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Explorando la metaestabilidad del Higgs en régimen de dominación cinética

Exploring Higgs metastability during kination

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Exploring Higgs metastability during kination

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We investigate an scenario where a non-minimally-coupled-to-gravity spectator Higgs field is responsible for reheating the Universe after inflation through a period of kinetic domination. The minimal reheating temperature at the onset of radiation and the stability of the EW vacuum impose restrictions on the allowed parameter space associated to the Standard Model Higgs field potential. In this work, we explore the impact of new physics through the inclusion of a sixth order operator in the Higgs effective potential and the subsequent new scenarios arising. We analyze the constraints on the new allowed parameter space and make a comparison with the standard case. Employing 3+1-dimensional classical lattice simulations with *CosmoLattice*, we compute the main cosmological observables to determine whether the previous results are robust under the inclusion of the sixth order operator or whether it introduces relevant new physics contributions.

I. INTRODUCTION

The energy scale in which the Standard Model (SM) becomes theoretically inconsistent is not fully determined. It is clear that physics beyond the Standard Model (BSM) must appear at energies above Planck scale ($\sim 10^{18}$ GeV) due to the presence of Landau poles [1] and the possible increasing influence of gravity [2]. However, there is no evidence that SM is a complete theory for any energy scale below.

Particularly, the Higgs sector has been experimentally explored at energies up to $\sim$TeV [3]. Theoretical inconsistencies arise at energies below Planck mass when considering the stability of the electroweak (EW) vacuum where the Higgs field resides today. This is caused by the running of the SM renormalized Higgs self-coupling $\lambda(\mu)$. It depends on several parameters of the model, specially on the top quark Yukawa coupling [4]. For a plausible range of their values, it might become negative at energies of order $\sim 10^{10}$ GeV [5]. This would lead to a two minimum Higgs potential where the EW vacuum is metastable and a phase transition to the true vacuum at higher energies would be possible. As this phase transition is clearly not observed today, the lifetime of the unstable vacuum must be longer than the lifetime of the Universe. One may consider whether the parameter space is restricted to values for which this condition holds, or whether new physics emerge at these scales. The latter would enforce the condition or even restore the convexity of the potential and stability of the EW vacuum.

The distance from the already experimentally explored energy range and reasonable inflationary scales ($\sim 10^{15}$ GeV) is orders of magnitude long. Together with the limited understanding of the Higgs boson behaviour and interactions at high energies, this makes us wonder whether physics BSM may have described the Higgs boson role in the early Universe. Different possibilities emerge: there are theoretical models where the Higgs field dominates and is responsible for inflation [6] and others exist where it is energetically subdominant. In particular, in the present work, we are going to focus on the second case, considering a spectator Higgs field in a Universe where a period of kinetic domination follows inflation. This scenario is typical of quintessential inflation models [7], whose main idea is considering that the inflaton field also acts as dark energy in the late Universe. Usual quintessential models introduce a stiff expansion era (*kination*) after the slow-roll phase. In this regime, the energy density of the inflaton field redshifts as $\rho_\phi \propto a^{-6}$, until radiation, scaling as $\rho_r \propto a^{-4}$, becomes dominant. Inflation models where the inflaton field oscillates around a minimum of potential generically propose reheating mechanisms based on parametric resonance [8]. Instead, quintessential inflation models suggest a reheating phase originated by gravitational production of particles. A description of this mechanism is provided in Ref. [9], where particle production occurs because of the transition between a quasi-de Sitter space in inflation to a Robertson-Walker Universe.

Here, we will consider the spectator Higgs field non-minimally coupled to gravity. Such a coupling is generically generated by renormalization in curved spacetime and it is therefore expected to be present [10]. As shown in Ref. [11], this non-minimal coupling to the Ricci scalar, together with a kination era after inflation, can trigger a Hubble-induced phase transition: a change of the sign of the Ricci scalar results in an effective tachyonic mass that leads to large vacuum expectation field values. The field reaches the minimum and oscillates around it until the symmetry is restored. This allows the Higgs scalar field to transfer its energy to the SM particles and, as presented in Ref. [12], ultimately to be responsible for reheating the Universe. The feasibility of this scenario is determined by the stability of the EW vacuum, the fulfillment of the lower bound restriction on the reheating temperature at the onset of radiation domination, and the validity of all approximations used. The Universe must ultimately reside in the EW vacuum, since it is the observed state today. This imposes some constraints on the allowed parameter space of the model, which have been investigated in Ref. [13] assuming no new physics are present at the energies considered. This limitation sets an upper bound on the temperature at the onset of radiation era. If the reheating temperature were sufficently high, thermal corrections to the potential could restore the stability of the EW vacuum [14]. However, the upper bound on the reheating temperature in this model prevents this phenomenon from occurring.

One may wonder whether including new physics at high energies may relax the previously mentioned constraints on the parameter space and whether this would allow higher reheating temperatures. In the present work, we will consider the previous scenario with an effective extension of the SM close to the Planck scale by including a sixth order operator. Among all operators allowed by SM gauge symmetries, we restrict ourselves to the purely Higgs-sector operator $(H^\dagger H)^3$, where H is the Higgs doublet. This operator can be viewed as a minimal proxy for new physics effects that modify the Higgs potential, neglecting operators involving derivatives or couplings to other SM fields.

New scenarios arise when considering the stability of the EW vacuum. A study of the allowed parameter space is needed to determine whether the constraints on the conditions at the onset of radiation change from the standard case. Such an analysis would reveal whether, otherwise, the model is robust under the introduction of the sixth order operator in the Higgs effective potential.

In this work, an extensive semi-analytical scanning of the parameter space will be carried out. We aim to provide a description of all the possible cosmological scenarios emerging from the inclusion of this operator in the effective theory. One of the main objectives is to characterize the regions in the parameter space leading to each scenario and to determine their cosmological viability: only those where the Universe ends up in the EW vacuum will be valid. We can already expect that an analytical treatment will not be sufficient, since the non-linearity of the field dynamics prevents us from analytically solving the equations of motion. For this purpose, 3+1-dimensional classical lattice simulations will be performed. Moreover, we aim to understand the impact of this sextic operator of the effective theory on the reheating process, which will require lattice simulations. One may anticipate that the broadening of the parameter space caused by the inclusion of this term could relax the upper bound on the reheating temperature. A new scanning of the parameter space will be carried out to determine whether the previous intuition is correct, computing afterwards this reheating temperature.

The work is structured as follows. In section II, we introduce the Lagrangian of the model and the main equations governing the evolution of the Higgs field. We also review the main results obtained in the absence of the sixth order operator. In section III, we present an exhaustive study of the parameter space and the scenarios its different regions lead to. Finally, in section IV, we perform lattice simulations to analyze the Higgs field dynamics and other cosmological observables of the model. The codes used may be found in GitHub.

II. THEORETICAL FRAMEWORK

A. Lagrangian

We are considering an early Universe where the SM particles reside in the EW vacuum, the Higgs doublet H is energetically subdominant and non-minimally coupled to gravity and there is a dominating inflaton scalar field ϕ non-coupled to SM particles or the Higgs doublet. The Higgs effective potential includes its usual SM self-coupling term $\sim (H^\dagger H)^2$ and a BSM sextic term $\sim \frac{(H^\dagger H)^3}{\Lambda^2}$ supressed at low energies. Λ is the new physics scale and cutoff of the effective field theory, above which the effective description ceases to be valid. This term corresponds to one of the dimension-six operators allowed by the SM gauge symmetries in the effective field theory expansion. The effective Lagrangian of this model is $\mathcal{L} = \frac{M_P^2}{2}R + \mathcal{L}_{\rm SM} + \mathcal{L}_\phi + \mathcal{L}_H$, where M_P is the Planck mass, $\mathcal{L}_{\rm SM}$ the SM Lagrangian without Higgs sector, $\mathcal{L}_\phi$ the inflaton Lagrangian and $\mathcal{L}_H$ the Higgs Lagrangian given by

$$\begin{aligned}\mathcal{L}_H =& g^{\mu\nu}(D_\mu H)^\dagger(D_\nu H) - \xi H^\dagger H R - \\ &-\lambda\left(H^\dagger H - \frac{v_{EW}^2}{2}\right)^2 - \frac{g^2}{\Lambda^2}\left(H^\dagger H\right)^3 .\end{aligned} \tag{1}$$

Here, v_{EW} is the value of the EW vacuum, R the Ricci scalar, ξ the non-minimal-coupling constant, λ the renormalized SM self-coupling constant and g^2 a perturbative sextic self-coupling. The last coupling is chosen positive since, in the absence of higher order self-interaction terms in the effective Lagrangian, a negative value of this coupling would lead to the appearance of a new minimum at super-Planckian field values, which would subsequently drive the potential towards arbitrarily large negative values. In the unitary gauge, $H = (0, h/\sqrt{2})^T$, assuming we are interested in energies much higher than the EW scale, $h \gg 246$ GeV, the Higgs Lagrangian can be written as

$$\mathcal{L}_h \approx \partial_\mu h \partial^\mu h - \xi R h^2 - \frac{\lambda}{4}h^4 - \frac{g^2}{8}\frac{h^6}{\Lambda^2}.$$

Using the zero-temperature renormalisation group, the renormalised Higgs quartic self-interaction λ depends on the scale μ given by $\mu^2 = h^2 + \mathcal{H}^2$. The largest Yukawa coupling of the Higgs field to SM particles is the one to the top quark. This leads to a strong dependence of the running of λ on the top quark mass, while contributions from couplings to other SM particles remain subdominant. To parametrize this dependence, we will use the fitting formula (2) derived in Ref. [13] for the three-loop $\overline{\rm MS}$ running, valid for $\mu \in [10^4{\rm GeV}, 10^{23}{\rm GeV}]$. $m_h = 125.20 \pm 0.11{\rm GeV}$ and $\alpha_s = 0,1180$ are fixed due to the small error they introduce.

$$\begin{aligned}\lambda_{fit}(\mu) =& \lambda_0 + c_1 \log\left(\frac{\mu}{\rm GeV}\right) + \\ &+ c_2 \log^2\left(\frac{\mu}{\rm GeV}\right) + c_3 \log\left(\log\left(\frac{\mu}{\rm GeV}\right)\right) .\end{aligned} \tag{2}$$

The dependence of the coefficients with the top quark mass m_t is given in Ref. [13]. It should be noted that this mass is not the top mass reconstructed using MonteCarlo, but the pole mass, whose value can differ from the former by O(1-2) GeV [15].

The Higgs effective mass depends on the Ricci scalar. In an expanding Friedmann-Lemaître-Robertson-Walker (FLRW) Universe dominated by ϕ it is $R = 3(1-3w_\phi)H^2$, with H the Hubble parameter and w_ϕ the equation of state parameter of the fluid ϕ. We will suppose the transition from inflation to kination to be instantaneous. During inflation, assuming the accelerated expansion takes place in a quasi-de Sitter Universe, $w_\phi \approx -1$, and the curvature term contributes as a large mass with $R \approx 12H^2$. When kination begins, $w_\phi = 1$, and the Ricci scalar becomes negative, $R = -6H^2$. This results in a tachyonic mass term that triggers the spontaneous breaking of Z_2 symmetry: $h = 0$ becomes a maximum and a time-dependent minimum appears at

$$h^2 = \lambda \frac{2\Lambda^2}{3g^2}\left(-1 + \sqrt{1 - \frac{6g^2}{\Lambda^2\lambda^2}\xi R}\right), \tag{3}$$

if $g \neq 0$ and, for $g = 0$, at $h_{min} = \sqrt{\frac{-2\xi R}{\lambda}}$. The former is precisely the expression obtained assuming $\Lambda^2 \gg \frac{g^2 \xi R}{\lambda^2}$ (which is natural if Λ is close to the Plank scale).

This leads to a second-order phase transition for the Higgs field, a Hubble-induced phase transition, forcing its evolution towards large expectation values. If no new physics take place, the allowed parameter space should be restricted to the region where the Higgs field resides in this minimum, even if the running of λ causes the appearance of another one. The time dependence of this vacuum expectation value ensures that the minimum shifts to smaller field values as the Hubble parameter decreases with time. It gradually approaches zero, until the symmetry is fully restored at the onset of radiation-dominated era, when $w_r = 1/3$ implies $R = 0$. During the process described above, particle production is required to fulfill the reheating temperature lower bound at the onset of radiation $\gtrsim$ 5 MeV, imposed by Big Bang Nucleosyntesis. In the next section, we discuss more in detail the dynamics of the Higgs field, the reheating process and the symmetry restoration.

B. Evolution of the Higgs field

Defining t_{kin} the time at the onset of radiation, $a_{kin} \equiv a(t_{kin})$, $H_{kin} \equiv H(t_{kin})$, $h_{kin} \equiv h(t_{kin}) \equiv \sqrt{6\xi} H_{kin}$ and under the assumption of instantaneous transition between inflation and kination, the evolution of the Hubble parameter is

$$H(t) = \frac{H_{kin}}{1 + 3H_{kin}(t - t_{kin})}. \tag{4}$$

It is useful to introduce the change of coordinates $Y = \frac{a}{a_{kin}}\frac{h}{h_{kin}}$, $\vec{y} = a_{kin}h_{kin}\vec{x}$, $z = a_{kin}h_{kin}\tau$, with τ the conformal time. The equation of motion for Y is

$$Y'' - \nabla^2 Y - M^2(z)^2 Y + \lambda Y^3 + \frac{3g^2 a_{kin}^2 h_{kin}^2}{4a^2\Lambda^2} Y^5 = 0, \tag{5}$$

where $M^2(z) = \frac{\nu^2 - 1/4}{(z+\nu)^2} = (4\nu^2 - 1)\mathcal{H}^2$ is the mass term, with the definition $\nu = \sqrt{\frac{3\xi}{2}}$ and where the conformal Hubble parameter is

$$\mathcal{H} = \frac{1}{2(z+\nu)}. \tag{6}$$

In Ref. [16], the previous equation without the term $\sim Y^5$ and $\lambda = cte > 0$ is extensively studied. In what follows, some useful results are summarized. A first approach is made neglecting the non linear term and obtaining the exact solution for the homogeneous equation. This linear regime is a valid approximation immediately after the beggining of kination. It allows to derive an approximated expression for the time evolution of the enhaced scalar modes amplitude and occupation number. As quantum fluctuations grow, the non-linear term begins to dominate the evolution and the linear approximation ceases to be valid. The backreaction time z_{br} is defined as the first time where $Y'' = 0$. It is stimated to coincide with the moment the quartic term starts to play an important role, slightly before the minimum is reached by the field. This backreaction time satisfies

$$\kappa_*^2 + 2Y^2(z_{br}) = M^2(z_{br}), \tag{7}$$

with $\kappa_* = 2\sqrt{2\nu+1}\mathcal{H}$ related to the two point correlation function obtained through the linear regime analysis. A fitting formula for the energy of the spectator field at that moment is also derived,

$$\rho_{br}(\lambda,\nu) = 16H_{kin}^4 \exp(\beta_1 + \beta_2\nu + \beta_3 \ln\nu), \tag{8}$$

with $\beta_1(\lambda) = -7.03 + 0.56\log\lambda$, $\beta_2(\lambda) = -0.06 + 0.04\log\lambda$, $\beta_3(\lambda) = 7.15 - 1.10\log\lambda$.

At the time z_{br}, the field has not reached the scales at which the sextic term becomes non-neglectable, so the previous expressions will be particularly relevant for the model considered in this work.

In the absence of sixth order contribution, a positive quartic self-interaction stabilizes the tachyonic potential. This marks the end of the exponential growth of fluctuations and the beginning of an oscillatory phase around the minimum of the effective potential. The oscillations persist as the Higgs field asymptotically approaches radiation-like equation of state. During this stage, the inflaton energy density decreases as $\rho_\phi \propto a^{-6}$ until radiation, evolving as $\rho_r \propto a^{-4}$, starts to dominate the energy budget of the Universe. This marks the onset of the radiation era. The radiation-dominated equation of state parameter, $w = 1/3$, results in a vanishing Ricci scalar $R = 0$. This restores the symmetry of the potential by removing the curvature-induced tachyonic mass term and ultimately drives the field back towards the EW vacuum. The temperature at the onset of radiation is

$$T_{ht} = \left(\frac{30\rho_{ht}}{\pi^2 g_*^{ht}}\right)^{1/4}, \tag{9}$$

and is called the reheating temperature. Here, $g_*^{ht} = 106.75$ is the SM number of relativistic degrees of freedom at energies $\gtrsim 100\,\text{GeV}$ and ρ_{ht} is the Higgs field energy density at z_{ht}, which coincides with the inflaton energy density at that time. The vast majority of particle production occurs in the tachyonic instability phase, while contribution of the oscillatory regime remains subdominant. This is the reason why the highest temperatures are reached for small λ and large H or ν values. For these parameter values, the minimum h_{min} is located at larger field values, and tachyonic instability regime lasts longer.

When the sextic term and the running of λ are introduced, different scenarios arise after back-reaction time and the following oscillatory regime is not ensured. This range of situations, which depend on the parameter space, has to be analyzed. Only the regions where this oscillatory phase around the time-dependent minimum, leading to radiation era, is guaranteed, can be regarded as viable descriptions for the history of our Universe.

III. CLASSIFICATION OF SCENARIOS

A. Higgs potential scenarios

In this section, we carry out a semi-analytic study of the parameter space. As mentioned previously, the $\lambda(\mu)$ running (2) may become negative at high energies for some values of the top quark mass. Ignoring the sixth order effects, this triggers the appearance of a barrier in the Higgs potential. The barrier is followed by a deeper stable vacuum, thereby making the time-dependent vacuum at lower energies metastable. If the Universe were to overcome the barrier, either classically or through quantum tunneling, it would end up trapped in this high energy vacuum, which is incompatible with the observed state of our Universe. At some point, quantum effects should be taken into account, as tunneling to the deep vacuum could be problematic [5]. To ensure the Universe resides today in the EW vacuum, lifetime of the metastable EW vacuum has to be longer than lifetime of the Universe. However, in this work, we will focus on characterizing the classical evolution of the field. We consider to be phenomenologically valid those scenarios where, either the EW vacuum is stable, or it is metastable but the Universe classically ends up there. We leave for future investigation the inclusion of quantum tunneling to evaluate the the stability of the classically-allowed parameter space.

We now consider the new possibilities for potential shapes arising when the sextic term is included. If its effects become significant before λ turns negative, this new term would restore the full convexity of the potential, recovering the scenario discussed in the previous section.

Conversely, if the sixth order contribution becomes dominant at higher energies, the potential still develops a new high energy vacuum, but the new term will shift it towards lower energies. Moreover, it could even reduce its depth. In the first case, one might conclude that all configurations in which the Higgs field overcomes the barrier should be discarded. However, it is worth analyzing whether thermal effects, taking into account the temperature of the Universe at this stage and the new vacuum scales, could significantly decrease the depth of the high energy vacuum. This could potentially restore the stability of the EW vacuum, allowing the Universe to tunnel back to it or even to return classically. For all these scenarios, estimating the reheating temperature will be crucial, not only for the consistency of the model itself but also to determine whether this region of the parameter space is physically valid. If the sextic term itself reduced the depth of the high energy minimum to the point of making it metastable, the Universe could return to the EW stable vacuum.

All these possible scenarios are illustrated in Fig. 1. Six different shapes of the Higgs potential are obtained by modifying m_t, H_{kin} and g, while keeping Λ and ν fixed. It is also set $H = H_{kin}$, so the potentials correspond to the beginning of kination. This is not a restriction in the sense that no other possibilities of shapes arise with the evolution of H. Nevertheless, when considering whether a given set of initial parameters is viable, it has to be taken into account that $z_{br} \sim \mathcal{O}(10)$ and therefore $H(z_{br})$ may be up to an order of magnitude smaller than H_{kin}. Consequently, the shape of the potential at that time can be completely different from the initial one.

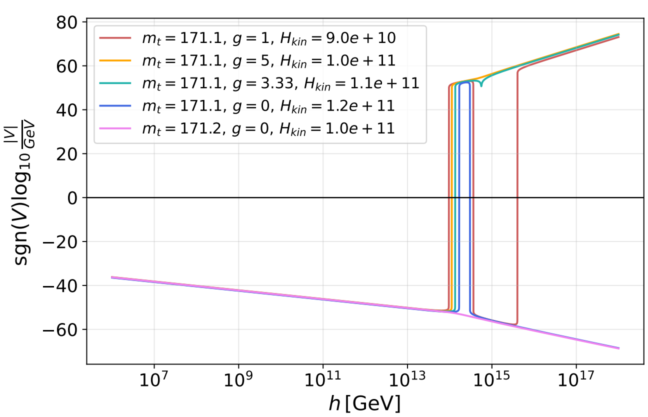

FIG. 1: Higgs potential at the onset of kination depending on m_t, g and H_{kin}. $\Lambda = 10^{17}$ GeV and $\nu = 10$ are fixed in all cases. The values of m_{top} and H_{kin} are given in GeV

.

It should be noted that, for $g \sim 1$, perturbation theory ceases to be valid. However, any change in g is equivalent to a rescaling of Λ, which is not fixed as long as Λ is up to a few orders of magnitude below Planck scale. Therefore, we will allow non-perturbative values of g. The reason behind this choice is purely practical, as it is more convenient to work with small changes in g than in Λ.

If the potentials represented above were those governing the dynamics at z_{br}, the different possible scenarios determine whether the sets of parameters are phenomenologically feasible. Regarding those in which $g = 0$, where there is no sixth order contribution, we have already analyzed both cases. It is clear that the trajectory without barrier (pink) is going to run towards arbitrarily negative values and is therefore forbidden. Meanwhile, the case in which the barrier exists but the second minimum is unbounded (blue) will depend on whether the field overcomes the barrier or not. Alternatively, the scenario in which $g \neq 0$ derives in the existence of only one minimum (orange) will always be valid in our set up. The one with a second bounded but stable minimum (red) will also depend on the overcoming of the barrier. Even if it happened, thermal contributions to its depth could be studied. Finally, in the case where the second minimum is metastable (green) and is reached by the field, quantum tunneling to the true vacuum should be analyzed. However, one may anticipate that the latter case is very fine-tuned just by observing the shape of the potential. The sextic contribution is so stiff that dominates nearly the moment it stops being negligible. This provokes the abrupt reduction of the width of the minimum when its depth decreases. It could strongly constrain the range of configurations where the Higgs field not only reaches the minimum but also gets trapped in it. This will be analyzed later through lattice simulations.

B. Analysis of the parameter space

As we have discussed above, it is crucial to understand the shape of the potential at the time z_{br} to qualitatively analyze the parameter space. To this end, we can evaluate the number of local extrema the potential presents at z_{br} for each set of parameters. In the cases with two minima, we will also evaluate if the field can classically cross the barrier. To archieve this, we compute its energy density at back-reaction time using Eq. (8) and compare it to the height of the barrier. The time z_{br} is numerically obtained by a consistent algorithm using the implicit equation (7), along with Eq (6) and Eq. (5), neglecting in the last one the term $\sim Y^5$. This is valid since the sextic term should not play any role in the dynamics at z_{br}.

Figures 2 and 3 show parameter scannings obtained by fixing three parameters in each case while varying the others. In those cases in which two vacua coexist and the first one is deeper than the second one, we distingüish two scenarios: those where the potential in the second minimum (V_2) is positive from the ones where it lies at negative values. The explanation for this choice comes from the fact that, as the first vacuum evolves towards 0, its depth decreases. Consequently, if the second minimum is negative, it eventually becomes stable and the field may be trapped there unless tunneling happens before. Therefore, those configurations with $V_2 > 0$ are always valid for our set up. Meanwhile, those with $V_2 < 0$ need further investigation on tunneling probabilities and lifetimes involved to ensure the whole Universe ends up in the correct vacuum.

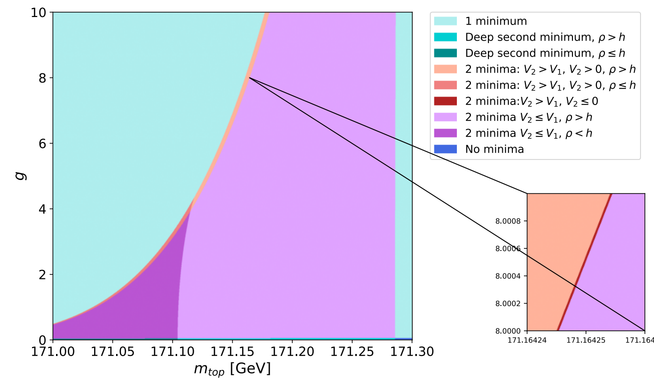

FIG. 2: Potential shape (m_{top},g) plane. The other parameters are fixed to: $H_{kin} = 10^{11}$ GeV, $\Lambda = 10^{17}$ Gev, $\nu = 10$

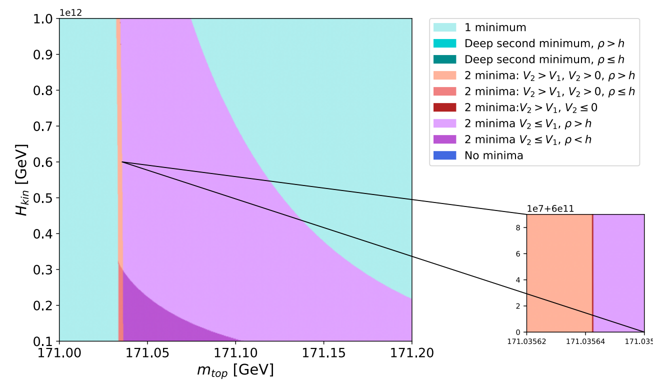

FIG. 3: Potential shape in (m_{top},H_{kin}) plane. The other parameters are fixed to: $g = 1$, $\Lambda = 10^{17}$ Gev, $\nu = 10$

It can be seen that the region with a high energy metastable vacuum and low energy stable one is relatively narrow. Nevertheless, the field overcomes the barrier in a part of it, so, for these sets of parameters, quantum tunneling to the EW vacuum seems to be possible. This would be a BSM scenario, which could never occur without the introduction of the sixth order operator to the effective Lagrangian. As we mentioned before, the existence of this region in the parameter space does not imply the feasibility of this scenario, since the time dependence of the Hubble parameter might lead the potential to evolve to another case before tunneling takes place. However, Fig. 3 shows that, once Λ, g and ν are fixed, the m_{top} values for which the previous scenario occurs depend very weakly on H_{kin}. Together with the fact that the time dependence of the potential shape is entirely encoded in the time dependence of H, this raises the question of how long such a configuration can persist. Figure 4 shows the evolution of the potential for parameters from the region of interest at different times. We can conclude that, if the field were to get trapped in the metastable minimum, it would remain there and quantum tunneling would need to be studied. The issue of whether it is possible for the field to get trapped will be analyzed in greater detail in the next section.

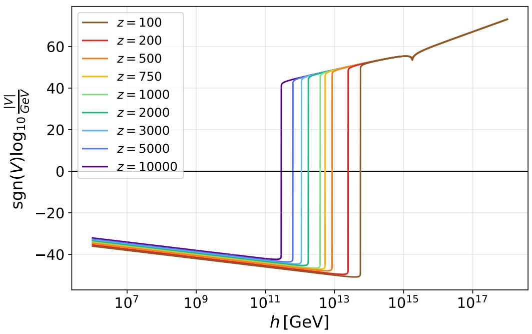

FIG. 4: Higgs potential at different times for $g = 1$, $\nu = 10$, $H_{kin} = 10^{11}$ GeV, $\Lambda = 10^{17}$ GeV, $m_{top} = 171.036$ GeV

It can also be observed in Figs. 2 and 3 that the region in which the second minimum takes negative values is extremely fine-tuned and will therefore not be regarded as a physically relevant scenario. Apart from this, we recover the different configurations discussed.

To summarize, those scenarios where the EW vacuum is stable or all those where it is not but the barrier is not crossed are valid in our set up, as they lead the Universe to reside in the EW vacuum and ultimately to the Universe we observe today. In the last case, quantum tunnelling could make them invalid if lifetime of the EW vacuum were not long enough. These correspond to a broad area in the figures (light blue, orange and dark purple). Although these scenarios are all considered acceptable possibilities for our Universe, their phenomenological implications differ from one another. Physical magnitudes such as reheating temperature, number of e-folds of kination or even gravitational wave background will not be the same for a Universe which is not energetic enough to overcome the barrier than for one which reaches a metastable high-energy vacuum and tunnels back.

Conversely, those scenarios where the high energy vacuum is stable and the barrier is crossed are not generically valid in our set up. However, the possibility of thermal corrections recovering the stability of the EW vacuum for some cases remains open. Furthermore, exceptions without taking this into account might still exist: the fact that the barrier is crossed does not necessarily imply it will not be crossed backwards. As we will illustrate later with an example, there are cases where the EW vacuum at z_{br} is metastable but it turns stable with time. These correspond to points in the light purple region in Figs. 2 and 3 which are close to the orange area.

This analysis of the parameter space allows us to qualitatively comprehend the scenarios each region might lead to. Nonetheless, as previously mentioned, the model is not so simple and we cannot precisely classify every possible situation. Moreover, it is required to derive the field dynamics from the equation of motion in order to compute physical magnitudes that differ from one valid scenario to another. However, the non-linearity of the dinamics prevent the equations of motion from being analytically solvable. Consequently, to study the field dynamics, we perform classiccal lattice simulations, which provide a way of introducing statistical fluctuations in the dynamics.

IV. LATTICE SIMULATIONS

In this section, we focus on analyzing the dynamics of the Higgs field and the reheating temperature that our model allows the Universe to reach. This is carried out through 3+1-dimensional classical lattice simulations. The main objective is to determine whether the introduction of the sextic term produces significant effects on the cosmological observables or whether the model remains otherwise robust under the new physics considered here. For this purpose, we use the code *CosmoLattice*, implementing an effective potential with the Higgs spectator field evolving in an expanding background with equation-of-state parameter $w = 1$. Regarding the numerical variables, the number of lattice points per spacial dimension is set to $N = 128$, so the length of a lattice cell is $\delta x = \frac{4\pi\nu}{N} = 0.9817$ for $\nu = 10$. The time step is set to $\delta t = 0.01$ to ensure stability ($\delta t/\delta x \ll 1/\sqrt{3}$). The smallest amplified momentum in the tachyonic band is $\kappa_{min} = \mathcal{H}_{kin}$[11], while the largest is $\kappa_{max} = \sqrt{4\nu^2 - 1}\mathcal{H}_{kin}$. Therefore, the lattice parameter κ_{IR} is set as $\mathcal{H}_{kin}$ while $\kappa_{UV} = \pi/\delta x \gg \kappa_{UV}$ to guarantee the relevant amplified modes are within the lattice range. The initial conditions for the field are stablished as $h(0) = h'(0) = 0$ and gaussian fluctuations.

A. Field distributions

In order to analyze whether the field evolves towards the second minimum or remains at the EW vacuum, we obtain snapshots of the lattice field configurations and construct the corresponding field value histograms. These allow us to identify whether the barrier is crossed in any region of the lattice. We perform, for a fixed value of m_{top}, a scan in H_{kin} across the orange region in Fig. 3 through simulations. The former is precisely the region where the high energy minimum is metastable. The purpose of this analysis is to determine whether the scenario in which part of the Universe resides in this metastable vacuum is physically viable. We take $m_{top} = 171.036$ GeV. The reason behind this choice resides in the fact that the second minimum is metastable for lower H_{kin} values within those used for Fig.3, whereas it becomes stable by increasing H_{kin}.

In Fig. 5 we observe histograms (in logarithmic scale) corresponding to values of H_{kin}, ranging from $5 \cdot 10^{11}$ GeV to $1.1 \cdot 10^{12}$ GeV in steps of 10^{11} GeV, at $z = 30$. The plot reveals that the barrier has not been crossed for $H_{kin} \leq 10^{12}$ GeV, while a part of the field overcomes the barrier for $H_{kin} = 1.1 \cdot 10^{12}$ GeV. This conclusion is not drawn exclusively from this specific time but from an analysis of histograms from various times.

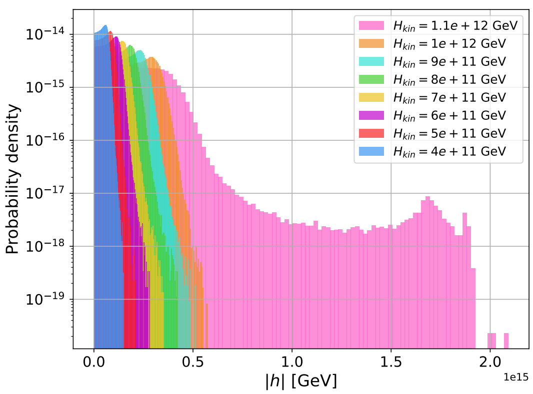

FIG. 5: Field distribution histograms for $g = 1$, $m_{top} = 171.036$ GeV, $\nu = 10$ and $\Lambda = 10^{17}$ GeV at $z = 30$

Once we have determined that the barrier is crossed for values of H_{kin} above $1.1 \cdot 10^{12}$ GeV, we can study if the field gets trapped in the high energy vacuum. Figures 6, 7 and 8 show histograms for $H_{kin} = 1.2 \cdot 10^{12}$ GeV at times $z = 70$, 100 and 300 respectively. The potential at the corresponding time is also displayed and we calculate the field value at which the barrier is located by obtaining the maximum of the potential. Subsequently, we compute the percentage of the lattice points on each side of the barrier, which is shown in the legend. At $z = 70$ (Fig. 6), we find that the field has overcome the barrier at a considerable portion of the lattice and the high energy minimum is stable. However, at $z = 100$ (Fig. 7), it has become metastable, while part of the lattice is still trapped in it. Finally, at $z = 300$ (Fig. 8), the field has returned to the EW vacuum. In this latter situation, the potential still presents a second minimum. Nonetheless, it is so shallow that the field, which is oscillating, returns classically to the EW vacuum.

This also represents an example of an scenario whose description at z_{br} does not allow to predict whether it is physically viable or not. The stability of the high energy minimum at that time and the field overcoming the barrier do not imply the field ends up trapped there. Therefore, the analysis made in the previous section to characterize the parameter space is not completely precise. Moreover, it cannot be carried out analytically, since cases like the former can only be classified as vi-

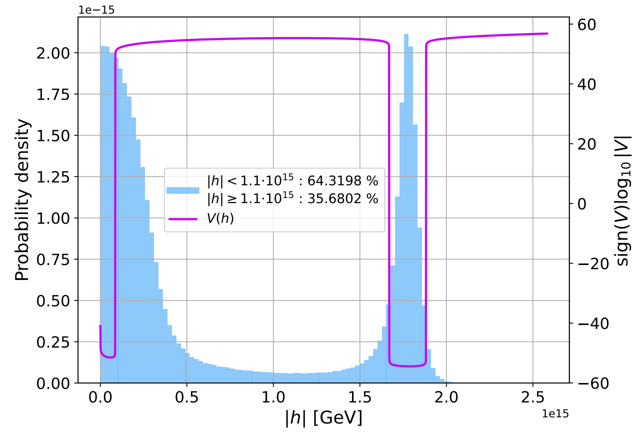

FIG. 6: Field distribution histogram (left y-axis) and Higgs potential (right y-axis) for $g = 1$, $m_{top} = 171.036$ GeV, $\nu = 10$, $\Lambda = 10^{17}$ GeV and $H_{kin} = 1.2 \cdot 10^{12}$ GeV at $z = 70$

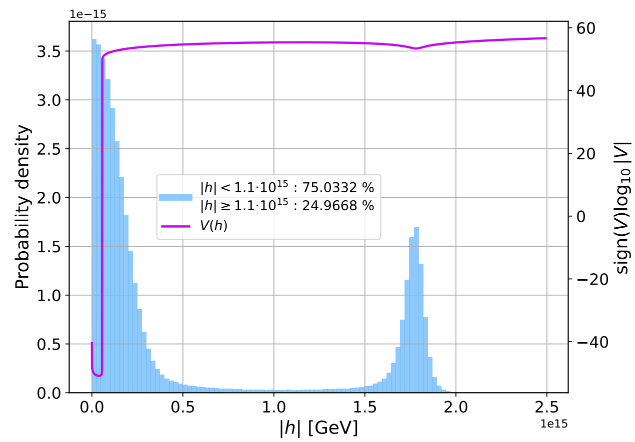

FIG. 7: Field distribution histogram (left y-axis) and Higgs potential (right y-axis) for $g = 1$, $m_{top} = 171.036$ GeV, $\nu = 10$, $\Lambda = 10^{17}$ GeV and $H_{kin} = 1.2 \cdot 10^{12}$ GeV at $z = 100$

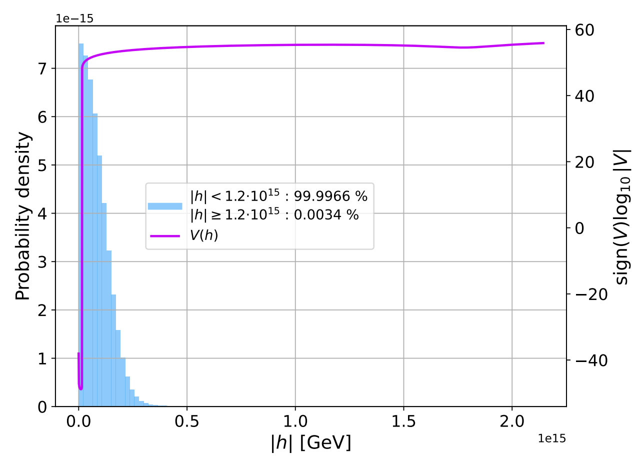

FIG. 8: Field distribution histogram (left y-axis) and Higgs potential (right y-axis) for $g = 1$, $m_{top} = 171.036$ GeV, $\nu = 10$, $\Lambda = 10^{17}$ GeV and $H_{kin} = 1.2 \cdot 10^{12}$ GeV at $z = 300$

able or non-viable through their study on the lattice.

The previous analysis can be carried out for different values of H_{kin}. If the field is more energetic at z_{br} (larger H_{kin}), its oscillations amplitude are broader and therefore it can be expected a greater probability of returning to the EW vacuum. However, when considering smaller values of H_{kin}, we could also expect the existence of an interval in which the field is sufficiently energetic to overcome the barrier but not enough to go back.

By repeating the process with values of H_{kin} between 10^{12} GeV and $H_{kin} = 1.1 \cdot 10^{12}$ GeV, in steps of 10^{10} GeV, we obtain the same results: for values above $H_{kin} = 1.07 \cdot 10^{12}$ GeV part of the lattice overcomes the barrier but the totality of it returns, whereas values below do not allow the field to cross the barrier anywhere. We may conclude that, if the field resides in a high energy metastable minimum for a set of parameters, the region corresponding to this scenario is extremely narrow. Nevertheless, the possibility of first-order phase transition through quantum tunneling to the EW vacuum in scenarios such as the one presented in Fig. 7 remains of interest. Although the field escapes classicaly the metastable vacuum within a finite time, quantum tunnelig could occur before this happens. Its study, which is left for future investigations, would require a comparison between quantum tunneling lifetimes and the timescales associated to its classical return to the EW vacuum. This purely new-physics scenario could entail a significant production of gravitational waves, which would distinguish it observationally from other viable scenarios.

B. Energy densities and reheating temperature

We previously mentioned the importance of obtaining the temperature at the onset of radiation era that our model allows the Universe to reach. It is given by (9). Since the majority of particle production takes place in the tachyonic instability phase, the highest temperatures are archived for those sets of parameters that make this phase last longer. In particular, larger values of m_{top}, ν, H_{kin}, Λ and lower values of g lead to potentials with the first minimum at higher field values. The weakest dependence is on ν and any variation of Λ can be encoded in a modification of g. Therefore, we will work for fixed values $\nu = 10$ and $\Lambda = 10^{17}$ GeV, and analyze the reheating temperature in the (g, m_{top}, H_{kin}) space.

However, for fixed values of g and Λ, a limit to the increase of m_{top}, ν and H_{kin} exists, since the field eventually crosses the barrier and becomes trapped in the high-energy stable vacuum. To determine the highest possible temperatures in those scenarios where the field classically settles in the EW vacuum, we perform a parameter scanning in (m_{top}, H_{kin}) plane, computing the position of the minimum h_{min} at z_{br}.

For each point, we run 3+1 dimensional classical lattice simulations while varying g to identify the smallest value for which the field ends up in the correct vacuum. Once the corresponding value of g has been determined for a given pair (m_{top}, H_{kin}), we compute h_{min} for that set of parameters. The results are shown in Fig. 9. The g scan is performed with resolutions $\delta g = 0.1$, $\delta m_{top} = 0.01$ GeV and $\delta H_{kin} = 10^{11}$ GeV. For intermediate points, we use the g value associated to the nearest point at larger (m_{top}, H_{kin}), ensuring that all the configurations used remain physically viable. These values of the sextic coupling are indicated on the figure. Each one is valid for the line on top of them and the space below and the left part, where no lines are shown, corresponds to $g = 0$.

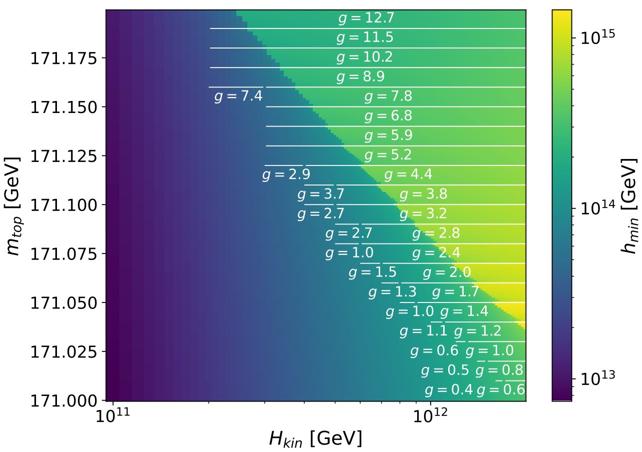

FIG. 9: Scanning of h_{min} in (m_{top}, H_{kin}) plane for the minimum physically viable value of g.

This scan enables us to draw an important conclusion: the introduction of the sextic term increases the maximum reheating temperature attainable within our model.

The strongest dependence of this temperature is the one on the value of H_{kin} and consequently, for a fixed m_{top} value, this new term broadens the allowed H_{kin} interval, leading to higher temperatures. For a given top quark mass, we find an essentially constant g, except for the point closest to the limit in some cases. We observe an abrupt change in the value of h_{min}, which does not take place directly when entering the region $g \neq 0$ but afterwards. This can be explained considering the shape of the potential. For the lowest H_{kin} that needs $g \neq 0$, the quartic term produces a barrier which is overcome by the field. Meanwhile, the sixth order term only reduces the width of the high energy minimum, which allows the field to cross the barrier backwards. Conversely, when H_{kin} continues increasing, the barrier disappears at a certain point. Whereas the minimum from the first case is determined by the quartic term, in the second case it is set by the sextic term and the quartic term contributes to the tachyonic mass. Together with the fact that the change in the shape of the potential is rather sharp within the variation of H_{kin}, this explains the abrupt change alteration of the location of the minimum. Once we enter the second region, this minimum, and therefore, the reheating temperature, becomes independent of H_{kin} and strongly dependent on g. Larger top quark masses require larger values of g, while the direct dependence on m_{top} is weaker. Consequently, the highest temperatures are reached for smaller values of m_{top}.

Having established the general behavior of the parameter space, we now turn to a more detailed analysis of the reheating temperature. We do this by performing long lattice simulations for specific parameter choices. The spectator field approximation does not allow us to obtain the energy density at reheating directly from the lattice. However, the Higgs field behaves asymptotically as radiation and we can definine z_{rad} as the time after which we consider equation-of-state parameter for the Higgs field $w_\phi \approx 1/3$. Then, it is straightforward to compute the reheating temperature. The criteria used to fix z_{rad} is similar to the one employed in Ref. [16]. We define $V = \frac{\lambda}{4}h^4 + \frac{g^2}{8}\frac{h^6}{\Lambda^2}$, $G = \frac{1}{2a}(\nabla h)^2$, $K = \frac{1}{2}h'^2$, $I = 3\xi\mathcal{H}h^2 + 6\xi\mathcal{H}hh' - \frac{\xi}{a^2}\nabla^2 h$. Here, K constitutes the kinetic energy density, V the potential without curvature term, G the gradient energy density and I the energy density interaction component associated to the curvature term of the potential. They verify $K + V + G + I = \rho_h$. The radiation-like asymptotic behaviour reflects in the energy configuration as a supression of the interaction term. This leads to an energy density $\rho \approx K + V + G$ and pression $p \approx K - \frac{1}{3}G - V$. Therefore, the condition $w_h \approx 1/3$ is equivalent to $K \approx G + 2V$. The oscillatory nature of the regime motivates us to impose the following criteria for z_{rad}: $\left\langle \frac{K(z_{rad})}{G(z_{rad})+2V(z_{rad})} \right\rangle_t = 1 \pm 0.05$, where we average over several oscillations around 1. Lattice simulations provide a direct value of $\rho_h^{rad} \equiv \rho_h(z_{rad})$. We can derive from this value the scale factor at the onset of radiation domination, a_{rh}, and $\rho_{rh} \equiv \rho_h(z_{rh}) = \rho_\phi(z_{rh})$. Since $\rho_\phi \propto a^{-6}$, $\rho_h \propto a^{-4}$ and $\rho_\phi(a_{kin}) = 3\overline{M_P}^2 H_{kin}^2$

$$\rho_{rh} = \frac{(\rho_h^{rad})^3 a_{rad}^{12}}{9\overline{M_P}^4 H_{kin}^4}, \qquad a_{rh} = \sqrt{\frac{3\overline{M_P}^2 H_{kin}^2}{\rho_h^{rad} a_{rad}^4}}. \tag{10}$$

From this figure, we can already observe what we expected from the previous study of the parameter-space: the total energy density increases with the inclusion of the sextic term, therby making it possible to have larger values of H_{kin}. The reheating temperature obtained from (10) for $H_{kin} = 1.2 \cdot 10^{11}$ GeV and $g = 0$ is $T_{rh} = 2.69 \cdot 10^{11}$GeV, while for $H_{kin} = 1.8 \cdot 10^{11}$ GeV and $g = 0.8$ we have $T_{rh} = 2.30 \cdot 10^{12}$GeV. The second which is almost an order of magnitude above the first one. Consequently, new physics introduced as a sixth order contribution to the Higgs potential allow the model to reach higher reheating temperatures compared to the SM Higgs potential. These temperatures are at least of order $\mathcal{O}(10^{11}\text{ GeV})$. However, a deeper analysis of the parameter space through long lattice simulations would be necessary to establish an upper bound. Figure 10 shows the evolution of $\langle\rho\rangle a^4$, where $\langle\cdot\rangle$ denotes the spatial average over the lattice. It is displayed for the value $m_{top} = 171.02$ GeV and three different choices of H_{kin}, using in each case the minimum value of g determined from the analysis made in Fig.9. The asymptotic radiation-like behaviour results in $\langle\rho\rangle a^4$ asymptotically approaching a constant value.

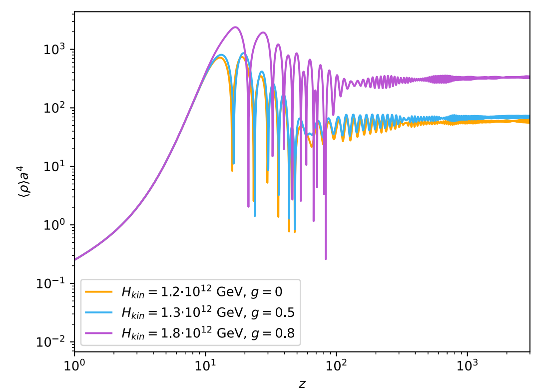

FIG. 10: Evolution of the energy density for $m_{top} = 171.02$ GeV, $\nu = 10$, $\Lambda = 10^{17}$ GeV, three different values of H_{kin} and their associated minimum physically viable value of g

Once estimated the order of magnitude of the reheating temperature, we can approach the issue of whether thermal corrections to the potential could reduce the depth of the high energy stable vacuum. These thermal corrections can be parametrized as $\Delta V(h,T) \approx 0.06h^2T^2e^{-\frac{h}{2\pi T}}$ [14] and then exponentially supressed at scales above $2\pi T$. Since the high energy vacuum is located, in all relevant cases, above $\mathcal{O}(10^{14})$ GeV, we can conclude that the temperatures attainable within the model are not sufficently high to have this effect on the Higgs potential. It is worth mentioning that the reheating temperature can only be computed through (9) in those scenarios where the field ends up in the EW vacuum. However, those situations of interest for thermal corrections imply the field trapped in the wrong vacuum. In these cases, the process to obtain thermal corrections is not straightforward, since the system is out of thermal equilibrium.

V. CONCLUSIONS

In this work, we have introduced a model based on a non-minimally-coupled-to-gravity Higgs field acting as a spectator field in a kinetic dominated era following the slow-roll phase of inflation. During a tachyonic instability phase at the onset of kination, the Higgs field is responsible for reheating the Universe. The main purpose was to study the impact of the inclusion of a sixth order operator $\sim (H^\dagger H)^3$ in the Higgs SM effective potential.

We have first established the different scenarios that emerge when considering this sextic contribution, studying their physical feasibility with the aim of determining the restrictions we should impose to the model parameters. By performing a scan of the parameter space, we have essentially found different situations depending on the number of vacua and the stability of the EW vacuum at the back-reaction time z_{br}. We have classified them as cosmologically viable scenarios if the Higgs field classically ends up in the EW vacuum, since it is the observed state of the Universe today. This is the case of potentials with only one minimum at z_{br} and all those others where a second exists but the barrier is not classically crossed. A particular case corresponding to a narrow region of the parameter space has also been found to be physically viable. It is a purely new physics scenario with a high energy metastable vacuum. It could show different cosmological properties than other viable scenarios, since quantum tunneling to the EW vacuum could occur and thereby produce graviational waves.

Other scenarios have been found which we cannot establish as cosmologically viable set ups. These correspond to all the cases where the EW vacuum is metastable and the field overcomes the barrier, getting trapped in a high-energy stable vacuum. However, we have determined these cases are not so easily classified: the sixth order contribution may allow scenarios where the field reaches this high-energy stable vacuum and returns classically to the EW one.

Motivated by the imposibility of further analysis without the study of field dynamics, we have carried out 3+1-dimensional classical lattice simulations, using the code *CosmoLattice*. These simulations were aimed at two purposes: first, to refine the previuos analysis of the scenarios across the parameter space, and second, to study the reheating process. In the first case, the evaluation of the posibility of the field trapped in a metastable high energy vacuum has shown it is a highly fine-tuned situation. The cases analyzed reveal a non-neglectable probability for the field to reach the high energy vacuum but also that it returns classically to the EW vacuum. However, this study has shown the existence of scenarios where the high energy vacuum is stable at z_{br} but turns metastable. In those cases, the field classically returns to the EW vacuum. Consequently, the parameter space cannot be scanned analytically with sufficient precision, and lattice simulations are necessary.

Regarding the reheating process, we finally have investigated the impact that the inclusion of the sextic term may have on the reheating temperature of the Universe. By performing an extensive parameter scanning, we concluded that the sixth order contribution allows the model to reach higher temperatures than the ones attainable with the SM Lagrangian. In particular, we have carried out long lattice simulations for a specific choice parameters. We have found out that the sixth order contribution may lead to temperatures at least an order of magnitude above the standard case.

This work provides several directions for future research. One important aspect is to investigate the new restrictions that quantum tunneling may impose on the classically-allowed parameter space. The classical characterization performed in this work does not ensure quantum tunneling from the EW vacuum to the high-energy stable one cannot occur. To this end, it is required to compute the lifetime of the EW vacuum and impose that it must be longer than the lifetime of the Universe. Another natural step would be to analyze in greater detail the case with a high energy metastable vacuum and the probability of a first order phase transition to the EW vacuum before the field returns classically. Regarding this scenario, it would also be worth assessing its contribution to the gravitational wave background.

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