Chapter 2
Physics Beyond the Standard Model

The SM has been demonstrated as successful by many measurements performed at high-energy experiments. In particular, the discovery of a new particle compatible with the SM scalar boson, considered as the cornerstone of the SM, has consecrated the theory. However there are strong indications that the SM is only a low-energy expression of a more global theory. If new physics shows up beyond the SM, it could be related to the scalar sector. Some motivations for the existence of BSM physics are detailed in Sect. 2.1, while two-Higgs-doublet models, supersymmetry including the minimal supersymmetric extension of the SM, and two-Higgs-doublet models extended with a scalar singlet, are presented in Sects. 2.2, 2.3 and 2.4 respectively. Section 2.5 discusses how to uncover a possibly extended scalar sector at the LHC, while the chapter ends, in Sect. 2.6, with a comparison between the reach of precision measurements and direct discovery of new scalars, in a simple benchmark scenario.

2.1 Motivations for New Physics

The existence of new physics beyond the SM [1, 2] is strongly motivated. Some motivations are based on direct evidence from observation, such as the existence of neutrino masses, the existence of dark matter and dark energy, or the matter-antimatter asymmetry, while others come from conceptual problems in the SM, such as the large number of free parameters, the "hierarchy problem" or the coupling unification. Each of these issues is shortly described in the next sections.

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2.1.1 Neutrino Masses

It is well-established by experiments with solar, atmospheric, reactor and accelerator neutrinos, that neutrinos can oscillate and change their flavor in flight [3, 4]. Such oscillations are possible if neutrinos have masses. Flavor neutrinos (ν_e, ν_μ, ν_τ) are then linear combinations of the fields of at least three mass eigenstate neutrinos ν_1, ν_2, ν_3. Only upper limits on the neutrino masses have been set as of now ($m_\nu <$ 2 eV), but the differences between the neutrino squared masses have been measured: $\Delta m^2_{12} = (7.53 \pm 0.18) \times 10^{-5}$ eV^2 and $\Delta m^2_{32} = (2.44 \pm 0.06) \times 10^{-3}$ eV^2 [5].

2.1.2 Dark Matter and Dark Energy

In 1933, Zwicky carried measurements of the velocities of galaxies in the Coma cluster, using the Doppler shift of their spectra [6]. With the virial theorem, he could relate these results to the total mass of the Coma cluster. Zwicky also measured the total light output of the cluster, and compared the ratio of the luminosity to the mass for the Coma cluster and for the nearby Kapteyn stellar system. The two orders of magnitude difference between both of them made him conclude that the Coma cluster contains some massive matter that does not radiate: dark matter. The ordinary matter that surrounds us and is described by the SM, only represents 5% of the mass/energy content of the universe. Astrophysical evidence indicates that dark matter contributes approximately to 27%, and dark energy to 68% of this content. Measurements of the temperature and polarization anisotropies of the cosmic microwave background (CMB) by the Planck experiment could determine a density of cold non-baryonic matter [7].

Nowadays little is also known about dark energy, which is responsible for the accelerated expansion of the universe.

2.1.3 Asymmetry Between Matter and Antimatter

It is believed that matter and antimatter were produced in exactly the same quantities at the time of the Big Bang. It is clear however that we are surrounded by matter, and a legitimate question is "How is it possible to explain this preponderance of matter over antimatter?". It is very unlikely that our matter-dominated corner of the universe is balanced by another corner of the universe dominated by antimatter, as this would have been seen as perturbations in the CMB. Sakharov, in 1967, identified the three mechanisms necessary to obtain a global matter/antimatter asymmetry [8]:

– Baryon and lepton number violation;
– Interactions in the universe out of thermal equilibrium at a given moment of the universe history;

– C- and CP-violation (the rate of a process $i \to f$ can be different from the CP-conjugate process $\tilde{i} \to \tilde{f}$).

The SM includes sources of CP-violation, through the residual phase in the CKM matrix, but they are in no way sufficient to explain the magnitude of the matter-antimatter asymmetry observed.

2.1.4 Free Parameters in the SM

The SM contains no less than nineteen free parameters, which can be taken as:

- 9 fermion masses ($m_e, m_\mu, m_\tau, m_u, m_d, m_c, m_s, m_t, m_b$);
- 3 CKM[1] mixing angles and 1 CP-violating phase;
- 1 electromagnetic coupling constant (g');
- 1 weak coupling constant (g);
- 1 strong coupling constant (g_S);
- 1 QCD vacuum angle;
- 1 vacuum expectation value (v);
- 1 mass for the scalar boson (m_H).

This large number of free parameters, especially in the scalar sector, could be an indication for the existence of a more general and elegant theory than the SM.

2.1.5 Hierarchy Problem

The so-called gauge hierarchy problem [9] is related to the huge energy difference between the weak scale and the Planck scale. The weak scale is given by the vev of the Brout-Englert-Higgs field, which is equal to approximately 246 GeV. Radiative corrections to the scalar boson squared mass, coming from its couplings to fermions and gauge bosons, and from its self-couplings, are quadratically proportional to the ultraviolet momentum cutoff Λ_{UV}, which is at least equal to the energy to which the SM is valid without any addition of new physics. If one considers that the SM is valid up to the Planck mass M_P, the quantum correction to m_H^2 is about thirty orders of magnitude larger than m_H^2, which implies that some extraordinary cancellation of terms should happen. Even if these corrections are absorbed in the renormalization process, some may find uncomfortable with this sensitivity to the details of high scales. This is also known as the naturalness problem of the H boson mass.

In particular, the correction to the squared mass of the scalar boson coming from a fermion f that couples directly to the scalar field ϕ with a coupling λ_f is:

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[1]The Cabibbo-Kobayashi-Maskawa (CKM) matrix contains information about the likelihood of weak decays with flavor changing in charged currents.

$$\Delta m_H^2 = -\frac{|\lambda_f|^2}{8\pi^2}\Lambda_{UV}^2. \quad (2.1)$$

Similarly, some corrections to the mass of the SM scalar boson also arise from scalars. In the case of a scalar particle S with a mass m_S and that couples to the Brout-Englert-Higgs field with a Lagrangian term $-\lambda_S|\phi|^2|S|^2$, the correction to the squared scalar boson mass is:

$$\Delta m_H^2 = \frac{\lambda_S}{16\pi^2}\left[\Lambda_{UV}^2 - 2m_S^2\ln(\Lambda_{UV}/m_S) + ...\right]. \quad (2.2)$$

Again, the correction term to the squared mass is much larger than the squared mass itself. BSM models that avoid this fine-tuning introduce new scalar particles at the TeV scale that couple to the scalar boson, in such a way as to cancel the Λ_{UV}^2 divergence.

Additionally, the large mass differences between fermions, related to Yukawa couplings that can differ by up to six orders of magnitude in the case of the electron and the top quark, constitute the fermion mass hierarchy problem.

2.1.6 Coupling Unification

One of the fundamental questions raised by the SM concerns the particular choice of the $SU(3) \times SU(2) \times U(1)$ symmetry group. Additionally, the three forces included in the SM, the weak, the electromagnetic and the strong forces, are treated separately. The intensity of the forces shows an apparent large disparity around the electroweak scale, but at higher energies their coupling constants tend to have comparable strengths. The electromagnetic and weak forces can be unified in a so-called electroweak interaction, but in the SM, the strong coupling constant does not meet the two other coupling constants at high energies. The running of the coupling constants can be modified by the addition of new particles, such as to reach a grand unification.

2.2 Two-Higgs-doublet Models

Two-Higgs-doublet models (2HDM) [10] are simple extensions of the SM. They introduce two doublets of scalar fields, which, after symmetry breaking lead to five physical states: two charged scalars $H^\pm$, one CP-odd pseudoscalar A and two neutral scalars h and H. Similarly to all models that have extra scalar singlets or doublets relative to the SM, 2HDM satisfy the condition $\rho = m_W/(m_Z\cos\theta_W) = 1$. 2HDM include the minimal supersymmetric extension of the SM (MSSM), which addresses the hierarchy problem and the coupling unification (see Sect. 2.3). Moreover 2HDM allow for the existence of additional CP-violation sources with respect to the SM, which could explain the baryon asymmetry in the universe. Finally, 2HDM are a simple extension of the SM in the scalar sector, and there is no strong motivation against adding an additional scalar doublet to the SM.

2.2.1 Formalism

The most general gauge invariant form of the scalar potential V in 2HDM can be written as [10, 11]:

$$\begin{aligned} V = m_{11}^2\phi_1^\dagger\phi_1 + m_{22}^2\phi_2^\dagger\phi_2 - m_{12}^2\left(\phi_1^\dagger\phi_2 + \phi_2^\dagger\phi_1\right) + \frac{\lambda_1}{2}\left(\phi_1^\dagger\phi_1\right)^2 \\ + \frac{\lambda_2}{2}\left(\phi_2^\dagger\phi_2\right)^2 + \lambda_3\phi_1^\dagger\phi_1\phi_2^\dagger\phi_2 + \lambda_4\phi_1^\dagger\phi_2\phi_2^\dagger\phi_1 \\ + \left\{\frac{\lambda_5}{2}\left(\phi_1^\dagger\phi_2\right)^2 + \left[\lambda_6\left(\phi_1^\dagger\phi_1\right) + \lambda_7\left(\phi_2^\dagger\phi_2\right)\right]\phi_1^\dagger\phi_2 + h.c.\right\}. \end{aligned} \tag{2.3}$$

Under the widely-used assumptions that CP is conserved in the scalar sector and not spontaneously broken, and that all quartic terms odd in either of the doublets are eliminated by discrete symmetries, the expression can be simplified as:

$$\begin{aligned} V = m_{11}^2\phi_1^\dagger\phi_1 + m_{22}^2\phi_2^\dagger\phi_2 - m_{12}^2\left(\phi_1^\dagger\phi_2 + \phi_2^\dagger\phi_1\right) + \frac{\lambda_1}{2}\left(\phi_1^\dagger\phi_1\right)^2 + \frac{\lambda_2}{2}\left(\phi_2^\dagger\phi_2\right)^2 \\ + \lambda_3\phi_1^\dagger\phi_1\phi_2^\dagger\phi_2 + \lambda_4\phi_1^\dagger\phi_2\phi_2^\dagger\phi_1 + \frac{\lambda_5}{2}\left[\left(\phi_1^\dagger\phi_2\right)^2 + \left(\phi_2^\dagger\phi_1\right)^2\right], \end{aligned} \tag{2.4}$$

where all the parameters are real.

The minima of the ϕ_1 and ϕ_2 doublets are respectively $\begin{pmatrix} 0 \\ v_1/\sqrt{2} \end{pmatrix}$ and $\begin{pmatrix} 0 \\ v_2/\sqrt{2} \end{pmatrix}$. The ratio of the vacuum expectation values of the two doublets is written as:

$$\tan\beta = \frac{v_2}{v_1}. \tag{2.5}$$

The squared mass matrix of the neutral scalars can be diagonalized to obtain the physical states h and H; the rotation angle performing the diagonalization is called α. The angle β defined in Eq. (2.5) can also be seen as the angle diagonalizing the squared mass matrices of the charged scalars and of the pseudoscalars (one massive pseudoscalar A and one massless Goldstone boson). Without loss of generality, β can be chosen in the first quadrant, whereas α is either in the first or in the fourth quadrant [12]. The meaning of the angles α and β is illustrated in Fig. 2.1.

The scalar couplings to gauge bosons and fermions in 2HDM can be expressed as a function of the two parameters α and $\tan\beta$. There are four types of 2DHM that do not violate flavor conservation in neutral current interactions; the condition

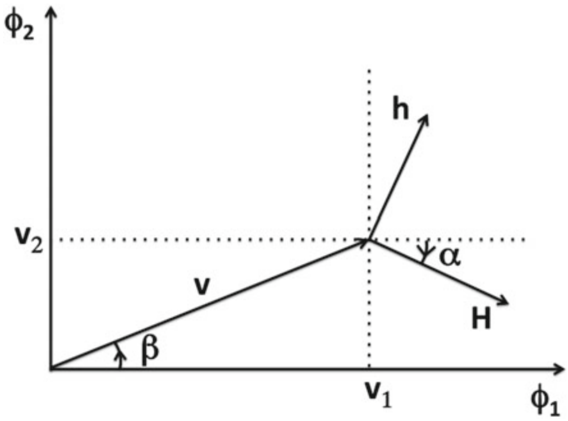

Fig. 2.1 Schematic view of the two angle parameters of 2HDM. The parameter tan β is the ratio between the vacuum expectation values of the two doublets ϕ_1 and ϕ_2, whereas α is the mixing angle between the neutral scalars. The vacuum expectation value $v = 246$ GeV can be decomposed in two components v_1 and v_2 along the doublets ϕ_1 and ϕ_2 respectively. Adapted from [13]

Table 2.1 Scalar doublet to which the leptons ℓ, up-type quarks u and down-type quarks d couple in the different types of 2HDM

	Type-1	Type-2	Type-3 (lepton specific)	Type-4 (flipped)
ℓ	ϕ_2	ϕ_1	ϕ_1	ϕ_2
u	ϕ_2	ϕ_2	ϕ_2	ϕ_2
d	ϕ_2	ϕ_1	ϕ_2	ϕ_1

for avoiding flavor changing neutral currents being that all fermions with the same quantum numbers couple to a same scalar multiplet [14]. A Z_2 symmetry ($\phi_1 \rightarrow +\phi_1$, $\phi_2 \rightarrow -\phi_2$) can be found to ensure that only these interactions exist. The Z_2 symmetry is softly broken if the term m_{12}^2 is non-zero [15]. As shown in Table 2.1, the difference between the four types lies in the doublets to which the charged leptons, up-type quarks and down-type quarks couple. Type-1 is the easiest of them and is the most SM-like: leptons, up-type quarks and down-type quarks all couple to the same doublet, ϕ_1. In type-2, the leptons and down-type quarks couple to ϕ_2, whereas up-type quarks couple to ϕ_1. In the so-called "lepton-specific" type-3, all quarks couple to ϕ_1 and leptons to ϕ_2. Finally in the flipped type-4 model, leptons and up-type quarks couple to ϕ_1 while down-type quarks couple to ϕ_2. In a general way, the intensity of the couplings in the different scenarios are functions of both α and β as presented in Table 2.2.

Table 2.2 Yukawa couplings of vector bosons V, up-type quarks u, down-type quarks d and leptons ℓ to the neutral scalars and pseudoscalar in the four types of 2HDM without flavor changing neutral currents. The Yukawa couplings to the charged scalars can be determined from the couplings of to the neutral pseudoscalar [10]

Particle	Coupling	Type-1	Type-2	Type-3 (lepton specific)	Type-4 (flipped)
h	g_{hVV}	$\sin(\beta-\alpha)$	$\sin(\beta-\alpha)$	$\sin(\beta-\alpha)$	$\sin(\beta-\alpha)$
	$g_{hu\bar{u}}$	$\cos\alpha/\sin\beta$	$\cos\alpha/\sin\beta$	$\cos\alpha/\sin\beta$	$\cos\alpha/\sin\beta$
	$g_{hd\bar{d}}$	$\cos\alpha/\sin\beta$	$-\sin\alpha/\cos\beta$	$\cos\alpha/\sin\beta$	$-\sin\alpha/\cos\beta$
	$g_{h\ell\bar{\ell}}$	$\cos\alpha/\sin\beta$	$-\sin\alpha/\cos\beta$	$-\sin\alpha/\cos\beta$	$\cos\alpha/\sin\beta$
H	g_{HVV}	$\cos(\beta-\alpha)$	$\cos(\beta-\alpha)$	$\cos(\beta-\alpha)$	$\cos(\beta-\alpha)$
	$g_{Hu\bar{u}}$	$\sin\alpha/\sin\beta$	$\sin\alpha/\sin\beta$	$\sin\alpha/\sin\beta$	$\sin\alpha/\sin\beta$
	$g_{Hd\bar{d}}$	$\sin\alpha/\sin\beta$	$\cos\alpha/\cos\beta$	$\sin\alpha/\sin\beta$	$\cos\alpha/\cos\beta$
	$g_{H\ell\bar{\ell}}$	$\sin\alpha/\sin\beta$	$\cos\alpha/\cos\beta$	$\cos\alpha/\cos\beta$	$\sin\alpha/\sin\beta$
A	g_{AVV}	0	0	0	0
	$g_{Au\bar{u}}$	$\cot\beta$	$\cot\beta$	$\cot\beta$	$\cot\beta$
	$g_{Ad\bar{d}}$	$-\cot\beta$	$\tan\beta$	$-\cot\beta$	$\tan\beta$
	$g_{A\ell\bar{\ell}}$	$-\cot\beta$	$\tan\beta$	$\tan\beta$	$-\cot\beta$

2.2.2 Decoupling and Alignment Limits

In the SM, there is one neutral scalar boson, while there are two (h and H) in 2HDM. The lightest neutral scalar of 2HDM is in all generality not identical to the one of the SM, which points to the possibility of determining with property measurements whether the new observed particle belongs to the SM or to 2HDM. As of now, the measured properties of the 125 GeV-state are all compatible with the SM hypothesis within experimental uncertainties. However the 2HDM hypothesis is not ruled out, as there are two important scenarios where the neutral h of 2HDM tends to be SM-like: the decoupling and the alignment limits [16, 17]:

- In the decoupling limit, the mass of the H, A and $H^{\pm}$ all are much larger than the h mass, which causes h to have SM-like couplings. Indeed, if there are two very different mass scales $m_L \ll m_S$ such that $m_h \simeq m_L$ and $m_H, m_{H^{\pm}}, m_A \simeq m_S$, a low mass effective theory can be derived and corresponds to the SM because the m_S-related effects have been integrated out. The decoupling limit implies $\cos(\beta-\alpha) \simeq 0$.

- In the alignment limit, the whole vacuum expectation value (246 GeV) lies in the neutral component of only one of the scalar doublets, and the mixing between the h and H states disappears, which causes one of the neutral mass eigenstates to align with the direction of the scalar field vacuum expectation value and to become SM-like. In this case, the H, A and $H^{\pm}$ particles are not necessarily heavy. The alignment limit corresponds to $\cos(\beta - \alpha) \simeq 0$, and is more general than the decoupling limit.

In Fig. 2.1, the equality $\cos(\beta - \alpha) = 0$, which is satisfied both in the decoupling and the alignment limits, corresponds to the alignment of the state h along the vacuum expectation value v.

2.2.3 *Light Scalars in 2HDM*

In 2HDM, in the alignment limit, one of the neutral scalars (h) can be SM-like, while the pseudoscalar A can be lighter than 125/2 GeV. In the case where the branching fraction of the SM-like scalar to two light pseudoscalars is limited ($\mathcal{B}(h \to AA)$ less than about 0.3), such a topology is still allowed by the limited precision measurements made on the 125-GeV state at the LHC. The branching fraction $\mathcal{B}(h \to AA)$ can be small in the alignment limit when the mass mixing parameter m_{12} has a modest value. Another case where $\mathcal{B}(h \to AA)$ is allowed to take small values with larger m_{12} is when $\sin(\beta + \alpha) \simeq 1$; this relation is compatible with the measured $h \to VV$ signal strength when $\tan\beta$ is large (>5). When $\sin(\beta + \alpha) \simeq 1$, $\sin\alpha$ has to be positive, which, in the type-2 of 2HDM, leads to a so-called "wrong sign" Yukawa coupling of the SM-like h boson to down-type quarks and leptons (see Table 2.2).

The production cross section of light pseudoscalars at the LHC can be large [18]. Figure 2.2 illustrates the viable production cross sections for the gluon-gluon fusion production (ggA) and the production in association with b quarks (bbA) of the pseudoscalar boson A, in type-1 and type-2 of 2HDM. The two scenarios that give small $\mathcal{B}(h \to AA)$ are shown. It can be seen that the largest cross section times branching ratio for A decays to tau leptons is achieved in 2HDM type-2 in the wrong-sign Yukawa coupling scenario.

2.3 Supersymmetry

Supersymmetry (SUSY) **golfand** [19, 20] is a symmetry that relates bosons and fermions. The SUSY operator Q, an anticommuting spinor, transforms a fermionic field F into a bosonic field B and vice-versa:

$$Q|B\rangle = F \text{ and } Q|F\rangle = B. \tag{2.6}$$

A new quantum number R can be introduced to enforce baryon number and lepton number conservation:

$$R = (-1)^{2S+3(B-L)}, \tag{2.7}$$

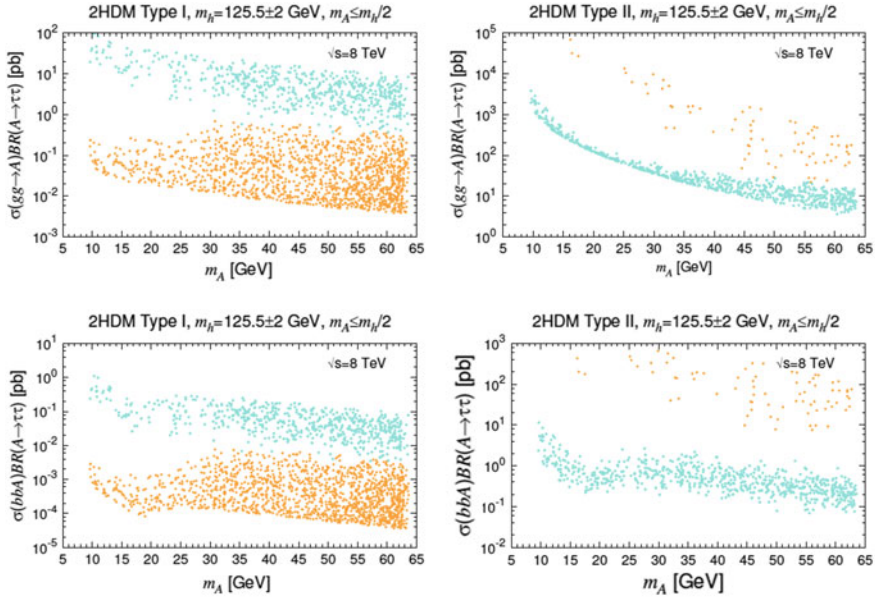

Fig. 2.2 Viable production cross sections for the ggA (top) and bbA (bottom) productions at a center-of-mass energy of 8 TeV, times the branching fraction for A decay to a pair of tau leptons in type-1 (left) and type-2 (right) 2HDM. The cyan points have $\sin(\beta-\alpha)\simeq 1$, $\cos(\beta-\alpha)>0$ and modest m_{12}, whereas orange points have $\sin(\beta+\alpha)\simeq 1$ and $\tan\beta>5$, and correspond to the wrong-sign Yukawa coupling scenario. The largest production cross sections times branching fraction are obtained in 2HDM type-2 with wrong-sign Yukawa couplings. [18]

where S, B and L are respectively the spin, lepton and baryon numbers. With this definition, all SM particles have $R=+1$ and all their superpartners have $R=-1$. R is usually assumed to be conserved in such a way as to forbid fast rates of proton decays. The R-parity conservation implies that supersymmetric particles can only be produced in pairs and that the lightest supersymmetric particle (LSP) is stable. This LSP is an excellent dark matter candidate.

One of the main motivations for the existence of SUSY is the solution to the hierarchy problem. Because fermion loops and boson loops have opposite signs, and SUSY associates a new boson to each fermion and vice-versa, the Λ_{UV}^2 terms in Eqs. 2.1 and 2.2 can exactly cancel for each fermion-scalar pair. Additionally, SUSY can unify the electromagnetic, weak and strong forces below the Planck scale, as illustrated in the case of the minimal supersymmetric extension of the SM (see Sect. 2.3.1) in Fig. 2.3.

Individual particles are grouped in supermultiplets. As the supersymmetric operators Q and $Q^\dagger$ commute with the generators of gauge transformations, particles sharing a same supermultiplet have the same electric charge, weak isospin and color degrees of freedom. In addition, because the same operators also commute with the squared-mass operator $-P^2$, the fermions and bosons in a same supermultiplet

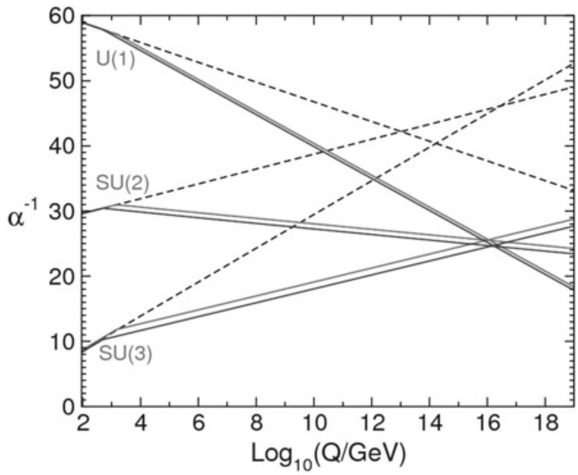

Fig. 2.3 Evolution of the $U(1)_Y$, $SU(2)_L$ and $SU(3)_C$ couplings in the SM (dashed lines), and in two MSSM scenarios (solid lines). Unlike the SM case, the three couplings can be unified at a high energy scale in the MSSM. [19]

should have the same mass. The last point is however known not to be true in reality, because superpartners at the electroweak scale would already have been observed at colliders. If superparticles are heavier than SM particles, which would explain why they have not been discovered yet, SUSY must be broken. For radiative corrections not to exceed typical scalar masses, the SUSY breaking scale should be limited to a few TeV.

2.3.1 Minimal Supersymmetric Extension of the SM (MSSM)

In the minimal supersymmetric extension of the SM (MSSM) [21–23], there exist two types of supermultiplets: vector supermultiplets, where a spin-1 vector boson is associated to a spin-1/2 Weyl fermion, and chiral supermultiplets, where a single Weyl fermion is associated to a complex scalar field. The particle content of the MSSM is shown in Tables 2.3 and 2.4. The superpartners of quarks are called squarks, while the superpartners of leptons are called sleptons. The bino, the neutral wino and the higgsinos mix to form four neutralinos ($\tilde{\chi}_1^0$, $\tilde{\chi}_2^0$, $\tilde{\chi}_3^0$ and $\tilde{\chi}_4^0$), while the winos and charged higgsinos mix to form four charginos ($\tilde{\chi}_1^\pm$ and $\tilde{\chi}_2^\pm$). Two chiral superfields, (H_1^+, H_1^0) and (H_2^0, H_2^-), with hypercharges +1/2 and -1/2 as seen in Table 2.3, are necessary to cancel chiral anomalies.

The MSSM is a particular case of 2HDM type-2. The main specificities in the scalar sector are that, in the MSSM, the mass of the lightest scalar is constrained by some upper bounds, the scalar self-couplings are specified, α and β are not independent from each other, and the decay of the charged scalars $H^\pm$ to a pseudoscalar A and a W boson is kinematically forbidden because $m_{H^\pm} \simeq m_A$ [10].

Table 2.3 MSSM chiral supermultiplets

Particles	Spin-0	Spin-1/2	$(SU(3)_C, SU(2)_L, U(1)_Y)$
Quark, Squark	$(\tilde{u}_{iL}, \tilde{d}_{iL})$	(u_{iL}, d_{iL})	(3,2,+1/6)
	$\tilde{u}^*_{iR}$	$u^\dagger_{iR}$	(3*,1,-2/3)
	$\tilde{d}^*_{iR}$	$d^\dagger_{iR}$	(3*,1,+1/3)
Lepton, Slepton	$(\tilde{e}_{iL}, \tilde{\nu}_{iL})$	(e_{iL}, ν_{iL})	(1,2,-1/2)
	$\tilde{e}^*_{iR}$	$e^\dagger_{iR}$	(1,1,+1)
H, Higgsino	(H_1^+, H_1^0)	$(\tilde{H}_1^+, \tilde{H}_1^-)$	(1,2,+1/2)
	(H_2^0, H_2^-)	$(\tilde{H}_2^0, \tilde{H}_2^-)$	(1,2,-1/2)

Table 2.4 MSSM vector supermultiplets

Particles	Spin-1/2	Spin-1	$(SU(3)_C, SU(2)_L, U(1)_Y)$
gluino, gluon	$\tilde{g}$	g	(8,1,0)
wino, W boson	$\tilde{W}^\pm, \tilde{W}^0$	$W^\pm, W^0$	(1,3,0)
bino, B boson	$\tilde{B}^0$	B^0	(1,1,0)

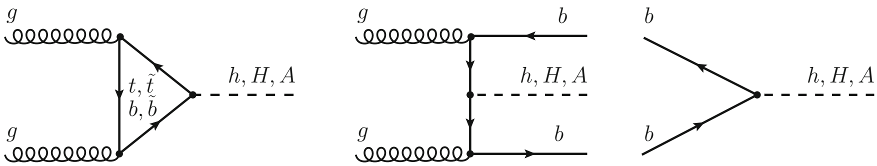

Fig. 2.4 Feynman diagrams for the production of neutral scalars in the MSSM, in gluon-gluon fusion (left), and production with b quarks (center and right). [24]

2.3.2 *Scalar Sector in the MSSM*

In the MSSM, neutral scalars $\Phi = H/A/h$ can be produced by two mechanisms: gluon-gluon fusion ($gg\Phi$) and production with b quarks ($bb\Phi$). Characteristic Feynman diagrams for such processes are shown in Fig. 2.4, where the $bb\Phi$ mechanism is shown in two different schemes of proton parton distribution functions. The $bb\Phi$ production cross section is increased at large $\tan\beta$ because of the enhanced Yukawa couplings to down-type fermions.

At tree level in the MSSM, the only two free parameters can be taken as the mass of the pseudoscalar, m_A, and $\tan\beta$. The masses of the neutral scalars and of the charged scalars, as well as the angle α can be expressed as [25]:

$$m^2_{h/H} = \frac{1}{2}\left(m_A^2 + m_Z^2 \mp \sqrt{(m_A^2 + m_Z^2)^2 - 4m_A^2 m_Z^2 \cos^2 2\beta}\right), \tag{2.8}$$

$$m^2_{H^\pm} = m_A^2 + m_W^2, \text{ and} \tag{2.9}$$

$$\tan 2\alpha = \tan 2\beta \left(\frac{m_A^2 + m_Z^2}{m_A^2 - m_Z^2} \right) \text{ with } -\frac{\pi}{2} \leq \alpha \leq 0. \tag{2.10}$$

This leads to:

$$m_h \leq \min(m_A, m_Z) \times |\cos 2\beta| \leq m_Z. \tag{2.11}$$

The mass of the lightest neutral scalar is thus inferior to the Z boson mass, which is excluded by LEP bounds and does not correspond to the observation of the 125-GeV scalar at the LHC. Fortunately, radiative corrections above tree level, essentially loop corrections due to top and stop quarks, enable the h scalar to be as heavy as approximately 135 GeV. In the case where m_A is much larger than the Z boson mass, the relations above give:

$$m_H \simeq m_{H^\pm} \simeq m_A \text{ and } \alpha \simeq \beta - \frac{\pi}{2}, \tag{2.12}$$

which is the decoupling limit as seen in Sect. 2.2.2.

The phenomenology of the scalar sector of the MSSM can be described by two parameters: the mass of the pseudoscalar m_A, and $\tan\beta$. It is generally assumed that $\tan\beta$ lies approximately between 1 and 60, where 60 is the ratio between the top quark mass and the bottom quark mass. Above tree level, more parameters appear and some benchmark scenarios fixing these parameters can be studied. It has been shown however that, taking into account the mass measured for the h boson, the MSSM can be approximately reparameterized as a function of m_A and $\tan\beta$, in the so-called hMSSM [26]. One can distinguish three regions of the parameter space, where the search strategies will differ: the low m_A, the high $\tan\beta$ and the low $\tan\beta$ regions. It has been shown in [26] that the full parameter space of the MSSM could be almost entirely covered in the search for additional scalars at the LHC at 14 TeV with a luminosity of 300 fb^{-1}, while a good part of the parameter space has already been explored in LHC searches at 7 and 8 TeV, as shown in Fig. 2.5.

Low m_A Region

At low m_A, the most powerful channel to search for an MSSM scalar sector is clearly $H^+ \to \tau^+\nu_\tau$ (and its charge conjugate decay). The limits shown in Fig. 2.5 correspond to the $t \to H^+b$ production, for charged scalar masses below the difference of the top quark and bottom quark masses.

High $\tan\beta$ Region

In the region of the parameter space where $\tan\beta$ is large, say $\tan\beta > 5$, the most sensitive final state to search for new heavy resonances Φ is by far $\Phi = A/H/h \to \tau\tau$. The reason for this is that the couplings to leptons and down-type quarks are enhanced with increased $\tan\beta$, because these particles couple to the second scalar doublet (see Table 2.1). In addition, for the same reason, the production cross section for the Φ resonance in association with bottom quarks is also enhanced at large $\tan\beta$. Even if the decay branching fraction of the resonance to bottom quarks remains larger (approximately nine times higher), the experimental difficulties, such as the

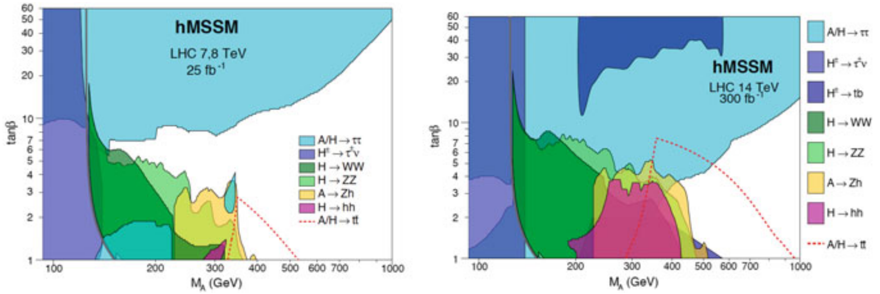

Fig. 2.5 Sensitivity of MSSM scalar searches at the LHC at 7 and 8 TeV in LHC Run-1 (left) and projection with 300 fb^{-1} of 14 TeV data collected at the LHC (right), in the context of the hMSSM parameterization. The $A \to t\bar{t}$ search (dashed red line) has not yet been performed at the LHC, and the sensitivity is predicted. The exclusion limit of the $A \to \tau\tau$ analysis around $m_A = 350$ GeV and $2 < \tan\beta < 4$ is due to the strong increase of the $gg \to A$ cross section at the $t\bar{t}$ threshold, coupled to the suppression of $A \to Zh$ decays and enhanced couplings to down-type quarks and leptons because $\tan\beta > 1$. The hMSSM scenario takes into account the mass measured for the new SM-like scalar. [26]

distinction between b jets and other flavor jets, make the channel $\Phi \to bb$ less sensitive. Finally, the channel $\Phi \to \mu\mu$ also has some potential, but is hurt by its low decay branching fraction: $\mathcal{B}(\Phi \to \mu\mu) \simeq \mathcal{B}(\Phi \to \tau\tau) \times m_\mu^2/m_\tau^2$.

Low $\tan\beta$ Region

The phenomenology in the low $\tan\beta$ region is much richer than in the high $\tan\beta$ region. Experimentally, one interesting decay of the A pseudoscalar to study is $A \to Zh$, in the intermediate mass range $m_Z + m_h < m_A < 2m_t$, as seen in Fig. 2.6. If the Z boson decays leptonically, it is possible to achieve a good background reduction, and the most favorable h decays in terms of branching fractions, $h \to bb$ and $h \to \tau\tau$, can be targeted. At higher m_A, the $A \to t\bar{t}$ channel opens, but, due to interference effects with the SM backgrounds, it is experimentally a difficult channel. The dominant H decay channel in the intermediate mass range $2m_h < m_H < 2m_t$ is $H \to hh$, whereas there are also non negligible contributions from $H \to WW$ and $H \to ZZ$. Outside of the low m_A region described previously, for $m_{H^\pm} > m_t + m_b$, the charged scalars $H^\pm$ almost exclusively decay to a top and a b quarks.

2.4 Two-Higgs-doublet Models + a Singlet

2.4.1 Introduction to 2HDM+S

The extensions of 2HDM where a complex scalar singlet is added to the already present scalar doublets, are called 2HDM+S. Because of the additional singlet, two new bosons are introduced. The next–to-minimal supersymmetric extension of the

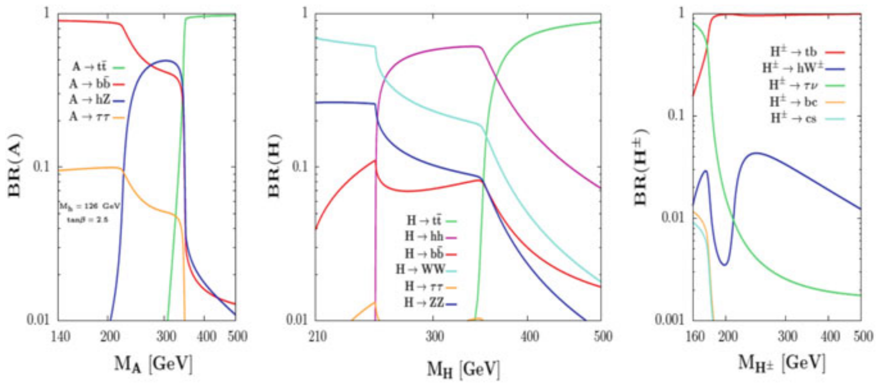

Fig. 2.6 Branching fractions of the A, H and $H^{\pm}$ scalars in the MSSM as a function of their masses, for $\tan\beta = 2.5$ and $m_h = 126$ GeV. [27]

SM (NMSSM) [28, 29] (for a review, see [30]) is a case of 2HDM+S type-2 and is the easiest extension of the MSSM. The supersymmetric potential in the MSSM contains a mass parameter μ in the expression $\mu\phi_1\phi_2$, which has to be at the SUSY breaking scale (m_{SUSY}). The fact that μ is at a scale well below the Planck scale without any theoretical reason, constitutes the so-called μ problem [31]. This problem disappears in the NMSSM, where an effective mass is generated via a coupling to the complex scalar field associated to the new singlet; this is a strong motivation for the existence of the NMSSM over the MSSM. Another motivation comes from the fact that new scalar particles contribute to the mass of the scalar boson h in the NMSSM, which removes the tensions in the MSSM originating from the large measured mass of the new particle.

2.4.2 Exotic Decays of the 125-GeV Particle

In 2HDM+S, the h boson, identified as the 125-GeV particle discovered in 2012, can decay exotically to non-SM particles. Even though decays of the h boson of 2HDM to non-SM particles are theoretically allowed, the 2HDM parameter space is by now extremely constrained by LHC searches. However, in 2HDM+S, a wide range of exotic decays of the 125-GeV boson is still allowed after the Run-1 of the LHC. The singlet added to 2HDM does not have Yukawa couplings of its own, and only couples to ϕ_1 and ϕ_2 in the potential, from which it inherits its couplings to SM fermions. To keep the scalar h SM-like, the mixing with the singlet S needs to be small. The imaginary part of the singlet gives rise to a pseudoscalar a (after a small mixing with the pseudoscalar A), and the real part to a scalar s (after mixing with H and h). Exotic decays of the type $h \to aa$, $h \to ss$ or $h \to Za$ are then possible.

In the pseudoscalar case, the light pseudoscalar a, mostly singlet-like, inherits its couplings to fermions from the heavy pseudoscalar A. As in the case of general 2HDM, the four types of 2HDM+S lead to different scenarios and give rise to many exploitable signatures for exotic h decays at the LHC:

- Type-1: All fermions couple to ϕ_1, and therefore the branching fractions of the pseudoscalar are proportional to those in the SM, without any dependence on $\tan\beta$.
- Type-2: Leptons and down-type quarks couple to the same doublet, as in the MSSM. For values of $\tan\beta$ larger than unity, the branching fractions of the pseudoscalar to leptons and down-type quarks are enhanced, which makes $h \to aa \to bbbb$, $h \to aa \to \mu\mu bb$ and $h \to aa \to \tau\tau bb$ interesting channels to search for exotic h decays.
- Type-3 (lepton-specific): Leptons couple to the ϕ_2 doublet contrarily to quarks, which means that the branching fractions for pseudoscalar decays to leptons increase for large values of $\tan\beta$. In this scenario with large $\tan\beta$, $h \to aa \to \mu\mu\tau\tau$ and $h \to aa \to \tau\tau\tau\tau$ are favoured channels when kinematically allowed.
- Type-4 (flipped): Pseudoscalar decays to leptons and up-type quarks are enhanced with respect to down-type quarks when $\tan\beta > 1$. In this case, $h \to aa \to \tau\tau cc$ and $h \to aa \to \tau\tau bb$ can be interesting channels.

2.5 Search for BSM Physics in the Scalar Sector

The scalar sector is a favored place to look for new physics, because it is described much less elegantly than the other parts of the SM as most free parameters of the theory are related to the scalar interaction. The existence of a Higgs-portal [32], where the scalar sector is the only one to interact with BSM physics, is strongly motivated. Complementary ways exist to point to the existence of BSM physics in the scalar sector:

1. Precision measurements of the properties of the 125-GeV scalar boson, that would reveal deviations from the SM;
2. Direct discovery of new scalar particles;
3. Discovery of BSM decays of the 125-GeV scalar boson;
4. Observation of BSM physics in signatures involving scalar bosons.

A review about the complementarity between precision measurements and direct searches in the MSSM can be found in [33]. The four points are detailed below:

- **Precision measurements:** The properties of the lightest scalar h of 2HDM can deviate from the properties of the SM scalar boson; precision measurements of the 125-GeV state therefore should make the distinction between 2HDM and SM possible. However, as seen in Sect. 2.2.2, many extended sector scalar models have a decoupling or an alignment limit, which makes the properties of the h boson of 2HDM very close to those of the SM scalar boson. Figure 2.7 illustrates the dependence of the production cross section of the MSSM h boson as a function

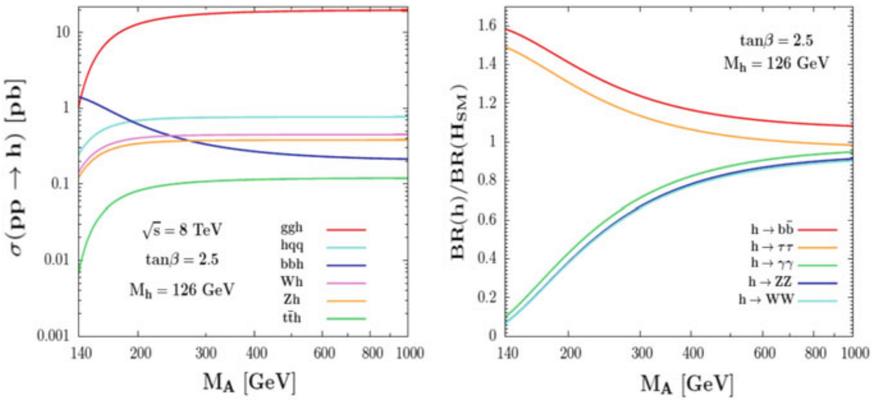

Fig. 2.7 Left: Production cross sections of the MSSM h boson at 8 TeV, for $\tan\beta = 2.5$. Right: Evolution of the decay branching ratios of the h scalars in the MSSM and SM as a function of the pseudoscalar mass m_A, for $\tan\beta = 2.5$. [27]

of the mass of the pseudoscalar A for a given $\tan\beta$, as well as the ratio between the decay branching fractions of the MSSM h and the SM scalar boson. As expected according to the decoupling limit, the branching fractions tend to be very similar in the two scenarios when the mass of the pseudoscalar increases, and a great precision is needed to highlight deviation from the SM. The measurement of the properties of the discovered 125 GeV-boson in the decay channel to tau leptons is presented in Chap. 7.

– **Direct discovery of new scalar particles:** Discovering new scalar particles would be a direct evidence of BSM physics. Many searches for extra scalars, in the context of general 2HDM or MSSM for example, are performed at the LHC. The search for a heavy neutral scalar decaying to a pair of tau-leptons is described in Chap. 13, while the search for the heavy pseudoscalar A of the MSSM, decaying to a Z and a SM-like h bosons is described in Chap. 10. Light bosons with a mass lighter than the Z boson could also be discovered at the LHC; Chap. 12 details the search for such a particle in its decays to tau leptons. Many other searches exist at the LHC but are not described in this thesis, such as the search for charged scalars.
– **Discovery of BSM decays of the 125-GeV scalar boson**: Motivations for the existence of exotic decays of the 125-GeV boson to non-SM particles are various [34, 35]. First, the SM scalar boson has an extremely narrow width ($\Gamma \simeq 4.07$ MeV) compared to its mass, because of the suppression of tree-level Yukawa couplings. Given that the coupling to two b quarks has the small value of approximately 0.02, the coupling, even small, to another light state could open non negligible decay modes. Second, the scalar sector could be a portal to new physics, which allows SM matter to interact with a hidden-sector matter. If there exists a Higgs portal, the hidden-sector matter does not have to be charged under SM forces. And finally, exotic scalar decays are a relatively simple extension of the SM, and are still allowed after all the measurements made during LHC Run-1. Indeed an upper

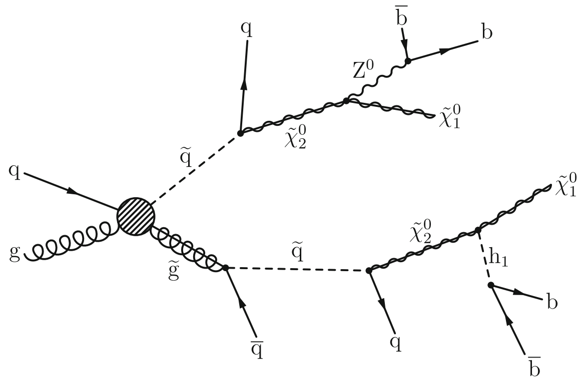

Fig. 2.8 Feynman diagram of quark-gluino production with subsequent cascade decays via neutralinos into the h_1 boson of the NMSSM. [37]

limit on the branching fraction of the 125-GeV boson to BSM particles can be set experimentally and, as of today, this upper limit leaves a large room for exotic decays. In particular, CMS measured $\mathcal{B}(H \to \text{BSM}) < 30$ at 95% CL, using all data collected during LHC Run-1 [36]. Projections for future LHC runs give a final precision on $\mathcal{B}(H \to \text{BSM})$ of the order of 10%, which still allows for non negligible decays of the 125-GeV particle. The variety of possible BSM decays is extremely large. A group of searches among others explores the possibility for the h boson to decay to invisible particles, resulting in missing transverse energy in the detector. A search for exotic decays with SM particles in the final state is presented in Chap. 11.

– **Observation of BSM physics in signatures involving scalar bosons:** Scalar bosons could be produced in some BSM physics processes. An example of such a process is a squark-gluino production with subsequent cascade decays via neutralinos into the h_1 boson of the NMSSM; the corresponding Feynman diagram is shown in Fig. 2.8. This method to look for an extended scalar sector is not explored in this thesis.

2.6 Precision Measurements Versus Direct Discovery in the Case of the SM+S

To compare the reach of precision measurements and direct discovery of new scalars, one can take the simple case where a real scalar singlet S is added to the usual scalar doublet ϕ of the SM [38]. After symmetry breaking, the singlet mixes with the doublet with a mixing angle α, resulting in two neutral mass eigenstates H_1 and H_2:

$$H_1 = \phi \cos\alpha + S \sin\alpha, \tag{2.13}$$

$$H_2 = -\phi \sin\alpha + S \cos\alpha, \tag{2.14}$$

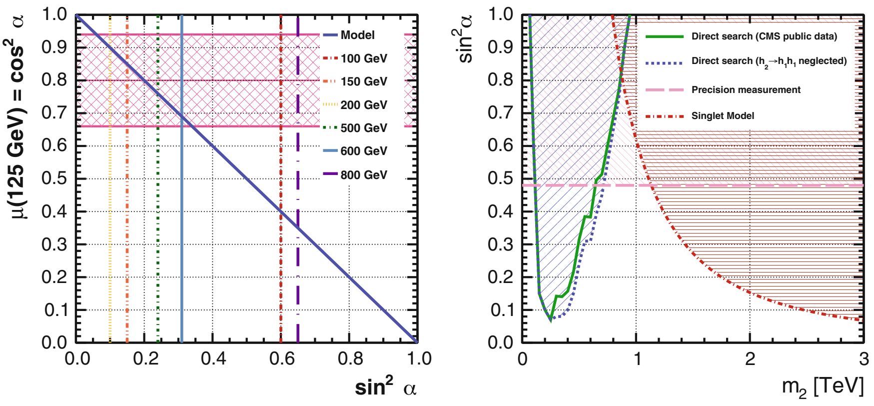

Fig. 2.9 Left: Comparison of the reach of the precision measurement of the signal strength of the 125 GeV state at 1σ and of the search for extra scalars at 2σ for different masses (100, 150, 200, 500, 600 and 800 GeV). The blue line corresponds to the SM+S benchmark: $\cos^2\alpha + \sin^2\alpha = 1$. Right: Comparisons of the constraints on the parameter $\sin^2\alpha$ set by direct searches for extra scalar states (blue curve if the decay $H_2 \to H_1 H_1$ is neglected, green curve otherwise), by indirect constraints on the measurement of the strength parameter of the 125 GeV state with a partial set of data collected by CMS during LHC Run-1, and by perturbative unitarity conditions, in the case of the SM+S model. Uncertainties are given at 2σ level. [38]

where H_1 has a mass m_1 = 125 GeV and H_2 a mass m_2 lower or greater than 125 GeV. The H_1 and H_2 states interact only via their ϕ component. The case where the new particle interacts only through mixing is studied here. While the decay modes and branching fractions of the new states are the same as predicted for the scalar boson in the SM at their respective masses m_1 and m_2, their production rates are however scaled by factors $\cos^2\alpha$ in the case of H_1 and $\sin^2\alpha$ in the case of H_2. The decay $H_2 \to H_1 H_1$ should be taken into account when kinematically allowed.

The SM+S benchmark model can be represented by the equality $\cos^2\alpha + \sin^2\alpha = 1$, as shown in Fig. 2.9 (left) with the continuous blue line. Precision measurements of the signal strength of the 125-GeV state give some constraints on the parameter $\cos^2\alpha$, which scales the H_1 production cross section; the measured signal strength, when lower than one, is equivalent to the measured $\cos^2\alpha$. The best-fit signal strength was measured to be $\hat{\mu} = 0.80 \pm 0.14$ (where the uncertainties correspond to one standard deviation) in 2013 [39]; this value is represented together with its uncertainty by the horizontal pink lines in Fig. 2.9 (left). In a first approximation, the uncertainty band can be doubled to correspond to uncertainties at the level of two standard deviations. On the other hand, searches for extra scalars with masses different from 125 GeV can set constraints on $\sin^2\alpha$; the upper limit on the signal strength at a given mass, if lower than one, is equivalent to an upper limit on $\sin^2\alpha$ at that same mass. Limits from searches for additional scalars with masses between 100 and 800 GeV [39–42] are shown in Fig. 2.9 (left) with the vertical dashed lines. This makes it possible to compare the reach of the two approaches, using the CMS data collected during Run-1 and presented at the HCP12 conference. It can be seen in the figure that in this particular benchmark model, in most part of the accessible mass range, namely between 125 and 600 GeV, the direct detection is a more powerful approach than precision measurements.

The limit on $\cos^2\alpha$ obtained by the precision measurements on the signal strength of the 125-GeV state can be converted to a limit on $\sin^2\alpha$ for all m_2. In the right part of Fig. 2.9, the upper limits on $\sin^2\alpha$ are superimposed to the constraint from the precision measurement (uncertainty at two standard deviations level), for all masses m_2 probed at the LHC. Also superimposed is the limit set by tree level unitarity constraints in SM+S, which play a role at large m_2.

2.7 Chapter Summary

The SM is known not to answer a series of fundamental questions, such as the hierarchy problem or the existence of dark matter. Many BSM models that address some of the SM issues predict the existence of more than one scalar particle. This is the case of the minimal supersymmetric extension of the SM (MSSM): it brings a solution to the hierarchy problem and the coupling unification, proposes a dark matter candidate, and introduces in total five scalar bosons. The MSSM is part of a more generic class of models, two-Higgs-doublet models (2HDM), which have five scalar particles and give rise to a large variety of phenomenological signatures. Three complementary manners to uncover an eventually exotic scalar sector are explored in this thesis:

– Precision measurements of the properties of the discovered boson, which may highlight deviations from the SM;
– Direct search for new scalar particles in specific models;
– Search for exotic decays of the discovered boson.

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