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Highlights of the LHC run 1 / Résultats marquants de la première période d'exploitation du GCH

Implications of the Higgs boson discovery

Les implications de la découverte du boson de Higgs

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ABSTRACT

With the discovery of the Higgs boson by the LHC in 2012, a new era started in which we have direct experimental information on the physics behind the breaking of the electroweak (EW) symmetry. This breaking plays a fundamental role in our understanding of particle physics and sits at the high-energy frontier beyond which we expect new physics that supersedes the Standard Model (SM). In this review we summarize what we have learned so far from LHC data in this respect. In the absence of new particles having been discovered, we discuss how the scrutiny of the properties of the Higgs boson (in search for deviations from SM expectations) is crucial as it can point the way for physics beyond the SM. We also emphasize how the value of the Higgs mass could have far-reaching implications for the stability of the EW vacuum if there is no new physics up to extremely large energies.

RÉSUMÉ

La découverte du boson de Higgs par les expériences du GCH en 2012 a ouvert une nouvelle ère, avec un accès expérimental direct à la dynamique responsable de la brisure de la symétrie électrofaible. Cette brisure de symétrie joue un rôle fondamental dans notre compréhension de la physique des particules et se situe à la limite des nos connaissances dans un domaine d'énergie au-delà duquel le modèle standard de la physique des particules devrait montrer ses limites. Dans cet article, nous résumons ce que les données du LHC nous ont d'ores et déjà appris. En l'absence de découverte de nouvelles particules, nous expliquons en quoi une étude méticuleuse des propriétés du boson de Higgs, et en particulier la recherche de déviations par rapport aux prédictions standards, est primordiale, puisqu'elle peut en effet indiquer comment dépasser ce modèle standard. Nous discutons aussi les implications de la valeur de la masse du boson de Higgs sur la stabilité du vide électrofaible dans l'hypothèse où le modèle standard reste valide jusqu'à des énergies extrêmement élevées.

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1. Introduction

"With great power comes great responsibility" says a good book [1] and a bad movie [2]. And logically "with great discoveries should come great measurements". The discovery of the Higgs boson [3] is definitively one of the greatest achievements in the recent history of fundamental sciences. The emerging understanding of the nature of this new particle is confirmed by various measurements of its properties, all consistent with the Brout–Englert–Higgs mechanism [4] and other properties of the Standard Model (SM).

The observation of the Higgs boson makes the questions about the dynamics at the origin of electroweak symmetry breaking more pressing. The most relevant and urgent issue now facing us concerns the structure of the newly discovered Higgs scalar. Are there additional states accompanying it? Is it elementary or composite? Could this really be the first elementary scalar observed in Nature, or could it just be a bound state arising from some novel strong dynamics, like a π or η in QCD? The answer to these questions will have profound implications on our picture of fundamental physics through its bearing on the hierarchy problem. Establishing, to the best of our experimental capability, that the Higgs boson is elementary, weakly coupled and solitary, would surely be shocking, but it may well start a revolution in the basic concepts of quantum mechanics and space-time. If instead deviations from the SM emerge in the dynamics of the Higgs, we will have to use them as a diagnostic tool of the underlying dynamics. A crucial part of this program is the identification of the smoking guns of compositeness in Higgs dynamics. Moreover, along this basic question there are more specific ones we can ask, related to the symmetry properties of the new state. For instance, it is essential to establish whether the new scalar is indeed "a Higgs" fitting into a SU(2) doublet and not some exotic impostor, like for instance a pseudo-dilaton. Although there is really no strong theoretical motivation for such an alternative, and so far the data disfavor it, it remains a logical possibility that can be tested and possibly ruled out. A perhaps more interesting question is whether the Higgs particle is just an ordinary composite, like a σ, or whether it is a pseudo-Nambu–Goldstone boson, like the π. The answer to this question will give us important clues on the high-energy/ultra-violet (UV) completion of the electroweak breaking dynamics.

Identifying the expected deviations in the Higgs couplings should be one priority for the next runs of the LHC and might call in addition for a dedicated program at future colliders.

In the absence of any hint of physics beyond the Standard Model (BSM), a blind extrapolation of the Standard Model to high energies reveals an intriguing possibility: that the electroweak vacuum might be unstable, albeit with an extremely large lifetime against vacuum decay. This results from the rather peculiar region of parameter space in which we might be living (under the rather strong assumption about the absence of new physics), very close to the boundary between stability and instability. The meaning and possible implications of such coincidence are far from clear.

2. Beyond the standard model implications

2.1. Naturalness as a guide to BSM physics

With the addition of the Higgs boson, the SM is now theoretically consistent at the perturbative level a priori up to very high scale, possibly the Planck scale of quantum gravity, see Section 3. With the discovery of the Higgs boson by the ATLAS and CMS experiments [3], the picture would be perfect were it not for the fact that the quantum corrections to the Higgs potential reveal a dramatic sensitivity to the details of the physics at very high energy, as if Newton would have realized that the exact value of the top quark mass plays a crucial role in the motion of the Moon around the Earth. This property goes against our intuition that physical phenomena at different scales decouple from each other. Concretely, the one-loop corrections to $M_{\rm h}$, the mass parameter in the Higgs potential, are depicted in Fig. 1 and amount to

$$\delta M_{\rm h}^2 = \left(\frac{1}{4}(9g^2+3g'^2) - 12y_{\rm t}^2 + 6\lambda\right)\frac{\Lambda^2}{32\pi^2} = \left(M_Z^2 + 2M_W^2 - 4M_{\rm t}^2 + M_H^2\right)\frac{3G_{\rm F}\Lambda^2}{16\sqrt{2}} \tag{1}$$

where Λ stands for the typical mass scale of any new threshold associated with new particles or new dynamics beyond the Standard Model, and g and g' are the $SU(2)_L$ and $U(1)_Y$ gauge couplings, $y_{\rm t}$ is the top Yukawa coupling and λ is the Higgs self-coupling, all these parameters being related to the masses of SM particles and the Fermi constant, $G_{\rm F}$. As an example, for a 10 TeV SM cutoff, the gauge, top and Higgs contributions to the Higgs mass squared corrections are respectively of the order of $(600\ \text{GeV})^2$, $-(1.5\ \text{TeV})^2$ and $(800\ \text{GeV})^2$, all quite far from what the Higgs mass should be. The SM particles give unnaturally large corrections to the Higgs mass: they destabilize the Higgs vacuum expectation value (vev) and tend to push it towards the UV cutoff of the SM. Some precise adjustment (fine-tuning) between the bare mass and the loop correction is needed to maintain the vev of the Higgs around the weak scale: given two large numbers, their sum/difference will naturally be of the same order unless these numbers are almost equal up to several significant digits. This is the so called hierarchy problem [5–7]. It is a generic technical problem in any theory involving elementary light scalar fields.

It is often argued that the quadratic divergences in the Higgs mass corrections have no meaning since they can be set to 0 in dimensional regularization. Hence the belief that there is no hierarchy problem. This is actually true in the SM (at least when gravity is ignored) which involves a single scale. The hierarchy problem exists only when multiple scales are present. The hierarchy problem can be seen when dealing with the renormalized running Higgs mass, see Ref. [8] for a

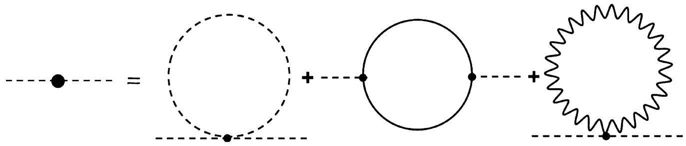

Fig. 1. One-loop corrections to the Higgs mass. The three diagrams are quadratically divergent and make the Higgs mass highly UV-sensitive.

Table 1
ATLAS and CMS fits of the Higgs mass (in GeV) in the $H \to \gamma\gamma$ and $H \to ZZ^* \to 2\ell^+2\ell^-$ channels and their combination. From Refs. [14].

Experiment	$H \to \gamma\gamma$	$H \to ZZ^* \to 4\ell$	Combined
ATLAS	$125.98 \pm 0.42(\text{stat.}) \pm 0.28(\text{syst.})$	$124.51 \pm 0.52(\text{stat.}) \pm 0.06(\text{syst.})$	$125.36 \pm 0.37(\text{stat.}) \pm 0.18(\text{syst.})$
CMS	$124.70 \pm 0.31(\text{stat.}) \pm 0.15(\text{syst.})$	$125.59 \pm 0.42(\text{stat.}) \pm 0.17(\text{syst.})$	$125.02 \pm 0.27(\text{stat.}) \pm 0.15(\text{syst.})$

recent narrative, and results from *finite* threshold corrections generated by new particles coupled with the Higgs boson: these corrections are proportional to the masses of these new particles and potentially drive the value of the Higgs mass itself to the largest energy scale unless they are precisely counter-balanced by the boundary value of the Higgs running mass at high energy.

Extra structures, particles and/or symmetries, are needed to stabilize the Higgs potential and to screen the radiative corrections. For instance one can make use of the following guidelines:

– The **spin trick** [9]: in general, a particle of spin s has $2s+1$ degrees of polarization with the only exception of a particle moving at the speed of light, in which case fewer polarizations may be physical. And conversely, if a symmetry decouples some polarization states then the particle will necessarily propagate at the speed of light and thus remain massless. For instance, gauge invariance ensures that the longitudinal polarization of a vector field is non-physical, and chiral symmetry keeps only one fermion chirality: both spin-1 and spin-1/2 particles are protected from dangerous radiative corrections. Unfortunately, this spin trick cannot be used for a spin-0 particle such as the SM Higgs scalar boson.
– The **Goldstone theorem** [10]: when a global symmetry is spontaneously broken, the spectrum contains a *massless* spin-0 particle. However, here again, it seems difficult to invoke this trick to protect the SM Higgs boson from radiative corrections since a Nambu–Goldstone boson can only have derivative couplings, unlike the Higgs field. Little Higgs models [11] have been constructed to circumvent these difficulties and they provide realistic examples of Higgs as a (pseudo-)Nambu–Goldstone boson.

In the late 60s, the Coleman–Mandula and Haag–Lopuszanski–Sohnius theorems [12] taught model-builders how to apply the spin trick to spin-0 particles: the four-dimensional Poincaré symmetry has to be enlarged. The first construction of this type consists in embedding the 4D Poincaré algebra into a super-algebra. Then the supersymmetry between fermions and bosons extends the spin trick to scalar particles: spin-0 particles inherit the chirality-protection of their spin-1/2 superpartners. Actually there exists an even simpler way of enlarging the Poincaré symmetry, which is going into extra dimension(s): the 5D Poincaré algebra obviously contains the 4D Poincaré algebra as a sub-algebra. After compactification of the extra dimensions, from a 4D dimensional point of view, the higher-dimensional gauge field decomposes into a 4D gauge field (the components along our 4D world) and 4D scalar fields (the components along the extra dimensions). The symmetry between vectors and scalars allows one to extend the spin trick to spin-0 particles.

Neither supersymmetry nor higher-dimensional Poincaré symmetry are exact symmetries of Nature. Therefore, if they ever have a role to play, they have to be broken. In order not to lose any of their benefits, this breaking has to proceed without reintroducing any strong UV dependence into the renormalized scalar mass: a *soft breaking* of these symmetries is needed.

The biggest surprise at the LHC run 1 is maybe that there has been no big surprise beyond the discovery of a light Higgs boson, i.e. that no hint of physics beyond the Standard Model has been observed. This situation puts the naturalness principle under stress, if not under siege, and cries out either for a rapid discovery at run 2 or for a change of paradigm in favor of alternatives such as the multiverse idea (for a recent account, see Ref. [13]).

2.2. The Higgs mass as a model-discriminator

The value of the Higgs boson mass opens many decay modes at a rate accessible experimentally. Two channels are particularly accurate in accessing the Higgs mass itself: $H \to \gamma\gamma$ and $H \to ZZ^* \to 2\ell^+2\ell^-$. Table 1 reports the ATLAS and CMS measurements, reaching a 0.2%-accuracy (to be compared to the 0.5%-accuracy of the top mass measurement).

As will be discussed in detail in Section 3, under the assumption that the SM laws govern Nature up to very high energy, the precise value of the Higgs mass has thrilling implications on the stability of the EW vacuum and hence the fate of our Universe.

The value of the Higgs mass also gives clues about the details of possible UV completions of the SM itself. This can be exemplified in the leading scenarios, namely the Minimal Supersymmetric Model (MSSM) and the Minimal Composite Higgs model (MCHM). In short, the Higgs mass is larger than what is typically expected in the MSSM and smaller than what is expected in the MCHM. At the classical/Born level, the mass of the lightest (SM-like) Higgs boson, in the MSSM, is bounded to be lower than the Z-boson mass since supersymmetry dictates the Higgs quartic to be fixed in terms of the gauge couplings. Some significant amount of radiative corrections, mostly from the top and stop sectors, are therefore called to raise the value of the Higgs mass. At one-loop, the Higgs mass can be approximated by

$$M_{\rm h}^2 \simeq M_Z^2 \cos^2 2\beta + \frac{3\sqrt{2} G_{\rm F} M_{\rm t}^4}{16\pi^2}\left[\log\frac{M_{\tilde{\rm t}}^2}{M_{\rm t}^2} + \frac{X_{\rm t}^2}{M_{\tilde{\rm t}}^2}\left(1-\frac{X_{\rm t}^2}{12 M_{\tilde{\rm t}}^2}\right)\right] \tag{2}$$

where $M_{\tilde{\rm t}}^2 = M_{Q_3} M_{U_3}$ is the geometric mean of the stop masses and $X_{\rm t}$ is the mixing between the two stops. Clearly, a Higgs boson as heavy as 125 GeV requires either heavy stops ($M_{\tilde{\rm t}} > 800$ GeV) and/or maximally mixed stops ($X_{\rm t} \simeq \sqrt{6} M_{\tilde{\rm t}}$), which brings back some amount of irreducible fine-tuning or call for non-trivial boundary conditions like non-universal gaugino masses at high-energy. Going beyond the minimal model, for instance by adding an extra gauge singlet, can easily help increasing the Higgs mass with significantly less amount of tuning, see for instance Ref. [15] for a recent discussion.

In the Minimal Composite Higgs models, the Higgs boson emerges from the strong sector as a pseudo-Nambu Goldstone boson. Therefore, the strong interactions themselves are not responsible for generating a potential for the Higgs field, that is generated only at the one-loop level from the interactions between the strong sector and the SM. Computing the details of the potential from first principles remains out of reach but it is possible [16], like what is done to compute the mass difference between the charged and neutral pions in QCD, to estimate the Higgs mass using general properties about the asymptotic behavior of correlators, i.e. imposing the saturation of the Weinberg sum rules with the first few light resonances, to obtain

$$M_{\rm h}^2 \simeq \frac{3 M_{\rm t}^2 M_{\rm Q}^2}{\pi^2 f^2} \tag{3}$$

where f is the scale of the strong interactions (the decay constant of the Higgs boson, the equivalent of f_π of the QCD pions) and $M_{\rm Q}$ is the typical mass scale of the fermion resonances (aka the top partners). This estimate can read as

$$M_{\rm Q} \simeq \left(\frac{M_{\rm h}}{125\ \text{GeV}}\right)\left(\frac{160\ \text{GeV}}{M_{\rm t}}\right)\left(\frac{f}{500\ \text{GeV}}\right) 700\ \text{GeV} \tag{4}$$

For a natural set-up ($v^2/f^2 \le 0.2$), we therefore expect some light top partners below one TeV. The discovery of such fermionic top-partners would be a first evidence of a strong dynamics at the origin of the breaking of the electroweak symmetry.

2.3. The Higgs profile as a probe of the deformations away from the SM

A dedicated study of the Higgs boson properties and couplings offers a way to infer what the structure of physics beyond the Standard Model can be. Natural models trying to give a rationale for why/how the Higgs mass is screened from high energy corrections at the quantum level generically predict some deviations in the Higgs couplings compared to the SM predictions of the order 1% to 100%. The current Higgs data accumulated at the LHC by the ATLAS and CMS collaborations already constrain the Higgs couplings to massive gauge bosons and to fermions not to deviate by more than 20–30% from the SM predictions, see Figs. 2 and 3.

In general, new physics can deform the SM in many ways but most of these deformations are already severely constrained by electroweak precision measurements or flavor data. Assuming flavor universality among the couplings between the Higgs boson and the SM fermions, it was shown [18,19] that eight directions among the leading CP-conserving deformations of the SM can be probed, at tree-level, only in processes with a physical Higgs boson. These deformation induce deviations in the Higgs couplings that respect the Lorentz structure of the SM interactions, or generate simple new interactions of the Higgs boson to the W and Z field strengths, or induce some contact interactions of the Higgs boson to photons (and to a photon and a Z boson) and gluons that take the form of the ones that are generated by integrating out the top quark. In other words, the Higgs couplings are described, in the unitary gauge, by the following effective Lagrangian [20,21]

$$\begin{aligned}
\mathcal{L} = {} & \kappa_3 \frac{m_H^2}{2v} H^3 + \kappa_Z \frac{m_Z^2}{v} Z_\mu Z^\mu H + \kappa_W \frac{2 m_W^2}{v} W_\mu^+ W^{-\mu} H \\
& + \kappa_{\rm g} \frac{\alpha_{\rm s}}{12\pi v} G_{\mu\nu}^{\rm a} G^{{\rm a}\mu\nu} H + \kappa_\gamma \frac{\alpha}{2\pi v} A_{\mu\nu} A^{\mu\nu} H + \kappa_{Z\gamma} \frac{\alpha}{\pi v} A_{\mu\nu} Z^{\mu\nu} H \\
& + \kappa_{VV} \frac{\alpha}{2\pi v}\left(\cos^2\theta_W Z_{\mu\nu} Z^{\mu\nu} + 2\, W_{\mu\nu}^+ W^{-\mu\nu}\right) H \\
& - \left(\kappa_{\rm t} \sum_{\rm f=u,c,t} \frac{m_{\rm f}}{v} \bar{f}_L f_R + \kappa_b \sum_{\rm f=d,s,b} \frac{m_{\rm f}}{v} \bar{f}_L f_R + \kappa_\tau \sum_{\rm f=e,\mu,\tau} \frac{m_{\rm f}}{v} \bar{f}_L f_R + h.c.\right) H
\end{aligned} \tag{5}$$

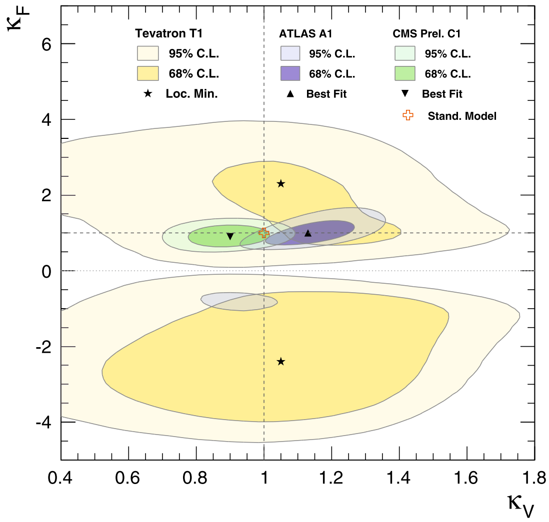

Fig. 2. (Color online.) Global fit of the ATLAS and CMS Higgs data. The Higgs couplings to massive EW gauge bosons and to fermions are rescaled from their SM values by the parameters κ_V and κ_f. From Ref. [17].

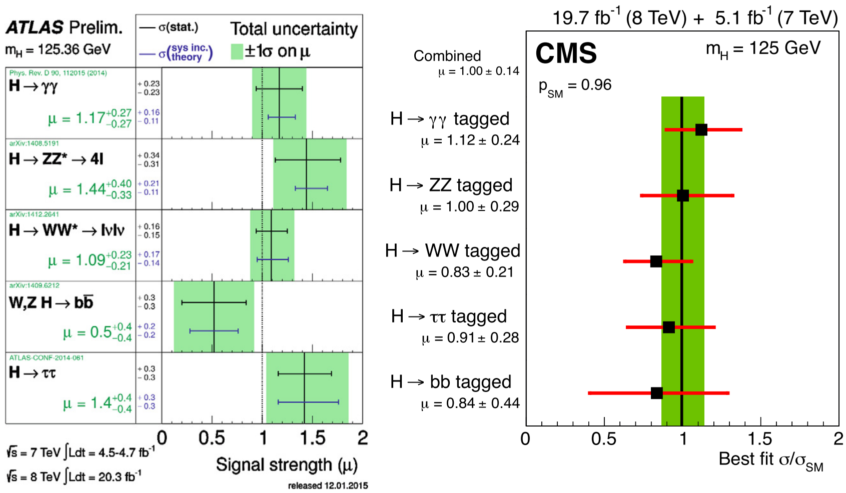

Fig. 3. (Color online.) Values of the best-fit signal strengths, μ (the ratio of the production cross section times the branching ratio normalized to their SM predictions), in the various Higgs channels.

In the SM, the Higgs boson does not couple with massless gauge bosons at tree level, hence $\kappa_g = \kappa_\gamma = \kappa_{Z\gamma} = 0$. Nonetheless, the contact operators are generated radiatively by loops of SM particles. In particular, the top quark gives a contribution to the three coefficients $\kappa_g, \kappa_\gamma, \kappa_{Z\gamma}$ that does not decouple in the infinite top mass limit. For instance, in that limit, the top quark contribution to $h \to gg/\gamma\gamma$ is effectively equivalent to the one obtained with $\kappa_\gamma = \kappa_g = 1$ [22].

The coefficient for the contact interactions of the Higgs boson to the W and Z field strengths is not independent but obeys the relation

$$(1 - \cos^4\theta_W)\kappa_{VV} = \sin 2\theta_W\, \kappa_{Z\gamma} + \sin^2\theta_W\, \kappa_\gamma \tag{6}$$

This relation is a general consequence of the so-called custodial symmetry [23]. When the Higgs boson is part of an $SU(2)_L$ doublet, the custodial symmetry could only be broken by a single operator at the level of dimension-6 operators and it is accidentally realized among the interactions with four derivatives, like the contact interactions considered. Custodial symmetry also implies

$$\kappa_Z = \kappa_W \tag{7}$$

leaving exactly eight free couplings [18,19]. Out of these eight coefficients, only κ_V can be indirectly constrained by EW precision data at a level comparable from the direct constraints from LHC Higgs data [24].

Table 2
Largest deviations in the Higgs couplings expected in a variety of new physics models. From Ref. [40].

Models	κ_f	κ_V	κ_g	κ_γ	$\kappa_{Z\gamma}$	κ_3
MSSM	✓					✓
NMSSM	✓	✓	✓	✓	✓	✓
MCHM aka PNGB Higgs	✓	✓			✓	✓
SUSY Composite Higgs	✓	✓				✓
Higgs as a Dilaton			✓	✓	✓	✓
Partly-Composite Higgs			✓	✓	✓	✓
Bosonic Technicolor						✓

The effective Lagrangian of Eq. (5) can be amended by six extra Higgs couplings that break the CP symmetry

$$\mathcal{L} = \tilde{\kappa}_g \frac{\alpha_s}{12\pi v} G^a_{\mu\nu} \tilde{G}^{a\mu\nu} H + \tilde{\kappa}_\gamma \frac{\alpha}{2\pi v} A_{\mu\nu} \tilde{A}^{\mu\nu} H + \tilde{\kappa}_{Z\gamma} \frac{\alpha}{\pi v} A_{\mu\nu} \tilde{Z}^{\mu\nu} H$$
$$- \left(i\tilde{\kappa}_t \sum_{f=u,c,t} \frac{m_f}{v} \bar{f}_L f_R + i\tilde{\kappa}_b \sum_{f=d,s,b} \frac{m_f}{v} \bar{f}_L f_R + i\tilde{\kappa}_\tau \sum_{f=e,\mu,\tau} \frac{m_f}{v} \bar{f}_L f_R + h.c. \right) H \quad (8)$$

where $\tilde{F}_{\mu\nu} = \epsilon_{\mu\nu\rho\sigma} F^{\rho\sigma}$ is the dual field-strength of $F_{\mu\nu}$. It is certainly tempting to consider new sources of CP violation in the Higgs sector, potentially bringing in one of the necessary ingredients for a successful baryogenesis scenario. The prospects for measuring at the LHC these CP violating sources in the $ht\bar{t}$, $h\tau\bar{\tau}$ and $h\gamma\gamma$ couplings have been studied in Refs. [25–27] respectively. It should however be noted [28] that these CP violating couplings would induce quark and electron electric dipole moments at one- (for $\tilde{\kappa}_\gamma$ and $\tilde{\kappa}_{Z\gamma}$) or two-loops (for $\tilde{\kappa}_f$). Unless the Yukawa couplings of the Higgs to the electron and light quarks are significantly reduced compared to their SM values, these constraints severely limit the possibility to observe any CP violating signal in the Higgs sector at the LHC.

The coefficient κ_3 can be accessed only through double Higgs production processes, hence it will remain largely unconstrained at the LHC and a future machine like an ILC or a future circular collider might be needed to pin down this coupling [29–31]. The LHC will also have a limited sensitivity on the coefficient κ_τ since the lepton contribution to the Higgs production cross section remains subdominant and the only way to access the Higgs coupling is via the $H \to \tau^+\tau^-$ and possibly $H \to \mu^+\mu^-$ channels. Until the associated production of a Higgs with a pair of top quarks is observed, the Higgs coupling to the top quark is only probed indirectly via the one-loop gluon fusion production or the radiative decay into two photons. However, these two processes are only sensitive to the two combinations $(\kappa_t + \kappa_g)$ and $(\kappa_t + \kappa_\gamma)$ and a deviation in the Higgs coupling to the top quark can in principle always be masked by new contact interactions to photons and gluons [32]. Recently, it was proposed [33] to study the hard recoil of the Higgs boson against an extra jet which provides a second scale above the Higgs mass to probe the effective field theory (EFT) structure and resolve this coupling degeneracy. The double Higgs production by fusion of gluons also effectively introduces a second mass scale and can be used to separate the top Yukawa coupling from the contact interaction to gluons or photons [31,34]. The off-shell Higgs production, e.g. in $gg \to h^* \to ZZ \to 4\ell$, is another obvious place to break this degeneracy of the couplings and to learn about the top Yukawa coupling [35]. Note that these three channels will require some large integrated luminosity, beyond the run 2 of the LHC, to compete with the still delicate $t\bar{t}h$ channel [36].

The operators already bounded by EW precision data modify in general the Lorentz structure of the Higgs couplings and hence induce some modifications of the kinematical differential distributions [37,38]. A promising way to have a direct access to the Wilson coefficients of these operators in Higgs physics is to study the VH associated production with a W or a Z at large invariant mass [37,39]. It has not been estimated yet whether the sensitivity on the determination of the Wilson coefficients in these measurements can compete with the one derived for the study of anomalous gauge couplings. In any case, these differential distributions could also be a way to directly test the hypothesis that the Higgs boson belongs to an $SU(2)_L$ doublet together with the longitudinal components of the massive electroweak gauge bosons.

Various dynamics produce different patterns among the Higgs coupling deviations. Table 2 summarizes the largest effects expected in popular classes of models of new physics addressing the hierarchy problem. The correlations among these deviations can thus reveal the nature of the dynamics above the weak scale while their magnitude will indicate the scale of this new dynamics.

3. Electroweak vacuum stability

As reviewed in previous sections, after the first LHC run we know the Higgs exists and is light, with mass $M_H \simeq 126$ GeV [3], it has SM-like properties (with some room for deviations [14,17]) and no trace of BSM physics has been found (with bounds on the mass scale of different BSM scenarios, supersymmetric or otherwise, of order the TeV). For those willing to hold on to naturalness, the hierarchy problem afflicting the breaking of the electroweak (EW) symmetry would imply that BSM physics is most probably around the corner, likely on the reach of the next LHC runs.

A different attitude is, however, possible: disregard naturalness as a requisite for the physics associated with the breaking of the EW symmetry and explore the possibility that the scale of new physics, Λ, might be as large as the Planck scale, M_{Pl}.

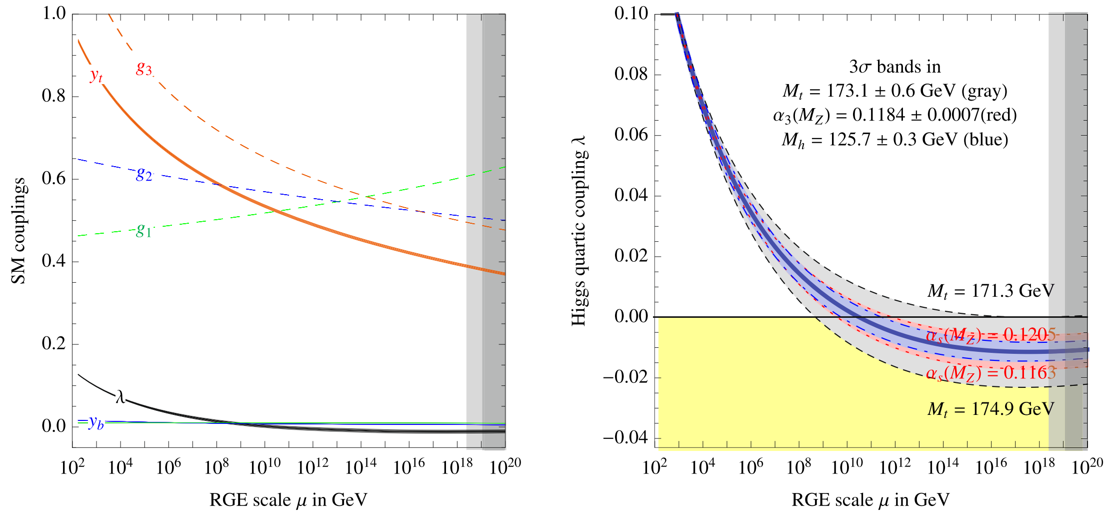

Fig. 4. (Color online.) Left: Evolution of SM couplings from the EW scale to $M_{\rm Pl}$. Right: Zoom on the evolution of the Higgs quartic, $\lambda(\mu)$, for $M_h = 125.7$ GeV, with uncertainties in the top mass, α_s and M_h as indicated. (Taken from Ref. [41]).

In this view, the SM should then describe physics in the vast range from M_W to $M_{\rm Pl}$. Fig. 4 (left) shows the running of the most relevant SM couplings extrapolated to very high energy scales using renormalization group (RG) techniques [41]. It shows the three $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge couplings getting closer in the UV but failing to unify precisely. It also shows how the top Yukawa coupling gets weaker in the UV (due to α_s effects, see below). The Higgs quartic coupling is also shown: it starts small at the EW scale, $\lambda(M_t) \sim 1/8$, because the Higgs turned out to be light, and gets even smaller at higher scales. The zoomed right plot in Fig. 4 shows that, in fact, λ does something interesting: it gets negative at around $\mu \sim 10^{10}$ GeV.

This steep behavior of λ is due to the effect of one-loop top corrections, which represent the dominant contribution to the beta function of λ, which describes the evolution of λ with scale: $\beta_\lambda \simeq d\lambda/d\log\mu = -6y_t^4/(16\pi^2)$, where μ is the renormalization scale and y_t is the sizable top Yukawa coupling. The dependence of β_λ on the fourth power of y_t explains the strong dependence of the running of λ on the top quark mass, as shown by the gray band in Fig. 4 (right), which corresponds to a 3σ variation of M_t around is central value (as indicated). The bigger (smaller) M_t is, the steeper (milder) the slope of the running λ.

There is a smaller dependence of the running of λ on the value of α_s, which affects β_λ indirectly through its effect on the running of y_t: $\beta_{y_t} \simeq dy_t/d\log\mu = y_t(9y_t^2/2 - 8g_s^2)/(16\pi^2)$, where g_s is the $SU(3)_C$ gauge coupling. This smaller effect is illustrated in the same Fig. 4 (right) by the thinner 3σ pink band, with higher (lower) α_s corresponding to softer (steeper) running. Finally, the thinnest band, in blue, corresponds to 3σ variations in the Higgs mass, as indicated. One also sees that λ flattens out after getting negative: in that range of scales, gauge couplings become comparable in size with y_t (see Fig. 4, left) and there is a cancellation leading to $\beta_\lambda \simeq 0$.

The trouble with λ becoming negative is that it causes an instability in the Higgs potential: at high field values the potential is dominated by the quartic term, and a good approximation to the full potential at some field value h requires couplings to be evaluated at a renormalization scale $\mu \sim h$. Therefore, at very high values of the Higgs field, the potential is $V(h \gg M_t) \simeq (1/4)\lambda(\mu = h)h^4$, which for $\lambda(h) < 0$ is much deeper than the EW vacuum. This instability problem caused by heavy fermions coupled with light scalars has been known for a long time [42] and has been investigated since then in the SM with increasing degree of precision [43], especially recently [41,44–47] when it became clear that the Higgs mass was in a critical region concerning the stability of the potential.

With the current precision in the Higgs and top masses and theoretical calculation of the stability bound, one concludes that (within theoretical assumptions about the absence of BSM physics) the EW vacuum would most likely be metastable. One should then worry about its lifetime against decay through quantum tunneling to larger field values.

The decay probability rate of the EW vacuum per unit time and unit volume [48] is $\sim h_I^4 \exp(-S_4)$, where h_I is the value of the field around the region of instability (the only relevant mass scale), and S_4 is the action of the 4-d Euclidean tunneling bounce solution, interpolating between the high field values and the EW phase. The simple analytical approximation $S_4 \simeq -8\pi^2/(3|\lambda(h_I)|)$ obtained for a negative-quartic potential $V \simeq -|\lambda(h)|h^4/4$, captures the main effect.[1] The logarithmic dependence of $\lambda(h)$ on its argument breaks the scale invariance of the classical potential and the tunneling occurs preferentially towards the scale h_I for which $\lambda(h)$ takes its minimum value (for which $\beta_\lambda = 0$). One gets for the decay rate

[1] The tunneling rate has been calculated beyond tree level, including the effect of fluctuations around the bounce solution in Ref. [49]. Gravitational effects, which have a negligible impact on the rate, were included in Ref. [50].

$dp/(dV\,dt) \sim h_I^4 \exp[-2600/(|\lambda|/0.01)]$. This has to be multiplied by the 4-d space-time volume inside our past light-cone: the fourth power of the age of the Universe $\sim\tau_U^4 \sim (e^{140}/M_{Pl})^4$. It is then clear that the exponential suppression of the decay rate [for the observed value $\lambda(h_I) \sim -0.01$] wins over the large 4-volume factor: the decay probability is extremely small, $p \ll 1$ or, in other words, the lifetime of the metastable EW vacuum τ_{EW} is extremely long, much larger than the age of the Universe. This conclusion would have been different for a smaller Higgs mass. In that case the running $\lambda(\mu)$ would be lower and could enter the region of negative $\lambda(\mu) < -0.05$ corresponding to a vacuum lifetime τ_{EW} smaller than τ_U (we will call this region of parameter space the instability region).

3.1. NNLO stability bound and implications

Fig. 4 shows that having an EW vacuum absolutely stable up to M_{Pl} requires values of M_t and α_s in some ∼2–3σ tension with their central experimental values. In the literature, this possibility has been discussed in terms of the so-called stability bound on M_h, which tells how heavy should M_h be to ensure a stable potential up to M_{Pl}. The state-of-the-art calculation of this stability bound [41,47] (at NNLO, or next-to-next-to leading-log order, see below) reads

$$M_h[\text{GeV}] > 129.6 + 2.0\,[M_t(\text{GeV}) - 173.35] - 0.5\left[\frac{\alpha_s(M_Z) - 0.1184}{0.0007}\right] \pm 0.3_{\text{th}} \tag{9}$$

The main uncertainty derives from the experimental uncertainty in the top mass measurement. The combination of the experimental measurements from Tevatron and LHC gives [51], $M_t = 173.34 \pm 0.27_{\text{stat}} \pm 0.71_{\text{syst}}$ GeV. The total 1σ error for M_t in Eq. (9) has been rounded up to 1 GeV and can be enlarged to allow for a somewhat larger theoretical error. The precise evaluation of the theoretical error in the top mass determination is a key issue for the improvement of this bound, see below. Second in importance is the error from $\alpha_s(M_Z) = 0.1184 \pm 0.0007$ [52]. Finally, the theoretical error is an estimate of higher order corrections, beyond NNLO. Such small error has been achieved only quite recently, with Refs. [41, 46,47] being the main contributors towards this goal.

In order to achieve this precision one has to calculate reliably the scalar potential in a wide range of field values, from the EW scale up to M_{Pl}. There are potentially large logs, $\log[h/M_t]$, that need to be resummed and this can be done using standard renormalization group techniques [53]. The ingredients for the NNLO calculation of the stability bound are: use the RG-improved two-loop effective potential [54], in which couplings are running with 3-loop beta functions [55] and use 2-loop matching [41,46,47] to relate λ and y_t to M_h and M_t.

To illustrate the need of such precise calculation of the stability bound given our precise knowledge of M_h and M_t, Fig. 5 shows the regions in (M_h, M_t) parameter space corresponding to an EW vacuum that is stable (green), metastable (lifetime $\tau_{EW} > \tau_U$, yellow) or unstable (with $\tau_{EW} < \tau_U$, red). The plots in Fig. 5 show the location of these regions in a LO, NLO and NNLO calculation, from top to down. The experimental ellipses for M_h and M_t are also shown.

This figure demonstrates that NNLO precision is crucial to answer questions about the stability of the EW vacuum. What about higher order (N^{3}LO) corrections? The NNLO plot shows also (dashed lines) the remaining error, obtained by combining in quadrature the (rather small) theoretical error expected from the non-inclusion of such higher order corrections and the uncertainty from α_s: clearly a definitive answer to the stability question requires a better knowledge of M_t rather than an even more refined theoretical calculation. In terms of the top mass, the stability bound reads [47]:

$$M_t < (171.36 \pm 0.15 \pm 0.25_{\alpha_s} \pm 0.17_{M_h})\ \text{GeV} = (171.36 \pm 0.46)\ \text{GeV} \tag{10}$$

where, in the last expression, the theoretical error is combined in quadrature with the indicated experimental uncertainties from α_s and M_h.

Concerning the impact of M_t on the stability bound, there is some controversy in the literature regarding the relationship between the top mass measured at the Tevatron and LHC and the top pole-mass. Although the naive expectation would assign an error of order Λ_{QCD} to the connection between these two numbers, a more drastic proposal has been advocated in Refs. [56]: to use instead the running top mass measured through the total production cross-section $\sigma(pp/p\bar{p} \to t\bar{t} + X)$ at Tevatron and the LHC, which allows for a theoretically cleaner determination of M_t. However, this leads to a value of M_t compatible with the Tevatron and LHC values but with an error which is a factor of 4 worst: $M_t = 173.3 \pm 2.8$ GeV [56]. Of course, if one is willing to downgrade the error on M_t in this way, there would still be room for absolute stability up to M_{Pl} by moving into the lower range for M_t. Clearly, a better understanding of the theoretical errors in the top mass determination would be desirable. See Refs. [57,58] for a review of the issues involved, current status and future expectations (presumably down to $\delta M_t \sim 500$–600 MeV at the LHC) concerning this important measurement.

Fig. 6, left plot, shows the different regions for stability of the EW vacuum (NNLO), with additional information on the scale of instability, in red dashed lines. The right plot shows the same stability regions [plus the region in which $\lambda(\mu)$ becomes non-perturbative below M_{Pl}] in a zoomed-out range for M_h and M_t. This version emphasizes that we seem to be living in a very untypical region of parameter space, really close to the boundary for absolute stability, in the narrow wedge for a long-lived EW vacuum. A complementary view of the same observation is Fig. 7, which plots the different regions for the Higgs potential behavior in the $(\lambda(M_{Pl}), y_t(M_{Pl}))$ plane (as these parameters should be more fundamental). The SM location in the narrow metastability wedge is indicated by an arrow, showing once again how atypical our universe looks like.

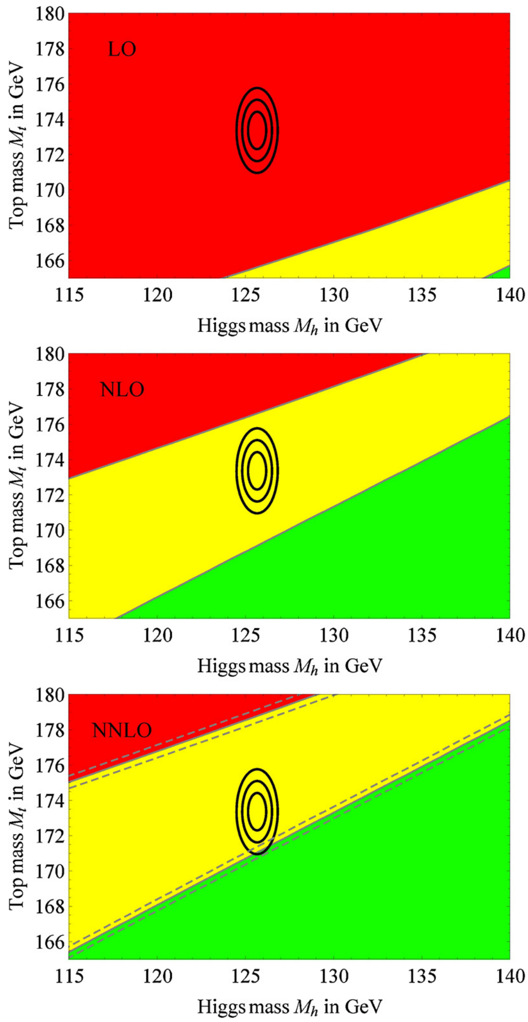

Fig. 5. (Color online.) Regions in the $(M_{\rm h}, M_{\rm t})$ parameter space corresponding to absolute stability (green), metastability with lifetime $\tau_{\rm EW} > \tau_{\rm U}$ (yellow), and instability, with $\tau_{\rm EW} < \tau_{\rm U}$, of the EW vacuum. The ellipses give the experimental values at 1, 2 and 3σ. The different versions correspond to progressively more precise calculations, from LO to NNLO as indicated.

This intriguing fact has triggered many speculations concerning its possible significance [41,46,47] including: high-scale Supersymmetry [59], enforcing $\lambda(\Lambda) = 0$ through $\tan\beta = 1$; IR fixed points of some asymptotically safe gravity [60], among other ideas (some predating the Higgs discovery [61]). Is $\lambda(M_{\rm Pl}) \simeq 0$ related to the fact that we also live very close to a second phase boundary, the one separating the EW broken and unbroken phases? This boundary is associated with the fact that the mass parameter in the Higgs potential, m^2, is extremely small in Planck units: $m^2/M_{\rm Pl}^2 \sim 0$. In this respect, it seems that the Higgs potential has a very particular form at the Planck scale, with both λ and m^2 being very small. In addition, also β_λ takes a special value $\simeq 0$ not far from $M_{\rm Pl}$. Why do EW parameters seem to take such intriguing values at the Planck scale, the scale of gravitational physics, which is totally unrelated to the EW scale? No compelling theoretical explanation has been offered so far.

3.2. Vacuum instability and physics beyond the standard model

Clearly, the intriguing "near-criticality" discussed in the previous section would be just a mirage if new BSM physics appears below $M_{\rm Pl}$ and modifies the running of $\lambda(\mu)$ significantly. On the other hand, the SM instability cannot be used by itself as a motivation for BSM, given the huge EW vacuum lifetime. Nevertheless, we do expect new physics BSM, e.g., to explain dark matter, neutrino masses or the matter–antimatter asymmetry and it is natural to ask how such physics could affect the near-criticality issue.

There are three possible impacts of BSM on the stability of the Higgs potential: *a)* it can make the stability worse; *b)* be irrelevant; or *c)* cure it. It is easy to find examples of the three options and we will take for illustration the case of type I see–saw neutrinos. In such scenario, neutrinos impact the running of $\lambda(\mu)$ through their Yukawa couplings, which scale

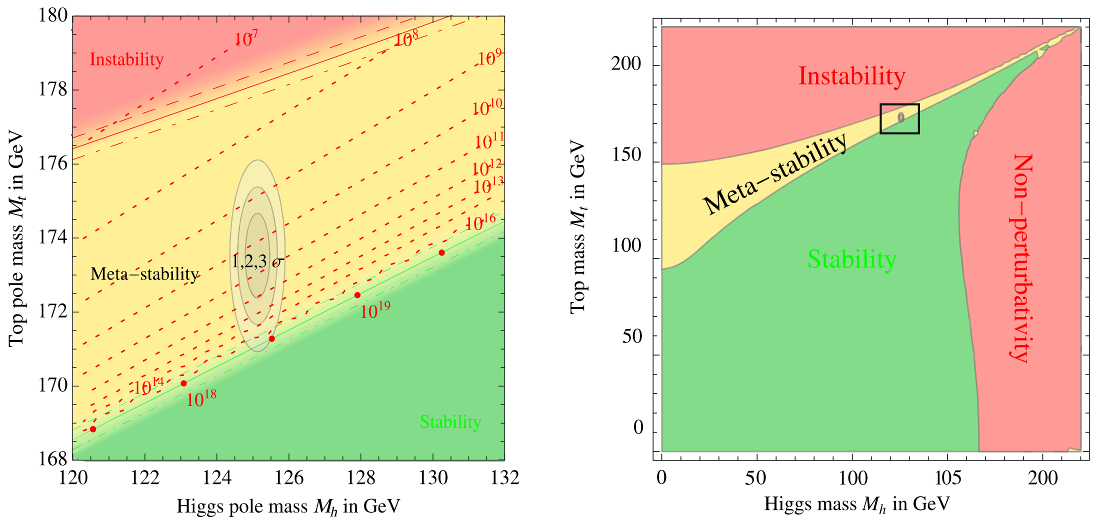

Fig. 6. (Color online.) Regions in the (M_h, M_t) parameter space corresponding to absolute stability (green), metastability with lifetime τ_{EW} longer than τ_U (yellow), and instability, with $\tau_{EW} < \tau_U$, calculated at NNLO. The ellipses give the experimental values at 1, 2 and 3 σ. The red–dashed lines in the zoomed-in version (left, from Ref. [47]) indicate the scale of instability, in GeV. The zoomed-out version (right, from Ref. [41]) also shows the region corresponding to non-perturbative Higgs quartic below M_{Pl}.

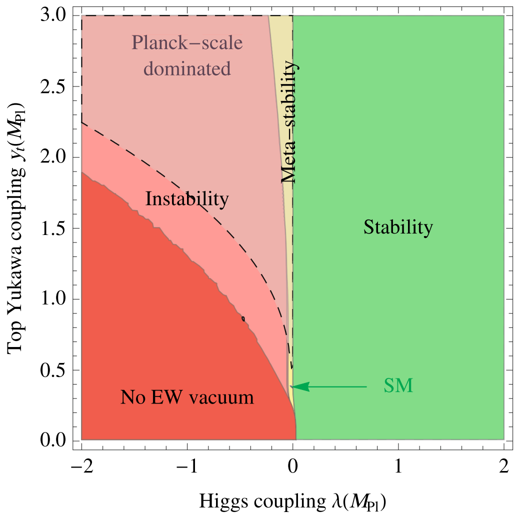

Fig. 7. Different regions for the Higgs potential behavior in the $(\lambda(M_{Pl}), y_t(M_{Pl}))$ plane. (The region labeled "Planck-scale dominated" corresponds to an instability scale beyond the Planck scale, for which (unknown) gravitational physics must play a crucial role in the potential structure.) Taken from [47].

like $y_\nu^2 \sim M_N m_\nu / v^2$, where m_ν is the mass of the lightest neutrinos, M_N the mass of the heavy right handed ones and $v = 246$ GeV is the Higgs vacuum expectation value.

a) For large enough M_N, the destabilizing effect of a large y_ν can worsen the instability, even reducing the vacuum lifetime below τ_U (if $\lambda(\mu)$ is driven below ~ -0.05). This would contradict our existence and can be used to set an upper bound on M_N, see Refs. [44,62].

b) For M_N smaller than this upper bound, of order $M_N \simeq 10^{13-14}$ GeV for $m_\nu \simeq 0$–1 eV, the neutrino Yukawas would be too small to alter significantly the running of λ and their presence would be irrelevant for the potential instability.

c) A see–saw scenario that cures the instability is easy to build using a powerful stabilization mechanism through a heavy singlet field S, with $\langle S \rangle \neq 0$, and coupled with the Higgs boson as $\lambda_{HS} S^2 |H|^2$ [63]. When S is decoupled, the low-energy λ is reduced by a negative threshold effect. The apparent instability of the potential is a mirage, as λ above the S threshold is larger than the naive SM extrapolation indicates. This mechanism is compatible with a see–saw scenario in which $M_N = \langle S \rangle$ is smaller than the SM instability scale $\sim 10^{10}$ GeV, and also satisfies the lower constraints on M_N from leptogenesis [63].

Obviously, other stabilization mechanisms exist, and almost all extensions of the SM at the TeV scale will modify the behavior (or very existence) of the Higgs field at high energies. In any case, potential stability (at least in the weak sense of demanding $\tau_{\rm EW} \gg \tau_{\rm U}$) can be used to constrain BSM models that do not guarantee (unlike Supersymmetry) a good UV behavior of the Higgs potential.

4. Conclusions

The first run of the LHC operations crowned the Standard Model as the successful description of the fundamental constituents of matter and their interactions to the tiniest details, from the QCD jet production over many orders of magnitude, to the multiple productions of electroweak gauge bosons as well as the production of top quarks. Undeniably, the Higgs boson discovery will remain the acme of the LHC run 1 and it has profound theoretical and phenomenological implications. In a few months from now, the second run of the LHC will start, with an increased center-of-mass energy, $\sqrt{s} = 13$ TeV. This run will provide us in the coming years with crucial experimental information on the physics behind the breaking of the electroweak symmetry and it carries the hopes to finally reveal the first cracks in the SM grounds. If naturalness turned out to be a good guide, the LHC should soon find new states and revolutionize the field. If we are not so lucky and such new states are too heavy for the LHC reach, we might still detect indirectly their presence through the deviations they can induce on the Higgs properties. Precise measurements of such properties are therefore crucial and could be extremely useful to guide future direct searches at higher energies, either at the LHC itself or at other future facilities.

The Higgs boson might also be a portal to a hidden sector whose existence is anticipated to account for the total matter and energy budget of the Universe. The Higgs boson could also be one key agent in driving the early exponentially growing phase of our Universe and thus allowing large scale structures to emerge from original quantum fluctuating seeds.

Whatever the outcome of the next LHC run, we are exploring new territory and living in exciting times!

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