

Higgs Stability and the 750 GeV Diphoton Excess

Alberto Salvio^{a,b}, Anupam Mazumdar^c

^a *Theoretical Physics Department, CERN, Geneva, Switzerland*

^b *Departamento de Física Teórica, Universidad Autónoma de Madrid
and Instituto de Física Teórica IFT-UAM/CSIC, Madrid, Spain*

^c *Consortium for Fundamental Physics, Lancaster University, Lancaster, LA1 4YB, UK*

Abstract

We study the implications of a possible unstable particle with mass $M_X < \text{TeV}$ for the Higgs stability, naturalness and inflation. We pay particular attention to the case $M_X \approx 750 \text{ GeV}$, suggested by recent results of ATLAS and CMS on diphoton final states, and work within the minimal model: we add to the Standard Model field content a pseudoscalar and a vector-like fermion carrying both color and electric charge. This can stabilize the electroweak vacuum without invoking new physics at very high energies, which would give an unnaturally large contribution to the Higgs mass. We also show that inflation can be obtained via a UV modification of General Relativity.

KEYWORDS: Large Hadron Collider, Higgs boson, Inflation

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1 Introduction

The experiments at the LHC have recently released the first studies at the highest energy scale ever reached in a collider, $\sqrt{s} = 13$ TeV. Besides confirming the Standard Model (SM) predictions in many observables, the ATLAS [1] and CMS [2] collaborations have reported an excess in the diphoton channel at 750 GeV, with a local significance of 3.6σ and 2.6σ , respectively, possibly due to a resonance, whose width Γ might be relatively large: ATLAS results suggest $\Gamma/M_X \approx 0.06$. However, the statistical preference over a narrow width is very small [3]; indeed, CMS suggests a narrow width. If a large width, $\Gamma/M_X \approx 0.06$, is assumed for CMS, then the significance decreases to 2.0σ . Note that this excess of di-photons at 750 GeV is not accompanied by any missing energy, leptons, i.e. $\ell^+\ell^-$, ZZ , or jets, nor that the $\sqrt{s} = 8$ TeV data showed any anomaly apart from a slight upward fluctuation at $m_{\gamma\gamma} \approx 750$ GeV.

Taking into account the Landau-Yang theorem [4, 5], this signal could be due to a boson with spin different from 1, decaying into 2 photons. Of course, the simplest option would be to assume the presence of a new scalar field X , a singlet under the SM gauge group with a mass $M_X \approx 750$ GeV. The presence of any new boson, which can couple to photons, would naturally beg plethora of questions relevant for beyond the Standard Model physics and cosmology. One of the important issues is the stability of the electroweak (EW) vacuum; within the SM the tussle between the top Yukawa and the Higgs self-coupling suggests a metastable vacuum for the Higgs, which has been confirmed at the 2.8σ level by computing two-loop corrections in Ref. [6, 7]. Indeed, the metastability of the EW vacuum can cause severe restrictions on the possible models of the early universe [8]. There are also other cosmological questions if one introduces a new scalar field: how would X couple to dark matter [9]? Or could X play the role of the inflaton, or the curvaton?

The aim of this paper is to understand how our picture of the early universe would respond to the introduction of a new SM gauge singlet scalar field with mass below 1 TeV, and decaying into photons at a high enough rate, if possible. As usual, scalar fields play a crucial role in inflation, acting as secondary fields whose perturbations can potentially seed large scale structures, or being

responsible for reheating the universe, or creating matter-anti-matter asymmetry (for a review see [10]).

In the simplest scenario, we wish to model the possible new spin-0 boson, X , as a real CP-odd scalar, which couples to an extra vector-like fermion Ψ , carrying color and electric charge, and the SM Higgs¹. Furthermore, we will demand that the X field does not develop a vacuum expectation value (VEV) at late times, otherwise it would break CP spontaneously. A natural question that we would like to answer is then the following: what would happen to the stability of the EW vacuum in the presence of another light scalar field X , which has couplings to the Higgs? Could the couplings be large enough, but still within perturbative limits, to explain the diphoton excess? The second paramount question which we pose is: could this light X be responsible for cosmic inflation like the SM Higgs with its non-minimal coupling to gravity?

Of course, there is the possibility that the diphoton excess will turn out to be a statistical fluctuation. Even in this case we believe our questions are interesting. One reason is that if the EW vacuum is stabilized by new physics that is not much above the TeV scale the Higgs mass is natural² [11], while if it is stabilized by much heavier physics a fine-tuning is required (modulo protection mechanisms such as supersymmetry, which so far have not been observed).

In Section 2 we introduce the model we will work with and discuss the typical values of the parameters that can account for the diphoton excess. We will study the modified renormalization group equations (RGEs) up to 1-loop due to X and Ψ couplings and the matching conditions due to the new masses in Section 3. In fact, the diphoton excess would require a large Yukawa coupling between X and Ψ , which contributes negatively to the self-coupling of X ; this coupling, however, is required to be positive to maintain the stability of the EW vacuum. As we shall see in Section 4, despite this fact, it is possible to stabilize the EW vacuum. In Section 4 we will also study the classical dynamics of gravity and the scalars, by considering the most general Lagrangian with operators of dimension up to 4, and investigate whether inflation can be realized. Both X and the SM Higgs can now couple to gravity, also via non-minimal interactions, which we will take into account in our analysis. We provide our conclusions in Section 5.

2 Minimal model

We consider the model with Lagrangian

$$\mathcal{L} = \sqrt{-g}(\mathcal{L}_{\text{SM}} + \Delta\mathcal{L} + \mathcal{L}_{\text{gravity}}), \quad (2.1)$$

The gauge group is the SM one: $G_{\text{SM}} \equiv \text{SU}(3)_c \times \text{SU}(2)_L \times \text{U}(1)_Y$. Also, $g_{\mu\nu}$ is the metric of the space-time and g its determinant. \mathcal{L}_{SM} is the SM Lagrangian (minimally coupled to gravity). $\Delta\mathcal{L}$ represents the beyond-the-SM terms in the Lagrangian due to the chosen model. We consider the following new physics.

- **A Dirac fermion.** (In Weyl notation) it is a pair of two-component fermions Ψ_1 and Ψ_2 in the following vector-like representation of G_{SM} : $\Psi_1 \sim (3, 1)_q$, $\Psi_2 \sim (\bar{3}, 1)_{-q}$. Namely, they

¹In this case the width is smaller than $\Gamma/M_X \approx 0.06$; we will ignore this possible issue here, given that such value has currently almost the same statistical significance as a narrow width.

²See also [12, 13] for an extension of these ideas to include gravity and inflation.

form a Dirac fermion in the fundamental representation of $SU(3)_c$, neutral under $SU(2)_L$ and with electric charge qe (which is left a priori as a free parameter).

- **A real pseudoscalar.** This scalar X , which may be identified with a possible resonance at ≈ 750 GeV, is CP-odd and neutral under G_{SM} .

Assuming renormalizability, the most general CP-even $\Delta\mathcal{L}$ is

$$\Delta\mathcal{L} = i \sum_{j=1}^2 \bar{\Psi}_j \not{D} \Psi_j - M_\psi (\Psi_1 \Psi_2 + \bar{\Psi}_1 \bar{\Psi}_2) + \frac{1}{2} (\partial X)^2 - \Delta V(H, X) - iy X (\Psi_1 \Psi_2 - \bar{\Psi}_1 \bar{\Psi}_2) + \dots \quad (2.2)$$

where H is the Higgs doublet,

$$\Delta V(H, X) \equiv \frac{m_X^2}{2} X^2 + \frac{\lambda_X}{4} X^4 + \frac{\lambda_{HX}}{2} (|H|^2 - v^2/2) X^2 \quad (2.3)$$

and the classical potential of the full model is

$$V(H, X) = \lambda_H (|H|^2 - v^2/2)^2 + \Delta V(H, X). \quad (2.4)$$

The dots in Eq. (2.2) represent extra Yukawa couplings and mixing terms between Ψ and the SM quarks, which are possible for special values of the charge of Ψ ; these couplings can make Ψ unstable. For example, for $q = 2/3$ one can write down the gauge invariant operator $H\Psi_2 Q$, where Q is the SM quark doublet. For simplicity we take these terms small enough that their contribution to the running of the relevant parameters is negligible³.

Finally, $\mathcal{L}_{\text{gravity}}$ includes the pure gravitational Lagrangian and the possible non-minimal couplings between gravity and the other fields, we consider

$$\mathcal{L}_{\text{gravity}} = -\frac{\bar{M}_{\text{Pl}}^2}{2} R - \Lambda - (\xi_H |H|^2 + \xi_X X^2/2) R + \alpha R^2, \quad (2.5)$$

where \bar{M}_{Pl} is the reduced Planck mass, Λ is a cosmological constant, α , ξ_H and ξ_X are real couplings and R is the Ricci scalar. We demand α to be positive, which ensures the stability of Minkowski's space-time. This $\mathcal{L}_{\text{gravity}}$ is the most general gravity Lagrangian with parameters of dimension of non-negative powers of energy modulo the Gauss-Bonnet and the Weyl squared terms. The Gauss-Bonnet term can be written as a linear combination of the R^2 and Weyl squared term plus a total derivative, which does not affect the following analysis. The Weyl squared term, being invariant under conformal transformations and vanishing on the flat space-time, does not change the classical treatment of inflation. We will assume that the same is true at quantum level.

Note that, even if we set $\alpha = 0$, $\xi_H = 0$ and/or $\xi_X = 0$, the corresponding operators (R^2 , $|H|^2 R$ and $X^2 R$) are generated by quantum corrections. Therefore, one should not regard the absence of such terms as a satisfactory option.

The requirement that $\Delta\mathcal{L}$ be CP-even simplifies considerably the model, avoiding terms that are odd under $X \rightarrow -X$ in the potential.⁴ This guarantees that the mixing between the Higgs and

³In Fig. 1 we will choose $q = 3/2$, which forbids the above-mentioned operator.

⁴CP is broken by small SM Yukawa couplings: the effect of this breaking on V is negligibly small for our analysis as suppressed by loop factors and the small value of such couplings.

X is small as suggested by data [3]. We introduce a colored fermion Ψ because it easily ensures (through its Yukawa coupling y) that the production of X is sizable in proton collisions at the LHC.

In addition to the SM parameters and q , this model has only 5 parameters if we ignore the couplings to gravity: M_X , M_ψ , λ_{HX} , λ_X and y . We take all of them to be real and positive.⁵

The EW symmetry breaking is triggered by the VEV $v \approx 246$ GeV of the neutral component of the Higgs doublet. We do not want to break CP spontaneously, therefore we take the VEV of X to be zero. Given that we assume $\lambda_{HX} > 0$ and of course

$$\lambda_H > 0, \quad \lambda_X > 0, \quad (2.6)$$

for these values of the scalar fields we reach⁶ the absolute minimum of V .

This model has been studied in [14] (see also [15] and [3]). The partial decay rates for $X \rightarrow gg$ and $X \rightarrow \gamma\gamma$ are

$$\Gamma_{gg} \equiv \Gamma(X \rightarrow gg) = M_X \frac{\alpha_3^2}{8\pi^3} \tau y^2 |\mathcal{P}(\tau)|^2, \quad \Gamma_{\gamma\gamma} \equiv \Gamma(X \rightarrow \gamma\gamma) = M_X \frac{9\alpha^2}{16\pi^3} q^4 \tau y^2 |\mathcal{P}(\tau)|^2 \quad (2.8)$$

where $\tau \equiv 4M_\psi^2/M_X^2$ and $\mathcal{P}(\tau) \equiv \arctan^2(1/\sqrt{\tau-1})$. The reported excess can be achieved for $\Gamma_{gg}/M_X \sim 10^{-3} - 10^{-6}$ and $\Gamma_{\gamma\gamma}/M_X \sim 10^{-6}$. Note that $y, q \sim 1$ and $M_\psi \sim \text{TeV}$ can account for the claimed excess.

3 RGE analysis and thresholds

Since we want to study the predictions of this model at energies much above the EW scale, up to the Planck scale, we need the complete set of RGEs. We use the modified minimal subtraction ($\overline{\text{MS}}$) renormalization scheme to define the renormalized couplings and their RGEs. Moreover, for a generic renormalized coupling g we write the RGEs as

$$\bar{\mu}^2 \frac{dg}{d\bar{\mu}^2} = \beta_g, \quad (3.1)$$

where $\bar{\mu}$ is the $\overline{\text{MS}}$ renormalization energy scale. The β -functions β_g can also be expanded in loops:

$$\beta_g = \frac{\beta_g^{(1)}}{(4\pi)^2} + \frac{\beta_g^{(2)}}{(4\pi)^4} + \dots, \quad (3.2)$$

where $\beta_g^{(n)}/(4\pi)^{2n}$ is the n -loop contribution.

⁵At least at the classical level; whether they remain positive after quantum corrections will be addressed later on.

⁶This is the case because

$$M_X^2 - \frac{v^2 \lambda_{HX}}{2} > 0, \quad (2.7)$$

for a large X squared mass, e.g. $M_X \approx 750$ GeV, and λ_{HX} in the perturbative regime.

Let us start from energies much above M_X and M_ψ . In this case we have

$$\begin{aligned}
\beta_{g_1^2}^{(1)} &= \frac{(24q^2 + 41)g_1^4}{10}, & \beta_{g_2^2}^{(1)} &= -\frac{19g_2^4}{6}, & \beta_{g_3^2}^{(1)} &= -\frac{19g_3^4}{3}, \\
\beta_{y_t^2}^{(1)} &= y_t^2 \left(\frac{9}{2}y_t^2 - 8g_3^2 - \frac{9g_2^2}{4} - \frac{17g_1^2}{20} \right), \\
\beta_{\lambda_H}^{(1)} &= \left(12\lambda_H + 6y_t^2 - \frac{9g_1^2}{10} - \frac{9g_2^2}{2} \right) \lambda_H - 3y_t^4 + \frac{9g_2^4}{16} + \frac{27g_1^4}{400} + \frac{9g_2^2g_1^2}{40} + \frac{\lambda_{HX}^2}{4}, \\
\beta_{\lambda_{HX}}^{(1)} &= \left(3y_t^2 - \frac{9g_1^2}{20} - \frac{9g_2^2}{4} + 6\lambda_H + 3\lambda_X + 6y^2 \right) \lambda_{HX} + 2\lambda_{HX}^2, \\
\beta_{\lambda_X}^{(1)} &= \lambda_{HX}^2 + 9\lambda_X^2 + 12y^2\lambda_X - 12y^4, \\
\beta_{y^2}^{(1)} &= y^2 \left(9y^2 - 8g_3^2 - \frac{18q^2g_1^2}{5} \right),
\end{aligned}$$

where g_3 , g_2 and $g_1 = \sqrt{5/3}g_Y$ are the gauge couplings of $SU(3)_c$, $SU(2)_L$ and $U(1)_Y$ respectively and y_t is the top Yukawa. The RGEs of the massive parameters can be neglected because in this step we are interested in energies several orders of magnitude above the EW scale.

In order to solve the RGEs above we need the values of the involved couplings at some reference energy, which we take to be the top mass M_t . As far as the SM couplings are concerned, we will use the central values given in Ref. [6]. Regarding the new couplings, let us first consider what happens at the mass threshold M_ψ . Following [16] we adopt the approximation in which the new Yukawa coupling run only above the corresponding mass thresholds; this is technically implemented by substituting $y \rightarrow y\theta(\bar{\mu} - M_\psi)$ on the right-hand side of the RGEs. An analogous substitution can be performed to take into account the scalar mass threshold M_X : we have to perform the substitutions $\lambda_X \rightarrow \lambda_X\theta(\bar{\mu} - M_X)$ and $\lambda_{HX} \rightarrow \lambda_{HX}\theta(\bar{\mu} - M_X)$ on the right-hand side of the RGEs. The situation is different from the case where the new heavy scalar acquires a VEV: in this case the Higgs quartic coupling gets a tree-level shift of order λ_{HX}/λ_X [17–19]. This is the result of integrating out the massive scalar degree of freedom at tree-level. The reason why this shift occurs is because setting the heavy scalar to zero is not consistent with the field equations. In the present case X does not develop a VEV and such shift does not occur.

4 Stability analysis, naturalness and inflation

Since we use the 1-loop RGEs, we approximate the Coleman-Weinberg [20] effective potential of this model with its RG-improved tree-level potential: we substitute the bare couplings in the classical potential with the corresponding running ones. The fact that the $\overline{\text{MS}}$ quantities are gauge invariant, as proved in [6, 21], guarantees that our results will not be affected by any gauge dependence.

The conditions to guarantee the stability of the EW vacuum at this level of approximation are given in (2.6), where λ_H and λ_X have to be thought of as the renormalized couplings.

The question of the stability of the EW vacuum has been previously addressed in other economic extensions of the SM. For example, singlet scalars were considered in [17, 22–24] and more elaborated possibilities were studied in [25].

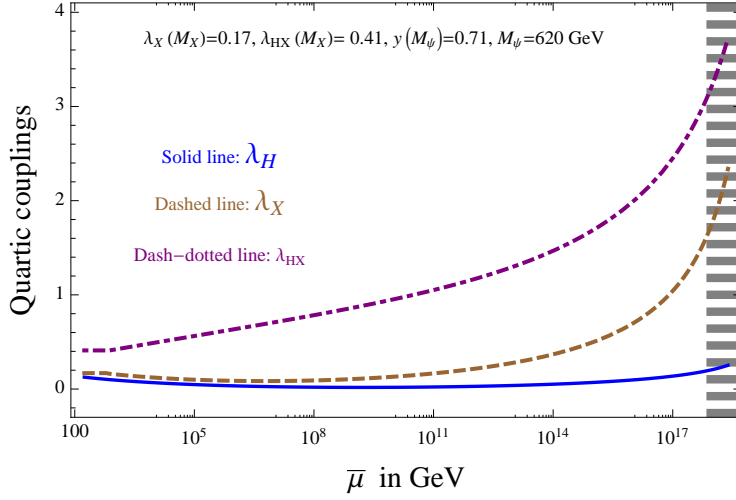


Figure 1: *RG evolution of the quartic couplings λ_H , λ_X and λ_{HX} for $M_X \approx 750$ GeV and $q = 3/2$. The stripes on the right indicate the region that is presumably dominated by Planck physics.*

The sizable value of y needed to reproduce the diphoton excess tends to make λ_X negative at high energies. However, one can compensate this effect via a sizable λ_{HX} , such that both λ_H and λ_X remain positive.

In Fig. 1 we show the evolution of the quartic couplings as a function of the renormalization scale. When the parameters are chosen appropriately (e.g. as in that figure), there are no Landau poles in the model, *all* couplings remain perturbative and the stability conditions are fulfilled (all couplings are positive) below the Planck scale. Note that the model recently proposed in [19], with $q = 2$, develops a Landau pole for g_1 five orders of magnitude below the Planck scale (and therefore is not perturbative) as it can be shown by solving the RGE of g_1 given above. The region with stripes on the right of Fig. 1 corresponds to the regime where Planck physics is expected to be dominant; the behavior of the curves there is thus presumably unreliable. The values of the partial decay widths corresponding to Fig. 1 are

$$\{\Gamma_{\gamma\gamma}/M_X, \Gamma_{gg}/M_X\} \approx \{1.3 \times 10^{-6}, 1.2 \times 10^{-5}\} \quad (4.1)$$

and are compatible with the reported signal. If one increases M_ψ up to the TeV one can still be (although barely) compatible with $\Gamma_{\gamma\gamma} \sim 10^{-6}M_X$, $\Gamma_{gg} \gtrsim 10^{-6}M_X$ and perturbativity and stability up to the Planck scale by choosing the remaining parameters appropriately:

$$\{\lambda_X(M_X), \lambda_{HX}(M_X), y(M_\psi), q\} \approx \{0.2, 0.4, 0.8, 3/2\}. \quad (4.2)$$

In Fig. 2 we provide a graphic representation of the phase diagram of the model. In the green region all couplings are perturbative and the quartic couplings fulfill the stability conditions, see Eq. (2.6). The non-perturbativity of the red region in Fig. 2 is triggered by Landau poles of the quartic couplings, not of the Yukawa y .

If such signal will persist, it could be interpreted as a manifestation of the physics needed to stabilize the EW vacuum. Nevertheless, if the signal will turn out to be a statistical fluctuation,

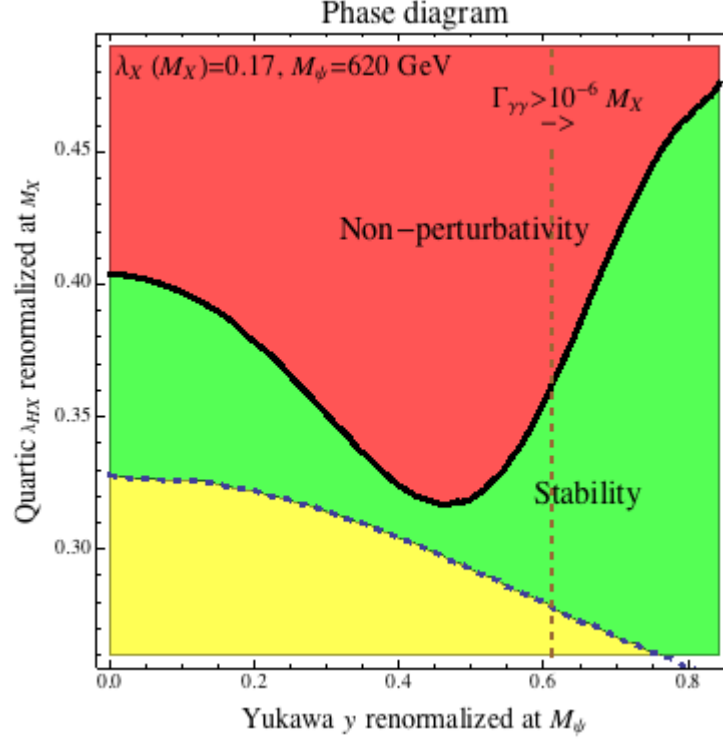


Figure 2: *Phase diagram of the model for $M_X \approx 750 \text{ GeV}$ and $q = 3/2$. We give the stability region (where all couplings are perturbative and the EW vacuum is stable) and the non-perturbativity region. The yellow region below the stability one does not satisfy the stability conditions. The part on the right of the vertical dashed line can account for the diphoton signal.*

we still regard our results interesting because they provide an example of new physics able to stabilize the EW vacuum without introducing unnaturally large contributions to the Higgs mass M_h . In fact in the present model the leading radiative correction to M_h^2 is

$$\delta M_h^2 \approx \frac{C_X \lambda_{HX}}{(4\pi)^2} M_X^2, \quad (4.3)$$

where C_X is a quantity of order one. In the present case (see e.g. Fig. 1) such correction is not exceeding the order of magnitude of M_h . Therefore a natural EW scale occurs: this is because all the new particles that have sizable couplings to the Higgs are not too heavy [11, 12, 26].

We now turn to inflation. Note that $\mathcal{L}_{\text{gravity}}$ is non-standard because there are non-minimal couplings between the scalar fields and there is an R^2 term. The latter term can be transformed into an additional scalar z and the non-minimal couplings can be transformed into a non-standard scalar kinetic and potential terms by going to the so-called Einstein frame: the final scalar-tensor Lagrangian is [13]

$$\mathcal{L}_{\text{st}} = \sqrt{-g_E} \left(-\frac{\bar{M}_{\text{Pl}}^2}{2} R_E + \mathcal{L}_{\text{kin}} - U \right), \quad (4.4)$$

where the scalar kinetic and potential terms are

$$\mathcal{L}_{\text{kin}} = \frac{3\bar{M}_{\text{Pl}}^2}{z^2} [(\partial h)^2 + (\partial X)^2 + (\partial z)^2] \equiv \frac{K_{ij}(\phi)}{2} \partial_\mu \phi^i \partial^\mu \phi^j \quad (4.5)$$

$$U(H, X, z) = \frac{36\bar{M}_{\text{Pl}}^4}{z^4} \left[V(H, X) + \frac{1}{16\alpha} \left(\frac{z^2}{6} - \xi_H 2|H|^2 - \xi_X X^2 - \bar{M}_{\text{Pl}}^2 \right)^2 \right] \quad (4.6)$$

and everything is computed with the metric

$$g_{\mu\nu}^E \equiv g_{\mu\nu} \times z^2 / (6\bar{M}_{\text{Pl}}^2). \quad (4.7)$$

Here we have introduced $\{\phi^1, \phi^2, \phi^3\} \equiv \{h, X, z\}$, K_{ij} is the field metric and h is the real scalar field associated with the physical Higgs. The minimum of U occurs for $z \approx \sqrt{6}\bar{M}_{\text{Pl}}$, $X = 0$ and $H \approx 0$ (having neglected v at the high inflationary scales).

The slow-roll conditions can then be written in a compact form:

$$\epsilon \equiv \frac{\bar{M}_{\text{Pl}}^2 U_{,i} U^{,i}}{2U^2} \ll 1, \quad (4.8)$$

$$\left| \frac{\eta^i_j U^{,j}}{U^{,i}} \right| \ll 1 \quad (i \text{ not summed}), \quad \text{where} \quad \eta^i_j \equiv \frac{\bar{M}_{\text{Pl}}^2 U^{,i}_{;j}}{U} \quad (4.9)$$

and we raise the indices with the inverse field metric K^{ij} ; also, $U_{,i} \equiv \partial U / \partial \phi^i$ and $U_{;i}$ is the covariant derivative built with K_{ij} . When these conditions are fulfilled the field equations are

$$\dot{\phi}^i \approx -\frac{U^{,i}}{3H}, \quad H^2 \approx \frac{U}{3\bar{M}_{\text{Pl}}^2}, \quad (4.10)$$

where a dot denotes the derivative w.r.t. the cosmic time t ; also $H \equiv \dot{a}/a$ and a is the Robertson-Walker scale factor.

The sizable quartic couplings in Fig. 1 at the inflationary scales suggest that inflation is driven by z rather than the Higgs or X . Indeed, sizable quartic couplings correspond to field directions with steep potential. In Fig. 3 we show that this is indeed the case: even if we start from initial values of h and X larger than z the dynamics converge to the line $h = X = 0$ and only there the slow-roll conditions in (4.8) and (4.9) are satisfied. The inflationary predictions are therefore those of Starobinsky's model [27], in agreement with current observations.

Even if one set initially $\alpha = 0$, Higgs inflation [28] or X -inflation would generically require large values of ξ , given the present bound on the tensor to scalar ratio $r \lesssim 0.1$ [29] and such large non-minimal couplings would source a large value of α modulo fine-tuning [30]. The same conclusion is reached even if one does not introduce X and the only scalar fields are h and z .

5 Conclusions

In this paper, we considered an extra CP-odd scalar X and an extra vector-like fermion Ψ carrying color and electric charge: this is the simplest model that can explain the excess in the diphoton channel reported by ATLAS and CMS at 750 GeV.

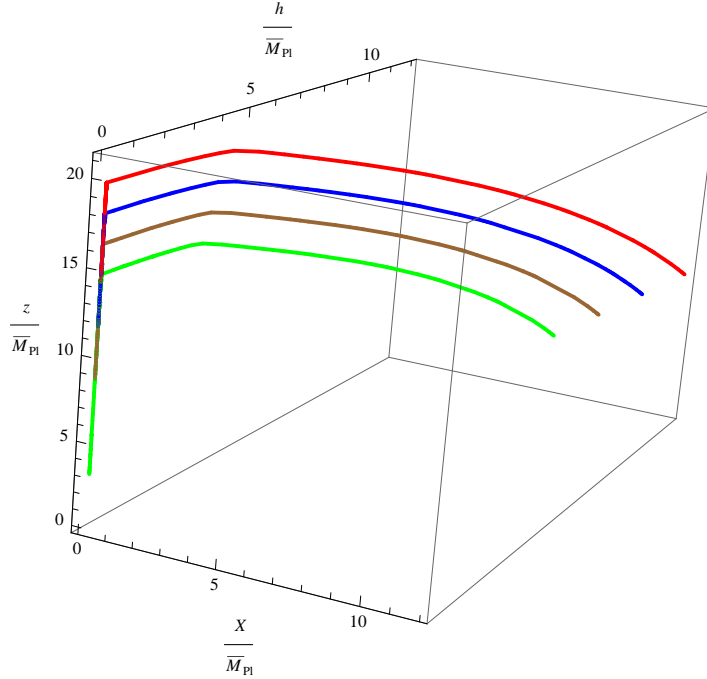


Figure 3: *Path in scalar field space obtained by solving the field equations in (4.10). Here we set $\xi_H = \xi_X = -1/6$, the value of α is fixed by the observed amplitude of scalar fluctuations (for a number of e -folds of about 60) and the remaining parameters are chosen like in Fig. 1.*

Our computations, valid up to one-loop in the RGEs, suggest that it is indeed possible to explain the excess in the diphoton channel, curing at the same time the SM instability of the EW vacuum and keeping *all* couplings perturbative up to the Planck scale. An example is provided in Fig. 1, which uses a parameter choice leading to the partial decay rates of X in Eq. (4.1). The phase diagram with the regions of stability (which corresponds to perturbative couplings and stable EW vacuum), non-perturbativity and unstable EW vacuum are given in Fig. 2. There, the part that can account for the diphoton signal is also provided.

Besides the excess in the diphoton channel, we believe that these results are interesting for another reason, independently on whether or not the ATLAS and CMS excess at 750 GeV will be confirmed by future analysis: they provide a SM extension that can stabilize the EW vacuum without invoking an unnaturally large contribution to the Higgs mass. Indeed, in the present model M_h is natural as the largest correction to it, given by Eq. (4.3), is never $\gg M_h$.

Moreover, we studied the classical dynamics of the model during inflation by considering the most general Lagrangian with operators of dimension up to 4 and including the new scalar X . Although, both X and/or the SM Higgs can potentially inflate the early universe and possibly explain the current CMB data, the sizable values of the quartic couplings needed to stabilize the EW vacuum forces inflation to be driven by the effective scalar z , which emerges from the R^2 operator: we found that during inflation, the classical slow-roll trajectory is dominated by z and both X and h roll down quickly within few Hubble times due to the sizable quartic couplings. Moreover, another reason favors z -inflation (also known as Starobinsky inflation [27]). Large

non-minimal couplings ξ_H and ξ_X to gravity are generically required in order to satisfy the right amplitude of scalar perturbations in the CMB and the current constraint on the tensor-to-scalar ratio; indeed, as analyzed before in Ref. [30], a large non-minimal coupling to gravity would yield higher order curvature corrections (e.g. R^2 terms) modulo fine-tuning.

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Note added. While we were finalizing this paper another work [31] that addresses the EW vacuum instability appeared in the arXiv. Moreover, other articles, which are related to the present analysis, appeared [32]. After our article was posted on the arXiv, Ref. [19] was updated admitting that the $q = 2$ case has a Landau pole much below the Planck scale.

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