Dynamical generation of the weak and dark matter scale

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Assuming that naturalness should be modified by ignoring quadratic divergences, we propose a simple extension of the standard model where the weak scale is dynamically generated together with an automatically stable vector. Identifying it as thermal dark matter, the model has one free parameter. It predicts one extra scalar, detectable at colliders, which triggers a first-order dark/electroweak cosmological phase transition with production of gravitational waves. Vacuum stability holds up to the Planck scale.

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I. INTRODUCTION

The discovery of the standard model (SM) scalar [1] together with negative results of searches for supersymmetry and for other solutions to the usual hierarchy problem [2] invite us to explore the idea that this paradigm needs to be abandoned or reformulated. One possibility is that naturalness still holds but in a modified version, namely under the assumption that the unknown cutoff at Planckian scales has the property that quadratic divergences vanish, like in dimensional regularization. Such modified "finite naturalness" was discussed in [3], showing that it is satisfied by the standard model and that new physics models motivated by data (about neutrino masses, dark matter (DM), QCD θ problem, inflation) can satisfy it.

In this paper we address what could be the dynamical origin of the weak scale and of the dark matter scale in the context of finite naturalness. Although not logically necessary, the extra hypothesis that mass terms are absent from the fundamental Lagrangian may be conceptually more appealing than just putting by hand small masses of the order of the electroweak scale. We adopt the following three guidelines:

(1) We assume that the SM and DM particles have no mass terms in the fundamental Lagrangian, and that their masses arise from some dynamical mechanism.

(2) We assume that the extended theory has the same automatic properties of the SM supported by data: accidental conservation of lepton and baryon number, of lepton flavor, etc.

(3) We assume that DM stability is one more automatic consequence of the theory.

The resulting model is presented in Sec. II and its phenomenology is explored in Sec. III. In Sec. IV we conclude.

II. THE MODEL

One simple model is obtained by merging previously proposed ideas that possess some of the properties 1, 2, 3: Refs. [4,5] for dynamical generation of the weak scale (see also [6] for related ideas), and Refs. [7,8] for automatic (accidental) DM stability.

The model has gauge group $U(1)_Y \otimes SU(2)_L \otimes SU(3)_c \otimes SU(2)_X$, namely the SM gauge group with an extra $SU(2)_X$. The field content is just given by the SM fields [singlets under $SU(2)_X$] plus a scalar S, doublet under the extra $SU(2)_X$ and neutral under the SM gauge group. The Lagrangian of the model is just the most general one, omitting the mass terms for the SM scalar doublet H ("Higgs" for short) and for the scalar doublet S, because we want to dynamically generate the weak and DM scales. Consequently, the scalar potential of the theory is

$$V = \lambda_H |H|^4 - \lambda_{HS} |HS|^2 + \lambda_S |S|^4. \qquad (1)$$

We now show how it can lead to dynamical symmetry breaking down to $U(1)_{\rm em} \otimes SU(3)_c$ such that, in unitary gauge, the scalar doublets H and S can be expanded in components h and s as

$$H(x) = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ v + h(x) \end{pmatrix}, \qquad S(x) = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ w + s(x) \end{pmatrix}, \qquad (2)$$

where $v \approx 246$ GeV is the usual Higgs vacuum expectation value (vev), and w is the vev that completely breaks $SU(2)_X$ giving equal masses $M_X = g_X w/2$ to all $SU(2)_X$ vectors. Symmetry breaking happens when [9]

$$4\lambda_H \lambda_S - \lambda_{HS}^2 < 0, \qquad (3)$$

a condition that can be dynamically verified at low energy because quantum corrections make λ_S smaller at low energy, as described by the beta function

$$\beta_{\lambda_S} \equiv \frac{d\lambda_S}{d\ln\mu} = \frac{1}{(4\pi)^2}\left[\frac{9g_X^4}{8} - 9g_X^2\lambda_S + 2\lambda_{HS}^2 + 24\lambda_S^2\right]. \qquad (4)$$

Unlike in the β function of λ_H, there is no negative Yukawa contribution: β_{λ_S} is definite positive and the gauge term makes λ_S negative at low energy. Thereby the dynamically generated hierarchy between $v \sim w$ and the Planck scale is exponentially large, of order $e^{\lambda_S/\beta_{\lambda_S}}$.

While the analysis of the full one-loop potential is somehow involved, a simple analytic approximation holds in the limit of small λ_{HS} (which will be phenomenologically justified *a posteriori*). In this limit the instability condition of Eq. (3) can be approximated as $\lambda_S < 0$ and the potential at one loop order can be approximated by inserting a running λ_S in the tree-level potential of Eq. (1):

$$\lambda_S \simeq \beta_{\lambda_S} \ln s/s_*, \tag{5}$$

where s_* is the critical scale below which λ_S becomes negative. The use of s_* is not an approximation but a convenient parametrization. Given that around $s \sim s_*$ the typical size of λ_S is β_{λ_S}, "small λ_{HS}" in Eq. (3) precisely means $\lambda_{HS}^2 \ll \lambda_H \beta_{\lambda_S}$. In this approximation, the potential is minimized as

$$v \simeq w\sqrt{\frac{\lambda_{HS}}{2\lambda_H}}, \qquad w \simeq s_* e^{-1/4}. \tag{6}$$

The scalar mass matrix at the minimum in the (h, s) basis is

$$v^2 \begin{pmatrix} 2\lambda_H & -\sqrt{2\lambda_H \lambda_{HS}} \\ -\sqrt{2\lambda_H \lambda_{HS}} & \lambda_{HS} + 2\beta_{\lambda_S}\lambda_H/\lambda_{HS} \end{pmatrix}. \tag{7}$$

Its eigenvalues m_1 and m_2 are

$$m_1^2 \simeq 2v^2\lambda, \qquad m_2^2 \simeq v^2 \frac{2\beta_{\lambda_S}\lambda_H}{\lambda_{HS}}, \tag{8}$$

where $\lambda \simeq \lambda_H - \lambda_{HS}^2/\beta_{\lambda_S}$, having included the next-to-leading term in the small λ_{HS} approximation (as discussed later, this small effect helps in keeping λ_H and the whole potential stable up to large Planckian field values [10]). The first state, h_1 can be identified with the Higgs boson with mass $m_1 \approx 125.6$ GeV [1]. The mixing angle, defined by $h_1 = h\cos\alpha + s\sin\alpha$, is given by

$$\sin 2\alpha = \frac{v^2\sqrt{8\lambda_H\lambda_{HS}}}{m_2^2 - m_1^2}, \quad \text{i.e.,} \quad \alpha \overset{m_1 \ll m_2}{\simeq} \frac{\lambda_{HS}^{3/2}}{\sqrt{2\lambda_H}\beta_{\lambda_S}}. \tag{9}$$

Due to mixing, the extra state $h_2 = s\cos\alpha - h\sin\alpha$ inherits the couplings to SM particles of the Higgs boson h, rescaled by the factor $\sin\alpha$. Note that m_1^2 can be rewritten as $m_1^2 \simeq w^2\lambda_{HS}$, showing explicitly that the SM scalar boson mass is induced by the λ_{HS} portal, proportionally to the $SU(2)_X$ gauge symmetry breaking scale w. Electroweak symmetry breaking (EWSB) does not need a large value for the portal coupling λ_{HS}, provided that w is large enough.[1] Effectively s acts as "the Higgs of the Higgs" and as "the Higgs of dark matter." Furthermore, the "Higgs of the s" is s itself: all scales are dynamically generated via dimensional transmutation.

As discussed in [7,15,16], the SU(2)X vectors with mass M_X are DM candidates, automatically stable thanks to the analogous of the accidental custodial symmetry of the SM.[2] Such symmetry can be violated only by nonrenormalizable dimension-6 operators. If suppressed by Planckian scales, such operators leave these hidden vector dark matter TeV-scale particles stable enough on cosmological time scales.

III. PHENOMENOLOGICAL ANALYSIS

The DM thermal relic abundance and DM indirect signals require the computation of σv, the nonrelativistic DM annihilation cross section times the relative velocity v. We compute them in the limit of small λ_{HS}, where only $SU(2)_X$ gauge interactions are relevant, and the mixing angle α is small, such that DM dominantly annihilates into s. The model gives rise to $VV \to ss$ annihilations and to $VV \to Vs$ semi-annihilations [7] (see also [15,16,19]). Averaging over initial spin and gauge components we find

$$\sigma v_{\rm ann} = \frac{11 g_X^2}{1728\pi w^2}, \qquad \sigma v_{\rm semi\text{-}ann} = \frac{g_X^2}{32\pi w^2}. \tag{10}$$

The cosmological DM abundance is reproduced as a thermal relic for

$$\sigma v_{\rm tot} = \sigma v_{\rm ann} + \frac{1}{2}\sigma v_{\rm semi\text{-}ann} \approx 2.3 \times 10^{-26}\ \frac{\rm cm^3}{\rm s}, \tag{11}$$

which means

$$g_X \simeq w/1.9\ \text{TeV}. \tag{12}$$

[1]This differs from the DM driven EWSB mechanism proposed in Ref. [11] in the framework of the inert Higgs doublet model (see also [12,13] with singlet scalars). In these models only the SM scalar field gets a vacuum expectation value, so that in order to compensate the negative top Yukawa coupling contribution to β_{λ_H}, EWSB requires either quite large new quartic couplings (that can become nonperturbative already at multi-TeV energies) or many extra scalars. This differs also from another DM induced EWSB framework based on dimensionful scalar parameters turning negative at low energy [14].

[2]DM vectors are a triplet, and thereby cannot decay into the bosons, which are singlets under the dark custodial symmetry. In this respect, it is important that the dark gauge group is $SU(2)_X$. In models that employ instead an Abelian U(1)X dark gauge group [7,17], vectors are unstable unless kinetic mixing with U(1)Y is forbidden. In models with a non-Abelian group larger than $SU(2)_X$ it is more difficult to find a scalar representation with dimension just above the dimension of the gauge group, which can fully break it (but possible if one makes extra assumptions on the structure of the vacuum [18]).

With this assumption, the model only has one free parameter: out of the four parameters λ, λ_{HS}, λ_S (or s_*) and g_X, three of them are fixed by the observed Higgs mass and vacuum expectation value, and by the cosmological DM abundance $\Omega_{\rm DM}$ [20]. We can view g_X or λ_{HS} as the only free parameter. They are related as

$$\lambda_{HS} \simeq 0.004/g_X^2. \tag{13}$$

The observables are predicted in terms of g_X as

$$\begin{aligned} m_2 &\simeq 165\ \text{GeV} g_X^3, \qquad \alpha \simeq 0.07/g_X^7,\\ M_X &\simeq 985\ \text{GeV} g_X^2,\\ \sigma_{\rm SI} &\simeq 0.7\times 10^{-45}\ \text{cm}^2[1-(m_1^2/m_2^2)]^2/g_X^{12}, \end{aligned} \tag{14}$$

where $\sigma_{\rm SI}$ is the DM direct detection cross section, computed below. These approximations hold in the limit of small λ_{HS}, which numerically means $\lambda_{HS} \ll \sqrt{\lambda_H \lambda_S} \approx 0.015$, i.e. $g_X \gg 0.5$. Figure 1 shows the same predictions, computed numerically without making the small λ_{HS} approximation.

A. The Higgs of the Higgs

The left panel of Fig. 1 shows the predictions for Higgs physics. We see that the new scalar h_2 cannot have a mass in the range between 100 and 140 GeV. Indeed, the two scalar states roughly have the same mass terms for $\lambda_{HS} \approx 0.005$;
however, due to the off-diagonal term in their mass matrix, the mass difference must be larger than

$$|m_1 - m_2| \cdots |m_{12}^2|/m_1 \approx m_1\sqrt{\lambda_{HS}/2\lambda_H} \approx 17\ \text{GeV}. \tag{15}$$

The extra state h_2 behaves as an extra Higgs boson with couplings rescaled by $\sin\alpha$. This means that it is a narrow resonance even if heavier than 1 TeV. For $m_2 < 2m_1$ the extra scalar behaves as a Higgs-like state with production cross section suppressed by $\sin^2\alpha$ (the decay mode into h_1 and an off-shell h_1 being subdominant), while for $m_2 > 2m_1$ the extra state also has a decay width into two Higgs,

$$\Gamma(h_2 \to h_1 h_1) = \frac{\lambda_{HS}^2}{32\pi}\frac{w^2}{m_2^2}\sqrt{m_2^2 - 4m_1^2}, \tag{16}$$

which contributes up to 20% to its total width, dominated by $h_2 \to WW, ZZ, t\bar{t}$. The shaded regions in the left panel of Fig. 1 are excluded by LEP (at small mass) and LHC (at large mass, $h \to WW$ searches are plotted as dashed curves and $h \to ZZ$ searches as dot-dashed curves). Future sensitivities are discussed in [21]. Present experimental searches for $h \to ZZ$ and for $h \to \gamma\gamma$ show some (nonstatistically significant) hint for an extra state at $m_2 \approx 143$ GeV [1].

The cross section for DM production at LHC (mediated by off-shell h_1 or h_2) can easily be negligibly small.

B. Direct dark matter signals

The spin-independent cross section for DM direct detection is [7]

$$\sigma_{\rm SI} = \frac{m_N^4 f^2}{16\pi v^2}\left(\frac{1}{m_1^2} - \frac{1}{m_2^2}\right)^2 g_X^2 \sin^2 2\alpha, \tag{17}$$

where $f \approx 0.295$ is the nucleon matrix element and m_N is the nucleon mass. The right panel of Fig. 1 shows the predictions for DM direct searches. Present direct detection constraints imply the bounds $\lambda_{HS} = 0.007$ (so that the approximation of small λ_{HS} holds in the phenomenologically interesting region), $m_2 = 70$ GeV, $w = 1.5$ TeV, $M_X = 560$ GeV, $g_X = 0.75$, and $\alpha = 0.5$. A value of $\lambda_{HS} = 0.007$ is also disfavored by LEP Higgs searches [left panel of Fig. 1]. Future experiments should be able to probe smaller values of λ_{HS} by improving the sensitivity to the extra Higgs cross section by 1–2 orders

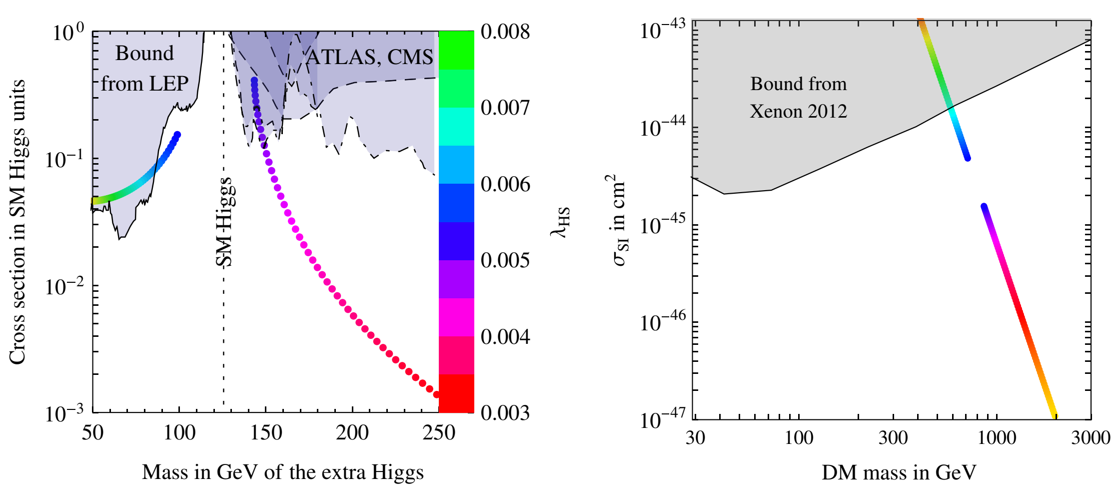

FIG. 1 (color online). Predicted cross sections for the extra scalar boson (left) and for DM direct detection (right) as functions of the only free parameter of the model λ_{HS}, varied as shown in the color legend.

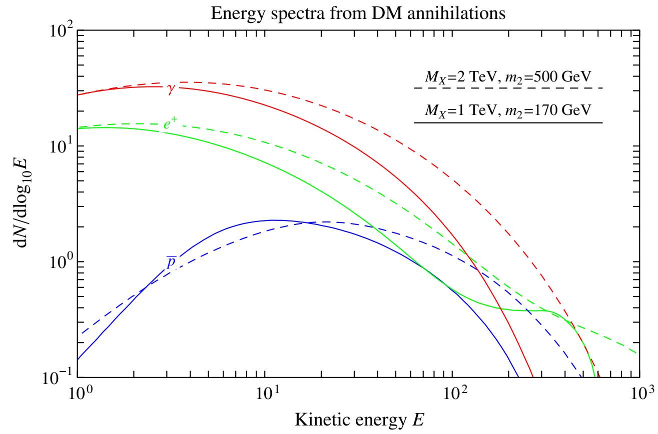

FIG. 2 (color online). Energy spectra of final-state antiprotons, positrons, and photons produced per DM annihilation for two different values of the DM mass: 1 TeV (continuous curves) and 2 TeV (dashed).

of magnitude and the sensitivity to $\sigma_{\rm SI}$ by 2–3 orders of magnitude before 2020.

Signals disappear in the limit of small λ_{HS}. However, λ_{HS} cannot be arbitrarily small, because Eq. (13) would imply a nonperturbatively large value for g_X.[3] In practice, values of λ_{HS} below about 10^{-3} imply a large value of the gauge coupling $g_X = 2$ and correspond to $m_2 = 1.3$ TeV, $M_X = 4$ TeV, and $\alpha = 0.0005$. This is enough to render the signals too small to be observed in forthcoming experiments.

C. Indirect dark matter signals

Concerning indirect DM signals, DM annihilates into ss and into s DM at leading order in the small mixing α. The total amount of s particles produced is not affected by the presence of coannihilations and is effectively dictated by the standard thermal cross section $\sigma v_{\rm tot}$ in Eq. (11). Next, via its mixing with the SM scalar, the s particles decay into SM particles, producing a detectable astrophysical flux of e^+, $\bar{p}$, γ, ν, $\bar{d}$ which can be computed using e.g. the public code in [22]. Figure 2 shows two examples of $\bar{p}$, e^+, γ spectra at production (astrophysical propagation effects are not included), which are typical of hadronic channels. Present data have a sensitivity to σv which, for the considered DM masses $M_X \approx$ TeV, is about 2 orders of magnitude above the predicted thermal cross section.

Furthermore, DM semi-annihilations (around the center of the Galaxy and of the Sun) produce a flux of DM particles with $E = M_X$. They could be detected by looking for hadronic showers produced by DM scatterings with matter. However, it is difficult to detect such DM particles, because their cross section on matter, $\sigma \sim g_X^4 m_N^2 \sin^2 2\alpha / 4\pi M_X^4$, is smaller than the analogous cross section of neutrinos on matter.

[3] If the g_X coupling becomes nonperturbative at a scale Λ_X larger than the scale where Eq. (3) is satisfied, the dark scalars and vectors confine [15]. In this case the $\lambda_{HS} S^\dagger S H^\dagger H$ term could also lead to EWSB by inducing a negative $\lambda_{HS}\Lambda_X^2 H^\dagger H$ mass term [15].

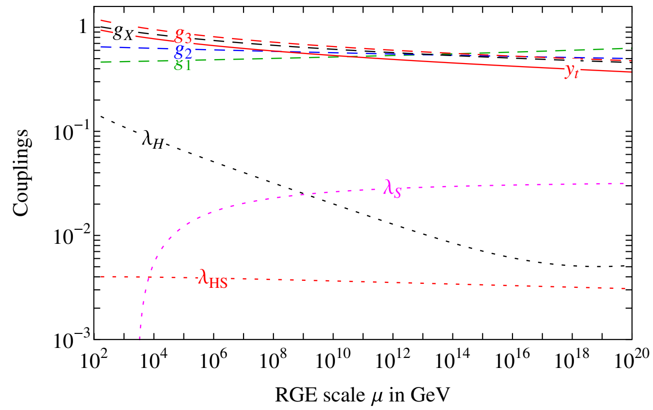

FIG. 3 (color online). Running of the model parameters up to the Planck scale for $g_X = 1$.

D. RGE and stability of the potential

Having determined the weak-scale values of the parameters, we now explore how the model can be extrapolated up to large energies, dynamically generating the weak scale. The renormalization group equations (RGE) for the model are those of the SM, with extra terms in the RGE for the quartic Higgs coupling,

$$(4\pi)^2 \frac{d\lambda_H}{d\ln\mu} = \left(12 g_t^2 - \frac{9g_1^2}{5} - 9g_2^2\right)\lambda_H - 6g_t^4 + \frac{27 g_1^4}{200} + \frac{9}{20} g_2^2 g_1^2 + \frac{9g_2^4}{8} + 24\lambda_H^2 + 2\lambda_{HS}^2, \quad (18)$$

and supplemented by the RGE for the extra couplings:

$$(4\pi)^2 \frac{dg_X}{d\ln\mu} = -\frac{43}{6} g_X^3 - \frac{1}{(4\pi)^2}\frac{259}{6} g_X^5 + \cdots, \quad (19a)$$

$$(4\pi)^2 \frac{d\lambda_{HS}}{d\ln\mu} = \lambda_{HS}\left(6y_t^2 - \frac{9g_X^2}{2} - \frac{9g_1^2}{10} - \frac{9g_2^2}{2} + 12\lambda_H + 12\lambda_S\right) - 4\lambda_{HS}^2, \quad (19b)$$

$$(4\pi)^2 \frac{d\lambda_S}{d\ln\mu} = -9g_X^2\lambda_S + \frac{9g_X^4}{8} + 2\lambda_{HS}^2 + 24\lambda_S^2. \quad (19c)$$

Figure 3 shows the resulting running of the couplings of the model up to the Planck scale for $g_X = 1$, which corresponds to $\lambda_{HS} = 0.004$. We notice that the model interprets the observed proximity of the QCD scale to the electroweak scale as due to a proximity between the strong gauge coupling g_3 and the dark gauge coupling g_X. Indeed, g_3 and g_X happen to have not only similar values at the weak scale, but also a numerically similar β function, such that all gauge couplings roughly reach a common value at large energies. At low energy g_X becomes large, of order one, triggering a negative λ_S and consequently dynamically generating the DM scale and the weak scale.

E. Dark/electroweak phase transition

The mechanism of dynamical scale generation implies a negative value of the cosmological constant (barring metastable minima). The contribution of the present model is $V_{\min} \simeq -w^4 \beta_{\lambda_S}/16$. Despite being suppressed by a one-loop factor, this contribution is larger by about 60 orders of magnitude than the observed value. Assuming that the cosmological constant problem is solved by a fine-tuning, we can proceed to study how the dark and electroweak phase transitions occur during the big bang.

We recall that the SM predicts a second-order phase transition where the Higgs boson starts to obtain a vacuum expectation value $v(T)$ at temperatures below $T_c^{\rm SM} \approx 170$ GeV and sphalerons decouple when $T_{\rm dec}^{\rm SM} \approx v(T_{\rm dec}^{\rm SM}) \approx 140$ GeV [23].

Within the present model, using again the small λ_{HS} approximation, the one-loop thermal correction to the potential is

$$V_T(s, h \approx 0) = \frac{9T^4}{2\pi^2} f\left(\frac{M_X}{T}\right) + \frac{T}{4\pi}[M_X^3 - (M_X^2 + \Pi_X)^{3/2}], \tag{20}$$

where $f(r) = \int_0^\infty x^2 \ln(1 - e^{-\sqrt{x^2+r^2}})dx$ and $\Pi_X = 11 g_X^2 T^2/6$ is the thermal propagator for the longitudinal X component which accounts for resummation of higher order daisy diagrams [12,13]. Equation (20) predicts that s and consequently h acquire a vacuum expectation value through a first-order phase transition. The critical temperature at which the two phases are degenerate is $T_c/M_X \simeq 0.37$, 0.42, 0.49, 0.75 for $g_X = 0.75$, 1, 1.2, 1.5, respectively.

A cosmological first-order phase transition occurring at temperatures around the weak scale generates gravitational waves at a potentially detectable level.

Their present peak frequency and energy density are [13,24]

$$f_{\rm peak} \approx 5\ {\rm mHz} \frac{\beta/H}{100} \frac{T_f}{1\ {\rm TeV}}, \qquad \Omega_{\rm peak} h^2 \approx 1.84 \times 10^{-6} \kappa^2 \left(\frac{\alpha}{1+\alpha}\right)^2 \frac{H^2}{\beta^2} \frac{v_b^3}{0.42 + v_b^2}, \tag{21}$$

where $T_f \lesssim T_c$ is the temperature at which the phase transition happens, α is the energy fraction involved in the first-order phase transition, H/β is the duration of the phase transition in Hubble units, $v_b = (\sqrt{1/3} + \sqrt{\alpha^2 + 2\alpha/3})/(1+\alpha)$ is the wall velocity, and $\kappa = (0.715\alpha + \sqrt{8\alpha/243})/(1 + 0.715\alpha)$ is the fraction of latent heat converted into gravitational waves. α and β are explicitly given by

$$\alpha \equiv \frac{\Delta V - d(\Delta V)/d\ln T}{\pi^2 g_* T^4/30}, \qquad \frac{\beta}{H} \equiv \frac{d(S_3/T)}{d\ln T}, \tag{22}$$

where S_3 is the action of the thermal bubble that determines the tunneling rate per space-time volume as $\Gamma \approx (S_3/2\pi T)^{3/2} T^4 e^{-S_3/T}$ and ΔV is the potential difference between the two minima. These quantities are evaluated at $T \approx T_f$, which is roughly determined as the temperature at which $S_3/T \approx 4 \ln M_{\rm Pl}/M_X \approx 142$. Given that S_3 scales as $1/\beta_{\lambda_S} \propto g_X^4$, the result depends strongly on g_X:

(i) For the critical value $g_X \approx 1.2$, one has $\alpha \approx 1$ in view of $T_f \approx 0.15 T_c$. Furthermore $\beta/H \approx 70$ such that $\Omega_{\rm peak} h^2 \approx 2 \times 10^{-11}$, which could be easily detected by planned experiments such as LISA. Notice that $T_f \approx M_X/13 \approx 110$ GeV is larger than the DM freeze-out temperature, $T_{\rm fo} \approx M_X/25$ (such that DM freeze-out is negligibly affected by the phase transition) and is smaller than $T_c^{\rm SM}$ (such that also the SM scalar boson is involved in the first-order phase transition).

(ii) For $g_X \gtrsim 1.2$ tunneling is faster and consequently T_f higher. For example, for $g_X \approx 1.5$ one has $T_f \approx 0.5 T_c \approx 0.4 M_X \approx 800$ GeV and $\alpha \approx 0.01$, $\beta/H \approx 200$ such that $\Omega_{\rm peak} h^2 \approx 10^{-19}$, which could be detected by futuristic experiments. The Higgs phase transition happens later independently of the dark phase transition.

(iii) For $g_X \lesssim 1.2$ thermal tunneling is slower than the Hubble rate such that the universe would enter into an inflationary stage, which presumably ends when the temperature cools down to the $SU(2)_X$ confinement scale Λ_X after $N \approx \ln s_*/\Lambda_X \approx 8\pi^2/7 g_X^2$ e-folds, next reheating the universe up to $g_* \pi^2 T_{\rm reh}^4/30 \approx \Delta V$ i.e. $T_{\rm reh} \approx M_X/9$. The baryon asymmetry gets suppressed by a factor $\approx e^{-3N}$, such that the model is excluded when this factor is smaller than the observed baryon asymmetry $\approx 10^{-9}$, provided that the baryon asymmetry cannot be regenerated at the weak scale.

The critical value of g_X could have an uncertainty of about 30%, given that higher order corrections to the thermal potential are suppressed by g_X/π.

Finally note that the extended dark/electroweak phase transition also occurs for a more general situation with e.g. positive H and S squared mass terms typically smaller than the v and w symmetry breaking scales.

IV. CONCLUSIONS

"Just so" comparable small masses for the Higgs boson and for dark matter (much smaller than the Planck scale) satisfy a reformulation of the naturalness concept, modified by assuming that quadratic divergences should be ignored and thereby named finite naturalness. Within this heretic point of view, it might be more satisfactory to find a dynamical explanation of the smallness and of the proximity of these low-energy scales. We here assumed that such masses vanish to start with, and that they originate from a physical mechanism that occurs at such energy scales.[4]

[4]This assumption is not demanded by finite naturalness. An alternative possibility compatible with this scenario is e.g. that gravitational loops of unobservable intermediate-scale particles (as suggested to us by A. Arvinataki, S. Dimopoulos, and S. Dubovsky) could generate small and comparable Higgs and scalar DM masses.

This is achieved by a simple extension of the standard model, which contains just one extra scalar doublet under one extra SU(2) gauge group. The extra vectors are automatically stable thermal DM candidates, just like the proton is automatically stable within the SM. The extra scalar doublet gives mass to such vectors (because of their gauge interactions), to the SM Higgs boson (because of a quartic coupling between them), and to itself (because its quartic couplings run negative at low energy). Thereby, all scales are related and exponentially suppressed with respect to the Planck scale.

As a function of only one free parameter, the model predicts the properties of the extra scalar [observable at collider experiments, see left panel of Fig. 1], of dark matter [observable in direct and indirect detection experiments, see right panel of Fig. 1]. The scalar potential of the model can be stable up to the Planck scale, even when the SM potential is unstable, namely for the present best-fit values of its parameters. In cosmology, the model predicts a first order phase transition with emission of gravitational waves at a possibly detectable level.

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