings (only ZhA,ZHA,$W^{\pm}H^{\pm}h/H/A$).
The couplings of h and H bosons to ZA and $W^{\pm}H^{\pm}$ are complementary :

$$\begin{aligned} G_{hAZ}^2 + G_{HAZ}^2 &= (4M_Z^2)^{-1} g_{H_{\rm SM}ZZ}^2 \\ G_{hH^{\pm}W^{\pm}}^2 + G_{HH^{\pm}W^{\pm}}^2 &= G_{AH^{\pm}W^{\pm}}^2 = (4M_W^2)^{-1} g_{H_{\rm SM}WW}^2 \end{aligned}$$

Yukawa couplings to fermions :

Originate from the superpotential W which leads to the Yukawa Lagrangian

$$\mathcal{L}_{\rm Yuk} = -\frac{1}{2}\sum_{ij}\left[\overline{\psi}_{iL}\frac{\partial^2 W}{\partial z_i \partial z_j}\psi_j + \text{h.c.}\right]$$

to be evaluated in terms of the scalar fields H_1 and H_2 :

$$\mathcal{L}_{\rm Yuk} \quad = \quad -\lambda_u[\bar{u}P_L u H_2^0 - \bar{u}P_L d H_2^+] - \lambda_d[\bar{d}P_L d H_1^0 - \bar{d}P_L u H_1^-] + \text{h.c.}$$

The fermion masses are generated when the neutral components of the Higgs fields acquire their vacuum expectation values and they are related to the Yukawa couplings by :

$$\lambda_u = \frac{\sqrt{2}m_u}{v_2} = \frac{\sqrt{2}m_u}{v\sin\beta} \quad , \quad \lambda_d = \frac{\sqrt{2}m_d}{v_1} = \frac{\sqrt{2}m_d}{v\cos\beta}$$

Expressing the fields H_1 and H_2 in terms of the physical fields, one obtains the Yukawa Lagrangian in terms of the fermion masses :

$$\begin{aligned}\mathcal{L}_{\rm Yuk} \;=\; & -\frac{g_2 m_u}{2M_W \sin\beta}\left[\bar{u}u(H\sin\alpha + h\cos\alpha) - i\bar{u}\gamma_5 u\, A\cos\beta\right] \\ & -\frac{g_2 m_d}{2M_W \cos\beta}\left[\bar{d}d(H\cos\alpha - h\sin\alpha) - i\bar{d}\gamma_5 d\, A\sin\beta\right] \\ & +\frac{g_2}{2\sqrt{2}M_W} V_{ud}\left\{H^+\bar{u}[m_d\tan\beta(1+\gamma_5) + m_u{\rm cot}\beta(1-\gamma_5)]d + {\rm h.c.}\right\}\end{aligned}$$

with V_{ud} the CKM matrix element (for quarks). The MSSM Higgs boson couplings to fermions are given by :

$$G_{huu} = i\frac{m_u}{v}\frac{\cos\alpha}{\sin\beta}\,, \qquad G_{Huu} = i\frac{m_u}{v}\frac{\sin\alpha}{\sin\beta}\,, \qquad G_{Auu} = \frac{m_u}{v}\cot\beta\,\gamma_5$$

$$G_{hdd} = -i\frac{m_d}{v}\frac{\sin\alpha}{\cos\beta}\,, \qquad G_{Hdd} = i\frac{m_d}{v}\frac{\cos\alpha}{\cos\beta}\,, \qquad G_{Add} = \frac{m_d}{v}\tan\beta\,\gamma_5$$

$$\begin{aligned}G_{H^+\bar{u}d} &= -\frac{i}{\sqrt{2}v}V_{ud}^*[m_d\tan\beta(1+\gamma_5) + m_u{\rm cot}\beta(1-\gamma_5)] \\ G_{H^-u\bar{d}} &= -\frac{i}{\sqrt{2}v}V_{ud}[m_d\tan\beta(1-\gamma_5) + m_u{\rm cot}\beta(1+\gamma_5)]\end{aligned}$$

–The couplings of $H^\pm$ have the same $\tan\beta$ dependence as the pseudoscalar A boson, for values $\tan\beta > 1$:

the A and $H^\pm$ couplings to isospin down–type fermions are enhanced,

the A and $H^\pm$ couplings to up–type fermions are suppressed.

–With a normalization factor $(i)g_2 m_f/2M_W = im_f/v$:

$$\begin{aligned}
g_{hbb} &= -\frac{\sin\alpha}{\cos\beta} = \sin(\beta-\alpha) - \tan\beta\cos(\beta-\alpha) \\
g_{htt} &= \frac{\cos\alpha}{\sin\beta} = \sin(\beta-\alpha) + \cot\beta\cos(\beta-\alpha) \\
g_{Hbb} &= \frac{\cos\alpha}{\cos\beta} = \cos(\beta-\alpha) + \tan\beta\sin(\beta-\alpha) \\
g_{Htt} &= \frac{\sin\alpha}{\sin\beta} = \cos(\beta-\alpha) - \cot\beta\sin(\beta-\alpha)
\end{aligned}$$

$\Rightarrow$ the $bb\,(tt)$ coupling of either the h or H boson are enhanced (suppressed) by a factor $\tan\beta$, depending on the magnitude of $\cos(\beta-\alpha)$ or $\sin(\beta-\alpha)$.

–The reduced pseudoscalar–fermion couplings are simply :

$$g_{Abb} = \tan\beta \;, \quad g_{Att} = \cot\beta$$

The trilinear and quartic scalar couplings :

Couplings between 3/4 Higgs fields can be obtained from V_H by performing derivatives :

$$\lambda_{ijk} = \left.\frac{\partial^3 V_H}{\partial H_i \partial H_j \partial H_k}\right|_{\langle H_1^0\rangle = v_1/\sqrt{2}, \langle H_2^0\rangle = v_2/\sqrt{2}, \langle H_{1,2}^{\pm}\rangle = 0}$$

$$\lambda_{ijkl} = \left.\frac{\partial^4 V_H}{\partial H_i \partial H_j \partial H_k \partial H_l}\right|_{\langle H_1^0\rangle = v_1/\sqrt{2}, \langle H_2^0\rangle = v_2/\sqrt{2}, \langle H_{1,2}^{\pm}\rangle = 0}$$

in units of $\lambda_0 = -iM_Z^2/v$:

$$\begin{aligned}
\lambda_{hhh} &= 3\cos 2\alpha \sin(\beta+\alpha) \\
\lambda_{Hhh} &= 2\sin 2\alpha \sin(\beta+\alpha) - \cos 2\alpha \cos(\beta+\alpha) \\
\lambda_{HHh} &= -2\sin 2\alpha \cos(\beta+\alpha) - \cos 2\alpha \sin(\beta+\alpha) \\
\lambda_{HHH} &= 3\cos 2\alpha \cos(\beta+\alpha) \\
\lambda_{HAA} &= -\cos 2\beta \cos(\beta+\alpha) \quad \lambda_{HH^+H^-} = -\cos 2\beta \cos(\beta+\alpha) + 2c_W^2 \cos(\beta-\alpha) \\
\lambda_{hAA} &= \cos 2\beta \sin(\beta+\alpha) \quad \lambda_{hH^+H^-} = \cos 2\beta \sin(\beta+\alpha) + 2c_W^2 \sin(\beta-\alpha)
\end{aligned}$$

The Higgs couplings to sfermions :

Come from the F terms,the D terms,$\mathcal{L}_{\rm tril.}$:

$$g_{\tilde{q}_i\tilde{q}'_j\Phi} = \sum_{k,l=1}^{2} (R^q)^{\rm T}_{ik}\; C^{kl}_{\Phi\tilde{q}\tilde{q}'} \left(R^{q'}\right)_{lj}$$

$$C_{h\tilde{q}\tilde{q}} = \begin{pmatrix} -\left(I_q^{3L} - Q_q s_W^2\right) M_Z^2 \sin(\beta+\alpha) + m_q^2 s_1^q & \frac{1}{2} m_q (A_q s_1^q + \mu s_2^q) \\ \frac{1}{2} m_q (A_q s_1^q + \mu s_2^q) & -Q_q s_W^2 M_Z^2 \sin(\beta+\alpha) + m_q^2 s_1^q \end{pmatrix}$$

$$C_{H\tilde{q}\tilde{q}} = \begin{pmatrix} \left(I_q^{3L} - Q_q s_W^2\right) M_Z^2 \cos(\beta+\alpha) + m_q^2 r_1^q & \frac{1}{2} m_q (A_q r_1^q + \mu r_2^q) \\ \frac{1}{2} m_q (A_q r_1^q + \mu r_2^q) & Q_q s_W^2 M_Z^2 \cos(\beta+\alpha) + m_q^2 r_1^q \end{pmatrix}$$

$$C_{A\tilde{q}\tilde{q}} = \begin{pmatrix} 0 & -\frac{1}{2} m_q \left[\mu + A_q (\tan\beta)^{-2I_3^q}\right] \\ \frac{1}{2} m_q \left[\mu + A_q (\tan\beta)^{-2I_3^q}\right] & 0 \end{pmatrix}$$

$$C_{H^\pm \tilde{t}\tilde{b}} = \frac{1}{\sqrt{2}} \begin{pmatrix} m_b^2 \tan\beta + m_t^2 \cot\beta - M_W^2 \sin 2\beta & m_b\,(A_b \tan\beta + \mu) \\ m_t\,(A_t \cot\beta + \mu) & m_t\, m_b (\tan\beta + \cot\beta) \end{pmatrix}$$

$$s_1^u = -r_2^u = \frac{\cos\alpha}{\sin\beta}\;,\quad s_2^u = r_1^u = \frac{\sin\alpha}{\sin\beta}\;,\quad s_1^d = r_2^d = -\frac{\sin\alpha}{\cos\beta}\;,\quad s_2^d = -r_1^d = \frac{\cos\alpha}{\cos\beta}$$

These couplings are potentially large since they involve terms $\propto m_t^2$ and $m_t A_t$.
In the case $\alpha = \beta - \frac{\pi}{2}$ (decoupling limit $M_A \gg M_Z$) :

$$
\begin{aligned}
g_{h\tilde{t}_1\tilde{t}_1} &= \cos 2\beta M_Z^2 \left[\frac{1}{2}\cos^2\theta_t - \frac{2}{3}s_W^2 \cos 2\theta_t\right] + m_t^2 + \frac{1}{2}\sin 2\theta_t m_t X_t \\
g_{h\tilde{t}_2\tilde{t}_2} &= \cos 2\beta M_Z^2 \left[\frac{1}{2}\sin^2\theta_t - \frac{2}{3}s_W^2 \cos 2\theta_t\right] + m_t^2 - \frac{1}{2}\sin 2\theta_t m_t X_t \\
g_{h\tilde{t}_1\tilde{t}_2} &= \cos 2\beta \sin 2\theta_t M_Z^2 \left[\frac{2}{3}s_W^2 - \frac{1}{4}\right] + \frac{1}{2}\cos 2\theta_t m_t X_t
\end{aligned}
$$

with $X_t = A_t - \mu \cot\beta$.

For large values of X_t $\Rightarrow$ enhance the $g_{h\tilde{t}_1\tilde{t}_1}$ coupling and make it larger than the top quark coupling of the h boson ($g_{htt} \propto m_t/M_Z$).

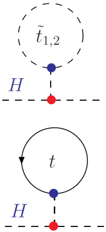
The upper bound on the lighter Higgs boson mass :

- The lighter CP–even Higgs boson should have a mass below M_Z.
- $M_h \simeq M_Z$, when $M_A > M_Z$ and $|\cos 2\beta| \simeq 1 \Rightarrow \beta \simeq \frac{\pi}{2} \Rightarrow$ large $\tan\beta$.
- The mixing angle $\alpha \simeq \frac{\pi}{2} - \beta \Rightarrow g_h$ are SM–like, $g_{huu} \simeq g_{hdd} \simeq g_{hVV} \simeq 1$. (the decoupling limit).
- Since the h boson is light it should have been observed at LEP2.
- It were not because of the radiative corrections which push its mass beyond the reach of LEP2.
- These radiative corrections can be very large since rather strong couplings, such as the Higgs couplings to the top quarks and to their spin–zero SUSY partners, are involved in the Higgs sector.

$\Rightarrow$ The RC due to top and stop quark loops should be incorporated in the MSSM Higgs sector!

In addition to the two–point functions including top and stop loops, one has also counterterm tadpole contributions :

$$\Delta M_h^2|_{\rm tad} = -\frac{3\lambda_t^2}{4\pi^2}\left[m_{\tilde t}^2 \log\left(\frac{\Lambda}{m_{\tilde t}}\right) - m_t^2 \log\left(\frac{\Lambda}{m_{\tilde t}}\right)\right]$$

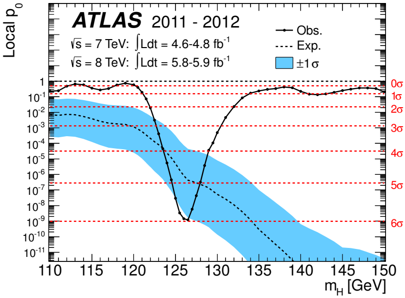


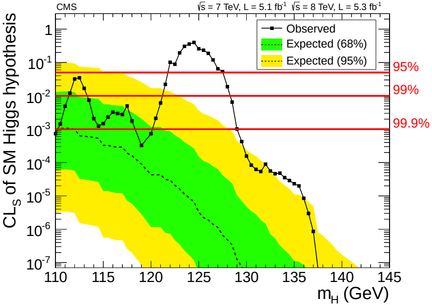
one obtains the total RC to the upper bound on M_h :

$$\Delta M_h^2 = \frac{3G_\mu}{\sqrt{2}\pi^2} m_t^4 \log\frac{M_S^2}{m_t^2}$$

The correction grows as : $\Delta M_h^2 \propto m_t^4$ and $\Delta M_h^2 \propto \log(m_{\tilde t}^2/m_t^2) \Rightarrow$ very large and shift M_h from M_Z to $M_h^{\rm max} \sim 140$ GeV.

In the MSSM : when the one–loop radiative corrections are included the h boson can be kinematically not accessible at LEP2 energies.





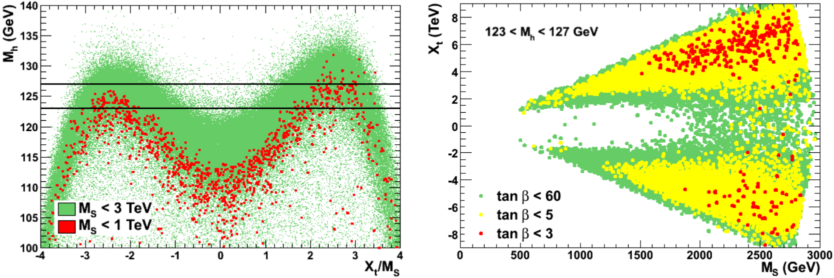
- In the SM, the Higgs mass is essentially a free parameter.
- In the MSSM, the lightest CP-even Higgs particle is bounded from above: $M_h^{max} \approx M_Z|cos2\beta|$+radiatives corrections $\lesssim 110 - 135 GeV$.
- Imposing M_h places very strong constraints on the MSSM parameters through their contributions to the radiative corrections.

The pMSSM(22 free parameters+SM) : most of the parameters have only a marginal impact on the MSSM Higgs masses.

Important parameters for MSSM Higgs mass : $\tan\beta$, M_A, $M_S = \sqrt{m_{\tilde{t}_1} m_{\tilde{t}_2}}$ and the mixing parameter in the stop sector $X_t = A_t - \mu \cot\beta$

$M_h^{\rm max}$ is obtained for the following choice :
i) the decoupling regime with a heavy pseudoscalar Higgs boson, $M_A \sim \mathcal{O}(\text{TeV})$;
ii) $\tan\beta$ large, $\tan\beta \gtrsim 10$;
iii) heavy stops, i.e. large M_S ($M_S = 3$ TeV as a maximal value);
iv) a stop trilinear coupling $X_t = \sqrt{6} M_S$, the so–called maximal mixing scenario.

We calculate the Higgs and superparticle spectrum in the MSSM with the full one–loop + dominant two–loop corrections.



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- the no–mixing scenario is entirely ruled out.
- only a small fraction of the typical-mixing scenario parameter space, with high tan β and M_S values, would survive.
- only the scenarios with large X_t/M_S values and, in particular, those close to the maximal mixing scenario $X_t/M_S \approx \sqrt{6}$ survive.

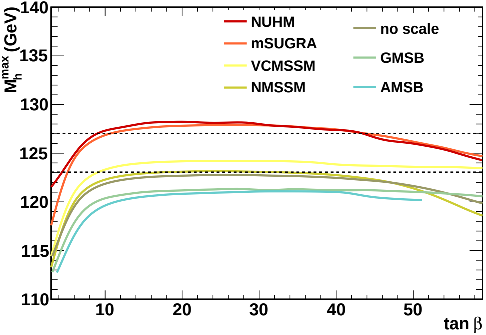
cMSSM scenarios : soft SUSY–breaking parameters obey universal boundary conditions at a high energy scale ⇒ reduce the number of input parameters.

These inputs are evolved via the MSSM renormalisation group equations down to the low energy scale.

3 classes of such models have been widely discussed in the literature :

- mSUGRA : SUSY–breaking is assumed to occur in a hidden sector which communicates with the visible sector only via flavour-blind gravitational interactions. Only 4 free parameters : tan β,m_0 (scalar masses), $m_{1/2}$(gaugino masses), A_0(trilinear scalar interactions), sign(μ).
- GMSB : SUSY–breaking is communicated to the visible sector via gauge interactions. The basic parameters : tan β,sign(μ),$M_{\rm mess}$(messenger field mass scale), $N_{\rm mess}$(the number of SU(5) representations of the messenger fields) and Λ(the SUSY–breaking scale in the visible sector)
- AMSB : SUSY–breaking is communicated to the visible sector via a super-Weyl anomaly. The basic parameters : tan β,sign(μ), m_0 (scalar masses at the GUT scale), $m_{3/2}$ (gravitino mass).

In the case of the mSUGRA scenario, we study 4 special cases:
– The no-scale scenario with the requirement $m_0 \approx A_0 \approx 0$. This model leads to a viable spectrum compatible with all present experimental constraints and with light staus for moderate $m_{1/2}$ and sufficiently high $\tan\beta$ values; the mass of the gravitino (the lightest SUSY particle) is a free parameter and can be adjusted to provide the right amount of dark matter.
– A model with $m_0 \approx 0$ and $A_0 \approx -\frac{1}{4} m_{1/2}$ which, approximately, corresponds to the constrained next-to–MSSM (cNMSSM) in which a singlet Higgs superfield is added to the two doublet superfields of the MSSM, whose components however mostly decouple from the rest of the spectrum. In this model, the requirement of a good singlino dark matter candidate imposes $\tan\beta \gg 1$ and the only relevant free parameter is thus $m_{1/2}$.
– A model with $A_0 \approx -m_0$ which corresponds to a very constrained MSSM (VCMSSM) for input values of the B_0 parameter close to zero.
– The non–universal Higgs mass model (NUHM) in which the universal soft SUSY–breaking scalar mass terms are different for the sfermions and for the two Higgs doublet fields. We will work in the general case in which, besides the four mSUGRA basic continuous inputs, there are two additional parameters which can be taken to be M_A and μ.



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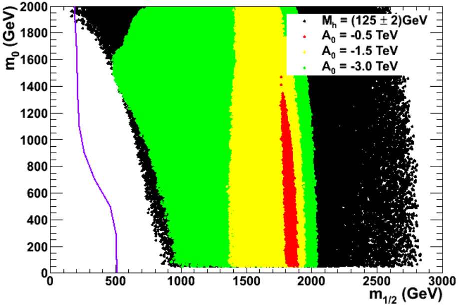
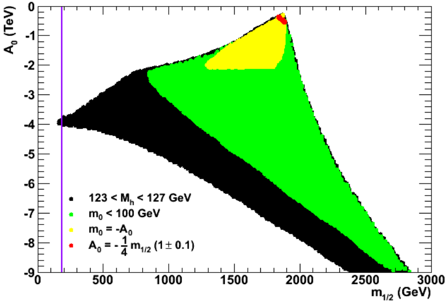
model	AMSB	GMSB	mSUGRA	no-scale	cNMSSM	VCMSSM	NUHM
$M_h^{\max}$	121.0	121.5	128.0	123.0	123.5	124.5	128.5

End of AMSB and GMSB in their minimal versions!

The upper bound on M_h in these scenarios can be understood by considering the allowed values of the trilinear coupling A_t(stop mixing parameter X_t) :

- In GMSB, one has $A_t \approx 0$ at relatively low scales and its magnitude does not significantly increase in the evolution down to the scale M_S; this implies that we are almost in the no–mixing scenario which gives a low value of M_h.
- In AMSB, one has a non-zero A_t; however, the ratio A_t/M_S with M_S determined from the overall SUSY breaking scale $m_{3/2}$ turns out to be rather small, implying again that we are close to the no–mixing scenario.
- Finally, in the mSUGRA model, since we have allowed A_t to vary in a wide range as $|A_0| \leq 9$ TeV, one can get a large A_t/M_S ratio which leads to a heavier Higgs particle. However, one cannot easily reach A_t values such that $X_t/M_S \approx \sqrt{6}$ so that we are not in the maximal–mixing scenario and the higher upper bound on M_h in the pMSSM is not reached.





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mSUGRA/cMSSM still survives, but only for negative values of A_0

Until now $M_S \lesssim 3$ TeV $\Rightarrow$ SUSY Higgses could be observed at the LHC.
BUT choice mainly dictated by fine–tuning considerations (rather subjective criterion).
If very large $M_S \Rightarrow$ no other scalar particle accessible at the LHC (except h)
$\rightarrow$ split SUSY scenario :

- The soft SUSY–breaking mass terms for the scalars are at $M_S \gg 1$ TeV (except 1 Higgs doublet).
- Gauginos and higgsinos, are left at the EWSB scale.
- The parameters : M_S, 1 Higgs mass, M_1, M_2, M_3, μ and $\tan\beta$.
- Boundary condition on the quartic Higgs coupling :

$$\lambda(M_S) = \frac{1}{4}\left[g^2(M_S) + g'^2(M_S)\right] \cos^2 2\beta \ .$$

- If the scalars are very heavy $\Rightarrow$ RC in the Higgs sector that are enhanced by large logarithms, $\log(M_{\rm EWSB}/M_S)$.
- One has to properly decouple the heavy states from the low-energy theory and resum the large log corrections. N. Bernal, A. Djouadi and P. Slavich, JHEP 0707 (2007) 016

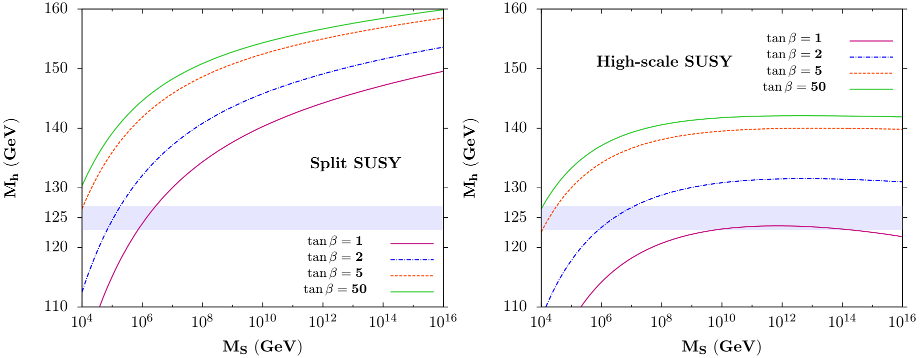
A radical attitude :

Assume that gauginos and higgsinos are at $M_S \rightarrow$ high–scale SUSY model :

- Abandon the SUSY solution to the fine-tuning problem.
- Abandon the solution to the dark matter problem (LSP).
- Abandon the unification of the gauge coupling constants.

- The matching of SUSY/EWSB is encoded in the Higgs quartic coupling λ.
- SUSY would still lead to a light Higgs boson whose mass will contain information on M_S and $\tan\beta$.

$\Rightarrow$ Still a trace of SUSY at low energy



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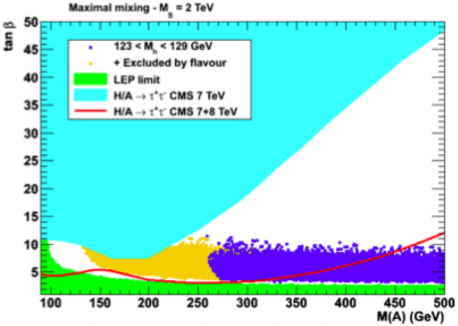
- The maximal M_h values are obtained at high $\tan\beta$ and M_S.
- Strong constraints on the parameters of these two models.
- $\tan\beta \approx 1$: very large M_S needed.
- At high $\tan\beta$: $M_S \approx 10^4$ GeV.
- Even in these extreme scenarios, SUSY should manifest itself at scales much below $M_{\rm GUT}$.

Conclusion before to go further :

- The Higgs sector can play an important role in constraining SUSY.
- Several constrained MSSM scenarios are ruled out because of a Higgs of 125 GeV.

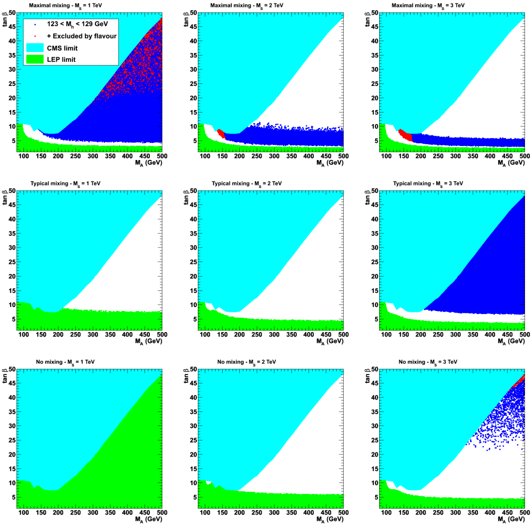
There are other constaints on pMSSM to be included :

- The non-observation of Higgs bosons at LEP2 excludes $\tan\beta \lesssim 3$ and $\tan\beta \approx 5$–10 for $M_A \lesssim 100$ GeV.
- The latest results of CMS on the search of resonances decaying into $\tau^+\tau^-$.
- The $B_s \to \mu^+\mu^-$ constraint.



A. Arbey, M. Battaglia, A. Djouadi, F.Mahmoudi, arXiv:1207.1348

$M_S = 1, 2, 3$ TeV for the maximal,typical,zero mixing scenarios :



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The various regimes of the pMSSM Higgs sector :

The decoupling regime :

- occurs for $M_A \gtrsim 300$ GeV for low $\tan\beta$ and $M_A \gtrsim M_h^{\max}$ for $\tan\beta \gtrsim 10$.
- One will have a SM–like Higgs boson $h \equiv H_{\rm SM}$.
- The heavier $H/A/H^{\pm}$ bosons, degenerated in mass, almost decouple.

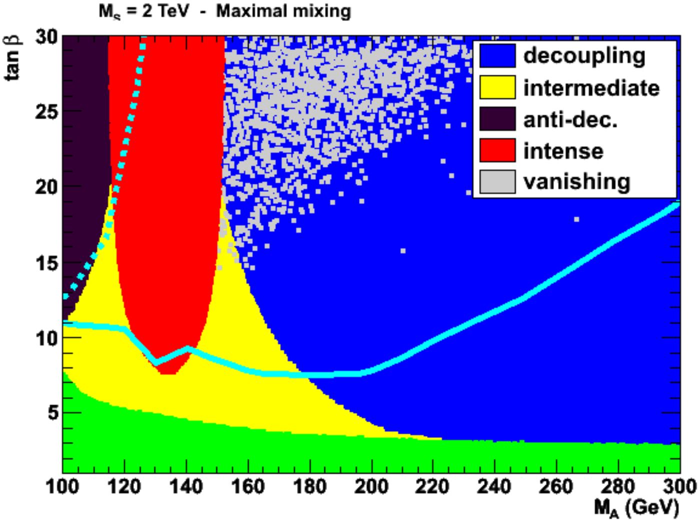
$\rightarrow$ The MSSM Higgs sector reduces effectively to the SM Higgs sector, but with a light h boson.

The anti–decoupling regime : occurs for $M_A \lesssim M_h^{\max}$, and is exactly opposite to the decoupling regime. The roles of the h and H bosons are reversed.

The intense–coupling regime : occurs when $M_A \approx M_h^{\max}$ then $M_h \sim M_H \sim M_A \sim M_h^{\max}$ (3 pseudo-scalar like Higgs particles).

The intermediate–coupling regime : occurs for $\tan\beta \lesssim$ 5–10 and $M_A \lesssim$ 300–500 GeV. Both CP–even Higgs bosons have non–zero couplings to gauge bosons and their couplings to down–type (up–type) fermions are not strongly enhanced (suppressed) since $\tan\beta$ is not too large.

The vanishing–coupling regime occurs for large values of $\tan\beta$ and intermediate to large M_A values as well as for specific values of the other MSSM parameters. The RC lead to a strong suppression of the couplings of one of the CP–even Higgs bosons to fermions or gauge bosons.



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In the pMSSM few parameters plays a significant role in the Higgs sector :
This mass matrix receives RC at higher orders :

$$\mathcal{M}^2 = \begin{bmatrix} \mathcal{M}^2_{11} + \Delta\mathcal{M}^2_{11} & \mathcal{M}^2_{12} + \Delta\mathcal{M}^2_{12} \\ \mathcal{M}^2_{12} + \Delta\mathcal{M}^2_{12} & \mathcal{M}^2_{22} + \Delta\mathcal{M}^2_{22} \end{bmatrix}$$

$\Delta\mathcal{M}^2_{ij}$: leading one–loop radiative corrections controlled by λ_t+subleading contributions at 1–loop+the leading logarithmic contributions at 2–loops :
Carena,Haber, 2003

$$\Delta\mathcal{M}^2_{11} = -\frac{v^2\sin^2\beta}{32\pi^2}\bar{\mu}^2\left[x_t^2\lambda_t^4(1+c_{11}\ell_S)+a_b^2\lambda_b^4(1+c_{12}\ell_S)\right]$$

$$\Delta\mathcal{M}^2_{12} = -\frac{v^2\sin^2\beta}{32\pi^2}\bar{\mu}\left[x_t\lambda_t^4(6-x_t a_t)(1+c_{31}\ell_S)-\bar{\mu}^2 a_b\lambda_b^4(1+c_{32}\ell_S)\right]$$

$$\Delta\mathcal{M}^2_{22} = \frac{v^2\sin^2\beta}{32\pi^2}\left[6\lambda_t^4\ell_S(2+c_{21}\ell_S)+x_t a_t\lambda_t^4(12-x_t a_t)(1+c_{21}\ell_S)-\bar{\mu}^4\lambda_b^4(1+c_{22}\ell_S)\right]$$

with : $\ell_S = \log(M_S^2/m_t^2)$, $\bar{\mu} = \mu/M_S$, $a_{t,b} = A_{t,b}/M_S$, $x_t = X_t/M_S$,
$c_{ij} = \frac{1}{32\pi^2}(t_{ij}\lambda_t^2 + b_{ij}\lambda_b^2 - 32g_3^2)$, $(t_{11}, t_{12}, t_{21}, t_{22}, t_{31}, t_{32}) = (12, -4, 6, -10, 9, -7)$,
$(b_{11}, b_{12}, b_{21}, b_{22}, b_{31}, b_{32}) = (-4, 12, 2, 18, -1, 15)$

A very simple and good approximation :

$$\Delta\mathcal{M}^2_{11} \sim \Delta\mathcal{M}^2_{12} \sim 0 \,,$$

$$\Delta\mathcal{M}^2_{22} = \epsilon \sim \frac{3\,\bar{m}_t^4}{2\pi^2 v^2 \sin^2\beta}\left[\log\frac{M_S^2}{\bar{m}_t^2} + \frac{X_t^2}{M_S^2}\left(1 - \frac{X_t^2}{12\,M_S^2}\right)\right]$$

The radiately corrected CP–even Higgs boson masses :

$$M^2_{h,H} = \frac{1}{2}(M_A^2 + M_Z^2 + \epsilon)\left[1 \mp \sqrt{1 - 4\frac{M_Z^2 M_A^2 \cos^2 2\beta + \epsilon(M_A^2 \sin^2\beta + M_Z^2\cos^2\beta)}{(M_A^2 + M_Z^2 + \epsilon)^2}}\right]$$

$$M_{H^\pm} = \sqrt{M_A^2 + M_W^2}$$

For $M_A \gg M_Z$:

$$M_h \overset{M_A \gg M_Z}{\rightarrow} \sqrt{M_Z^2\cos^2 2\beta + \epsilon\sin^2\beta}\left[1 + \frac{\epsilon M_Z^2\cos^2\beta}{2M_A^2(M_Z^2 + \epsilon\sin^2\beta)} - \frac{M_Z^2\sin^2\beta + \epsilon\cos^2\beta}{2M_A^2}\right]$$

$$M_H \overset{M_A \gg M_Z}{\rightarrow} M_A\left[1 + \frac{M_Z^2\sin^2 2\beta + \epsilon\cos^2\beta}{2M_A^2}\right] \,,\quad M_{H^\pm} \overset{M_A \gg M_Z}{\rightarrow} M_A\left[1 + \frac{M_W^2}{2M_A^2}\right]$$

In exact decoupling $M_A/M_Z \to \infty$: $M_H = M_A = M_{H^+}$ and :

$$M_h \equiv M_h^{\max} = \sqrt{M_Z^2\cos^2 2\beta + \epsilon\sin^2\beta}$$

If we consider only $\Delta\mathcal{M}^2_{22}$ which involves the by far dominant stop,top sector correction,

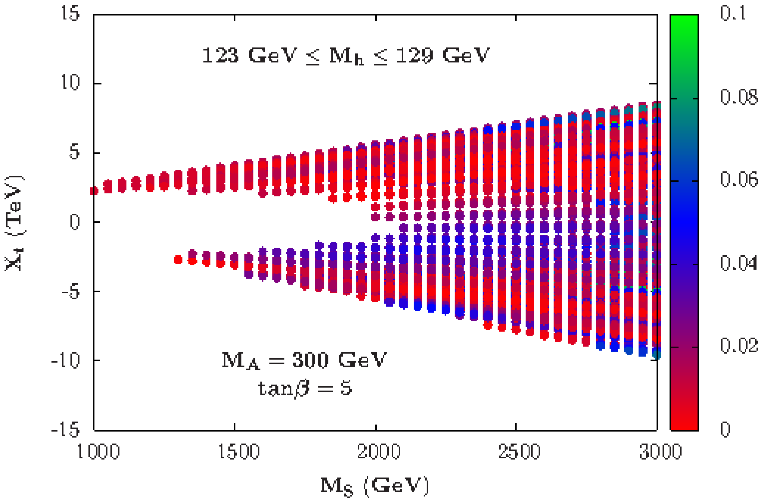
$\Delta\mathcal{M}^2_{22}$ (M_S!) is fixed by M_h :

$$\epsilon = \frac{-M_A^2 M_h^2 + M_h^4 - M_h^2 M_Z^2 + M_A^2 M_Z^2 \cos^2 2\beta}{M_h^2 - M_Z^2 \cos^2\beta - M_A^2 \sin^2\beta}$$

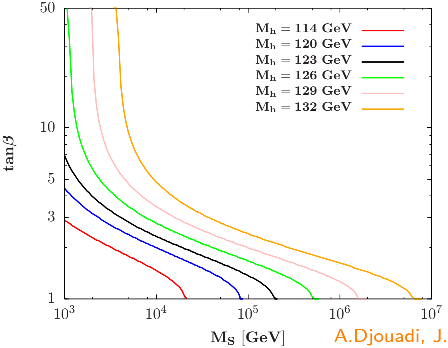
The RC in the MSSM are fixed by the constraint $M_h = 125$ GeV!

$\Rightarrow$ we can simply write the MSSM parameters in terms of M_A,$\tan\beta$ and M_h :

$$\begin{aligned} M_H^2 &= M_A^2 + M_Z^2 + \epsilon - M_h^2 \\ \alpha &= -\arctan\left(\frac{M_A^2 + M_Z^2}{M_A^2 \tan\beta - 2M_h^2/\sin 2\beta + M_Z^2 \cot\beta}\right) \\ &\ldots \end{aligned}$$


123 GeV ≤ M_h ≤ 129 GeV
M_A = 300 GeV
tanβ = 5
X_t (TeV)
M_S (GeV)
15
10
5
0
-5
-10
-15
1000
1500
2000
2500
3000
0.1
0.08
0.06
0.04
0.02
0

- Large value of M_h+ nonobservation of superparticles at the LHC $\Rightarrow$ suggest a high M_S.
- $\tan\beta \lesssim 3$ usually "excluded" by LEP2: $M_h \gtrsim 114$ GeV with $M_S \sim 1$ TeV. But we can be more relaxed: $M_S \gg M_Z \Rightarrow \tan\beta \approx 1$ could be allowed!



A.Djouadi, J.Q.,arXiv:1304.1787

- We turn $M_h \sim M_Z|\cos 2\beta| + RC$ to $RC = 126GeV - f(M_A, \tan\beta)$ ie. we trade the RC with the measured M_h
  $\Rightarrow$ MSSM with only 2 inputs at HO: M_A,$\tan\beta$ a model indep. effective approach!

Φ	$g_{\Phi\bar{u}u}$	$g_{\Phi\bar{d}d}$	$g_{\Phi VV}$	$g_{\Phi AZ}/g_{\Phi H^+W^-}$
h	$\cos\alpha/\sin\beta$	$-\sin\alpha/\cos\beta$	$\sin(\beta-\alpha)$	$\propto\cos(\beta-\alpha)$
H	$\sin\alpha/\sin\beta$	$\cos\alpha/\cos\beta$	$\cos(\beta-\alpha)$	$\propto\sin(\beta-\alpha)$
A	$\cot\beta$	$\tan\beta$	0	$\propto 0/1$

The decoupling limit is controlled by $g_{HVV}=\cos(\beta-\alpha)$:

$$g_{HVV} \overset{M_A\gg M_Z}{\longrightarrow} \chi \equiv \frac{1}{2}\frac{M_Z^2}{M_A^2}\sin 4\beta - \frac{1}{2}\frac{\epsilon}{M_A^2}\sin 2\beta \quad \to 0$$

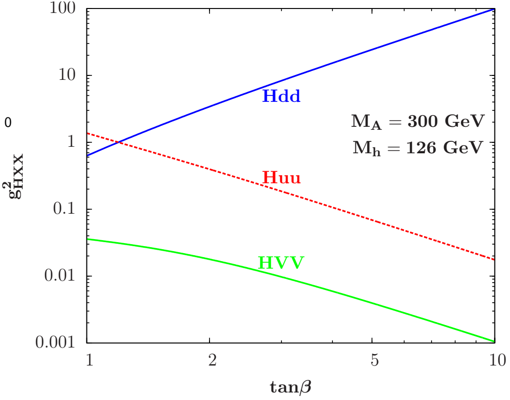
Tree–level part: doubly suppressed in both the $\tan\beta \gg 1$ and $\tan\beta \sim 1$ cases.

$$\sin 4\beta = \frac{4\tan\beta(1-\tan^2\beta)}{(1+\tan^2\beta)^2} \to \begin{cases} -4/\tan\beta & \text{for } \tan\beta \gg 1 \\ 1-\tan^2\beta & \text{for } \tan\beta \sim 1 \end{cases} \to 0$$

The radiative part : behave as $-\epsilon/M_A^2 \times \cot\beta$, also vanishes at high $\tan\beta$ values $\Rightarrow$ the decoupling limit $g_{HVV}\to 0$ is reached very quickly at high $\tan\beta$, as soon as $M_A \gtrsim M_h^{\max}$. Instead, for $\tan\beta \approx 1$, this radiatively generated component is maximal.

$$\begin{array}{llll} g_{huu} & \overset{M_A\gg M_Z}{\longrightarrow} & 1+\chi\cot\beta & \to 1 \\ g_{hdd} & \overset{M_A\gg M_Z}{\longrightarrow} & 1-\chi\tan\beta & \to 1 \\ g_{Huu} & \overset{M_A\gg M_Z}{\longrightarrow} & -\cot\beta+\chi & \to -\cot\beta \\ g_{Hdd} & \overset{M_A\gg M_Z}{\longrightarrow} & +\tan\beta+\chi & \to +\tan\beta \end{array}$$

g_{HVV} is non–zero at low $\tan\beta$. The g_{Htt} and g_{Att} is significant at low $\tan\beta$ values, $g_{Htt}^2 \gtrsim 0.1$ for $\tan\beta \lesssim 3$. $\Rightarrow$ $H/A/H^\pm$ bosons can have sizable couplings to top quarks and massive gauge bosons if $\tan\beta \sim 3$.



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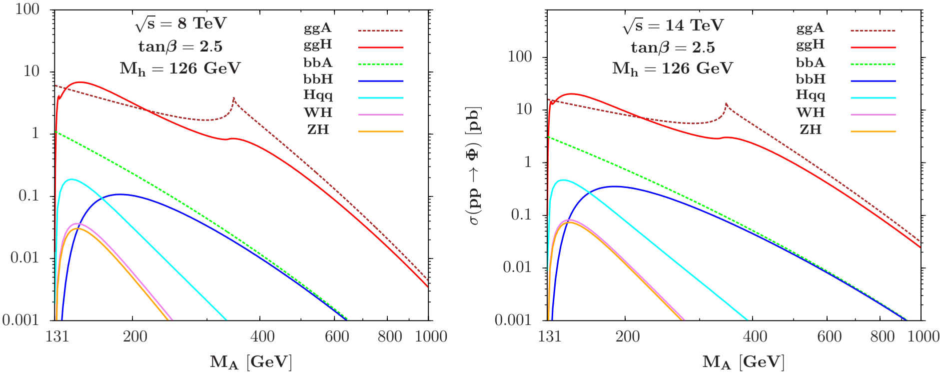
The high $\tan\beta$ *regime ("simple" pattern)* :

- couplings of non–SM like Higgs bosons to b quarks and to τ leptons are strongly enhanced, to top quarks and massive gauge bosons suppressed.
- $\text{BR}(H/A \to \tau^+\tau^-) \approx 10\%$ and $\text{BR}(H/A \to b\bar{b}) \approx 90\%$.
- $\text{BR}(H^\pm \to \tau\nu_\tau) \approx 100\%$ for $M_{H^\pm} \gtrsim m_t + m_b$,
  otherwise : $\text{BR}(H^\pm \to tb) \approx 90\%$
- The processes $gg/q\bar{q} \to b\bar{b} + H/A$, must be considered.
- all the other production channels are irrelevant.
- The most powerful search channel at the LHC :

$$pp \to gg + b\bar{b} \to H/A \to \tau^+\tau^-$$

The phenomenology of $A, H, H^{\pm}$ is richer at low $\tan\beta$ and leads to a production and decay pattern that is slightly more involved than in the high $\tan\beta$ regime :



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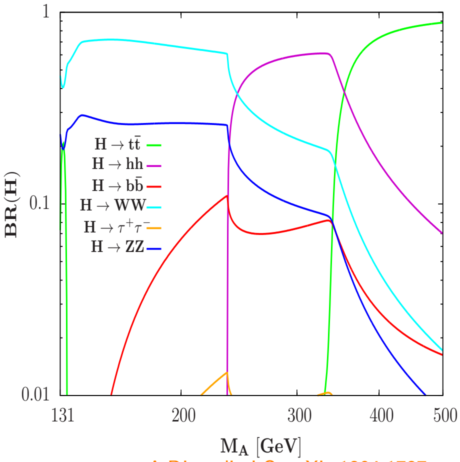
Decay pattern of the heavier MSSM Higgses in the low tan β *regime:*

We assume that Higgs decays into SUSY particles are kinematically closed.

- Above $t\bar{t}$ threshold, $H \to t\bar{t}$ dominate.
- Below $t\bar{t}$ threshold, $H \to WW$ and ZZ with substantial rates.
- $2M_h \lesssim M_H \lesssim 2m_t$, very interesting channel $H \to hh$.



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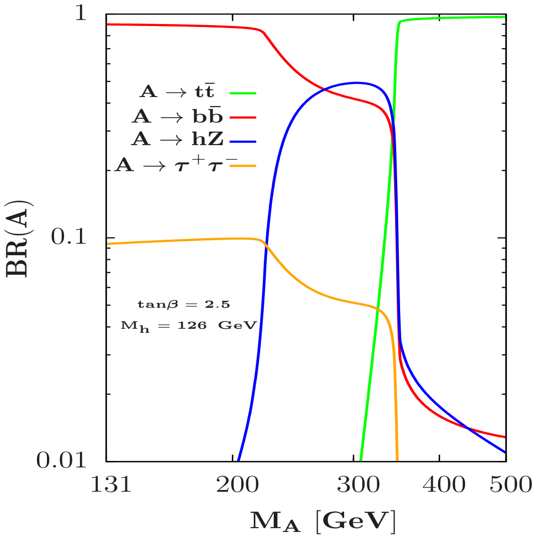
- Above the $t\bar{t}$ threshold, $A \to t\bar{t}$ dominate.
- For $M_A \gtrsim M_h + M_Z$ GeV, $A \to hZ$ with a significant rate below the $t\bar{t}$ threshold.
- BR($A \to \tau\tau$) $\gtrsim$ 5% up to $M_A \approx 2m_t$.



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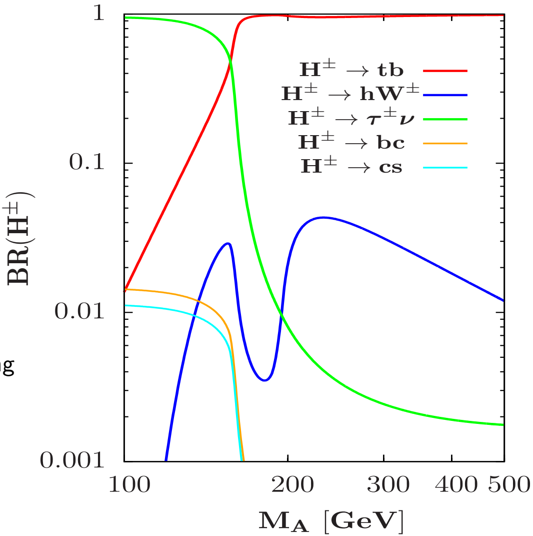
- $H^+ \to t\bar{b}$ dominant for $\tan\beta \lesssim 3$.
- $\text{BR}(H^+ \to t\bar{b}) \gtrsim 1\%$ for $M_{H^\pm} \approx 130$ GeV.
- $H^+ \to Wh$ important.
- $M_{H^\pm} \lesssim 170$ GeV, $H^+ \to c\bar{s}$ and $H^+ \to c\bar{b}$ (flavor changing mode) have rates $\sim 1\%$ ($\approx 10\%$, for $\tan\beta \sim 1$).



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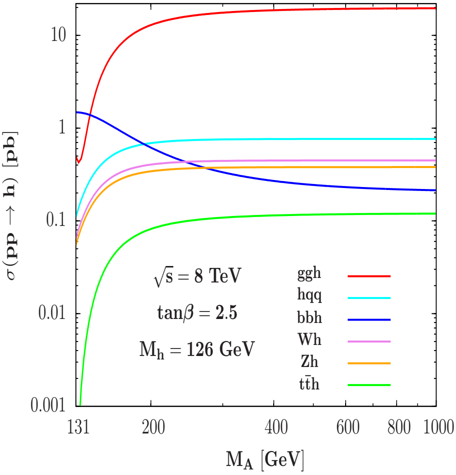
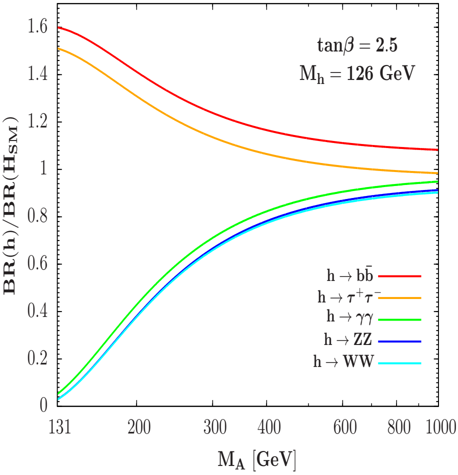
The case of the h boson :





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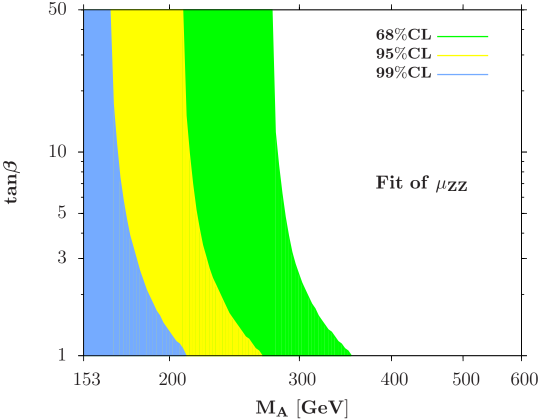
Constraints from the h boson rates :

- $\mu_{XX} = \frac{\sigma(H)BR(H\to XX)}{\sigma(H)_{SM}BR_{SM}(H\to XX)}$
- cleanest observable : the signal strength μ_{ZZ} in the search channel $h \to ZZ$.
- The combination of ATLAS and CMS data in the ZZ : $\mu_{ZZ} = 1.10 \pm 0.22 \pm 0.20$ (25 fb^{-1}, 7+8 TeV data).
- Constrains the couplings of the h state in $[\tan\beta, M_A]$ parameter space :



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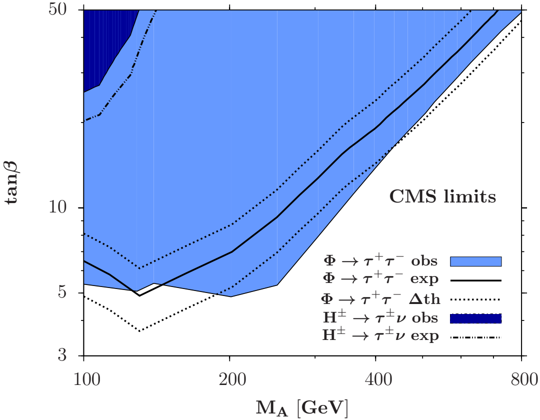
Constraints from the heavier Higgs searches at high $\tan\beta$:

- $\sigma(pp \to H/A) \times BR(H/A \to \tau\tau)$ is very robust $\Rightarrow$ exclusion limit almost model independent.
- CMS $H/A \to \tau\tau$ analysis : constraint very restrictive for $M_A \lesssim 250$ GeV, excludes $\tan\beta \gtrsim 5$.
- Caveat : ATLAS/CMS constraint for a specific benchmark : $X_t/M_S = \sqrt{6}$ and $M_S = 1$ TeV;
- exclusion limit can be obtained in any MSSM scenario with the only assumption that SUSY particles are too heavy to affect $\sigma(pp \to H/A) \times \mathrm{BR}(H/A \to \tau\tau)$ by more than 25%, (th. uncertainty).



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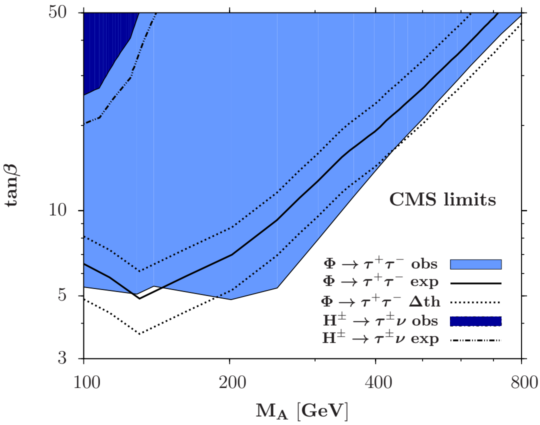
- Constraint from $H^- \to \tau\nu$ ($t \to H^+ b \to \tau\nu b$). ATLAS & CMS results with ≈ 5 fb^{-1} at $\sqrt{s} = 7$ TeV.
- Constraint only $M_A \lesssim 150$.
- The search is sensitive to the very high tan β region which is completely excluded by the $\tau\tau$ search (need much more data).



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Important remarks:

The LEP2 constraint excludes the entire low $\tan\beta$ regime, $\tan\beta \lesssim 3$, and at low $M_A \approx 100$ GeV, $\tan\beta$ values up to $\tan\beta \approx 10$.

We have removed the region excluded by the bound $M_h \gtrsim 114$ GeV from negative Higgs searches at LEP2 :

- the "observation" constraint $123\ \mathrm{GeV} \lesssim M_h \lesssim 129\ \mathrm{GeV}$ is by far stronger.
- if $M_S = 1$ TeV and maximal stop mixing : $\tan\beta \lesssim 5$ and $\tan\beta \gtrsim 20$ for any M_A value would be excluded simply by requiring $123\ \mathrm{GeV} \lesssim M_h \lesssim 129\ \mathrm{GeV} \Rightarrow$ excluded regions would be completely different for other M_S and X_t.
- The LEP2 M_h constraint can be simply evaded for any value of $\tan\beta$ or M_A by assuming large $M_S \Rightarrow$ open the very interesting low $\tan\beta$ region

Low $\tan\beta$ region can be probed in a model independent way by Higgs search channels involving the $H, A, H^\pm$ bosons....

Extrapolation to the low $\tan\beta$ *region* :

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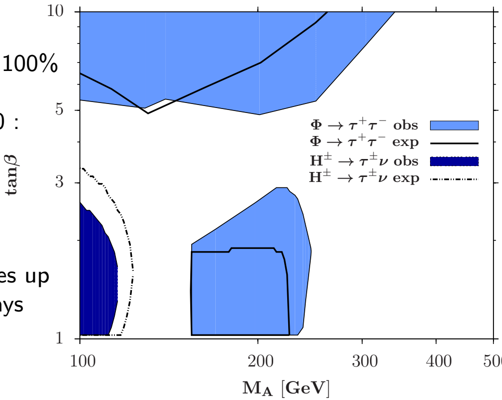
- BR($t \to bH^+$) significant at low $\tan\beta$ ($g_{tbH^+} \propto \bar{m}_t/\tan\beta$) + BR($H^\pm \to \tau\nu \sim 100\%$ $\Rightarrow pp \to t\bar{t}$ with $t \to bH^+ \to b\tau\nu$ are comparable for $\tan\beta \approx 3$ and $\tan\beta \approx 30$: probe the low $\tan\beta$.
- $pp \to H/A \to \tau\tau$ useful at low $\tan\beta$: for $\tan\beta \approx 1$, $gg \to H/A$ large ($g_{H/Att} \propto \bar{m}_t/\tan\beta$). $A \to \tau^+\tau^-$ stays significant for M_A values up to the $t\bar{t}$ threshold $\Rightarrow gg \to A \to \tau\tau$ stays large : CMS search limit is effective and excludes low $\tan\beta$, for $M_A \approx 350$ GeV.



Low $\tan\beta$ areas,thought to be buried under the LEP2 exclusion bound on M_h, are now open territory for heavy MSSM Higgs hunting!
This can be done not only in these 2 channels but also in a plethora of channels...

The **H** $\to$ **WW**, **ZZ** *channels* :

- The $H \to WW$ particularly useful in the region $160 \lesssim M_H \lesssim 180$ GeV where the branching ratio is close to 100%.
- The search modes that are most useful at relatively low M_H should be $pp \to H \to ZZ \to 4\ell^{\pm}$ and $pp \to H \to WW \to 2\ell 2\nu$: When the 2 processes are combined, the sensitivity is an order of magnitude larger than for the SM Higgs for masses below 400 GeV and one can thus afford a substantial reduction of the couplings g_{Htt} and g_{HVV} which should allow to probe $\tan\beta$ values significantly higher than 1.
- For $M_H \gtrsim 300$ GeV, one could also add the $pp \to H \to ZZ \to 2\ell 2q, 2\nu 2q, 2\ell 2\nu$ and $pp \to H \to WW \to \ell\nu 2q$ channels to increase the statistics (recently done by CMS).
- One difference with the SM Higgs : the SM Higgs has a large total width at high masses ($M^3_{H_{\rm SM}}$) & the MSSM H boson remains narrow as the coupling g_{HVV} is suppressed $\Rightarrow$ reduce the continuum ZZ background by selecting smaller bins for M_{ZZ}.

The $H/A \to t\bar{t}$ channels :

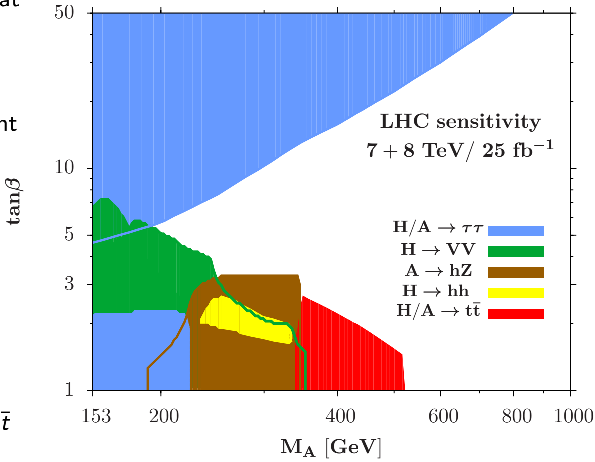
- This search channel has not been considered in the case of the SM Higgs boson for two reasons :
  1. For $M_{H_{\rm SM}} \gtrsim 350$ GeV, $H_{\rm SM} \to WW, ZZ$ still relevant and largely dominate over the $H_{\rm SM} \to t\bar{t}$.
  2. The continuum $t\bar{t}$ background was thought to be overwhelmingly large as it had to be evaluated in a large mass window because of the large Higgs total width.
- Situation very different in MSSM : total width for heavy H and A states are much smaller, less than $\lesssim 20$ GeV for any $\tan\beta \gtrsim 1$ value for $M_{H,A} \lesssim 500$ GeV and grow linearly with the Higgs masses beyond this value $\Rightarrow$ one can integrate the $t\bar{t}$ continuum background in a smaller invariant mass bin and significantly enhance the signal to background ratio.
- BR$(H/A \to t\bar{t}) \approx 100\%$ for $\tan\beta \lesssim 3$ if kinematically open.

Search for $H/A \to t\bar{t}$ will certainly be more favorable for the MSSM at low $\tan\beta$ than in the SM.

The **A → Zh** *channel* :

- $\sigma(gg \to A) > \sigma(gg \to H_{SM})$ at low tan β.
- For $M_h + M_Z \lesssim M_A \lesssim 2m_t$ BR$(A \to hZ)$ is large for tan $\beta \approx 3$ and largely dominant for tan $\beta \approx 1 \Rightarrow$ for $M_A = 210$–350 GeV $\sigma(gg \to A \to hZ)$ should be very high at low tan β.

The **H → hh** *channel* : very large production rates in the low tan β regime in the mass range 250 GeV $\lesssim M_H \lesssim$ 350 GeV when the decay channels $H \to hh$ is kinematically open and the $H \to t\bar{t}$ mode is closed.



Conclusion:

- After the observation of the 126 GeV SM–like Higgs boson by the ATLAS and CMS collaborations, the next challenge at the LHC should be to search for new phenomena beyond the SM. This can be done not only by refining the precision determination of the properties of the observed Higgs particle to pin down small deviations of its couplings from the SM expectations, but also by looking for the direct production of new states.
- We have considered the production of the heavier H, A and $H^{\pm}$ bosons of the MSSM at the LHC, focusing on the low $\tan\beta$ regime.
- We have first shown that this area of the MSSM parameter space, which was long thought to be excluded, is still viable provided that the SUSY scale is assumed to be very high, $M_S \gtrsim 10$ TeV.
- We have used a simple but not too inaccurate approximation to describe the radiative corrections to the Higgs sector, in which the unknown scale M_S and stop mixing parameter X_t are traded against the measured M_h.

Conclusion ...

- The neutral H/A states can still be produced with large rates, and they will decay into a variety of interesting final states such as $H \to WW, ZZ$, $H \to hh$, $H/A \to t\bar{t}$, $A \to hZ$.
- Interesting decays can also occur in the case of the charged Higgs bosons, e.g. $H^+ \to hW, c\bar{s}, c\bar{b}$.
- These modes come in addition to the two channels $H/A \to \tau^+\tau^-$ and $t \to bH^+ \to b\tau\nu$ which are currently being studied by ATLAS and CMS and which are very powerful in constraining the parameter space at high $\tan\beta$ values and, as is shown here, also at low $\tan\beta$ values.
- We have shown that already with the current LHC data at $\sqrt{s} = 7+8$ TeV, the area with small $\tan\beta$ and M_A values can be probed by simply extrapolating to the MSSM Higgs sector the available analyses in the search of the SM Higgs boson.
- All this promises a very nice and exciting program for Higgs searches at the LHC in both the present and future runs!

Merci !