The hierarchy problem and the cosmological constant problem in the Standard Model

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Abstract

We argue that the SM in the Higgs phase does not suffer form a "hierarchy problem" and that similarly the "cosmological constant problem" resolves itself if we understand the SM as a low energy effective theory emerging from a cut-off medium at the Planck scale. We discuss these issues under the condition of a stable Higgs vacuum, which allows to extend the SM up to the Planck length. The bare Higgs boson mass then changes sign below the Planck scale, such the the SM in the early universe is in the symmetric phase. The cut-off enhanced Higgs mass term as well as the quartically enhanced cosmological constant term trigger the inflation of the early universe. Reheating follows by the heavy Higgses decaying predominantly into top–anti-top pairs, which at this stage are sill massless. Preheating is suppressed in SM inflation as in the symmetric phase bosonic decay channels are absent at tree level. The coefficients of the shift between bare and renormalized Higgs mass $m_{H0}^2 - m_H^2 = \delta m^2 = \frac{\Lambda_{\rm Pl}^2}{16\pi^2}\, C(\mu)$ as well as of the shift between bare and renormalized vacuum energy density $\rho_{\Lambda 0} - \rho_\Lambda = \delta\rho_{\rm vac} = \frac{\Lambda_{\rm Pl}^4}{(16\pi^2)^2}\, X(\mu)$ exhibit close-by zeros $C(\mu) = 0$ at about $\mu_0 \approx 1.4 \times 10^{16}$ GeV or $\mu_0' \approx 7.7 \times 10^{14}$ GeV after Wick rearrangement $C(\mu) \to C'(\mu) = C(\mu) + \lambda(\mu)$ and $X(\mu) = \frac{1}{8}\,(2\,C(\mu) + \lambda(\mu)) = 0$ at $\mu_{\rm CC} \approx 3.1 \times 10^{15}$ GeV. The zero of $C(\mu)$ triggers the electroweak phase transition with $C(\mu) < 0$ in the low energy Higgs phase and $C(\mu) > 0$ in the symmetric phase above the transition point. Since inflation tunes the total energy density to take the critical value of a flat universe $\Omega_{\rm tot} = \rho_{\rm tot}/\rho_{\rm crit} = \Omega_\Lambda + \Omega_{\rm matter} + \Omega_{\rm radiation} = 1$ it is obvious that Ω_Λ today is of order $\Omega_{\rm tot}$ given that $1 > \Omega_{\rm matter}, \Omega_{\rm radiation} > 0$ which saturate the total density to about 26% only, the dominant part being dark matter(21%). Obviously, the SM Higgs system initially provides a

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huge **dark energy** density and the resulting inflation is taming the originally huge cosmological constant to the small value observed today, whatever its initial value was, provided it was large enough to trigger inflation. While laboratory experiments can access physics of the broken phase only, the symmetric phase above the Higgs transition point is accessible though physics of the early universe as it manifests in cosmological observations. The main unsolved problem remains the origin of dark matter.

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1. Introduction

The discovery of the Higgs boson [1, 2] by the ATLAS [3] and CMS [4] experiments at CERN revealed a very peculiar value for the Higgs boson mass, just in a very narrow window which allows to extrapolate the SM way up to the Planck scale [5]. ATLAS and CMS results therefore may "revolution" particle physics in an unexpected way, namely showing that the SM has higher self-consistency (conspiracy) than expected and previous arguments for the existence of new physics may turn out not to be compelling. Also the absence so far of any new physics signal at the LHC may indicate that commonly accepted expectations may not be satisfied. On the one hand it seems to look completely implausible to assume the SM to be essentially valid up to Planck energies, on the other hand the high tide of speculations about physics beyond the SM have been of no avail. One also has to keep in mind that precision tests of the SM already revealed a test in depth of its quantum structure, besides large corrections form the running fine structure constant $\alpha(s)$, the running of the strong coupling $\alpha_s(s)$ and the large top Yukawa $y_t^2(s)$ effect as contributing to the $\rho = G_{\rm NC}/G_\mu(0)$ parameter, subleading corrections amount to a 10 σ deviation form the SM leading order effects predictions. Thus the SM is on very solid grounds better than everything else we ever had.

On the other hand the view that the SM is a low energy effective theory of some cutoff system at the Planck energy scale $M_{\rm Pl}$ appears to be consolidated. A crucial point is that $M_{\rm Pl}$ providing the scale for the low energy expansion in powers $E/M_{\rm Pl}$ is exceedingly high, very far from what we can see! A dimension 6 operator at LHC energies is suppressed by $(E_{\rm LHC}/\Lambda_{\rm Pl})^2 \approx 10^{-30}$.

This seems to motivate a change in paradigm from the view that the world looks simpler the higher the energy to a more natural scenario which understands the SM as the "true world" seen from far away, with symmetries emerging from not resolving the details.

The methodological approach for constructing low energy effective theories we have learned from Ken Wilson's [6, 7, 8] investigations of condensed matter systems and his insight that critical long distance phenomena are governed by emergent quantum field theories. As I will argue in the following, cut-offs in particle physics are important to understand early cosmology [8], such as inflation, reheating, baryogenesis and all that [9, 10, 11, 12, 13]. As in condensed matter physics the connection between macroscopic long distance physics (at laboratory scales) and the microscopic underlying cut-off system (high energy events as they were natural in the early universe) turn out to have a physical meaning.

In this context naturalness arguments play an important role. The SM's naturalness problems and fine-tuning problems have been made conscious by G. 't Hooft [14] long time ago as a possible problem in the relationship between macroscopic phenomena which follow from microscopic laws (a condensed matter system inspired scenario), soon later the "hierarchy problem" had been dogmatized as a kind of fundamental principle. In fact the hierarchy problem of the SM seems to be the key motivation for all kind of extensions of the SM. It is therefore important to reconsider the "problem" in more detail.

One of my key points concerns the different meaning of a possible hierarchy problem has in the symmetric and in the broken phase of the SM. In order to understand the point we have to remember why we need the Higgs in the SM. The Higgs is necessary to get a renormalizable low energy effective electroweak theory [15]. Interestingly, one scalar particle is sufficient to solve the renormalizability problems arising form each of many different massive fields in the SM, of which each causes the problem independently of the others. The point is that this one particle has to exhibit as many new forces as there are individual massive states [16]. All required new interactions are in accordance with the SM symmetry structure in the symmetric phase as we know. The taming of the high energy behavior of course requires the Higgs boson to have a mass in the ballpark of the other given heavier SM states, if it would be much heavier it would not serve its purpose in the low energy regime. Note that the Higgs boson has to cure the unphysical mass effects for the **given** gauge boson masses M_W, M_Z and fermion masses M_f[1], via adequate *Higgs exchange forces*, where the coupling strength is proportional to the mass of the massive field coupled. A very heavy Higgs eventually would decouple and thus miss to restore renormalizability of the massive vector-boson gauge theory. Interestingly, in the symmetric phase the SM gauge-boson plus chiral fermions sector is renormalizable without the Higgs-boson and Yukawa sectors and scalars are not required at all to cure the high energy behavior, because it is renormalizable on its own structure. Therefore, in the symmetric phase the mass-degenerate Higgs fields in the complex Higgs doublet can be as heavy as we like. Since unprotected by any symmetry, naturally we would expect the Higgses indeed to be very heavy. Indeed, the "origin" of the Higgs mass is very different in the broken phase, where the mass is generated by the Higgs mechanism [1, 2] also for the Higgs itself ($m_H^2 = \frac{1}{3}\,\lambda\, v^2$), and in the symmetric phase, where is is dynamically generated by the Planck medium, as we will argue below. Therefore, the usual claim that the SM requires to be extended in such a way that quadratic divergences are absent has no foundation. Purely formal arguments based on perturbative counterterm adjustments do not lead any further.

The hierarchy problem in particular addresses the presence of quadratic ultraviolet (UV) divergences related with the SM Higgs mass term. Infinities in physical theories are the result of idealizations and show up as singularities in a formalism or in models. UV singularities in general plague the precise definition as well as concrete calculations in quantum field theories (QFT). A closer look usually reveals infinities to parametrize our ignorance or mark the limitations of our understanding or knowledge. One particular consequence of UV divergences in local QFTs is that a vacuum energy is ill-defined as it is associated with quartically divergent quantum fluctuations.

This is another indication which tells us that local continuum QFT has its limitation and that the need for regularization is actually the need to look at the true system behind it. In fact the cut-off system is more physical and does not share the problems with infinities which result from the idealization. In any case the framework of a renormalizable QFT, which has been extremely successful in particle physics up to highest accessible energies, is not able to give answers to the questions related to vacuum energy and hence to all questions related to dark energy, accelerated expansion and inflation of the

[1]We denote on-shell masses by capital, $\overline{\rm MS}$ masses by lower case letters as in Ref. [8]

universe.

It is thus natural to consider the Standard Model to be what we observe as the low energy effective SM (LEESM), the renormalizable tail of the real cutoff system sitting at the Planck scale. As a consequence all properties required by renormalizability, gauge symmetries, chiral symmetry, anomaly cancellation naturally emerge as a consequence of the low energy expansion. The infinite tower of higher order operators becomes invisible, and only a few operators are effectively observable, which makes the world look much simpler. In reality infinities are replaced by eventually very large but finite numbers, and I will show that sometimes such huge effects are needed to understand the real world. I will argue that cutoff enhanced effects are responsible for triggering the Higgs mechanism not very far below the Planck scale and the inflation of the early universe.

The Planck medium, *the ether*, is characterized by a fundamental cutoff $\Lambda_{\rm Pl}$ or equivalently the Planck mass $M_{\rm Pl}$ which derive from the basic fundamental constants, the speed of light c characterizing special relativity, the Planck constant $\hbar$ intrinsic to quantum physics and Newton's constant G_N the key parameter of gravity. **Unified** they provide an intrinsic length $\ell_{\rm Pl}$, the Planck length, which also translates into the Planck time $t_{\rm Pl}$ and the Planck temperature $T_{\rm Pl}$.

The history of our universe we can trace back 13.7 billion years close to the Big Bang, when the expansion of the universe was ignited in a "fireball", an extremely hot and dense state when all structures and at the end all atoms, nuclei and nucleons were disintegrated to a world of elementary particles only. Besides the missing cold dark matter (DM), one of the last piece which was missing in the SM, the Higgs boson, now seems to provide a new milestone in our understanding of the dynamics of the very early universe.

I think questions concerning the early universe can be addressed only in the LEESM "extension" of the SM as such, given by a local QFT supplied by cutoff effects in a minimal way. As we know, in a renormalizable QFT all renormalized quantities as a function of the renormalized parameters and fields in the limit of a large cut-off are finite and devoid of any cut-off relicts! Here, it is adequate to remember the Bogoliubov-Parasyuk renormalization theorem which states that renormalized Green's functions and matrix elements of the scattering matrix (S-matrix) are free of ultraviolet divergences. It implies that in the low energy world cut-off effects are not accessible to experiments and a "problem" like the hierarchy problem is not a statement which can be checked to exist as a observable conflict.

To my knowledge the only non-perturbative definition of a renormalizable local quantum field theory is the possibility to put in on a lattice. This again may be taken as an indication that the need for a cut-off actually is an indication that the cutoff exists in the real(er) world. In this sense the lattice QFT is the true(er) system than its continuum tail. Of course, there are many ways to introduce a cut-off and actually we cannot know what the cutoff system looks truly. This is not a real problem if we are interested in the long range patterns mainly, the only thing we have to care is that the underlying system is in the *universality class* of the SM.

2. The hierarchy problem revisited

The hierarchy problem cannot be addressed within the renormalizable and renormalized SM, which is what we can confront with experiments. In this framework all independent parameters are free and have to be supplied by experiment.

In the LEESM "extension" of the SM bare parameter turn into physical parameters of the underlying cut-off system as the "true world" at short distances. Then the hierarchy problem is the problem of "tuning to criticality" which concerns the relevant operators of dimension < 4, in particular the mass terms. In the symmetric phase of the SM, where there is only one mass (the others are forbidden by the known chiral and gauge symmetries), the one in the potential of the Higgs doublet field, the fine tuning to criticality has the form

$$m_0^2(\mu_1 = M_{\rm Pl}) = m^2(\mu_2 = M_H) + \delta m^2(\mu_1, \mu_2) \; ; \;\; \delta m^2 = \frac{\Lambda_{\rm Pl}^2}{16\pi^2}\, C(\mu) \qquad (1)$$

with a coefficient typically $C = O(1)$. To keep the renormalized mass at some small value, which can be seen at low energy, the bare m_0^2 has to be adjusted to compensate the huge number δm^2 such that about 35 digits must be adjusted in order to get the observed value around the electroweak scale. Is this a real problem?

One thing is apparent: our fine-tuning relation exhibits quantities at very different scales, the renormalized one at low energy and the bare one at the Planck scale. In the LEESM both are observable, in principle. In fact, if we consider a renormalization condition like (1), our presumed fine-tuning relation, if we want to test it experimentally, it is not possible to test it by low energy experiments only. Low energy experiments only allow us to test relations between measurable renormalized quantities.While in a renormalizable theory, relations between measurable quantities are devoid of any cutoff effects, this changes when we perform high energy experiments at a scale sensitive to the cutoff. Although such experiments are not possible with down to earth accelerators, the expansion of the universe has provided us a scan from Planck energies down the 3 °K of the Cosmic Microwave Background (CMB) [17] radiation. Thus, in the early universe a relation like (1) has a direct physical meaning. In the SM at low energies we are in the broken phase, where $m_{H0} \approx \delta m_H$ is huge negative, when looked at in a formal perturbative bookkeeping. However, this is not something we can test by observation. If we want to test it we have to go to the short distance scale, which however automatically flips the sign of δm_H^2 and we automatically end up in the symmetric phase, where the relation gets a different meaning. The key observation is that the running SM parameters conspire in such a way that the Higgs mass counterterm as well as the vacuum energy counterterm exhibit a zero, which provides a matching point between the bare short distance world and the renormalized low energy world. While the low energy world is what we access by laboratory experiments, the high energy world is what has shaped the early universe with the observable consequences we see in cosmology, specifically in the CMB, Baryogenesis, Nucleosynthesis and structure formation. We thus are able to learn more about the short distance structure by making use of the early universe as a natural accelerator.

In the Higgs phase, there is no hierarchy problem [18] (see also [19]). It is true that in the relation $m_{H0}^2 = m_H^2 + \delta m_H^2$ both m_{H0}^2 and δm_H^2 are many many orders of magnitude larger than m_H^2 . However, in the broken phase one automatically obtains $m_H^2 \propto v^2(\mu_0)$, which is $O(v^2)$ not $O(M_{\rm Pl}^2)$, irrespective of what the, at this scale unobservable, objects of the bare theory are. Thus, in the broken phase the Higgs is naturally light. That the Higgs mass likely is $O(M_{\rm Pl})$ in the symmetric phase is what promotes the Higgs to a candidate for the inflaton.

One indeed can avoid artificial large numbers to show up by choosing μ_0 as a renormalization point, where $\delta m^2 = 0$ and $m^2(\mu_0) = m_0^2(\mu_0)$ and after the EW phase transition and corresponding vacuum rearrangement $m_H^2(\mu_0) = 2m^2(\mu_0) = \frac{1}{3}\lambda(\mu_0)v^2(\mu_0)$ and then get the physical Higgs mass by standard RG running and matching. Certainly not the most practical way to implement M_H as a physical input parameter. The point is that in principle it is possible circumvent fine-tuning. We also note that unlike in regularized renormalizable QFT thinking, m_0^2 is not a given basic parameter to be adjusted by renormalization. In the LEESM the bare mass $m_0^2(\mu)$ as an effective mass, dynamically generated in the Planck medium, is obviously also a running mass. Not far below the scale μ_0 the universe undergoes the EW phase transition (a point of no-analyticity) and the Higgs mass is generated by a different mechanism: the Higgs mechanism, the Higgs mass being given now by $m_H^2 = \frac{1}{3}\,\lambda\,v^2$ after vacuum rearrangement [20]. In the broken phase the hierarchy problem is a pseudo problem.

In the broken phase, characterized by the non-vanishing Higgs field vacuum expectation value (VEV) $v(\mu) \neq 0$, all the masses are determined by the well known mass-coupling relations

$$\begin{aligned} m_W^2(\mu^2) &= \frac{1}{4}\,g^2(\mu^2)\,v^2(\mu^2)\,;\;\; m_Z^2(\mu^2) = \frac{1}{4}\,(g^2(\mu^2)+g'^2(\mu^2))\,v^2(\mu^2)\,; \\ m_f^2(\mu^2) &= \frac{1}{2}\,y_f^2(\mu^2)\,v^2(\mu^2)\,;\;\; m_H^2(\mu^2) = \frac{1}{3}\,\lambda(\mu^2)\,v^2(\mu^2)\,. \end{aligned} \tag{2}$$

Here we consider the parameters in the $\overline{\text{MS}}$ renormalization scheme, μ is the $\overline{\text{MS}}$ renormalization scale, which we have to identify with the energy scale of the physical processes or equivalently with the corresponding temperature in the evolution of the universe. The RG equation for $v^2(\mu^2)$ follows from the RG equations for masses and massless coupling constants using one of these relations. As a key relation we use [21]

$$\mu^2 \frac{d}{d\mu^2} v^2(\mu^2) = 3\,\mu^2 \frac{d}{d\mu^2}\left[\frac{m_H^2(\mu^2)}{\lambda(\mu^2)}\right] \equiv v^2(\mu^2)\left[\gamma_{m^2} - \frac{\beta_\lambda}{\lambda}\right], \tag{3}$$

where $\gamma_{m^2} \equiv \mu^2 \frac{d}{d\mu^2} \ln m^2$ and $\beta_\lambda \equiv \mu^2 \frac{d}{d\mu^2}\lambda$. We write the Higgs potential as $V = \frac{m^2}{2}\,H^2 + \frac{\lambda}{24}H^4$, which fixes our normalization of the Higgs self-coupling. When the m^2-term changes sign and λ stays positive, we know we have a first order phase transition. Funny enough, the Higgs get its mass from its interaction with its own condensate! and thus gets masses in the same way and in the same ballpark as the other SM species. As mentioned before the Higgs mass cannot by much heavier than the other heavier particles if renormalizability is to be effective at low and moderate energies. The interrelations (2) also show that for fixed v, as determined by the Fermi constant $G_\mu = 1/(\sqrt{2}\,v^2)$, the Higgs cannot get too heavy if perturbation theory should remain applicable. An extreme point of view claims that naturalness requires all particles to have masses $O(M_{\rm Pl})$ i.e. $v = O(M_{\rm Pl})$. This would mean that the symmetry is not restored at the cutoff scale and the notion of spontaneous symmetry breaking (SSB) would be obsolete as a concept! The SM's successful structure relies on a symmetric Lagrangian, and a ground state which breaks the symmetry of the Lagrangian. The ground state is not residing at the cutoff scale and breaks the symmetry in such a way that the UV structure is not affected. In view that $v \equiv 0$ above the EW phase transition point, why should it be natural to expect that v jumps from 0 to $O(M_{\rm Pl})$ during the phase transition. We also note that as $v \to 0$ all masses vanish, with the exception of the Higgs mass which acquires a large value in the symmetric phase (see below).

The Higgs VEV v is an **order parameter** resulting form long range **collective behavior** and can be as small as we like. Prototype is the magnetization M as a function of temperature T in a ferromagnetic spin system, where $M = M(T)$ and actually $M(T) \equiv 0$ for $T > T_c$ and furthermore $M(T) \to 0$ as $T \lesssim T_c$. For a direct non-perturbative check in case of the SM, one would put the SM in the unitary gauge on a lattice and simulate its long range properties. The Higgs boson VEV is then a well defined physical order parameter. Difficulties related to Elitzur's theorem [22] thus can be avoided.

Small $v/M_{\rm Pl} \ll 1$ just means we are close to just below a $2^{\rm nd}$ order phase transition point, which is not unnatural if we take into consideration that long range behavior of condensed matter systems are effective quantum field theories in a vicinity of second order phase transition points [6].

In the mass renormalization relation (1) the renormalized mass measures the distance from the critical bare mass m_{0c} for which the renormalized m is zero: thus $m^2 = m_0^2 - m_{0c}^2$. A particle is seen at low energy only if it is light. In the symmetric phase (short distance regime) is is natural to have both m_0^2 and δm^2 large, but why $m_0^2 \approx \delta m^2$ to such high precision?

At very high energy we see the bare system and the Higgs field is a **collective field** which acquires its effective mass via radiative effects $m_0^2 \approx \delta m^2$ near below $M_{\rm Pl}$. In particle physics a radiatively induced mass in known from the Coleman-Weinberg mechanism [23], now in the symmetric phase and applied to the Planck medium. Such mechanism, which is natural in this context, eliminates a possible fine-tuning problem at all scales. There are many examples in condensed matter systems, like the effective mass of the photon in the superconducting phase (Meissner effect) or the effective mass of the effective field which encodes the spin-singlet electron (Cooper) pairs in the Ginzburg-Landau model [24] of superconductivity. The latter directly corresponds to the Abelian Higgs model.

3. Running SM parameters trigger the Higgs mechanism

We remind that all dimensionless couplings satisfy the same renormalization group (RG) equations in the broken and in the unbroken phase and are not affected by any power cutoff dependencies. The evolution of SM couplings in the $\overline{\mathrm{MS}}$ scheme up to the Planck scale has been investigated in Refs. [5, 25, 26, 27, 28, 29, 30, 31, 32, 33, 34] recently, and has been extended to include the Higgs VEV and the masses in Refs. [35, 8]. Except for g', which increases very moderately, all other couplings decrease and stay positive up to the Planck scale. This strengthens the reliability of perturbative arguments and reveals a stable Higgs potential up to the Planck scale [35, 8]. While most analyses [26, 27, 28, 31, 34] are predicting that for the given Higgs mass value vacuum stability is nearby only (meta-stability), and actually fails to persist up to the Planck scale, our evaluation of the matching conditions yields initial $\overline{\mathrm{MS}}$ parameters at the Z boson mass scale which evolve preserving the positivity of λ. Thereby the critical parameter is the top quark Yukawa coupling, for which we find a slightly lower value. My $\overline{\mathrm{MS}}$ input at M_Z is [8] $g_3 = 1.2200$, $g_2 = 0.6530$, $g_1 = 0.3497$, $y_t = 0.9347$ and $\lambda = 0.8070$. At M_{Pl} I get $g_3 = 0.4886$, $g_2 = 0.5068$, $g_1 = 0.4589$, $y_t = 0.3510$ and $\lambda = 0.1405$. In view of the fact that the precise meaning of the experimentally extracted value of the top quark mass is not free of ambiguities, usually it is identified with the on-shell mass M_t (see e.g. [35] and references therein), it may be premature to claim that instability of the SM Higgs potential is a proven fact already. I also think that the implementation of the matching conditions is not free of ambiguities, while the evolution of the couplings over many orders of magnitude is rather sensitive to the precise values of the initial couplings. Accordingly, all numbers presented in this article depend on the specific input parameters adopted, as specified in Ref. [35, 8]. In case the Higgs self-coupling has a zero $\lambda(\mu^2) = 0$, at some critical scale μ_c below M_{Pl}, we learn from Eq. (3), or more directly from $v(\mu^2) = \sqrt{6m^2(\mu^2)/\lambda(\mu^2)} \stackrel{\lambda\to+0}{\to} \infty$ that the SM looses it meaning above this singular point oft non-analyticity.

Running couplings can affect dramatically the quadratic divergences and the interpretation of the hierarchy problem. Quadratic divergences have been investigated at one loop in Ref. [36] (see also [37, 38, 39]), at two loops in Refs. [40, 42, 41]. At n loops the quadratic cutoff dependence is of the form

$$\delta m_H^2 = \frac{\Lambda^2}{16\pi^2}\, C_n(\mu) \tag{4}$$

where the n-loop coefficient only depends on the gauge couplings g', g, g_3, the Yukawa couplings y_f and the Higgs self-coupling λ. Neglecting the numerically insignificant light fermion contributions, the one-loop coefficient function C_1 may be written as

$$C_1 = 2\,\lambda + \frac{3}{2}\,g'^2 + \frac{9}{2}\,g^2 - 12\,y_t^2 \tag{5}$$

and is uniquely determined by dimensionless couplings. The latter are not affected by quadratic divergences such that standard RG equations apply. Surprisingly, as first pointed out in Ref. [42], taking into account the running of the SM couplings, the coefficient of the quadratic divergences of the bare Higgs mass correction can vanish at some scale. In contrast to our evaluation Hamada et al. actually find the zero to lie above the Planck scale. In our analysis, relying on matching conditions for the top quark mass analyzed in [35], we get a scenario where $\lambda(\mu^2)$ stays positive up to the Plank scale and looking at the relation between the bare and the renormalized Higgs mass we find C_1 and hence the Higgs mass counterterm to vanish at about $\mu_0 \sim 1.4 \times 10^{16}$ GeV, not very far below the Planck scale. The next-order correction, first calculated in Refs. [40, 41] and confirmed in [42] read

$$\begin{aligned} C_2 \;\; = \;\; & C_1 + \frac{\ln(2^6/3^3)}{16\pi^2}\,[18\,y_t^4 + y_t^2\,(-\frac{7}{6}\,g'^2 + \frac{9}{2}\,g^2 - 32\,g_s^2) \\ & -\frac{87}{8}\,g'^4 - \frac{63}{8}\,g^4 - \frac{15}{4}\,g^2 g'^2 + \lambda\,(-6\,y_t^2 + g'^2 + 3\,g^2) - \frac{2}{3}\,\lambda^2]\,, \end{aligned} \tag{6}$$

numerically does not change significantly the one-loop result. The same results apply for the Higgs potential parameter m^2, which corresponds to $m^2 \hat{=} \frac{1}{2}\,m_H^2$ in the broken phase. For scales $\mu < \mu_0$ we have δm^2 large negative, which is triggering spontaneous symmetry breaking by a negative bare mass $m_0^2 = m^2 + \delta m^2$, where m again denotes the renormalized mass. At $\mu = \mu_0$ we have $\delta m^2 = 0$ and the sign of δm^2 flips, implying a phase transition to the symmetric phase. Finite temperature effects, which must be included in a realistic scenario, turn out not do to change the gross features of our scenario [8].

4. The SM cosmological constant and dark energy

It is crucial that in the early universe both terms in the Higgs potential are positive in order to condition slow-roll inflation during long enough time.

In fact the quadratically and quartically cutoff enhanced terms in the Higgs potential enforce the condition $\frac{1}{2}\dot{\phi}^2 \ll V(\phi)$ and given the Higgs boson pressure $p_\phi = \frac{1}{2}\dot{\phi}^2 - V(\phi)$ and the Higgs energy density $\rho_\phi = \frac{1}{2}\dot{\phi}^2 + V(\phi)$, we arrive at the equation of state $w = p/\rho \approx -1$ characteristic for dark energy and the equivalent cosmological constant (CC) (see e.g. [43, 44, 46, 45] and references therein). A first measurement of the dark energy equation of state $w = -1.13^{+0.13}_{-0.10}$, has been obtained by the Planck mission [47] recently.

A key point is that in the LEESM scenario the vacuum energy is a calculable quantity. In the symmetric phase $SU(2)$ symmetry implies that $\Phi^+\Phi$ is a singlet such that the invariant vacuum energy is given just by simple Higgs loops [20]

$$\langle 0|\Phi^+\Phi|0\rangle = \frac{1}{2}\langle 0|H^2|0\rangle \equiv \frac{1}{2}\,\Xi\,;\;\; \Xi = \frac{\Lambda_{\rm Pl}^2}{16\pi^2}\,. \tag{7}$$

A Wick type of rearrangement of the Lagrangian then provides a CC represented by $V(0) = \langle V(\phi)\rangle = \frac{m^2}{2}\,\Xi + \frac{\lambda}{8}\,\Xi^2$ and a mass shift $m'^2 = m^2 + \frac{\lambda}{2}\,\Xi$. For our values of the $\overline{\rm MS}$ input parameters the zero in the Higgs mass counter term gets shifted as $\mu_0 \approx 1.4\times 10^{16}$ GeV $\to \mu_0' \approx 7.7\times 10^{14}$ GeV . We notice that the SM predicts huge CC at $M_{\rm Pl}$: $\rho_\phi \simeq V(\phi) \sim 2.77\,M_{\rm Pl}^4 \sim 6.13\times 10^{76}$ GeV^4 exhibiting a very weakly scale dependence (running couplings) and we are confronted with the question how to get ride of this huge quasi-constant? An intriguing structure again solves the puzzle. The effective CC counterterm has a zero, which again is a point where renormalized and bare quantities are in agreement:

$$\rho_{\Lambda 0} = \rho_\Lambda + \frac{M_{\rm Pl}^4}{(16\pi^2)^2}\,X(\mu) \tag{8}$$

with $X(\mu) \simeq \frac{1}{8}\,(2C(\mu) + \lambda(\mu))$ which has a zero close to the zero of $C(\mu)$ when $2\,C(\mu) = -\lambda(\mu)$. Note that $C(\mu) = -\lambda(\mu)$ is the shifted Higgs transition point.

Again we find a matching point between low energy and high energy world: $\rho_{\Lambda 0} = \rho_\Lambda$ where the memory of the quartic Planck scale enhancement gets lost, as it should be as we know that the low energy phase does not provide access to cutoff effects.

At the Higgs transition as soon as $m'^2 < 0$ for $\mu < \mu_0'$ the vacuum rearrangement of Higgs potential takes place. As a result at the minimum ϕ_v of the potential we should get $V(0) + V(\phi_v) \sim (0.002~{\rm eV})^4$ about the observed value of today's CC. How can this be? Indeed, at the zero of $X(\mu)$ we have $\rho_{\Lambda 0} = \rho_\Lambda$ and one may expect that like the Higgs boson mass another free SM parameter is to be fixed by experiment here. One might expect ρ_Λ to be naturally small, since the $\Lambda_{\rm Pl}^4$ term is nullified at the matching point. Note that the huge cutoff prefactors act as amplifiers of small changes in the effective SM couplings. But how small we should expect the low energy effective CC to be? In fact, in the LEESM scenario neither the Higgs mass nor the CC are really free parameters in the low energy world. The Higgs mass, more precisely the Higgs self coupling, has to be constrained to a window where the Higgs potential remains stable up to the Planck scale, and the CC which triggers inflation gets tuned down by inflation to lie in the ballpark of the critical density of a flat universe.

5. Inflation and reheating

In contrast to standard scenarios of modeling the evolution of the early universe, SM cosmology is characterized by the fact that almost everything is known, within uncertainties of the parameters and perturbative approximations. In LEESM cosmology the form of the potential is given by the bare SM Higgs potential $V(\phi) = \frac{m^2}{2}\,\phi^2 + \frac{\lambda}{24}\,\phi^4$, the parameters are known, calculable in terms of the low energy parameters, the only unknown is the magnitude of the Higgs field. The latter must be large – trans-Planckian – in order to get the required number of *e-folds* N (expansion factor $a(t_e)/a(t_i) = \exp H\,(t_e - t_i) = \exp N$, where $a(t)$ is the Friedmann-Robertson-Walker radius of the universe, t_i the begin of inflation and t_e the end of inflation and H the Hubble constant). For our set of $\overline{\rm MS}$ input parameters we require $\phi_0 = \phi(\mu = M_{\rm Pl}) \approx 4.5\,M_{\rm Pl}$. At start the slow-roll condition $V(\phi) \gg \frac{1}{2}\,\dot{\phi}^2$ is well satisfied, by the fact that in the symmetric phase the mass term as well as $V(0) = \langle V(\phi)\rangle$ are huge. Because of the large initial field strength ϕ_0, however, the interaction term is actually dominating for a short time after the initial Planck time $t_{\rm Pl}$. The field equation $\ddot{\phi} + 3H\dot{\phi} = -V'(\phi)$ then predicts a dramatic decay of the field, $\phi(t) = \phi_0\,{\rm e}^{E_0\,(t-t_0)}$ with $E_0 = \sqrt{2\lambda}/(3\sqrt{3}\ell) \approx 4.3 \times 10^{17}$ GeV , $V_{\rm int} \gg V_{\rm mass}$ and shortly after $E_0 = m^2/(3\ell\sqrt{V(0)}) \approx 6.6 \times 10^{17}$ GeV , $V_{\rm mass} \gg V_{\rm int}$ [$\ell^2 = 8\pi G_N/3$], such that in almost no time, still under slow-roll conditions, the mass term dominates and for what follows the field equation is a damped harmonic oscillator. The universe thus undergoes an epoch of Gaussian inflation [48] before the oscillations set in. In the symmetric phase the four Higgses are very heavy and rather unstable. The Higgses decay predominantly (largest Yukawa couplings) into as yet massless top–anti-top pairs and lighter fermion–anti-fermion pairs $H \to t\bar{t}$, $b\bar{b}$, $\cdots$ thereby reheating the young universe which just had been cooled down dramatically by inflation. Preheating is suppressed in SM inflation as in the symmetric phase bosonic decay channels like $H \to WW$ and $H \to ZZ$ are absent at tree level. The CP violating decays $H^+ \to t\bar{d}$ [rate $\propto y_t y_d V_{td}$] $H^- \to b\bar{u}$ [rate $\propto y_b y_u V_{ub}$] likely are important for baryogenesis. After the electroweak phase transition which closely follows the Higgs transition where m^2 in the Higgs potential changes sign, the now heavy top quarks decay into normal matter as driven by CKM [49] couplings and phase space. At these scales the $B + L$ violating dimension 6 operators [50, 51, 52] can still play a key role for baryogenesis and together with decays like $t \to d e^+ \nu$ provide CP violating reactions during a phase out of thermal equilibrium. For details see [8, 18, 20].

6. Remark on Trans-Planckian Higgs fields

If the SM Higgs is the inflaton, sufficient inflation requires trans-Planckian magnitude Higgs field at the Planck scale. At the Planck scale the low energy expansion obviously gets obsolete and likely we cannot seriously argue with field monomials and the operator hierarchy appearing in the low energy expansion. What is important is that the field is decaying very fast. Formally, given a truncated series of operators in the potential, the highest power is dominating in the trans-Planckian regime. One then expects that for some time the ϕ^4 term of the potential is dominating, the decay of the field is then exponential, for higher dimensional operators it is faster than exponential, such that the field very rapidly reaches the Planck- and sub-Planck regime. This means that the mass term is dominating after a very short period and before the kinetic term becomes relevant and slow-roll inflation ends. So fears that in low energy effective scenarios with trans-Planckian fields higher order operators would mess up things are not in any sense justified[2]. Obviously,

[2]The constructive understanding of LEETs we have learned from Ken Wilson's renormalization semi-group, based on integrating out short distance fluctuations. This produces all kinds, mostly of irrelevant higher order interactions. A typical example is the Ising model, which by itself seen as the basic microscopic system has simple nearest neighbor interactions only and by the low energy expansion develops a tower of higher order operators, which at the short distance scale are simply absent altogether. Such operators don't do any harm at the intrinsic short distance scale.

without the precise knowledge of the Planck physics, very close to the Planck scale we never will be able to make a precise prediction of what is happening. This however seems not to be a serious obstacle to quantitatively describe inflation and its properties as far as they can be accessed by observation. The LEESM scenario in principle predicts not only the form of the effective potential not far below the Planck scale but also its parameters and the only quantity not fixed by low energy physics is the magnitude of the field at the Planck scale. We also have shown that taking into account the running of the parameters is mandatory for understanding inflation and reheating and all that.

Trans-Planckian fields are not unnatural in a low energy effective scenario as the Planck medium exhibits a high temperature and temperature fluctuations make higher excitations not improbable. While the Planck medium will never be accessible to direct experimental tests, a phenomenological approach to constrain its effective properties is obviously possible, especially by CMB data [53].

In the extremely hot Planckian medium, the Hubble constant in the radiation dominated state with effective number $g_*(T) = g_B(T)+\frac{7}{8}\,g_f(T) = 102.75$ of relativistic degrees of freedom is given by $H = \ell\sqrt{\rho} \simeq 1.66\ (k_BT)^2\,\sqrt{102.75}$ $M_{\rm Pl}^{-1}$, at Planck time $H_i \simeq 16.83\,M_{\rm Pl}$ such that a Higgs field of size $\phi_i \simeq 4.51\,M_{\rm Pl}$, is not unexpected and could as well also be larger.

Often it is argued that trans-Planckian field are unnatural in particular in a LEET scenario. I cannot see any argument against strong fields and LEET arguments (ordering operators with respect to a polynomial expansion and their dimension) completely loose their sense when $E/\Lambda_{\rm Pl} \gtrsim 1$.

Provided the Higgs field decays fast enough towards the end of inflation we expect the mass term to be dominant such that a Gaussian fluctuation spectrum prevails. The quasi-constant cosmological constant $V(0)$ at these times mainly enters the Hubble constant H and does not affect the fluctuation spectrum.

7. The self-organized cosmological constant

In principle, like the Higgs mass in the LEESM, also ρ_Λ is expected to be a free parameter to be fixed by experiment. This we would expect in a standard Einstein-Friedman cosmology not exhibiting a large CC term. However, the situation is different if an inflaton exists providing an appropriate CC. Inflation is fine-tuning the total energy density to unity $\Omega_{\rm tot} = 1$, in units of the critical density defined as a boarder density between a positively curved universe exhibiting sufficient mass to stop matter to escape for ever and a negative curved universe with too little mass to stop the expansion. Thus inflation means that the universe after the inflation phase has a particular energy density $\rho_0 = \rho_{EdS}$ of a flat Einstein-de Sitter universe. This is remarkable as the value is fixed irrespective of the initial energy density. If there is too much "mass" space adjusts itself such that the same universal density is reached after the inflation epoch. This also solves the cosmological constant problem of the SM. The typical problem is that in general one gets a CC which is way too big and this looks to be a tremendous fine-tuning problem. For the SM this concerns the contribution to the vacuum density via the Higgs VEV in the broken phase, as well as the contributions from spontaneous breakdown of chiral symmetry, which are much to big and even of wrong sign. However, if inflation is at work, the final vacuum density is fixed irrespective of it initial density. Given that $\Omega_{\rm tot} = \Omega_\Lambda + \Omega_{\rm DM} + \Omega_{\rm BM} + \Omega_{\rm rad} = 1$ with $1 > \Omega_{\rm DM} > \Omega_{\rm BM} > \Omega_{\rm rad} > 0$ we know that Ω_Λ must be of order $\Omega_{\rm tot}$. As a non-vanishing $\rho_{\Lambda 0}$ is needed to provide inflation in any case, it is not unlikely that the other components contributing to the total energy density do not saturate the bound. This means that the fine-tuning is dynamically enforced by inflation and the value $\Lambda_{\rm obs} = \Omega_\Lambda \times \kappa \simeq 1.6517 \times 10^{-56}$ cm^{-2} $[\kappa = \frac{8\pi G_N}{3c^2}]$ is actually not far from the critical total energy density of the flat universe $\rho_{0,\rm crit} = \rho_{\rm EdS} = \frac{3H_0^2}{8\pi G_N} = 1.878 \times 10^{-29}\, h^2\, \mathrm{gr/cm}^3$ where H_0 is the present Hubble constant, and $h = 0.67 \pm 0.02$ its value in units of 100 km s^{-1} Mpc^{-1}. Today's dark energy density is $\rho_{0\Lambda} = \Omega_\Lambda\, \rho_{0,\rm crit}$, with $\Omega_\Lambda = 0.74(3)$. While $\Omega_{\rm rad}$ and very likely $\Omega_{\rm BM}$ are essentially SM predictions if we include the $B + L$ violating dimension 6 four-fermion operators, $\Omega_{\rm DM}$ is the only missing piece which remains an open problem and definitely requires additional beyond the SM physics. This also concerns contributions from quark and possible gluon condensates, which we do not explicitly consider here.

So the solution of the cosmological constant problem is a dynamical consequence of inflation and overlarge looking values at earlier epochs in the evolution of the universe just means that the CC is a time-dependent dynamical quantity. Provided the SM for the specific conspiring input parameters yields a stable Higgs potential, inflation and the CC itself are SM ingredients leading to a highly self-consistent conspiracy which shapes the universe.

This scenario only seems to be natural if the running of the SM couplings is not affected by physics beyond the SM in a substantial way. This

LEESM scenario does not exclude the existence of new physics, however, possible new effects should be natural in the low energy effective framework. Still axions are a very good candidate to solve the strong CP problem and eventually provide the missing dark matter. Similarly, it would be natural that the right-handed singlet neutrinos are Majorana fermions, which naturally would exhibit a large Majorana mass and trigger a sea-saw mechanism which would explain the lightness of the neutrinos. Also an unbroken confining $SU(4)$ sector could be there forming bound mesons which could constitute the missing dark matter. Like normal baryonic matter is essentially hadronic binding energy, bound in fermions, also dark matter could be condensed energy, bound in bosons. In contrast a supersymmetric or grand unified extension of the SM would not fit into the picture. Surprisingly also a fourth family would completely deteriorate our scenario. I think two points are very much in favor of a change of the game; Higgs vacuum stability or very close to stability and why should we need two different scalar field, one for the Higgs mechanism and one as an inflaton, if one can do what we need? I agree that it is against any reasonable expectation to believe that the SM should hold up to the Big Bang. However, fact is that also cosmological and astrophysical observations have not given any definite hint for non-SM physics with the big exception of dark matter.

8. Summary

The Higgs has two different functions in our world: 1) it has to render the effective low energy electroweak theory (massive vector-boson and fermion sector) renormalizable. In the broken low energy phase the Higgs acquires the vacuum condensate which provides masses to all massive fields including the Higgs boson itself. Key point are the many new Higgs-exchange forces necessary to render the low energy amplitudes renormalizable. 2) in the symmetric phase the four very heavy Higgses generate a huge dark energy, which causes inflation. After inflation has ended and we are out of equilibrium the Higgses are decaying predominantly into the heaviest fermions pairs which provides the reheating of the inflated universe. The universe cooling further down then pushes the universe into the Higgs phase, where the particles acquire their masses. The predominating heavy quarks decay into the light ones which later form the baryons and normal matter. This scenario is possible because of the quadratically enhanced Higgs boson mass and the quartically enhanced dark energy, which show up in the symmetric phase of the SM before the Higgs transition. The existence of such relevant operator effects in my opinion are supported by observation, in particular by observed inflation patterns, meaning that the hierarchy as well as the cosmological constant "problems" reflect important properties of the SM needed to understand the evolution of the early universe (for different opinions see [54, 55, 56, 57, 58]). Consolidation of our bottom-up path to physics near the Planck scale will sensibly depend on progress in high precision physics around the EW scale v. Especially, Higgs and a top-pair factories will play a key role in this context.

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