What can the Higgs tell us about UV physics?

A. K. Knochel

Institut für Theoretische Teilchenphysik und Kosmologie, RWTH Aachen, Aachen, Germany
E-mail: knochel@physik.rwth-aachen.de

After the discovery of a Higgs-like boson with a mass of $m_h \approx 125$ GeV at the LHC, we can now attempt to draw conclusions about physics beyond the Standard Model. I argue that there are several hints towards new physics at intermediate scales $\Lambda \gtrsim 10^8$ GeV. I review a class of stringy models with intermediate scale SUSY which relate the observed Higgs mass to symmetries of the Higgs sector. I then discuss radiative corrections to $m_h$, unification, dark matter and the possibility of classically unstable UV completions in these models.

Keywords: Higgs; SUSY; Vacuum Stability; String Phenomenology

1. Introduction

The properties of the new particle [1, 2] at $m \approx 125$ GeV discovered by the ATLAS and CMS experiments at the CERN LHC match within experimental uncertainties those predicted by the minimal Higgs sector of the Standard Model [3–6]. These experimental uncertainties are already small enough to be a nontrivial test of the SM Higgs hypothesis. Alternative spin assignments as well as parity assignments are experimentally disfavored, and those couplings of the new boson to SM states which have already been measured agree with predictions closely enough to warrant the name “Higgs boson”. We therefore have an indirect (albeit model-dependent) measurement of the last unknown parameter of the minimal Standard Model - the quartic Higgs coupling. For the sake of this talk, I will assume that the new boson is indeed the SM Higgs in the sense that deviations from the SM Higgs sector are suppressed by a large new physics scale significantly above a TeV.

In absence of evidence for other new physics at the LHC8 and other colliders, and only indirect or unspecific experimental evidence from other observations, one can ask how strongly the physics at the TeV scale really deviates from the electroweak Standard Model. What is the scale at which a radical departure from the minimal SM is to be expected, and of what type is this departure? Since before the definite discovery of the new boson, it has been noted by several authors (see for example [7–

* Talk given at XXIX-th International Workshop on High Energy Physics at IHEP, Protvino.
Fig. 1. The 2-loop renormalization group flow of the Higgs quartic coupling in the Standard Model for values of the top quark mass of $m_t = 170.7, 172.9, 175$ from upper to lower line. The quartic coupling vanishes at or near the Planck scale only for very low values of the top quark mass. Around the current PDG central value, the sign change takes place at intermediate scales. 10]) that the top quark Yukawa coupling and the Higgs quartic coupling, interpreted within the SM, take very peculiar values in nature: we live very close to a “critical line” in the $m_h - m_t$ plane which separates the parameter region of absolute vacuum stability (high Higgs masses and low top masses) from a region of instability in which the SM predicts our vacuum $\langle h \rangle \sim 175$ GeV to be metastable or even unstable at cosmological time scales. Since this analysis is predicated on taking the SM Higgs sector at face value up to very high energy scales with only mild modifications, and since we still appear to be in the region of sufficient stability on cosmological time scales, no absolutely imperative conclusions can be drawn from it. However, in light of absence of evidence for other new physics at the LHC, nature’s location in the $m_h - m_t$ plane could be taken as a hint that the SM Higgs sector might remain essentially unmodified up to scales far beyond a TeV. If significant modifications of the minimal Higgs sector such as scalar singlet extensions, 2HDM (SUSY or not), or compositeness exist at the TeV scale, the stability diagram is meaningless and our position on the “would-be” critical line a mere coincidence.

When interpreted within the intermediate scale SUSY scenario which we propose [11–14], the vanishing of the quartic coupling at this mass scale can be explained from stringy symmetries of the Higgs sector, and a connection of the observed Higgs mass is established to other phenomena in nature which point towards intermediate scales of new physics such as neutrino masses (and possibly leptogenesis), axion dark matter and gauge unification. The preferred axion decay constant for dark matter is around $f_a \approx 10^{12}$ GeV, where higher values can be accommodated if the initial misalignment angle is $\theta \ll 1$, and smaller ones if there are other DM sources. These models are currently being tested, e.g. by the ADMX experiment. Gauge unification can be easily achieved in our scenario without low scale SUSY, for example in the presence of GUT breaking gauge flux in type IIB compactifications. In Fig. 2 of [13]
this coincidence of scales is nicely illustrated for a specific model. I argue that after LHC8, there should be renewed efforts to think about UV completions at intermediate scales and what they might entail for (stringy) model building, cosmology, HEP and dark matter experiments.

2. Shift symmetric Higgs sectors

Recently, we have proposed [11, 12] how the projected vanishing of the quartic coupling at intermediate scales may be connected to stringy UV completions. They predict a vanishing tree level Higgs quartic coupling at the soft breaking scale due to an approximate shift symmetry

\[ H_u \rightarrow H_u + c, \quad H_d \rightarrow H_d - c \]  

in the Higgs sector. It restricts the leading order lowest dimension Kähler potential to be of the form

\[ K \sim f(X)|H_u + \bar{H}_d|^2 \]  

where \( f(X) \) encodes the moduli dependence of the Kähler function. An immediate consequence of this is that \( \tan \beta = 1 \) and the SM Higgs doublet lies along a flat direction of the electroweak D-term

\[ V \sim \frac{g_1^2 + g_2^2}{8} (|H_u|^2 - |H_d|^2)^2 + \ldots . \]

There is therefore no quartic self coupling at tree level.

Such shift symmetries have been known for quite some time to appear in heterotic orbifold compactifications [15–18] and simpler field theoretic models [19, 20], where they are essentially a remnant of higher-dimensional gauge invariance. In [12] we argue that in type II models, analogous situations can arise not only for type IIA Wilson line Higgs sectors but also for a type IIB bulk Higgs on D7 branes. In the remainder of this talk I want to concentrate on certain field-theoretic aspects of these models, and in particular on the effective field theories below the compactification scale.

3. Radiative Corrections to the Weak Scale and the Higgs Mass

The hierarchy problem is not obviously present in the SM in regularization/renormalization schemes such as \( \overline{\text{MS}}/\overline{\text{DR}} \) or functional renormalization group

\[ \text{It was since proposed [13] to realize this situation using an approximate } \mathbb{Z}_2 \text{ parity rather than shift symmetries.} \]

\[ \text{However, one has to be careful since it depends on the details of the compactification whether a shift symmetry is actually realized in terms of the variables of the 4D Kähler potential [12]. For example, both components of the complex Wilson line moduli on D7 branes transform nonlinearly under certain gauge transformations, which would naively entail that they drop out of the Kähler potential entirely if it were shift-symmetric with respect to both.} \]
techniques which avoid the introduction of a hard scale-invariance breaking cut-off [21, 22]. It reappears when new particles coupling to the SM exist far beyond the TeV scale. It is conceivable that the hierarchy problem is somehow remedied at this high scale (in contrast to SUSY which must be present far below the high scale in order to work as a remedy), but no mechanism is presently known to us. The relevant scale in the shift symmetric SUSY scenarios is \( \Lambda_{\lambda=0}/4\pi \ll M_{Pl} \), giving us a fine tuning measure which is large but nevertheless up to \( \sim 23 \) orders of magnitude less severe than the naive cutoff-based estimate in the SM, \( M_{Pl}^2/m_W^2 \). Since we have a theory prediction for the quartic coupling, once the electroweak scale is set to the measured value, the Higgs mass is fixed as well. We now want to consider the radiative corrections to this ratio \( m_h/m_W \), i.e. to the physical Higgs mass.

The SM effective potential for the Higgs and consequently the relation between the \( MS \) quartic coupling and \( m_h \) as well as the running of the quartic coupling are well known to NNLO. We are now concerned with the corrections to the quartic coupling at the high scale of new physics. There are two main contributions: i) corrections to \( \tan \beta = 1 \) and thus to the tree level quartic coupling; ii) radiative corrections to the quartic coupling itself. The former is suppressed by one additional loop factor, but it can be log-enhanced by large hierarchies between the soft scale and the string compactification scale, and therefore competes with the 1-Loop radiative corrections.

i) **Corrections to** \( \tan \beta = 1 \) or equivalently to \( \cos 2\beta = 0 \) arise when the “shift-symmetric” Higgs mass matrix \( m_1^2 = m_2^2 = B\mu \) with an exactly massless eigenstate receives radiative corrections which destroy this degeneracy. This is generally the case if the top mass comes from \( W \sim H_u QT \) at the renormalizable level. One can control this radiative violation of shift symmetry by dialing the soft breaking parameters in order to obtain an \( O(100) \) GeV eigenvalue. However, the resulting Higgs mass matrix will generally yield \( \cos 2\beta = \epsilon \), where \( \epsilon \ll 1 \) depends on the details of our parameter choice. We can give a good estimate of its magnitude. The resulting tree level quartic coupling is [11]

\[
\delta \lambda_{SV}(m_S) \sim C \left[ \frac{g_2^2 + g_1^2}{8} \right] \left[ \frac{6\tilde{Y}_t^2}{16\pi^2} \log \left( \frac{m_S}{m_C} \right) \right]^2.
\]

where \( C \) is an \( O(1) \) constant, and \( m_S \) and \( m_C \) are the soft breaking and compactification scales.

ii) **Corrections to** \( \lambda \) at the high scale arise from loop diagrams with four external Higgs fields and heavy internal lines. We operate in the decoupling limit of the MSSM, where the masses of the extended Higgs sector and the superpartners are spread around the soft scale. In the limit \( \cos 2\beta \ll 1 \), the resulting corrections to \( \lambda \) are given by [12]

\[
\delta \lambda = \frac{3\tilde{Y}_t^4}{16\pi^2} \left[ \frac{X_i^2}{m_i^2} \left( 1 - \frac{X_i^2}{12m_i^2} \right) + 2 \log \left( \frac{m_i}{m_S} \right) \right] - \frac{1}{16\pi^2} \tilde{b}_\lambda \log \frac{m_A}{m_S} + \frac{\tilde{b}_\lambda}{16\pi^2} \left[ \log \frac{\mu}{m_S} + \frac{(r-1)(r+1)^2 + 2(r-3)r^2 \log r}{2(r-1)^3} \right].
\]
where $\tilde{b}_\lambda = \frac{1}{2}(g_1^2 - 2g_1^2 g_2^2 - 3g_2^4)$, $M_\mu = M_\mu = \frac{M_\lambda}{\mu} \equiv r$, $m_\chi \equiv \max(\mu, M_\lambda)$. Knowing these corrections allows us to define an effective SUSY scale at leading log precision,

$$m_{S}^{\text{eff}} = \left[ m_{\tilde{b}_\lambda}^{\lambda/3} M_{t}^{4s_{t}^{4}} M_{\chi}^{4\tilde{b}_\lambda/3} \right]^{1/(\tilde{b}_\lambda + 8y_t^4)}.$$  \hspace{1cm} (6)

Since $y_t \approx 1/2$ at high scales, the corrections to $m_h/m_W$ are much smaller than in TeV SUSY models. We find that they can be either positive or negative and are typically below 1 GeV unless one happens to be in a “worst-case” region. This is illustrated in Figure 2 for both types of corrections. One sees from the large sensitivity of the new physics scale to $m_t$ that a more precise experimental and theoretical determination of the $\overline{MS}$ top mass can be crucial for our understanding of UV physics.

Fig. 2. The impact of squark decoupling corrections to the quartic Higgs coupling (left) and shift/exchange symmetry violation (right) on the physical Higgs mass. The narrow dark(broad light) bands are for $X^2_{\tilde{t}} = m_{\tilde{t}}^2 (6m_{\chi}^2)$ for the decoupling contributions from top partners, and $m_C = 10^2 m_\phi (m_{\text{mssm}^2})$ for the shift symmetry violation. The top quark masses are $m_t = 175.5, 173.5, 171.5$ from upper (red) to lower (green) band. The scale $m_S$ should be understood as the effective SUSY scale.

4. UV completions with $\lambda < 0$

A universal feature of the string models we consider here is the appearance of an extended SUSY sector at some scale $m_C > m_S$. The 4D D-Term then becomes one component of an extended scalar potential. As the minimal example, we consider an $\mathcal{N} = 2$ sector, where the $D$ field is part of a triplet $\bar{P}$. The usual MSSM physics is recovered by decoupling two of these fields in an $\mathcal{N} = 1$ supersymmetric fashion. One finds that this decoupling is not exact in the presence of soft SUSY breaking.
In the simplest case, the resulting quartic potential from these effects is given by

\[ V_{\Lambda=M} = \kappa^2 \frac{m_s^2}{M^2} |H_u H_d|^2. \]  

(7)

where \( M \) is the scale of extended SUSY, and \( m_s^2 \) the soft breaking parameter. An interesting consequence is that a negative mass squared parameter will lead to a quartic (non-tachyonic!) instability. One can perform a field theoretic matching of such an “unstable” UV theory to the SM by introducing a suitable IR cutoff. This is discussed in detail in [12]. This raises interesting issues for future research. It shows that the soft scale can be in the unstable regime and therefore higher than naively expected from the Higgs mass measurement. Might hierarchies \( m_S \ll m_C \) and \( m_S \gg \text{TeV} \) be connected to vacuum stability at cosmological time scales? Does the Higgs field still prefer our false weak scale vacuum after inflation in such scenarios?

Acknowledgements

I would like to thank the organizers of the XXIX-th International Workshop on High Energy Physics in Protvino and the IHEP theory division for their hospitality.

References

[1] S. Chatrchyan et al., Phys.Lett. B716, 30 (2012).
[2] G. Aad et al., Phys.Lett. B716, 1 (2012).
[3] A. Falkowski, F. Riva and A. Urbano (2013).
[4] T. Plehn and M. Rauch, Europhys.Lett. 100, p. 11002 (2012).
[5] A. Djouadi and G. Moreau (2013).
[6] T. Corbett, O. Eboli, J. Gonzalez-Fraile and M. Gonzalez-Garcia, Phys.Rev. D87, p. 015022 (2013).
[7] J. Elias-Miro, J. R. Espinosa, G. F. Giudice, G. Isidori, A. Riotto et al., Phys.Lett. B709, 222 (2012).
[8] M. Holthausen, K. S. Lim and M. Lindner, JHEP 1202, p. 037 (2012).
[9] G. Degrassi, S. Di Vita, J. Elias-Miro, J. R. Espinosa, G. F. Giudice et al., JHEP 1208, p. 098 (2012).
[10] M. Cabrera, J. Casas and A. Delgado, Phys.Rev.Lett. 108, p. 021802 (2012).
[11] A. Hebecker, A. K. Knochel and T. Weigand, JHEP 1206, p. 093 (2012).
[12] A. Hebecker, A. K. Knochel and T. Weigand, Nucl.Phys. B874, 1 (2013).
[13] L. E. Ibanez, F. Marchesano, D. Regalado and I. Valenzuela, JHEP 1207, p. 195 (2012).
[14] L. E. Ibanez and I. Valenzuela, JHEP 1305, p. 064 (2013).
[15] G. Lopes Cardoso, D. Lüst and T. Mohaupt, Nucl.Phys. B432, 68 (1994).
[16] I. Antoniadis, E. Gava, K. Narain and T. Taylor, Nucl.Phys. B432, 187 (1994).
[17] A. Brignole, L. E. Ibanez, C. Munoz and C. Scheich, Z.Phys. C74, 157 (1997).
[18] A. Brignole, L. E. Ibanez and C. Munoz, Phys.Lett. B387, 769 (1996).
[19] K.-w. Choi, N.-y. Haba, K.-S. Jeong, K.-i. Okumura, Y. Shimizu et al., JHEP 0402, p. 037 (2004).
[20] F. Brümmer, S. Fichet, A. Hebecker and S. Kraml, JHEP 0908, p. 011 (2009).
[21] W. A. Bardeen (1995).
[22] M. Shaposhnikov and C. Wetterich, Phys.Lett. B683, 196 (2010).