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Cosmological Signature of the Standard Model Higgs Vacuum Instability: Primordial Black Holes as Dark Matter

J. R. Espinosa,[1,2] D. Racco,[3] and A. Riotto[3]
[1]Institut de Física d'Altes Energies (IFAE), The Barcelona Institute of Science and Technology (BIST), Campus UAB, 08193 Bellaterra, Barcelona, Spain
[2]ICREA, Institució Catalana de Recerca i Estudis Avançats, Passeig de Lluís Companys 23, 08010 Barcelona, Spain
[3]Département de Physique Théorique and Centre for Astroparticle Physics (CAP), Université de Genève, 24 quai E. Ansermet, CH-1211 Geneva, Switzerland

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For the current central values of the Higgs boson and top quark masses, the standard model Higgs potential develops an instability at a scale of the order of 10^{11} GeV. We show that a cosmological signature of such instability could be dark matter in the form of primordial black holes seeded by Higgs fluctuations during inflation. The existence of dark matter might not require physics beyond the standard model.

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Introduction.—It has been known for a long time that the standard model (SM) Higgs potential develops an instability at large field values [1,2]. For the current central values of the Higgs boson and top quark masses, the quartic coupling λ in the Higgs potential becomes negative for Higgs field values $\gtrsim 10^{11}$ GeV, making our electroweak vacuum not the one of minimum energy. While some take this as motivation for the presence of new physics to change this feature, this is not necessarily a drawback of the SM. Indeed, our current vacuum is quite stable against both quantum tunneling in flat spacetime and thermal fluctuations in the early Universe [2,3].

The situation is different during inflation [4]. If the effective mass of the Higgs field is smaller than the Hubble rate H during inflation, quantum excitations push the Higgs field away from its minimum. The classical value of the Higgs field randomly walks receiving kicks $\sim \pm (H/2\pi)$ each Hubble time and can surmount the potential barrier and fall deep into the unstable side of the potential [4–6]. At the end of inflation, patches where this happened will be anti–de Sitter regions, and they are lethal for our Universe as they grow at the speed of light [7]. One can derive upper bounds on H, which depend on the reheating temperature $T_{\rm RH}$ and on the Higgs coupling to the scalar curvature or to the inflaton [7,8].

The upper bound on H depends on $T_{\rm RH}$ because, for sufficiently large values of $T_{\rm RH}$, patches in which the Higgs field probes the unstable part of the potential can be recovered thanks to the thermal effects after inflation. Indeed, the mass squared of the Higgs field receives a positive correction proportional to T^2 in such a way that in those would-be dangerous regions the Higgs field can roll back down to the origin and be safe.

The physical implications of living in a metastable electroweak vacuum are fascinating and have far-reaching consequences for cosmology. This has triggered much activity in a field that involves inflationary dynamics, the physics of preheating, the interplay between Higgs properties, and observables of cosmological interest, etc. In spite of this richness, a word of warning is in order: the energy scale of this physics is very high and we have no smoking-gun signature (comparable to proton decay for GUTs) that the electroweak vacuum metastability is actually realized in nature (with the exception of the vacuum decay itself).

One reasonable question to ask is how can we probe, even if indirectly, the SM Higgs vacuum instability. In this short note we argue that there might be a cosmological signature of the SM vacuum instability: the very presence of dark matter (DM) in our Universe. We argue that the origin of DM does not need physics beyond the SM: DM may be due to primordial black holes seeded by the perturbations of the Higgs field generated during the last stages of inflation. The black holes may provide the seeds for structure formation [9,10].

The picture is the following. During inflation, there are patches where the Higgs field has been pushed by quantum fluctuations beyond the potential barrier and is classically rolling down the slope. Higgs fluctuations do not contribute significantly to the total curvature perturbation ζ, which is ultimately responsible for the anisotropies in the Cosmic

Microwave Background (CMB). Higgs perturbations instead grow to large values in the last e folds of inflation, which are irrelevant for observations in the CMB. When inflation ends and reheating takes place, these regions are rescued by thermal effects and the Higgs field rolls down to the origin of its potential. At later times, the Higgs perturbations reenter inside the Hubble radius, and they provide high peaks in the matter power spectrum that give rise to primordial black holes (PBH). We show that these PBHs can provide the DM we see in the Universe today.

Within a more anthropic attitude, one could say that the electroweak SM instability is beneficial to our own existence as DM is necessary to form structures. In the absence of other DM candidates, the SM would be able to provide the right DM abundance. As discussed below, although the parameter choices needed for PBH formation might seem fine-tuned, they would be anthropically motivated. In particular, this mechanism offers an anthropic explanation of why the electroweak vacuum is metastable (but near critical, very close to being stable).

The dynamics during inflation.—We are agnostic about the details of the model of inflation and the origin of the curvature perturbation responsible for the CMB anisotropies, which we call $\zeta_{\rm st}$. This $\zeta_{\rm st}$ might be caused by a single degree of freedom [11] or by another mechanism such as the curvaton [12]. Also, we take a constant Hubble rate H during inflation and suppose that it ends going through a period of reheating characterized by a reheating temperature $T_{\rm RH}$. Of course, one can repeat our calculations for a preferred model of inflation. We suppose that H is large enough to have allowed the SM Higgs field to randomly walk above the barrier of its potential to probe the potentially dangerous unstable region. As a representative value we take $H \simeq 10^{12}$ GeV.

Despite the Higgs field's negative potential energy, this region keeps inflating as long as the total vacuum energy during inflation is larger, that is, for

$$3H^2 m_P^2 \gtrsim \frac{\lambda}{4} h_c^4, \tag{1}$$

where h_c is the Higgs field's classical value and $m_P = 2.4 \times 10^{18}$ GeV is the reduced Planck mass. The equation of motion of the classical value of the SM Higgs field is

$$\ddot{h}_c + 3H\dot{h}_c + V'(h_c) = 0, \tag{2}$$

where, as usual, dots represent time derivatives and primes field derivatives. For the sake of simplicity, from now on we will approximate the potential as

$$V(h_c) = -\frac{1}{4}\lambda h_c^4, \tag{3}$$

with $\lambda > 0$ running logarithmically with the field scale. During inflation, λ should in fact be evaluated at a scale μ given by $\mu^2 \simeq h_c^2 + H^2$ [6]. A typical value (for $h_c \gtrsim 10^{12}$ GeV) is $\lambda \simeq 10^{-2}$. In order to make any prediction deterministic and not subject to probability arguments, we are interested in the regime in which the dynamics of the zero mode of the Higgs field is dominated by the classical motion rather than by the randomness of the fluctuations. We require therefore that in a Hubble time, $\Delta t = 1/H$, the classical displacement of the Higgs field

$$\Delta h_c \simeq -\frac{V'(h_c)}{3H^2}, \tag{4}$$

is larger (in absolute value) than the quantum jumps

$$\Delta_q h \simeq \pm \left(\frac{H}{2\pi}\right). \tag{5}$$

This implies that, inside the inflating region, h_c must be bounded from below

$$h_c^3 \gtrsim \frac{3H^3}{2\pi\lambda}. \tag{6}$$

We call t_* the initial time at which the Higgs field starts its classical evolution. In this estimate, we assume that the motion of the Higgs field is friction dominated, that is $\ddot{h}_c \lesssim 3H\dot{h}_c$. This is true as long as $h_c^2 \lesssim 3H^2/\lambda$. If so, the Higgs field is slowly moving for a sufficient number of e folds. The evolution of the classical value of the Higgs field is

$$h_c(N) \simeq \frac{h_e}{(1 + 2\lambda h_e^2 N/3H^2)^{1/2}}, \tag{7}$$

where we have introduced the number of e folds until the end of inflation N and denoted by h_e, the value of the classical Higgs field at the end of inflation.

Meanwhile, Higgs fluctuations are generated. Perturbing around the slowly-rolling classical value of the Higgs field and accounting for metric perturbations as well, the Fourier transform of the perturbations of the Higgs field satisfy the equation of motion (in the flat gauge)

$$\delta\ddot{h}_k + 3H\delta\dot{h}_k + \frac{k^2}{a^2}\delta h_k + V''(h_c)\delta h_k = \frac{\delta h_k}{a^3 m_P^2}\frac{d}{dt}\left(\frac{a^3}{H}\dot{h}_c^2\right), \tag{8}$$

where a is the scale factor and the last term accounts for the backreaction of the metric perturbations. Driven by the Higgs background evolution in the last e folds of inflation, the Higgs perturbations grow significantly after leaving the Hubble radius. The reason is the following. Having numerically checked that the last term in Eq. (8) is negligible, the Higgs perturbations and $\dot{h}_c$ solve the same equation on scales larger than the Hubble radius $k \ll aH$, as can be seen by taking the time derivative of Eq. (2). Therefore the two quantities must be proportional to each other during the evolution and on super-Hubble scales

$$\delta h_k = C(k)\dot{h}_c(t). \tag{9}$$

Matching at Hubble crossing $k = aH$, this super-Hubble solution for δh_k, with its standard wave counterpart on sub-Hubble scales, implies that

$$C(k) = \frac{H}{\dot{h}_c(t_k)\sqrt{2k^3}}, \tag{10}$$

where t_k is the time when the mode with wavelength $1/k$ leaves the Hubble radius. The growth of δh_k is therefore dictated by the growth of $\dot{h}_c$. These Higgs perturbations will be responsible for the formation of PBHs. In fact, we should deal with the comoving curvature perturbation ζ, which is gauge-invariant and reads (still in the flat gauge)

$$\zeta = H\frac{\delta\rho}{\dot{\rho}} = \frac{\dot{\rho}_{\rm st}}{\dot{\rho}}\zeta_{\rm st} + \frac{\dot{\rho}_h}{\dot{\rho}}\zeta_h = \frac{\dot{\rho}_{\rm st}}{\dot{\rho}}\zeta_{\rm st} + H\frac{\delta\rho_h}{\dot{\rho}}, \tag{11}$$

where ζ_h is the Higgs perturbation. We assume $\zeta_{\rm st}$ is conserved during inflation on super-Hubble scales and, for simplicity, that there is no energy transfer with Higgs fluctuations. In the curvaton model, for instance, $\zeta_{\rm st}$ could even be zero on large scales during inflation.

Using Eqs. (2) and (8) (again with the negligible last term dropped), one then obtains

$$\delta\rho_h(k \ll aH) = \dot{h}_c\delta\dot{h}_k + V'(h_c)\delta h_k = -3HC(k)\dot{h}_c^2. \tag{12}$$

Since $\dot{\rho}_h = \dot{h}_c(\ddot{h}_c + V'(h_c)) = -3H\dot{h}_c^2$, one can easily show (and we have checked it numerically) that during inflation and on super-Hubble scales, ζ_h reaches the plateau

$$\zeta_h(k \ll aH) = H\frac{\delta\rho_h}{\dot{\rho}_h} = HC(k) = \frac{H^2}{\sqrt{2k^3}\dot{h}_c(t_k)}. \tag{13}$$

This is the quantity that gives the largest contribution to ζ in the last few e folds before the end of inflation.

Dynamics after inflation: reheating.—At the end of inflation, the vacuum energy that has driven inflation gets converted into thermal relativistic degrees of freedom—a process dubbed reheating. For simplicity, we suppose that this conversion is instantaneous, in such a way that the reheating temperature is $T_{\rm RH} \simeq 0.5(Hm_{\rm P})^{1/2}$, obtained by energy conservation, and taking the number of relativistic degrees of freedom to be about 10^2. For our representative value of $H = 10^{12}$ GeV, we obtain $T_{\rm RH} \simeq 10^{15}$ GeV. Because of the thermal effects, the Higgs potential receives thermal corrections such that the potential is quickly augmented by the term [7]

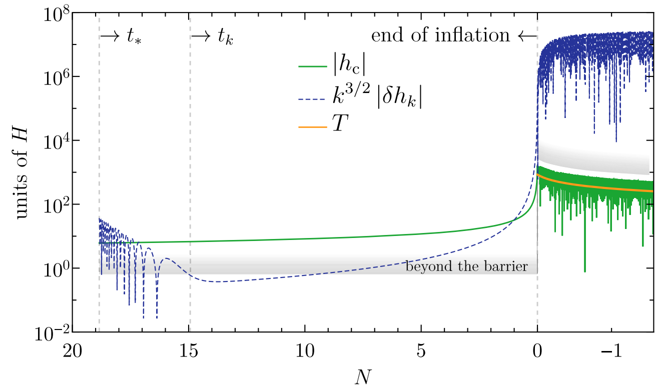

FIG. 1. The evolution of H, T, h_c, and δh_k during the last e folds of inflation, for $k = 50a(t_*)H$ where t_* is defined to be the time when h_c starts its classical evolution. The region of h_c beyond the top of the potential barrier is shaded gray.

$$V_T \simeq \frac{1}{2}m_T^2h_c^2, \qquad m_T^2 \simeq 0.12T^2e^{-h_c/(2\pi T)}, \tag{14}$$

(a fit that is accurate for $h \lesssim 10T$ in the region of interest and includes the effect of ring resummation). If the maximum temperature is larger than the value of the Higgs field $h_{\rm e}$ at the end of inflation, or more precisely if $T_{\rm RH}^2 \gtrsim \lambda h_e^2$, the corresponding patch is thermally rescued and the initial value of the Higgs field immediately after the end of inflation coincides with $h_{\rm e}$. The classical value of the Higgs field starts oscillating around the origin, see Fig. 1. The Higgs fluctuations oscillate as well, with the average value remaining constant and the amplitude slowly increasing for a fraction of e folds. At the same time, the curvature perturbation, with power spectrum $\mathcal{P}_\zeta = k^3/(2\pi^2)|\zeta_k|^2$, given in Fig. 2, gets the largest contribution from the Higgs fluctuations. After inflation, the long wavelength Higgs perturbations decay after several oscillations into radiation curvature perturbation which, being radiation now the only component, will stay constant on super-Hubble scales. We have taken the Higgs damping rate to be $\gamma_h \sim 3g^2T^2/(256\pi m_T) \sim 10^{-3}T$ [13] (where g is the

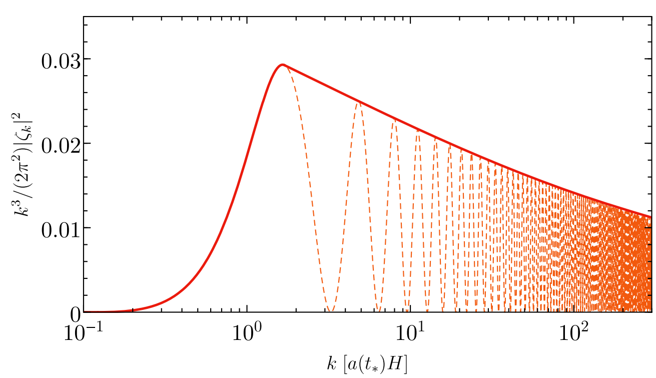

FIG. 2. The power spectrum $\mathcal{P}_\zeta$, shown as the envelope of different modes (averaged over their uncorrelated time phases).

$SU(2)_L$ coupling constant). This value has been derived by noticing that for a thermal Higgs mass $m_T \simeq 0.34T$, the one-loop absorption and direct decay channels for quarks and gauge bosons are forbidden, and the damping occurs through two-loop diagrams involving gauge bosons. Therefore, we have evaluated the value of the curvature perturbation after a fraction of an e fold.

Generation of primordial black holes.—After inflation, the Hubble radius grows, and the perturbations generated during the last e folds of inflation are the first to reenter the horizon. If they are large enough, they collapse to form PBHs almost immediately after horizon reentry, see Ref. [14] and references therein. A given region collapses to a PBH if the density contrast (during the radiation era) $\Delta(\vec{x}) = (4/9a^2H^2)\nabla^2\zeta(\vec{x})$ is above a critical value Δ_c; typically, $\Delta_c \sim 0.45$ [15]. As a result, in order to obtain a significant number of PBHs, the power spectrum on small scales must be sizeable. The mass of a PBH at formation, and corresponding to the density fluctuation leaving the Hubble radius N e folds before the end of inflation, is about [16] $M \simeq (m_P^2/H)e^{2N}$. We first define the variance of the density contrast

$$\sigma_\Delta^2(M) = \int_0^\infty d\ln k W^2(k,R)\mathcal{P}_\Delta(k), \tag{15}$$

where $W(k,R)$ is a Gaussian window function smoothing out the density contrast on the comoving horizon length $R \sim 1/aH$. The mass fraction $\beta(M)$ of the Universe, which ends up into PBHs at the time of formation t_M, is

$$\beta(M) = \int_{\Delta_c}^\infty \frac{d\Delta}{\sqrt{2\pi}\sigma_\Delta} e^{-\Delta^2/2\sigma_\Delta^2} \simeq \frac{\sigma_\Delta}{\Delta_c\sqrt{2\pi}} e^{-\Delta_c^2/2\sigma_\Delta^2}, \tag{16}$$

The total contribution of PBHs at radiation-matter equality (t_{eq}) is obtained by integrating the fraction $\beta(M, t_{eq}) = a(t_{eq})/a(t_M)\beta(M)$ [9]

$$\Omega_{PBH}(t_{eq}) = \int_{M_{ev}(t_{eq})}^{M(t_{eq})} d\ln M\beta(M, t_{eq}), \tag{17}$$

where $M_{ev}(t_{eq}) \simeq 10^{-21}\ M_\odot$ is the lower mass that has survived evaporation at equality, and $M(t_{eq})$ is the horizon mass at equality (which, for our purposes, can be taken equal to infinity).

Figure 3 shows the resulting mass spectrum of PBHs at their formation time. The position of the peak in the PBH mass spectrum is set by the mode k_* that exits the Hubble radius during inflation, when the Higgs zero mode starts its classical evolution. To be on the safe side, we ask that the interesting range of PBH masses is large enough to avoid the bounds from evaporating PBHs by the present time. This requires the dynamics to last about 17 e folds before the Higgs field hits the pole in Eq. (7). Interestingly, this can be achieved in the SM for realistic values of the Higgs

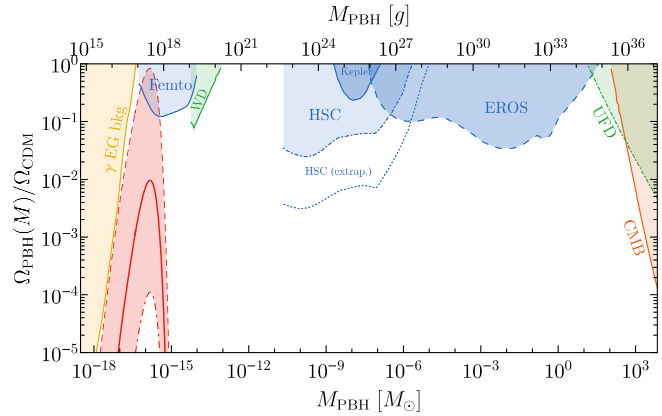

FIG. 3. The spectrum of PBHs at formation generated by the mechanism we discuss (solid red refers to $S_3 = 0$, and dashed lines to $S_3 = \pm 1$), superimposed with the experimental constraints on monocromatic PBH spectra (from Ref. [17] and references therein): in yellow, the observations of extra-galactic γ-ray background; in blue, femto-, micro- and milli- lensing observations from Fermi, Eros, Kepler, Subaru HSC; in green, dynamical constraints from White Dwarves and Ultra-Faint Dwarf galaxies; in orange, constraints from the CMB.

field and top quark masses and α_s: in our numerical example we use $M_h = 125.09$ GeV, $M_t = 172$ GeV, and $\alpha_s = 0.1184$.

In our findings, we have not included the fact that the mass of the PBH is not precisely the mass contained in the corresponding horizon volume, but in fact, it obeys a scaling relation with initial perturbations [16] or the fact that the threshold is shape-dependent [18]. Furthermore, we have not accounted for the fact that the threshold amplitude and the final black hole mass depend on the initial density profile of the perturbation [19]. We estimate that the first two effects change the abundance by order unity. The third effect would require a thorough study of the spatial correlation of density fluctuations. Nevertheless, we have included in Fig. 3 the possible effect of non-Gaussianity in the PBH mass function. To estimate the impact of non-Gaussianity is not an easy task, as one needs to evaluate the second-order contribution to the comoving curvature perturbation ζ_2. A rough estimate based on Ref. [20] gives $\zeta_2 = \mathcal{O}(1)\zeta_1^2$, and therefore, we include two bands corresponding to $S_3 = \pm 1$ in Fig. 3, where $S_3 = \langle\Delta^3\rangle/\sigma_\Delta^4$ is the skewness that appears in the modification of the arguments of the exponential in Eq. (16) via the shift $\nu^2 \to \nu^2\{1 - S_3(\sigma_\Delta/3)[\nu - 2 - (1/\nu^2)]\}$, with $\nu = \Delta_c/\sigma_\Delta$ [21]. The shift in the final abundance is not negligible, but we stress that there will be a set of parameters in our model that can provide the right final abundance. We also stress that the primordial abundance of PBHs depends in a very sensitive way on the value of t_*, keeping all of the other parameters fixed. This does not come as a surprise, as the function $\beta(M)$ is exponentially sensitive to ν. In this sense, the anthropic argument based on the necessity of having

DM would justify a tuned initial PBH abundance. As a final warning, one should keep in mind that splitting the metric into background and perturbations might be questionable for large perturbations.

From the time of equality to now, the PBH mass distribution will slide to larger masses due to merging. While the final word can only be said through *N*-body simulations, one can expect merging to shift the spectrum to higher masses even by orders of magnitude [22] and to spread the spectrum while maintaining the abundance. Accretion, on the other hand, increases both the masses and the abundance of PBHs as DM. On the other side, both merging and accretion help to render the PBHs more long-living. To roughly account for an increase of the current abundance by a representative factor 10^2 because of accretion, we have properly set the abundance at formation time to be $\Omega_{\rm PBH}/\Omega_{\rm DM} \sim 10^{-2}$ (higher values can be achieved). It would be certainly interesting to analyze these issues in more detail and account for the fact that the abundance of PBH has to be of the right magnitude during standard Big Bang nucleosynthesis.

Conclusions.—If the scenario we have presented were in fact realized in nature, we can highlight three points as the most relevant. First, the SM would be capable of explaining DM by itself (supplemented by a period of inflation that is well motivated by other reasons). This has a double side: the SM provides a DM candidate in the form of PBHs and also provides the mechanism necessary to create the PBH seeds during inflation via the quantum fluctuations of the Higgs field in the unstable part of the Higgs potential. Both aspects (DM candidate and the PBH generation mechanism) go against the common lore that physics beyond the SM are needed. In fact, if this scenario were correct, the Higgs field would not only be responsible for the masses of elementary particles but also for the DM content of our Universe. Second, the PBH generation mechanism gives an anthropic handle on the Higgs field near criticality, which would be explained as needed to get sufficient DM, so that large enough structures can grow in the Universe. Finally, the PBHs responsible for DM would represent a conspicuous cosmological signature of the actual existence of an unstable range in the Higgs potential at large field values.

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