

The mystery of the Higgs particle/Le mystère de la particule de Higgs

Higgs boson properties in the Standard Model and its supersymmetric extensions

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Abstract

We review the realization of the Brout–Englert–Higgs mechanism in the electroweak theory and describe the experimental and theoretical constraints on the mass of the single Higgs boson expected in the minimal Standard Model. We also discuss the couplings of this Higgs boson and its possible decay modes as functions of its unknown mass. We then review the structure of the Higgs sector in the minimal supersymmetric extension of the Standard Model (MSSM), noting the importance of loop corrections to the masses of its five physical Higgs bosons. Finally, we discuss some non-minimal models. *To cite this article: J. Ellis et al., C. R. Physique 8 (2007).*

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Résumé

Propriétés du boson de Higgs dans le Modèle Standard et ses extensions supersymétriques. Nous examinons le mécanisme de Brout–Englert–Higgs et décrivons les limites expérimentales et théoriques de la masse de l’unique boson de Higgs attendu dans le Modèle Standard minimal. Nous discutons également les couplages de ce boson de Higgs et ses modes de désintégration en fonction de sa masse inconnue. Nous examinons ensuite la structure du secteur de Higgs dans l’extension supersymétrique minimale du Modèle Standard (MSSM), en soulignant l’importance des corrections induites par les boucles aux masses de ses cinq bosons de Higgs physiques. Enfin, nous examinons quelques modèles non-minimaux. *Pour citer cet article : J. Ellis et al., C. R. Physique 8 (2007).*

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1. Introduction

The key properties of any elementary particle are its mass, its spin and its couplings to other particles. Owing to its very specific theoretical rôle [1–3], the single Higgs boson [3] of the Standard Model of strong and electroweak interactions [4] has some very characteristic properties. In particular, its spin is zero, unlike any other elementary

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particle ever observed, and its couplings to other particles are proportional to their masses. The mass dependence of its couplings is related to the special rôle that this scalar field has in generating the masses of the other elementary particles.

On the other hand, the mass of the Higgs boson is largely unconstrained in the Standard Model, although there are upper bounds derived from unitarity [5]. More recently, significant lower limits have been provided by direct experimental searches at LEP [6], and precision electroweak data seem to prefer a specific portion of the mass range allowed by these upper and lower limits [7]. Moreover, if one requires the Standard Model to remain valid up to high energies, there are upper and lower limits [8] on the possible mass of the Higgs boson. Nevertheless, within the Standard Model the mass of the Higgs boson remains an unknown quantity, though it may be more constrained in certain extensions such as supersymmetry. As a result, knowing the couplings of the Standard Model Higgs boson enables one to predict its decay properties only as functions of its unknown mass.

Later we use the standard field-theoretical point of view to discuss the properties of the Standard Model Higgs boson. However, they can be understood qualitatively using very simple arguments, without appealing to the details of its field-theoretical formulation. In order to be compatible with unitarity, the cross sections for particle scattering are bounded, and cannot grow without limit. Moreover, a field theory is capable of making many predictions, over a wide range of energy scales and in terms of a finite number of input parameters, only if it is renormalizable, i.e., if the divergences encountered in calculating quantum loop diagrams can be absorbed in the definitions of the input parameters. The requirement of renormalizability imposes further restrictions on scattering cross sections.

The scattering cross sections in the Standard Model are bounded suitably at high energies only if the Higgs boson is included, and this requirement determines its properties uniquely [9]. For example, if one considers the annihilation of a fermion–antifermion pair into a pair of W^\pm bosons, it is well known that the high-energy behaviour of the cross-section is largely tamed by cancellations between direct-channel photon- and Z -exchange diagrams with the crossed-channel neutrino-exchange diagram, an effect measured at LEP [10]. However, there is in principle a residual divergence proportional to the product of the fermion and W masses, which occurs only in the direct-channel spin-0 partial wave.¹ The only way to remove this residual divergence is to include the direct-channel exchange of a spin-0 particle whose couplings to fermions and W bosons are determined by their masses [9]. Similar divergences would occur in $W^+W^- \rightarrow W^+W^-$ scattering, a process that will be observable at the LHC. Again, the only way to tame this divergence within the Standard Model is to include the exchange of a scalar Higgs boson with a coupling to the W boson that is proportional to its mass [9].

The Higgs sector of the Standard Model that is described in the next Section of this article is the simplest field-theoretical realization of these cancellations, and it is known to be part of a renormalizable theory [11]. After formulating the theory, we then discuss the couplings and decay modes of the single Higgs boson of the Standard Model [12,13], as well as theoretical, phenomenological, experimental and cosmological bounds on its possible mass. The following Section extends this discussion to supersymmetric extensions of the Standard Model at the Fermi scale of weak interactions, as motivated by requiring [14] the gauge hierarchy to be natural [15]. We discuss supersymmetric extensions of the Standard Model, mainly the minimal one (MSSM), but also some non-minimal versions.

There are five physical Higgs bosons in the MSSM, three of them neutral and two charged. All of them have spin 0, and the couplings of the neutral ones obey sum rules that reflect how they share the mass-generating task of the single Higgs boson of the Standard Model, and the subtle interplay of gauge symmetry and supersymmetry [13]. The masses and couplings of these Higgs bosons are calculable in terms of underlying model parameters, and receive important quantum corrections [16–18]. Nevertheless, at least in the MSSM, the mass of the lightest Higgs boson is tightly circumscribed to be less than about 130 GeV, and its couplings are generally predictable. As a result, in quite a large part of the MSSM parameter space, its properties are expected to be quite similar to those of a light Standard Model Higgs boson. If a light Higgs boson is found at the LHC, it will be a challenge to distinguish between the Standard Model and its minimal supersymmetric extension, unless one or more heavier Higgs bosons or some supersymmetric particles are also found. On the other hand, if no light Higgs boson is found at the LHC, it would be difficult to maintain faith in supersymmetry at the weak scale as a solution to the gauge-hierarchy problem. Because of space constraints, the present article is a concise review. For more extensive recent reviews on the same subject, see e.g. [19].

¹ This was numerically insignificant at LEP, because of the very small electron mass. However, it would be a very large effect in $\bar{\tau}\tau \rightarrow W^+W^-$ annihilation.

2. Standard Model

2.1. General properties

In this section, we review the Brout–Englert–Higgs mechanism [2,3] in the minimal version of the Standard Model (SM from now on). An SU(2)-doublet scalar field ϕ is introduced, which is allowed to transform inhomogeneously under the action of the SU(2) factor of the SM gauge group:

$$\phi = \begin{pmatrix} \phi_1 \\ \phi_2 \end{pmatrix}, \quad \phi \rightarrow e^{ig\xi^a \frac{\sigma^a}{2}} \left(\phi + \frac{v}{\sqrt{2}} \right) - \frac{v}{\sqrt{2}}, \quad v = \begin{pmatrix} v_1 \\ v_2 \end{pmatrix} \quad (1)$$

where g is the SU(2) coupling constant, ξ^a are real transformation parameters, σ^a ($a = 1, 2, 3$) are the Pauli matrices, and v_1 and v_2 are scalar constants. Such an inhomogeneous transformation rule is possible only for a scalar field ϕ : otherwise, the values of v_1, v_2 would depend on the choice of reference frame. The value of the hypercharge of ϕ is fixed by the requirement that it transform homogeneously under ordinary electric charge transformations that correspond to the subgroup U(1)_{em}. This amounts to requiring

$$e^{ie\alpha Q} \frac{1}{\sqrt{2}} \begin{pmatrix} v_1 \\ v_2 \end{pmatrix} - \frac{1}{\sqrt{2}} \begin{pmatrix} v_1 \\ v_2 \end{pmatrix} = 0 \quad (2)$$

or equivalently

$$\begin{pmatrix} Q_1 & 0 \\ 0 & Q_2 \end{pmatrix} \begin{pmatrix} v_1 \\ v_2 \end{pmatrix} = \begin{pmatrix} 1/2 + Y/2 & 0 \\ 0 & -1/2 + Y/2 \end{pmatrix} \begin{pmatrix} v_1 \\ v_2 \end{pmatrix} = \begin{pmatrix} 0 \\ 0 \end{pmatrix} \quad (3)$$

where $Q_{1,2}$ are the electric charges of $\phi_{1,2}$ in units of the proton charge e , and we have used $Q = T_3 + Y/2$. The system in Eq. (3) has non-vanishing solutions only for $Y = \pm 1$. Without loss of generality, we shall adopt the first choice, $Y = +1$, which gives $Q_1 = 1, Q_2 = 0$. The solution of Eq. (3) is in this case $v_1 = 0, v_2 \equiv v$, and we can also assume without loss of generality that v is real and positive.

The covariant derivative

$$D_\mu \left(\phi + \frac{v}{\sqrt{2}} \right) = \left(\partial_\mu - ig \frac{\sigma^a}{2} W_\mu^a - ig' \frac{1}{2} B_\mu \right) \left(\phi + \frac{v}{\sqrt{2}} \right) \quad (4)$$

transforms as an ordinary doublet of SU(2), and mass terms for the gauge vector boson fields arise in the term $|D_\mu(\phi + v/\sqrt{2})|^2$, with $m_W^2 = g^2 v^2/4, m_Z^2 = (g^2 + g'^2)v^2/4$, whereas the photon is massless.

The form of the scalar potential is uniquely determined by renormalizability and gauge invariance, together with the requirement that $\phi = 0$ correspond to a minimum:

$$V(\phi) = \lambda \left[\left(\phi + \frac{v}{\sqrt{2}} \right)^\dagger \left(\phi + \frac{v}{\sqrt{2}} \right) - \frac{v^2}{2} \right]^2 \quad (5)$$

Three of the four scalar degrees of freedom in the doublet ϕ are unphysical, because of the gauge symmetry; they do not appear in S -matrix elements as asymptotic states. Their presence as intermediate states is needed in order to cancel unphysical singularities in the gauge boson propagators, but they may be removed from the spectrum by choosing a unitary gauge. The fourth degree of freedom is instead a physical one; we shall denote the corresponding real scalar field by H . It corresponds to a massive, spinless neutral particle, the Higgs boson [3], with squared mass $m_H^2 = 2\lambda v^2$.

The two constants v and λ are fundamental parameters of the theory, whose values must be extracted from experiment. The value $v \simeq 250$ GeV can be obtained from the measured value of the Fermi constant, $G_F/\sqrt{2} = g^2/(8m_W^2) = 1/(2v^2)$, whereas the value of λ (and hence of m_H) is still unknown.

Fermion mass terms arise from Yukawa interactions with the Higgs doublet. We do not reproduce here the details of this mechanism, but we do recall that, whereas the Yukawa couplings of the SM fermions to the Higgs doublet are in general non-diagonal in the space of electroweak doublets of fermions, as are the corresponding couplings of the W^\pm gauge bosons, the couplings of the neutral Z boson and the single physical Higgs boson of the SM are diagonal in flavour space at the tree level, and flavour-changing phenomena induced by neutral currents are strongly suppressed in loop amplitudes [20], as required by experiment. Violation of CP invariance can also be implemented with three or more fermion families [21], using just one Higgs doublet [22].

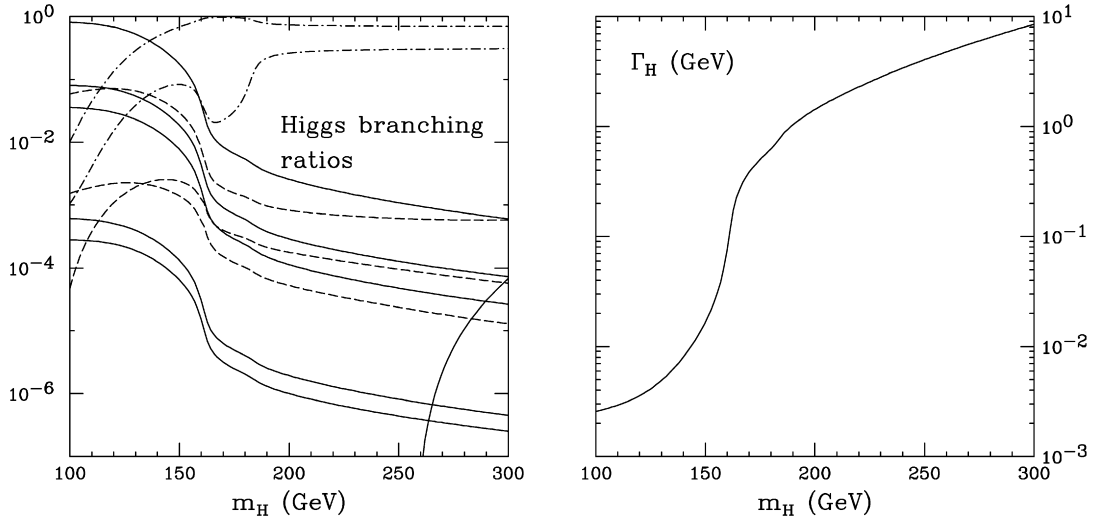


Fig. 1. Left: branching ratios of the different Higgs decay modes. Solid: $H \rightarrow f\bar{f}$ ($f = t, b, c, \tau, s, \mu$); dot-dashed: WW, ZZ ; dashed: $gg, \gamma\gamma, Z\gamma$. Right: total width of the Higgs boson as a function of its mass. This figure was obtained with the program HDECAY [25].

The couplings and self-couplings of the Higgs boson are contained in the following interaction terms of the SM Lagrangian:

$$\mathcal{L}_{\text{Higgs}} = \left(m_W^2 W^{\mu+} W_{\mu}^- + \frac{1}{2} m_Z^2 Z^{\mu} Z_{\mu} \right) \left(\frac{H^2}{v^2} + \frac{2H}{v} \right) - \frac{H}{v} \sum_f m^f \bar{\psi}^f \psi^f - \lambda v H^3 - \frac{1}{4} \lambda H^4 \quad (6)$$

where the ψ^f are Dirac fields for massive fermions, and the m^f are the corresponding masses. We see that the theory contains couplings of the Higgs boson to the W, Z bosons that are proportional to $m_V^2/v \sim gm_V$ ($V = W, Z$), and Yukawa couplings to fermions that are proportional to $m^f/v \sim gm^f/m_V$, as argued earlier on the basis of unitarity and renormalizability.

It follows that, for $m_H < 2m_W$, the Higgs boson decays mainly into pairs of heavy fermions ($b\bar{b}$ or $\tau^+\tau^-$), while for higher values of m_H , the W^+W^- and ZZ channels open and become dominant [5]. At even larger values of m_H , the $t\bar{t}$ mode is also available, though it never becomes as important as the decays into pairs of gauge bosons. The rare decay into a pair of photons [12] is also potentially important for detection at the LHC, where the coupling to gluon pairs [23] may be important for Higgs production [24]. Both these couplings are induced at the one-loop level by loops of heavy electromagnetically or colour charged particles. Below the W^+W^- threshold, the gg branching ratio is comparable to those for the $\tau^+\tau^-$ and $c\bar{c}$ channels. The branching ratios for the various decay modes are displayed in the left panel of Fig. 1, as functions of the Higgs mass. The total width of the Higgs boson is shown in the right panel of Fig. 1 as a function of the Higgs mass. It is below 1 GeV for $m_H \lesssim 200$ GeV, but becomes comparable to m_H for $m_H \sim 1$ TeV.

2.2. Theoretical bounds

In principle, the self-interaction coupling constant λ of the Higgs scalar can take any (real and positive) value; thus, the mass of the physical Higgs boson is not fixed by the theory. However, as anticipated in the Introduction, some information about the value of m_H can be obtained by means of theoretical considerations.

Perturbative unitarity [5] provides an upper bound on the Higgs boson mass, much in the same way as in the case of the W boson. The process to be considered is the elastic scattering of weak vector bosons with longitudinal polarizations; the unitarity bound on the leading-order amplitude is respected, including the contribution of Higgs boson intermediate states, provided the value of the Higgs mass does not exceed 800–1000 GeV [5]. A similar bound can be obtained by considering Higgs decays, whose rates for WW and ZZ final states grow as m_H^3 for large m_H , and by requiring that the total width Γ_H is smaller than m_H . Note that a value of m_H close to the unitarity bound

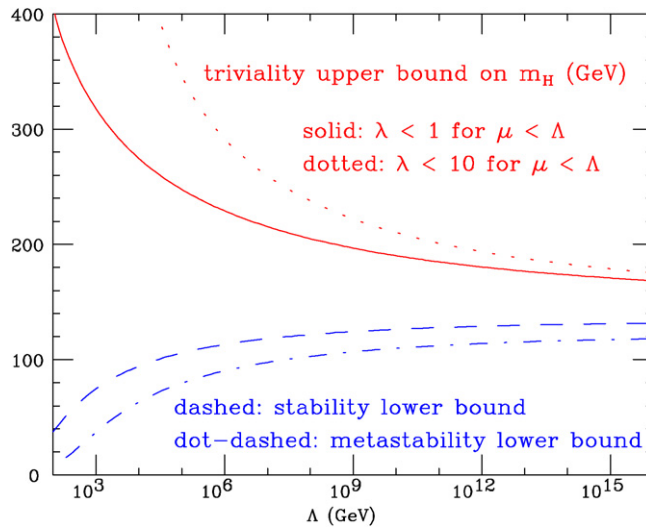


Fig. 2. Theoretical bounds on the Higgs boson mass as a function of the energy scale [26].

would imply a large value of the coupling λ : non-perturbative phenomena would in this case become important, and the whole mechanism should be reconsidered in this light, since a perturbative approach becomes unreliable.

A second class of theoretical bounds arises from the study of the behaviour of the quartic coupling λ under the effect of renormalization [8]. For sufficiently small values of the Higgs boson mass, $m_H^2 \simeq 2\lambda(v)v^2$, the solution of the renormalization-group equation for λ is a decreasing function of the renormalization scale μ for μ of the order of the weak scale, and eventually becomes negative at a value Λ which depends on the initial condition for λ at the weak scale, $\mu = v$: the smaller $\lambda(v)$, the smaller Λ . If $\lambda(\mu) < 0$, the effective potential is not bounded from below, and the ground state becomes unstable. Therefore, the theory is only consistent at energy scales $\mu \lesssim \Lambda$. Correspondingly, m_H has a Λ -dependent lower bound, which is shown in Fig. 2, taken from a recent analysis [26]. A slightly less restrictive lower bound is obtained by allowing metastability of the ground state, with the constraint that its lifetime be larger than the age of the universe. This metastability bound is also shown in Fig. 2.

The argument can be reversed: some new degrees of freedom would have to show their effects at energy scales of the order of Λ , in order to restore the stability of the ground state. The value of Λ is therefore of great importance, since it is related to the energy scale at which we should expect non-standard phenomena to take place.

For larger values of $m_H^2 \simeq 2\lambda(v)v^2$, the running coupling $\lambda(\mu)$ increases at larger scales, and has a simple (Landau) pole in $\log \mu$, similarly to what happens in electrodynamics. This behaviour is easily seen in the perturbative solution for the running coupling, and is confirmed by non-perturbative (lattice) computations (see, e.g., [27] and references therein). The theory cannot then be consistent up to arbitrarily large energies. Thus, one may view the SM as an effective theory, valid up to some energy scale Λ , such that $\lambda(\mu)$ remains small for $\mu \lesssim \Lambda$. Since the larger the value of $\lambda(v)$, the smaller the value of Λ , a Λ -dependent upper bound on m_H is obtained. This bound (usually referred to as the triviviality bound) is also shown in Fig. 2, for two different choices of the maximum value for λ .

We see from Fig. 2 that if the SM is assumed to be valid up to very large energy scales, of the order of the extrapolated unification scale for the gauge coupling constants [28], a relatively small allowed range for the Higgs boson mass is obtained.

2.3. Direct searches and constraints from precision measurements

Most of our present experimental knowledge about the SM Higgs boson comes from the study of e^+e^- collisions performed at LEP and the SLC between 1988 and 2000. As concerns the Higgs boson, the results of this enormous amount of experimental work can be summarized as follows. No direct evidence for the existence of the SM Higgs has been produced. This allows one to set a lower limit on the Higgs mass of 114.4 GeV, mainly based of the non-observation of Higgs bosons in association with a Z^0 [12,29,13], followed by the decay of the Higgs into a heavy fermion–antifermion pair [6].

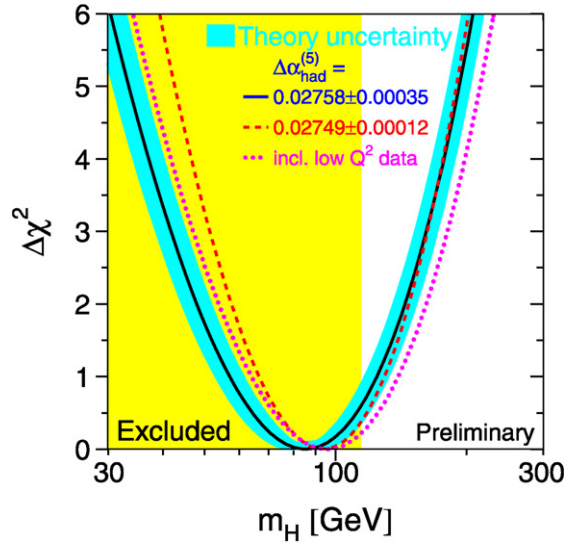


Fig. 3. The value of $\Delta\chi^2 = \chi^2 - \chi_{\min}^2$ as a function of m_H , used as a fit parameter, for a global fit to precision observables in the SM [7].

On the other hand, the value of the Higgs boson mass affects the SM predictions for the observables measured in experiments at LEP, the SLC, the Tevatron collider and elsewhere, through radiative corrections. The impact of the Higgs boson is relatively small; furthermore, the dependence of the precision observables on the Higgs mass is only logarithmic at one loop (as opposed, for example, to the dependence on the top quark mass, which is quadratic) [30].

However, thanks to the great accuracies achieved by the electroweak experiments (some of the observables have been measured with an accuracy of the order of 0.1%), it is possible to perform a global fit [7] using the SM Higgs mass as a free parameter, thus obtaining a preferred value for m_H and the corresponding uncertainty.² The result of such a global fit is displayed in Fig. 3, where the value of $\Delta\chi^2 = \chi^2 - \chi_{\min}^2$ is shown as a function of m_H . The solid black curve and the dotted red curve correspond to different evaluations of the renormalization of the electromagnetic coupling due to light quarks, and the dashed purple line includes more low- Q^2 data. (For colours see the web version of this article.) The blue band represents the uncertainty due to higher-order effects. At the $\Delta\chi^2 = 1$ level, the global electroweak fit yields [7] $\log_{10} m_H(\text{GeV}) = 1.93^{+0.16}_{-0.17}$, or

$$m_H = 85^{+39}_{-28} \text{ GeV} \quad (7)$$

which is consistent with the bounds discussed in the previous section. It should be noted, however, that the value of m_H preferred by the precision observables is slightly below the present exclusion limit.

The fit is reasonably good, but not excellent: the value of χ^2 per degree of freedom is 17.8/13, corresponding to a probability of 17%. This poorness of fit is mainly due to a marginal discrepancy between forward-backward asymmetries for hadronic final states, which favour m_H around 400 GeV, and leptonic asymmetries, which prefer a lower value [31]. This effect is illustrated in the left panel of Fig. 4, where the computed value of the weak mixing angle is shown as a function of m_H , with the top mass fixed at its central value or at its lower and upper bounds. The mass of the W boson also indicates a relatively small value of the Higgs mass, as one can see in the right panel of Fig. 4, where we show the computed value of the W boson mass as a function of m_H . The measured value of m_W is also shown, with its present error band. Indeed, within the SM, this single observable would prefer a value of m_H below the direct LEP exclusion limit.

Higgs searches at the Tevatron are not sensitive yet to the small cross-sections valid for the SM Higgs [33], but the situation may change with the accumulation of sufficient integrated luminosity. On the other hand, the LHC should be sensitive to the SM Higgs in its entire possible mass range [34]. We note here only that the most important SM Higgs

² It should be recalled that a similar procedure was followed with the top quark mass as a free parameter, before the direct observation of the top quark at the Tevatron. The fitted value of m_t was very close to the observed value of about 175 GeV.

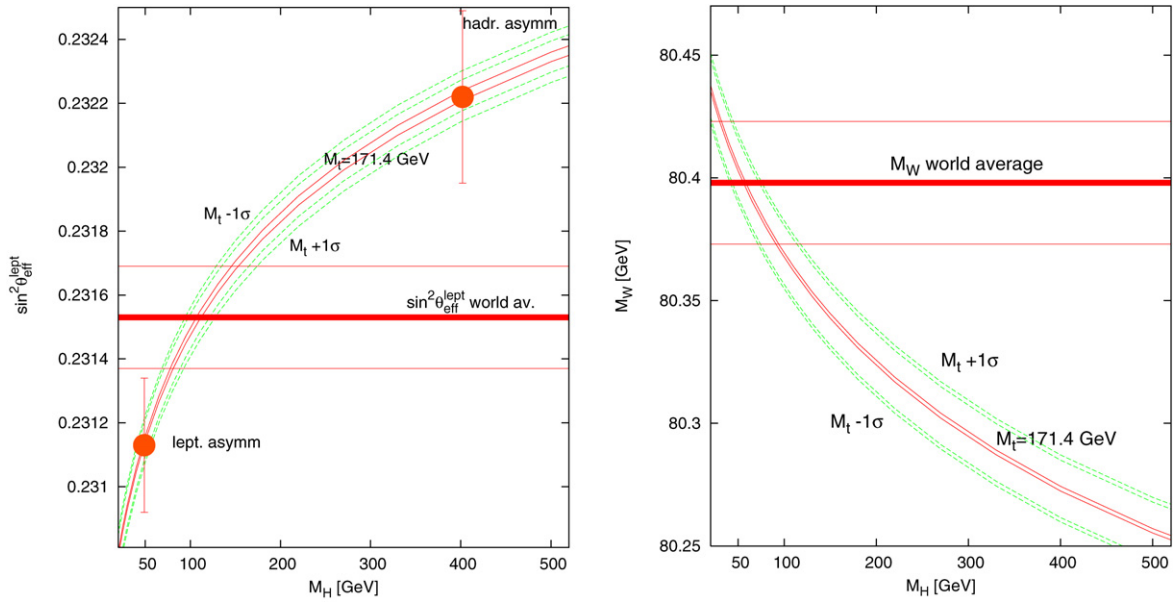


Fig. 4. Impact of different observables on the fit to m_H (updated from [32]).

production mechanisms are expected to be gluon-gluon [24] and WW/ZZ fusion [35]. Production in association with a W/Z boson [29,36] or a top-antitop pair [37] is suppressed but may still play an important role for detection.

3. Supersymmetric Models

3.1. Motivations

There are deep and general reasons to envisage supersymmetric extensions of the SM (for a recent review with references to the original literature, see, e.g., [38]). The most general symmetries of local relativistic quantum field theories include supersymmetry, and the link it provides between bosons and fermions is the only type of symmetry that has not yet been discovered in the observed fundamental interactions. Supersymmetry also appears to be an essential feature of string theory, our best candidate for a unified quantum theory of all interactions. Finally, supersymmetry provides, in its linear realization, a rationale for the existence of elementary scalars, since they are combined with chiral fermions in the same supersymmetric representations.

The failure to find evidence for supersymmetry in experiments performed so far implies that it must be broken. Arguments that are not fundamental, but nevertheless very plausible, hint at an effective supersymmetry-breaking scale very close to the present experimental bounds, with the superpartners of the SM particles having masses around the Fermi scale of weak interactions, $G_F^{-1/2} \simeq 300$ GeV. The first reason is the so-called hierarchy problem [15]: in the SM, there is no symmetry protecting the Higgs mass parameter from quantum corrections quadratic in the ultraviolet cut-off scale of the theory, Λ . The scale invariance of the classical equations of motions, recovered in the limit of vanishing Higgs mass parameter, is expected to be broken not only by SM quantum effects, but also by the new physics at the cut-off scale. Therefore, extrapolating the SM to scales Λ much higher than the Fermi scale requires an increasingly unnatural fine-tuning of parameters. Supersymmetric field theories possess very special quantum properties, including the absence of most quadratic and some logarithmic quantum corrections [39]. Consequently, weak-scale supersymmetry would suppress the sensitivity to the ultraviolet scale, and hence might explain the stability of the hierarchy [14] between the Fermi scale and the Planck scale, $M_P = 1/\sqrt{8\pi G_N}$. This feature enables supersymmetry also to provide a framework for understanding the dynamical origin of the hierarchy. Additionally, the renormalization-group evolution of the gauge couplings [28] is changed in theories with weak-scale supersymmetry [40], which facilitates the grand unification of the gauge couplings measured in low-energy experiments [41]. Supersymmetry at the electroweak scale also predicts a relatively light Higgs boson, as favoured by the precision electroweak data [7]. Moreover, in many supersymmetric models, the Lightest Supersymmetric Particle (LSP), typically

a weakly-interacting neutral spin-1/2 fermion called a neutralino, is stable and a suitable candidate for the dark matter inferred from astrophysical and cosmological data [42].

In the absence of a satisfactory dynamical understanding of supersymmetry breaking at a fundamental level (which cannot avoid addressing the vacuum energy problem, with its unexplained huge hierarchy between the vacuum energy scale, 10^{-3} – 10^{-4} eV, and the Planck scale), most phenomenological aspects of supersymmetry are usually discussed within the minimal supersymmetric extension of the Standard Model (MSSM) [43,44]. This is the most economical and predictive supersymmetric extension of the SM at the weak scale, but is certainly not the only possibility. In the following we describe in some detail the properties of the MSSM Higgs sector, and conclude with some comments on non-minimal models.

3.2. The MSSM Higgs sector

A general property of any (renormalizable) supersymmetric extension of the SM is the presence of at least two Higgs doublets [43], which leads to an extended Higgs sector. This is because, in the linear realization of $N = 1$ supersymmetry, spin-0 particles come in chiral supermultiplets, whose physical degrees of freedom are a complex spin-0 boson and a two-component spin-1/2 fermion. Trying to build a realistic supersymmetric extension of the SM with just one Higgs chiral supermultiplet H would immediately lead to a triple problem. First, there would be a massless charged fermion in the spectrum, since charged particles have Dirac masses, involving an even number of two-component spinors. Furthermore, since in supersymmetric models Yukawa couplings are determined by the superpotential, an analytic function of the chiral superfields, it would be impossible to write down superpotential couplings giving masses to all charged quarks and leptons. Finally, the spin-1/2 component of H would contribute to the chiral anomaly, whose cancellation is required for the quantum consistency of the theory. As a result, any (renormalizable) supersymmetric extension of the SM must contain at least two Higgs doublets with opposite weak hypercharge, denoted by

$$H_1 \equiv \begin{pmatrix} H_1^0 \\ H_1^- \end{pmatrix} \sim (1, 2, -1), \quad H_2 \equiv \begin{pmatrix} H_2^+ \\ H_2^0 \end{pmatrix} \sim (1, 2, +1) \quad (8)$$

where the $SU(3) \times SU(2) \times U(1)$ quantum numbers are given in brackets in an obvious notation. In contrast with a generic two-Higgs-doublet model, supersymmetry ensures that, in the MSSM, the H_1^0 has tree-level Yukawa couplings only with charge-1/3 quarks and charged leptons, and the H_2^0 only with charge-2/3 quarks, enforcing the natural suppression of flavour-changing neutral currents [45] in the limit of exact supersymmetry.

For the moment, we concentrate on the MSSM, which is defined by the following properties: (i) the minimal gauge group, $SU(3) \times SU(2) \times U(1)$; (ii) the minimal particle content, three generations of quarks and leptons and two Higgs doublets plus their superpartners; (iii) an exact discrete R-parity, which guarantees baryon- and lepton-number conservation in renormalizable interactions since $R = +1$ for SM particles and Higgs bosons whereas $R = -1$ for their superpartners; (iv) supersymmetry breaking parametrized by explicit but soft breaking terms, comprising gaugino and scalar mass terms and trilinear scalar couplings, without additional CP-violating phases beyond those of the Yukawa couplings.

The tree-level potential of the MSSM depends as follows on the two Higgs doublets in Eq. (8):

$$V_0 = m_1^2 |H_1|^2 + m_2^2 |H_2|^2 + m_3^2 (H_1 H_2 + h.c.) + \frac{g^2}{8} (H_2^\dagger \sigma^a H_2 + H_1^\dagger \sigma^a H_1)^2 + \frac{g'^2}{8} (|H_2|^2 - |H_1|^2)^2 \quad (9)$$

where m_1^2, m_2^2, m_3^2 are mass parameters, g and g' are the gauge coupling constants of $SU(2)$ and $U(1)$, respectively. Eq. (9) displays a crucial difference between the SM and MSSM potentials: in the SM the quartic coupling is an arbitrary parameter λ , proportional to the SM Higgs mass whereas in the MSSM, because of supersymmetry and in spite of the presence of a second Higgs doublet, the quartic scalar couplings in V_0 are related to the electroweak gauge couplings.

For suitable values of the mass parameters, V_0 has a minimum for $\langle H_1^0 \rangle = v_1 \neq 0$, $\langle H_2^0 \rangle = v_2 \neq 0$, where it is not restrictive to take v_1 and v_2 real and positive. The combination $v^2 = v_1^2 + v_2^2$ controls the Fermi scale and the weak boson masses, and is fixed by experimental data. Fermion masses are proportional to v_2 for charge-2/3 quarks, and to v_1 for charge-1/3 quarks and charged leptons. Neutrino masses can be accounted for by modifications of the

MSSM that do not affect the present discussion. Of the eight real degrees of freedom of Eq. (8), three are the would-be Goldstone bosons that provide the longitudinal components of the massive gauge bosons. The physical spectrum contains then three neutral states, two of which are CP-even (h, H) and one CP-odd (A), and two charged states (H^\pm).

3.3. Tree-level masses and couplings

At the classical level, all masses and couplings of the MSSM Higgs sector depend only on measured SM parameters and two more independent parameters: the latter are usually taken to be $\tan \beta \equiv v_2/v_1$ and one mass parameter, for example m_A . Although m_A is essentially unconstrained, apart from naturalness arguments suggesting that it should not be much larger than the Fermi scale, the range of $\tan \beta$ favoured by model calculations is $1 < \tan \beta < m_t/m_b$, where m_t and m_b are the running top and bottom masses evaluated near the Fermi scale. The remaining tree-level mass eigenvalues are then given by:

$$m_{H^\pm}^2 = m_W^2 + m_A^2, \quad m_{h,H}^2 = \frac{1}{2} \left[m_A^2 + m_Z^2 \mp \sqrt{(m_A^2 + m_Z^2)^2 - 4m_A^2 m_Z^2 \cos^2 2\beta} \right] \tag{10}$$

leading to stringent inequalities such as $m_W, m_A < m_{H^\pm}$, $m_h < m_Z |\cos 2\beta| < m_Z < m_H$, and $m_h < m_A < m_H$, holding at the classical level.

Similarly, all tree-level Higgs boson couplings can be computed in terms of the mixing angle α that is required to diagonalize the mass matrix of the neutral CP-even Higgs bosons, and is given by

$$\cos 2\alpha = -\cos 2\beta \frac{m_A^2 - m_Z^2}{m_H^2 - m_h^2}, \quad -\frac{\pi}{2} < \alpha \leq 0 \tag{11}$$

For example, the couplings of the three neutral MSSM Higgs bosons to vector-boson and fermion pairs are easily obtained from the SM Higgs couplings, by multiplying the latter by the α - and β -dependent factors summarized in Table 1. The remaining tree-level Higgs boson couplings in the MSSM can be found, for example, in [13]. An interesting situation in parameter space is the so-called decoupling limit: when $m_A^2 \gg m_Z^2$, the lightest neutral Higgs boson h behaves much as the SM Higgs boson, with $\alpha \sim (\beta - \pi/2)$, whereas the (H, A, H^\pm) are much heavier and nearly degenerate, forming an isospin doublet that decouples at sufficiently low energy.

3.4. Radiative corrections

An important consequence of the structure of V_0 in Eq. (9) is the existence of at least one neutral CP-even Higgs boson with mass smaller than m_Z (h) or very close to it (H), and significantly coupled with the W and Z bosons. This initially raised the hope that a crucial test of the MSSM Higgs sector could be performed at LEP. However, as reviewed below, the inclusion of radiative corrections to the MSSM Higgs sector [16–18] drastically changes the picture by increasing m_h , and it was not possible to complete the test of the MSSM at LEP, because its maximum centre-of-mass energy was $\sqrt{s} = 209$ GeV, not much above $2m_Z$.

Radiative corrections to the MSSM Higgs sector are dominated by loop diagrams involving virtual top quarks (t) and their superpartners, the spin-0 stop squarks, denoted by (\tilde{t}_1, \tilde{t}_2) in the basis of definite mass. The leading one-loop effects can easily be computed from the effective potential, and used for a first estimate of the one-loop corrected mass eigenvalues and of the mixing angle α in the neutral CP-even sector. Neglecting mixing in the stop mass matrix,

Table 1
Correction factors, relative to their SM values, for the couplings of the MSSM neutral Higgs bosons to fermion and vector boson pairs, expressed in terms of the angle β that describes the ratio of MSSM Higgs vacuum expectation values and the angle α that describes mixing between the CP-even neutral Higgs bosons

$d\bar{d}, l^+l^-$	$u\bar{u}$	W^+W^-, ZZ	
h	$-\sin \alpha / \cos \beta$	$\cos \alpha / \sin \beta$	$\sin(\beta - \alpha)$
H	$\cos \alpha / \cos \beta$	$\sin \alpha / \sin \beta$	$\cos(\beta - \alpha)$
A	$-i\gamma_5 \tan \beta$	$-i\gamma_5 \cot \beta$	0

working in the decoupling limit $m_A^2 \gg m_Z^2$, and assuming $m_{\tilde{t}_1}^2, m_{\tilde{t}_2}^2 \gg m_t^2$ and $\tan \beta \ll m_t/m_b$, the most important correction to the lightest Higgs boson mass takes the very simple form [16–18]:

$$\Delta m_h^2 \simeq \frac{3g^2 m_t^4}{16\pi^2 m_W^2} \log \frac{m_{\tilde{t}_1}^2 m_{\tilde{t}_2}^2}{m_t^4} \quad (12)$$

An alternative way to understand the size of this radiative correction is to consider an effective theory in which all the heavy MSSM particles, including the stop squarks, the extra Higgs bosons and the top quark, have been integrated out: in this case, the quartic Higgs coupling in the low-energy effective theory gets large positive contributions from one-loop stop and top diagrams, which translate into a corresponding increase in the Higgs boson mass. Mixing in the stop sector can lead to further, large positive contributions to m_h^2 . At one-loop order, and again for simplicity in the decoupling limit, one finds, in addition to (12):

$$(\Delta m_h^2)_{\text{mix}} \simeq \frac{3g^2 m_t^2 s_{2\theta}^2 (m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2)}{32\pi^2 m_W^2} \left[\log \frac{m_{\tilde{t}_1}^2}{m_{\tilde{t}_2}^2} + \frac{s_{2\theta}^2 (m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2)}{4m_t^2} \left(1 - \frac{1}{2} \frac{m_{\tilde{t}_1}^2 + m_{\tilde{t}_2}^2}{m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2} \log \frac{m_{\tilde{t}_1}^2}{m_{\tilde{t}_2}^2} \right) \right] \quad (13)$$

where θ is the mixing angle in the stop mass matrix and $s_{2\theta}^2 \equiv [\sin(2\theta)]^2$.

Over the years, extensive efforts have been devoted to progressive refinements of these and other radiative corrections to the MSSM Higgs sector, with special emphasis on the prediction for m_h . These include resummation of large logarithms using one- and two-loop renormalization group equations and calculation of all the most important two-loop contributions. The status of the two-loop calculations is presented in [38]: the most important ones are $O(\alpha_s \alpha_t)$ corrections, which tend to decrease the upper bound on m_h , and $O(\alpha_t^2)$ corrections, which can partially compensate the previous effect. The most important two-loop calculations are implemented in [46], and we note that the leading three-loop corrections have recently been shown to be small [47]. Phenomenological analyses within the MSSM point at a typical upper bound $m_h^{\text{max}} \sim 130$ GeV as seen in Fig. 5: calculations performed within specific models for supersymmetry breaking tend to produce lower values of m_h^{max} , whereas stretching the model parameters beyond their natural values, and taking the uncertainty in the top quark mass into full account, can produce slightly higher values of m_h^{max} . Radiative corrections also affect the MSSM Higgs boson couplings. For moderate values of $\tan \beta$, the leading corrections can be included by using tree-level formulae with running couplings, evaluated at an appropriate scale, and the loop-corrected values of the mixing angles α and β . For large values of $\tan \beta$, however, there can also be important threshold corrections to the Higgs couplings to the bottom/sbottom sector, with strong dependences on the model parameters.

3.5. Signals and phenomenology

The experimental signatures of the MSSM Higgs bosons depend on their branching ratios and on their production cross sections at high-energy colliders. A systematic discussion can be very complicated, due to the large number of parameters involved, already at the classical level and even more so after the inclusion of radiative corrections. As for the branching ratios, we content ourselves by mentioning some general features that may distinguish the lightest MSSM Higgs boson from the SM one, when the remaining MSSM Higgs bosons are significantly heavier than the Z but not yet completely decoupled. Being typically lighter than 130 GeV, h mostly decays into fermion pairs. Since $\tan \beta > 1$ in general, the branching ratios for h decays into $b\bar{b}$ and $\tau^+\tau^-$ tend to be enhanced with respect to a SM Higgs boson of the same mass. Decays of h into AA , or into final states with an even number of supersymmetric particles, may be important in the regions of parameter space where they are kinematically allowed.

In e^+e^- collisions, the main production mechanisms for the MSSM neutral CP-even Higgs bosons are $e^+e^- \rightarrow hZ(HZ)$ and $e^+e^- \rightarrow hA(HA)$. Vector-boson fusion can be important for centre-of-mass energies much larger than the Higgs mass. This mechanism did not play an important role at LEP, but it could play a role at a future high-energy linear collider. At LEP, the relevant final states were those generated by the intermediate states hZ and hA , which play a complementary role, since their cross-sections are proportional to $\sin^2(\beta - \alpha)$ and $\cos^2(\beta - \alpha)$, respectively. Extensive searches have been carried out by the four LEP collaborations and are summarized in the reports of the Higgs Working group (for the most recent one at the time of this writing, see [48]). Values of m_h and m_A below 92–93 GeV are typically excluded at the 95% c.l. over most of the MSSM parameter space. The limit on m_h gradually approaches that in the SM in the decoupling limit.

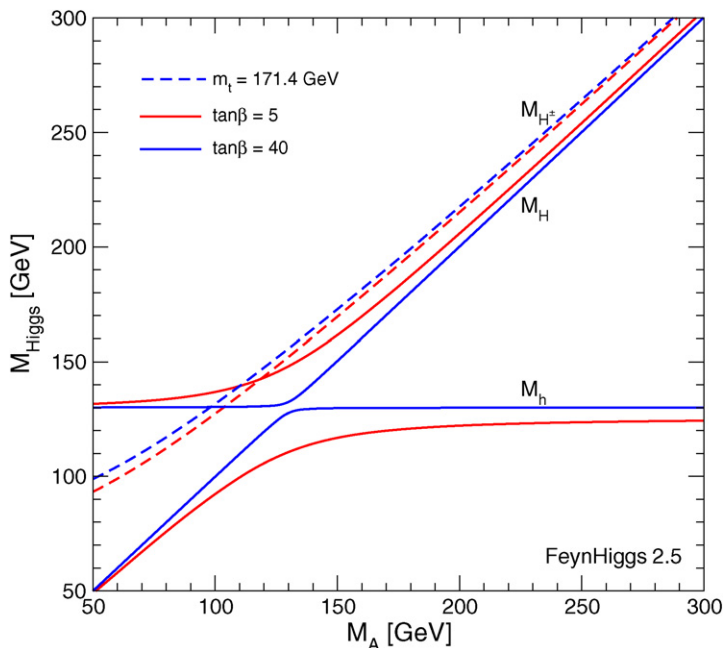


Fig. 5. Calculations of the lightest Higgs boson h , the heavier CP-even neutral Higgs boson H and the charged Higgs bosons H^\pm as functions of the mass of the CP-odd Higgs boson A , as calculated in the MSSM with a squark mass of 1 TeV, a gluino mass of 800 GeV, a wino mass of 200 GeV, $A = \mu/\tan\beta + 2000$ GeV and $\mu = 200$ GeV, assuming $m_t = 171.4$ GeV, provided by S. Heinemeyer using the first of Refs. [46].

The present Tevatron searches for the lightest MSSM Higgs boson [33] are marginally sensitive to the region of very large $\tan\beta \sim m_t/m_b$, when the Yukawa couplings to b quarks and τ leptons are strongly enhanced with respect to the SM, but there are also strong constraints from rare B decays. Studies [34] indicate that the LHC will be sensitive to most of the MSSM parameter space, but distinguishing the light MSSM Higgs from the SM Higgs will not always be possible, in the absence of signals for other MSSM particles.

3.6. Non-minimal models

While CP cannot be spontaneously broken by the tree-level MSSM Higgs potential, there is still the possibility of radiatively induced CP-violating effects in the MSSM Higgs sector [49], coming from explicit CP-violating phases in other sectors of the MSSM. This leads to further complications in the discussion of the MSSM Higgs searches, since all three neutral Higgs bosons can mix.

In the constrained MSSM with universal soft supersymmetry-breaking parameters, there are just two additional explicit CP-violating phases, namely the phases of the (common) gaugino mass $m_{1/2}$ and soft supersymmetry-breaking trilinear coupling A relative to the Higgs mixing parameter μ . Without loss of generality, one may assume that μ is real, and term the CP-violating phases $\theta_{1/2}$ and θ_A . Their first effect on the Higgs sector of the MSSM is the generation of off-diagonal mass terms mixing the CP-even Higgs bosons h, H with the CP-odd boson A [50], leading to a general 3×3 Higgs-boson mass matrix. Each of the off-diagonal CP-violating scalar-pseudoscalar mixing contributions to the neutral MSSM mass-squared matrix contains terms scaling qualitatively as

$$M_{SP}^2 \sim \frac{m_t^4 \sin\theta_A |\mu| |A_t|}{v^2 32\pi^2 M_{\text{SUSY}}^2} \left(6, \frac{|A_t|^2}{M_{\text{SUSY}}^2}, \frac{|\mu|^2}{\tan\beta M_{\text{SUSY}}^2}, \frac{2 \cos\theta_A |\mu| |A_t|}{M_{\text{SUSY}}^2} \right) \quad (14)$$

where A_t is the effective trilinear coupling associated with the top squarks, and M_{SUSY} is an average stop mass. For $\theta_A \sim \pi/2$ and $|\mu|, |A_t| > M_{\text{SUSY}}$, M_{SP}^2 could be of order M_Z^2 , though one should take into account the stringent constraints coming from electric dipole moments [51]. The gaugino mass phase $\theta_{1/2}$ also has an influence at the two-loop level, via the gluino mass. In addition to inducing three-way mixing between the neutral Higgs bosons, the phases θ_A and $\theta_{1/2}$ also introduce CP-violating effects in the Higgs couplings. For codes to evaluate such effects, see [52] and the first paper in Ref. [46].

There is a sum rule requiring the sum of the squares of the three couplings of neutral Higgs bosons to vector boson pairs to equal the square of the single Higgs–vector–vector coupling in the Standard Model, whereas this squared coupling is shared between just the two CP-even Higgs bosons in the CP-conserving case displayed in Table 1. Observable manifestations of CP violation in the MSSM Higgs sector have been considered at LEP, the LHC and other accelerators [50]. It is possible that the lightest neutral Higgs boson may be lighter than the LEP lower limit obtained within the Standard Model or the CP-conserving MSSM, since it may have a weaker coupling to vector-boson pairs [53].

Extensions of the MSSM can be introduced, where the Higgs sector is further enlarged and the Higgs masses are less constrained. We mention here a single example, the so-called Next-to-Minimal Supersymmetric Standard Model (NMSSM), whose Higgs sector includes not only two Higgs doublets, but also an additional singlet [43]. The phenomenology of the NMSSM Higgs sector was studied in [54] (for an updated account of later developments, see also [50]). Such an extension may slightly decrease the level of fine-tuning required to reconcile the present stringent lower bounds on supersymmetric particle and Higgs boson masses with the measured value of the Fermi scale. With an additional singlet, the quartic Higgs coupling becomes an independent parameter, as in the SM, thus the NMSSM is less predictive than the MSSM: the MSSM upper limit on the lightest Higgs boson mass is somewhat relaxed in the NMSSM, becoming similar to the triviality bound of the SM. However, in the NMSSM there is less need to push the mass parameters controlling radiative corrections to extreme values, in order to satisfy the present experimental bounds.

4. Prospects

The Higgs boson may be not only the capstone of the Standard Model, but also provide the opening to a whole new world beyond the Standard Model. Much is known about its possible properties in both the Standard Model and many supersymmetric scenarios. The stage is now set for the LHC to prove all or most of these ideas wrong.

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