750 GeV diphoton resonance and inflation

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An $\mathcal{O}(750)$ GeV Resonance and Inflation

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We study the possibility of a heavy scalar or pseudoscalar in TeV-scale beyond the Standard Model scenarios being the inflaton of the early universe in light of the recent $\mathcal{O}(750)$ GeV diphoton excess at the LHC. We consider a scenario in which the new scalar/pseudoscalar couples to the Standard Model gauge bosons at the loop-level through new massive Standard Model charged vector-like fermions with/without dark fermions. We calculate the renormalization group running of both the Standard Model and the new scalar couplings, and present two different models that are perturbative and with a stabilized vacuum up to near the Planck scale. Thus, the Standard Model Higgs and this possible new resonance may still preserve the minimalist features of Higgs inflation.

I. INTRODUCTION

One of the recurrent lessons in physics is that simplicity beats complexity. While theoretical reasonings (such as the hierarchy problem and naturalness) and observations (e.g., dark matter) have motivated many beyond the Standard Model (BSM) scenarios, the search for new physics has so far been elusive. One may wonder if we have seen the beginning of a desert in that the new physics that completes the Standard Model (SM) only appears at higher energies currently not within reach.

Very recently the CMS and ATLAS collaboration have reported a roughly 3σ enhancement in the $\gamma\gamma$ spectrum at $\mathcal{O}(750)$ GeV, based on the first $\sim 3/\text{fb}$ of data from Run 2 of the LHC at $\sqrt{s} = 13$ TeV [1, 2]. This potential new resonance can be interpreted as either a spin zero or a spin two particle, as the Laudau-Yang theorem states that an on-shell spin-one particle cannot decay to diphoton. A spin-two particle would yield an interesting $\gamma\gamma$ distribution that is peaked in the beam directions. The possibility that the candidate resonance is a KK-graviton is not excluded but is quite constrained by the Run 1 data. For concreteness, we focus on the interpretation of this candidate resonance as a spin zero particle, which can be a scalar or a pseudoscalar. The 8 TeV run of the LHC did not report a signal, which hints at the possibility that the production rate might have a steeper energy dependence. All these hints, taken together, suggest that this new resonance couples to the SM gauge bosons including photons and gluons only through higher dimensional operators. The resonance can be produced by $gg \to \Phi \to \gamma\gamma$ though $q\bar{q}$ annihilation can also contribute. Moreover, the ATLAS collaboration has enough events to show a larger width of the resonance, which suggests the presence of some other decay channels besides diphoton. Since there are no excesses shown in the other channels, one tempting explanation is that this new resonance decays into the dark sector.

This possible excess has already sparked numerous theoretical explorations [3–7]. To explain this potential new resonance, two prevailing ideas are: (1) to stay within weakly coupled theories and introduce vector-like SM charged fermions, or (2) to invoke strong dynamics as in composite Higgs models. While we maintain our healthy skepticism towards the reality of this signal, we think it worthwhile to explore the cosmological implications of this potential new resonance while awaiting further data from the LHC to either confirm or falsify its existence. Though previous works [3–7] provide various explanations for the possible excess, the $X(750)$ resonance and the additional new particles are introduced without necessarily a purpose. Here, we would like to see, if we take Occam's razor seriously, what correlated statements in cosmology can we make.

Inflation is the leading paradigm of early universe cosmology. An inflationary universe can be realized with a scalar/pseudo scalar (known as the inflaton) with a sufficiently flat potential. Moreover, this inflaton should couple to the SM fields somehow in order for reheating to occur. It is in this minimalist spirt that Higgs inflation [8], in which the SM Higgs field was identified with the inflation (with non-minimal coupling to gravity), was proposed. Here, we examine how the new resonance, if confirmed, may alter this minimalist scenario. We focus on the weak coupling explanation of the resonance, as it is challenging (if at all possible) to accommodate inflation concretely with strongly coupled theories (see e.g., [9–11]). Moreover, a strongly coupled sector typically comes with a plethora of particles, far from the minimalist approach we opt to take. In the weakly coupled scenarios, the new resonance can be a candidate for the inflaton, while the additional vector-like SM charged fermions are necessarily there for the inflaton to reheat the SM sector. In a sense, our scenario is a minimal model to accommodate the success of both the SM and cosmic inflation, should the $X(750)$ resonance is confirmed. We studied the perturbativity and the stability of the SM vacuum in light of the new particles and the associated new couplings.

The paper is organized as follows. We discuss some general aspects of this new resonance in Section II. We present some models with/without dark fermions and solve the corresponding renormalization group (RG) equations in Section III. The parameter spaces of the initial conditions that preserve perturbativity and vacuum stability of the models are presented. We conclude in Section IV. For completeness, we also include in an appendix the complete RG equations for the models explaining the $X(750)$ resonance considered in [3].

II.   GENERAL ASPECTS OF THE NEW RESONANCE

The CMS and ATLAS excess suggests a cross section that:

$$\sigma(pp \to \Phi) \times BR(\Phi \to \gamma\gamma) \gtrsim 2fb \tag{1}$$

with the mass $M_\Phi \approx 750$ GeV of a new spin 0 particle Φ. As discussed in the introduction, we consider the five-dimensional operators with a SM singlet scalar/pseudoscalar that couples to the SM gauge bosons. For the scalar case, $\Phi = S$, the effective couplings are given by

$$L_{\text{effective}} \supset \frac{S}{4}(-g_{s\gamma}F_{\mu\nu}F^{\mu\nu} - g_{s\gamma}G_{\mu\nu}G^{\mu\nu}), \tag{2}$$

whereas for the pseudoscalar case, $\Phi = P$, we have

$$L_{\text{effective}} \supset \frac{P}{4}(-g_{p\gamma}F_{\mu\nu}\tilde{F}^{\mu\nu} - g_{p\gamma}G_{\mu\nu}\tilde{G}^{\mu\nu}). \tag{3}$$

Here $G_{\mu\nu}$ and $F_{\mu\nu}$ are the Standard Model color $SU(3)_C$ and $U(1)_{em}$ field strength. The gluon couplings account for the gluon-gluon fusion production of this new resonance, and the photon couplings give rise to the dominant decay channels to diphoton. Notice here there are also likely decay modes of S/P into dibosons (γZ, ZZ, WW) because the dimensional 5 operators with these gauge fields are also allowed. Currently an excess does not show up in the other channels, therefore they are constrained by the Run 1 and 2 data [29].

It is well-known that the new vector-like fermion at the TeV scale can give rise to these effective operators. In fact, such fermions also appear in BSM extensions that address other pressing questions like flavor physics or dark matter issues [15, 16]. Let us introduce here N_ψ vector-like fermions $\psi = (\psi_L, \psi_R)^T$ in the 3 representation of the color $SU(3)_C$ and with an electromagnetic charge Q. Their Yukawa couplings to S and the mass term are

$$L \supset -\lambda_{S\psi\psi}S\bar{\psi}\psi - M_\psi\bar{\psi}\psi \,. \tag{4}$$

Here and in what follows we focus on the scalar case, $\Phi = S$, for concreteness, but the results for the pseudoscalar case should be qualitatively the same. We also assume suppressed Yukawa couplings between vector-like quarks and the SM fermions. By integrating out the heavy fermion we arrive at the effective couplings:

$$g_{sg} = \frac{N_\psi\lambda_{S\psi\psi}\alpha_s}{3\pi\sqrt{2}M_\psi} \,, \tag{5}$$

$$g_{s\gamma} = \frac{\sqrt{2}N_\psi\lambda_{S\psi\psi}Q^2\alpha}{\pi M_\psi} \,. \tag{6}$$

The LHC measurement production times branching ratio is $\sigma(pp \to \Phi) \times BR(\Phi \to \gamma\gamma) \propto g_{sg}^2 g_{s\gamma}^2/(8g_{sg}^2 + g_{s\gamma}^2)$, with $g_{sg}/g_{s\gamma} \gg 1$. Taking $\alpha_s(\text{TeV}) \sim 0.09$ from PDG [17], we have

$$\left(\frac{2N_\psi\lambda_{S\psi\psi}Q^2\alpha}{\pi}\frac{1\text{TeV}}{M_\psi}\right)^2 \gtrsim \frac{2}{13000} \,. \tag{7}$$

III. SPECIFIC MODELS OF THE NEW RESONANCE FOR PERTURBATIVITY AND VACUUM STABLIZATION

In this section we explore models which are perturbative and stable up to near the Planck scale, 10^{17}GeV, while at the same time explaining the diphoton excess. As we have discussed, the diphoton excess motivates the existence of a new spin 0 particle and new vector-like massive fermions with SM charges. The couplings of this sector at the electroweak scale is constrained by Eq. (7). In addition, we here would like to include dark fermions which have no SM charges,

suggested by a possibility that the dark matter sector couples to the new spin 0 particle. Our matter content then includes the SM particles, the new scalar S, the new SM charged vector-like fermions, and the dark fermions.

Let us summarize our notation for the couplings. First, the scalar potential is given by [30]

$$V = -\mu^2|H|^2 + \lambda(|H|^2)^2 + m^2 S^2 + \frac{\lambda_S}{4!}S^4 + \frac{\kappa}{2}|H|^2 S^2 \,. \tag{8}$$

We introduce N_ψ vector-like fermions ψ and N_D dark fermions D [31]. We assume that the vector-like fermions are in the 3 representation of $SU(3)_C$, and have an $SU(2)_L$ isospin S_ψ and a hypercharge Y. On the other hand, the dark fermions have no SM charges. Their Yukawa couplings to S are denoted by $\lambda_{S\psi\psi}$ and λ_{SDD}, respectively. More explicitly, $L \supset -\lambda_{S\psi\psi}S\bar{\psi}\psi - \lambda_{SDD}S\bar{D}D$. Their RG equations are then given by

$$16\pi^2 \frac{d\lambda_S}{dt} = 3\lambda_S^2 + 12\kappa^2 + 24N_\psi(2S_\psi + 1)\lambda_{S\psi\psi}^2\lambda_\psi - 144(2S_\psi + 1)N_\psi\lambda_{S\psi\psi}^4 + 8N_D\lambda_{SDD}^2\lambda_S - 48N_D\lambda_{SDD}^4 \,,$$

$$16\pi^2 \frac{d\lambda_{S\psi\psi}}{dt} = (3 + 6N_\psi(2S_\psi + 1))\lambda_{S\psi\psi}^3 - 8g_3^2\lambda_{S\psi\psi} - 6Y^2 g_Y^2\lambda_{S\psi\psi} - 6S_\psi(S_\psi + 1)g_2^2\lambda_{S\psi\psi} + 2N_D\lambda_{SDD}^2\lambda_{S\psi\psi} \,,$$

$$16\pi^2 \frac{d\lambda_{SDD}}{dt} = (3 + 2N_D)\lambda_{S\psi\psi}^3 + 6N_T(2S_\psi + 1)\lambda_{SDD}\lambda_{S\psi\psi}^2 \,,$$

$$16\pi^2 \frac{dg_3}{dt} = (\text{SM part}) + \frac{2}{3}N_\psi(2S_\psi + 1)g_3^3 \,,$$

$$16\pi^2 \frac{dg_2}{dt} = (\text{SM part}) + 4N_\psi\frac{S_\psi(S_\psi + 1)(2S_\psi + 1)}{3}g_2^3 \,,$$

$$16\pi^2 \frac{dg_Y}{dt} = (\text{SM part}) + 4Y^2 N_\psi(2S_\psi + 1)g_Y^3 \,,$$

$$16\pi^2 \frac{dy_t}{dt} = (\text{SM part}) \,,$$

$$16\pi^2 \frac{d\lambda}{dt} = (\text{SM part}) + \frac{1}{2}\kappa^2 \,,$$

$$16\pi^2 \frac{d\kappa}{dt} = \kappa\left(4\kappa + 12\lambda + \lambda_S + 6y_t^2 + 12N_\psi(2S_\psi + 1)\lambda_{S\psi\psi}^2 + 4N_D\lambda_{SDD}^2 - \frac{3}{2}g_Y^2 - \frac{9}{2}g_2^2\right) \,, \tag{9}$$

where g_3, g_2, and g_Y are the gauge couplings for $SU(3)_C$, $SU(2)_L$, and $U(1)_Y$, respectively. $t = \ln\mu$ is the renormalization scale, y_t is the top Yukawa coupling and λ is the Higgs self-coupling. The calculation here is based on [19]. The SM part and initial conditions can be found in, e.g. [20]. In the rest of this section, based on the above RG equations, we investigate two types of models with (a) top-like vector-like fermions and dark fermions (b) vector-like fermions with exotic SM charges.

(a) Top-like vector-like fermions and dark fermions:

Let us start with a model with N_T vector-like fermions T with $(3,1,2/3)$ under $SU(3)_C \times SU(2)_L \times U(1)_Y$ and N_D dark Dirac fermions D. This class of models is given by $\psi = T$, $S_T = 0$, and $Y = 2/3$ in the previous RG equations. Note that these dark fermions can be the dark matter of the universe. They can annihilate mainly through $DD \to S/P \to gg$ in the early universe. Because in our case the dark fermion masses are not much constrained from perturbativity and vacuum stablization (as long as they are not far from the $\mathcal{O}(\text{TeV})$ range), we can easily achieve the correct relic abundance.
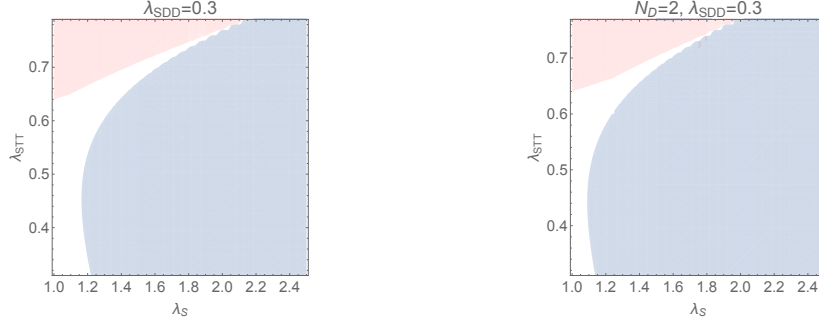
FIG. 1. The blue and red regions are not perturbative and stable up to 10^{17}GeV, respectively. In the left(right) panel, we take $M_T = 1000$GeV, $\lambda_{SDD} = 0.3, N_D = 1(2)$ and $N_T = 1$.
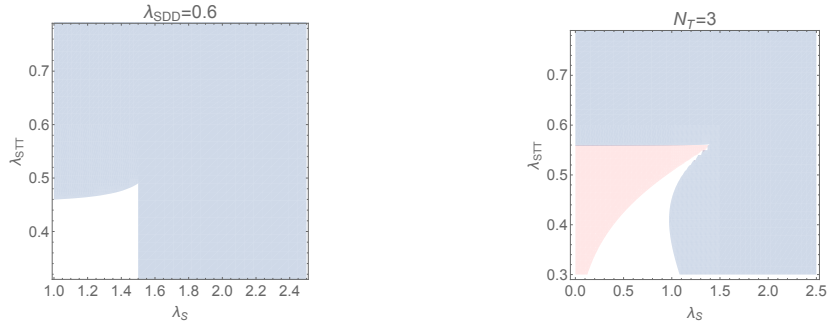

FIG. 2. Left: Same as Fig. 1, but taking $\lambda_{SDD} = 0.6$. Right: We take $N_T = 3$ and $\lambda_{SDD} = 0$.

In Fig. 1 and Fig. 2, we show the region where this model becomes perturbative and stable up to 10^{17}GeV for $N_T = 1$. The blue and red regions correspond to non-perturbative and unstable regions, respectively. Unfortunately, we cannot find a region realizing Eq. (7), which requires $\lambda_{S\psi\psi} \gtrsim \mathcal{O}(1)$. Increasing λ_{SDD} and N_D make the problem worse as in the right panel of Fig. 1 and the left panel of Fig. 2. The dark sector can therefore be an obstruction to realizing a "desert" up to high scale generically. On the other hand, if we increase the number of N_T, the situation becomes better, because the constraint (7) on the Yukawa coupling λ_{STT} at the electroweak scale is relaxed. However, it is not sufficient to explain the diphoton signal, see the right panel of Fig. 2.

(b) Exotic charge fermion:

The difficulty of the previous model comes from the necessity of large coupling for one top-like vector-like fermion. From Eq. (7), we see that we can take smaller λ_S by considering vector-like fermion with a larger charge. However, if the representation is too large, the gauge coupling hits the Landau pole much below the Planck scale. By solving the one loop renormalization group equation of the gauge couplings, the position of the pole is

$$\Lambda_{g_i} = M_\psi \exp\left(\frac{8\pi^2}{b_i}\frac{1}{g_i^2(M_\psi)}\right)$$
$$= M_\psi \left(\frac{M_t}{M_\psi}\right)^{b_{i,\mathrm{SM}}/b_i} \exp\left(\frac{8\pi^2}{b_i}\frac{1}{g_i^2(M_t)}\right), \tag{10}$$

$(SU(3)_C, SU(2)_L, U(1)_Y)$	$Q_{\max}$	Λ_{g_Y} [GeV]	Λ_{g_2} [GeV]	$\sigma \times BR > 2\text{fb}$	$\Lambda_{\lambda_{S\psi\psi}} < 10^{17}\text{GeV}$
$(3, 5, 0)$	2	8.9×10^{40}	1.3×10^{5}	$\lambda_{S\psi\psi} > 0.18 \left(\frac{M_\psi}{700\text{GeV}}\right)$	–
$(3, 4, 1/2)$	2	7.6×10^{26}	6.1×10^{7}	$\lambda_{S\psi\psi} > 0.32 \left(\frac{M_\psi}{700\text{GeV}}\right)$	–
$(3, 3, 1)$	2	4.7×10^{16}	1.1×10^{20}	$\lambda_{S\psi\psi} > 0.35 \left(\frac{M_\psi}{700\text{GeV}}\right)$	–
$(3, 2, 3/2)$	2	2.1×10^{13}	–	$\lambda_{S\psi\psi} > 0.35 \left(\frac{M_\psi}{700\text{GeV}}\right)$	–
$(3, 1, 2)$	2	1.8×10^{14}	–	$\lambda_{S\psi\psi} > 0.44 \left(\frac{M_\psi}{700\text{GeV}}\right)$	–
$(3, 3, 2/3)$	5/3	1.8×10^{24}	1.1×10^{20}	$\lambda_{S\psi\psi} > 0.53 \left(\frac{M_\psi}{700\text{GeV}}\right)$	$\lambda_{S\psi\psi} < 0.75$
$(3, 2, 7/6)$	5/3	3.4×10^{17}	–	$\lambda_{S\psi\psi} > 0.54 \left(\frac{M_\psi}{700\text{GeV}}\right)$	$\lambda_{S\psi\psi} < 0.81$
$(3, 1, 5/3)$	5/3	2.3×10^{17}	–	$\lambda_{S\psi\psi} > 0.62 \left(\frac{M_\psi}{700\text{GeV}}\right)$	$\lambda_{S\psi\psi} < 1.02$

TABLE I. The possible $SU(2)_L \times U(1)_Y$ charge. $Q_{\max}$ represents the maximum electromagnetic charge among the multiplet. Only the particle with $Q_{\max} \leq 5/3$ can be perturbative up to 10^{17}GeV. Here we take $M_\psi = 700\text{GeV}$.
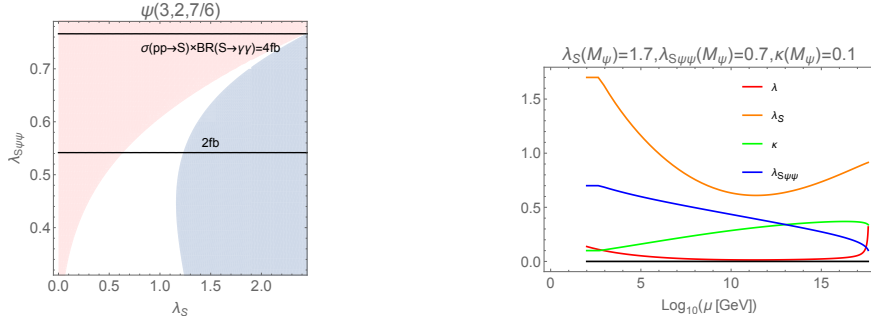

FIG. 3. Left: The blue and red regions are not perturbative and stable up to 10^{17}GeV, respectively. The white region is allowed. $\kappa = 0$ and $M_\psi = 700\text{GeV}$ is taken. Right: The running of scalar and Yukawa couplings. All couplings are perturbative and stable up to 10^{17}GeV.

where $i = Y, 2$, and $b_{i,\text{SM}}$ and b_i are the coefficients of the beta function of g_i in the SM and SM with exotic quark, respectively. We summarize the position of the pole in Table I, from which we can see that only $Q \leq 5/3$ is consistent with perturbativity up to 10^{17}GeV. The position of the Landau pole of $\lambda_{S\psi\psi}$, $\Lambda_{\lambda_{S\psi\psi}}$, is also shown. We note that $\Lambda_{\lambda_{S\psi\psi}}$ is independent of κ and λ_S. Therefore, as simple successful examples, we consider vector-like fermions whose charges are (3,2,7/6) and (3,1,5/3) under $SU(3)_C \times SU(2)_L \times U(1)_Y$, and show that these models indeed explain the diphoton excess and can be valid up to near the Planck scale. The mass bound on a charge 5/3 particle is about 900 GeV [21], if ψ mainly decays into top, while the bound is 690GeV if the main decay mode is light quark [22].

We plot the region excluded by perturbativity and stability up to 10^{17}GeV, in Fig. 3 with $M_\psi = 700\text{GeV}$, $S_\psi = 2$ and $Y = 7/6$. The inclusion of κ does not change the allowed region very much. It is found that only very small part in the white region has $\sigma(pp \to S) \times BR(S \to \gamma\gamma) \sim 2\text{fb}$. In Fig. 4, we also show the
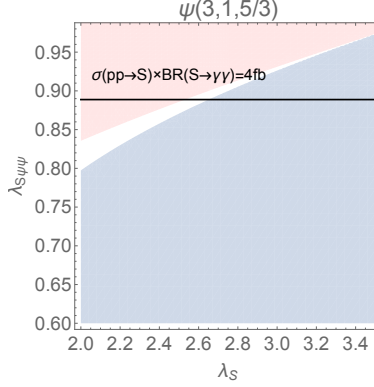
FIG. 4. Same as Fig. 3, but $S_\psi = 0$, $Y = 5/3$.

viable region with $S_\psi = 0, Y = 5/3$. Although the region is small, the consistent region survives.

IV. CONCLUSION

Motivated by the recent $\mathcal{O}(750)$ GeV diphoton excess at LHC, we have studied the RG running in the scenario where a new spin 0 particle couples to the SM gauge bosons at the loop-level through new massive SM charged vector-like fermions. For both models with and without dark fermions, we have explicitly shown that there exists a certain parameter range which accommodates the stability and perturbativity up to 10^{17} GeV. We illustrated two types of workable models with vector-like fermions with exotic SM charges which may explain the current diphoton excess at the same time. Our result thus suggests a possibility that, other than the X(750) resonance and the associated vector-like fermion in the $\mathcal{O}(\text{TeV})$ range, there is a "desert" up to near the Planck scale.

In the same spirit of Higgs inflation, it is natural to look for the inflaton among the two (pseudo)scalars, i.e., the SM Higgs and the new spin 0 particle, of the minimal setups we proposed. In the standard Higgs inflation, the action for the inflaton φ in the Jordan frame is given by

$$S_J \simeq \int d^4 x \left[-\frac{M_{\text{Pl}}^2}{2} R - \xi \varphi^2 R + \frac{1}{2} (\partial \varphi)^2 - \frac{\lambda_{\text{inf}}}{4} \varphi^4 \right], \tag{11}$$

which is characterized by the inflaton quartic coupling λ_{inf} and the nonminimal coupling ξ to the scalar curvature. In the region $|\xi| \gg 1$, however, the action in the Einstein frame is reduced to [32]

$$S_E \simeq \int d^4 x \left[-\frac{M_{\text{Pl}}^2}{2} R + \frac{1}{2} (\partial \chi)^2 - \frac{M_{\text{Pl}}^2}{4} \frac{\lambda_{\text{inf}}}{\xi^2} \left(1 + \exp \left(-\frac{2\chi}{\sqrt{6} M_{\text{Pl}}} \right) \right) \right], \tag{12}$$

which depends only on the ratio $\lambda_{\text{inf}}/\xi^2$. It can directly be related to the length of inflationary period, so that we have, e.g., $\xi = 4.5 \times 10^4 \sqrt{\lambda_{\text{inf}}}$ for 60 e-folds. Since both the SM Higgs and the new spin 0 particle may have the quartic coupling of the order $\lesssim \mathcal{O}(1)$, either can be used as the inflaton φ by introducing an appropriate size of the nonminimal coupling ξ. The primordial tilt n_s and the tensor-to-scalar ratio r are the same as the standard Higgs inflation, $n_s \simeq 0.97$ and $r \simeq 0.0033$ for the e-folding number $N = 60$, so they

agree well with the current constraints by Planck [26] and BICEP2/Keck Array [27]. This is a conclusion of this paper.

One important question to be investigated next is whether the inflaton is the SM Higgs or the new spin 0 particle. The main differences between those particles are their couplings to the SM particles and the dark fermions (if exist). Such a difference may leave some interesting imprints, e.g., on the reheating process and the dark matter abundance. Moreover, both scalar fields may play the role of the inflaton and thus our setups naturally suggest a realization of multifield inflation. It would be interesting to explore this possibility. The isocurvature remnants would be a key feature of this class of models and would be useful to constrain details of the model such as the reheating process and the initial condition of the inflationary trajectory [28][33]. Another interesting direction would be to embed inflation into other scenarios explaining the current diphoton excess, such as models in which the new spin 0 particle is identified with th dilaton or an axion. To explain the diphoton excess, the decay constant of the spin 0 particle is of the order of the electroweak scale. At least naively, the dilaton would go into the strong coupling regime at the inflation scale. An electroweak scale decay constant is also not suited for the simplest form of axion inflation. It seems therefore not easy to realize inflationary models such as dilaton inflation or axion inflation, and explain the diphoton excess simultaneously. However, it would be interesting to explore this direction further. We hope to address these issues elsewhere.

Note: While this work was completed, other papers [5–7] on electroweak vacuum stablization appeared on the arXiv. Our analysis concerns more general set-ups and the implications to inflation. We would like to also point out that the references [5, 6] did not take into account the Landau pole of gauge couplings g_Y and g_2, which is appropriately taken into account in our analysis. See TABLE I. The model in [7] contains vector-like fermions with the electromagnetic charge 3/2, which may form rather exotic mesons with a fractional electromagnetic charge, whereas the electromagnetic charge of mesons in our models are always integers.

[1] CMS Collaboration , CMS-PAS-EXO-15-004.

[2] The ATLAS collaboration, ATLAS-CONF-2015-081.

[3] J. Ellis, S. A. R. Ellis, J. Quevillon, V. Sanz and T. You, arXiv:1512.05327 [hep-ph].

[4] K. Harigaya and Y. Nomura, arXiv:1512.04850; Y. Mambrini, G. Arcadi and A. Djouadi, arXiv:1512.04913; S. Fichet, G. von Gersdorff and C. Royon, arXiv:1512.05751 [hep-ph]. A. Angelescu, A. Djouadi and G. Moreau, arXiv:1512.04921; Y. Nakai, R. Sato and K. Tobioka, arXiv:1512.04924 [hep-ph]. S. D. McDermott, P. Meade and H. Ramani, arXiv:1512.05326; R. S. Gupta, S. Jäger, Y. Kats, G. Perez and E. Stamou, arXiv:1512.05332; C. Petersson and R. Torre, arXiv:1512.05333; E. Molinaro, F. Sannino and N. Vignaroli, arXiv:1512.05334; R. Franceschini *et al.*, arXiv:1512.04933 [hep-ph]. G. M. Pelaggi, A. Strumia and E. Vigiani, arXiv:1512.07225 [hep-ph]. M. Son and A. Urbano, arXiv:1512.08307 [hep-ph]. S. Kanemura, K. Nishiwaki, H. Okada, Y. Orikasa, S. C. Park and R. Watanabe, arXiv:1512.09048 [hep-ph]. X. J. Bi *et al.*, arXiv:1512.08497 [hep-ph]. B. Dutta, Y. Gao, T. Ghosh, I. Gogoladze and T. Li, arXiv:1512.05439; F. Goertz, J. F. Kamenik, A. Katz and M. Nardecchia, arXiv:1512.08500 [hep-ph]. S. Di Chiara, L. Marzola and M. Raidal, arXiv:1512.04939 [hep-ph]. S. Knapen, T. Melia, M. Papucci and K. Zurek, arXiv:1512.04928 [hep-ph]. P. S. B. Dev, R. N. Mohapatra and Y. Zhang, arXiv:1512.08507 [hep-ph]. J. Chakrabortty, A. Choudhury, P. Ghosh, S. Mondal and T. Srivastava, arXiv:1512.05767 [hep-ph]. A. Ahmed, B. M. Dillon, B. Grzadkowski, J. F. Gunion and Y. Jiang, arXiv:1512.05771 [hep-ph]. T. Higaki, K. S. Jeong, N. Kitajima and F. Takahashi, arXiv:1512.05295 [hep-ph]. B. Bellazzini, R. Franceschini, F. Sala and J. Serra, arXiv:1512.05330 [hep-ph]. R. Franceschini *et al.*, arXiv:1512.04933 [hep-ph]. J. Jaeckel, M. Jankowiak and M. Spannowsky, Phys. Dark Univ. **2**, 111 (2013) doi:10.1016/j.dark.2013.06.001 [arXiv:1212.3620 [hep-ph]]. J. J. Heckman, arXiv:1512.06773 [hep-ph]. H. Han, S. Wang and S. Zheng, arXiv:1512.07992 [hep-ph]. A. Alves, A. G. Dias and K. Sinha, arXiv:1512.06091 [hep-ph]. W. Chao, arXiv:1512.08484 [hep-ph]. C. Cai, Z. H. Yu and H. H. Zhang, arXiv:1512.08440 [hep-ph]. Q. H. Cao, Y. Liu, K. P. Xie, B. Yan and D. M. Zhang, arXiv:1512.08441 [hep-ph]. S. Chakraborty, A. Chakraborty and S. Raychaudhuri, arXiv:1512.07527 [hep-ph]. U. K. Dey, S. Mohanty and G. Tomar, arXiv:1512.07212 [hep-ph]. M. Backovic, A. Mariotti and D. Redigolo, arXiv:1512.04917 [hep-ph]. M. T. Arun and P. Saha, arXiv:1512.06335 [hep-ph]. Q. H. Cao, Y. Liu, K. P. Xie, B. Yan and D. M. Zhang, arXiv:1512.05542; A. Kobakhidze, F. Wang, L. Wu, J. M. Yang and M. Zhang, arXiv:1512.05585; P. Cox, A. D. Medina, T. S. Ray and A. Spray, arXiv:1512.05618; S. V. Demidov and D. S. Gorbunov, arXiv:1512.05723; W. Chao, R. Huo and J. H. Yu, arXiv:1512.05738; D. Curtin and C. B. Verhaaren, arXiv:1512.05753; L. Bian, N. Chen, D. Liu and J. Shu, arXiv:1512.05759; P. Agrawal, J. Fan, B. Heidenreich, M. Reece and M. Strassler, arXiv:1512.05775; C. Csaki, J. Hubisz and J. Terning, arXiv:1512.05776; D. Aloni, K. Blum, A. Dery, A. Efrati and Y. Nir, arXiv:1512.05778; E. Gabrielli, K. Kannike, B. Mele, M. Raidal, C. Spethmann and H. Veermäe, arXiv:1512.05961; R. Benbrik, C. H. Chen and T. Nomura, arXiv:1512.06028; J. S. Kim, J. Reuter, K. Rolbiecki and R. R. de Austri, arXiv:1512.06083; E. Megias, O. Pujolas and M. Quiros, arXiv:1512.06106; L. M. Carpenter, R. Colburn and J. Goodman, arXiv:1512.06107; W. Chao, arXiv:1512.06297; C. Han, H. M. Lee,

M. Park and V. Sanz, arXiv:1512.06376; S. Chang, arXiv:1512.06426; I. Chakraborty and A. Kundu, arXiv:1512.06508; R. Ding, L. Huang, T. Li and B. Zhu, arXiv:1512.06560; H. Han, S. Wang and S. Zheng, arXiv:1512.06562; X. F. Han and L. Wang, arXiv:1512.06587; D. Bardhan, D. Bhatia, A. Chakraborty, U. Maitra, S. Raychaudhuri and T. Samui, arXiv:1512.06674; T. F. Feng, X. Q. Li, H. B. Zhang and S. M. Zhao, arXiv:1512.06696; W. S. Cho, D. Kim, K. Kong, S. H. Lim, K. T. Matchev, J. C. Park and M. Park, arXiv:1512.06824; D. Barducci, A. Goudelis, S. Kulkarni and D. Sengupta, arXiv:1512.06842; O. Antipin, M. Mojaza and F. Sannino, arXiv:1512.06708; F. Wang, L. Wu, J. M. Yang and M. Zhang, arXiv:1512.06715; J. Cao, C. Han, L. Shang, W. Su, J. M. Yang and Y. Zhang, arXiv:1512.06728; F. P. Huang, C. S. Li, Z. L. Liu and Y. Wang, arXiv:1512.06732; X. J. Bi, Q. F. Xiang, P. F. Yin and Z. H. Yu, arXiv:1512.06787; L. Berthier, J. M. Cline, W. Shepherd and M. Trott, arXiv:1512.06799; M. Bauer and M. Neubert, arXiv:1512.06828; M. Chala, M. Duerr, F. Kahlhoefer and K. Schmidt-Hoberg, arXiv:1512.06833; S. M. Boucenna, S. Morisi and A. Vicente, arXiv:1512.06878; C. W. Murphy, arXiv:1512.06976; A. E. C. Hernandez and I. Nisandzic, arXiv:1512.07165; J. de Blas, J. Santiago and R. Vega-Morales, arXiv:1512.07229; P. S. B. Dev and D. Teresi, arXiv:1512.07243; W. C. Huang, Y. L. S. Tsai and T. C. Yuan, arXiv:1512.07268; S. Moretti and K. Yagyu, arXiv:1512.07462; K. M. Patel and P. Sharma, arXiv:1512.07468; M. Badziak, arXiv:1512.07541; W. Altmannshofer, J. Galloway, S. Gori, A. L. Kagan, A. Martin and J. Zupan, arXiv:1512.07616; M. Cvetič, J. Halverson and P. Langacker, arXiv:1512.07622; J. Gu and Z. Liu, arXiv:1512.07624; B. C. Allanach, P. S. B. Dev, S. A. Renner and K. Sakurai, arXiv:1512.07645; H. Davoudiasl and C. Zhang, arXiv:1512.07672; N. Craig, P. Draper, C. Kilic and S. Thomas, arXiv:1512.07733; K. Das and S. K. Rai, arXiv:1512.07789; K. Cheung, P. Ko, J. S. Lee, J. Park and P. Y. Tseng, arXiv:1512.07853; J. Liu, X. P. Wang and W. Xue, arXiv:1512.07885; J. A. Casas, J. R. Espinosa and J. M. Moreno, arXiv:1512.07895; L. J. Hall, K. Harigaya and Y. Nomura, arXiv:1512.07904.

[5] M. Dhuria and G. Goswami, arXiv:1512.06782.

[6] J. Zhang and S. Zhou, arXiv:1512.07889 [hep-ph].

[7] A. Salvio and A. Mazumdar, arXiv:1512.08184 [hep-ph].

[8] F. L. Bezrukov and M. Shaposhnikov, Phys. Lett. B **659** (2008) 703 doi:10.1016/j.physletb.2007.11.072, arXiv:0710.3755 [hep-th]; F. Bezrukov and M. Shaposhnikov, JHEP **0907** (2009) 089 doi:10.1088/1126-6708/2009/07/089, arXiv:0904.1537 [hep-ph]; Y. Hamada, H. Kawai, K. y. Oda and S. C. Park, Phys. Rev. Lett. **112** (2014) 24, 241301 doi:10.1103/PhysRevLett.112.241301, arXiv:1403.5043 [hep-ph]; F. Bezrukov and M. Shaposhnikov, Phys. Lett. B **734** (2014) 249 doi:10.1016/j.physletb.2014.05.074, arXiv:1403.6078 [hep-ph]; Y. Hamada, H. Kawai, K. y. Oda and S. C. Park, Phys. Rev. D **91** (2015) 053008 doi:10.1103/PhysRevD.91.053008, arXiv:1408.4864 [hep-ph].

[9] P. Channuie, J. J. Joergensen and F. Sannino, JCAP **1105**, 007 (2011) doi:10.1088/1475-7516/2011/05/007 [arXiv:1102.2898 [hep-ph]].

[10] L. Anguelova, Nucl. Phys. B **843**, 429 (2011) doi:10.1016/j.nuclphysb.2010.10.007 [arXiv:1006.3570 [hep-th]].

[11] H. Jiang, T. Liu, S. Sun and Y. Wang, arXiv:1512.07538 [astro-ph.CO].

[12] G. Aad *et al.* [ATLAS Collaboration], JHEP **1512**, 055 (2015) doi:10.1007/JHEP12(2015)055 [arXiv:1506.00962 [hep-ex]].

[13] B. Allanach, F. S. Queiroz, A. Strumia and S. Sun, arXiv:1511.07447 [hep-ph].

[14] B. C. Allanach, B. Gripaios and D. Sutherland, Phys. Rev. D 92, no. 5, 055003 (2015) arXiv:1507.01638 [hep-ph].

[15] S. Sun, D. B. Kaplan and A. E. Nelson, Phys. Rev. D **87**, no. 12, 125036 (2013) doi:10.1103/PhysRevD.87.125036 [arXiv:1303.1811 [hep-ph]]. S. Sun, arXiv:1411.0131 [hep-ph]. G. Moreau, Phys. Rev. D **87**, no. 1, 015027 (2013) doi:10.1103/PhysRevD.87.015027 [arXiv:1210.3977 [hep-ph]]. N. Bonne and G. Moreau, Phys. Lett. B **717**, 409 (2012) doi:10.1016/j.physletb.2012.09.063 [arXiv:1206.3360 [hep-ph]]. A. Angelescu, A. Djouadi and G. Moreau, arXiv:1510.07527 [hep-ph].

[16] A. Alves, A. Berlin, S. Profumo and F. S. Queiroz, JHEP **1510**, 076 (2015) doi:10.1007/JHEP10(2015)076 arXiv:1506.06767 [hep-ph].

[17] K. A. Olive *et al.* [Particle Data Group Collaboration], Chin. Phys. C **38**, 090001 (2014). doi:10.1088/1674-1137/38/9/090001

[18] N. Craig, H. K. Lou, M. McCullough and A. Thalapillil, arXiv:1412.0258 [hep-ph]. N. Craig, M. Farina, M. McCullough and M. Perelstein, JHEP **1503**, 146 (2015) doi:10.1007/JHEP03(2015)146 [arXiv:1411.0676 [hep-ph]]. S. Sun, arXiv:1510.02309 [hep-ph]. A. Aravind, M. Xiao and J. H. Yu, arXiv:1512.09126 [hep-ph].

[19] T. P. Cheng, E. Eichten and L. F. Li, Phys. Rev. D **9** (1974) 2259. doi:10.1103/PhysRevD.9.2259

[20] Y. Hamada, H. Kawai and K. y. Oda, Phys. Rev. D **87** (2013) 5, 053009 [Phys. Rev. D **89** (2014) 5, 059901] doi:10.1103/PhysRevD.87.053009, 10.1103/PhysRevD.89.059901 [arXiv:1210.2538 [hep-ph]]; D. Buttazzo, G. Degrassi, P. P. Giardino, G. F. Giudice, F. Sala, A. Salvio and A. Strumia, JHEP **1312** (2013) 089 doi:10.1007/JHEP12(2013)089 [arXiv:1307.3536 [hep-ph]].

[21] CMS Collaboration [CMS Collaboration], CMS-PAS-B2G-15-006.

[22] G. Aad *et al.* [ATLAS Collaboration], Phys. Rev. D **92** (2015) 112007 doi:10.1103/PhysRevD.92.112007 [arXiv:1509.04261 [hep-ex]].

[23] C. P. Burgess, H. M. Lee and M. Trott, JHEP **0909**, 103 (2009) doi:10.1088/1126-6708/2009/09/103 [arXiv:0902.4465 [hep-ph]].

[24] F. Bezrukov, A. Magnin, M. Shaposhnikov and S. Sibiryakov, JHEP **1101**, 016 (2011) doi:10.1007/JHEP01(2011)016 [arXiv:1008.5157 [hep-ph]].

[25] J. McDonald, arXiv:1604.01711 [hep-ph].

[26] P. A. R. Ade *et al.* [Planck Collaboration], arXiv:1502.01589 [astro-ph.CO].

[27] P. A. R. Ade *et al.* [BICEP2 s Collaboration], arXiv:1510.09217 [astro-ph.CO].

[28] Y. Hamada, T. Noumi, G. Shiu and S. Sun, *in progress.*

[29] The diboson resonance (WW,ZZ,WZ) around 2TeV reported earlier this year [12] by ATLAS might have some correlation to the current excess in the diphoton channel. See, e.g., [13, 14] for the physics

argument for the earlier diboson excess.

[30] This simple scalar extension of the SM has also appeared in the study of current and future collider probes to new physics and dark matter portals, see e.g. [18].

[31] We would like to thank A.Cohen for pointing it out that an $S|H|^2$ term can also be present, unless one considers a Z_2 symmetry acting both on S and fermions. However $S|H|^2$ is a dimension-3 soft operator which does not contribute much to the RG running in the UV limit.

[32] As pointed out in [23, 24], the cutoff scale derived from the tree-level unitarity is generically around the background inflaton value during inflation. Although it might suggest that non-renormalizable operators could spoil the flatness of the inflationary potential (a common feature for large field inflation models in general), we assume, as in the conventional Higgs inflation scenarios, that the inflaton potential is protected by a scale symmetry in the Jordan frame which as required by the unknown UV complete theory. See also a recent paper [25] for a possibility that the existence of a new scalar may help to push up the cutoff scale and resolve this issue.

[33] Notice that if only one of the two scalars plays the role of the inflaton, the inflationary dynamics is reduced to the single field one and no significant isocurvature remnants are generated during inflation.