Thanks to 2D and maybe even beyond:

115 GeV and 140 GeV almost Standard Model Higgs without problems

Marko B Popovic *

(Cambridge, MA, June 4, 2010)

Abstract

I address 3 key points: (1) the 2D fundamental theory as potential resolution to the 4D hierarchy and vacuum energy problems, (2) a likely QCD influence on top condensation, and (3) the two regions of preferred Higgs mass, in the vicinity of $116.5\ GeV$ and $140.5\ GeV$, with associated high energy models.

I show that the SM in 2D could simultaneously satisfy (a) complete radiative generation of the Higgs mass via top loop and (b) cancelation of the remaining leading order corrections to scalar propagator. The Higgs mass, parameterized with k=1 (2), in the leading order is $113.0 \pm 1.0\ GeV$ ($143.4 \pm 1.3\ GeV$).

I show that SM top condensation is consistent with the QCD and Higgs mediated top-anti top interactions. I predict the QCD fine structure constant up to precision better than 2% in the leading order with the mean value only 0.25% away from the world average value.

The SM driven theory at energies larger than the Higgs Mass Zero Crossing (HMZC) scale, $\sim\Lambda_{EWSB}$, necessary includes tachyonic Higgs, see Popovic 2001. Here, I map the SM physical Higgs mass to the low energy HMZC scale ($0.8 - 1.8\ TeV$). "Long lived" SM, see text, necessitates a) Higgs heavier than $137.0 \pm 1.8\ GeV$ due to the vacuum stability, b) Higgs lighter than $171 \pm 2\ GeV$ due to the perturbativity and c) Higgs lighter than $146.5 \pm 2\ GeV$ such that there is a single HMZC scale at energies smaller than the Planck mass. I present candidate $m_H = 138.1 \pm 1.8\ GeV$ for the SM valid up to an energy scale, nearly equal Planck mass, obtained from a conjecture which minimize the parameters of the Higgs potential. I show that the Minimal Supersymmetric SM is less unnatural than the SM for $m_H \leq 120.9 \pm 0.9\ GeV$. I introduce a class of models potentially *exactly* removing tachyons. I analyze class of Composite Particles Models (CPM), see Popovic 2002, where top quark is composite, composed of 3 fundamental fermions, and Higgs scalar is composite, composed of 2 fundamental fermions, with $m_H \cong \frac{2}{3} m_t = 115.4 \pm 0.9\ GeV$.

KEY WORDS: Standard Model, Higgs, Tachyon, electroweak, 2D, HMZC, HMNZ^2, top condensate, dynamical symmetry breaking, renormalization, composite top, composite Higgs, CPM, vacuum stability. (48 Pages, 12 Figures, 1 Table, 139 References)

*Please address all correspondence to:
Marko B Popovic, PhD,
115 Harvard St #3, Cambridge, MA 02139, USA
Email: mbpopovic@gmail.com ;
URL: http://web.media.mit.edu/~marko/ ,
http://walltrust.com/MarkoBPopovic/index.htm

1. Introduction

The Standard Model (SM) [1-21] of particle physics has been verified during the last four decades with an unprecedented accuracy. The Higgs scalar particle is the last ingredient that still has not been experimentally confirmed. Compared to all other particles Higgs is expected to be a very unique one.

Traditionally, Higgs is anticipated to acquire a non-zero *vacuum expectation value* (VEV) everywhere in the 4D space-time. Moreover, by coupling to itself and other particles Higgs is expected to generate its mass as well as masses for all other known massive particles. The Higgs particle is, therefore, expected to bear most of the responsibility for the mass generation in the known Universe.

Furthermore, during an early period of our Universe, Higgs VEV is expected to "break" the original electroweak symmetry mediated by the four kinds of lights, one associated with the "hypercharge" gauge symmetry and another three associated with the "weak" gauge symmetries, at an energy scale Λ_{EWSB}, corresponding to the average collision center of mass energies. As Universe cools down, non-zero Higgs VEV is effectively low-pass filtering massless bosons to the single kind of light. Hence, photon remains massless while Z and $W^{\pm}$ gauge bosons acquire masses $M_Z(\Lambda)$ and $M_W(\Lambda)$ for $\Lambda < \Lambda_{EWSB}$.

Higgs is indirectly or directly anticipated to be the root of majority of physical phenomena. Hence, it is no surprise that some popular media call it the "God's particle". In a sense Higgs is a modern times version of the *ether* idea that was, hence, just partially removed by the Einstein's Special Relativity [22].

The Large Electron Positron (LEP) particle accelerator near Geneva observed number of suspicious events [23-24] in the vicinity of $115\ GeV/c^2$, at the center of mass energies a bit above $\sqrt{s} \cong 206\ GeV/c^2$, just before the accelerator was shut down in 2000. Now, after 10 years, the LEP's successor, the Large Hadron Collider (LHC), [25] the biggest ever endeavor in the particle physics research, is again collecting the high energy data (on the order of TeV and beyond). As I discuss later on, it is quite certain that new physics phenomena, beyond current physics dogma [1-21], will be observed relatively soon.

First, I give an overview of the traditional set of problems associated with the SM Higgs scalar particle: hierarchy and vacuum energy. I then address the renormalized SM at high energies; under a few reasonable assumptions, the renormalized SM high energy behavior may imply important limits on the SM physical Higgs mass. I carefully map the physical Higgs mass with the low energy Higgs Mass Zero Crossing (HMZC) scale $\sim\Lambda_{EWSB}$ [26] at which renormalized effective Higgs mass is zero; effective Higgs particle goes from regular massive particle $m_H^2 \geq 0$ at small energies to tachyon degree of freedom $m_H^2 < 0$ at large energies. Because there are two HMZC branches per Higgs mass only branch with a correct crossing should be considered affiliated with the scale of the electroweak symmetry breaking Λ_{EWSB}. I also explain why is $\Lambda_{HMZC}\sim\Lambda_{EWSB}$ and why tachyonic Higgs within a non-zero VEV theory corresponds to an ordinary, i.e. non-tachyonic, scalar during early stages of our Universe.

In Section 2, I obtain the Higgs mass range for the Minimal Supersymmetric SM [27-29], see also [30-31], for which the MSSM is less unnatural theory than the SM at low energies.

In Section 3, I explain why two particular Higgs mass regions, in the vicinity of $113\ GeV$ and $143\ GeV$ may be favored on the theoretical grounds based on an analysis of the quantum corrections related to the 2D Higgs scalar, which might be thought of as a low energy effective composite field.

In Section 4, I show that top quark condensate formation may be consistent with interplay between the QCD gluon and Higgs mediated top interactions, which is not to be confused with the QCD driven top quark infrared fixed point, e.g. see [32]. This is good news, as this symmetry breaking defining principle, or better, symmetry breaking contributing principle, may span vast energy scales in a natural fashion.

In Section 5, I hypothesize that 4D theory may be only an effective theory, which corresponds to more fundamental 2D theory, as this can solve serious problems that particle physics faces today. This could support that either a) 4D electroweak symmetry breaking is governed by 2D electroweak symmetry breaking and 4D couplings or b) that 4D theory is just an effective realization of 2D theory where dimensionality of space-time is considered less as a premise and more as a consequence of the fundamental 2D theory.

In Section 6, I discuss the nature of electroweak phase transition at energy Λ_{EWSB} and I motivate mostly second order (continuous) quantum phase transition by sketching several simplified models. Here, in my terminology, renormalized theory does not abruptly change the parameters and degrees of freedom across the low energy scale affiliated with the second order electroweak symmetry breaking. I also explain why the hierarchy problem is actually a rather benign problem within the minimal SM with composite Higgs in 4D at energies larger than the electroweak breaking scale; i.e. theory may span across vast energy scales. I illustrate this potential high energy SM behavior with particularly interesting Higgs mass candidate, $\sim 138\ GeV$, where SM is assumed to be valid up to an composite energy scale, nearly equal Planck mass, obtained from a conjecture which minimize the parameters of the Higgs potential [33]. Moreover, I address the Planck scale adaptation of the Coleman-Weinberg conjecture. I also hypothesize on how to overcome the vacuum energy obstacle.

In Section 7, I discuss the nature of electroweak phase transition and I motivate mainly first order (discontinuous) quantum phase transition by sketching several simplified models. Here, in my terminology, renormalized theory does abruptly change the parameters and degrees of freedom across the low energy scale affiliated with the first order electroweak symmetry breaking. These models are mainly related to the top quark sector and they deal with external particle degrees of freedom within 2D and 4D space-times as well as with degrees of freedom (color, flavor etc) within the internal space. I introduce a class of models that potentially may *exactly* remove the tachyon solution at high energies. I analyze class of models, see [33], where top quark is composite, composed of 3 fundamental fermions, and Higgs scalar is composite, composed of 2 fundamental fermions, with $m_H \cong \frac{2}{3} m_t = 115.4 \pm 0.9\ GeV$.

In the conclusion, I summarize and discuss findings as well as present the best Higgs mass candidates in the vicinity of $116.5\ GeV$ and $140.5\ GeV$; these are in good agreement with the predictions in [33].

I also “prove” that LHC, since last month, is already making its mark in history by stepping distinctively outside of the region covered by well established phenomenology. Even if SM is valid description below and above the HMZC scale, never in the past have particle physicists dealt with an effective theory that includes tachyons, i.e. particles with negative effective mass squared. These are exactly the characteristic of the SM driven theory above the HMZC scale unless there is some other yet unknown physics. While tachyon theories are commonly addressed in the context of string theory and cosmology there is an *alarming lack* of literature and ongoing research effort among the rest of physics community.

In this manuscript I tried to present material with enough clarity and detail so that graduate physics students in their first years and maybe even a few advanced senior physics major students could easily follow, understand, and reproduce all major points. I apologize to reader if I failed to succeed in my goal.

2. Current state of affairs

2.1. **Problems with current model**

Traditionally, there are two main problems with the SM Higgs model: (1) *Hierarchy* (or fine tuning problem/naturalness) and (2) *Vacuum energy problem.*

Hierarchy is usually associated with the idea that there are likely two important energy scales separated by many orders of magnitude. One is the electroweak symmetry breaking scale Λ_{EW} on the order of magnitude or so of the VEV, $v_{EW} = 2.462 \cdot 10^2\ GeV$, (the low energy HMZC scale hereafter) and the other one is the Planck mass energy scale, $M_{Pl} \sim 10^{19}\ GeV$, at which quantum physics is traditionally expected to be strongly entangled with gravity. i.e. with the dynamics of the 4D curved space-time described by the General Relativity [34-35]. Therefore, the hierarchy problem is how to connect these theories at two largely separated scales within single theoretical framework, which is expressed in the spirit of the effective theory and Wilson's approach to renormalization theory [36].

Traditionally, one of the main obstacles is the presence of divergences (without the high-energy cut-off that would be infinities) that grows quadratically with energy scale. The scale-renormalized Higgs mass grows quadratically (see Appendix) and if Higgs mass in the vicinity of each of the two important scales is expected to be on the order of that energy scale the parameters of the theory might need to be *fine-tuned*; a slight change of the parameters at one scale causes large changes at the other scale (see Fig 4). Clearly if there is fine-tuning [37] without explanation then such a model is considered *unnatural*.

Vacuum energy problem is among other caused by the non-zero VEV of the Higgs scalar field traditionally expected to span the entire 4D space-time. This, however, implies a huge energy density everywhere and, hence, an enormously large space-time curvature. Therefore, the Higgs mechanism is by many orders of magnitude inconsistent with our everyday physical reality [38-39].

Similarly, if universe is described by an effective local quantum field theory down to the Planck scale, then one would expect a cosmological constant of the order of M_{Pl}^4. The measured cosmological constant is smaller than this by a factor of 10^{-120}. This discrepancy is termed "the worst theoretical prediction in the history of physics!" [40].

I address both of these problems throughout this paper.

2.2. **Standard Model at high energies as function of the Higgs mass**

The effective parameters of the SM theory may be explored at high energies according to the renormalization group flow, i.e. according to the effective theory and Wilson's approach to renormalization theory [36]. The parameters of practical importance here are gauge couplings, top Yukawa coupling, and finally, parameters defining Higgs potential energy density

$$V = -\frac{m_H^2}{4}|\Phi|^2 + \frac{\lambda}{8}|\Phi|^4 \quad (1)$$

where m_H^2 is the Higgs mass squared and λ is the Higgs scalar quartic coupling. Here, I overview previous results [26, 33] on the vacuum stability [41-48], the perturbativity [49-50] and the HMZC scale [26, 33] affiliated with the electroweak phase transition. Details of these calculations are provided in Appendix.

In parallel, two independent techniques were utilized: (1) the $\overline{MS}$ scheme [51], applied to the effective potential [52] analysis [53-55], and (2) the Euclidean hard cut-off scheme, applied to the generalized original Veltman's approach [56], subsequently confirmed by Osland and Wu [57] and completed with the logarithmic divergences by Ma [58], see Appendix.

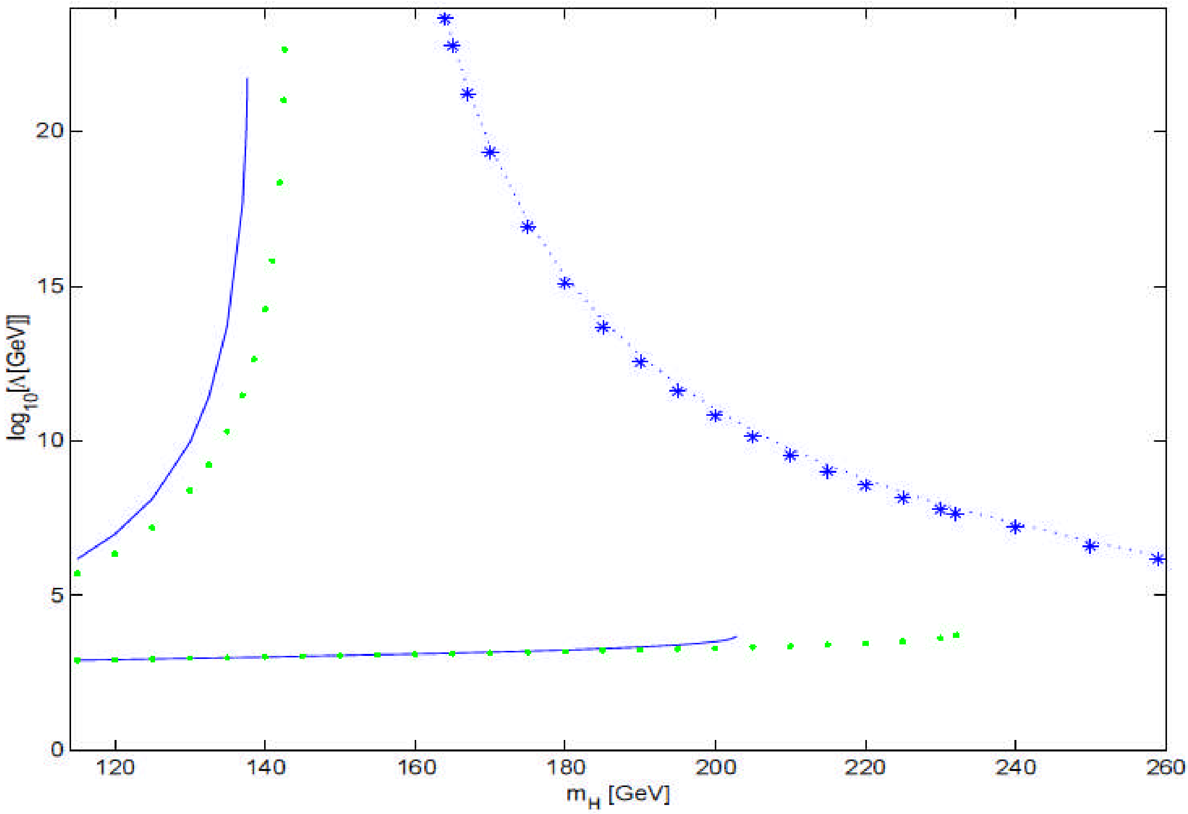


Figure 1 The SM high energy curves as function of the SM Higgs mass [26]. The vacuum stability curve (upper left corner) as obtained in the hard-cutoff method (dotted) and the $\overline{MS}$ scheme (solid). The perturbativity curve (upper right corner) as obtained in the hard-cutoff method (dotted) and the $\overline{MS}$ scheme (solid). The HMZC curve [26] (bottom line), $\sim\Lambda_{\mathrm{EWSB}}$, as obtained in the hard-cutoff method (dotted) and the $\overline{MS}$ scheme (solid).

For the Higgs mass smaller than $137.0 \pm 1.8\, GeV$ [26] there is an energy scale smaller than the Planck scale at which unacceptable deeper minima of the SM effective potential occur; this is usually referred to as the stability criteria [41-48]. For the Higgs mass larger than $171 \pm 2\, GeV$ [26] there is an energy scale smaller than the Plank scale at which the Higgs scalar quartic coupling λ reaches the Landau pole (essentially blows up) and the Higgs scalar sector becomes strongly coupled; this is usually referred to as the perturbativity criteria [49-50]. A more conservative estimate may include an additional $o(3\, GeV)$ uncertainty [26] for the Higgs mass, in response to the requirement that effective potential must be renormalization scale independent; see Appendix for details.

Finally, for the Higgs mass smaller than $203^{+14}_{-3}\, GeV$ [26] there is so-called Higgs Mass Zero-Crossing (HMZC) scale, introduced by Popovic 2001, at which renormalized effective Higgs mass is zero; effective Higgs goes from regular massive particle $m_H^2 \geq 0$ at small energies to tachyon degree of freedom $m_H^2 < 0$ at large energies. Because there are two HMZC branches per Higgs mass only branch with a correct crossing should be considered affiliated with the scale of the electroweak symmetry breaking Λ_{EWSB}.

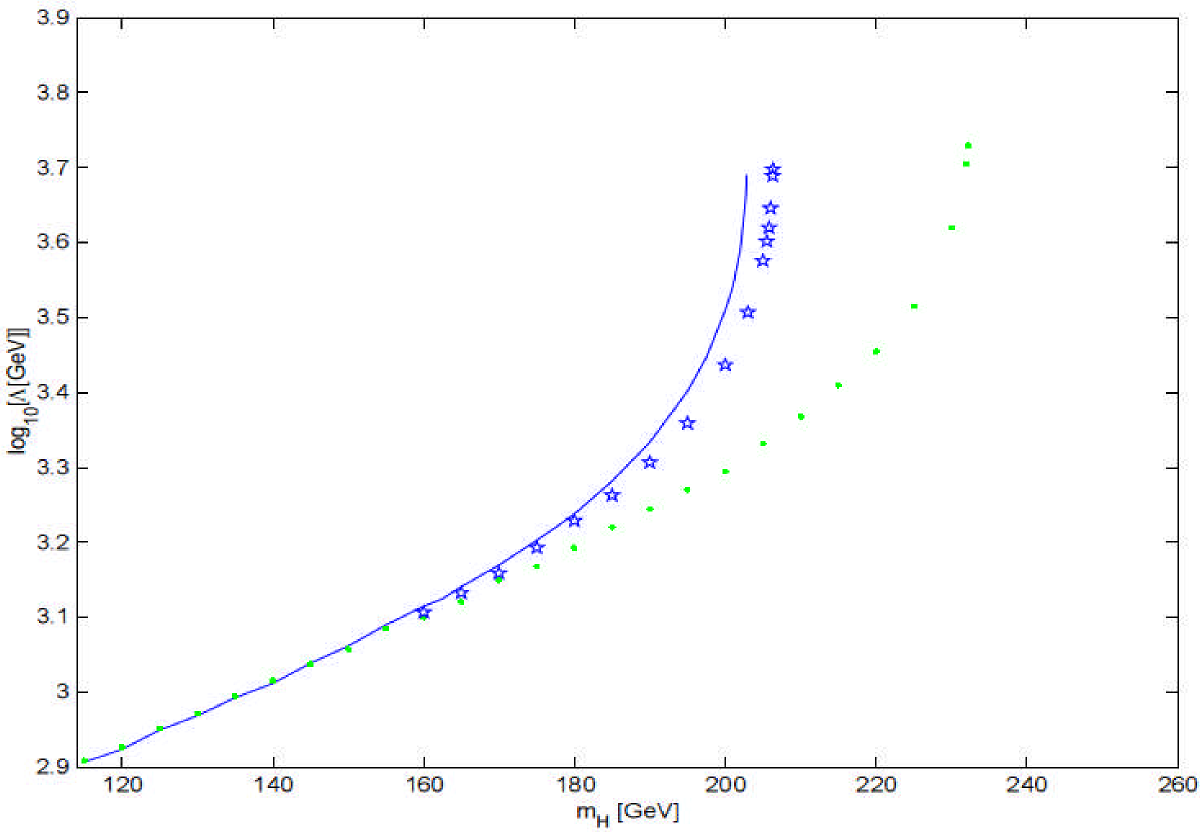


Figure 2 The Higgs Mass Zero Crossing (HMZC) curve as function of the SM physical Higgs mass in the hard-cutoff method (dotted), the $\overline{MS}$ scheme (solid) and with matching as in [50] (pentagrams) [26].

Q: Why is the low energy Λ_{HMZC} scale affiliated with the electroweak symmetry breaking Λ_{EWSB} scale?

A: The electroweak symmetry breaking is thought here as semi-classical phase transition at non-zero temperature. By preparing classical system (raising temperature, density etc) the average CM collision energies can be brought to the HMZC energy $\sim\Lambda_{\mathrm{EWSB}}$, which then creates condition for the classical phase transition, which probably happened in the very early Universe [59]. Hence, by going back in time, the vacuum structure/state of today's Universe with non-zero VEV transitions to one with zero VEV; in the zero VEV Universe the tachyonic Higgs corresponds to an ordinary, non-tachyonic particle.

The stability, perturbativity and low energy HMZC curves are shown in Fig 1-2 and Higgs mass "running" for several Higgs masses is shown in Fig 3 [26, 33]. This was obtained with $m_t = 175\,GeV$ and $\alpha_S(M_Z) = 0.1182$; see [26]. Due to variability [26] in the Higgs mass $\delta m_H[GeV] \cong 1.4\,\delta m_t[GeV] - 360\delta\alpha_S$, the current world average top quark mass from combined CDF and DØ analysis, $m_t = 173.1 \pm 1.3\,GeV$ [60], introduces a small shift and implies existence of the HMZC scale for the Higgs mass below $200.3^{+14}_{-3}\,GeV$.

The low energy HMZC branch is in the range $800\,GeV$ and $5\,TeV$ and high energy HMZC branch is larger than the Planck mass for the physical Higgs mass smaller than $146.0 \pm 2.0\,GeV$ (33). The variability in the Higgs mass, here $\delta m_H[GeV] \cong -0.24\,\delta m_t[GeV]$, introduces a small shift, and the upper limit on the physical Higgs mass is now $146.5 \pm 2.0\,GeV$. Again, a more conservative estimate may include additional $o(3\,GeV)$ uncertainty [26] for the Higgs mass, in response to the requirement that effective potential must be renormalization scale independent, see Appendix.

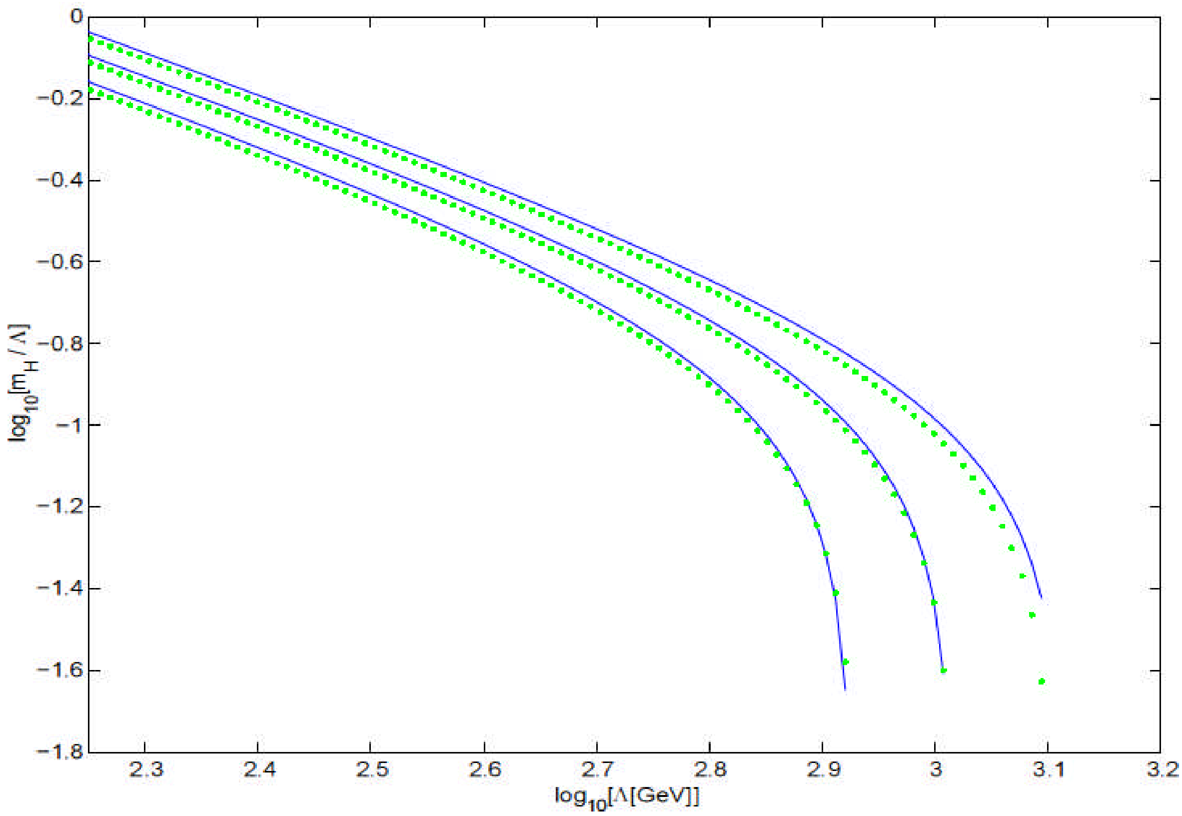


Figure 3 The logarithm of the effective Higgs mass rescaled with energy as function of the energy scale for several SM physical Higgs masses (from left to right 120, 140 and 160 GeV) in the hard-cutoff method (doted) and in the $\overline{MS}$ scheme (solid).[26]

Therefore, if the SM is a valid description of Nature all the way up to the Planck scale, where effective potential corresponds to an unbroken electroweak symmetry, then stability curve and condition requiring a single HMZC below the Planck scale limit the Higgs masses to a very tight window of roughly $142 \pm 6\ GeV$ with an electroweak phase transition scale roughly in the range $1 - 1.15\ TeV$.

However, the first order electroweak phase transition, taking place at the HMZC scale, may introduce an abrupt change of the parameters and the degrees of freedom of the theory. Therefore, there may be two very different descriptions below and above the HMZC scale that could make the above considerations, for energies larger than the low energy HMZC scale $\sim\Lambda_{EWSB}$, inappropriate.

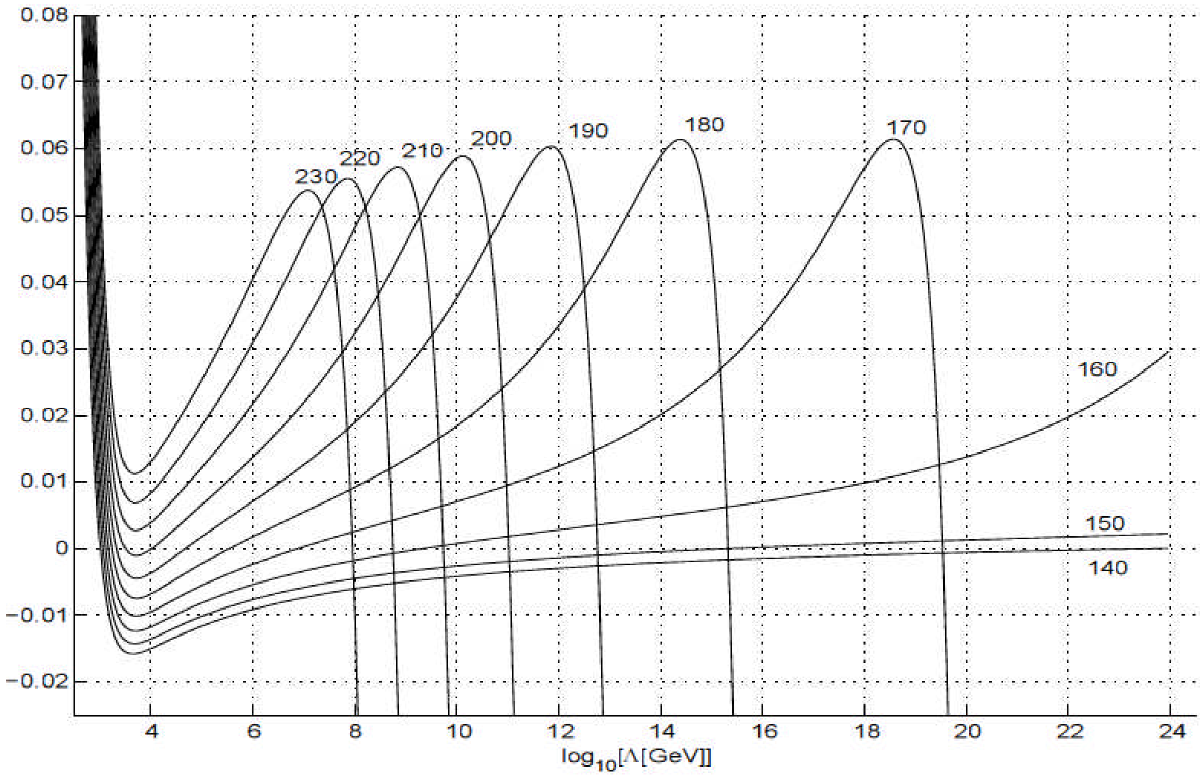


Figure 4 The SM dimensionless parameter μ, i.e. effective Higgs mass squared divided by the energy scale squared, as function of the energy scale for ten physical Higgs masses (from left to right 140, 150, 160, 170, 180, 190, 200, 210, 220 and 230 GeV) in dimensional $\overline{MS}$ regularization.[33]

2.3. <u>Two Standard Misconceptions</u>

The quantity $m_H^2(\Lambda^2)$, introduced above, is sometimes even described as a quantity that as a "matter of principle" [61] cannot be calculated. This however suggests that hierarchy/fine-tuning [37] problem cannot be quantified (i.e. ill-posed problem)? A "matter of principle", I believe, refers to the assumption that the calculation of $m_H^2(\Lambda^2)$ needs to be performed within the specific regularization scheme; therefore if one chooses two different and supposedly equally good regularization schemes, one might get two completely different and supposedly equally good answers with unclear physical significance.

The way I think about $m_H^2(\Lambda^2)$ completely disagrees with the above statement. The effective Higgs mass should communicate actual measurable physical effects at large collision CM energies. Hence, one need to make sure that, first, the regularization method is correctly used, second, the appropriate quantity is interpreted as the effective Higgs mass squared and third, the result has clear physical significance.

Results obtained in Fig 4 are identical in both the SM dimensional $\overline{MS}$ regularization and in the Veltman's hard-cutoff method, the two most popular and most reliable approaches, to a very high precision with relatively small numerical processing error. The quantity $f = \frac{dm_H^2(\Lambda)}{d\Lambda^2}$ and all higher derivatives, i.e. $f^{(n)} = d^n f/d(ln\Lambda^2)^n$, are completely expressible as polynomials in the SM couplings and can be entangled at the one loop level [48] with clear physical meaning without regularization artifacts. This is not possible at the two-loop level, but it introduces almost negligible effect [26, 33].

The interested reader should be able to easily reproduce results in Fig 2-4 that clearly show transition from regular to tachyon effective Higgs particle at HMZC energies within the LHC reach.

Another potential misconception is the interpretation of the tachyonic Higgs and classical vacuum.

The Higgs mechanism is typically addressed in the graduate physics textbooks, first, by introducing a scalar field obeying the Klein-Gordon field equation with $V = \frac{m_{scalar}^2}{4}|\Phi_{scalar}|^2$, characterized by the zero scalar VEV and with a "correct" sign for the mass term, i.e. $m_{scalar}^2 > 0$. Next step, then, is to analyze a more intricate potential, e.g. as one in Equ (1), that has a non-zero scalar VEV for $m_H^2,\ \lambda > 0$.

Standard misconception is that physics of these two theories can be directly related by renormalization; the high energy physics, with a "correct" sign, $m_{scalar}^2 > 0$, and zero VEV, can transition to the low energy physics with tachyonic $m_{scalar}^2 = -m_H^2 < 0$, corresponding to the broken electroweak symmetry.

Well, that is not exactly right! Physics of those two theories cannot be directly related by the renormalization group flow. As discussed in Appendix, it is the zero-temperature effective potential V_{eff}, Equ (A3), and not some particular values of the running effective parameters m_H^2 and λ, that defines the vacuum structure / state of our today's Universe. If minima of V_{eff} are away from zero, the electroweak symmetry is *broken* and non-zero VEV is characteristic of the effective theory *at all energy scales*.

If Higgs is considered to be a regular particle, i.e. with a positive mass squared, at low energies, then Higgs at energies larger than the HMZC scale, with a negative mass squared, must be considered to be a tachyon, the reason being the vacuum state of the world we live in. As discussed, in the zero VEV Universe, i.e. very early in the history, the tachyonic Higgs is just an ordinary, non-tachyonic particle.

The assumption contrived here is that the LHC experiment will not change the vacuum structure / state of the theory. As generally accepted, changing the vacuum state of the theory in practice may not be a particularly wise thing to do; catastrophic false vacuum scenario has been addressed by Coleman [62] and Callan and Coleman [63]. Many authors also addressed the metastable vacuum [64-73].

2.4. Preferred Higgs mass for Minimal Supersymmetry

The HMZC scale can distinguish more meaningful from less meaningful Minimal Supersymmetric Standard Model (MSSM) [27-29], see also [30-31], at low energies. By using the approximate relation [74-75] for the radiatively corrected Higgs mass in the MSSM, $m_H^2 \le M_Z^2 + \frac{3G_F}{\sqrt{2}\pi^2} m_t^4 ln\left(\frac{m_T^2}{m_t^2}\right)$, one finds that the MSSM decoupling scale, defined here by the stop mass, m_T , is smaller than the HMZC scale for Higgs lighter than $m_H \le 127.0\ GeV$ [26] for top quark mass $m_t = 175\ GeV$; this translates to $125.9\ GeV$ for the current world average $m_t = 173.1\ GeV$ [60] and $124.9\ GeV$ for $m_t = 171.4\ GeV$ [76]. This result is in agreement with findings in [77-79] with the delimiting scale approximately $900\ GeV$ (in agreement with the standard $1\ TeV$ scale [80]). However, more "pedantic" result [74-75], for the radiatively corrected Higgs mass in the MSSM, $m_H^2 \le M_Z^2 + \frac{3g_W^2 M_Z^4}{16\pi^2 M_W^2}\left[\frac{2m_t^4 - m_t^2 M_Z^2}{M_Z^4} ln\left(\frac{m_T^2}{m_t^2}\right) + \frac{m_t^2}{3M_Z^2}\right]$, yields even tighter delimiting scale, $\cong 860\ GeV$ GeV, corresponding to the $m_H \le 122.0\ GeV$[26] for top quark mass $m_t = 175\ GeV$; which translates to $120.9\ GeV$ for $m_t = 173.1\ GeV$ and $120.0\ GeV$ for $m_t = 171.4\ GeV$. This gives a combined estimate of $m_H \le 120.9 \pm 0.9\ GeV$. Again, a more conservative estimate may include additional $o(3\ GeV)$ uncertainty [26] for the Higgs mass, see Appendix.

Also, for $m_H \cong 160.0\ GeV$ the MSSM decoupling scale becomes ~10 times larger than the HMZC scale [26] corresponding to theory that is ~ 100 times more finely tuned, i.e. unnatural than SM.

2.5. SM Higgs Mass: Direct Searches and Indirect limits from the electroweak precision data

Back in 2000, based on LEP2 data, ALEPH reported an excess of about three standard deviations, suggesting the production of a SM Higgs boson with mass $\sim 115\ GeV$ [81]. The combined analysis by ALEPH, DELPHI, L3, and OPAL could not either confirm or exclude the $\sim 115\ GeV$ Higgs; instead that analysis placed a current 95% C.L. lower bound of $114.4\ GeV$ for the mass of the SM Higgs boson based on direct searches [82]; see Particle Data Group review [83] for related references.

A global fit to the precision electroweak data, accumulated in the last decade at LEP, SLC, Tevatron and elsewhere [84], gives $m_H = 76 + 33 - 24\ GeV$, or $m_H < 144\ GeV$ at 95% C.L. [84].

However if the direct LEP search limit of $m_H > 114.4\ GeV$ is taken into account, an upper limit of $m_H < 182\ GeV$ at 95% C.L. is obtained for the SM Higgs mass [83].

Therefore, global fit and direct searches suggest the SM Higgs within the HMZC regime, see Fig 4.

3. 4D SM Higgs mass from 2D considerations

In this Section, I explain why two particular Higgs mass regions, centered at $113\ GeV$ and $143\ GeV$, may be favored on the theoretical grounds based on an analysis of the leading quantum corrections in 2D.

At energies smaller that the low energy HMZC scale, the electroweak g_Y, g_W and top quark Yukawa g_t couplings' "running" is very slow compared to the "running" of the dimensionless mass parameter

$\mu = {m_H^2(\Lambda)}/{\Lambda^2}$. Also, otherwise infinite 4D scalar quadratic divergences are finite integrals in 2D. I use both of these properties to hypothesize on the physical Higgs mass.

The following analysis is independent of the physics at energies larger than the HMZC scale and assumes that there is no "new physics" beyond the SM at energies smaller than the HMZC scale.

Consider top loop contribution to the SM Higgs scalar propagator with the SM particles: Higgs, Z and W bosons and top quark, in the loops. For details regarding the loop calculation in 2D see Appendix.

I split the top loop to two, x and y, contributions. Part proportional to x exactly cancels the Higgs and gauge boson loops while piece proportional to y equals radiatively generated Higgs mass, see Fig 5.

Therefore, this is analogous to absence of the bare mass term; Higgs mass is explicitly generated by the top loop contribution. As I discuss later, cancelation of leading "divergences" may point out to important relationships between the physical quantities. Here, I assume that x and y add to z equals 1 or $\frac{2m_t^2}{v_{EW}^2}$.

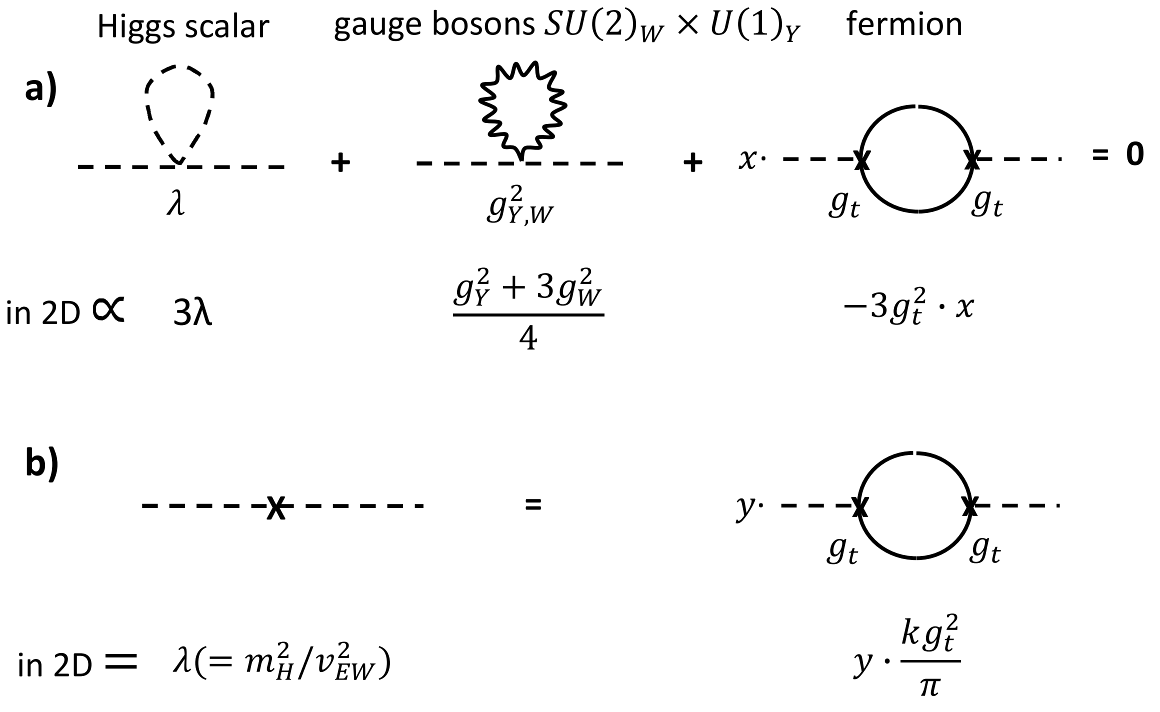


Figure 5 2D SM a) cancelation of leading quantum corrections and b) radiatively generated Higgs mass.

Top loop in 2D is exactly solved as

$$top\ loop = \frac{kg_t^2}{\pi}\ , \quad (2)$$

result that may be compared with the famous Schwinger result for 2D massive QED with photon vacuum polarization generating mass $\frac{e}{\sqrt{\pi}}$ for the gauge boson [85]; see also [86] and references therein. Factor of 2 is expected here, due to the scalar nature of interaction, as top spins may point inward or outward. However, I leave an explicit dependence on the relevant phase space parameterized with $k = 1\ (2)$.

This is a system with three unknowns x, y and λ, and three equations,

$$3\lambda + \frac{g_Y^2+3g_W^2}{4} = 3g_t^2 x \ , \quad x + y = z \quad \text{and} \quad \lambda = \frac{kg_t^2}{\pi} y \ , \tag{3 a,b,c}$$

leading to an unique solution

$$3\lambda + \frac{g_Y^2+3g_W^2}{4} = 3g_t^2\left(z - \frac{\lambda\pi}{kg_t^2}\right) \rightarrow \sqrt{\lambda} = \sqrt{\frac{g_t^2\, z - \frac{g_Y^2+3g_W^2}{12}}{1+\frac{\pi}{k}}}\ , \tag{4}$$

or in terms of the physical Higgs mass

$$m_H = \sqrt{\frac{6m_t^2\, z - M_Z^2 - 2M_W^2}{3\left(1+\frac{\pi}{k}\right)}}\ . \tag{5}$$

In the range of the top quark mass, $m_t = 170 - 175\ GeV$, the above result varies as

$$m_H = \begin{cases} 110.7 - 114.4\ \ (107.7 - 115.0)\ GeV\ for\ \ z = 1\ or\ (\frac{2m_t^2}{v_{EW}^2}),\ k = 1 \\ 140.5 - 145.2\ (136.7 - 146.0)\ GeV\ for\ z = 1\ or\ (\frac{2m_t^2}{v_{EW}^2}),\ k = 2 \end{cases} . \tag{6}$$

For the world average top quark mass, $m_t = 173.1 \pm 1.3\ GeV$ [60], I obtain

$$m_H = \begin{cases} 113.0\ (112.3) GeV\ for\ z = 1\ or \left(\frac{2m_t^2}{v_{EW}^2}\right),\ k = 1\ \rightarrow\ y = \mathbf{0.669} \\ 143.4\ (142.5) GeV\ for\ z = 1 or \left(\frac{2m_t^2}{v_{EW}^2}\right),\ k = 2\ \rightarrow\ y = \mathbf{0.539} \end{cases} . \tag{7a}$$

Whereas for the several years older value, $m_t = 171.4 \pm 2.1\ GeV$ [76], I obtain

$$m_H = \begin{cases} 111.9\ (110.0) GeV\ for\ z = 1\ or \left(\frac{2m_t^2}{v_{EW}^2}\right),\ k = 1\ \rightarrow\ y = \mathbf{0.669} \\ 142.0\ (139.6) GeV\ for\ z = 1 or \left(\frac{2m_t^2}{v_{EW}^2}\right),\ k = 2\ \rightarrow\ y = \mathbf{0.539} \end{cases} . \tag{7b}$$

The above results should be corrected by the effects of the "running" electroweak g_Y, g_W and top quark Yukawa g_t couplings between m_H and $\Lambda_{HMZC} \cong 1\,TeV$. This will be addressed in detail elsewhere; preliminary analysis suggests that this correction is on the order of maximally a few percents.

There is an additional uncertainty for the physical Higgs mass, due to the finite cut-off; the top loop is exactly solved in Equ (2) with an infinite cut-off. For example, for the k=1 branch, the finite cut-off scale equals the HZMC scale $\sim 10^{2.9} \cong 800\ GeV$ [26] at which the effective Higgs mass is zero for the physical

Higgs mass $m_H = 115.0\,GeV$. Therefore, the cut-off effects are $\leq {m_H^2}/{\Lambda_{cut-off}^2}$ and may introduce additional uncertainty on the order of $\frac{3}{8\pi^2}2\% \cong 0.01\%$ for the physical Higgs mass obtained above.

Hence, the k=1 mass branch embraces the late LEP Higgs signal candidate [23-24, 81-83]. If the SM smoothly expands above the HZMC scale, within the second order phase transition, there is certainly another cut-off scale anywhere between $\sim 1\,TeV$ and ~$500\,TeV$ due to the vacuum stability limit [26].

The k=2 branch has the HZMC scale at $\sim 10^{3.03} \cong 1100\,GeV$ for the physical Higgs mass $m_H = 142.0\,GeV$. It is worth noting that the k=2 mass branch masses are in the center of the Higgs mass range favored by the combined electroweak precision data global fit and direct searches [83-84], see Section 2.5. If the SM smoothly expands above the HZMC scale, within the second order phase transition, there is, interestingly, no stability and perturbativity constraining scales below the Planck mass. Furthermore, both SM dimensionless renormalized parameters, λ and $= {m_H^2(\Lambda)}/{\Lambda^2}$, are independently approaching zero in the close vicinity of the Planck mass energy scale. Their deviations from zero are smallest for $m_H = 137.6\,GeV$ [33] and $m_t = 175\,GeV$, see Fig 8; this value is slightly shifted to $m_H = 138.1\,GeV$ due to corrections associated with the current top quark mass world average, $m_t = 173.1 \pm 1.3\,GeV$.

4. Top condensation consistent with gluon and Higgs scalar mediation

Here, I briefly introduce the dynamical mass generation involving top condensation; see for example [87] for more complete review. I also investigate whether the SM $\bar{t}t$ channel is repulsive or attractive as a necessary SM condition for, as anticipated here, an almost-loose bound state.

4.1. <u>Top condensation and "new physics" model building: brief overview</u>

The concept of the dynamical mass generation and spontaneous symmetry breaking, potentially explaining the electroweak symmetry breaking in the particle physics, was built upon the pioneering work on the "microscopic" theory of superconductivity by Bardeen, Cooper and Schrieffer [88]. This concept was revisited, further advanced, and introduced in the high energy particle physics by Nambu and Jona-Lasinio [89-91] as the NJL model. The NJL model was further advanced by Hill [92], Miransky et al. [93] and Bardeen et al. [94]. The assumption there is that strong effective 4-fermion interactions may trigger the top quark condensation, hence, introducing the composite effective Higgs scalar that has exactly the right quantum numbers to break the electroweak symmetry in a dynamical manner. In a difference to the Technicolor [95-96, 37] models where technifermions condense, it is the SM top quark degrees of freedom here that are anticipated to be responsible for the electroweak symmetry breaking.

It was also shown that there is no fundamental theoretical obstacle that would prohibit the composite effective Higgs particle to completely mimic the SM fundamental Higgs particle at low energies [97].

The minimal model, which attempted to incorporate the SM, as an effective low energy theory, was proposed by Bardeen et al. [94]. However, this model predicted too large top quark mass in the close vicinity of the SM renormalization group, QCD driven, top quark infrared fixed point $\sim 230\,GeV$ [92, 94].

It was observed that smaller top quark masses generally could not provide enough electroweak breaking VEV to create the appropriate masses for the W and Z bosons. That can be also verified, for example, with the Pagels Stokar relationship [98] or with gap equations in the gauged NJL model [99-100] where one obtains the f_t, i.e. the top analog of the pion decay constant, f_π, that seems to be too small. The observed mismatch between the SM fermion and boson masses motivated the Topcolor model [101], with a new strong interactions singling out the top quark, as well as the class of models which combined Topcolor with Technicolor within a model building effort termed the Topcolor assisted Technicolor, TC^2 [102-110]. Another related approach, the Topcolor Seesaw [111-113] models applied a seesaw type of mixing among the "new" fermions, either the weak singlets [111-112] or the weak doublets [113], with a goal to lower the dynamically generated quark mass. Finally, some of the model building efforts attempted to incorporate the physics of extra dimensions with the top condensation [114-115].

Without experimental data that may directly confirm or reject particular theoretical concepts, the majority of above models should be considered as quite attractive and viable possibilities, though most likely highly incomplete. For example, they may require additional model structure to generate the realistic particle mass spectrum. On those line, experience with the Extended Technicolor [116-117] was that a lot of thought has to be given to the flavor changing neutral currents [116-117], unwanted contributions to R_b [118], excessive isospin violation [119] etc.

4.2. Top condensation within the Standard Model considerations

Here, I assume that there is an underlying dynamics that correlates $\bar{t}t$ values and orientations among different space-time points as well as among different momentum eigenstates. However, I assume that this underlying dynamics is completely (or almost completely) expressible with the SM degrees of freedom, thought here to represent the low energy effective theory in Wilson's approach [36].

Hence, I first investigate whether SM $\bar{t}t$ channel is repulsive or attractive as a necessary condition for an almost-loose bound state. I postpone more advanced analysis to Sections 6 -7. Physics presented here is different from analysis that linked the top quark mass with the QCD driven infrared fixed point [92, 94].

Consider $\bar{t}t$ scattering in the Euclidean space and ignore chiralities of the incoming and outgoing particles while assuming that left and right handed tops are equally represented within particle and antiparticle solution. The main interaction channels are gluon and Higgs exchange with identical number of topologically different Feynman diagrams. The weak interactions are absent as interacting particles have opposite chiralities and the hypercharge interactions are zero due to the equal sharing conjecture.

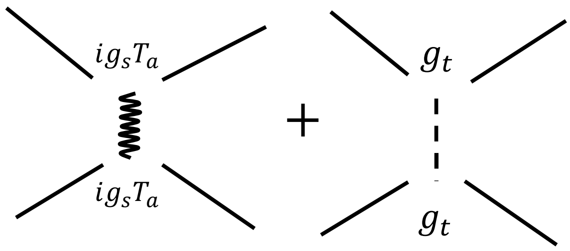
I now assume that strong QCD interactions proportional to $-g_{QCD}^2\, T_{aij}T_{akl}$, where $a = 1 \dots 8,\ i,j = 1,2,3$ and summation over repeated indices is implied, are exactly balanced with the Yukawa forces due to the virtual Higgs particle exchange proportional to g_t^2 as condition for the loose bound state; see Fig 6. Hence, the back of the envelope calculation suggests

$$2\frac{g_{QCD}^2}{6} = \frac{g_t^2}{2} \quad or \quad \alpha_S = \frac{3}{2}\frac{g_t^2}{4\pi} \tag{8}$$

where I used $(T^a)_{ij}(T^a)_{kl} = \frac{1}{2}\left(\delta_{il}\delta_{kj} - \frac{1}{N}\delta_{ij}\delta_{kl}\right)$ for $SU(N)$ groups, see for example [120], where $N = 3$.

Result in Equ (8) is in an excellent agreement with the standard estimate of the strong running coupling constant [83, 121]. The above result predicts $\alpha_s = 0.1181 \pm 0.0018$ for the world average top quark mass $m_t = 173.1 \pm 1.3\ GeV$ [60] where uncertainty is therefore solely due to the top quark mass. This value is to be compared with the current world average value $\alpha_s \cong 0.1184 \pm 0.0007$ at $s = M_Z^2$ [83, 121].

Even if equal distribution assumption is ignored and hypercharge interactions are taken into account that lowers the above result by only order $\frac{3}{2} Y_L Y_R \frac{g_Y^2}{g_t^2} = \frac{3}{2}\frac{1}{6}\frac{2}{3}\frac{g_Y^2}{g_t^2} \sim 2\%$. And if this difference is to be accounted for example by difference between $\alpha_s(M_Z^2)$ and $\alpha_s(m_H^2)$ then even in the next to leading order correction it appears that lighter Higgs, e.g. in the vicinity of $m_H \sim 115\ GeV$ with $\alpha_s(m_H^2 = 115\ GeV) \sim 0.105$, would be more preferred than the heavier one, $m_H \sim 145\ GeV$. The lighter Higgs tends to be more compatible with the first order quantum electroweak phase transition, see Section 7.



$$\frac{1}{3} g_s^2 \times 2 \qquad -g_t^2 \qquad = \mathbf{0} \rightarrow \alpha_s = \frac{3}{2}\alpha_t = \frac{3 g_t^2}{8\pi} = 0.1181 \pm 0.0018$$

World average $\alpha_s \cong 0.1184 \pm 0.0007$

Figure 6 Finely balanced interplay between the QCD gluon and Higgs scalar mediated top anti-top interactions.

If above observations are correct then interplay between logarithmically running top Yukawa coupling constant and logarithmically running strong coupling constants may indeed define the low energy HZMC scale, at which the top condensate forms and breaks the electroweak symmetry within second (first) order phase transition as addressed in Section 6 (7). Anyhow, this is good news as this symmetry breaking principle, or maybe only contributing principle, may span vast energies in a natural fashion.

Finally this interplay may be completely local and bound to the finite volumes surrounding propagating particles [33]. For the loose bound state it takes small or no energy to locally order the background condensate field. If this is a local, dynamical process it may clearly resolves the *vacuum energy problem*.

Could solution for hierarchy and vacuum energy problems come from the fundamental 2D theory?

5. The unbearably speculative lightness of being 2 dimensional

If Higgs mechanism [12-21] is confined to the propagator 2D space-time then the non-zero VEV is not required to span the entire 4D space-time. That would clearly eradicate both the hierarchy and vacuum energy problems. In 2D, the leading quantum corrections are only logarithmically divergent and the scalar VEV, confined to propagator 2D space associated with the propagating particles in 4D would imply only a reasonably large space-time curvature. Hence, the electroweak symmetry breaking defined by 2D propagator physics and 4D couplings might be an attractive option. If the HMZC scale is rather small then the theory does not have to be substantially fine tuned. I addressed these ideas in the past [33].

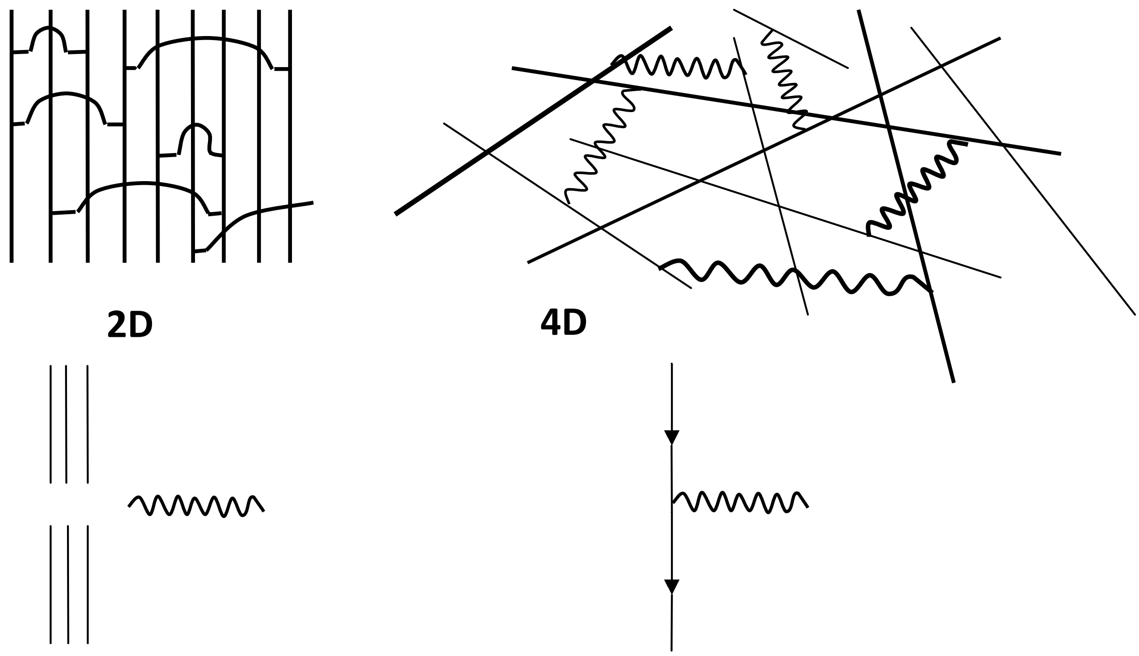


Figure 7 Generating 4D theory out of more fundamental 2D theory.

Extending these lines of thoughts even further, it would be an interesting option to interpret the 4D notion of our ordinary space-time not as a premise but rather as consequence of a more fundamental theory in 2D; see Fig 7. Again, this radical concept would render vacuum energy and hierarchy problems non-existent and it could have a deep impact on the current notion of gravity [34-35].

If 2D fermions have internal degrees of freedom that in combination with external degrees of freedom transform in the "right" way under the 4D Poincaré group, i.e. the inhomogeneous Lorentz group, then there is no *apriori* reason why this theory cannot be interpreted as 4D. For example, the internal degrees of freedom would contribute to the 4D 4-momentum which may or may not be "on shell".

To illustrate this idea I sketch the simplest possible example. Consider that 2D fermions are described by two "flavors", A and B. Furthermore, imagine that 4D fermion (dimension 3/2) may be constructed out of 4D scalar field (dimension 1) and 2D fermion (dimension ½). Moreover, assume that 4D scalar may be interpreted as 2D condensate composed of left and right moving A or B or their linear combination. On these lines, the 4D vector boson, e.g. the transversal spin 1 component, may be interpreted as linear combination of two 2D fermions moving in the same direction. And the 2D vertex may appear as 4D if there is additional phase space attached to the interacting 2D particle. For example that phase space may be a consequence of one $SU(2)_A$ or $SU(2)_B$ symmetry that may be related to $SU(2)_L$ or $SU(2)_R$; after all the Lorentz group is related to $SL(2,C)$, which is $SU(2)_L \times SU(2)_R$, see [122]. Finally, Higgs mechanism may be confined to 2D and described by the non-zero condensate VEV.

As discussed in Section 4, condensation is most likely closely related to the interplay of "new physics", e.g. 4 fermion interactions which are *renormalizable* interactions in 2D, with QCD. Furthermore, as discussed in Section 7.1, one may naturally expect a bound state built out of three fermions to balance electroweak gauge bosons' and Higgs loops at the energies smaller than the electroweak symmetry breaking scale while potentially *exactly* removing the tachyon solution at all larger energies.

Interpretation that space-time is 2D and that it only appears being 4D where final dimensionality enters more as a consequence of the original 2D theory than as premise is certainly intriguing. Details of this model and similar models are worth investigating; I will address them elsewhere.

Why is this concept extremely fascinating? Well, apart from almost shocking idea that our ordinary 4D space-time might be compactly "written" in 2D, this theory could have many interesting implications on physics. As previously discussed [33] this may completely remove the vacuum energy problem as only negligibly small portion of the 4D space-time has actual non-zero VEV and it could make the hierarchy problem benign as 4D quadratic divergences are exchanged with 2D logarithmic running that can naturally span vast energy scales. Finally, it will be very interesting to investigate what may be the consequences on the Einstein's notion of gravity [34-35].

Therefore I hypothesize that 4D theory may be only an effective theory which corresponds to more fundamental 2D theory as this can solve hierarchy and vacuum energy problems that particle physics faces today. This could be to the extent that (1) 4D electroweak symmetry breaking is governed by 2D electroweak symmetry breaking and 4D couplings or (2) that 4D theory is effective theory completely described by 2D theory, where dimensionality of space-time enters less as a premise and more as a consequence of the fundamental 2D theory. In 2D leading quantum corrections are only logarithmically divergent and scalar VEV may be confined only to propagator 2D space associated with propagating particles in 4D; i.e. equivalent to compactifying the Higgs *ether* from entirety of 4D to just a small subset of that space. Complete removal of the Higgs *ether* is dynamical symmetry breaking as discussed above.

6. Higgs at very high energies and second order phase transition

In this section I hypothesize on physics that may be responsible for the $142 \pm 6\ GeV$ masses, as obtained in the Section 2, or k=2 Higgs mass branch $137 - 146\ GeV$, as obtained in Section 3. I show that the heavier Higgs solution is likely associated with the "desert scenario", or "long lived" SM, within the second order electroweak phase transition. Here, in my terminology, renormalized effective theory does not abruptly change the parameters and degrees of freedom across the low energy HMZC scale affiliated with the second order electroweak symmetry breaking.

As I explain next, it is possible that the SM with composite Higgs in 4D may span vast energy scales. Then, the hierarchy problem seems to be a rather benign problem for $m_H^2(\Lambda^2) < 0$.

6.1. Two special solutions affiliated with the Planck mass scale

Within the entire high-energy SM effective theory spectrum there is a single region where both dimensionless parameters $\mu \equiv m_H^2(\Lambda^2)/\Lambda^2$ and λ almost coincide with zero value, see Fig 8. Interestingly enough this is in the vicinity of the Planck mass, obtained as a consequence and not as a premise. This solution to conjecture that minimize the parameters of the Higgs potential is obtained for the physical Higgs mass centered at $m_H = 137.6\ GeV$ [33]; this result, however, is slightly shifted to $m_H = 138.1\ GeV$ to accommodate for the current top quark mass world average. This solution overlaps with both the $142 \pm 6\ GeV$ Higgs mass range, obtained in Section 2, and the k=2 mass branch, obtained in Section3.

While it is traditionally anticipated that μ should run quadratically, the actual SM μ for the $m_H = 138\ GeV$ solution, Fig 8, runs logarithmically at energies larger than the HMZC scale. Moreover, already based on visual inspection, it appears that SM dimensionless parameters, λ and μ, are directly related.

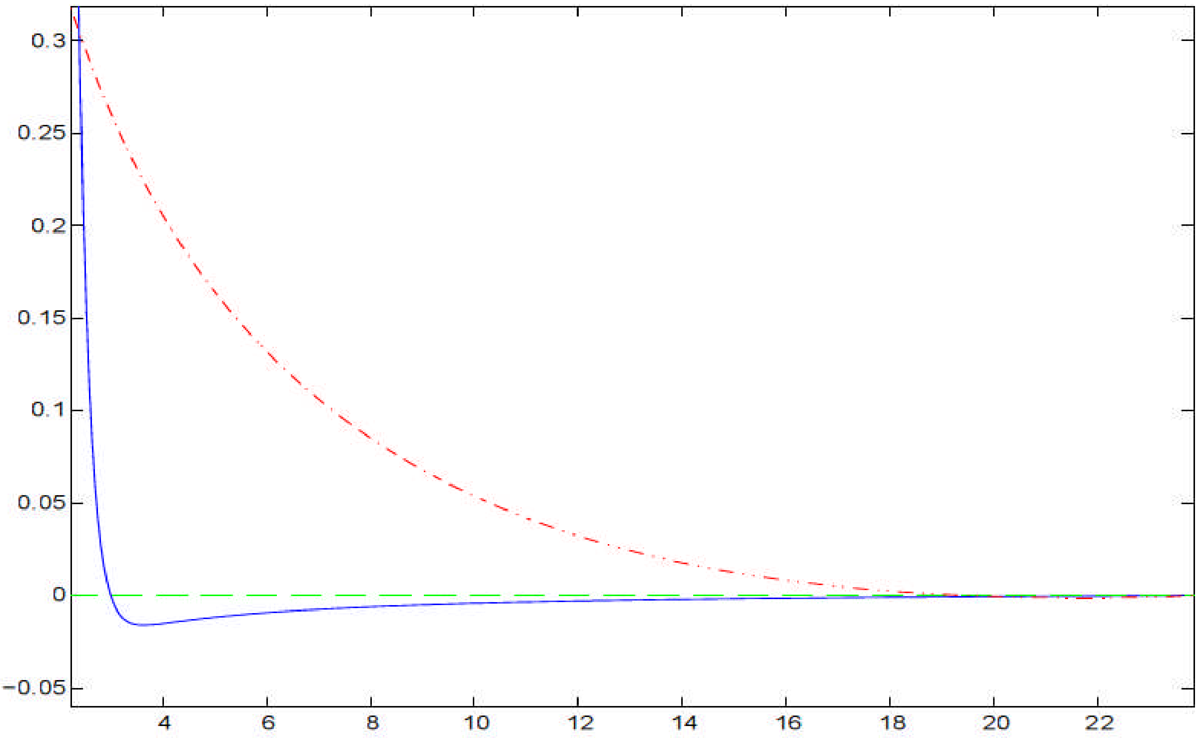


Figure 8 The running of the SM Higgs quartic coupling λ and dimensionless parameter $\mu \equiv m_H^2(\Lambda^2)/\Lambda^2$, i.e. rescaled effective mass squared, for the SM Higgs candidate, physical mass $m_H \cong 138\ GeV$, for the physics originating at the very high energy corresponding to roughly the Planck mass [33].

The parameter μ at the Planck scale is exactly equal zero for $m_H \approx 146.5\ GeV$; this is the Planck mass adaptation of the Coleman-Weinberg (CW) conjecture where bare mass (though and not $m_H(M_{Pl})$ as here) is zero and the electroweak breaking is governed by the quantum corrections [52]; see[26, 33].

6.2. Composite Higgs from very high energies

As I show next, the *slow* “running” or better “walking” should be expected from the composite Higgs for positive λ and negative $m_H^2(\Lambda^2)$. If one ignores the higher order corrections the running has simple functional dependence supporting a “long lived” solution

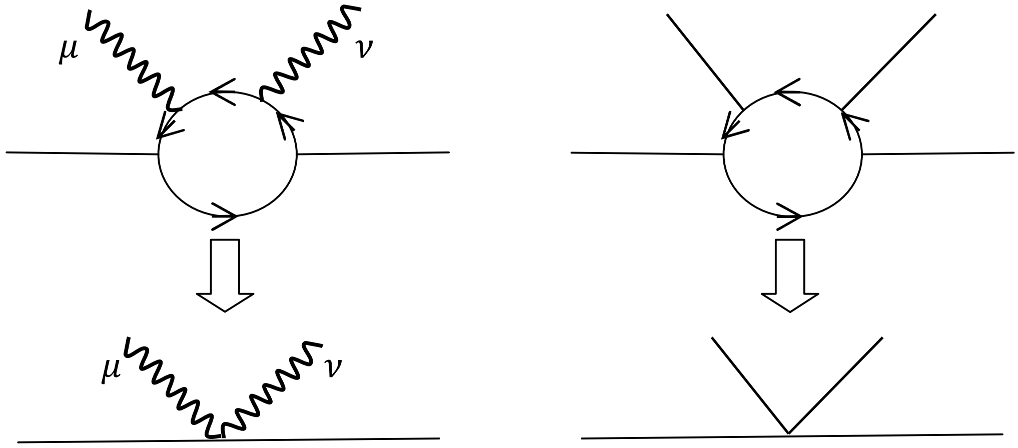
$$\mu \propto -\sqrt{\lambda}\ . \tag{9}$$

Main idea here is that Higgs field might be an “almost” fundamental field generated in the proximity of the Planck scale or some other high energy scale Λ_{high}, i.e. effective field composed out of fermion degrees of freedom with assumed zero potential energy density. Composite Higgs in the context of top condensate electroweak breaking has been addressed in the past [94]. Also, it has been shown that theory with composite Higgs may indeed mimic the minimal SM with fundamental Higgs scalar field at low energies [97]. However, as I show here, there are still a few constraining features which if not satisfied may tell apart elemental from the composite Higgs particle.

Just beneath the high energy scale Λ_{high} the Higgs scalar acquires correct couplings to gauge boson fields through field renormalization and top Yukawa coupling renormalization, see Fig 9, such that

$$-I(\Lambda) g_{t*}^2 {Z_\Phi}^{-1} = 1 \tag{10}$$

where g_{t*} is the bare top Yukawa coupling and top loop integral equals $-I(\Lambda) g^{\mu\nu}$ where $\Lambda \le \Lambda_{High}$.



$$\boldsymbol{\mathcal{L}_{\Phi^2 G^\mu G^\nu} = +g_{Y,W}^2 g_{\mu\nu} \Phi^2 G_{Y,W}^\mu G_{Y,W}^\nu = g_{Y,W}^2 (-1) g_{\mu\nu} g_{t*}^2 I \Phi_*^2 G_{Y,W}^\mu G_{Y,W}^\nu}$$

$$\boldsymbol{\mathcal{L}_{\Phi^4} = -\lambda \Phi^4 = g_{t*}^4 I \Phi_*^4}$$

$$\to -\mathbf{g}_{t*}^2 \boldsymbol{I \Phi}_*^2 = \boldsymbol{\Phi}^2 \,\&\, \boldsymbol{I} < 0 \to \mathbf{g}_{t*}^2 = -\frac{\boldsymbol{Z}}{\boldsymbol{I}} \,\&\, \boldsymbol{\Phi} = \sqrt{\boldsymbol{Z}} \boldsymbol{\Phi}_* \to -\boldsymbol{\lambda} = \frac{\mathbf{1}}{\boldsymbol{I}} \to \boxed{\boldsymbol{\lambda} = \frac{\mathbf{g}_{t*}^2}{\boldsymbol{Z}} = \boldsymbol{g}_t^2 > 0}$$

Figure 9 Top loop induced generation of the Higgs scalar electroweak couplings and quartic coupling.

Similarly, the Higgs scalar acquires quartic coupling subject to

$$I(\Lambda) g_{t*}^4 Z_\Phi^{-2} = -\lambda(\Lambda) \tag{11}$$

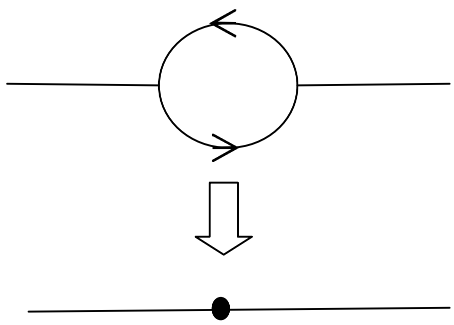
where top loop integral equals $I(\Lambda)$ proportional to $ln(\Lambda^2)$ in 4D.

Interested reader may prove that $I(\Lambda)$ is common for both loops.

By dividing Equ (11) with square of Equ (10) one obtains a *positive* effective quartic coupling $\lambda(\Lambda)$, i.e.

$$\lambda = -\frac{1}{I(\Lambda)} = g_{t*}^2 {Z_\Phi}^{-1} = g_t^2(\Lambda) > 0\ . \qquad (12)$$

The above solution should go to zero with $\Lambda \to \Lambda_{High}$ and should have a zero potential energy density at Λ_{high}. Below this scale coupling runs "logarithmically" and is renormalized according to the SM renormalization flow. However, being the characteristic of composite theory it nonetheless conspires to mainly reproduce the leading order term Equ (12). It would be interesting to investigate which solution within the $k = 2$ Higgs mass branch expresses this property the most; I will address that elsewhere.



$$\mathcal{L}_{\Phi^2} = m_H^2 \Phi^2 = -g_{t*}^2 I' \Phi_*^2$$

$$\to \ m_H^2 = \frac{g_{t*}^2 I'}{Z} < 0$$

$$m_H^4 = \lambda\,(I')^2$$

$$m_H^4 = \lambda\,(\text{"} \propto \Lambda^2 \text{"})^2 \ \to \ \boxed{\mu \propto -\sqrt{\lambda}}$$

Figure 10 Radiative generation of the Higgs scalar mass from top loop.

Consider next the composite dimensionless Higgs mass squared μ which is radiatively generated through top loop, see Fig 10. This solution should have a zero value at the high energy scale and subsequently, just beneath that scale, smoothly gains its *negative***,** therefore tachyonic, effective mass squared, as a consequence of the minus sign associated with the fermion loop.

According to Wigner [123], the space-like, negative mass squared, particles have non-compact little groups; their spin is not described by rotation group $SU(2)$ and in difference to massless particles their "spin" may be continuous parameter. This might be related to the extended Higgs sectors I introduce in Sections 6-7. But, the general concept of Higgs tachyon solution requires much better understanding. Hence, I hope that this paper will motivate more focused research effort in that direction.

The top loop $I'(\Lambda)$ in 4D is negative and proportional to quadratic term Λ^2

$$I'(\Lambda^2) \sim \Lambda^2 \ \to \ \mu(\Lambda^2) \propto -\sqrt{\lambda\,(\Lambda^2)}. \qquad (13)$$

Therefore, as anticipated in Equ (9), the dimensionless mass squared, μ, in the leading order is expected to be proportional to the square root of the scalar quartic coupling which runs only "logarithmically"; meaning no traditional hierarchy problem for the minimal SM with composite Higgs in 4D!

After many orders of magnitude the higher order corrections should overcome the composite Higgs mass leading order, and μ should finally reach the zero value corresponding to the low energy HMZC scale -- beginning of the electroweak phase transition. See Section 4 for interpretation of the interplay between g_t , g_{QCD} and their effect on μ. Finally, shortly beneath the HMZC scale, the renormalization flow drives μ to the intersection with λ at which point the smooth second order electroweak phase transition is completed with the correct value of the Higgs' VEV. Short running below the HMZC scale is a natural consequence of, now, positive Higgs mass squared and the renormalization flow at low energies.

Therefore, $m_H = 138.1 \pm 1.8\, GeV$ might be a good candidate for the composite Higgs mass, high energy fundamental scale placed in the vicinity of the Planck scale and the electroweak symmetry breaking scale $\cong 1050\, GeV$ within the second order phase transition.

It seems to me that effort to quantify deviation from Equ (13) across $142 \pm 5\, GeV$ mass range, and across energies smaller than the Planck mass, can be worthwhile; I will address that elsewhere.

As shown in Section 4, the SM scale at which tops condense may coincide with the low energy HMZC scale. This may be interpreted with a dual model in which, instead of an almost "fundamental" Higgs particle, one considers the original high energy model structure. The unknown elements of that model, e.g. four-fermion interactions, extra or less dimensions etc., that single out the dominant top quark, hence, conspires, in a natural fashion and in accord with the QCD, to create the electroweak symmetry breaking condensate at just the "right" low energy scale corresponding to the low energy HMZC scale.

In a summary, the $142 \pm 5\, GeV$ solutions, supporting the composite Higgs originating at very high energy scales, and mimicking the minimal SM Higgs at low energies, are favored for the following reasons:

1) These are the values in the center of the currently favored SM Higgs mass range as obtained by the combined electroweak precision data global fit and Higgs direct searches; see Section 2.5.

2) The leading divergences cancel out with the consistent value of radiatively generated Higgs mass; see k=2 branch solutions, presented in Section 3.

3) The Higgs and gluon mediated top interactions might satisfy condition for the $\bar{t}t$ loose bound state at low energies; see Section 4.

4) There are no vacuum stability and perturbativity constraints at energies smaller than the Planck mass and there is a single HMZC scale (i.e. no multiple HMZC scales) in the same energy range; see Section 2.

5) condition $\mu, \lambda = 0$ at Λ_{high} is *most closely* satisfied for $\Lambda_{high} \sim M_{Planck}$ and $m_H = 138.1\, GeV$

6) condition $\mu(\Lambda) < 0,\ \lambda(\Lambda) > 0$ is satisfied for $\Lambda < \Lambda_{high}$ in the vicinity of $\Lambda_{high} \sim M_{Planck}$.

7) hierarchy problem for $m_H^2(\Lambda^2) < 0,\ \lambda(\Lambda) > 0$ and composite $\mu(\Lambda^2) \propto -\sqrt{\lambda(\Lambda^2)}$ is rather benign.

8) Finally there is a way, that I show next, to directly match the heavy 4D Higgs with light k=1 2D branch solutions, presented in Section 3, implications of which are discussed in Section 7.

6.3. Model $3^{-1/2}$ within the second order phase transition

Consider 4D Lagrangian density

$$\mathcal{L} = \sum_k \left\{ \left[D_\mu \Phi_k \right]^+ \left[D^\mu \Phi_k \right] + \sum_i g_i \overline{\Psi}_i \Phi_k \Psi_i \right\} - \frac{m_H^2}{4} \sum_k \Phi_k^+ \Phi_k + \frac{\lambda}{8} \left(\sum_k \Phi_k^+ \Phi_k \right)^2 \qquad (14)$$

where $k = 1 \dots 3$ and i counts fermions. Assume that each scalar field Φ develops a non-zero vacuum expectation value equal to $\langle |\Phi| \rangle = \langle |\Phi_{SM}| \rangle / \sqrt{3}$. Hence, the fermion Yukawa coupling is $g_i = g_{SMi}/\sqrt{3}$ and the scalar - gauge bosons coupling is the same as in the SM.

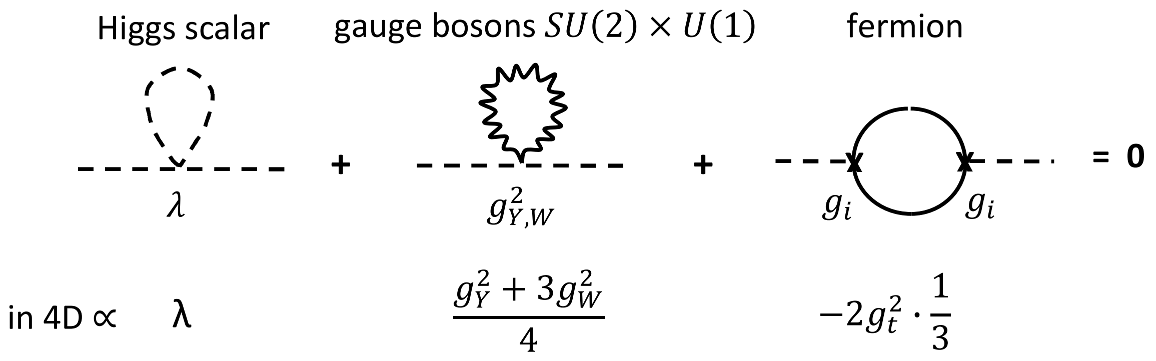


Figure 11 Cancelation of leading quantum corrections for Higgs scalar propagator in 4D Model $\mathbf{3^{-1/2}}$.

The conditions for cancellation of quadratic divergences, see Appendix, in the scalar propagators are

$$\lambda + \frac{g_Y^2 + 3g_W^2}{4} - \frac{2g_t^2}{3} = 0 \quad \rightarrow \quad \sqrt{\lambda} = \sqrt{\frac{2g_t^2}{3} - \frac{g_Y^2 + 3g_W^2}{4}}. \qquad (15)$$

Or in terms of observed masses, the Higgs mass is

$$m_H = \sqrt{\frac{4}{3} m_t^2 - M_Z^2 - 2M_W^2} = 138.3 \;\; GeV \, for \, m_t = 174 \, GeV. \qquad (16)$$

This result was obtained in [33] and it is in good agreement with the Higgs mass obtained in the previous subsection. In the range $m_t = 173.1 \pm 1.3 \, GeV$ one obtains a slightly smaller value $m_H = 136.8 \pm 2.2 \, GeV$.

If one now tries to reproduce the result on radiatively generated Higgs mass one obtains

$$\lambda v_{EW}^2 = \left(1 - \frac{2}{3}\right) \cdot top\ loop \;\; \text{or} \; top\ loop = 3\lambda v_{EW}^2 = 3 \frac{m_H^2}{v_{ew}^2} v_{EW}^2 = 0.93 v_{EW}^2 \, , \qquad (17)$$

or prediction for the "top loop" in 4D. It is interesting that for the central value within the k=2 branch

$$top\ loop = 1 \cdot v_{EW}^2 \;\; \text{for} \;\; m_H = 142.5 \, GeV. \qquad (18)$$

As discussed in Section 4, an interesting interplay between QCD and Higgs scalar exchange may be potential reason for just the "right" HMZC scale, i.e. the $\sim 1\ TeV$ electroweak symmetry breaking scale.

This model in notation of Section 3 has $x_1 = x(k=1) = 0.331$ and $y_1 = y(k=1) = 0.669$ that are characteristics of the k=1 branch; however, in 4D, this corresponds to the k=2 Higgs mass range.

For $m_t = 173.1\ GeV$ scaling between Equ (15) and (3) for k=1 corresponds to $\cong {}^{1.4676}/_3$ and $\cong {}^{2}/_{3}$ for scalar and top loop respectively.

The "curious" number connecting 2D (k=1) and 4D theories can be expressed as $\frac{\lambda_{4D}}{\lambda_{2D}} = 1.4676 = \left[\frac{2}{3} - 3\left(x_1 - \frac{y_1}{\pi}\right)\right] / \left(\frac{y_1}{\pi}\right)$. Interested reader should be able to easily derive this relationship.

I use this scaling in Section 7, to show how to conserve theory structure, in particular cancelation of leading divergences, across space-times with different dimensions, below and above the HMZC scale.

6.4. ***Model $\pi^{-1/2}$ in 2D***

As next excursion into "new physics" consider a model embedded in 2D with Lagrangian density

$$\mathcal{L}_0 = \int_0^\pi d\theta \left\{ \left[D_\mu \Phi(\theta) \right]^+ \left[D^\mu \Phi(\theta) \right] + \sum_i g_i \bar{\Psi}_i \Phi(\theta) \Psi_i \right\}, \qquad (19)$$

whereas the Lagrangian density describing fermions' interaction with gauge bosons is

$$\mathcal{L}_1 = \sum_i \bar{\Psi}_i D_\mu \gamma^\mu \Psi_i\ , \qquad (20)$$

where $D_\mu = \boldsymbol{\partial}_\mu + i g_Y Y B_\mu + i g_W \vec{T} \vec{W}_\mu$. The scalar self-interaction is of the form

$$\mathcal{L}_2 = \int_0^\pi d\theta \left\{ -\frac{m_H^2}{4} |\Phi(\theta)|^2 + \frac{\lambda}{8} |\Phi(\theta)|^2 \int_0^\pi d\theta' \, |\Phi(\theta')|^2 \right\} \qquad (21)$$

Each scalar field $\Phi(\theta)$ obtains a non-zero VEV equal to $\langle |\Phi(\theta)| \rangle = \langle |\Phi_{SM}| \rangle / \sqrt{\pi}$ and

$$\langle |\Phi|^2 \rangle = \int_0^\pi d\theta \, \langle \Phi(\theta) \rangle^2 = \int_0^\pi d\theta \, \left(\frac{\langle \Phi_{SM} \rangle}{\sqrt{\pi}} \right)^2 = \langle \Phi_{SM} \rangle^2 \qquad (22)$$

provides the appropriate masses for the Z and W gauge bosons. Similarly, the fermion Yukawa coupling must be $g_i = g_{SM}/\sqrt{\pi}$ and the scalar - gauge bosons coupling is the same as in the SM.

Clearly, the idea is that maybe

$$\langle Re(\Phi) \rangle = \int_0^\pi d\theta \, \langle \Phi(\theta) \rangle cos\theta = \int_0^\pi d\theta \, \frac{\langle \Phi_{SM} \rangle}{\sqrt{\pi}} cos\theta = 0 \qquad (23)$$

may provide alternative explanation for the vacuum energy problem.

The condition for cancellation of quadratic divergences for each scalar field $\Phi(\theta)$ is shown in Fig 12.

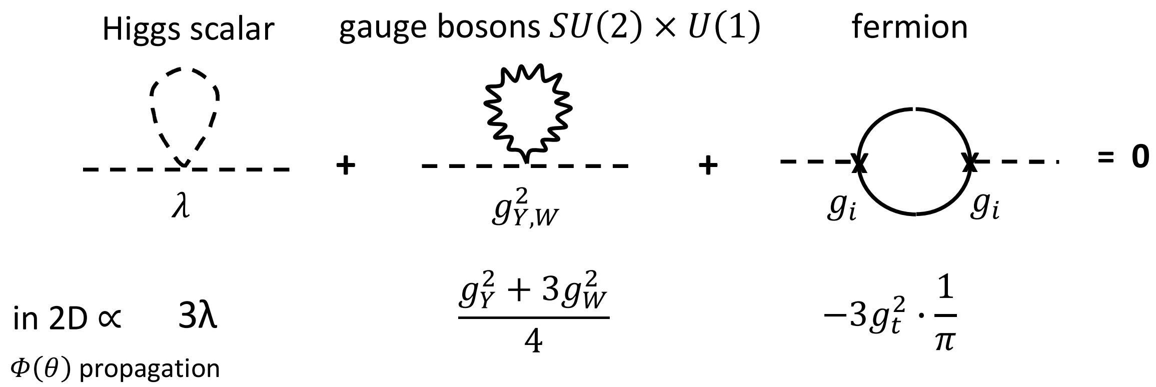


Figure 12 Cancelation of leading quantum corrections for Higgs scalar propagator in 2D Model $\boldsymbol{\pi^{-1/2}}$.

The above result suggests that

$$\sqrt{\lambda} = \sqrt{\frac{1}{3}\left(\frac{3}{\pi}g_t^2 - \frac{g_Y^2+3g_W^2}{4}\right)} \rightarrow m_H = \sqrt{\frac{1}{3}\left(\frac{6}{\pi}m_t^2 - M_Z^2 - 2M_W^2\right)}\text{, i.e.} \qquad (24)$$

$$m_H = 109.5 \pm 2.9\ (5.2)\ GeV \text{ for } m_t = 171.4 \pm 3\ (4)\sigma\ \ GeV, \sigma = 1.3\ GeV. \qquad (25)$$

The Higgs mass may be obtained from the radiatively generated mass, as in Equ (3) with $z, k = 1$ and $y = 1 - \frac{1}{\pi}$,

$$m_H = m_t\sqrt{\frac{2}{\pi}\left(1 - \frac{1}{\pi}\right)} = m_t 0.6588\ \cong m_t\frac{2}{3} = 115.4 \pm 0.9\ GeV. \qquad (26)$$

7. Higgs within the first order phase transition

In this section, I hypothesize on physics that may be responsible for the k=1 branch Higgs masses, introduced in Section 3. I show that lighter Higgs mass solution is likely associated with the first order (discontinuous) electroweak phase transition. In my terminology, renormalized theory does abruptly change the parameters and degrees of freedom across the low energy HMZC scale affiliated with the first order electroweak symmetry breaking. I sketch class of models that may *exactly* remove tachyon solution at high energies. Hence, I introduce more appropriate name for that transitional scale.

I analyze model, proposed by Popovic 2002 [33], where top quark is composite, composed of 3 fundamental fermions, and Higgs scalar is composite, composed of 2 fundamental fermions. I dub class of models, with $m_H \cong \frac{2}{3}m_t = 115.4 \pm 0.9\ GeV$as generic feature, the Composite Particles Model (CPM).

7.1. Composite Particles Model (CPM) and HMNZ^2 scale

The low energy Higgs field may be created as a consequence of the electroweak symmetry breaking within the first order phase transition. The "new physics", = a more fundamental physics, may conspire to promote effective scalar field with propagator free of leading divergences and with $\mu, \lambda = 0$ *at, just above* or maybe even *at all energies* above the low energy HMZC scale.

Calculation in 2D, Equ (3a) with $x = 1$, then leads to

$$g_{t*} = \sqrt{\frac{w}{12}(g_Y^2 + 3g_W^2)} \ , \tag{27}$$

i.e. the top Yukawa coupling at the HMZC scale and where $w = 1 \ or \ \frac{2}{3}$ is introduced to account for the correct counting of the gauge boson polarization degrees of freedom. In 2D one expect $w = 1$ for the longitudinal degree of freedom and in 4D one expect $w = \frac{2}{3}\ (1)$ for two transversal polarizations (all three polarizations). Traditionally, the gauge boson in 4D acquires the longitudinal polarization via the electroweak symmetry breaking; see for example [120].

Possibility that there are no tachyons *at all energies* above the HMZC scale $\sim \Lambda_{EWSB}$ but rather just massless particles whose exact zero mass is established by defining principle expressed in Equ (27) or Equ (3a) with $x = 1$ ($y = 0$ or $\sqrt{\lambda} = 0$ or $m_H = 0$) seems worthwhile investigating. This would require a more appropriate name, the *Higgs Mass Non-Zero to Zero* (HMNZ^2) transitional scale instead of the HMZC scale. But, Higgs mass, then, needs to be stabilized *at all energies* larger than the HMNZ^2 scale!

Massless limit, within a non-zero VEV theory, I believe, should correspond to an ordinary, massless scalar during early stages of our Universe, within a zero VEV theory. Therefore, if there are no other fundamental mass terms except Planck mass, M_{Pl}, the Universe dynamics should be parameterized with single mass parameter and set of couplings running naturally slow.

Interested reader may show that it is not possible to exactly remove tachyon solution with unassisted SM at high energies. While quadratic and *all* logarithmic divergences in the scalar propagator cancel out at the HMZC scale thanks to Equ (27) and $\mu, \lambda = 0$, the gauge couplings and parameters defining Higgs potential are still running. Hence, nothing prohibits divergences to reappear at slightly larger energies.

However, let us still assume that Equ (27) may indeed describe the *decoupling limit* where scalar sector effectively decouples and tachyon solution is exactly removed.

Let us imagine that Equ (27) is satisfied at very low energies with SM values of g_Y, g_W. One obtains

$$g_{t*} \cong 0.34\sqrt{w} \sim \frac{g_t}{3}\sqrt{w} \ . \tag{28}$$

In 4D, see Equ A5, the above results, i.e. Equ (27-28), translate to

$$w\frac{g_Y^2+3g_W^2}{4} - 2g_{t*}^2 = 0 \ \rightarrow \ g_{t*} = \sqrt{\frac{w}{8}(g_Y^2 + 3g_W^2)} \ , \tag{29a}$$

$$g_{t*} \cong 0.42\sqrt{w} \sim g_t\sqrt{\frac{w}{6}} = \frac{g_t}{3}\sqrt{\frac{3w}{2}} \ . \tag{29b}$$

As anticipated, Equ (27) and Equ (29a) with $g_{t*2D} = g_{t*4D}$ agree for $w = 1$ in 2D and $w = \frac{2}{3}$ in 4D, independent of the actual values of g_Y , g_W.

That is promising but there are at least three problems: (1) SM λ is not equal zero at low energies, (2) no SM fermion has that Yukawa coupling and (3) that Yukawa coupling is not strong enough to balance the QCD forces, i.e. no suitable condensate that would break the electroweak symmetry, see Section 4.

One may compare the above situation with discussion in Section 4.1; dynamical generation of the appropriate masses for Z and W, e.g. based on the Pagels-Stokar relationship [98] or gap equations in the gauged NJL model [99-100], requires a large dynamically generated fermion mass; i.e. problem that motivated TC^2 [102-110] and Top Seesaw [111-113]. Unfortunately, the above situation suggests an even worse mismatch between the dynamically generated boson and fermion degrees of freedom!

But there is a cure.

First and second problems are indicative that this needs to be the first order phase transition where parameters and degrees of freedom of the theory abruptly change across the HMZC scale. Second and third problems are suggestive that this might be interplay with at least three "fundamental" fermions forming a composite fermion, identified as top quark, with characteristics that provides condition for formation of top condensate responsible for the electroweak symmetry breaking, see section 4. This conclusion is anticipated by both 2D and 4D theories as result of Equ (28) and (29b).

If creation of composite or bound state is thought of to be a non-local phenomenon, as it should be, the Yukawa coupling should be thought to be an additive quantity that adds to ~one as in SM.

If there are 3 scalar fields within condensate of mass $2m_t$, the scalar mass, in the non-relativistic limit, is

$$m_H = \frac{2}{3} m_t = 115.4 \pm 0.9\ GeV \qquad (30)$$

which then appears, according to 2D considerations in Section 3, as the k=1 Higgs mass branch physics within the first order phase transition at $\sim 10^{2.9} \cong 800\ GeV$ [26].

In this picture, one expects the QCD assisted with "new" physics to create (1) Higgs, "meson"-like particle consisting of two "fundamental" fermions, (2) top quark, "baryon"-like particle consisting of three "fundamental" fermions and (3) top condensates breaking the electroweak symmetry at the HMZC or HMNZ^2 scale. These lines of thoughts have been proposed in [33].

And Higgs mass in Equ (30) in the non-relativistic limit probably need to be corrected by only a relatively small amount; after all, the strong QCD interactions are not that strong above the M_Z!

I call this class of models, with Equ (30) as common feature, the Composite Particles Model (CPM).

Therefore, in this picture, Equ (27) may describe the *decoupling limit* in a sense that there is no scalar field above the HMNZ^2 scale and therefore no tachyons, *exactly*, either; hence, similar to models with the strong interaction mediated dynamical electroweak symmetry breaking [83-84, 37]. Higgs mediated top-anti top interaction in Section 4 is identified here as more complex dynamics associated with "new" physics responsible for CPM. Hence, top condensation may be natural consequence of QCD and "new" physics logarithmic running, i.e. resolving hierarchy problem and removing tachyons at once.

Alternative interpretation could be that there is still a composite scalar field but it is effectively decoupled from the electroweak sector of the theory.

Consider $\Phi\Phi \to GG$ scattering at high energies above the HMNZ^2 scale where G symbolizes electroweak gauge boson. By rewriting longitudinal gauge bosons in terms of Goldstone bosons and interpreting them as the fermion composites one may show that Lagrangian term for $\Phi\Phi \to GG$ scattering vanish if Equ (27) is satisfied. The right-handed fermion within one composite scalar field may couple with 3 left-handed fermions, either from another scalar field or from the two longitudinal gauge bosons, and the other right-handed fermion, belonging to other composite scalar field, may couple with other two left-handed fermions. That is factor of 6. The extra factor of 2 is due to chiral symmetry which is unbroken at all energies above the HMNZ^2 scale. This will be addressed in more detail elsewhere.

The composite fermions have been addressed in the past [124], though in different context than here or [33]. More recently, the possibility that fermions may be composite is discussed as one way of preventing the excessive FCNC in the "Littlest Higgs" model [125]. The "Little Higgs" [125-128] may be composite if the sigma field is a condensate of strongly interacting fermions; hence the fermions may be composite with masses protected by approximate global symmetries [124-125].

The models with composite scalars, with masses sensitive to the quadratic quantum corrections, have been proposed in the past [129-132]. Recent model building effort [125, 133] in the context of $SU(5)/SO(5)$ breaking pattern [131-132] and Higgs thought of as a light pseudo-Goldstone boson claimed that Higgs mass may be stabilized against radiative corrections. That is accomplished with approximate global symmetries, involving new heavy particles, that are imposed to soften the cutoff-dependence [125] and with TC^2 top-color interactions and conjecture that "we live in a region of the explicit chiral symmetry breaking interaction parameter space that lies between successive electroweak symmetry breaking phase transitions — at which m_H and VEV must vanish" [133].

It would be interesting to investigate if these mechanisms may be compatible with CPM.

Should there be cancelation of quadratic divergences at energies smaller than the HMZC scale? Or alternatively, what are degrees of freedom that might entangle that behavior at low energies?

Consider again the 2D theory with the phase transition corresponding to $g_{t*} \sim \frac{g_t}{3} \to Ag_{t*} \sim \frac{g_t}{3}A$, $m_H = 0 \to m_H = \frac{2}{3}m_t$ and $\lambda = 0 \to \lambda = \frac{m_H^2}{v_{EW}^2} = \frac{4m_t^2}{9v_{EW}^2} = \frac{2g_t^2}{9}$ and require cancelation above and below the HMZC, or HMNZ^2, scale. in 2D, The leading divergence cancelation in 2D, Equ (A6), then leads to

$$\frac{g_Y^2+3g_W^2}{4} = 3g_t^2\left(\frac{1}{3}\right)^2 \quad \text{and} \quad 3\lambda + \frac{g_Y^2+3g_W^2}{4} = \frac{6g_t^2}{9} + 3g_t^2\left(\frac{1}{3}\right)^2 = 3g_t^2\left(\frac{1}{3}\right)^2 A^2 \qquad (31)$$

where Equ (30) was assumed and factor A was introduced as a free parameter to be determined. I find

$$A^2 = 3 \quad (32)$$

which is the k=1 solution with $y = 0.66$ introduced in Section 3. Therefore, the above electroweak phase transition is simply $x = 1 \rightarrow x = 0.33$ transition, in the notation of Section 3, and taking place at HMZC, or HMNZ^2, scale $\sim\Lambda_{EW} \sim 10^{2.9} \cong 800\ GeV$ [26].

After inspection, one finds the top quark mass as in [33],

$$g_t = \sqrt{3\ \frac{g_Y^2+3g_W^2}{4}} \cong 1.025 \ \rightarrow\ m_t = \sqrt{\frac{3}{2}(M_Z^2 + 2M_W^2)} = 178.51\ GeV. \quad (33)$$

What about the 4D cancelations?

Well, left hand expression in (31) is identical in 2D and 4D for $w = \frac{2}{3}$. And 4D analogy of the right hand expression in Equ (31) maybe *do not need* to be satisfied!

If Higgs mechanism takes place only in 2D then VEV do not need to be non-zero everywhere in space-time which clearly removes both the hierarchy and vacuum energy problems. Therefore electroweak symmetry breaking may be determined by 2D propagator physics and 4D couplings. For low energy HMZC or HMNZ^2 scale there should be no substantial fine tuning. I revisit that in Section 7.4, where I show that it is possible, after all, to cancel the leading divergences and conserve theory structure across the space-times with different dimensions, both below and just above the HMZC, or HMNZ^2, scale.

Next, I investigate the "new" physics imprints that could reproduce the above CPM structure by sketching several models that deals with external particle degrees of freedom within the fundamental 2D and 4D space-times as well as with degrees of freedom (e.g. color, flavor) in the internal space.

7.2. $3^{-1/2}$ model with flavor

Consider that right handed "up" quarks create condensates in a generation universal manner, i.e. each of the right handed "up" quarks couples with each of the left handed "up" quarks. Each of the nine condensates is assigned scalar field Φ_{ij} where $i,j = 1,2,3$ and, finally, each scalar acquires identical vacuum expectation value $v_{EW}/3$. However, there exists only one non-zero fermion mass eigenstate which correspond to identically populated superposition of the left handed and the right handed "original" states. That massive fermion is identified as the top quark. The top condensate will have non-zero VEV that is three times larger than the "original" condensate value,

$$\langle\bar{\Psi}_t\Psi_t\rangle = \langle\frac{1}{\sqrt{3}}(\bar{\Psi}_1 + \bar{\Psi}_2 + \bar{\Psi}_3)\frac{1}{\sqrt{3}}(\Psi_1 + \Psi_2 + \Psi_3)\rangle = 3\langle\bar{\Psi}_i\Psi_j\rangle \quad (34)$$

where $i,j = 1,2,3$. And the other two fermion mass eigenvalues are zero.

The top quark however couples with *three* physical scalar fields below the HMZC scale, each of which is superposition of the scalar fields corresponding to the condensates with common left-handed partner

$$\Phi_i = \sqrt{\Phi_{i1}^2 + \Phi_{i2}^2 + \Phi_{i3}^2} \leftrightarrow \bar{\Psi}_{iL}\Psi_{jR} + \bar{\Psi}_{jR}\Psi_{iL} \text{ where } j = 1,2,3. \quad (35)$$

Superposition which mixes the left-handed partners is meaningless due to $SU(2)$ rotations that mix the final fermion mass eigenstates. Each of the three fields Φ_i where $= 1,2,3$, acquires $^{v_{EW}}/_{\sqrt{3}}$, whereas the physical top quark couples to each of them with coupling $3^{-1/2}$. Interested reader may note that this physics may be described with the effective Lagrangian density in Equ (14) addressed in Section 6.3.

If one assumes CPM non-relativistic limit $m_H = \frac{2}{3}m_t$, which I base on discussion in Section 7.1, and consider leading divergences in 2D with $w = 1$ and $\frac{g_t}{3} \rightarrow \frac{g_t}{\sqrt{3}}$, $\lambda = 0 \rightarrow \frac{2}{9}$ i.e. $m_H = 0 \rightarrow \frac{2}{3}m_t \in (115.4, 119.0)GeV$ while crossing the HMZC, or HMNZ^2, scale from higher to lower energies, one finds the conditions for cancelation for scalar fields Φ_{ij} and Φ_i respectively to be

$$\frac{g_Y^2+3g_W^2}{4} = 3g_t^2\left(\frac{1}{3}\right)^2 \text{ and } 3\lambda + \frac{g_Y^2+3g_W^2}{4} = \frac{6g_t^2}{9} + \frac{g_Y^2+3g_W^2}{4} = 3g_t^2\left(\frac{1}{\sqrt{3}}\right)^2 \quad (36)$$

i.e. two numerically identical expressions corresponding to two completely different physical interpretations. Hence, one obtains the prediction for the top quark Yukawa coupling and mass as [33]

$$g_t = \sqrt{3\,\frac{g_Y^2+3g_W^2}{4}} \cong 1.025 \rightarrow m_t = \sqrt{\frac{3}{2}(M_Z^2 + 2M_W^2)} = 178.51\ GeV. \quad (37)$$

In 4D with $= \frac{2}{3}$, left hand side from Equ (36) takes the form

$$w(M_Z^2 + 2M_W^2) = 4m_t^2 \cdot \left(\frac{1}{3}\right)^2 \rightarrow m_t = 178.51\ GeV. \quad (38)$$

Hence, one again obtains Equ (37). The right hand side in Equ (36) in 4D with $w = \frac{2}{3}$ (1) takes the form

$$m_H^2 + w(M_Z^2 + 2M_W^2) = 4m_t^2\left(\frac{1}{\sqrt{3}}\right)^2 \rightarrow m_H = 160.6\ GeV\ (136.8\ GeV) \quad (39)$$

for $m_t = 173.1\ GeV$. This translates to $m_H = 168.3\ GeV\ (145.8\ GeV)$ for $m_t = 178.1\ GeV$; result that is clearly inconsistent with the $m_H = \frac{2}{3}m_t$ premise. But as discussed above, Equ (39) maybe *do not need* to be satisfied at energies smaller than the low energy HMZC, or HMNZ^2, scale.

The $w = 1$ solution introduced above appears as the k=2 branch Higgs mass in the close vicinity to two special Planck mass affiliated solutions, discussed in Section 6.1. The *approximate* $\mu, \lambda = 0$ solution [33] is obtained for the top quark mass world average, $m_t = 173.1\ GeV$, while Higgs within the Planck mass version of the Coleman-Weinberg conjecture [52] is obtained for the non-relativistic limit model prediction $m_t = 178.1\ GeV$ where m_t is obtained from the Z and W gauge boson masses alone.

The $w = \frac{2}{3}$ mass, Equ (39), extends beyond the $147\ GeV$ limit, requiring a single HMZC scale at energies smaller than the Planck mass, addressed in Section 2, but it is below the $171\ GeV$ perturbativity limit.

It would be rather out of the ordinary situation if the late LEP suspicious signals [81] corresponds to important 2D Higgs dynamics centered at $m_{H2D} \cong 117\, GeV$, whereas the LHC discovers Higgs 4D dynamics centered at $m_{H4D} \cong 164\ (141)\ GeV$ for $w = \frac{2}{3}\ (1)$ polarizations, as described above.

I revisit the potential interplay between the 2D and 4D Higgs dynamics in Section 7.4.

7.3. $3^{-1/2}$ model with color

Consider model with 9 scalars that mix fermions with different colors, but not flavor, at energies slightly larger than the HMZC, or HMNZ^2, scale. The fermion top Yukawa coupling equals $\frac{g_t}{3}$. However, at energies slightly smaller than the HMZC, or HMNZ^2, scale each top color couples with color specific Higgs scalar particle and there are three of those. Hence, each scalar field acquires $v_{EW}/\sqrt{3}$ and again the physical top quark couples to each of them with coupling equal to $g_t 3^{-1/2}$. This model structure, which is clearly numerically identical to those presented in Sections 6.3 and 7.2, was proposed in [33].

If one again assumes non-relativistic CPM limit $m_H = \frac{2}{3} m_t$, which I base on discussion in Section 7.1, and consider leading divergences, one should discover, as in Section 7.2, that this is the first order phase transition where $\frac{g_t}{3} \to \frac{g_t}{\sqrt{3}}$, $\lambda = 0 \to \frac{2}{9}$ i.e. $m_H = 0 \to \frac{2}{3} m_t \in (115.4, 119.0) GeV$. This transition is consistent with result of analysis that led to Equ (31-32).

I now consider quadratic divergences in 2D with $w = 1$ above and below the HMZC or HMNZ^2 scale. Condition for cancelation above the HMZC or HMNZ^2 scale for 9 scalar fields

$$\frac{g_Y^2 + 3g_W^2}{4} = 3g_t^2 \left(\frac{1}{3}\right)^2 \tag{40}$$

i.e. there is single color combination in the fermion loop and fermion coupling is $g_t/3$.

Condition for cancellation beneath the HMZC or HMNZ^2 scale for 3 color specific scalar fields is

$$3\lambda + \frac{g_Y^2 + 3g_W^2}{4} = \frac{6g_t^2}{9} + \frac{g_Y^2 + 3g_W^2}{4} = 3g_t^2 \left(\frac{1}{\sqrt{3}}\right)^2 \tag{41}$$

where there is again single color specific combination in the fermion loop but coupling is now $g_t 3^{-1/2}$.

After inspection one should discover that Equ (40) and Equ (41) are numerically identical.

Again, one obtains prediction on the top quark Yukawa coupling [33] that is identical to Equ (37) in Section 7.2., result that is 4 "$world$" σ (less than 3%) away from the world average top quark mass [60].

Reproducing the above analysis in 4D would suggest the same result as in Equ (38-39), see Section 7.2.

7.4. *2D and 4D models with movers*

Imagine for moment that 2D "fundamental" fermions are simply x, y, and z movers. The 4D top quark appears as composite of all three species. Consider next the 4D collision of top anti-top at small CM momentum. The collision requires pairing of orthogonal directions where $2 \cdot \frac{2}{3} \cdot m_t$ evaporates into the vacuum and propagator mass appears only as $\frac{2}{3} \cdot m_t$. Therefore, the 2D “fundamental” fermion couples with strength as in Equ (28). I also assume that phase space for “top loop” at energies smaller than the HMZC has an additional factor $b = 3 \text{ or } \pi$ to account for the orthogonal 2D space.

Therefore “new” physics, defining the “original” coupling g_{t*}, transforms, at energies slightly smaller than the HMZC, or HMNZ^2, scale to the minimal SM with the broken electroweak symmetry, as parameterized with a set of the effective parameters (and reproducing the CPM structure)

$$g_{t*} \to g_{t*}\sqrt{b}\,,\ m_H = 0 \to m_H = \frac{2}{3}m_t \text{ and } \lambda = 0 \to \lambda = \frac{m_H^2}{v_{EW}^2} = \frac{2g_t^2}{9}\,. \tag{42}$$

Again, consider the leading divergences in 2D and require that gauge boson loops are canceled at energies slightly larger than the HMZC, or HMNZ^2, scale, whereas both the gauge and Higgs boson loops are cancelled at the energies smaller than the HMZC, or HMNZ^2, scale, i.e. according to Equ (A6),

$$3\lambda + \frac{g_Y^2+3g_W^2}{4} = \frac{6g_t^2}{9} + 3g_{t*}^2 = 3g_{t*}^2 b \tag{43}$$

$$\to \quad \frac{g_t}{g_{t*}} = 3\sqrt{\frac{1}{2}(b-1)} \;=\; 3.00\ (3.10) \text{ for } b = 3\ \ (\pi). \tag{44}$$

Now I will match the 4D model with the above 2D model structure. I use scaling obtained in Section 6.3., i.e. $\frac{\lambda_{4D}}{\lambda_{2D}} = \left[\frac{2}{3} - 3\left(x_1 - \frac{y_1}{\pi}\right)\right] / \left(\frac{y_1}{\pi}\right) \cong 1.47$. Therefore, Equ (43-44) translate in 4D to

$$1.47\lambda + \frac{g_Y^2+3g_W^2}{4} = 1.47\frac{2g_t^2}{9} + \frac{2g_{t*}^2}{w} = 2g_{t*}^2 b \tag{45}$$

$$\to \quad \frac{g_t}{g_{t*}} = 3\sqrt{\frac{b-w^{-1}}{1.47}} \;=\; 3.03\ (3.17) \text{ for } w = \frac{2}{3} \text{ and } b = 3\ \ (\pi). \tag{46}$$

And only the $w = \frac{2}{3}$ solution is consistent with the premise with 3 “fundamental” fermions whose Yukawa couplings add linearly. This is as expected, see Equ (29). As it is “traditionally” anticipated, the gauge bosons in 4D have two transversal polarizations in the unbroken electroweak phase and subsequently gain one additional, longitudinal, polarization in the broken electroweak phase.

If there are π “fundamental” fermions then Equ (44) and (46) suggest consistent structure for $b = \pi$.

Therefore, discovered theory has electroweak symmetry breaking with abrupt change of parameters defined by Equ (42) with cancelation of quadratic divergences in both 2D and 4D, with correct counting of the gauge bosons polarizations and almost consistent ratio equal 3 between the physical and “bare” top Yukawa couplings: 3.03 and 3.00 in 4D and 2D respectively.

Hence, this model structure may exactly remove tachyons both in the fundamental 2D and 4D theories!

I was able to retain the 4D cancellation thanks to the appropriately understood scaling between the 2D and 4D theories. Clearly, this may be applied to models with entirely different physical interpretation.

The Higgs mass and top quark mass in 2D model equal

$$m_H = \frac{2}{3} m_t\,(115.4, 119.0) GeV, m_t \in \left(173.1, \sqrt{\frac{3}{2}(M_Z^2 + 2M_W^2)} = 178.5\ GeV\right) \quad (47)$$

corresponding to the phase transition at $\sim 10^{2.9} \cong 800\ GeV$ [26,33]. If predicted top quark mass is scaled back to the world average mass and predicted Higgs mass is scaled by the same factor, one obtains

$$m_H \rightarrow \frac{173.1}{178.5} 115.4\ GeV = 111.9\ GeV. \quad (48)$$

Similarly, the Higgs mass and top quark mass in 4D model equal

$$m_H = \sqrt{1.47 \frac{2g_t^2}{9}} v_{EW} = \sqrt{1.47 \frac{2}{9} 3.03^2 g_{t*}^2} v_{EW} = \sqrt{1.47 \frac{2}{9} 3.03^2 \frac{g_Y^2 + 3g_W^2}{12}} v_{EW},$$

$$m_H = \sqrt{1.47 \frac{2}{9} 3.03^2 \frac{M_Z^2 + 2M_W^2}{3}} = 145.8\ GeV,$$

$$m_t = \sqrt{3.03^2 \frac{M_Z^2 + 2M_W^2}{6}} = 180.3\ GeV. \quad (49)$$

with the electroweak phase transition scale roughly in the range $1 - 1.15\ TeV$.

If predicted top quark mass is scaled back to the world average top quark mass and predicted Higgs mass is scaled by the same amount, then one obtains

$$m_H \rightarrow \frac{173.1}{180.3} 145.8\ GeV = 140.0\ GeV \quad (50)$$

in the close vicinity of the $m_H = 138.1\ GeV$ solution, see [33], discussed in Section 6.1.

It would be rather out of the ordinary situation if the late LEP suspicious signals [81] corresponds to important 2D Higgs dynamics centered at $m_{H2D} \cong 117\ GeV$, whereas the LHC discovers Higgs 4D dynamics centered at $m_{H4D} \cong 140\ GeV$, for $w = \frac{2}{3}$, as described above.

8. Conclusion

With LHC collecting the high energy data it is expected that our understanding of Nature will dramatically advance in the near future. Paradigm shifts in physics always generated many important new technologies with a vast range of practical applications. Hence, there is a good chance that society will again benefit greatly from this largest physics research endeavor ever undertaken by mankind.

Why do particle physicists expect a huge paradigm shift in the next year or two? Well, there is a single particle anticipated by the current dogma [1-21] awaiting to be discovered. And it seems certain that SM Higgs should be within the reach of LHC. But it is also generally accepted that current dogma is incomplete and incorrect; the reason being hierarchy [37] and vacuum energy problems [38-40].

Will LHC address physics outside of the current dogma? As I show in this paper, it must. Even in the next to the "worst case" scenario, where Higgs is discovered and there is nothing else to surprise us at small energies, there is still an energy scale, HMZC $\sim \Lambda_{EWSB}$ [26, 33], within the LHC reach where current dogma itself suggests that effective Higgs particle should transition from standard particle, positive mass squared, to tachyon, negative mass squared, degree of freedom. Never in history have particle physicists dealt with anything similar. Last month, LHC reached the CM energies ($\sim 3.5\,TeV$) that are within the HMZC range with a goal to embrace the entire HMZC range soon. Therefore, yes, the LHC is just beginning to make its mark in history by stepping distinctively outside of the known physics territories.

By preparing classical system (raising temperature, density etc), the average CM collision energies can be brought to the HMZC energy $\sim\Lambda_{\mathrm{EWSB}}$; corresponding to condition for the classical phase transition which probably happened in the very early Universe [59], though in opposite direction and without the actual change of the vacuum structure / state. As discussed, by going back in physical time the vacuum structure/state of today's Universe with non-zero VEV transitions to one with zero VEV. And in the zero VEV Universe the tachyon Higgs is just an ordinary non-tachyon scalar particle. Therefore, we might learn a lot about the actual electroweak phase transition by studying physics at the HMZC scale.

While tachyon theories are often addressed in the context of string theory and cosmology, I find that there is an *alarming lack* of literature and ongoing research effort among the rest of particle physics community. I hope that this paper will motivate more focused research effort in that direction.

In this paper I carefully map the physical Higgs mass with the low energy HMZC scale $\sim\Lambda_{EWSB}$ based on my earlier work back in 2001 [26] and 2002 [33]. As shown here, the HMZC scale exist for the Higgs mass lighter than approximately 200 GeV; range that is also strongly favored by the electroweak precision data and direct Higgs searches [83]; see Section 2.5. The general HMZC scale range is $10^{2.9} - 10^{3.7}\,GeV$ ($800\,GeV - 5\,TeV$). For the Higgs masses within the $114.4 < m_H < 182\,GeV$ range [83], preferred by global fit and direct searches, the HMZC scale range is $800\,GeV - 1.8\,TeV$, see Fig 2.

If SM is expected to be valid at all energies smaller than the Planck mass, i.e. the SM "desert" or "long lived" scenario, then renormalization group flow implies that Higgs must be heavier than $137.0 \pm 1.8\,GeV$, based on the vacuum stability limit [26, 41-48, 83]; otherwise, unacceptable deeper minimum of the effective potential occurs. Similarly, the SM renormalization group flow implies that Higgs mass must be lighter than $171 \pm 2\,GeV$, based on the perturbativity limit [26, 49-50, 83]; otherwise, scalar self interactions diverge, i.e. strongly coupled Higgs sector cannot be described with the perturbation theory. Finally, because there are generally two HMZC scales per physical Higgs mass, the condition that there is a single HMZC scale [26, 33] at energies smaller than the Planck mass puts an upper limit on the Higgs mass equal to $146.5 \pm 2\,GeV$. Therefore, if SM is a valid description of Nature at all energies below the Planck scale, where it has the effective structure of an unbroken electroweak symmetry, then the stability curve and condition that there is a single HMZC scale at all energies smaller than the Planck scale limits the SM Higgs mass to a very tight range of roughly $142 \pm 6\,GeV$ with a corresponding electroweak phase transition scale roughly in the range $1 - 1.15\,TeV$.

However, traditionally, there is no strong enough reason to expect SM to be a valid theory at all energies smaller than the Planck scale. Just to the contrary, hierarchy problem [37] affiliated with the presence of leading quantum corrections associated with quadratic divergences in the scalar propagator suggests that theory cannot be valid across vast energy scales unless there is a defining principle that provides explanation which goes beyond the current dogma. Otherwise, the theory quickly, i.e. already after couple of magnitudes in energy, becomes unnaturally finely tuned.

In Section 2.4, I present an analysis which compare the HZMC scale with stop mass to show that MSSM [27-31] is less unnatural than SM at low energies for $m_H \leq 120.9 \pm 0.9\ GeV$.

In Section 3, I investigate the leading SM quantum corrections to the scalar propagator in the 2D theory. I show that one could simultaneously satisfy (1) complete radiative generation of the Higgs mass through top loop and (2) complete cancelation of the remaining leading quantum corrections to scalar propagator. There is unique solution for the Higgs mass. This solutions is parameterized with k=1 or 2 and corresponding SM solutions in the zeroth order are $113.0 \pm 1.0\ GeV$ and $143.4 \pm 1.3\ GeV$. I would not be very surprised if the actual Higgs mass happens to be within the $\pm 5\ GeV$ range.

It is worth noting that the k=1 branch almost embraces the late LEP suspicious signals [23-24, 81-83] in the vicinity of $115\ GeV/c^2$ whereas the k=2 branch almost embraces the “long lived” SM solutions, I discuss in Section 6.1, in the context of the SM renormalization flow all the way up to the Planck mass.

The Technicolor theories [95-96, 37] were introduced to address the hierarchy problem [37]; main idea was that there may be additional forces beyond the SM dogma that could provide glue for fermions to bond and create effective scalar field or appropriate set of Goldstone bosons [13-14] that would break the electroweak symmetry at low energies. The entire point was that gauge couplings associated with these forces would run logarithmically and therefore provide natural explanation for large hierarchy.

In Section 4, I show that top condensate formation may be consistent with interplay between the QCD gluon and Higgs scalar mediated top anti-top interactions. Vice versa, starting with this as a premise I predict the fine structure constant of the strong QCD interactions up to precision better than 2% in the leading order calculation and explain how to reach even better agreement. Interestingly, the predicted mean value is only 0.25% away from the world average value. Therefore it seems that dynamical top condensation may indeed be a viable option with either fundamental or composite Higgs.

In Section 5, I hypothesize that 4D theory may be only an effective theory which corresponds to more fundamental 2D theory as this can solve hierarchy and vacuum energy problems that particle physics faces today. This could be to the extent that (1) 4D electroweak symmetry breaking is governed by 2D electroweak symmetry breaking and 4D couplings or (2) that 4D theory is effective theory completely described by 2D theory, where dimensionality of space-time enters less as a premise and more as a consequence of the fundamental 2D theory. In 2D leading quantum corrections are only logarithmically divergent and scalar VEV may be confined only to propagator 2D space associated with propagating particles in 4D; i.e. equivalent to compactifying the Higgs *ether* from entirety of 4D to just a small subset of that space. Complete removal of the Higgs *ether* is likely dynamical symmetry breaking. In this paper I discuss 2D fundamental theory with effective scalar field as well as dynamical symmetry breaking.

As previously noted [33] there is one solution of the SM renormalization flow in the vicinity of $138.1\ GeV$ Higgs that is a very distinctive one. Both effective Higgs mass and quartic scalar coupling conspires to be almost zero at the same high energy scale. This high energy scale happens to be the Planck mass scale. Furthermore, for both μ and λ the SM renormalization flow is rather logarithmic and already based on visual inspection, see Fig 8, it appears that these quantities are directly related. Another distinctive solution, centered at $146.5\ GeV$, corresponds to the Planck scale version of the CW conjecture [52].

In Section 6, I address the composite Higgs built out of top quark degrees of freedom at very high energies. I show that top-built composite Higgs has effective mass squared proportional to the square root of quartic coupling also equal to the top Yukawa coupling and, therefore, naturally seems to imply a "long lived" solution. This observation is bound to energies larger than the low energy HMZC scale. The fact that Higgs mass squared is positive at low energies likely means that quantum corrections due to additional dynamics (most likely affiliated with QCD and top condensation physics) overcome the leading order prediction of the negative Higgs mass squared within the top-quark built composite Higgs model. While explaining hierarchy, this model in 4D doesn't explain the vacuum energy problem and therefore the hypothesis about 2D fundamental theory still holds. In a sense, it should not be surprising that this solution corresponds to the k=2 branch I discovered within the 2D theory considerations.

In Sections 6 and 7, I address the possible physics realizations beyond the standard dogma [1-21]. I introduce several conservative and a few radical models that deal with both external and internal degrees of freedom. Section 6 is mostly concerned with the "second order phase transition," where, in my terminology, renormalized theory doesn't abruptly change the parameters and degrees of freedom across the low energy HMZC scale. Whereas, Section 7 is mostly concerned with the "first order phase transition," where, in my terminology, renormalized theory does abruptly change the parameters and degrees of freedom across the low energy HMZC scale.

In Section 7, a new class of models is introduced within the first order phase transition that are neither Supersymmetric [27-31], Technicolor [95-96, 37], Topcolor [101], TC^2 [102-110], Top Seesaw [111-113] or Little Higgs [125-128] alike. All divergences (logarithmic terms are also zero) between fermions and electroweak gauge bosons loops in the Higgs propagator at energies in the vicinity of the HMZC scale are exactly cancelled therefore *exactly* removing the tachyon solution. Removing tachyon *at all energies larger than* the HMNZ^2 scale however requires dynamical symmetry breaking and composite Higgs beyond SM, see Section 7.1. Standard leading order cancelation is obtained at energies smaller than the HMZC or HMNZ^2 scale. Different levels of physical granularity below and above Λ_{EW} are described with the first order electroweak phase transition $x = 1 \to x = 0.33$ in the notion of Section 3, or $g_{t*} \sim \frac{g_t}{3} \to \sqrt{3} g_{t*} \sim \frac{g_t}{\sqrt{3}}$, $m_H = 0 \to m_H = \frac{2}{3} m_t$ and $\lambda = 0 \to \lambda = \frac{m_H^2}{v_{EW}^2} = \frac{4m_t^2}{9v_{EW}^2} = \frac{2g_t^2}{9}$ in the SM notation. A new class of models, dubbed Composite Particles Model (CPM) has 3 fundamental fermions creating a composite particle that is identified as top quark, whereas 2 fundamental fermions create Higgs. The non-relativistic limit in the relatively "weak" QCD regime suggests $m_H = \frac{2}{3} m_t = 115.4 \pm 0.9\ GeV$.

The Higgs compositeness within CPM appears very different from the Higgs compositeness in the context of $SU(5)/SO(5)$ breaking pattern [131-132]. However, it would be interesting to investigate if there is a way to connect these two composite frameworks.

Model	m_t	m_H	$m_H(GeV)$
k=2 branch	$173.1 \pm 1.3\ GeV$	$\sqrt{\frac{6m_t^2\,(z)-M_Z^2-2M_W^2}{3\left(1+\frac{\pi}{2}\right)}}, z=1\ (\frac{2m_t^2}{V_{EW}^2})$ **(2D)**	143.4 ± 1.3 (142.5 ± 2.4)
Long lived @ Planck	$173.1 \pm 1.3\ GeV$	$\cong 138.5\ GeV$ from μ & λ **(4D)** $\cong 146.0\ GeV$ from μ **(4D)**	138.1 ± 1.8 146.5 ± 2.0
$\mathbf{3^{-1/2}}$ **(k=2)**	$173.1 \pm 1.3\ GeV$	$\sqrt{\frac{4}{3}m_t^2 - M_Z^2 - 2M_W^2}$ **(4D)**	136.8 ± 2.2
k=1 branch	$173.1 \pm 1.3\ GeV$	$\sqrt{\frac{6m_t^2\,(z)-M_Z^2-2M_W^2}{3(1+\pi)}}, z=1\ (\frac{2m_t^2}{V_{EW}^2})$ **(2D)** and scaled to 4D with $1.4676 = \left[\frac{2}{3} - 3\left(x_1 - \frac{y_1}{\pi}\right)\right] / \left(\frac{y_1}{\pi}\right)$	113.0 ± 1.0 (112.3 ± 1.9) 136.9 ± 1.3 (136.0 ± 2.3)
Movers **2D** **4D**	$\sqrt{3\frac{3}{2\sqrt{2}}\frac{M_Z^2+2M_W^2}{6}}$ **(4D)** $\sqrt{\frac{3}{2}(M_Z^2+2M_W^2)}$ **(2D)**	$\sqrt{\frac{1}{3}\frac{M_Z^2+2M_W^2}{3}}$ **(4D)** $\frac{2}{3}m_t$ **(2D)**	145.3 (140.0) 115.4 (111.9) 119.0 (114.6)
$\mathbf{3^{-1/2}}$ **(k=1)** **Flavor** **Color**	$\sqrt{\frac{3}{2}(M_Z^2+2M_W^2)}$ **(2D & 4D)**	$\frac{2}{3}m_t$	115.4 ± 0.9 119.0 115.4 ± 0.9 119.0
$\boldsymbol{\pi^{-1/2}}$	$173.1 \pm 1.3\ GeV$	$\sqrt{\frac{1}{3}\left(\frac{6}{\pi}m_t^2 - M_Z^2 - 2M_W^2\right)}$ **(2D)** & $m_t\sqrt{\frac{2}{\pi}\left(1-\frac{1}{\pi}\right)} \cong \frac{2}{3}m_t$ **(2D)**	109.5 ± 1.4 114.1 ± 0.9

Table 1 The Higgs mass predictions. If Higgs mass below $114\ GeV$ is excluded the obtained values are roughly centered at $116.5\ GeV$, typically associated with the first order phase transition with a single exception, and at $140.5\ GeV$, typically associated with the second order phase transition with a single exception.

The summary of the Higgs mass predictions is presented in Table 1. The obtained values are roughly centered at $116.5\ GeV$, typically associated with the first order phase transition with one exception, and at $140.5\ GeV$, typically associated with the second order phase transition with a single exception.

The theme, which clearly stands out for the first order phase transition, is too dominant top loop; the expected balance between electroweak gauge bosons and top loops in the unbroken phase exists for $\cong 3$ times smaller physical top Yukawa coupling. This observation persist in both 2D and 4D theories. This creates an even stronger mismatch between dynamically generated boson and fermion degrees of freedom than what can be observed by Pagels-Stokar [86] or gap equations in the gauged NJL model [99-100] formalism, the problem that motivated TC^2 [102-110] and Top Seesaw [111-113] models.

One could resolve this problem, and possibly remove tachyons, by introducing additional phase space in the scalar sector hence, lowering the top Yukawa coupling. Or, one could consider CPM in which 3 fundamental fermions create a composite particle that is identified as top quark whereas 2 fundamental fermions create Higgs. The non-relativistic limit in the relatively "weak" QCD regime then suggests $m_H = \frac{2}{3} m_t = 115.4 \pm 0.9\ GeV$. Furthermore, in the zeroth order one could relate the top quark mass to the values of Z and W boson masses alone and obtain prediction of the top quark mass as in Popovic 2001 [26]. This prediction is less than 3% larger than the current world average top quark mass [60].

I would not be surprised if both Higgs and top quark masses in the next to leading order would vary by up to several percents. Furthermore, adding additional structure that could, for example, explain the masses of other particles, could probably account for additional ~5%. Therefore more detailed investigation of this class of models is necessary.

It rarely happens that over constrained system provide such good matches with observations.

In Section 6.3, I obtain an important scaling between the 2D and 4D theories. Thanks to that scaling, the theory structure was conserved across space-times with different dimensions as well across two different regimes below and above the HMZC scale, see section 7.4.

Finally, the general concept of Higgs tachyon solution above the HMZC scale requires much better understanding. According to Wigner [133], the space-like negative mass squared particles have non-compact little groups so their spin is not described by rotation group $SU(2)$ and in difference to massless particles their "spin" may be continuous parameter. Maybe that fact could provide partial explanation of the observed mismatch between the broken and unbroken electroweak phases across the HMZC scale.

In this paper, I point out to (1) the importance of 2D theory in relation to the current problems that particle physics faces, (2) an interesting connection between QCD and top condensation, and (3) two regions of theoretically preferred Higgs mass with accompanying models, see also [33]. I present class of composite models with dynamical symmetry breaking, which I dub the Composite Particles Model (CPM), which may potentially exactly remove the tachyon solution. Finally, I map the physical Higgs mass with HMZC scale $\sim \Lambda_{EWSB}$ based on my earlier work in 2001 [26] and 2002 [33].

As shown the LHC experiment is already stepping outside the known physics territories and one should definitively expect answers that go beyond the current dogma [1-21] in the very near future.

Appendix

Here, I overview details of calculations [26, 33] regarding vacuum stability [41-48, 83], perturbativity [49-50, 83] and the scale $\Lambda_{HMZC}\sim\Lambda_{EWSB}$ affiliated with the physical electroweak phase transition.

In parallel, two independent techniques were utilized: the $\overline{MS}$ scheme [51], applied to the effective potential [52] analysis [53-55], and the Euclidean hard cut-off scheme, applied to the generalized original Veltman's approach [56], confirmed by Osland and Wu [57] and added with logarithmic divergences by Ma [58].

Using a notation adopted here, the tree level potential of the neutral component of the Higgs scalar doublet is

$$V(\Phi) = -\frac{m_H^2}{4}\Phi^2 + \frac{\lambda}{8}\Phi^4 \tag{A1}$$

where $m_H^2 = V^{(2)}|_{\langle\Phi\rangle}$ is the tree level Higgs mass squared and $\langle\Phi\rangle = v_{EW} = 246.2\ GeV$. Hence, the running effective potential is defined as

$$V(\Phi_R) = -\frac{m_H^2(\Lambda\sim\Phi_R)}{4}\Phi_R^2 + \frac{\lambda(\Lambda\sim\Phi_R)}{8}\Phi_R^4 \tag{A2}$$

with the running effective parameters m_H^2 and λ. The connection with effective action at zero external momentum (i.e. the effective potential V_{eff}) is then

$$V_{eff}(\Phi_{cl}) = V(\Phi_R) \tag{A3}$$

where Φ_{cl} is the classical field (on which the generating functional of 1-Particle-Ireducible Green functions depends) corresponding to the running field Φ_R. Obviously, it is the zero-temperature effective potential V_{eff}, and not some particular values of the running effective parameters m_H^2 and λ, that defines the vacuum structure of the theory. If the minimum of V_{eff} is away from zero, the electroweak symmetry is broken.

The original Veltman's approach [56] can be used to describe the running effective potential in Eq. (A2). Veltman reached the conclusion that by redefining mass terms and fields (i.e. running), the SM Lagrangian with one-loop corrections (and with only quadratic divergences considered) may be brought in a gauge invariant fashion to the same form as the tree level Lagrangian. This result was confirmed by Osland and Wu [57] and refined with additional logarithmic terms at the one-loop level by Ma [58] inR_ξ. gauge. Moreover, the running of all the couplings of interest was included [134]. In addition, the higher-loop contributions to quadratic running were calculated in recursive manner [135, 136].

The running effective Higgs mass squared at one-loop level satisfies [56, 58]

$$\frac{dm_H^2}{d\Lambda^2} = \frac{3g_W^2}{64\pi^2 M_W^2}\left(m_H^2 + 2M_W^2 + M_Z^2 - 4\sum_f \frac{n_f}{3} m_f^2\right) \tag{A4}$$

$$+\frac{3g_W^2}{64\pi^2 M_W^2}\frac{m_H^2}{2\Lambda^2}\left(m_H^2 - 2M_W^2 - M_Z^2 + 2\sum_f \frac{n_f}{3} m_f^2\right),$$

where $n_f = 3$ (1) for quarks (leptons). First term corresponds to the famous quadratic divergence. Using the tree level SM relations, the first term may be rewritten in terms of the gauge couplings g_Y, g_W, Yukawa couplings g_f's, and quartic coupling λ as

$$\frac{dm_H^2}{d\Lambda^2} = -\frac{1}{32\pi^2}\left(12g_t^2 - 6\lambda - \frac{9}{2}g_W^2 - \frac{3}{2}g_Y^2\right) + \cdots . \tag{A5}$$

In 2D theory, where both space-time dimensions and Dirac trace were set to 2 (they were both set to 4 in the original calculation [14]), the Higgs mass running in the leading order [33] is proportional to

$$\frac{dm_H^2}{d\Lambda^2} \propto \left(3m_H^2 + 2M_W^2 + M_Z^2 - 6\sum_f \frac{n_f}{3} m_f^2\right), \tag{A6}$$

$$\frac{dm_H^2}{d\Lambda^2} \propto \left(3g_t^2 - 3\lambda - \frac{3}{4}g_W^2 - \frac{1}{3}g_Y^2\right),$$

with logarithmic running in the leading order.

Conversely, in the $\overline{MS}$ scheme, the scalar mass squared $m^2(t)$ in 4D is running just logarithmically!

Does this mean that the SM in the $\overline{MS}$ scheme does not suffer from the quadratic divergences embodied in the running effective Higgs mass squared?

The answer is no, and to show this, it is useful to consider the effective potential V_{eff}. Following the approach in [45-47] in the $\overline{MS}$ scheme and in the 't Hooft-Landau gauge, the renormalization group improved one-loop effective potential [53-55] is

$$V_{eff} = V_0 + V_1 \text{ where } V_0 = -\frac{1}{2}m^2(t)\Phi_R^2(t) + \frac{1}{8}\lambda(t)\Phi_R^4(t) \tag{A7}$$

$$\text{and } V_1 = \sum_{i=1}^{5} \frac{n_i}{64\pi^2}[k_i\Phi_R^2(t) - k_i']^2\left[log\frac{k_i\Phi_R^2(t) - k_i'}{\mu^2(t)} - c_i\right] + \boldsymbol{\Omega(t)}.$$

The values of the parameters n, k, k', and c here are same as in [45-47]. Contribution to the cosmological constant is denoted by $\boldsymbol{\Omega}$ and is assumed irrelevant for the current calculation (though it may have huge importance for the physics of early Universe [59]).

Classical and running Higgs fields are related as

$$\Phi_R(t) = exp\left[-\int_0^t \gamma(t')dt'\right]\Phi_{cl}\ , \tag{A8}$$

$$\text{where } \gamma(t') = \frac{3}{16\pi^2}\left[g_t^2(t') - \frac{1}{4}g_Y^2(t') - \frac{3}{4}g_W^2(t')\right]$$

is the anomalous dimension, and g_t is the top Yukawa coupling.

By using Equ (A2-A3) the running effective Higgs mass squared is extracted from Eqs. (A7-A8) as

$$m_H^2 = -\frac{4V_{eff}(\Phi_{cl}) - \frac{\lambda}{2}\Phi_R^4}{\Phi_R^2} \ . \qquad \text{(A9)}$$

This is a quantity in the $\overline{MS}$ scheme that should be compared with the quadratically unstable Higgs mass squared, Eq. (A4), in the method developed from the original Veltman's approach!

As discussed in [26, 33] results are identical in both the SM dimensional $\overline{MS}$ regularization and in the Veltman's hard-cutoff method, the two most popular and most reliable approaches, to a very high precision with relatively small numerical processing error.

In this study the running Higgs masses squared as obtained from the $\overline{MS}$ effective potential approach [53-55] as well as from the Euclidean hard cut-off generalized Veltman's approach are analyzed in the similar manner: at the one-loop level with the logarithmic terms included and with running of all the couplings of interest at the two-loop level, i.e. in the next-to-leading-log (NTLL) level approximation.

In the original study [26, 33] the strong coupling and the top pole mass were $\alpha_S = 0.1182$ and $m_t = 175\, GeV$ respectively. Here, the results are recalculated for the current world average value $m_t = 173.1\, GeV$ [60].

The matching condition for the running top Yukawa coupling is identical to the one in [45-47]. The one-loop level matching conditions for m^2 and λ in the $\overline{MS}$ effective potential approach, and m_H^2 and λ in the hard cut-off generalized Veltman's approach were obtained from the standard requirement that

$$V_{eff}^{(1)}|_{\langle\Phi\rangle} = 0 \ \text{ and } \ V_{eff}^{(2)}|_{\langle\Phi\rangle} = m_H^2(0) = m_H^2(pole) - \Delta\boldsymbol{\Pi}\left(m_H^2(pole)\right) \qquad \text{(A10)}$$

where $\Delta\boldsymbol{\Pi}\left(m_H^2(pole)\right) = \boldsymbol{Re}\left[\boldsymbol{\Pi}\left(m_H^2(pole)\right) - \boldsymbol{\Pi}(0)\right]$ with $\boldsymbol{\Pi}$ being the renormalized self-energy of the Higgs boson. The reader is directed to reference [45-46] for more details. That approach has been closely followed here with the main results reconfirmed with a very good precision.

In the case of the Euclidean hard cut-off generalized Veltman's approach, the main difference in the matching procedure is in different form of the running effective potential Eq. (A2). The running of the Higgs mass squared is given by Eq. (A4-A5). Although lengthy, the matching procedure is rather trivial.

After the matching conditions are properly set, the gauge and Yukawa (top is only relevant) couplings are set to run at the two-loop level [53, 137-138].

The variations in the physical Higgs mass upper bound as obtained in [26] from the existence of both HMZC scales are due to the variation in α_S and m_t in the linear approximation as

$$\delta m_H[GeV] \cong 1.40\ \delta m_t[GeV] - 360\delta\alpha_S \ \text{ and} \qquad \text{(A11)}$$

$$\delta m_H[GeV] \cong 1.80\ \delta m_t[GeV] - 70\delta\alpha_S$$

in the $\overline{MS}$ scheme and generalized Veltman's method, respectively. The errors are determined in a rather conservative manner. They are obtained from the separate variations of the matching conditions, in the range from the one-loop level matching to the tree level matching, for the three main parameters: m_H^2 (or m^2), λ and g_t. Top quark Yukawa coupling matching condition is mainly responsible for the upper errors. The quartic coupling matching condition tends to bring the results of the two methods closer. The results are rather insensitive to the variations in the mass squared matching conditions. Following the same logic as in [45-47] for the $\overline{MS}$ scheme, the O(3 GeV) uncertainty is also incorporated in response to the requirement that the effective potential V_{eff}, must be renormalization scale independent. The separate errors on the physical Higgs mass have been added in quadrature.

Here I used two completely different regularization methods and I show that the results are essentially identical. How is that possible?

Once again, it is important to make a distinction between the $\overline{MS}$ parameters m and m_H. The $\overline{MS}$ mass parameter m, has intrinsic logarithmic running. Whereas m_H is a quantity derived from the $\overline{MS}$ parameter m and obtained from the one loop corrections to the effective potential. And it is the parameter m_H defined by the tree-level form of the effective potential which runs quadratically.

In the hard cutoff Euclidean regularization scheme the integrals of the type

$$\int d^d p, \;\; \int \frac{d^d p}{p^2}, \;\; \dots \tag{A12}$$

are nonzero while in dimensional regularization due to the dilatation property they are identically zero. As can be shown the two regularization methods however agree in the logarithmic terms. This interplay may be easily seen if one dimensionally continues and regularizes propagators by the method of Pauli-Villars [139]. Then one finds [74] for example for d < 4

$$\int \frac{d^d q}{(2\pi)^2} \frac{1}{q^2 (p+q)^2} \sim \frac{(p^2)^{(d/2)-2} - \Lambda^{d-4}}{8\pi^2 (4-d)} \tag{A13}$$

Fixing the cutoff and taking the limit d=4 one obtains $ln(\Lambda)$ instead of a pole at d-4. Vice versa, by fixing d-4 and taking the cutoff to infinity one obtains the continuation of the initial integral.

Acknowledgement

I thank Bo Morgan for editing part of this manuscript. I am indebted to many great people who taught me how to think about the physics and beyond.

References

[1] S. L. Glashow, Nucl. Phys. 22, 579-588. (1961)

[2] S. Weinberg, Phys. Rev. Lett. 19, 1264-1266. (1967)

[3] A. Salam, Proc. of the 8th Nobel Symposium on `Elementary Particle Theory, Relativistic Groups and Analyticity', Stockholm, Sweden, 1968, edited by N. Svartholm, p. 367-377. (1969)

[4] G. 't Hooft, Nucl. Phys. B33, 173-199. (1971)

[5] G. 't Hooft, Nucl. Phys. B35, 167-188. (1971)

[6] G. 't Hooft and M. J. G. Veltman, Nucl. Phys. B44, 189-213. (1972)

[7] D. J. Gross and F. Wilczek, Phys. Rev. Lett. 30, 1343-1346. (1973)

[8] D. H. Politzer, Phys. Rev. Lett. 30, 1346-1349. (1973)

[9] D. J. Gross and F. Wilczek, Phys. Rev. D8, 3633-3652. (1973)

[10] D. J. Gross and F. Wilczek, Phys. Rev. D9, 980-993. (1974)

[11] D. H. Politzer, Phys. Rep. 14 (1974) 129-180.

[12] C. N. Yang and R. L. Mills, Phys. Rev. 96, 191-195. (1954)

[13] J. Goldstone, Nuovo Cim. 19, 154-164. (1961)

[14] J. Goldstone, A. Salam, and S. Weinberg, Phys. Rev. 127, 965-970. (1962)

[15] P. W. Higgs, Phys. Lett. 12, 132-133. (1964)

[16] F. Englert and R. Brout, Phys. Rev. Lett. 13, 321-322. (1964)

[17] G. S. Guralnik, C. R. Hagen, and T. W. B. Kibble, Phys. Rev. Lett. 13, 585-587. (1964)

[18] J. D. Bjorken and S. L. Glashow, Phys. Lett. 11, 255-257. (1964)

[19] P. W. Higgs, Phys. Rev. 145, 1156-1163. (1966)

[20] T. W. B. Kibble, Phys. Rev. 155, 1554-1561. (1967)

[21] S. Weinberg, Phys. Rev. Lett. 27, 1688-1691. (1971)

[22] A. Einstein, Annalen der Physik 17: 891–921. (1905)

[23]CERN Press release

<http://press.web.cern.ch/Press/PressReleases/Releases2000/PR08.00ELEPRundelay.html> (2000)

[24] P. Renton, Nature 428, 141-144 (2004)

[25] LHC website <http://public.web.cern.ch/public/en/LHC/LHC-en.html> (2010)

[26] M. B. Popovic, Boston University High Energy Physics, BUHEP-01-15, hep-ph/0106355 (2001)

[27] J. Wess and B. Zumino, Nucl. Phys. B70, 39. (1974)

[28] H. P. Nilles, Phys. Reports, **110**, 1 (1984).

[29] H. E. Haber and G. L. Kane, Phys. Reports, **117**, 75 (1985).

[30] J. Wess and J. Bagger, Supersymmetry and Supergravity, Princeton University Press. (1983)

[31] S. Ferrara ed., Supersymmetry , North Holland and World Scientific. (1987)

[32] C. T. Hill, FERMILAB-PUB-80-097-T, Dec 1980. 44pp. Published in Phys.Rev.D24, 691. (1981)

[33] M. B. Popovic, Harvard University Theory Physics, HUTP-02/A012, hep-ph/0204345 (2002)

[34] A. Einstein, "Die Feldgleichungen der Gravitation", Sitzungsberichte der Preussischen Akademie der Wissenschaften zu Berlin: 844–847.(1915)

[35] A. Einstein, Annalen der Physik 49, (1916)

[36] K. G. Wilson and J. Kogut, Phys. Reports, 12C, 75. (1974)

[37] L. Susskind, Phys. Rev. D20, 2619. (1979)

[38] F. Wilczek, “What is time?” talk at MIT, March 2010. (2010)

[39] J. Hogan, Nature 448 (7151), 240–245. (2007)

[40] M. P. Hobson, G. P. Efstathiou, and A. N. Lasenby (2006), “General Relativity: An introduction for physicists” (Reprinted with corrections 2007 ed.). Cambridge University Press. p. 187. (2007)

[41] N. Cabibbo, L. Maiani, G. Parisi, and R. Petronzio, Nucl. Phys. B158, 295. (1979)

[42] M. Lindner, Z. Phys. C31, 295. (1986)

[43] M. Sher, Phys. Rept. 179, 273. (1989)

[44] G. Altarelli and G. Isidori, Phys. Lett. B337, 141. (1994)

[45] J. A. Casas, J. R. Espinosa, and M. Quiros, Phys. Lett. B342, 171. (1995)

[46] J. A. Casas, J. R. Espinosa, and M. Quiros, Phys. Lett. B382, 374. (1996)

[47] M. Quiros, in “Perspective on Higgs Physics II”, G. Kane, ed., World Scientific, Singapore. (1998)

[48] Yu. F. Pirogov and O. V. Zenin, Eur. Phys. J. C10, 629. (1999)

[49] P. Q. Hung and G. Isidori, Phys. Lett. B402, 122. (1997)

[50] T. Hambye and K. Riesselmann, Phys. Rev. D55, 7255. (1997)

[51] W. A. Bardeen, A. J, Buras, D. W. Duke and T. Muta, Phys. Rev. D18, 3998. (1978)

[52] S. Coleman and E. Weinberg, Phys. Rev. D7, 1888. (1973)

[53] C. Ford, D. R. T. Jones, P. W. Stephenson and M. B. Einhorn, Nucl. Phys. B395, 17. (1993)

[54] M. Bando, T. Kugo, N. Maekawa and H. Nakano, Phys. Lett. B301, 83. (1993)

[55] M. Bando, T. Kugo, N. Maekawa and H. Nakano, Prog. Theor. Phys. 90, 405. (1993)

[56] M. Veltman, Acta Phys. Pol. B12, 437. (1981)

[57] P. Osland and T. T. Wu, Z. Phys. C55, 585. (1992)

[58] E. Ma, Phys. Rev. D47, 2143. (1993)

[59] E. W. Kolb and M. S. Turner, “The Early Universe”, Addison-Wesley Publishing Company. (1990)

[60] arXiv:0903.2503v1 [hep-ex] (2009)

[61] American Linear Collider Working Group, Linear Collider Physics Resource Book for Snowmass 2001, section 2.4.1, pp 17-20. (2001)

[62] S. Coleman, "Fate of the false vacuum: Semiclassical theory". Phys. Rev. D15: 2929–36. (1977)

[63] C. Callan and S. Coleman, "Fate of the false vacuum. II. First quantum corrections". Phys. Rev. D16: 1762–68. (1977)

[64] M. Stone, "Lifetime and decay of excited vacuum states". Phys. Rev. D 14: 3568–3573. (1976)

[65] P.H. Frampton, "Vacuum Instability and Higgs Scalar Mass". Phys. Rev. Lett. 37: 1378–1380. (1976)

[66] M. Stone, "Semiclassical methods for unstable states". Phys.Lett. B 67: 186–183. (1977)

[67] P.H. Frampton, "Consequences of Vacuum Instability in Quantum Field Theory". Phys. Rev. D15: 2922–28. (1977)

[68] S. Coleman and F. De Luccia, "Gravitational effects on and of vacuum decay". Physical Review D21: 3305. (1980)

[69] A. H. Guth, "The Inflationary Universe: A Possible Solution to the Horizon and Flatness Problems". Phys. Rev. D23: 347. (1981)

[70] S. W. Hawking and I. G. Moss, "Supercooled phase transitions in the very early universe". Phys. Lett. B110: 35–8. (1982)

[71] A. Linde, "A New Inflationary Universe Scenario: A Possible Solution Of The Horizon, Flatness, Homogeneity, Isotropy And Primordial Monopole Problems". Phys. Lett. B108: 389. (1982)

[72] M.S. Turner and F. Wilczek, "Is our vacuum metastable?" Nature 298: 633–634. (1982)

[73] P. Hut and M.J. Rees, "How stable is our vacuum?" Nature 302: 508–509. (1983)

[74] J. F. Gunion, H. E. Haber, G. Kane, and S. Dawson, “The Higgs Hunter’s Guide,” Addison-Wesley Publishing Company, Redwood City, CA. (1990)

[75] S. Dawson, Four lectures given at NATO Advanced Study Institute on Techniques and Concepts of High Energy Physics, July 11-22, 1996, St. Croix, Virgin Islands [hep-ph/9612229]. (1996)

[76] arXiv:hep-ex/0608032v1 [hep-ex] (2006)

[77] M. Carena, J. R. Espinosa, M. Quiros, and C. E. M. Wagner, Phys. Lett. B335, 209. (1995)

[78] M. Carena, M. Quiros, and C. E. M. Wagner, Nucl. Phys. B461, 407. (1996)

[79] H. E. Haber, R. Hempfling, and A. H. Hoang, Z. Phys. C75, 539. (1997)

[80] R. Barbieri and G. Giudice, Nucl. Phys. B306, 63. (1988)

[81] ALEPH Collab., Phys. Lett. B526, 191. (2002)

[82] ALEPH, DELPHI, L3, and OPAL Collaborations, The LEP Working Group for Higgs Boson Searches, Phys. Lett. B565, 61. (2003)

[83] C. Amsler et al. (Particle Data Group), Physics Letters B667, 1 (2008) and 2009 partial update for the 2010 edition (URL: http://pdg.lbl.gov). (2010)

[84] LEP Electroweak Working Group, status of September 2007, http://lepewwg.web.cern.ch/LEPEWWG/; M. Grunewald, arXiv:0709.3744v1(2007); J. Erler and P. Langacker, Electroweak Model and Constraints on New Physics, in [71]. (2008)

[85] J. Schwinger, Phys. Rev. 128, 2435. (1962)

[86] J. Zinn-Justin, “Quantum Field Theory and Critical Phenomena” 2nd ed., Clarendon Press, Oxford (1993).

[87] C. T. Hill and E. H. Simmons, Physics Reports 381: 235–402. (2003)

[88] J. Bardeen, L. N. Cooper, and J. R. Schrieffer, Phys. Rev. 106, 162 – 164. (1957)

[89] Y. Nambu and G. Jona-Lasinio, Physical Review 122: 345–358. (1961)

[90] Y. Nambu and G. Jona-Lasinio, Physical Review 124: 246–254. (1961)

[91] Y. Nambu, in Z. Adjduk, S. Pokorski, and A. Trautman. Proceedings of the Kazimierz 1988 Conference on New Theories in Physics. XI International Symposium on Elementary Particle Physics. pp. 406–415. (1989)

[92] C. T. Hill, FERMILAB-PUB-80-097-T, Dec 1980. 44pp. Published in Phys.Rev.D24, 691. (1981.)

[93] V.A. Miransky, M. Tanabashi, and K. Yamawaki, DPNU-89-08, Jan 1989. 20pp. Published in Mod.Phys.Lett.A4:1043. (1989)

[94] W. A. Bardeen, C. T. Hill, and M. Lindner, FERMILAB-PUB-89-127-T, Jul 1989. 43pp. Published in Phys.Rev.D41, 1647. (1990)

[95] S. Weinberg, Phys. Rev. D 13, 974. (1976)

[96] S. Weinberg, Phys. Rev. D 19, 1277. (1979)

[97] A. Hasenfratz, P. Hasenfratz, K. Jansen, J. Kuti and Y. Shen, Nucl. Phys. B 365, Issue 1, 79-97. (1991)

[98] H. Pagels and S. Stokar, Phys. Rev. D 20, 2947. (1979)

[99] W. A. Bardeen, C. N. Leung, and S. T. Love, Phys. Rev. Lett. 56, 1230. (1986)

[100] C. N. Leung, S. T. Love and W. A. Bardeen, Nucl. Phys. B273, 649. (1986)

[101] C. T. Hill, Phys.Lett.B266, 419-424. (1991)

[102] C. T. Hill Phys. Lett. B345, 483-489, (1995)

[103] K. Lane and E. Eichten, Phys. Lett. B 352, 382. (1995)

[104] K. D. Lane, Phys. Rev. D 54, 2204. (1996)

[105] G. Buchalla, G. Burdman, C. T. Hill, and D. Kominis, Phys. Rev. D 53, 5185. (1996)

[106] R. S. Chivukula, E. H. Simmons, and J. Terning, Phys. Lett. B331, 383. (1994)

[107] R. S. Chivukula, E. H. Simmons, and J. Terning, Phys. Rev. D53, 5258. (1996)

[108] E. H. Simmons, Phys. Rev. D55, 5494. (1997)

[109] R. S. Chivukula, A. G. Cohen and E. H. Simmons, Phys. Lett. B380, 92. (1996)

[110] M. B. Popovic and E. H. Simmons, Phys. Rev. D 58, 095007. (1998)

[111] B. A. Dobrescu and C. T. Hill, Phys. Rev. Lett. 81, 2634. (1998)

[112] R. S. Chivukula, B. A. Dobrescu, H. Georgi, and C. T. Hill, Phys. Rev. D 59, 075003. (1999)

[113] M.B. Popovic, Phys. Rev. D 64, 035001. (2001)

[114] N. Arkani-Hamed, H.-C. Cheng, B. A. Dobrescu, and L. J. Hall, Phys. Rev. D 62, 096006. (2000)

[115] T. Appelquist, H.-C. Cheng, and B. A. Dobrescu, Phys. Rev. D 64, 035002. (2001)

[116] S. Dimopoulos and L. Susskind, Nucl. Phys. B155, 237. (1979)

[117] E. Eichten and K. Lane, Phys. Lett. 90B, 125. (1980)

[118] R. S. Chivukula, S. B. Selipsky, and E. H. Simmons, Phys. Rev. Lett. 69, 575. (1992)

[119] T. Appelquist, M. J. Bowick, E. Cohler, and A. I. Hauser, Phys. Rev. Lett. 53, 1523. (1984)

[120] M. E. Peskin and D. V. Schroeder, An Introduction to Quantum Field Theory, Addison Wesley Publishing Company. (1995)

[121] S. Bethke, arXiv:0908.1135 [hep-ph]. (2009)

[122] L. H. Ryder, Quantum Field Theory, 2nd edition, Cambridge University Press. (1996)

[123] E.P. Wigner, Annals of Mathematics, 40, 149. (1939)

[124] G. ’t Hooft, Naturalness, chiral symmetry, and spontaneous chiral symmetry breaking. Lecture given at Cargese Summer Inst., Cargese, France, Aug 26 - Sep 8, 1979. (1979)

[125] N. Arkani-Hamed, A. G. Cohen, E. Katz, and A. E. Nelson, Harvard University Theory Physics, HUTP-02/A017, hep-ph/0206021, JHEP 07, 034. (2002)

[126] N. Arkani-Hamed, A. G. Cohen, and H. Georgi, Phys. Lett. B513, 232–240. (2001)

[127] N. Arkani-Hamed, A.G. Cohen, T. Gregoire, E. Katz, A.E. Nelson, and J.G. Wacker, JHEP 08, 021. (2002)

[128] M. Schmaltz, Nucl. Phys. Proc. Suppl. 117 40–49. (2003)

[129] D. B. Kaplan, H. Georgi, and S. Dimopoulos, Phys. Lett. B136 187. (1984)

[130] H. Georgi, D. B. Kaplan, and P. Galison, Phys. Lett. B143 152. (1984)

[131] H. Georgi and D. B. Kaplan, Phys. Lett. B145 216. (1984)

[132] M. J. Dugan, H. Georgi, and D. B. Kaplan, Nucl. Phys. B254 299. (1985)

[133] K. Lane and A. Martin, Phys. Lett. B635, 118-122. (2006)

[134] M. Chaichian, R. G. Felipe and K. Huitu, Phys. Lett. B363, 101. (1995)

[135] M. Einhorn and D. R. T. Jones, Phys. Rev. D46, 5206. (1992)

[136] C. Kolda and H. Murayama, JHEP 0007, 035. (2000)

[137] M. E. Machacek and M. T. Vaughn, Nucl. Phys. B222, 83. (1983)

[138] M. E. Machacek and M. T. Vaughn, Nucl. Phys. B236, 221 (1984)

[139] W. Pauli and F. Villars, Rev. Mod. Phys. 21, 434. (1949)