Cosmological Higgs-Axion Interplay for a Naturally Small Electroweak Scale

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Recently, a new mechanism to generate a naturally small electroweak scale has been proposed. It exploits the coupling of the Higgs boson to an axionlike field and a long era in the early Universe where the axion unchains a dynamical screening of the Higgs mass. We present a new realization of this idea with the new feature that it leaves no sign of new physics at the electroweak scale, and up to a rather large scale, 10^9 GeV, except for two very light and weakly coupled axionlike states. One of the scalars can be a viable dark matter candidate. Such a cosmological Higgs-axion interplay could be tested with a number of experimental strategies.

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Introduction.—Our understanding of nature is based on the empirical evidence that natural phenomena taking place at very different energy or distance scales do not influence each other. The parameters of an effective theory are *natural* if they do not require any special tuning of the parameters of the theory at higher energies. Wilson and Kogut [1] and 't Hooft [2] gave a quantitative meaning to this naturalness principle by demanding that all dimensionless parameters controlling the different effective theories should be of order unity unless they are associated to the breaking of a symmetry. The Higgs boson mass and the value of the cosmological constant have been long recognized as two notorious challengers of this naturalness principle. Supersymmetry or Higgs compositeness are two prime examples of models trying to associate the Higgs mass to a small symmetry breaking. Recently, however, a radically new approach to explain the smallness of the Higgs mass has been proposed [3], reminiscent of the relaxation mechanism of Ref. [4] proposed for explaining the smallness of the cosmological constant (see Refs. [5,6] for similar previous ideas). Technically, the relaxation mechanism of Ref. [3] is based on the cosmological interplay between the Higgs field h and an axionlike field ϕ, arising from the following three terms of the scalar effective potential:

$$V = \Lambda^3 g\phi - (\Lambda^2 - g\Lambda\phi)\frac{h^2}{2} + \epsilon\frac{h^n}{\Lambda_c^{n-4}}\cos(\phi/f) + \cdots, \tag{1}$$

where Λ is the UV cutoff scale of the model, while $\Lambda_c \lesssim \Lambda$ is the scale at which the periodic $\cos(\phi/f)$ term originates, and n is a positive integer. The first term is needed to force ϕ to roll down in time, while the second one corresponds to a Higgs mass-squared term with a (positive) dependence on ϕ such that different values of ϕ scan the Higgs mass over a large range. Finally, the third term plays the role of a potential barrier for ϕ, dependent on h, necessary to stop the rolling of ϕ once electroweak symmetry breaking (EWSB) occurs.

At the classical level, the proposed mechanism can be understood in the following way. Assuming that ϕ starts, at the beginning of the inflationary epoch, at a very large value $\phi \gtrsim \Lambda/g$, it will slow roll until it takes the critical value $\phi_c = \Lambda/g$, at which the Higgs mass-squared term becomes zero. From this time on, the Higgs mass becomes negative, and it is energetically favored to turn on the Higgs field. This raises the third term of Eq. (1) up to the point at which ϕ stops rolling. For $g \ll 1$, this occurs for a Higgs value v obeying $g\Lambda^3 \simeq \epsilon\Lambda_c^{4-n} v^n/f$. We can have $v \ll \Lambda$ by taking g small enough, which is technically natural as g defines the spurion that breaks the symmetry $\phi \to \phi + 2\pi f$. Therefore, this mechanism potentially offers a new solution to the hierarchy problem. We will refer to this as the cosmological Higgs-axion interplay (CHAIN) mechanism.

For $n = 1$, the third term of Eq. (1) is linear in h, implying that $\epsilon\Lambda_c^3$ must arise from a source of EWSB other than the Higgs field. This can be the QCD quark condensate $\langle q\bar{q}\rangle \sim \Lambda_{\rm QCD}^3$, as proposed in Ref. [3]. In this case, $\Lambda_c \sim \Lambda_{\rm QCD}$ and $\epsilon \sim y_u$, where y_u is the up-quark Yukawa coupling. This model, however, predicts too large a value for the QCD θ angle, in conflict with neutron electric dipole moment constraints. A possible way to fix this problem was explained in Ref. [3], but it requires a low cutoff scale, $\Lambda \lesssim 30$–1000 TeV.

For $n = 2$ on the contrary, the $h^2\cos(\phi/f)$ term in $V(\phi, h)$ can arise from the electroweak-invariant term $|H|^2\cos(\phi/f)$, where H is the Higgs doublet, and therefore no extra source of EWSB is needed beyond the standard model (SM) Higgs field. Nevertheless, at the quantum level, the term $\epsilon\Lambda_c^4\cos(\phi/f)$ can now be induced (just by closing H in a loop), which could stop the ϕ evolution much before the Higgs field turns on. Therefore, if we want the CHAIN mechanism to work, we must demand $\Lambda_c \sim v$ that implies new physics not far away from the weak scale and thus introduce a "coincidence problem." It is important to notice that this new physics would not be responsible for keeping the Higgs boson light, unlike in all other efforts to explain the hierarchy problem, but for generating the periodic term of Eq. (1).

The aim of his Letter is to offer an existence proof that it is indeed possible to devise a model that dynamically generates a large mass gap between the Higgs mass and the new-physics threshold. The proposed model will not have a coincidence problem as the only new-physics scale will be associated with $\Lambda \sim \Lambda_c \gg v$. For this to work, we need to make the terms, like $\epsilon\Lambda_c^4 \cos(\phi/f)$ or $\epsilon\Lambda_c^3 g\phi\cos(\phi/f)$ obtained by closing a Higgs loop, smaller than the term $\epsilon\Lambda^2|H|^2\cos(\phi/f)$. For this purpose, another slow-rolling field, σ is introduced and coupled to $\cos(\phi/f)$. During its cosmological evolution, σ will take a value such that $\sigma\cos(\phi/f)$ will cancel the quantum generated dangerous terms. When this occurs, ϕ will be free to move, tracking σ downhill. Only when the h-dependent term turns on, ϕ will stop tracking σ. This way, the cutoff scale can be pushed up to $\Lambda \sim 10^9$ GeV. The only new states, ϕ and σ, will have masses below the weak scale, but they will be very weakly coupled to the SM, making them very difficult to detect at present and future experiments. Interestingly, as we will see, they could provide the source of Dark Matter (DM) in the Universe. Additionally, for moderate values of $\Lambda \sim 10^4$ GeV, our model improves two aspects of the mechanism in Ref. [3]: (i) the required field excursion $\Delta\phi \sim \Lambda/g$ can be even sub-Planckian, and (ii) the required number of e-folds during inflation does not need to be extremely large.

Double scanner mechanism.—The key new ingredient of our proposal, with respect to Ref. [3], is a second scanning field, that we call σ. The full potential, up to terms of order ϵ, g_σ, and g, is given by

$$V(\phi,\sigma,H) = \Lambda^3(g\phi + g_\sigma\sigma) - \Lambda^2\left(\alpha - \frac{g\phi}{\Lambda}\right)|H|^2 + \lambda|H|^4 + A(\phi,\sigma,H)\cos(\phi/f), \quad (2)$$

where

$$A(\phi,\sigma,H) \equiv \epsilon\Lambda^4\left(\beta + c_\phi\frac{g\phi}{\Lambda} - c_\sigma\frac{g_\sigma\sigma}{\Lambda} + \frac{|H|^2}{\Lambda^2}\right), \quad (3)$$

with $0 < g, g_\sigma, \epsilon \ll 1$, while α, β and c_ϕ, c_σ are $O(1)$ positive coefficients. For a partial UV completion of this model, see Ref. [7].

From Eqs. (2) and (3), we see that ϕ scans the Higgs mass, while σ scans $A(\phi,\sigma,H)$, the overall amplitude of the oscillating term. This dependence of $A(\phi,\sigma,H)$ on σ and H is crucial for our CHAIN mechanism to work, while the other terms in Eq. (3) are added since they are anyway generated at the quantum level (by loops of H). The potential in Eq. (2) is stable under quantum corrections in the small-coupling limit.

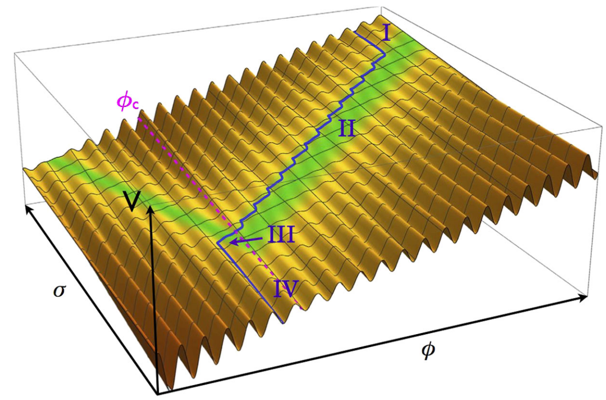

FIG. 1 (color online). Scalar potential in the $\{\phi,\sigma\}$ plane. The band without barriers is in green while the barriers getting high (er) are dark(er) brown. The blue line shows a possible slow-roll cosmological trajectory of the fields during inflation.

Inflation is needed, as in Ref. [3], to provide the friction that makes the fields slow roll and reach the desired minimum. The time evolution of σ is quite trivial, as for $\epsilon \ll 1$, it simply slides down $\sigma(t) = \sigma_0 - g_\sigma\Lambda^3 t/(3H_I)$. In the cosmological evolution of ϕ, we can distinguish four stages, depicted in Fig. 1.

(I) At the start of inflation we assume $\phi \gtrsim \Lambda/g$ and $\sigma \gtrsim \Lambda/g_\sigma$ such that the Higgs mass-squared term is positive and the amplitude $|A|$ is of order $\epsilon\Lambda^4$. The field ϕ is stuck in some deep minimum coming from the $A\cos(\phi/f)$ term of Eq. (2), while the Higgs field value is zero.

(II) As σ evolves down, the amplitude A decreases until the point where the steepness of $A\cos(\phi/f)$ along the ϕ direction is smaller than the slope coming from the linear term of Eq. (2), and ϕ can start to move down. The bumps from $A\cos(\phi/f)$ are very small and, for $g_\sigma \lesssim g$, ϕ goes down tracking σ: $\phi(t) \simeq \text{const} + c_\sigma g_\sigma\sigma(t)/(c_\phi g)$, which is the solution of $A \approx 0$.

(III) When ϕ crosses the critical value $\phi_c \equiv \alpha\Lambda/g$, the Higgs mass-squared term becomes negative, turning on H. This gives a positive contribution to the amplitude A and modifies the gradient $d\phi/d\sigma$ inside the tracking band forcing ϕ to exit.

(IV) Finally, ϕ gets stuck in a minimum from $A\cos(\phi/f)$ as in the model of Ref. [3]. The field σ continues going down, making A grow until σ finds its own minimum.

The tracking region of Fig. 1 corresponds to values of ϕ, called ϕ_*, for which the steepness from $A\cos(\phi/f)$ is smaller than the steepness from the $\Lambda^3 g\phi$ term of Eq. (2). These are determined by $A(\phi_*,\sigma,h(\phi_*)) \lesssim gf\Lambda^3$. In order for $\phi(t)$ to track down $\sigma(t)$, or what is equivalent, for $\phi(t)$ to stay in the ϕ_* region until reaching ϕ_c, the gradient of the dynamical trajectory in the $\{\phi, \sigma\}$ plane inside the ϕ_* interval, $d\phi(t)/d\sigma(t) = g/g_\sigma$, should be larger than the gradient $d\phi_*/d\sigma$ of the tracking band itself. This condition simply reads: $c_\phi g^2 > c_\sigma g_\sigma^2$. On the contrary, once $\phi \leq \phi_c$, we must demand ϕ to exit the tracking band to get trapped in some vacuum precisely as needed to explain the smallness of the electroweak scale. This is happening provided that $[c_\phi - 1/(2\lambda)]g^2 < c_\sigma g_\sigma^2$.

Let us emphasize that the CHAIN mechanism described above works independently of ϕ_i, the initial condition for ϕ, provided only that ϕ_i at the initial time t_i lies in the region $\phi_c < \phi_i < \phi_*(t_i)$, which is a natural and sizable range of the available field space. The cosmological evolution just described is purely classical. Quantum fluctations, governed by the Fokker-Planck equation, will give corrections that, however, do not spoil the solution of the hierarchy problem [7].

Consistency requirements for a small weak scale.—The cosmological evolution of our model can be broadly described by two external quantities fixed by the inflaton sector: H_I, the value of the Hubble parameter during inflation, and N_e, the number of e-folds. In order to provide a natural solution to the hierarchy problem, we require the following.

(1) *Dangerous quantum corrections to the potential are kept small*. Terms like $\epsilon^2\Lambda^4\cos^2(\phi/f)$ or $\epsilon^2\Lambda^3 g\phi\cos^2(\phi/f)$ are generated at the quantum level and their amplitudes cannot be canceled by σ simultaneously to $A\cos(\phi/f)$. They could give a barrier to ϕ at values that can be above the critical ϕ_c. To make sure that they remain subdominant to the Higgs barrier of Eq. (2), we must demand

$$\epsilon \lesssim v^2/\Lambda^2. \tag{4}$$

This condition also ensures that the contribution to the Higgs mass coming from the $\epsilon\Lambda^2|H|^2\cos(\phi/f)$ term in the potential is at most of electroweak size and does not spoil the tracking behavior.

(2) *ϕ must be trapped by the Higgs barrier.* The nonzero Higgs field must be responsible for stopping ϕ from sliding any longer. This requirement fixes the electroweak scale in terms of microscopic parameters:

$$v^2 \simeq \frac{g\Lambda f}{\epsilon}. \tag{5}$$

(3) *Inflation is independent of the ϕ and σ evolution*. The typical energy density carried by ϕ and σ should remain smaller than the inflation scale; i.e.,

$$\frac{\Lambda^2}{M_P} \lesssim H_I, \quad \text{with} \quad M_P \simeq 2.4\times 10^{18}\ \text{GeV}. \tag{6}$$

In addition, the two fields ϕ and σ should be slowly rolling during inflation, which requires $g_\sigma\Lambda, g\Lambda \lesssim H_I$.

(4) *Classical rolling dominates over quantum jumping*. During the cosmological evolution, the quantum fluctuations of the fields, typically of size H_I, should remain smaller than the classical field displacements over one Hubble time; i.e., for the case of σ [3],

$$H_I^3 \lesssim g_\sigma\Lambda^3. \tag{7}$$

(5) *Inflation lasts long enough for complete scanning*. The range scanned by ϕ and σ during inflation must be of order or larger than Λ/g and Λ/g_σ, respectively. This is ensured by requiring a long enough period of inflation, namely,

$$N_e \gtrsim \frac{H_I^2}{g_\sigma^2\Lambda^2}. \tag{8}$$

Combining these various consistency conditions, we obtain that the couplings g_σ and g are bounded to the interval $\Lambda^3/M_P^3 \lesssim g_\sigma \lesssim g \lesssim v^4/(f\Lambda^3)$. Since f cannot be much smaller than Λ, as this latter is the scale at which the $\cos(\phi/f)$ term is generated, we obtain an upper bound on the cutoff of our model $\Lambda \lesssim (v^4 M_P^3)^{1/7} \simeq 2\times 10^9$ GeV.

Cosmological signatures.—The new-physics scale of our model can be as large as $\Lambda \sim 10^9$ GeV, and, therefore, we do not expect any new state around the weak scale. Only the two additional scalars σ and ϕ are lighter than, or at most around, the weak scale. They are very weakly coupled to the SM states and have some phenomenological impact through astrophysical and cosmological effects only.

(1) *Properties of ϕ and σ.*—After the slow-rolling process ends and σ settles in a minimum, no cancellation is expected in the $A(\phi,\sigma,H)$ amplitude, so that $A(\phi,\sigma,H)\sim\epsilon\Lambda^4$. The mass of ϕ is thus controlled by $A\cos(\phi/f)$ and can be estimated as

$$m_\phi^2 \sim \frac{\epsilon\Lambda^4}{f^2} \sim g\frac{\Lambda^5}{fv^2} \lesssim v^2. \tag{9}$$

For σ we expect that higher-order terms in $g_\sigma\sigma/\Lambda$, not shown for simplicity in Eq. (2), give it a mass of order

$$m_\sigma^2 \sim g_\sigma^2\Lambda^2 \ll m_\phi^2. \tag{10}$$

Contours of constant m_ϕ and m_σ are shown in Fig. 2.

These two scalars interact with the SM particles mainly through a mass mixing with the Higgs boson. The corresponding mixing angles can be estimated as

$$\theta_{\phi h} \sim \frac{g\Lambda v}{m_h^2}, \qquad \theta_{\sigma\phi} \sim \frac{g_\sigma f v^2}{\Lambda^3},$$
$$\theta_{\sigma h} \sim \max\left\{\theta_{\sigma\phi}\theta_{\phi h}, \frac{g^2}{16\pi^2}\frac{g_\sigma\Lambda^7}{f^2v^3m_h^2}\right\}. \tag{11}$$

The first contribution in $\theta_{\sigma h}$ arises at tree level, whereas the second one originates from a ϕ loop. The scalar potential Eq. (2) also gives rise to interactions between ϕ and the Higgs boson, not suppressed by the small mixing angle $\theta_{\phi h}$, that are of order

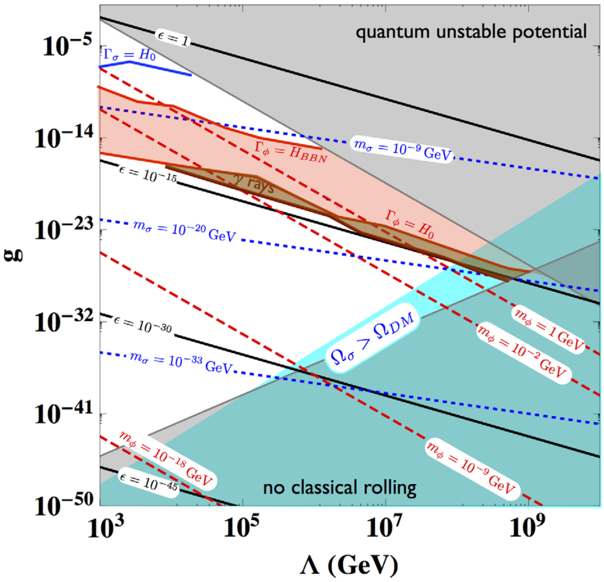

FIG. 2 (color online). Parameter space for a successful solution of the hierarchy problem ensured by the cosmological evolution of the fields ϕ and σ. We have taken $\Lambda = f$ and $g_\sigma/g = 0.1$.

$$\phi\phi hh\colon\ \epsilon\Lambda^2/f^2, \qquad \phi\phi h\colon\ \epsilon v\Lambda^2/f^2, \tag{12}$$

and will play an important role in the thermal production of ϕ. The decays of ϕ and σ are mediated by the mixing with the Higgs boson, and thus the widths are given by

$$\Gamma_\phi \sim \theta^2_{\phi h}\Gamma_h(m_\phi), \qquad \Gamma_\sigma \sim \theta^2_{\sigma h}\Gamma_h(m_\sigma), \tag{13}$$

where $\Gamma_h(m_i)$ is the SM Higgs boson width evaluated at $m_h = m_i$ [8]. As shown in Fig. 2, there is a sizable part of the parameter space in which ϕ is cosmologically unstable ($\Gamma_\phi > H_0$, where H_0 is the present Hubble value) but sufficiently long-lived to decay after big bang nucleosynthesis (BBN) [$\Gamma_\phi < H_{\rm BBN} \equiv H(T = 1\ {\rm MeV})$]. This region of the parameter space can be constrained by cosmology. On the other hand, σ is cosmologically stable in most of the relevant parameter space, and can decay within the age of the Universe only in a small corner of the parameter space.

(2) *Abundances of ϕ and σ from vacuum misalignment.*—If after inflation and reheating (assumed to be taking place at temperatures above the EWSB scale) the fields ϕ and σ end up displaced from their minima, they will fall towards them, oscillating around them if their lifetimes are large. The energy density stored in the field oscillations behaves like cold DM and can potentially overclose the Universe today or dissociate light elements if the decay takes place during or after BBN. We expect that during inflation σ slowly rolled down to its global minimum, somewhere in its $\sim\Lambda/g_\sigma$ range, as this requires a number of e-folds similar to the N_e estimated in Eq. (8). Because of quantum effects, σ reached the minimum with a spread $\sqrt{N_e}H_I$. The typical displacement from the minimum at the end of inflation is, therefore, $(\Delta\sigma)_{\rm ini} \sim \sqrt{N_e}H_I$, corresponding to an energy density of the order $\rho^\sigma_{\rm ini} \sim m^2_\sigma(\Delta\sigma)^2_{\rm ini} \sim H^4_I$. The energy density stored in σ oscillations today, relative to the critical energy density, is then $\Omega_\sigma \gtrsim (H_I M_P/\Lambda^2)^4(4\times 10^{-28}/g_\sigma)^{3/2} \times (\Lambda/10^8\ {\rm GeV})^{13/2}$. The bound to avoid Universe overclosure translates into a lower bound for g_σ as a function of Λ. It is shown in Fig. 2 in the case $H_I = \Lambda^2/M_P$. It is interesting that σ can be a good DM candidate in certain regions of the allowed parameter space, in particular, at large Λ. For certain values of m_σ, there can be other cosmological constraints. For example, for $\Omega_\sigma \gtrsim \Omega_{\rm DM}/20$, the mass range 10^{-32} eV $\lesssim m_\sigma \lesssim 10^{-25.5}$ eV is excluded by structure formation [9], while masses around $m_\sigma \sim 10^{-11}$ eV may be constrained by black hole superradiance [10]. Interestingly, for the particular case $m_\sigma \sim 10^{-23}$ eV, σ could be searched for by the SKA pulsar timing array experiment [11]. There are ways to go around these bounds, for instance, by assuming a late entropy production after σ has started to oscillate, as can occur if reheating is a very slow process such that $T_{RH} < T^\sigma_{\rm osc}$ (with $T^i_{\rm osc} \sim \sqrt{m_i M_P}$) [12].

For ϕ, the initial energy density arising from its displacement due to quantum spreading was at most $\rho^\phi_{\rm ini} \sim H^4_I$, that, since $m_\phi \gg m_\sigma$ and then $T^\phi_{\rm osc} \gg T^\sigma_{\rm osc}$, gives today a completely negligible effect.

(3) *Thermal production of ϕ.*—Thermal production of ϕ arises mainly from the $\phi\phi hh$ coupling of Eq. (12), leading to double production from the thermal bath via $h + h \to \phi + \phi$. At $T \gtrsim m_h$, this double production cross section is estimated to be $\langle\sigma_A {\rm v}\rangle \sim \epsilon^2(\Lambda^4/f^4)/T^2$. This implies that ϕ can reach thermal equilibrium only for T in the interval $[m_h, \epsilon^2 M_P(\Lambda/f)^4]$, in which the ϕ production rate is faster than the rate of expansion. This region corresponds roughly to the area above the $\Gamma_\phi = H_{\rm BBN}$ line of Fig. 2, so we conclude that in most of the parameter space ϕ never thermalizes.

The number density of ϕ produced thermally is $Y_\phi(T) \sim 10^{-4}\epsilon^2\Lambda^4 M_P/(f^4 T)$, where $Y_\phi = n_\phi/s$ and s is the entropy per comoving volume. The ϕ production is maximal at $T \sim m_h$. In the parameter region where ϕ is cosmologically stable, the contribution of ϕ to DM today is $\Omega_\phi \sim m_\phi Y_\phi s_0/\rho_c$ (s_0 is the present entropy density) and it varies from $\Omega_\phi \lesssim 10^{-4}$ along the line $\Gamma_\phi = H_0$ to $\Omega_\phi \lesssim 10^{-10}$ for $\Gamma_\phi \simeq 10^{-10}H_0$.

(4) *Constraints from BBN and gamma-ray observations.*—There is a region of parameter space in which ϕ is not cosmologically stable and decays after BBN. This is problematic if the decay of ϕ injects into the thermal bath an energy per baryon $E_{p.b} \gtrsim O$ (MeV), leading to a modification of the predictions for the abundances of the light elements. Since $E_{p.b} \sim m_\phi Y_\phi n_\gamma/n_b$, this results in the bound $m_\phi Y_\phi \lesssim 10^{-12}$ GeV, which, however, could be weakened sensitively depending on the precise value of the lifetime [13]. In addition, the cosmic microwave background (CMB) constrains lifetimes $\sim[10^{10}\text{–}10^{13}]$ s

for $E_{p.b}$ down to O (eV). Therefore, we expect that most of the region of the parameter space delimited by the lines $\Gamma_\phi = H_{\rm BBN}$ and $\Gamma_\phi = H_0$ in Fig. 2 is excluded. A dedicated analysis would be needed to derive the precise excluded regions, but it is beyond the scope of this Letter.

On the other hand, for regions in which the ϕ lifetime is larger than the age of the Universe, there are strong constraints coming from decays generating a distortion in the galactic and extragalactic diffuse x-ray or gamma-ray background. In particular, sub-GeV DM decaying into photons should satisfy $\tau_{\rm DM} \gtrsim 10^{27}$ s [14]. Since the gamma-ray flux scales as $d\Phi_\gamma/dE \propto Y_\phi \Gamma_\phi$, we can translate this bound into $\tau_\phi > 10^{27}$ s $\times\, \Omega_\phi/\Omega_{\rm DM}$. This excludes the thin brown band of Fig. 2.

We stress that the cosmological constraints derived above can be evaded if the temperature of the Universe never reaches m_h, in which case the thermal production of ϕ is suppressed.

Conclusion.—We provided an existence proof of a model, based on the idea of Ref. [3], that can naturally accommodate a small electroweak scale without requiring visible new physics at present and far future colliders. The model is based on the cosmological evolution of the Higgs field and two axionlike states whose backreactions lead to a naturally small electroweak scale. The only new physics of the model consists of these two scalars, ϕ and σ, that in most of the parameter space are very light and very weakly coupled to the SM. Therefore, strategies to detect them are completely different, as they do not require powerful high-energy colliders but dedicated searches in the sub-GeV regime.

Interestingly, the lightest state, σ, could be a DM candidate. The field ϕ cannot contribute to more than $\Omega_\phi \lesssim 10^{-10}$. For this maximum value, it may still be detected in gamma-ray observations from its late decay.

Part of the parameter space of our model can be tested through observations of the diffuse gamma-ray backgrounds, black hole superradiance, and even in pulsar timing arrays. In addition, there is a rather rich BBN and CMB phenomenology which motivates a more thorough study. Unfortunately, fifth-force signals and equivalence principle violations in intermediate mass ranges seem too small to be seen in the near future.

The ideas proposed in Ref. [3] and pursued here represent a new twist in the long and fruitful history of the interplay between particle physics and cosmology. While in the past particle physics has been a crucial ingredient to understand the cosmological history of our Universe, if these new ideas were correct, cosmological evolution would be a key ingredient in the understanding of some key parameters of particle physics.

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