Stability and UV completion of the Standard Model

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Abstract – The knowledge of the electroweak vacuum stability condition is of the greatest importance for our understanding of beyond Standard Model physics. It is widely believed that new physics that lives at very high-energy scales should have no impact on the stability analysis. This expectation has been recently challenged, but the results were controversial as new physics was given in terms of non-renormalizable higher-order operators. Here we consider for the first time new physics at extremely high-energy scales (say close to the Planck scale) in terms of renormalizable operators, in other words we consider a sort of toy UV completion of the Standard Model, and definitely show that its presence can be crucial in determining the vacuum stability condition. This result has important phenomenological consequences, as it provides useful guidance in studying beyond Standard Model theories. Moreover, it suggests that very popular speculations based on the so-called “criticality” of the Standard Model do not appear to be well founded.

Introduction. – The analysis of the electroweak (EW) vacuum stability condition is of the greatest importance for present research in theoretical and experimental particle physics, as they are crucial in guiding experimental and theoretical search for physics beyond the Standard Model (BSM). It is well known that, due to the top loop corrections, the Higgs effective potential $V(\phi)$ bends down for values of $\phi > v$, where $v \sim 246$ GeV is the location of the EW minimum, and develops a second minimum at a very large value $\phi_{\text{min}}^{(2)}$. The potential $V(\phi)$ is obtained by considering SM interactions only [1–11], and depending on the values of the Higgs and top masses, $M_H$ and $M_t$, the second minimum can be higher or lower than the EW one.

When $V(\phi_{\text{min}}^{(2)}) > V(v)$, the EW minimum is the stable vacuum state, while for $V(\phi_{\text{min}}^{(2)}) < V(v)$, it is a metastable state, in other words a false vacuum, and in order to decide the stability condition of such a state we have to calculate its lifetime $\tau$. Figure 1 shows the usual stability phase diagram in the $M_H$-$M_t$ plane. The stability, instability, and metastability regions are respectively for: $V(\phi_{\text{min}}^{(2)}) < V(\phi_{\text{min}})$; $V(\phi_{\text{min}}^{(2)}) < V(v)$ and $\tau < T_U$ ($T_U =$ age of the Universe); $V(\phi_{\text{min}}^{(2)}) < V(v)$ and $\tau > T_U$.

When $V(\phi_{\text{min}}^{(2)}) < V(v)$, the instability scale $\phi_{\text{inst}}$ of the Higgs potential is the value of the field where $V(\phi_{\text{inst}}) = V(v)$, and $V(\phi) < V(v)$ for $\phi > \phi_{\text{inst}}$. For the present central experimental values of the Higgs and top masses, $M_H \sim 125.09$ GeV and $M_t \sim 173.34$ GeV [12,13], it turns out that $\phi_{\text{inst}} \sim 10^{11}$ GeV $\gg v$, $\phi_{\text{min}}^{(2)} \sim 10^{30}$ GeV, and $\tau$ is much larger than $T_U$.

Naturally, new physics interactions are expected to have an effect long before the scale $\phi_{\text{min}}^{(2)} \sim 10^{30}$ GeV is reached. In this respect, it has to be noted that the usual analysis outlined above is done under the following assumptions: i) new physics shows up only at very high-energy scales, possibly close to the Planck scale; ii) despite the presence of these new physics interactions, $\tau$ can be calculated with the potential obtained with SM interactions only [8,14]. It is argued in fact that the relevant scale for tunneling is the instability scale $\phi_{\text{inst}} \sim 10^{11}$ GeV, and that the contribution to the tunneling rate coming from very high scale physics ($\gg \phi_{\text{inst}}$) should be suppressed (decoupling) [14].

Parametrizing new physics interactions in terms of few higher-order non-renormalizable operators, it has been shown that the presence of new physics can strongly modify the stability analysis of the EW vacuum, even if it occurs at very high-energy scales [15–17]. However it was observed that these results might depend on the parametrization of new physics in terms of non-renormalizable operators, and that it could be suspected that, when the infinite tower of higher dimensional (new physics) operators of the renormalizable UV completion of
the SM is considered, this effect should disappear, so that the expected decoupling of high-energy from low-energy physics is recovered. Actually, it was even suggested that this effect could take place above the physical cutoff, where the control of the theory is lost [18].

In this letter we consider for the first time a fully renormalizable model, that serves as a toy renormalizable UV completion of the SM, where new physics interactions live at scales much higher than the instability scale $\phi_{\text{inst}}$, and perform the stability analysis of the EW vacuum for this model. Before entering the details of the model, we begin with anticipating the main results of the present work. On the one hand we will be able to definitely show that the stability condition of the EW vacuum strongly depends on new physics even if it lives at very high-energy scales, on the other hand we will understand why the expected decoupling [14] does not take place.

As stressed above, understanding whether or not new physics at very high energy scales has an impact on the stability condition of the EW vacuum is of crucial importance for present research of fundamental physics beyond the SM, and it is the main motivation for the present work.

The model. – In order to illustrate the phenomenon discussed above in a simple framework, we consider a toy renormalizable UV-completion of the SM that is certainly unrealistic but very appropriate to the purposes of the present work. New physics at very high-energy scales is inserted by adding to the SM a scalar field $S$ and a fermion field $\psi$ that interact in a simple way (see below) with the Higgs field $\phi$, with masses $M_S$ and $M_f$ well above the instability scale, $M_S, M_f \gg \phi_{\text{inst}}$. Apart from the kinetic terms, the additional terms in the Lagrangian are

$$\Delta \mathcal{L} = \frac{M_S^2}{2} S^2 + \frac{\lambda_S}{4} S^4 + g_S \phi^2 S^2 + M_f \bar{\psi} \psi + g_f \bar{\psi} \psi.$$  \hspace{1cm} (1)

To understand how a new physics Lagrangian of this kind can arise in a physical setup, we note that the large mass term $M_f$ corresponds to a sort of heavy right-handed “neutrino” in the framework of a see-saw mechanism. While the corresponding light “neutrino” is totally harmless for the stability of the EW vacuum, the heavy “neutrino” can play an important role in destabilizing the vacuum. The scalar field $S$ counterbalance the destabilizing effect of $\psi$. Models with new scalar fields coupled to the Higgs (although admittedly unrealistic) have already been used to provide stabilization of the Higgs effective potential [19,20].

A relevant issue related to the study of physical observables involved in the vacuum stability analysis is the gauge dependence of the SM effective potential (and more generally of effective potentials) away from the extrema [21–23]. In particular, absolute stability bounds on the Higgs mass (that are formally gauge independent) turn out to be gauge dependent at any order of perturbation theory. Only when a consistent resummation is considered the result can be made gauge independent, and this provides an improvement for mass bounds [23]. Moreover, the gauge dependence of the vacuum instability scale has been investigated, with the range of uncertainty identified [22], and it is not known if it is possible to calculate this quantity in a gauge-independent manner.

Finally, for the main quantity of interest to us, namely the decay rate from the false to the true vacuum, again we know that it is a formally gauge-invariant quantity. In a truncated perturbative expansion, however, order-by-order gauge independence can possibly be achieved only after resumming the appropriate terms, as it was done for the energies at the minima of the effective potential [21]. Having these warnings in mind and waiting for improvements on these gauge dependence issues, we proceed now with the analysis of our model following the usual pattern.

For the purposes of the present work, it is sufficient to consider the impact of these additional terms on the Higgs effective potential $V(\phi)$ at the one-loop level only. We do not need a better level of precision as we are not interested in extracting numbers but we only want to illustrate the destabilization effect that arises from very high-energy physics. The one-loop contribution to $V(\phi)$ from these terms is

$$V_1(\phi) = \frac{(M_S^2 + 2g_S \phi^2)}{64\pi^2} \left[ \ln \left( \frac{M_S^2 + 2g_S \phi^2}{M_S^2} \right) - \frac{3}{2} \right]$$

$$- \frac{(M_f^2 + g_f^2 \phi^2)}{16\pi^2} \left[ \ln \left( \frac{M_f^2 + g_f^2 \phi^2}{M_S^2} \right) - \frac{3}{2} \right], \hspace{1cm} (2)$$

where the renormalization scale is $\mu = M_S$ (see below for comments on this choice).

According to the usual analysis based on the decoupling argument [8,14], these new physics interactions at very high-energy scales ($M_S, M_f \gg \phi_{\text{inst}} \sim 10^{11}$GeV) should have no impact on the stability analysis. We now
investigate this question by considering two choices for the parameters of our toy UV-completion of the SM. In both of the examples that we are going to consider a second minimum deeper than the EW one is formed, and so we calculate the EW (false) vacuum lifetime.

**Results.** We now impose to the modified potential

\[ V(\phi) = \frac{1}{4} \phi^4 + V_1(\phi) \] (as usual the quadratic term can be neglected as we consider very high values of \( \phi \)) the matching conditions at the threshold scale \( M_f \) so that the SM Higgs potential is recovered for values of \( \phi < M_f \). The EW vacuum lifetime \( \tau \) is then given by [17]

\[ \tau = \min_{\mu} \left( \frac{1}{T^2_{\mu} \mu^2} \exp \left( \frac{8g^2}{3|h^2_{\mu}| + 4V_1(\mu)/\mu^4} \right) \right), \] (3)

where \( \lambda_{SM}(\mu) \) is the running quartic coupling, and \( V_1(\phi) \) is nothing but the additional contribution (2) to the Higgs potential with the \( \phi^2 \) and \( \phi^4 \) terms subtracted.

To be more precise, \( V_1(\phi) \) is obtained by expanding the potential \( V_1(\phi) \) in powers of \( \phi \), so to isolate first the constant, the \( \phi^2 \) and the \( \phi^4 \) terms. Then at the threshold scale \( M_f \) we require: i) that the renormalized cosmological constant \( \Lambda \), that is given by the sum of all the constant terms (those coming from the SM potential and those coming from \( V_1(\phi) \)) vanishes, \( \Lambda(\mu = M_f) = 0 \); ii) that the renormalized mass term, given by the sum of all the coefficients of \( \phi^2 \), and identified with the SM mass parameter \( m^2_{SM}(\mu = M_f) \) at the scale \( M_f \), vanishes: \( m^2_{SM}(\mu = M_f) \sim 0 \) (more precisely we neglect this term to a very high degree of accuracy for the large values of \( \phi \) considered); iii) that the renormalized quartic coupling, given by the sum of all the coefficients of \( \phi^4 \), is identified with the SM quartic coupling at the scale \( M_f \), \( \lambda_{SM}(\mu = M_f) \). In other words, at the scale \( M_f \) this coefficient is matched with the value of the quartic coupling obtained by considering the running of the renormalization group equations for the SM couplings alone.

Note that the above requirements for the renormalized cosmological constant and mass are well known features. In fact, apart from the fine tuning problem, for the renormalized \( \Lambda \) we can practically consider that \( \Lambda(\mu = 0) \sim \Lambda(\mu = M_f) \sim 0 \). The same is true for the mass term. Again apart from the fine tuning, for the renormalized mass term we can take \( m^2(\mu = 0) \sim m^2(\mu = M_f) \sim 0 \). This simply means that we neglect the \( \phi^2 \) term as compared to the \( \phi^4 \) and other terms for these large values of \( \phi \), and also that the running of the renormalized mass is totally harmless in this respect. For the quartic coupling we have a true matching condition. In fact we require that at the threshold scale \( \mu = M_f \) the quartic coupling coincides with \( \lambda_{SM}(\mu = M_f) \), i.e., with the value that is obtained by running the renormalization group equations for the SM couplings only. Practically starting from the scale \( M_f \) the potential has also the contribution coming from \( V_1(\phi) \) (more precisely from \( V_1(\phi) \)). All this is well summarized in the form of the complete potential, \( V_{tot}(\phi) = \frac{1}{4} \lambda_{SM}(\phi) \phi^4 + V_1(\phi) \), that is used in eq. (3).

We also have to note that the SM quartic coupling at very high values of the running scale experiences a plateau. In particular \( \lambda \) has practically the same value in the whole range \([M_f, M_S] \). Therefore, even though strictly speaking we should consider \( \lambda_{SM}(M_f) \) as the threshold value, we can equally well write \( \lambda_{SM}(M_S) \), and this is what we do in the following. Moreover, in the range \([M_f, M_S] \) we have also runned the new RG equations that contain the beta functions modified by the presence of new couplings, starting with the threshold conditions that at the scale \( M_f \) the couplings take their SM values (obtained by solving the RG equations for the SM couplings only). We found that the changes in the running of the SM couplings due to these modified beta functions are totally negligible.

Before going on with the calculation of \( \tau \) for our toy model, it is worth to remind that when the presence of these new physics interactions (usually assumed to be harmless for calculating \( \tau \)) is neglected, for the present central values of the Higgs and top masses, \( M_H = 125.09 \text{ GeV} \) and \( M_t = 173.34 \text{ GeV} \) [12,13], the calculation gives

\[ \tau \sim 10^{6000} T_U. \] (4)

This result is the basis for the so-called metastability scenario, according to which although the EW minimum is a metastable state (and then a false vacuum), as its lifetime turns out to be much larger than the age of the universe, we may well live in such a state.

Figure 1 shows the full stability analysis done under the assumption that new physics at scales \( \gg \phi_{\text{inst}} \) has no impact on the stability condition of the EW vacuum [8–11]. The black dot corresponds to the tunneling time of eq. (4). The ellipses give the one, two and three sigma experimental uncertainties.

We move now to the computation of the EW vacuum lifetime for our model. More precisely, we are considering two specific examples for different choices of the model parameters. In our first example we choose \( M_S = 1.2 \cdot 10^{18} \text{ GeV}, M_f = 0.6 \cdot 10^{17} \text{ GeV}, g_S(M_S) = 0.97, g_f^2(M_S) = 0.48, \lambda_{SM}(M_S) = -0.0151 \). Note that the latter is the value of the running quartic coupling \( \lambda_{SM}(\mu) \) at the scale \( \mu = M_S \), obtained by considering the RG equations for the SM coupling constants and the boundary conditions at the next-to-next to leading order [7,10,24].

Concerning the choice of the renormalization scale \( \mu \), note that for our potential the one-loop corrections turn out to be small and practically independent on \( \mu \) as long as this scale is chosen within the range \([10^{16} \text{ GeV}, M_S \sim 10^{18} \text{ GeV}] \), and that the value of \( \mu \) that gives the minimum of eq. (3) turns out to be very close to \( M_S \), thus justifying our choice \( \mu = M_S \). Note also that, at the order of approximation that we are considering, \( \lambda_S \) plays no role.

For the values of the parameters given above, the Higgs effective potential \( V(\phi) \) develops a new minimum, lower than the EW one, at \( \phi_{\text{min}} \sim 0.4 \cdot 10^{19} \text{ GeV} \). To investigate the stability condition of the EW vacuum, we have then to
calculate its lifetime $\tau$. For the present central experimental values of the Higgs and top masses ($M_H = 125.09 \text{ GeV}$ and $M_t = 173.34 \text{ GeV}$) we find

$$\tau \sim 10^{180} T_U. \quad (5)$$

This result has to be compared with the tunneling time of eq. (4), obtained by considering the SM potential alone (no new physics included). Although for this example the tunneling time is still much higher than the age of the Universe, eq. (5) gives a result that is greatly different from the one of eq. (4).

Let us consider now another example, with a different value of $M_f$. More precisely, let us take $M_S = 1.2 \cdot 10^{18} \text{ GeV}$, $M_f = 2.4 \cdot 10^{15} \text{ GeV}$, $g_S(M_S) = 0.97$, $g_0^2(M_S) = 0.48$, and $\lambda_{SM}(M_S) = -0.0151$. Again for the present central experimental values of $M_H$ and $M_t$ we find

$$\tau \sim 10^{-65} T_U. \quad (6)$$

In this case, the situation is more dramatic than in the previous example: the tunneling time turns out to be much smaller than the age of the Universe. If realistic, this model could not be accepted.

The results (4), (5), and (6) contain a very important lesson. Despite of the fact that in our toy model new physics lives at very high-energy scales, scales much higher than the instability scale $\phi_{\text{inst}}$, the expectation that the tunneling time should be insensitive to it, in other words that the result shown in eq. (4) should not be modified by the presence of new physics at high energies, is not fulfilled.

From eqs. (5) and (6) we also see that while $M_S$ is kept fixed, the impact on the tunneling time is more pronounced for the lower scale. The reason is that when the matching for the quartic coupling to the SM value is done at the lower scale $M_f^{(1)}$ ($< M_f^{(2)}$), the destabilization of the SM potential occurs before the case when it is done at the higher one $M_f^{(2)}$, and the distance between the destabilizing scale $M_f$ and the rescuing scale $M_S$ is larger for $M_f^{(1)} < M_f^{(2)}$. This is not necessarily the general case, as there are other parameters that determine whether the potential is more or less destabilized. For values of $M_f$ lower than the ones that we have considered, the impact of the running coupling could be more important, thus changing the above picture. However this is not our concern here, as what we want to show is that the expected universality (decoupling) in the calculation of the tunneling time does not apply.

In general new physics fermionic degrees of freedom coupled to the Higgs boson tend to destabilize the EW vacuum, while bosonic degrees of freedom tend to stabilize it. The global result for the stability condition of the EW vacuum then depends on this competition, in particular on the values of the coupling constants involved and on the distance between the fermionic and bosonic relevant scales.

Clearly the crucial point is to understand why the decoupling argument fails. The reason is that the decoupling theorem applies when we calculate scattering amplitudes at energies $E$ lower than $M_S$ and $M_f$. For such physical quantities the contributions from this high-energy new physics is certainly suppressed by factors as $E/M_S$ and $E/M_f$ to the appropriate power.

However in our case we are calculating the tunneling time $\tau$, and tunneling is a non-perturbative phenomenon. More precisely, the bulk of the contribution to $\tau$ is given by the exponential that appears in eq. (3), whose argument is the action calculated for the bounce solution to the euclidean Euler-Lagrange equation of motion [25]. Technically, this is the saddle point for the path integral involved in the calculation of $\tau$, and for such tree level contribution no suppression factors of the kind $(E/M_S)^n$ and $(E/M_f)^n$ can ever appear: if the Higgs effective potential is modified by the presence of terms as those of eq. (2), the new bounce is certainly different from the one obtained when this term is absent. The action is modified and (once exponentiated) gives rise to a value for $\tau$ that can be enormously different from the one obtained when new physics is not considered.

This is a central result of the present work. With the help of a fully renormalizable toy UV completion of the SM, we have shown that the EW vacuum lifetime strongly depends on new physics even if the latter lives at very high-energy scales, much higher than the instability scale $\phi_{\text{inst}} \sim 10^{11} \text{ GeV}$.

This result is at odds with a widely diffused expectation that is based on a naive application of the decoupling argument, and definitely shows that the fact that the stability condition of the vacuum depends on physics that lives at very high-energy scales is not due to an illegitimate extrapolation of the theory beyond its validity, as it was previously thought [18]. On the contrary, it is an illegitimate application of the decoupling argument to a phenomenon to which it cannot be applied, namely the (non-perturbative) tunneling phenomenon, that leads to the expectation that physics at scales much higher than the instability scale $\phi_{\text{inst}}$ should have no impact on the stability condition.

Going back to the examples above, we note that for other choices of $M_f$ and $M_S$ we would obtain other values of $\tau$. In particular, for certain values of $M_f$ and $M_S$, the tunneling time can be even greater than the SM value. This can only occur if both scales ($M_f$ and $M_S$) lie below the value taken by the bounce calculated (with SM interactions only) at its center. On the contrary, if $M_S$ lies above this latter scale, the tunneling time turns out to be less or at best equal to the SM result.

We are now ready to proceed with our analysis. Figure 1 shows the stability diagram in the $M_H - M_t$ plane obtained under the assumption (decoupling argument) that the stability analysis does not depend on high-energy physics. The examples considered above, with the results (4), (5) and (6), indicate that we should expect that the whole
Fig. 2: This figure shows the stability diagram for the EW vacuum in the $M_H-M_t$ plane when the toy (very high-energy) UV completion of the SM given in eq. (1) is considered. The Higgs effective potential is modified by the presence of the additional term (2). Here $M_S = 1.2 \cdot 10^{18}$ GeV, $M_t = 0.6 \cdot 10^{17}$ GeV, $g_S = 0.97$, $g_t^2 = 0.48$. As in fig. 1, the $M_H-M_t$ plane is divided into three sectors: absolute stability, metastability and instability regions. The presence of the term (2) causes the lowering of the instability and stability lines. However, in this case the instability line goes below the experimental point, signaling that the model, for these values of the parameters, cannot be considered as a viable UV completion of the SM. The grey lines are for reference and are the stability and instability lines of fig. 1.

The stability phase diagram actually depends on new physics, even if it lives at very high-energy scales.

The dashed and the dash-dotted lines of fig. 1 are named the stability line and the instability line, respectively. The first one is obtained for those values of $M_H$ and $M_t$ such that the two minima are at the same height, the latter is obtained for the case when $V(\phi_{\text{min}}^{(2)}) < V(\nu)$ and $\tau = T_U$.

Let us repeat now the stability analysis when (1) is added to the SM Lagrangian, so that the term (2) is added to the Higgs effective potential. In fig. 2, the analysis is performed for the values of the parameters considered in our first example, namely $M_S = 1.2 \cdot 10^{18}$ GeV, $M_t = 0.6 \cdot 10^{17}$ GeV, $g_S(M_S) = 0.97$, $g_t^2(M_S) = 0.48$, $\lambda_{\text{SM}}(M_S) = -0.0151$.

We note that the instability line moves downwards. This result had to be expected from the previous results (4) and (5) for the tunneling time. In fact, we obtained $\tau \sim 10^{180} T_U$ for the “UV completed” Higgs potential and $\tau \sim 10^{690} T_U$ for the SM Higgs potential. It is clear that in the case of the UV completed potential, the experimental point (black dot) must be closer to the instability line than in the case of the unmodified potential. The grey lines of fig. 2 are the old instability and stability lines for the the unmodified Higgs potential (see fig. 1).

Actually, another important effect is that even the stability line moves downwards (see fig. 2). When it was thought that a decoupling effect assured that new physics at high scales could not modify this diagram, many speculations were triggered by the fact that the experimental point (black dot in the figure) $M_H \sim 125.09$ GeV and $M_t \sim 173.34$ GeV lies “close”, within 2–3 sigma, to the stability line.

In this respect, it was suggested that more refined measurements of $M_t$ and $M_H$ should allow to determine whether the EW vacuum is a stable or a metastable state. Some authors even went to the point to consider the closeness of the experimental point to the stability line as the most important message from LHC [11], speculating on this closeness and elaborating on it for model building [18].

The results presented in this letter show that the stability condition of the EW vacuum is much more sensitive to high-energy new physics than to the values of the Higgs and top masses. Therefore, more refined measurements of $M_t$ and $M_H$, that are clearly very important on their own, will not allow to determine the stability condition of the EW vacuum.

Moreover, speculations and model building inspired by the so-called “criticality condition”, the closeness of the experimental point to the stability (also called critical) line, are actually unjustified. As we have seen, new physics even if it lives at very high energies (we certainly expect new physics at least at very high energies, maybe Planck scale) can enormously modify the distance between the experimental point and the critical line.

Finally, in fig. 3 the stability diagram for our model with the values of the parameters considered in our second example ($M_S = 1.2 \cdot 10^{18}$ GeV, $M_t = 2.4 \cdot 10^{15}$ GeV, $g_S(M_S) = 0.97$, $g_t^2(M_S) = 0.48$, $\lambda_{\text{SM}}(M_S) = -0.0151$) is shown. The stability and instability lines move downwards as for the previous case. In this case, however, the tunneling time for the experimental point is much shorter than the age of the Universe, see eq. (6), and in fact we see that the experimental point is now inside the instability region.

This result means that the model with these values of the parameters cannot be considered as a viable UV completion of the SM. Naturally, as we have stressed several times, we are considering a toy and by no means realistic extension of the SM. However the lesson of this result is clear. We have shown that for the stability analysis we cannot rely on the decoupling of high-energy physics,
therefore motivated and/or “realistic” candidate BSM theories that contain very high-energy new physics have to pass a sort of “stability test”; only those with a stable or metastable (but with $\gamma > T_U$) vacuum can be considered.

Conclusions. – With the help of a renormalizable toy extension of the SM, we have definitely shown that new physics interactions, even if they live at energies much higher than the scale where the Higgs potential becomes unstable (the so-called instability scale $\phi_{\text{inst}} \sim 10^{11}$ GeV), have strong impact on the stability condition of the EW vacuum. With respect to previous analyses, here new physics interactions are given in terms of a fully renormalizable theory rather than with the help of higher-order non-renormalizable operators, and this makes the conclusions of the present work really robust. These results have far-reaching phenomenological consequences, providing very useful guidance for BSM model building. In particular, they show that speculations based on the so-called “criticality” are not well founded.

REFERENCES

[1] Cabibbo N., Maiani L., Parisi G. and Petronzio R., Nucl. Phys. B, 158 (1979) 295.
[2] Flores R. A. and Sher M., Phys. Rev. D, 27 (1983) 1679; Lindner M., Z. Phys., 31 (1986) 295; Sher M., Phys. Rep., 179 (1989) 273; Lindner M., Sher M. and Zaglauer H. W., Phys. Lett. B, 228 (1989) 139.
[3] Ford C., Jones D. R. T., Stephenson P. W. and Einhorn M. B., Nucl. Phys. B, 395 (1993) 17.
[4] Sher M., Phys. Lett. B, 317 (1993) 159.
[5] Altarelli G. and Isidori G., Phys. Lett. B, 337 (1994) 141.
[6] Casas J. A., Espinosa J. R. and Quirós M., Phys. Lett. B, 342 (1995) 171; 382 (1996) 374.
[7] Bezrukov F., Kalmykov M. Yu., Kniehl B. A. and Shaposhnikov M., JHEP, 10 (2012) 140.
[8] Isidori G., Ridolfi G. and Strumia A., Nucl. Phys. B, 609 (2001) 387.
[9] Elias-Miro J., Espinosa J. R., Giudice G. F., Isidori G., Riotto A. and Strumia A., Phys. Lett. B, 709 (2012) 222.
[10] Degrazii G., Di Vito S., Elias-Miro J., Espinosa J. R., Giudice G. F., Isidori G. and Strumia A., JHEP, 08 (2012) 098.
[11] Buttazzo D., Degrazii G., Giardino P. P., Giudice G. F., Sala F., Salvio A. and Strumia A., JHEP, 12 (2013) 089.
[12] ATLAS Collaboration and CMS Collaboration (Aad G. et al.), Phys. Rev. Lett., 114 (2015) 191803.
[13] ATLAS, CDF, CMS and D0 Collaborations, First combination of Tevatron and LHC measurements of the top-quark mass, arXiv:1403.4427 [hep-ex].
[14] Espinosa J. R., Giudice G. F. and Riotto A., JCAP, 05 (2008) 002 (arXiv:0710.2484 [hep-ph]).
[15] Branchina V. and Messina E., Phys. Rev. Lett., 111 (2013) 241801 (arXiv:1307.5193 [hep-ph]).
[16] Branchina V., Messina E. and Platania A., JHEP, 09 (2014) 182 (arXiv:1407.4112 [hep-ph]).
[17] Branchina V., Messina E. and Sher M., Phys. Rev. D, 91 (2015) 013003 (arXiv:1408.5302 [hep-ph]).
[18] Espinosa J. R., Giudice G. F., Morgeante E., Riotto A., Senatore L., Strumia A. and Tetrads N., arXiv:1505.04825 [hep-ph].
[19] Hung P. Q. and Sher Marc, Phys. Lett. B, 374 (1996) 138.
[20] Casas J. A., Di Clemente V. and Quiros M., Nucl. Phys. B, 581 (2000) 61.
[21] Andreassen A., Frost W. and Schwartz M. D., Phys. Rev. Lett., 113 (2014) 241801 (arXiv:1408.0292 [hep-ph]).
[22] Di Luizo L. and Mhaila L., JHEP, 06 (2014) 079 (arXiv:1404.7450 [hep-ph]).
[23] Andreassen A., Frost W. and Schwartz M. D., Phys. Rev. D, 91 (2015) 016009 (arXiv:1408.0287 [hep-ph]).
[24] Mhaila L. N., Salomon J. and Steinhausser M., Phys. Rev. Lett., 108 (2012) 151602; Chetyrkin K. and Zoller M., JHEP, 06 (2012) 033; Bednyakov A. V., Pikelner A. F. and Velizhanin V. N., Phys. Lett. B, 723 (2013) 336; Chetyrkin K. G. and Zoller M. F., JHEP, 04 (2013) 091; 09 (2013) 155; Kniehl B. A. and Pikelner A. F. and Veretin O. L., Nucl. Phys. B, 896 (2015) 19.
[25] Coleman S., Phys. Rev. D, 15 (1977) 2929; Callan C. G. and Coleman S., Phys. Rev. D, 16 (1977) 1762.