

# Light Higgs from Pole Attractor

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We propose a new way of explaining the observed Higgs mass, within the cosmological relaxation framework. The key feature distinguishing it from other scanning scenarios is that the scanning field has a non-canonical kinetic term, whose role is to terminate the scan around the desired Higgs mass value. We propose a concrete realisation of this idea with two new singlet fields, one that scans the Higgs mass, and another that limits the time window in which the scan is possible. Within the provided time period, the scanning field does not significantly evolve after the Higgs field gets close to the Standard Model value, due to particle production friction.

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## I. INTRODUCTION

One of the main remaining puzzles of the Standard Model (SM), the Higgs mass, led physicists to search for heavy ElectroWeak (EW) charged new physics at the TeV scale, as predicted by various scenarios, such as supersymmetry and composite Higgs. An alternative approach to the problem, named cosmological relaxation [1] (see [2–15] for subsequent developments), does not, *a priori* require this to be the case and makes the new physics either too heavy and beyond the reach of the current colliders, or very light and very weakly coupled. Given this difference, it seems especially important to examine theoretically this new concept to the greatest possible extent.

The key ingredient of cosmological relaxation scenarios is the coupling of the Higgs to a new spin-zero field, the relaxion. This coupling induces the Higgs mass dependence on the relaxion field value. Cosmological evolution of the latter then leads to the Higgs mass scan, starting from some generic large value, down to the much smaller value which is currently observed. The scan is stopped in the right place due to a back-reaction of the Higgs on the relaxion evolution. Existing realizations of this mechanism feature a scalar potential characterised by two hierarchically different periods. The larger period is needed for the complete Higgs mass scan and the smaller one allows to settle the final Higgs mass at the EW scale. The task of producing a UV completion for such a potential is very non-trivial and requires a dedicated extra model engineering [16–19]. Interestingly, another known type of scanning scenarios, proposed in [20, 21], does not require this feature. Instead of producing the short-period potential barriers for the scanning field, the whole scanning sector effectively decouples from the Higgs sector close to the SM Higgs mass. In the same spirit acts the mechanism proposed in this work.

We will examine a possibility of the Higgs mass scan termination by a non-canonical kinetic term of the relaxion field  $\phi$ . For this to happen we will assume that the field-dependent pre-factor of  $(\partial_\mu\phi)^2$  starts diverging when the Higgs mass approaches its SM value. Enhancement of the kinetic term coefficient then results in the effective suppression of the relaxion potential and its coupling to the Higgs boson. With enough suppression the scan can slow down to an unobservably small speed. Relaxion field gets frozen around the value which gives the desired Higgs mass, which we will call an attractor point. Along this paper we will discuss concrete ways of implementing this idea. The main model-building challenge lies in finding a proper way to connect the value of the relaxion kinetic term with the Higgs mass. We start with the most straightforward realization featuring a Higgs-dependent kinetic term  $\sim 1/h^n$ . It turns out to be incapable of producing a light SM-like Higgs, but its detailed analysis proves to be useful in explaining some basic features of the mechanism and, more importantly, constructing realistic models. In the following we present one such model, featuring an extra scalar field  $\chi$  which controls the pole of the relaxion kinetic term.

The structure of this paper is the following. In Sec. II we start by introducing a toy model, in Sec. III we describe a more complex set-up, with two singlet fields, discuss its one-loop structure and its evolution before and after the Higgs mass scan. The details of the scan are discussed in Sec. IV, preceded by a brief review of the particle production friction. Finally we discuss our results in Sec. V.

## II. TOY MODEL

### A. Main Idea

As usual in the scanning scenarios, we promote the Higgs mass to a field-dependent variable by coupling the Higgs to another field, a spin-zero SM singlet  $\phi$ . While we assume the Higgs potential to take a generic form, controlled by a cutoff  $\Lambda$ , the interactions of the  $\phi$  field are kept under control by imposing the shift symmetry  $\phi \rightarrow \phi + c$ , which is weakly broken by a dimensionless parameter  $\kappa$ . The leading terms of the resulting scalar potential are

$$V = \kappa\Lambda\phi h^2 - \kappa\Lambda^3\phi - \Lambda^2 h^2 + \lambda h^4, \quad (\text{II.1})$$

which makes the Higgs mass parameter to depend on the  $\phi$  vacuum expectation value ( $vev$ )

$$m_h^2 = -\Lambda^2 + \kappa\Lambda\phi. \quad (\text{II.2})$$

Here and in the following we use  $h^2$  for  $h^\dagger h$ , and for conciseness omit most of the order-one factors, as well as the Higgs quartic coupling constant  $\lambda$  and the quartic coupling term itself. In Eq. (II.2) we fixed the  $\phi$ -independent part of the Higgs mass term to be negative, and the starting  $\phi$  value is chosen to be less than  $\Lambda/\kappa$ , ( $\phi = 0$  for simplicity) such that the Higgs  $vev$  is initially of a cutoff size<sup>1</sup>. To make the Higgs  $vev$  decrease and approach the SM value, we fixed the sign of the leading term of the  $\phi$  potential  $\kappa\Lambda^3\phi$  so that  $\phi$  increases with time. Notice that we have not required any fine tuning of the theory parameters. As for the initial conditions, if we assume a uniform distribution of the relaxion field values over different space points in the beginning of the process, only an order-one fraction of them will give the  $\phi$  value below  $\Lambda/\kappa$ . But despite not being able to assure the needed initial value, we are satisfied with an order-one probability for it. As the inflation stretches away the field inhomogeneities, we can assume our initial condition  $\phi = 0$  to hold soon after the beginning of inflation everywhere in the given causally connected part of the universe. All the discussion so far closely followed the original relaxion proposal [1], up to the sign of the initial Higgs mass.

As was anticipated in the introduction, the core of our scenario is the mechanism allowing to stop the scanning when the Higgs mass approaches the SM value, using the diverging  $\phi$  kinetic term. For this toy model we will simply take the kinetic lagrangian

$$\mathcal{L}_{\text{kin}} = \frac{\Lambda_k^{2n}}{h^{2n}}(\partial_\mu\phi)^2 + (\partial_\mu h)^2 \quad (\text{II.3})$$

where  $\Lambda_k$  is a mass dimension-one parameter and  $n$  is some positive power. For now we will not give any comments on possible UV completions, and go on with the analysis. As can be immediately read-off the Eq. (II.3), the pre-factor of the  $\phi$  kinetic term starts exploding upon approaching to  $h = 0$ . This means that every additional unit  $\phi$  variation takes more and more time, and at some point the  $\phi$  evolution effectively stops, with a Higgs  $vev$  and mass being close to zero. Interestingly, thus generated attractor point,  $\phi = \Lambda/\kappa$  and  $h = 0$ , does not correspond to neither local or global minimum of the  $\phi$  potential. Clearly, in order to reproduce the SM we need the attractor to be around the SM Higgs  $vev$   $h = v$  and not at zero. In this introductory section we will however limit ourselves with a less realistic but simpler case.

The metric on the field space described by the kinetic terms (II.3) is not flat, therefore we are not able to canonically normalize both fields in all the time points simultaneously. To have a first glance on the details of their evolution we will integrate out the Higgs field using  $h^2 \rightarrow (\Lambda^2 - \kappa\Lambda\phi)^2$ . We thus arrive at the one-field Lagrangian

$$\mathcal{L}_\phi = \frac{\Lambda_k^{2n}}{(-\kappa\Lambda\phi)^n}(\partial_\mu\phi)^2 + \kappa\Lambda^3\phi \quad (\text{II.4})$$

where, for simplicity, we performed a shift  $\kappa\Lambda\phi - \Lambda^2 \rightarrow \kappa\Lambda\phi$ . In new notations the attractor point simply corresponds to  $\phi = 0$  and the initial  $\phi$  value is negative. We can now switch to a new canonically normalized field  $\hat{\phi}$  so that

$$\frac{\Lambda_k^n}{(-\kappa\Lambda\phi)^{n/2}}\partial\phi = \pm\partial\hat{\phi} \quad (\text{II.5})$$

Depending on the power  $n$ , we choose the following redefinitions (omitting obvious constant factors)

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<sup>1</sup> We will assume some, at least small, separation between the initial Higgs mass and the cutoff to make our EFT analysis meaningful.

<sup>2</sup> This substitution is not always valid, as the Higgs potential can also receive contributions from the  $\phi$  kinetic term. The regime of its validity will be discussed later on.

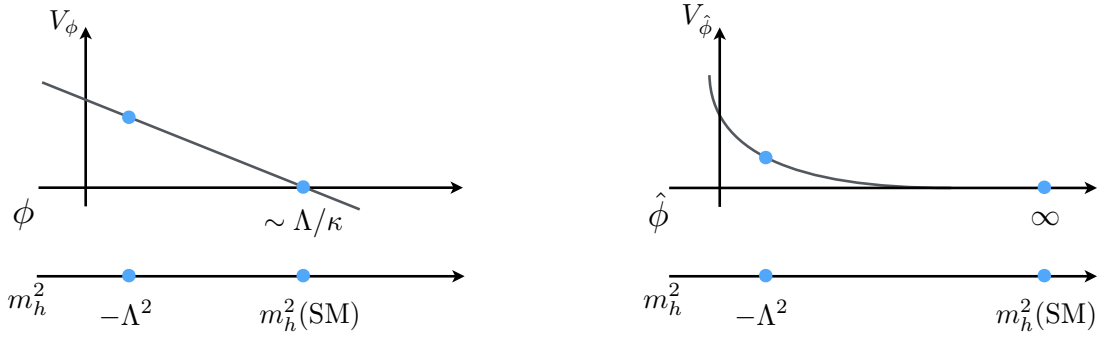


FIG. 1: Schematic representation of the scanning field potential and the corresponding Higgs mass, in terms of the initial field  $\phi$  and after its canonical normalization, as explained in the text.

$n = 1$	$-\phi \sim \hat{\phi}^2$	$(-\infty, 0) \rightarrow (-\infty, 0)$
$n = 2$	$-\phi \sim \exp[-\hat{\phi}]$	$(-\infty, 0) \rightarrow (-\infty, \infty)$
$n > 2$	$-\phi \sim \hat{\phi}^{\frac{1}{1-n/2}}$	$(-\infty, 0) \rightarrow (0, \infty)$

where in the last column we showed how the interval  $\phi \in (-\infty, 0)$  maps onto the canonical field  $\hat{\phi}$ . It shows that the  $n = 1$  case is special as the attractor point is mapped onto a finite value of  $\hat{\phi}$ . This potentially brings a problem once we try to move the attractor point away from  $h = 0$  to some finite non-zero value. The simplest way to do so is by changing the  $\phi$  kinetic term to  $(\partial_\mu \phi)^2 / (h^2 - \Delta^2)^n$ . Now, for  $n = 1$  the  $h^2 = \Delta^2$  corresponds to a finite  $\hat{\phi}$ . Hence it can be reached and also overshoot, making the  $\phi$  kinetic term negative.

For  $n \geq 2$ , on the other hand, the attractor point is mapped to infinity and hence  $\hat{\phi}$  will be eternally approaching it. This is because the stretching of the  $\phi$  field outruns the  $\phi$  time variation as we approach the attractor along the chosen trajectory. The resulting evolution of the  $h - \phi$  system for  $n \geq 2$  is schematically depicted on Fig. 1. It is interesting to notice that the same behaviour which we observe for  $\hat{\phi}$  after setting the Higgs into minimum of its potential, is used in various models of pole inflation (see e.g. [22] and references therein), with a difference that the inflaton field evolution goes away from what we call the attractor point.

We will now perform a slight refinement of the previous analysis, which will highlight a few additional features of the toy model. In particular we will see why the toy model can not be used to produce a light SM-like Higgs boson.

## B. Non-Perturbativity

Let us jump straight to the major phenomenological problem of this model, namely the exploding  $h - \phi$  coupling. In order to make it apparent we will split  $\phi$ ,  $\dot{\phi}$  and  $h$  into classical and fluctuation parts by

$$\phi = \phi_0 + \delta\phi, \quad \dot{\phi} = \dot{\phi}_0 + \delta\dot{\phi}, \quad h = h_0 + \delta h, \quad (\text{II.6})$$

where the zero subscript denotes classical background at some time  $t_0$ . On top of this we will locally (at the time point  $t_0$ ) canonically normalize the  $\phi$  field fluctuations  $\delta\phi \rightarrow \hat{\phi} h_0^n / \Lambda_k^n$ . After these manipulations the kinetic terms of the field fluctuations are contained in

$$\mathcal{L}_{\text{kin}} \supset \frac{h_0^{2n}}{h^{2n}} (\partial_\mu \hat{\phi})^2 + (\partial_\mu h)^2 \quad (\text{II.7})$$

while the scalar potential (II.1) becomes, omitting fluctuation-independent terms

$$V = (\kappa \Lambda \phi_0 - \Lambda^2) h^2 + \kappa \Lambda \frac{h_0^n}{\Lambda_k^n} \hat{\phi} h^2 - \kappa \Lambda^3 \frac{h_0^n}{\Lambda_k^n} \hat{\phi} \quad (\text{II.8})$$

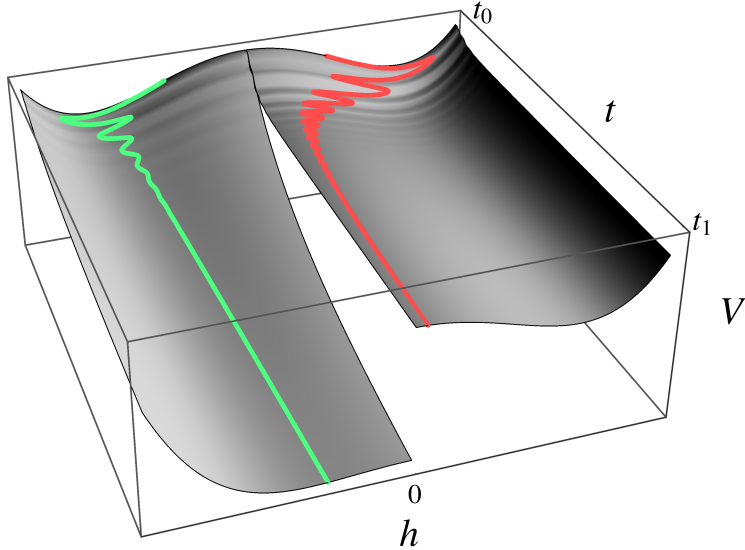


FIG. 2: Example of the Higgs field evolution without locking (left part, green line) and with locking (right part, red line). Values of the  $\hat{\phi}_0$ -independent part of the scalar potential  $V$  are shown in grey. In the first case the Higgs field follows the  $V$  minimum which evolves towards  $h^2 = 0$ . While in the second case the Higgs  $vev$  is driven towards zero, with the  $V$  minimum frozen close to the initial value.

We can now for instance estimate the amplitude of the Higgs decay into two  $\hat{\phi}$ 's. Close to the attractor ( $h_0^2 \rightarrow 0$ ) the  $\hat{\phi}$  potential is negligible and we can treat  $\hat{\phi}$  as massless. The  $h - \hat{\phi}$  interaction then arises from the first term of Eq. (II.7). The resulting non-SM Higgs decay amplitude

$$A_{h \rightarrow \hat{\phi}\hat{\phi}} \sim \frac{m_h^2}{h_0} \quad (\text{II.9})$$

is expected to be sizeable, far beyond the experimental bounds, and moreover ill-behaving close to the attractor.

### C. Locking of the pole field

Even though the previously discussed problem rules out the toy model, we will still make use of it to explain the locking of the  $h$  field. It will be useful in the following as it also applies to any other field which produces the kinetic pole. After singling out the background component of the  $\phi$  field (II.6), the kinetic lagrangian (II.3) also generates a  $\dot{\phi}_0$ -dependent term contributing to the Higgs potential

$$\delta V = -\frac{\Lambda_k^{2n}}{h^{2n}} \dot{\phi}_0^2 \quad (\text{II.10})$$

which is minimized at  $h^2 = 0$ , thus competing with the rest of the Higgs potential which prefers  $h^2 = \Lambda^2 - \kappa\Lambda\phi$ . In order to understand when this extra term becomes important we need to find  $\dot{\phi}_0$  by solving the equation of motion (*e.o.m.*) following from the lagrangians (II.1) and (II.3)

$$\ddot{\phi}_0 + \left( 3H + n \frac{\partial_t(-h^2)}{h^2} \right) \dot{\phi}_0 = \frac{h^{2n}}{\Lambda_k^{2n}} V'_\phi, \quad (\text{II.11})$$

where  $H$  is the Hubble parameter which appears after accounting for a metric expansion of the Universe. Besides the usual Hubble friction the equation contains a friction-like term  $\sim \partial_t h^2$  coming from the non-canonical form of the kinetic term. In order to estimate the maximal  $\dot{\phi}_0$  we will consider the slow-roll limit,

*i.e.*  $\ddot{\phi}_0$  negligible compared to other terms. The maximal value of  $\dot{\phi}_0$  is achieved when the friction is minimised, hence determined mostly by the irreducible Hubble expansion contribution

$$\dot{\phi}_{0\max} \simeq \frac{h^{2n}}{3H\Lambda_k^{2n}} V'_\phi \quad (\text{II.12})$$

Now substituting  $\dot{\phi}_{0\max}$  into Eq. (II.10) we can estimate the maximal  $\partial\delta V/\partial h$  and conclude that it is negligible compared to the potential (II.1) if

$$\frac{(\kappa\Lambda)^2}{H^2} \frac{\Lambda^4}{\Lambda_k^4} \frac{h^2}{m_h^2} \frac{h^{2(n-2)}}{\Lambda_k^{2(n-2)}} < 1. \quad (\text{II.13})$$

In the opposite case  $\delta V$  drives the Higgs *vev* to zero, thus terminating  $\phi$  evolution independently of the  $\Lambda^2 - \kappa\Lambda\phi$  value. We will call this termination process locking. The evolution of the Higgs field with and without locking is shown on Fig. (2). Therefore  $\delta V$  can significantly distort the evolution of the fields, and requires a special attention when considering this type of models. For the realistic model of the next section we will have to forbid this behaviour for the relaxation to work.

### III. FORMULATION OF A TWO-FIELD MODEL

The toy model analysed in Sec. II was shown to fall into a strongly coupled regime close to the attractor, thus failing to reproduce the Standard Model. We will now show that a tractable realistic model of pole attractor can be constructed using one additional spin-zero field  $\chi$  which controls the relaxation kinetic term. The main goal of this section is to define the general structure of the two field ( $\phi$  and  $\chi$ ) model, while the detailed analysis and numerical results will be presented in Sec. IV.

#### A. Formulation

The discussion of Section II suggests that the Higgs field can not be simply put in the  $\phi$  field kinetic term denominator. Hence we will introduce another spin-zero singlet field  $\chi$  to produce a pole in the  $\phi$  kinetic term. This will allow for more freedom in choosing the pole field properties, in particular, we would like to make the  $\chi$  kinetic term to have the same type of pole as the one of  $\phi$ . In the following we will consider

$$\mathcal{L}_{\text{kin}} = \frac{\Lambda_k^2}{\chi^2} \{(\partial_\mu\phi)^2 + (\partial_\mu\chi)^2\} \quad (\text{III.1})$$

This solves the problem of strong coupling pointed out in Sec. II B. We recall that it arises from the ill-behaved expansion of the Higgs field around the classical value  $h_0 \rightarrow 0$

$$\frac{(\partial_\mu\phi)^2}{h^{2n}} \rightarrow \sum_k \left(\frac{\delta h}{h_0}\right)^k \frac{(\partial_\mu\phi)^2}{h_0^{2n}} \quad (\text{III.2})$$

Resulting interaction terms remain divergent even after  $\phi$  is normalized canonically, as each  $\phi$  removes only  $n$  powers of  $h_0$ . Thus we get rid of this problem by switching to  $\chi$ , which gets normalized as well, absorbing the remaining poles. The resulting theory is tractable within the usual perturbative approach. Notice that from this argument we are not strictly required to have the same order of poles for  $\chi$  and  $\phi$ , but we will stick to this option for definiteness. Additionally, having even-order poles we are safe from the problem of negative kinetic terms. The structures of the type (III.1) are also interesting as they are frequently used and motivated in supergravity models of inflation (see e.g. [23] for a review). In particular, the kinetic lagrangian (III.1) can be associated to the Kähler potential  $K \sim -\Lambda_k^2 \log[\Phi + \bar{\Phi}]$  of a chiral superfield  $\Phi$ . The scanning fields then are linked to the bosonic components of  $\Phi$  as  $\text{Re}\Phi \sim \chi$  and  $\text{Im}\Phi \sim \phi$ .

The next step would be to construct a mechanism which blocks the  $\phi$  evolution as the Higgs mass becomes small. To achieve this, we will build a model that gives the following behaviour. First, both  $\chi$  and  $\phi$  roll down

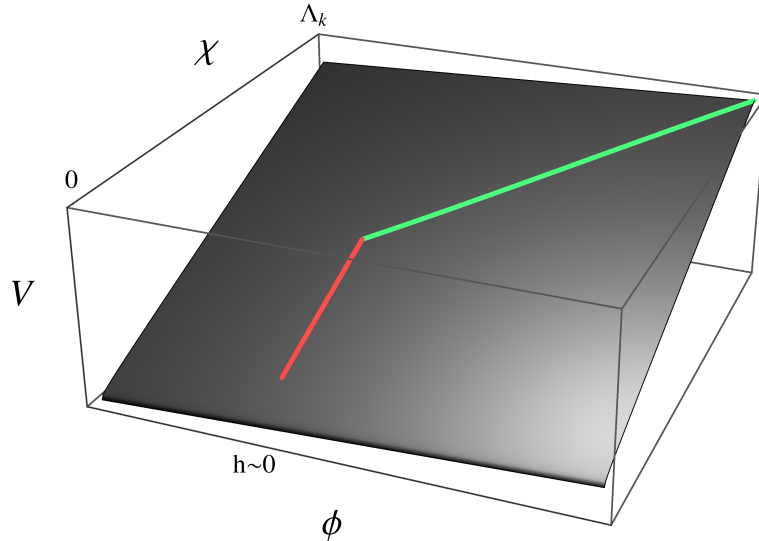


FIG. 3: Schematic plot of the  $\phi - \chi$  system evolution. The change between the initial (green) and final (red) regimes happens around  $h \sim 0$  due to appearance of the particle production friction of the  $\phi$  field.

their potential from the beginning of inflation, and the  $\phi$  field scans the Higgs mass the same way as before. Second, as the Higgs  $vev$  reaches the SM value, the relative speed  $\dot{\phi}/\dot{\chi}$  drops by a large factor. Therefore during the rest of the evolution until  $\chi$  gets close to the pole, the field  $\phi$  displaces by a much smaller amount than it did before the speed drop. As  $\chi$  approaches the pole it blocks the  $\phi$  evolution. The time dependence of the scanning fields is schematically shown on Fig. 3. In order to produce the relative speed drop we will use the Higgs-dependent particle production friction. For instance, if  $\phi$  has a particle friction, its time variation will be limited in the simplest case by  $\dot{\phi} \sim Hf$ , where  $H$  is the Hubble parameter and  $f$  is a mass scale suppressing the particle production. In the absence of sizeable particle production one will have instead a larger  $\dot{\phi}$  controlled by the Hubble friction  $\dot{\phi} \sim V'_\phi/H$ . We postpone the detailed discussion of different types of friction to Section IV A and continue with the general description of the model.

Within the chosen approach, i.e. changing friction, we need to require that

- the active scanning region (when  $\chi$  is away from the pole) is long enough so the Higgs mass can be completely scanned;
- after the Higgs mass reaches the SM value,  $\dot{\phi}/\dot{\chi}$  has to decrease by at least a factor of  $v^2/\Lambda^2$ , so that  $h$  mass is not changed significantly afterwards.

In the following we will present a model where the  $\dot{\phi}/\dot{\chi}$  drop originates solely from the growth of  $\phi$  friction around  $h = v$ , while the  $\chi$  friction is insensitive to  $h$ . The sole purpose of  $\chi$  is thus to provide a limited time window for the scan, before it gets close to zero. This construction looks very different from the toy model. But the underlying principles, broadly defined, are similar – the kinetic pole slows down  $\phi$  when  $h$  approaches  $v$ . The difference is that in the two-field model this back-reaction on  $\phi$  is delayed by the time  $\chi$  needs to fall to the pole. This results in a certain amount of residual  $\phi$  displacement which is suppressed by the high friction.

To fix the conventions we present the general lagrangian of our model, omitting for the moment the terms relevant for the particle production

$$\mathcal{L} = \frac{\Lambda_k^2}{\chi^2} \{(\partial_\mu \phi)^2 + (\partial_\mu \chi)^2\} + \kappa_\phi \Lambda^3 \phi + \kappa_\chi \Lambda^3 \chi \quad (\text{III.3})$$

$$- \kappa_h \Lambda h^2 \phi + \Lambda^2 h^2, \quad (\text{III.4})$$

where  $\kappa$ 's are dimensionless parameters related to shift symmetry breaking. Hence one can expect to have a freedom in making them small, and thus control the size of the scalar potential. One however naturally

expects  $\kappa_\phi \gtrsim \kappa_h/(4\pi)^2$ . Moreover, as we will discuss in the next section,  $\kappa_\chi = 0$  does not guarantee a shift symmetry for  $\chi$ , as the kinetic lagrangian is not shift invariant. We will assume that the fields are evolving in the range  $\phi \in (0, \Lambda/\kappa_\phi)$  and  $\chi \in (-\Lambda_k, 0)$ . The choice of signs of different terms in the lagrangian (III.3) already suggests that we will exploit the scanning with an initially large Higgs *vev*.

## B. One loop potential

We will now discuss quantum corrections to the evolution of the scanning fields. The questions discussed in the remainder of this section do not affect the details of the Higgs scanning (for it see Section IV) and are related to the periods before and after the Higgs mass scan. They are important though for a general consistency of the mechanism and also specify the requirements to possible UV completions. This discussion will be qualitative and we will only obtain the general form of the most important quantum corrections omitting, in particular, terms dependent on  $\dot{\phi}, \dot{\chi} \ll \Lambda^2$ . Before computing the one-loop correction let us first switch to the field  $\hat{\chi}$  defined as  $\chi = -\Lambda_k \exp[-\hat{\chi}/\Lambda_k]$ . The field range  $\chi \in (-\Lambda_k, 0)$  is then mapped onto  $\hat{\chi} \in (0, \infty)$ . After this redefinition we obtain

$$\mathcal{L}_\chi = \frac{\Lambda_k^2}{\chi^2} (\partial_\mu \chi)^2 + \kappa_\chi \Lambda^3 \chi \quad \rightarrow \quad (\partial_\mu \hat{\chi})^2 - \kappa_\chi \Lambda^3 \Lambda_k e^{-\hat{\chi}/\Lambda_k} + \frac{\Lambda^4}{(4\pi)^2} \hat{\chi}/\Lambda_k \quad (\text{III.5})$$

$$\mathcal{L}_\phi = \frac{\Lambda_k^2}{\chi^2} (\partial_\mu \phi)^2 + \kappa_\phi \Lambda^3 \phi \quad \rightarrow \quad e^{2\hat{\chi}/\Lambda_k} (\partial_\mu \phi)^2 + \kappa_\phi \Lambda^3 \phi \quad (\text{III.6})$$

Notice the presence of a term linear in  $\hat{\chi}$ . It appears due to the quartically divergent integral from the change of the path integral measure

$$\mathcal{D}\chi \exp[iS] \rightarrow \mathcal{D}\hat{\chi} \exp[iS + \text{Tr} \log[\partial\chi/\partial\hat{\chi}]]. \quad (\text{III.7})$$

Now computing the one-loop effective potential we find

$$V^{(1\text{-loop})} \sim \frac{\Lambda^4}{(4\pi)^2} \log \left[ \Lambda^2 e^{2\hat{\chi}/\Lambda_k} + \dots \right], \quad (\text{III.8})$$

which also depends linearly on  $\hat{\chi}$  for large  $\hat{\chi}$ . Comparing the differentials of the tree and loop level potentials of Eq. (III.5) with respect to  $\hat{\chi}$

$$\frac{\partial_{\hat{\chi}} V_{\text{tree}}}{\partial_{\hat{\chi}} V_{\text{loop}}} = \frac{(4\pi)^2 \kappa_\chi \Lambda_k}{\Lambda} e^{-\hat{\chi}/\Lambda_k} \quad (\text{III.9})$$

we find that the tree-level potential can only be relevant for  $\hat{\chi} \lesssim \Lambda_k$  and  $(4\pi)^2 \kappa_\chi \Lambda_k > \Lambda$ . We will assume that these requirements are satisfied during the active phase of the Higgs mass scan, as we would like to have a control over the  $\chi$  potential during it. The former constraint  $\hat{\chi} \lesssim \Lambda_k$  is also needed to ensure that the  $\phi$  evolution during the scan is unsuppressed and weakly dependent on  $\chi$ . We can add here another condition that  $\Lambda_k$  has to satisfy, namely  $\kappa_\chi \chi < \Lambda$  and consequently  $\kappa_\chi \Lambda_k < \Lambda$ , to provide a convergence of the  $\kappa_\chi \chi/\Lambda$  expansion of our effective field theory. These constraints lead to  $\Lambda_k \sim \Lambda/\kappa_\chi$ , which we will assume in the following.

The one loop correction only becomes important for  $\hat{\chi} > \Lambda_k$ , *i.e.* when the  $\phi$  kinetic term is already enhanced and the Higgs mass scan is mostly ended. In the following we will discuss the effects of this correction on the evolution outside the scanning window, namely on setting the initial and final conditions for the scan.

## C. Final vacuum after the scan

Depending on its sign the correction (III.8) would either block the  $\chi$  movement to the pole and thus spoil the mechanism, or make it move towards the pole even faster. Assuming the latter to be the case, the  $\chi$



potential becomes unbounded from below. Any phenomenologically viable UV completion of this type of model will therefore be required to contain a mechanism regularizing the scalar potential in the vicinity of  $\chi = 0$ . This can be done for instance by adding to the potential an extra piece with a different functional form than (III.8) to balance it. As another option we could shift the kinetic pole by a small constant

$$\frac{\Lambda_k^2}{\chi^2 + \epsilon^2} \{(\partial_\mu \chi)^2 + (\partial_\mu \phi)^2\}. \quad (\text{III.10})$$

This shift defines the maximal enhancement of the kinetic term, and thus the minimal slope of the  $\phi$  potential and the time variation of the Higgs mass. It is thus limited by

$$\delta h^2 \simeq \kappa_h \Lambda \delta \phi \simeq \kappa_h \kappa_\phi \Lambda^4 \left( \frac{\epsilon}{\Lambda_k} \right)^2 t^2 \ll v_{\text{SM}}^2 \quad (\text{III.11})$$

which for the current age of the Universe  $t \sim 10^{41} \text{ GeV}^{-1}$  gives  $\kappa(\epsilon/\Lambda_k) < 10^{-39}/(\Lambda/\text{GeV})^2$ . Such a correction will not affect any details of the scanning mechanism and hence we will not discuss it any further. As a consequence of such a regularisation,  $\chi$  can actually reach the minimum of its potential and stop its evolution there.

#### D. Initial conditions

The kinetic poles make the volume of the field space increasingly “stretched” upon approaching to  $\chi = 0$ . Thus, assuming uniformly distributed initial values of the renormalized fields over different patches of the universe before inflation, we would find that most of the patches have  $\chi$  which is very close to zero, almost completely blocking any possible Higgs mass scan. Needed for a successful scan  $|\chi| \sim \Lambda_k$ , instead, would correspond to a very tuned, non-typical initial condition. We would like to emphasize that this problem only arises if  $\chi$  field values are indeed distributed with a weight defined by the size of the kinetic terms. To verify this assumption we would need to know the exact UV completion. It can also be the case that the UV complete theory automatically sets initial values of  $\chi$  sufficiently far from the pole. We will now show that we do not necessarily need to rely on this latter possibility and there can be ways to successfully complete the scan even with the uniformly distributed values of the renormalized fields.

One of the ways to solve this issue would be in adding the second kinetic pole at  $\chi \lesssim -\Lambda_k$ . It would stretch the field space at large  $|\chi|$  and produce the second attractor value for the initial conditions. It will now be equally probable to start around the first or the second pole. Further, we require that the slope of the  $\chi$  potential around this new pole repulses  $\chi$  away from it towards zero. In this way, once  $\chi$  starts its evolution close to the new pole, it will unavoidably pass the region  $\chi \sim \Lambda_k$  where the scan can happen, and then evolve to zero. If the scalar potential in the vicinity of poles is determined by its one-loop expression, its monotonic decrease with  $\chi$  requires

$$V^{(1\text{-loop})} \sim i \int \frac{d^4 k}{(2\pi)^4} \log \left[ \frac{1}{(\chi - c_i)^2} + \frac{1}{\chi^2} \right] \rightarrow \begin{cases} -\Lambda^4 \log(\chi - c_i)^2 & \chi \sim c_i \\ \Lambda^4 \log \chi^2 & \chi \sim 0 \end{cases} \quad (\text{III.12})$$

where  $-c_i > 0$  is the position of the extra pole. This sign-changing behaviour may be achieved if the cutoff physics is sensitive to the *vev*'s of the  $\chi$ ,  $\phi$  and  $h$  fields, whose values change by an amount comparable to the cutoff during the evolution from one pole to another. An interesting consequence of such a construction is that  $\chi$  becomes almost completely decoupled from all the other fields in the beginning and in the end of its evolution. It is only active in the window around  $|\chi| \sim \Lambda_k$  when the Higgs mass scan happens.

Alternatively, we could use the slowly varying  $\chi$  field as the dominant source of inflation, in this case  $\chi$  far from the minimum of its potential, and from the pole, would be a natural initial condition. In this case there is no need for the second pole, but a detailed analysis of such a possibility lies beyond the scope of this paper.

## IV. RELAXATION IN THE TWO-FIELD MODEL

### A. Review of Particle Production Friction

As we have estimated in Sec. III for the two-field attractor, we need to produce an order  $v^2/\Lambda^2$  drop of the ratio  $\dot{\phi}/\dot{\chi}$  when the Higgs mass approaches the SM value. This section is dedicated to the brief review of the process allowing for this drop – the particle production friction. The results given here are mostly based on works [13–15, 24] where the particle friction was applied to the relaxion models, and the original model of inflation with a particle production [25]. We would like to emphasize that in the following we will be relying on analytic estimates of the relaxation dynamics. A comprehensive numerical study, while being important, lies beyond the scope of this paper. The results presented in this section will be applied to the two-field attractor dynamics in Sec. IV B.

We will consider an abelian field  $A_\mu$  with a mass  $m_A$  coupled to one of the scanning fields (*e.g.*  $\phi$  for definiteness) by means of an interaction

$$\mathcal{L} \supset \frac{\phi}{f} F_{\mu\nu} \tilde{F}^{\mu\nu} \quad (\text{IV.1})$$

where  $F_{\mu\nu}$  is the corresponding field strength tensor and  $\tilde{F}^{\mu\nu}$  its dual. In the time-dependent  $\phi$  background the transverse components of the  $A$  field can acquire exponentially growing modes, draining the  $\phi$  field kinetic energy, the process called particle friction. To see how it appears we first write down the solutions of the *e.o.m.* for two transverse polarisations of  $A$ , derived using the WKB approximation [24]

$$A[k]_\pm \sim \frac{1}{\sqrt{\omega_\pm}} e^{-i \int d\tau \omega_\pm} \quad \text{with} \quad \omega_\pm^2 = k^2 \mp ak\dot{\phi}/f + a^2 m_A^2 + a^2 \Pi_t[\omega_\pm, k] \quad (\text{IV.2})$$

where  $\pm$  stand for right and left helicity,  $a$  is a scale factor of the expanding universe,  $\tau$  is a conformal time  $ad\tau = dt$ ,  $k$  is a 3-momentum. The approximation (IV.2) is valid for  $|\partial_\tau \omega/\omega^2| < 1$ . Given that we are looking for exponentially growing gauge fields, one can end up in a space filled with a plasma of particles charged under  $A$ . Therefore in Eq. (IV.2) we have also included the thermal correction to the dispersion relation  $\Pi_t$  [27, 28]

$$\Pi_t[\omega, k] = m_D^2 \frac{\omega}{k} \left( \frac{\omega}{k} + \frac{1}{2} \left( 1 - \frac{\omega^2}{k^2} \right) \ln \frac{\omega + k}{\omega - k} \right) \quad (\text{IV.3})$$

with  $m_D^2 = g_A^2 T_p^2/6$  defining the Debye mass of a plasma with a temperature  $T_p$ . If the dispersion relation (IV.2) allows for imaginary  $\omega \equiv i\Omega$ , the vector field can experience exponential growth with time. Let us first notice that  $\Pi_t$  is a positive function for complex  $\omega$ , therefore the existence of complex solutions of (IV.2) for one of the two polarisations requires

$$a\dot{\phi} > 2m_A f. \quad (\text{IV.4})$$

Where without loss of generality we have assumed  $\dot{\phi} > 0$ . We further notice that in the  $\Omega \sim k$  limit  $\Pi_t$  saturates around  $m_D^2$  while for  $\Omega \ll k$  we obtain  $\Pi_t \sim m_D^2 |\Omega/k|$ . Therefore, once the condition (IV.4) is satisfied, the maximal  $\Omega$ , and hence maximal instability, is given by

$$\Omega_{\max} \sim \begin{cases} a\dot{\phi}/f & \text{for } a\dot{\phi} > 2m_D f \\ a\dot{\phi}/f \frac{(\dot{\phi}/f)^2}{m_D^2} & \text{for } a\dot{\phi} \ll m_D f \end{cases} \quad (\text{IV.5})$$

which shows that the instability growth becomes weaker in plasma. In all the cases the instability is maximized around

$$k_{\max} \simeq a\dot{\phi}/2f \quad (\text{IV.6})$$

We hence identified three regimes of  $A_\mu$  evolution: with no exponential growth, with a fast growth, and with a growth slowed down by the high temperature plasma,  $T_p \gtrsim \dot{\phi}/f$ . In case of growing instability we can expect

$$\langle F_{\mu\nu}\tilde{F}_{\mu\nu} \rangle \sim \langle F_{\mu\nu}F_{\mu\nu} \rangle \sim \langle A_\mu^2 \rangle \sim H^4 \exp \left[ 2 \int d\tau \Omega_{\max} \right]. \quad (\text{IV.7})$$

Notice that the time of the efficient exponential growth is limited by roughly one Hubble time, as after that the produced gauge field modes become significantly red-shifted. In Appendix VI we collect the precise expressions for the quantities listed in Eq. (IV.7). Thus produced energy density stored in the gauge field can lead to several effects.

- First, growing gauge field modes back-react on the rolling field  $\phi$ . The  $\phi$  *e.o.m.* reads

$$\ddot{\phi} + 3H\dot{\phi} - \frac{1}{f}\langle F_{\mu\nu}\tilde{F}^{\mu\nu} \rangle + V'_\phi = 0. \quad (\text{IV.8})$$

For the sake of this section we forget about the non-canonical kinetic terms, as the scan happens around  $|\chi| \sim \Lambda_k$ , *i.e.* in the regime where the kinetic terms are not significantly enhanced. Hence our results will remain parametrically accurate. In case of negligible  $\langle F_{\mu\nu}\tilde{F}^{\mu\nu} \rangle$  the  $\phi$  evolution is driven by the slope of the potential  $V$  and the Hubble friction, with a maximal speed defined by

$$3H\dot{\phi} + V'_\phi = 0 \quad \Rightarrow \quad \dot{\phi} = -\frac{V'_\phi}{3H}. \quad (\text{IV.9})$$

If instead the term  $\langle F_{\mu\nu}\tilde{F}^{\mu\nu} \rangle$  dominates, one can reach a stationary regime with the evolution defined by the last two terms of Eq. (IV.8):

$$-\frac{1}{fa^4}\langle F_{\mu\nu}\tilde{F}^{\mu\nu} \rangle + V'_\phi = 0. \quad (\text{IV.10})$$

Using the exact dependence of  $\langle F_{\mu\nu}\tilde{F}^{\mu\nu} \rangle$  on  $\dot{\phi}$  we can estimate the maximal  $\dot{\phi}$  as

$$\dot{\phi} \simeq fH \left( \frac{\max[H, T_p]}{H} \right)^{2/3} \quad (\text{IV.11})$$

To find out which of the two types of friction dominates we simply need to check which of the two gives the minimal  $\dot{\phi}$ .

- Second, if the gauge field  $A_\mu$  couples to the Higgs boson at tree level, through  $g_A^2 A^2 h^2$ , it can lead to a restoration of the electroweak symmetry due to the effective temperature which we define as  $T_{\text{eff}}^2 \equiv \langle A_\mu^2 \rangle$ . Therefore we do not need the Higgs and the gauge field to enter thermal equilibrium for this mass correction to appear, and  $T_{\text{eff}}$  simply describes the classical value the  $A$  field.
- Finally, the energy density stored in the gauge bosons can thermalise, leading to a creation of a thermal plasma. This plasma can then slow down the gauge field growth as described above. For thermalization to occur we need it to be faster than the Hubble expansion rate. We will consider two plasma production channels, a perturbative pair production and a non-perturbative Schwinger production.

Efficient perturbative pair production of charged fermions happens if the typical gauge quanta energy is higher than the fermion mass  $m_f$ , and if the pair production cross section, enhanced by the large multiplicity of gauge quanta  $N_\gamma$ , is higher than the Hubble expansion rate [26]

$$\Omega > m_f \quad \text{and} \quad N_\gamma > \frac{128\pi}{N_f g_A^4}, \quad (\text{IV.12})$$

where  $N_f$  is number of produced fermionic degrees of freedom and  $g_A$  is a gauge coupling. The plasma temperature can be estimated as an order-one fraction of the overall energy released by the rolling  $\phi$  field in one Hubble time,  $T_p^4 \sim V'_\phi \dot{\phi} H^{-1}$ .

The non-perturbative Schwinger production is characterized by

$$\Gamma = \frac{(g_A E)^2}{4\pi^3 H^3} \exp\left[-\frac{\pi m_f^2}{g_A E}\right], \quad (\text{IV.13})$$

where  $E$  stands for a modulus of the electric field analogue for  $A$ . For this production channel to be efficient one needs

$$g_A E \gg m_f^2, H^2. \quad (\text{IV.14})$$

And the maximal plasma temperature can be estimated as an order-one fraction of energy stored in the electric field  $T_p^4 \simeq E^2$  [15].

Once one of the fermion production channels opens,  $A$  field modes can thermalize. Thermalization will happen if there is enough time for  $A$  modes to interact with the plasma before they exit the horizon. Concretely, we require the mean free path of gauge quanta to be less than the Hubble length [15]

$$g_A^4 T_p > (4\pi)^2 H. \quad (\text{IV.15})$$

We now have at our disposal three different friction regimes – Hubble friction, thermal particle friction, non-thermal particle friction – with a potentially significant relative speed change. We have also identified the criteria leading to switching between different regimes, which one needs to satisfy when the Higgs mass gets close to its SM value.

## B. Details of Evolution and Numerical Results

For a concrete realization of the general scanning scheme presented in Section III A we choose the case with a constant  $\chi$  friction and a  $\phi$  friction increasing around the SM Higgs  $vev$ . The friction is produced by the term [13]

$$\frac{\phi}{f}(g_2^2 W_{\mu\nu} \tilde{W}^{\mu\nu} - g_1^2 B_{\mu\nu} \tilde{B}^{\mu\nu}) \quad (\text{IV.16})$$

where  $W$  and  $B$  correspond to the SM weak and hypercharge gauge bosons and  $g_{1,2}$  are their gauge couplings. The combination of the field strengths in Eq. (IV.16) is chosen so that it does not contain a photon, which is massless during the whole scan, hence not sensitive to the Higgs mass change. Starting at  $h \sim \Lambda$  with no particle friction we want to switch it on at  $h \sim v$  due to decreased gauge boson mass.

Notice that non-abelian gauge bosons develop the so-called magnetic mass in plasma  $m_M \sim g_2^2 T_p$  [29], which blocks the instability development. This can happen to the  $W$  and even the  $Z$  boson, as in the broken EW phase it contains  $W_3$ . The way out in this case would be a restoration of the EW symmetry, making the abelian hypercharge boson a mass eigenstate. As will be discussed later, the restoration of the EW symmetry is also necessary for another unrelated reason.

As for  $\chi$ , we can either leave it with the Hubble friction only, or assign it some other particle friction. In the latter case the easiest would be to couple it to a dark photon  $\chi X_{\mu\nu} \tilde{X}^{\mu\nu} / f_X$ .

We will assume the inflation to happen in the background of the relaxation process. One important advantage of this is the absence of  $\chi$  locking (see below), which requires a sizable Hubble scale. Another advantage is that inflation continues after the relaxation and washes away all its byproducts, such as thermal plasma, which simplifies the phenomenology. We are now ready to consider the Higgs mass scan in detail and list the conditions needed for the described mechanism to work.

1. Initially EW gauge bosons are too heavy to be produced. The particle friction turns on when the Higgs

field approaches the SM value <sup>3</sup>

$$\frac{\dot{\phi}_{\text{in}}}{f} \simeq 2m_W, \quad (\text{IV.17})$$

with the initial velocity  $\dot{\phi}_{\text{in}}$  being defined by the Hubble friction

$$\dot{\phi}_{\text{in}} = \frac{V'_\phi}{3H}. \quad (\text{IV.18})$$

2. As soon as the condition (IV.17) is satisfied, we need to restore the EW symmetry. This is necessary to produce an abrupt drop of  $m_W$  to allow for a quick  $\dot{\phi}$  decrease <sup>4</sup>. The same condition makes the abelian gauge boson  $B$  a mass eigenstate, allowing the associated instability to develop without creation of the magnetic mass [13]. This drop can be produced immediately by  $T_{\text{eff}} \sim A$ , or, a bit later by plasma if it forms. We find that once the condition (IV.17) is fulfilled, the  $A$  contribution dominates over the others, therefore we require

$$\langle A^2 \rangle \simeq \frac{\langle F\tilde{F} \rangle}{H^2} \simeq \frac{V'_\phi f}{H^2} \gtrsim v^2/g^2 \quad (\text{IV.19})$$

3. The resulting  $\dot{\phi}$  drop should be at least  $\sim v^2/\Lambda^2$ . From Eq. (IV.11) we obtain

$$\frac{\dot{\phi}_{(\text{fin})}}{\dot{\phi}_{(\text{in})}} \lesssim \frac{v^2}{\Lambda^2} \Rightarrow \frac{fH (\max[T_p, H]/H)^{2/3}}{V'_\phi/3H} \lesssim \frac{v^2}{\Lambda^2} \quad (\text{IV.20})$$

4. In the absence of plasma, the equilibrium speed  $\dot{\phi} \simeq fH$  is reached, otherwise  $\dot{\phi}$  can be higher, as the plasma decreases the friction efficiency. The fermion plasma can be formed in two ways, either from perturbative or non-perturbative production. Let us now consider under what conditions these channels can be active. In case if both production channels are efficient, the one which gives the higher plasma temperature will dominate. For simplicity we will assume that  $W$  gauge bosons are never exponentially produced because of the magnetic mass. In this way we will obtain a conservative estimate of the allowed parameter space.

- To allow for a perturbative production of fermions at  $v = 0$  and a plasma temperature  $T_p$  we need (see Eq. (IV.12) and (VI.2))

$$\exp \left[ \frac{8}{\pi^2} \frac{H^4}{m_{\text{D}}^4} \xi^6 \right] \gg 4\pi^2 \frac{H^{4/3}}{m_{\text{D}}^{4/3}} \frac{128\pi}{N_f g^4} \quad (\text{IV.21})$$

where  $N_f \sim \mathcal{O}(100)$  counts the number of SM fermionic degrees of freedom,  $m_{\text{D}}^2 = g^2 T_p^2/6$  and  $\xi \equiv \frac{\dot{\phi}}{fH}$  is defined from the balance  $V'_\phi \simeq F\tilde{F}/f$  as (for  $F\tilde{F}$  see Eq. (VI.2))

$$\exp \left[ \frac{8}{\pi^2} \frac{H^4}{m_{\text{D}}^4} \xi^6 \right] \simeq \frac{V'_\phi f}{\frac{1}{2\pi^3} (H/m_{\text{D}})^2 H^4 \xi^7} \simeq \frac{2\pi^3 V'_\phi f}{H^{4/3} m_{\text{D}}^{8/3}} \quad (\text{IV.22})$$

<sup>3</sup> If the friction switches on earlier, the  $\dot{\phi}$  will slowly decrease, tracking  $m_W$  evolution, until  $m_W$  gets to  $H$ , where the equilibrium  $\dot{\phi} \sim fH$  is reached. Generically,  $\dot{\phi}$  will remain at  $fH$  for roughly the same amount of time as it had in the  $m_W$  tracking regime. During this time the Higgs mass will still be scanned. To make the residual Higgs mass change negligible,  $H$  is forced to be very small and, consequently,  $A^2$  to be too high, washing away the sensitivity to  $h \sim 10^2$  GeV.

<sup>4</sup> Otherwise  $\dot{\phi}$  will track the decreasing  $m_W$  until it falls below  $H$ , resulting in  $v \sim H$  and  $\dot{\phi} \sim vf$ . So, in order to produce a large  $\dot{\phi}$  drop, the particle friction will have to start at  $\dot{\phi} \gg vf$ , which can not work, see previous footnote.

Thus, comparing the expressions (IV.21) and (IV.22) we conclude that the plasma temperature can not grow above

$$T_{\max}^{4/3} \simeq \frac{V'_\phi f}{H^{8/3}} \frac{6^{2/3} N_f g^{8/3}}{256} \quad (\text{IV.23})$$

as the plasma will stop being produced otherwise. At the same time plasma can not go above the equilibrium temperature, defined such that order-one fraction of all the gauge field energy gets thermalized. The equilibrium plasma temperature can be estimated as the total energy density lost by the rolling  $\phi$  field in one Hubble time  $T_p^4 \sim \delta\rho \sim V'_\phi \dot{\phi}/H$

$$T_{\text{eq}}^4 \simeq \frac{m_{\text{D}}^{2/3}}{H^{2/3}} (V'_\phi f). \quad (\text{IV.24})$$

- For the Schwinger production, the fermion plasma is produced if  $gE > m_f^2$  ( $m_f = 0$  in the unbroken phase) and (see Eq. (VI.3) for  $E$ )

$$(gE)^2 \gg 4\pi^3 H^4 \Rightarrow g^2 \frac{H^{2/3}}{m_{\text{D}}^{2/3}} (V'_\phi f) \gg 4\pi^3 H^4 \quad (\text{IV.25})$$

which defines the maximal plasma temperature above which the production stops

$$T_{\max}^{2/3} \simeq \frac{6^{1/3} g^{4/3} V'_\phi f}{4\pi^3 H^{10/3}} \quad (\text{IV.26})$$

At the same time the plasma temperature is also limited by the energy stored in the electric field. We can find this equilibrium temperature from  $T_p^2 \simeq E$

$$T_{\text{eq}}^4 \simeq \frac{H^{2/3}}{m_{\text{D}}^{2/3}} (V'_\phi f) \quad (\text{IV.27})$$

- Presence of thermal back-reaction on gauge bosons and the Higgs field in both cases requires  $g^4 T_p \gtrsim (4\pi)^2 H$ .
- Finally, the gauge boson effects should maintain unbroken EW symmetry. Hence either  $\min[T_{\max}, T_{\text{eq}}] > v$  is satisfied for at least one of the production modes, or the  $A^2$  (Eq. (VI.5)) correction to the Higgs mass has to remain large

$$A^2 > v^2/g^2 \Rightarrow \frac{H^{2/3}}{m_{\text{D}}^{2/3}} \frac{V'_\phi f}{H^2} > v^2/g^2 \quad (\text{IV.28})$$

It is easy to check that this is always the case if the other constraints are satisfied.

5. The scanning window exists long enough to scan the entire Higgs mass range. The typical time  $\phi$  needs for it is

$$t_\phi \sim \frac{\Lambda/\kappa_h}{\dot{\phi}} \sim \frac{H^2}{\kappa_\phi \kappa_h \Lambda^2} \frac{1}{H}, \quad (\text{IV.29})$$

while the time it takes  $\chi$  to pass its natural field range is

$$t_\chi \simeq \frac{\delta\chi}{\dot{\chi}} \simeq \left( \frac{H}{\kappa_\chi \Lambda} \right)^2 \frac{1}{H} \quad \text{or} \quad \left( \frac{\Lambda}{\kappa_\chi f_X} \right) \frac{1}{H}. \quad (\text{IV.30})$$

We therefore arrive at

$$\kappa_\chi \sim (\kappa_\phi \kappa_h)^{1/2} \quad \text{or} \quad \kappa_\chi \sim \kappa_\phi \kappa_h \frac{\Lambda^3}{f H^2} \quad (\text{IV.31})$$

depending on which type of friction we choose for it. From the expression (IV.29) we can also derive the constraint on the duration of inflation. The minimal required number of e-folds is given by the pre-factor of  $1/H$ .

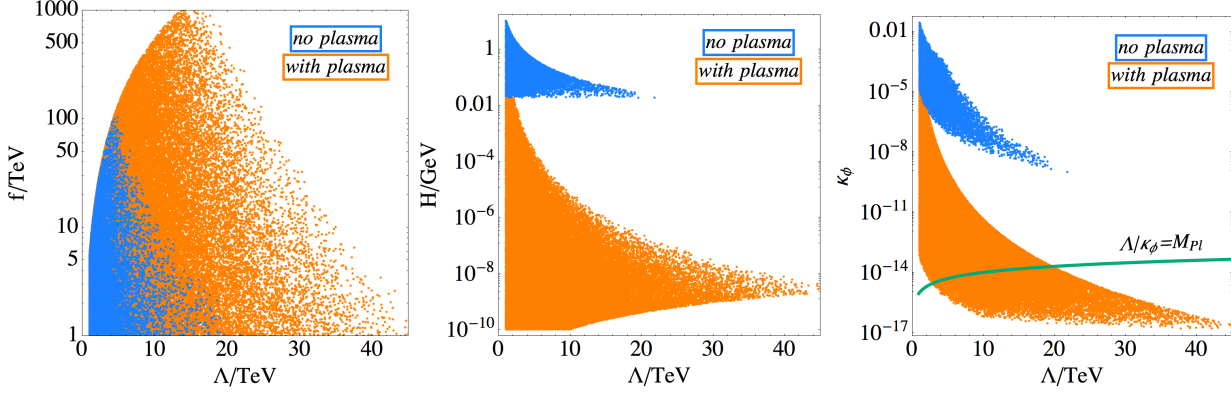


FIG. 4: Results of the numerical parameter space scan, in terms of  $\Lambda$ ,  $f$ ,  $H$  and  $\kappa_\phi$ . Blue points correspond to evolution without plasma production, orange – with plasma production. The green line on the right plot shows minimal  $\kappa_\phi$  below which the  $\phi$  field excursion becomes trans-Planckian.

6. The velocity-dependent contributions to the  $\chi$  potential will tend to push it to the poles, ruining its monotonic decrease (if we assume two poles as in Sec. III D), and significantly decreasing the scan window (if  $\chi$  manages to start moving towards zero). To prevent  $\chi$  locking we need the term  $\frac{\Lambda^2}{\chi^2}\dot{\phi}^2$  to be subdominant compared to  $\kappa_\chi\Lambda^3\chi$ .

The maximal size of the former is reached for  $\dot{\phi} \simeq \kappa_\phi\Lambda^3/3H$ . Therefore we need

$$\frac{\partial_\chi(\frac{\Lambda^2}{\chi^2}\dot{\phi}^2)}{\partial_\chi(\kappa_\chi\Lambda^3\chi)} = \frac{\kappa_\phi^2\Lambda^2}{9H^2} \ll 1 \quad (\text{IV.32})$$

where we took  $\chi \sim \Lambda_k \sim \Lambda/\kappa_\chi$ .

7. The Hubble expansion-induced field jumps are either absent or negligible compared to classical field evolution

$$v > H, \quad (\text{IV.33})$$

$$\delta\phi_{1H} \sim \dot{\phi}/H > H \Rightarrow f\max[T_p, H]^{2/3} > H^{5/3}, \quad (\text{IV.34})$$

$$\delta\chi_{1H} \sim \dot{\chi}/H > H \Rightarrow \kappa_\chi\Lambda^3 > H^3 \text{ or } f_X > H, \quad (\text{IV.35})$$

where  $1H$  subscripts stand for a classical displacement in one Hubble time.

8. Finally we have the constraint on the vacuum energy density, ensuring that inflation is not affected by our mechanism

$$H^2 M_{\text{Pl}}^2 > \Lambda^4 \quad (\text{IV.36})$$

The main result of this section is the scatter plots on Fig. 4, showing the values of the main relevant parameters which satisfy the listed above constraints. For all the points with the plasma production the main production channel is the perturbative production. The maximal allowed cutoff scale  $\Lambda$  is around 50 TeV, while without trans-Planckian field excursions it decreases to  $\sim 20$  TeV. The friction without plasma is generated for a moderate (for this kind of scenarios) Hubble scale  $\mathcal{O}(0.1)$  GeV, while in the presence of plasma  $H$  needs to be several orders of magnitude lower, with the maximal cutoff reached for  $H \sim 10^{-8}$  GeV.



## V. SUMMARY AND FUTURE DIRECTIONS

We have discussed a new scenario with the dynamical Higgs mass scan, within a class of models pioneered in [20, 21] and [1]. The novelty of our proposal lies in employing a non-canonical kinetic term of the scanning field, leading to its decoupling from the Higgs sector in the end of the scan. This scenario therefore does not require Higgs-dependent barriers for the relaxion (and hence new light EW-charged states), neither, more generally, the presence of two periods in the relaxion potential.

We have presented a particular realization of the pole attractor idea, in which the scanning field can evolve only during a limited time window during inflation. After that the second field, controlling the kinetic terms, reaches the pole value and blocks the scan. Starting with a large Higgs  $v_{ev}$ , the scanning field first evolves quickly until the Higgs mass gets close to the SM value. At this point the scanning field evolution is slowed down by a dissipation of energy into SM gauge bosons, hence it displaces by a very small distance before the scanning window closes. We have identified two viable regions of the parameter space, with and without a back-reacting plasma. The latter allows for a higher inflationary Hubble parameter, but a slightly lower maximal cutoff. The maximal cutoff of the Higgs sector  $\sim 50$  TeV is well above the reach of current and near future collider experiments. It is nevertheless not restricted to be that far and can reside as low as the current lower bounds on new physics allow. The relaxion sector fields couplings to the visible sector are exponentially suppressed making it very difficult to observe their direct effect.

All the models discussed in this work need the initial Higgs  $v_{ev}$  to be large, or in other words the Higgs mass squared to be negative. One can however think of modifications suitable to accommodate a vanishing initial Higgs field value  $h = 0$  and a large positive mass squared. These should include a change of signs in the scalar potential to provide a scan in the right direction. For the toy model, in order to prevent having a singularity in the relaxion kinetic term one would need to shift the pole away from  $h = 0$  as discussed in Sec. II A, which in any case is needed to stop the scan at  $h = v \neq 0$ . For the two field attractor one can think of assigning the Higgs-dependent particle production friction to  $\chi$  instead of  $\phi$ , so that initially  $\chi$  is slowed down by SM gauge boson production at  $v = 0$ , and then the particle friction disappears and  $\chi$  falls to zero, blocking the scan. The simplest implementation of this last mechanism however leads to a relatively low cutoff, at the TeV scale, so further refinements would be needed in order to increase it.

We would also like to mention that there is another potential way of implementing the two-field scenario, without using the concept of limited scanning window and changing friction. This would be more in accord to the spirit of the toy model, where the pole depends on the higgs vev. One can assign to  $\chi$  a potential with a minimum (local or global) away from zero, which disappears or displaces to zero when the Higgs mass reaches the SM value. The simplest realization would be a scalar potential with a tadpole  $V_\chi = h^2\chi + \Lambda^2\chi^2$ , which however has the problem that the quantum correction  $\sim \Lambda^4 \log \chi$  quickly becomes more important than the Higgs-dependent tadpole. It would be interesting to further investigate this model building direction.

The proposed mechanism has a lot in common with the attractor models of inflation, sharing some structural features and, possibly, can find UV completions in analogues type of theories. To this end, the existence of an appropriate UV completion for our scenario presents an important question for further studies. In particular because the behaviour close the attractor relies on certain crucial assumptions about the UV physics features.

Finally, we would like to mention that a distinctive phenomenological feature of this type of scenarios is a slow change of the theory parameters with time, as the scan never completely stops. While we do not necessarily need this time variation to be detectable with the current experiments, its observation would be an interesting hint for this mechanism, especially given the limited direct experimental access to the relaxion sector, whose couplings to the SM particles are exponentially suppressed.

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## VI. APPENDIX: PARTICLE PRODUCTION

The combinations of the field strengths can be rewritten in terms of electric and magnetic fields as [25]

$$\langle FF \rangle \sim \langle E^2 \rangle + \langle B^2 \rangle, \quad \langle F\tilde{F} \rangle \sim \langle EB \rangle. \quad (\text{VI.1})$$

At zero temperature, they take the form [25]

$$\langle EB \rangle \simeq 10^{-4} H^4 \xi^{-3} \exp[2\pi\xi], \quad (\text{VI.2})$$

$$\langle E^2 \rangle \simeq 10^{-4} H^4 \xi^{-3} \exp[2\pi\xi] \simeq \langle EB \rangle, \quad (\text{VI.3})$$

$$\langle B^2 \rangle \simeq 10^{-4} H^4 \xi^{-5} \exp[2\pi\xi] \simeq \xi^{-2} \langle EB \rangle, \quad (\text{VI.4})$$

where we have defined  $\xi \equiv \frac{\phi}{fH}$ . Using the results of [25] we can also estimate

$$\langle A^2 \rangle \sim \langle EB \rangle / H^2. \quad (\text{VI.5})$$

And the number of gauge field quanta in one Hubble volume as the total energy density  $\sim k^2 A^2$  over the energy per particle  $\Omega$  [26]

$$N_\gamma \simeq 10^{-4} \exp[4.5\xi]. \quad (\text{VI.6})$$

In plasma we have [15]

$$\langle EB \rangle \simeq \frac{1}{2\pi^3} \frac{H^2}{m_D^2} H^4 \xi^7 \exp\left[\frac{8}{\pi^2} \frac{H^4}{m_D^4} \xi^6\right], \quad (\text{VI.7})$$

$$\langle E^2 \rangle \simeq \frac{1}{\pi^4} \frac{H^4}{m_D^4} H^4 \xi^9 \exp\left[\frac{8}{\pi^2} \frac{H^4}{m_D^4} \xi^6\right] \simeq \frac{H^2}{m_D^2} \xi^2 \langle EB \rangle \simeq \frac{H^{2/3}}{m_D^{2/3}} \langle EB \rangle, \quad (\text{VI.8})$$

$$\langle B^2 \rangle \simeq \frac{1}{4\pi^2} H^4 \xi^5 \exp\left[\frac{8}{\pi^2} \frac{H^4}{m_D^4} \xi^6\right] \simeq \frac{m_D^2}{H^2} \xi^{-2} \langle EB \rangle \simeq \frac{m_D^{2/3}}{H^{2/3}} \langle EB \rangle, \quad (\text{VI.9})$$

where in the last column we used  $\xi \sim (m_D/H)^{2/3}$ , which is needed for the gauge field back-reaction on  $\phi$  to be sizeable. We can also derive from the results of [15]

$$\langle A^2 \rangle \sim \frac{1}{4\pi^2} \xi^3 H^2 \exp\left[\frac{8}{\pi^2} \frac{H^4}{m_D^4} \xi^6\right] \sim \frac{H^{2/3}}{m_D^{2/3}} \langle EB \rangle / H^2, \quad (\text{VI.10})$$

and estimate the particle number

$$N_\gamma \sim \frac{1}{4\pi^2} \xi^2 \exp\left[\frac{8}{\pi^2} \frac{H^4}{m_D^4} \xi^6\right] \sim \frac{1}{4\pi^2} \frac{\langle EB \rangle}{H^4} \frac{H^{4/3}}{m_D^{4/3}}. \quad (\text{VI.11})$$

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