PHYSICAL REVIEW D **93,** 103531 (2016)

Inflaton dark matter from incomplete decay

Mar Bastero-Gil,[1,*] Rafael Cerezo,[1,†] and João G. Rosa[2,‡]

[1]*Departamento de Física Teórica y del Cosmos, Universidad de Granada, Granada 18071, Spain*
[2]*Departamento de Física da Universidade de Aveiro and CIDMA, Campus de Santiago, 3810-183 Aveiro, Portugal*

(Received 4 February 2015; published 25 May 2016)

We show that the decay of the inflaton field may be incomplete, while nevertheless successfully reheating the Universe and leaving a stable remnant that accounts for the present dark matter abundance. We note, in particular, that since the mass of the inflaton decay products is field dependent, one can construct models, endowed with an appropriate discrete symmetry, where inflaton decay is kinematically forbidden at late times and only occurs during the initial stages of field oscillations after inflation. We show that this is sufficient to ensure the transition to a radiation-dominated era and that inflaton particles typically thermalize in the process. They eventually decouple and freeze out, yielding a thermal dark matter relic. We discuss possible implementations of this generic mechanism within consistent cosmological and particle physics scenarios, for both single-field and hybrid inflation.

DOI: 10.1103/PhysRevD.93.103531

I. INTRODUCTION

Inflaton and dark matter candidates in particle physics models share several common features, both being typically assumed to be weakly interacting and neutral fields. The inflaton scalar field must interact weakly with itself and other degrees of freedom in order to ensure the required flatness of the associated scalar potential, which could be spoiled by large radiative corrections [1]. Similarly, dark matter particles should form a stable nonrelativistic and nonluminous fluid at late times that accounts for the observed galaxy rotation curves [2] and the large scale structure in the Universe as inferred from cosmic microwave background [3,4] and weak-lensing observations [5,6] (for a recent review see [7]). Both inflation and dark matter are essential features in the modern cosmological paradigm and cannot be accounted for within the present framework of the Standard Model (SM) of particle physics. It is therefore interesting to consider the possibility that the same field accounts for both accelerated expansion in the early Universe and the hidden matter component at late times.

Scalar fields have the interesting property of mimicking fluids with different equations of state depending on the kinematical regime considered. For a homogeneous scalar field ϕ with potential $V(\phi)$, we have

$$\rho_\phi = \frac{1}{2}\dot{\phi}^2 + V(\phi), \qquad p_\phi = \frac{1}{2}\dot{\phi}^2 - V(\phi). \tag{1.1}$$

Hence, on the one hand a slowly varying field, $\dot{\phi}^2/2 \ll V(\phi)$ acts as an effective cosmological constant, which is the regime typically considered in canonical inflationary models. On the other hand, a field oscillating about the minimum of its potential where $V(\phi) \simeq m_\phi^2\phi^2/2$ behaves as nonrelativistic matter, with $\langle\dot{\phi}^2/2\rangle = -\langle V(\phi)\rangle$ such that $p_\phi \ll \rho_\phi$ [8]. These two regimes will generically be present in inflationary potentials, which further suggests a common framework for inflation and dark matter. The main difficulty in achieving such a unified description lies, however, in the fact that inflation must end with a transition to a radiation-dominated Universe, in order to recover the standard "big bang" evolution at least before the freeze-out of light nuclear abundances takes place [9]. An efficient transfer of energy between the inflaton field and radiation generically requires the former to decay into light degrees of freedom following the period of inflationary slow roll [10], even though other nonperturbative processes such as parametric resonance amplification could contribute significantly to the reheating process [11]. A plausible alternative is also that nonequilibrium dissipative processes continuously source a radiation bath during inflation that smoothly takes over as the field exits the slow-roll regime, a possibility generically known as warm inflation [12,13].

Nevertheless, efficient reheating does not imply the complete decay of the inflaton field and the possibility that a stable remnant is left after inflation has been considered in a few studies in the literature. In [14] the decay of the inflaton is eventually blocked through plasma effects, although there are concerns about the stability of the candidate in such scenarios. In [15] the reheating process is not efficient enough, and therefore a later period of thermal inflation is invoked. In [16] it is assumed that the sole nonperturbative decay of the inflaton is effective enough to reheat the universe, an assumption that is not clear given the current understanding of the process. In [17]

*mbg@ugr.es
†cerezo@ugr.es
‡joao.rosa@ua.pt

the scenarios for inflaton dark matter considered rely on nonminimal couplings to gravity. In [18] nonstandard kinetic terms for the scalar fields are considered, aiming to the unification of inflation, dark matter and dark energy, however the reheating process after inflation is not discussed. In this work, we propose a concrete realization of this generic idea in quantum field theory, where the decay of the inflaton is truly incomplete, occurring only for a finite period after the end of the slow-roll regime, and the properties of dark matter are independent of whether the decay is perturbative or nonperturbative.

We describe a simple mechanism where, through a standard Yukawa interaction, the inflaton can only decay into fermion pairs while it is oscillating about the minimum of its potential with a sufficiently large amplitude. At late times, the decay becomes kinematically forbidden and an appropriate discrete symmetry ensures the full stability of the inflaton particles. Moreover, we show that, for a successful reheating of the Universe, the produced fermions scatter off the inflaton particles in the oscillating scalar field and lead to its evaporation. The thermalized inflaton particles eventually decouple from the radiation bath and their abundance freezes out, similarly to other WIMP candidates [19]. We thus suggestively denote this as the "WIMPlaton" scenario.

We analyze different possible implementations of this generic mechanism, within both single-field and hybrid inflation. In the latter case, reheating is achieved through the (complete) decay of additional waterfall fields, which allows for different phenomenological realizations, but the incomplete decay of the inflaton is nevertheless required to ensure a period of radiation domination.

This work is organized as follows. In the next section we discuss the minimal model of inflaton dark matter from incomplete decay, starting with its basic properties and post-inflationary dynamics and then discussing the possibility of condensate evaporation leading to the WIMPlaton scenario. In Sec. III we discuss different phenomenological realizations of the generic mechanism that allow for the generation of SM particles in the thermal bath after inflation. We discuss how the post-inflationary dynamics and the predictions for inflaton dark matter change in the presence of an additional waterfall field in Sec. IV, and explore consistent embeddings within different inflationary scenarios in Sec. V. Finally, we summarize our main results and conclusions in Sec. VI.

II. MINIMAL MODEL

A. Basic properties and dynamics

The minimal model for inflaton dark matter considers a single dynamical (real) scalar field, the inflaton ϕ, with a potential energy $V(\phi)$ such that a period of slow roll can occur for some field range. We take the inflaton field to be coupled to fermion fields ψ_+ and ψ_- through standard Yukawa terms and impose a discrete symmetry[1] [20] $C_2 \subset \mathbb{Z}_2 \times S_2$ on the Lagrangian such that the scalar inflaton transforms under the $\mathbb{Z}_2$ group as $\phi \to -\phi$ and the fermions are simultaneously interchanged by the permutation symmetry $\psi_+ \leftrightarrow \psi_-$. This yields for the resulting Lagrangian density:

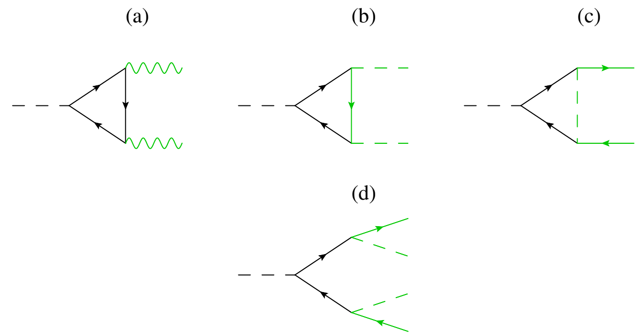

FIG. 1. Feynman diagrams for the two-body decay of the inflaton into (a) gauge bosons, (b) light scalars and (c) light fermions, induced at the one-loop level by gauge and Yukawa interactions of the $\psi_\pm$ fermions. In (d) we also show the four-body decay of the inflaton induced by the exchange of virtual $\psi_\pm$ modes with Yukawa interactions with other light species. For clarity, all light fields are represented by green lines.

$$\begin{aligned}\mathcal{L} = & \frac{1}{2}\partial^\mu\phi\partial_\mu\phi - V(|\phi|) \\ & + \bar{\psi}_+(i\gamma^\mu\partial_\mu - m_f)\psi_+ + \bar{\psi}_-(i\gamma^\mu\partial_\mu - m_f)\psi_- \\ & - h\phi\bar{\psi}_+\psi_+ + h\phi\bar{\psi}_-\psi_-, \end{aligned} \tag{2.1}$$

where we note that, as a result of the discrete symmetry, the two fermions have the same tree-level mass m_f but opposite Yukawa coupling to the inflaton field. The action of the discrete symmetry is restricted to the inflaton-fermion sector, such that all other fields, including the SM fields, are invariant under this symmetry.

We assume that the minimum of the potential lies at $\phi = 0$, where the discrete symmetry is preserved and only the Yukawa terms above can lead to inflaton decay, since no other linear terms are allowed. If $m_f > m_\phi/2$, where m_ϕ denotes the inflaton mass at the minimum, these decays are kinematically blocked and the inflaton is completely stable. Note that, in the absence of the discrete symmetry, inflaton decay could proceed through off-shell fermion modes, as illustrated with a few examples in Fig. 1. However, the contributions of ψ_+ and ψ_- cancel out if the discrete symmetry is exact. The stable inflaton particles may then contribute to the present dark matter abundance.

Away from the origin, however, the discrete symmetry is broken, leading to a mass splitting of the fermions:

[1] This subgroup contains only elements that transform simultaneously under $\mathbb{Z}_2$ and S_2.

$$m_{\pm} = |m_f \pm h\phi|. \qquad (2.2)$$

Hence, although the inflaton is stable at late times, it may decay into the fermions while it is oscillating about the minimum following the slow-roll period, namely for field values satisfying

$$|m_f \pm h\phi| < m_\phi/2. \qquad (2.3)$$

This implies, in particular, that decay is kinematically allowed for field amplitudes $|\phi| \gtrsim m_f/h$. The partial decay widths associated to the two fermionic decay channels are then given by

$$\Gamma_{\pm} = \frac{h^2}{8\pi} m_\phi \left(1 - \frac{4m_{\pm}^2}{m_\phi^2}\right)^{3/2}, \qquad (2.4)$$

with $\Gamma_\phi = \Gamma_+ + \Gamma_-$. Note that due to the opposite sign of the Yukawa couplings, the inflaton will alternately decay into each fermion species as it oscillates between negative and positive values. The inflaton equation of motion as it oscillates about the minimum of its potential is then given by

$$\ddot{\phi} + 3H\dot{\phi} + \Gamma_\phi \dot{\phi} + m_\phi^2 \phi = 0, \qquad (2.5)$$

which, upon multiplying by $\dot{\phi}$ is equivalent to

$$\dot{\rho}_\phi + 3H(\rho_\phi + p_\phi) = -\Gamma_\phi \dot{\phi}^2. \qquad (2.6)$$

The term on the right-hand side gives the rate at which energy density is transferred from the inflaton field into the fermions $\psi_\pm$. Let us assume that the fermions quickly thermalize, an assumption that we will check *a posteriori*, in the process exciting g_* relativistic degrees of freedom and forming a radiation bath at temperature T, with energy density $\rho_R = (\pi^2/30) g_* T^4$. Note that since the fermion masses oscillate due to the varying inflaton field, they only contribute periodically to the number of relativistic degrees of freedom. These may also include the inflaton for $T \gtrsim m_\phi$ and other species such as the SM particles. Note that the latter must be excited before the cosmological synthesis of light nuclear elements takes place, as we discuss below in more detail. For simplicity, we consider a fixed value of g_*, which is not a bad approximation since, as we will show later on, our results exhibit only a mild dependence on this parameter.

Energy conservation then implies that the energy lost by the inflaton field in Eq. (2.6) is gained by the radiation bath, which then follows the dynamical equation:

$$\dot{\rho}_R + 4H\rho_R = \Gamma_\phi \dot{\phi}^2. \qquad (2.7)$$

Since both the inflaton and the radiation contribute to the energy density in the Universe, we may write the Friedmann equation as

$$H^2 = \frac{\rho_\phi + \rho_R}{3M_P^2} = \frac{\frac{1}{2}\dot{\phi}^2 + \frac{1}{2}m_\phi^2\phi^2 + \rho_R}{3M_P^2}. \qquad (2.8)$$

Equations (2.5), (2.7) and (2.8) then form a complete set of differential equations that can be solved for given choices of the parameters (m_ϕ, m_f, h) and initial conditions. The nonlinearity of the equations precludes a complete analytical description of the problem, although as we discuss below a few simple analytical estimates can be performed, but it is straightforward to integrate them numerically. We show in Fig. 2 an example of the results obtained for a given choice of parameters that illustrates the main dynamical features.

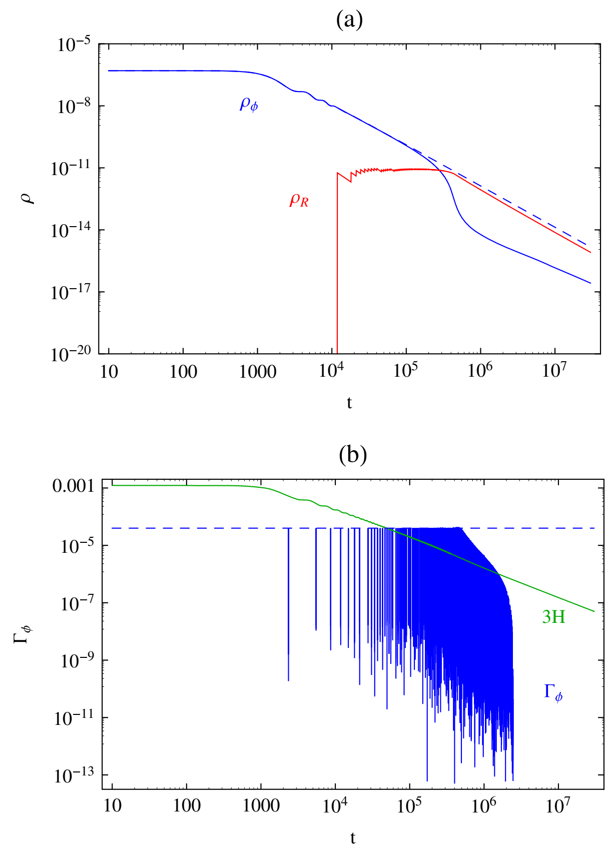

FIG. 2. Results of the numerical integration of the inflaton-radiation dynamical equations for $m_\phi = 10^{-3} M_P$, $m_f = 0.55 m_\phi$ and $h = 1$, showing the time evolution of (a) the inflaton (solid blue curve) and radiation (solid red curve) energy densities; and (b) the inflaton decay width (solid blue curve) compared to the Hubble parameter (solid green curve). The blue dashed curves in (a) and (b) give the evolution of the inflaton energy density in the absence of decay and the maximum value of the decay width, respectively. All quantities are given in Planck units such that $M_P = 1$.

As one can see in this figure, the inflaton field begins to oscillate about the origin with frequency m_ϕ for $t \sim m_\phi^{-1}$ and behaves initially as cold dark matter, $\rho_\phi \propto a^{-3}$, being the dominant energy component such that $a \propto t^{2/3}$ and $H = 2/3t$. While decay is blocked before the onset of oscillations, since $m_\pm \sim h|\phi| \gg m_\phi$, it becomes kinematically allowed as soon as the field goes through the origin. Since in this example the amplitude of the field oscillations largely exceeds the tree-level masses, with $h\Phi \gg m_f$, decay occurs for two narrow field ranges close to and on both sides of the origin. The decay width then corresponds to the series of periodic narrow peaks shown in Fig. 2(b), with maximum value $\Gamma_\phi^{\max} = (h^2/8\pi)m_\phi$. While initially $\Gamma_\phi^{\max} \ll 3H$, such that the inflaton's energy density remains essentially unaffected by the decay into fermions, the source term in the radiation equation quickly becomes significant, leading to a jump in the value of the radiation energy density. The latter remains approximately constant until the inflaton's energy density is sufficiently redshifted. When they become comparable in magnitude, the inflaton effectively decays and radiation takes over as the dominant component. As the field amplitude decreases, the maximum decay width becomes progressively smaller until decay is finally blocked. The oscillating inflaton then becomes stable and once more behaves as cold dark matter, eventually taking over the radiation as the dominant component at later times.

A peculiar dynamical feature of the evolution is the approximate constancy of the radiation energy density achieved just after the first few oscillations. This is inherent to the fact that inflaton decay occurs in short bursts in each oscillation, which does not occur if the decay were allowed for all field values. Before the effect of decay into fermions becomes significant, for $t \gtrsim m_\phi^{-1}$ the inflaton behaves as a damped harmonic oscillator with

$$\phi(t) \simeq \Phi(t)\sin(m_\phi t + \alpha_\phi), \qquad \Phi(t) = \sqrt{\frac{8}{3}}\frac{M_P}{m_\phi}\frac{1}{t}, \tag{2.9}$$

where α_ϕ is a phase depending on the initial field and velocity values. For $h\Phi \gg m_f$, one can easily see that decay into each fermion is allowed during a short period $\tau_d \simeq (h\Phi)^{-1} \ll 2\pi/m_\phi$ as the field goes through the origin, which occurs twice every oscillation period. Since the average decay width in this interval can be taken as $\Gamma_\phi^{\max}/2$ and the field velocity is $\dot\phi \simeq m_\phi\Phi$, every half period the radiation energy density increases due to inflaton decay by an amount:

$$\Delta\rho_R^{\rm decay} \simeq \frac{\Gamma_\phi^{\max}}{2}(m_\phi\Phi)^2(2\tau_d) \simeq \frac{h}{8\pi}m_\phi^3\Phi, \tag{2.10}$$

where we have taken into account the decays into both ψ_+ and ψ_-. When decay is forbidden, radiation simply redshifts with expansion, which counteracts the enhancement due to decay by an amount:

$$\Delta\rho_R^{\rm Hubble} \simeq -4H\rho_R(\pi/m_\phi). \tag{2.11}$$

Since $H = 2/3t = (m_\phi/\sqrt{6}M_P)\Phi$, it is easy to see that the amount of radiation produced by inflaton decay is eventually compensated by Hubble expansion to yield a constant energy density. Equating $\Delta\rho_R^{\rm decay} = -\Delta\rho_R^{\rm Hubble}$ then gives

$$\rho_R \simeq \frac{\sqrt{6}h}{32\pi^2}m_\phi^3 M_P, \tag{2.12}$$

which is in very good agreement with the numerical simulations. Noting that, from Eq. (2.7), $\dot\rho_R \simeq 0$ implies $4H\rho_R \simeq \Gamma_\phi\dot\phi^2 \simeq \Gamma_\phi\rho_\phi$, we see that $\Gamma_\phi \sim 3H$ for $\rho_R \sim \rho_\phi$, so that as observed numerically the inflaton energy density is only reduced significantly when it becomes comparable to the radiation energy density.

From Eq. (2.12) we can determine the temperature of the radiation bath, which remains approximately constant up to inflaton-radiation equality and thus constitutes the reheating temperature:

$$\begin{aligned} T_R &\simeq \left(\frac{15\sqrt{6}}{16\pi^4}\right)^{1/4} g_*^{-1/4} h^{1/4}\left(\frac{m_\phi}{M_P}\right)^{3/4} M_P \\ &\simeq 2.7\times 10^6 g_*^{-1/4} h^{1/4}\left(\frac{m_\phi}{1\ \mathrm{TeV}}\right)^{3/4}\ \mathrm{GeV}. \end{aligned} \tag{2.13}$$

From this we conclude, in particular, that the reheating temperature is generically larger than the inflaton mass for

$$m_\phi < 5.6\times 10^{16}\frac{h}{g_*}\ \mathrm{GeV}. \tag{2.14}$$

B. Condensate evaporation: The WIMPlaton scenario

Although the oscillating inflaton field becomes stable at late times and could thus account for dark matter in the right parametric regimes, the analysis above neglects the important effect of interactions between the field and the fermions resulting from its decay.

The classical inflaton field corresponds to a collective state of zero-momentum scalar bosons, assuming that no large field inhomogeneities are formed at the end of the slow-roll inflationary regime. Inflaton particles in this condensate can interact with the fermions that result from its decay and, in particular, these fermions can scatter some of the bosons out of the condensate and promote them to higher-momentum states that become part of the thermal bath. These correspond to scattering processes $\psi_\pm\langle\phi\rangle \to \psi_\pm\phi$, where we denote by $\langle\phi\rangle$ and ϕ scalar particles in the zero-momentum condensate and in higher-momentum modes, respectively, and which are mediated through both s- and t-channel fermion exchange. Moreover, these processes may occur as soon as the field begins oscillating and decay into $\psi_\pm$ becomes kinematically allowed, potentially leading to the evaporation of the condensate and the transfer of the inflaton particle number into the thermal bath.

As we have seen earlier, soon after the onset of inflaton oscillations, the temperature of the thermal bath rises sharply to a value that remains approximately constant until inflaton-radiation equality and that corresponds to the reheating temperature in Eq. (2.13). Assuming that this temperature exceeds the inflaton and fermion (tree-level) masses and that local thermal equilibrium is quickly achieved, we may then take the phase-space distributions for inflaton and fermion species in the thermal bath to be the relativistic Bose-Einstein and Fermi-Dirac distributions, respectively. Taking into account the above scatterings and the inverse processes, the net condensate evaporation rate is given by [21]

$$\Gamma_{\rm evap} = \frac{1}{n_\phi}\int\prod_{i=1}^{4}\frac{d^3\mathbf{p}_i}{(2\pi)^3 2E_i}(2\pi)^4\delta^4(p_1+p_2-p_3-p_4) \times|\mathcal{M}|^2[f_1f_2(1+f_3)(1-f_4) - f_3f_4(1+f_1)(1-f_2)], \quad (2.15)$$

where $\mathcal{M}$ is the scattering amplitude for $\langle\phi\rangle(p_1)\psi_\pm(p_2)\leftrightarrow\phi(p_3)\psi_\pm(p_4)$ and f_i the corresponding phase-space distribution factors. Since the condensate is inherently characterized by large occupation numbers $f_1\gg 1$, we obtain to leading order for $T\gg m_\phi, m_\pm$:

$$\Gamma_{\rm evap}\simeq\frac{h^4}{12\pi^3}\left(1+\log\left(\frac{T}{m_\phi}\right)\right)T. \quad (2.16)$$

In the radiation era, $H\simeq(\pi/\sqrt{90})g_*^{1/2}T^2/M_P$, such that the evaporation rate should eventually catch up with Hubble expansion. Note, however, that once the inflaton or fermions become nonrelativistic, this rate becomes Boltzmann suppressed. Condensate evaporation will then occur for $T\gtrsim m_\phi\sim m_f$ provided that

$$\left.\frac{\Gamma_{\rm evap}}{H}\right|_{T=m_\phi}\simeq 10^{13}h^4 g_*^{-1/2}\left(\frac{1\ \text{TeV}}{m_\phi}\right)\gtrsim 1. \quad (2.17)$$

The condensate evaporation rate is comparable to the rate of other four-body processes, such as inflaton or fermion annihilation, so that for sufficiently large Yukawa couplings we expect the condensate to be destroyed and lead to a thermalized bath of inflaton and fermion particles, as well as additional species. Evaporation thus increases the radiation energy density and hence its temperature up to a maximum value:

$$T_R^{\rm max} = \left(\frac{90}{\pi^2}\right)^{1/4}g_*^{-1/4}\sqrt{M_P m_\phi} \simeq 8.5\times10^{10}g_*^{-1/4}\left(\frac{m_\phi}{1\ \text{TeV}}\right)^{1/2}\ \text{GeV}, \quad (2.18)$$

which corresponds to the limiting case where evaporation occurs already at the onset of field oscillations.

Inflaton particles will be kept in thermal equilibrium through annihilation and elastic scattering processes. Once these become inefficient, the abundance of inflaton particles will freeze out, as for other conventional WIMP dark matter candidates. Assuming this occurs when both the inflaton and the fermions are nonrelativistic, the relevant (fermion t-channel) annihilation cross section is given by

$$\sigma_{\phi\phi}\simeq\frac{h^4}{8\pi m_\phi^2}, \quad (2.19)$$

which is independent of the fermion mass in this limit. Following the standard calculation for the thermal relic abundance of a decoupled nonrelativistic species, we obtain for the inflaton mass

$$m_\phi\simeq 1.4h^2\left(\frac{\Omega_{\phi0}h_0^2}{0.1}\right)^{1/2}\left(\frac{g_{*F}}{10}\right)^{1/4}\left(\frac{x_F}{25}\right)^{-3/4}\text{TeV}, \quad (2.20)$$

where g_{*F} denotes the number of relativistic degrees of freedom at freeze-out and $x_F = m_\phi/T_F$, with T_F denoting the freeze-out temperature.

We can now perform a consistency check of our analysis. Combining Eqs. (2.17) and (2.20), the oscillating inflaton field will evaporate and yield the correct dark matter relic abundance for

$$h\gtrsim 10^{-7}\left(\frac{\Omega_{\phi0}h_0^2}{0.1}\right)^{1/4}\left(\frac{g_{*F}}{10}\right)^{3/8}\left(\frac{x_F}{25}\right)^{-3/8}. \quad (2.21)$$

For smaller values of the Yukawa coupling, the oscillating inflaton condensate could in principle survive, but this would yield a too low reheating temperature ($\lesssim$100 MeV) for big bang nucleosynthesis (BBN) to take place [22], according to Eq. (2.13). Although a more detailed analysis of the Boltzmann equation determining the evolution of the inflaton condensate may be required, this estimate indicates that in the physically interesting parameter range, where the inflaton could consistently account for the present dark matter abundance while satisfying the BBN constraint, condensate evaporation is most likely inevitable.

This then gives us a more realistic dynamical picture of what we suggestively denote as the WIMPlaton scenario. After inflation, the scalar inflaton begins oscillating about the minimum of its potential, decaying into fermions in short bursts every oscillation. These may thermalize and excite other degrees of freedom in the plasma, and scatter off the inflaton particles in the condensate, leading to its evaporation. Both decay and evaporation increase the relative abundance of radiation and decrease the amplitude of oscillations, until eventually radiation becomes dominant and inflaton decay is no longer kinematically allowed. The stable inflaton particles remain in thermal equilibrium until the temperature drops below their mass and they decouple from the plasma, their frozen abundance yielding the inferred dark matter component of our present Universe.

Thanks to the separation of scales of the evaporation and decoupling processes, the properties of the dark matter candidate depend on the latter rather than on the details of the inflaton decay. The crucial point in the mechanism is that some decay of the inflaton is allowed to trigger evaporation, and that the decay will be blocked eventually, leading to a stable candidate. In the above analysis we have taken the inflaton decay to occur "perturbatively," through the two-body decay into fermion pairs. It may, however, be the case that nonperturbative effects are dominant during the first stages of inflaton oscillations. Preheating will produce fermions with a nonequilibrium number density n_f [23–25]. Similarly to our discussion above, these fermions will scatter off the inflaton particles in the oscillating condensate, and from Eq. (2.15) we obtain for the evaporation rate

$$\Gamma_{\rm evap} \simeq \langle\sigma v\rangle n_f \sim \frac{h^4 n_f}{m_f^2}, \qquad (2.22)$$

where we have approximated the averaged scattering cross section by $\langle\sigma v\rangle \sim h^4/m_f^2$. Fermions are produced efficiently during preheating up to a maximum comoving wave number k_s, with the distribution rapidly falling exponentially for $k > k_s$ and $n_f(k < k_s) \simeq 1/2$ [23]. The value of k_s depends on the resonance parameter $q(t)$:

$$q(t) = \frac{h^2\Phi(t)^2}{\omega_\phi^2}, \qquad (2.23)$$

where ω_ϕ is the frequency and $\Phi(t) \simeq \Phi_0/a(t)^{3/2}$ the amplitude of the inflaton condensate oscillations at the end of inflation (the subscript "0" denotes initial values for preheating). Notice that, for realistic models of inflation consistent with observations (see Sec. V), the frequency ω_ϕ can be typically much larger that the inflaton mass $m_\phi \sim m_f \sim O(1\ \text{TeV})$. Therefore, we will have preheating of "light" fermions, with modes being excited for

$$\frac{k^2}{\omega_\phi^2} \leq \frac{a(t)^2 q(t)^{1/2}}{\pi}. \qquad (2.24)$$

Once the resonance parameter $q(t)$ becomes of order unity, parametric resonance continues in the narrow resonance regime, which is equivalent to the perturbative regime [26]. That is, from that point onwards, our perturbative calculation holds. For the fermions produced during the broad resonance, the final value of k_s is given:

$$\left(\frac{k_s}{a_{\rm end}}\right)^2 = \frac{\omega_\phi^2}{\pi}, \qquad (2.25)$$

where $a_{\rm end}$ is the scale factor when $q(t_{\rm end}) \simeq 1$. We can then easily get an estimation for the number density of fermions at the end of the broad parametric resonance:

$$n_f = \frac{1}{2\pi^2 a^3}\int dk k^2 n_f(k) \sim \frac{\omega_\phi^3}{4\pi^3\sqrt{\pi}}, \qquad (2.26)$$

and the evaporation rate:

$$\Gamma_{\rm evap} \sim \frac{h^4}{m_f^2}\frac{\omega_\phi^3}{4\pi^3\sqrt{\pi}}. \qquad (2.27)$$

We need to compare this rate with the Hubble parameter at this time:

$$H = \frac{\rho_\phi^{1/2}}{\sqrt{3}M_P} \simeq \frac{\omega_\phi^2/h}{\sqrt{3}M_P}, \qquad (2.28)$$

and therefore

$$\left.\frac{\Gamma_{\rm evap}}{H}\right|_{\rm preh} \sim \frac{h^5}{4\pi^3}\sqrt{\frac{3}{\pi}}\frac{\omega_\phi M_P}{m_\phi^2} \sim 4.5\times 10^{28} h^5 \left(\frac{\omega_\phi}{M_P}\right)\left(\frac{1\ \text{TeV}}{m_\phi}\right)^2. \qquad (2.29)$$

For example, for typical values at the end of chaotic inflation, $\Phi_0 \sim M_P$, $\omega_\phi \sim \lambda^{1/2} M_P \sim 10^{-7} M_P$, we have

$$\left.\frac{\Gamma_{\rm evap}}{H}\right|_{\rm preh} \sim h^5 \times 10^{21}\left(\frac{1\ \text{TeV}}{m_f}\right)^2, \qquad (2.30)$$

which is larger than one unless the Yukawa coupling is $h \lesssim 10^{-4}$. In the opposite limit, $\omega_\phi \simeq 1$ TeV, the evaporation rate at the end of preheating is large enough for $h \geq 5\times 10^{-3}$.

We thus see that also in the case of nonperturbative inflaton decay we expect the oscillating condensate to evaporate unless the Yukawa coupling is very suppressed, generically leading to the WIMPlaton scenario, with the details of the dark matter particle being determined by the process of decoupling described above.

III. REHEATING THE STANDARD MODEL

In the previous sections we have described the basic generic features of the WIMPlaton scenario. We will now turn to explicit particle physics realizations of this mechanism, exploring different possibilities for the connection between the inflaton-fermion sector and the SM degrees of freedom.

Quarks, leptons and gauge bosons must be generated in the thermal bath produced by inflaton decay and evaporation before the temperature drops below 100 MeV, so that the successful synthesis of light elements may take place. Since the inflaton is typically taken as a gauge singlet, the structure of the Yukawa interactions implies that $\psi_\pm$ must be nonchiral fermions (either Dirac or Majorana), as opposed to the known SM fermions, being thus unlikely that they are explicitly charged under the SM gauge group. There are nevertheless several possibilities for exciting the SM degrees of freedom in the plasma either before or after radiation comes to dominate. In the following we will consider two different possibilities for the nature of this hidden fermion sector and we will also consider the related possibility that inflaton dark matter is only metastable.

A. Fermion decay

The simplest possibility is perhaps that of unstable $\psi_\pm$ fermions decaying into a light scalar and a light fermion, for which the decay width is

$$\Gamma_{\psi_\pm} = \frac{h_f^2}{16\pi} m_\pm, \tag{3.1}$$

where h_f denotes the associated coupling constant. Note that this is computed in the fermions' rest frame, whereas in the plasma's frame an additional time dilation factor $m_\pm/T$ suppresses the decay for relativistic fermions. For a significant part of the energy density in radiation to be transferred into these degrees of freedom, the fermions must decay before they become nonrelativistic. Thus, requiring $\Gamma_{\psi_\pm} \gtrsim H$ for $T \gtrsim m_f$ in the radiation era (where $m_\pm \simeq m_f$), we obtain the following bound on the coupling:

$$h_f \gtrsim 8 \times 10^{-8} g_{*f}^{1/4} \left(\frac{m_f}{1\ \mathrm{TeV}}\right)^{1/2}, \tag{3.2}$$

where g_{*f} is the number of relativistic degrees of freedom at $T = m_f$. In addition, for this to happen before BBN we require $m_f \gtrsim 100$ MeV. If the inflaton and fermion masses are comparable, this corresponds to $h \gtrsim 0.01$ according to Eq. (2.20). These light degrees of freedom may correspond to SM particles if, for example, the fermions coupled to the inflaton field correspond to a pair of degenerate sterile neutrinos, which are singlets under the SM gauge group and may decay into a Higgs-lepton pair through Yukawa terms of the form $h_f H \bar{l} \psi_\pm$. Note that this requires $m_f > m_H = 125$ GeV [27] and hence $m_\phi < 250$ GeV, which is compatible with the present dark matter abundance for couplings $h \lesssim 0.7$ from Eq. (2.20).

B. Mili-charged fermions

If the fermions are stable, another possibility for reheating the SM degrees of freedom is through efficient annihilation. A possible scenario is for the fermions to be charged under a hidden $U(1)_X$ gauge group, which may mediate fermion scatterings and thus improve the thermalization efficiency. This $U(1)_X$ hidden photon may be kinetically mixed with the SM photon or hypercharge gauge boson Y^μ through a term of the form $F_X^{\mu\nu} F_{\mu\nu}^Y$, which may be generated radiatively if there are fields charged under both gauge groups or simply via gravitational interactions, as happens e.g. in string theory. Diagonalization of the gauge kinetic terms then induces a small electric charge for the fermions $\psi_\pm$, such that they may annihilate into SM charged particles via s-channel photon exchange, $\psi_\pm \psi_\pm \to \gamma \to qq,\ ll$. At high temperature the annihilation cross section is given by the Thomson scattering formula and the corresponding interaction rate for relativistic species is then given by

$$\Gamma_{th} \simeq \frac{4\zeta(3)}{\pi} \epsilon^2 \alpha^2 N_{\mathrm{ch}} T, \tag{3.3}$$

where ϵ is the "mili-charge" of the fermions $\psi_\pm$, α is the fine-structure constant and N_{ch} is the effective number of charged species in the final state, which for the full SM is $N_{\mathrm{ch}} = 20/3$. In the radiation era, since $H \propto T^2$ annihilation becomes more efficient at smaller temperatures. Once $T \lesssim m_f$, however, the fermion abundance becomes Boltzmann suppressed and annihilations can no longer be efficient. Thus, requiring that SM species are excited before $\psi_\pm$ become nonrelativistic, we obtain the following bound on the mili-charge:

$$\epsilon \gtrsim 5 \times 10^{-7} g_{*f}^{1/4} \left(\frac{N_{\mathrm{ch}}}{20/3}\right)^{-1/2} \left(\frac{\alpha^{-1}}{128}\right) \left(\frac{m_f}{1\ \mathrm{TeV}}\right)^{1/2}. \tag{3.4}$$

This bound is not very stringent for fermion masses in the GeV–TeV range, where the main constraints come from (i) direct collider searches (including the LHC) yielding $\epsilon \lesssim 0.1$ for $1\ \mathrm{GeV} \lesssim m_f \lesssim \mathrm{few} \times 100$ GeV, and (ii) indirect bounds from the CMB anisotropy spectrum, which yield $\epsilon \lesssim 10^{-4}$ for few $\times$ 100 GeV $\lesssim m_f \lesssim$ few $\times$ TeV based on the effect of mili-charged particles on the baryon-photon oscillations (for a gauge coupling $g_X = 0.1$). Note that for masses above the TeV range, mili-charged particles may give a too large contribution to the dark matter abundance in the Universe, and more stringent bounds on ϵ apply in this case (see [28,29] and references therein). This thus constitutes a promising scenario with potential for experimental probing in the near future.

C. Metastable inflaton

The discrete $C_2 \subset \mathbb{Z}_2 \times S_2$ symmetry protects the inflaton from decaying at late times, thus constituting a viable dark matter candidate. One can consider, however, scenarios where this symmetry is broken and the inflaton is only metastable, with a lifetime larger than the age of the Universe, $t_0 \sim 14$ Gyrs. As an example, we have considered the case where $\psi_\pm$ are unstable, decaying into a light fermion and scalar, which induce diagrams (b)–(d) in Fig. 1. Since these processes have comparable magnitudes, we have computed the inflaton decay width for the one-loop process (b), where it decays into two light scalars. For concreteness we consider the case where $\psi_\pm$ are slightly nondegenerate with $m_- = m_f$ and $m_+ = m_f(1+\Delta)$, for $\Delta \ll 1$. This gives for $m_\phi = 2m_f$:

$$\Gamma_\phi^{(b)} \simeq \frac{h^2 h_f^4}{64\pi^5} m_\phi \Delta^2, \tag{3.5}$$

which is more suppressed for larger values of m_f.

For m_ϕ yielding the correct relic abundance to account for dark matter and using the lower bound on h_f obtained in Eq. (3.2), we obtain the following upper bound on the inflaton lifetime:

$$\tau_\phi < 15.7 g_{*f}^{-1} \left(\frac{m_\phi}{10\ \mathrm{GeV}}\right)^{-4} \left(\frac{\Delta}{0.01}\right)^{-2}\ \mathrm{Gyrs}, \tag{3.6}$$

where we have taken the reference values for the present dark matter abundance and freeze-out parameters in Eq. (2.20). Hence, we conclude that significant violations of the discrete symmetry can still yield a sufficiently long-lived inflaton for 100 MeV $\lesssim m_\phi \lesssim$ 10 GeV, such that the fermions $\psi_\pm$ decay before BBN and while they are still relativistic. For heavier inflaton particles, values of Δ below the percent level are required, signaling that the discrete symmetry most hold to a high degree of accuracy in this regime.

It is thus clear from the examples above that if the inflaton field can only decay to the $\psi_\pm$ fermions for a finite period following the end of the slow-roll regime, becoming (meta)stable at late times, it may account for the dark matter in the Universe while allowing for successful reheating of the SM particles and setting the necessary conditions for BBN.

IV. HYBRID MODEL

A. Basic properties and dynamics

An alternative framework to the one considered in the previous sections is the supersymmetric hybrid model [30]. In this scenario, the inflaton decay products need not include or interact with the SM degrees of freedom, since the additional waterfall sector can be responsible for reheating after inflation [31,32].

For the same reasons exposed in the minimal model, the symmetry $C_2 \subset \mathbb{Z}_2 \times S_2$ is imposed on the superfield containing the inflaton and all the superfields that it couples directly to. Hence, for the superpotential to be invariant under the action of this group, the inflaton is, as before, coupled to a pair of superfields $Y_\pm$ which contain the fermions $\psi_\pm$ with masses m_f, and to a waterfall sector with a pair of superfields $X_\pm$. The group C_2 simultaneously changes $\Phi \to -\Phi$ and interchanges the superfields $Y_+ \leftrightarrow Y_-$ and $X_+ \leftrightarrow X_-$.

The discrete symmetry forbids the linear term in the superpotential that is typically considered in SUSY hybrid inflation to generate the constant vacuum energy driving accelerated expansion. This may nevertheless be generated either by a D-term contribution [33], or through a non-vanishing F-term coming from a SUSY breaking sector [34]. In addition, in order to ensure the stability of the inflaton at late times, so that it may account for dark matter, the C_2 symmetry must be preserved in the ground state, implying equal vacuum expectation values for the scalar components of both waterfall fields. One possibility to satisfy this condition and simultaneously generate a constant vacuum energy is to introduce an additional "driving" superfield, Z, which is not charged under the discrete symmetry and is coupled to the waterfall sector, along the lines proposed in [35]. We thus consider a superpotential of the form

$$\begin{aligned} W = {} & \frac{g}{2}\Phi(X_+^2 - X_-^2) + \frac{h}{2}\Phi(Y_+^2 - Y_-^2) + \frac{m_f}{2}(Y_+^2 + Y_-^2) \\ & + \frac{\kappa}{2} Z(X_+^2 + X_-^2 - M^2) \\ & + \frac{h_\chi}{2}(X_+ + X_-)Q^2 + \cdots, \end{aligned} \tag{4.1}$$

where M is a constant mass scale and we have included a coupling between the waterfall superfields and additional chiral superfields Q which yield their decay products. The dots indicate additional terms that may be added, involving the inflaton and the superfield Z. In particular, if the scalar component of the latter has a sufficiently large mass, either from superpotential terms, soft masses from SUSY breaking in other sectors or nonminimal terms in the Kähler potential, its expectation value will be set to zero both during and after inflation. The global minimum of the scalar potential will then lie along the real direction $\langle X_+\rangle = \langle X_+\rangle \equiv \chi/\sqrt{2}$, which preserves the discrete symmetry, and the scalar potential relevant for the inflationary and post-inflationary dynamics has the usual hybrid form:

$$V(\phi,\chi) = \frac{\kappa^2}{4}(\chi^2 - M^2)^2 + \frac{g^2}{2}\phi^2\chi^2 + \cdots, \tag{4.2}$$

where $\phi = \sqrt{2}\langle\Phi\rangle$ is the real inflaton scalar field. We recover the usual SUSY hybrid case for $\kappa = g/\sqrt{2}$ and for simplicity we will consider this parametric regime, although our analysis can be extended to the generic case.

Inflation takes place for amplitudes of the inflaton field larger than a critical value, $\phi > \phi_c = M/\sqrt{2}$, such that the waterfall field is held at the origin $\chi = 0$. As the inflaton rolls towards its minimum at $\phi = 0$, its amplitude falls below the critical value and the waterfall field can roll to its true vacuum at $\chi = M$, thus ending inflation. After that point, both fields start to oscillate around its respective minima, triggering the process of reheating the Universe into a radiation era.

In this scenario, the inflaton cannot decay into either the bosonic or fermionic components in the waterfall sector due to kinematical blocking. This is easy to see at the global minimum, where $m_\phi = m_\chi = gM$, but extends to all field values. As in the minimal model, the inflaton can decay into the $Y_\pm$ fields and the decay will be incomplete for $m_f > m_\phi/2 = gM/2$. For simplicity, we assume that the scalar components of the $Y_\pm$ fields acquire large soft masses from SUSY breaking and focus on the fermionic decay channels, noting that the inclusion of both channels will not change our conclusions significantly.

The waterfall fields will decay into the Q sector fields and, for similar reasons, we include only the fermionic decay channels in this case as well. We assume that these fields are light, eventually leading to the complete decay of the waterfall sector and reheating the Universe. Note that neither the inflaton nor the waterfall field can be completely stable in order to reheat the Universe, since they carry a comparable amount of the energy density after inflation. In particular, if the inflaton were completely stable and behaved as dark matter at all times, the decay of the waterfall field would only convert at most half of the total energy density into radiation. The incomplete decay of the inflaton will then reduce its abundance and hence allow for an efficient reheating once the waterfall field decays. For the inflaton to decay incompletely before the waterfall field, we require $h \gtrsim h_\chi$, which is the parametric regime on which we will focus henceforth.

The evolution equations driving the post-inflationary dynamics of the inflaton and waterfall fields, as well as their decay products which we assume to quickly thermalize, are then given by

$$\ddot{\phi} + 3H\dot{\phi} + g^2\chi^2\phi = -\Gamma_\phi\dot{\phi}, \tag{4.3}$$

$$\ddot{\chi} + 3H\dot{\chi} + \frac{g^2}{2}(\chi^2 - M^2 + 2\phi^2)\chi = -\Gamma_\chi\dot{\chi}, \tag{4.4}$$

$$\dot{\rho}_R + 4H\rho_R = \Gamma_\chi\dot{\chi}^2 + \Gamma_\phi\dot{\phi}^2, \tag{4.5}$$

where the decay width of the inflaton is $\Gamma_\phi = \Gamma_+ + \Gamma_-$, with $\Gamma_\pm$ given by Eq. (2.4), while the decay width of the waterfall field is given by

$$\Gamma_\chi = \frac{h_\chi^2 m_\chi}{8\pi}, \tag{4.6}$$

where we neglect the Q fermion masses, which we have checked numerically to be a good approximation for couplings $h_\chi \lesssim 0.1$ in the parameter range of interest to our discussion.

Equations (4.3), (4.4) and (4.5), together with the Friedmann equation,

$$H^2 = \frac{\rho_\phi + \rho_\chi + \rho_R}{3M_P^2} = \frac{\frac{1}{2}\dot{\phi}^2 + \frac{1}{2}\dot{\chi}^2 + V(\phi,\chi) + \rho_R}{3M_P^2}, \tag{4.7}$$

form a complete set of differential equations that can be solved numerically given a set of parameters (g, M, m_f, h, h_χ) and initial conditions, as illustrated in Fig. 3.

At the beginning of the evolution, both the inflaton and the waterfall field mimic a pair of coupled matter fluids which give a roughly equal contribution to the total energy density. After the incomplete decay of the inflaton becomes efficient, its energy density is transferred into the radiation bath, while the waterfall field evolves as an effective noninteracting matter field.

Since radiation is diluted more quickly than matter by the cosmological expansion, the Universe experiences an era of matter domination until the waterfall field effectively decays and reheats the Universe. We then enter the standard radiation-dominated era of big bang cosmology, with the oscillating inflaton field remnant behaving as a cold dark matter component.

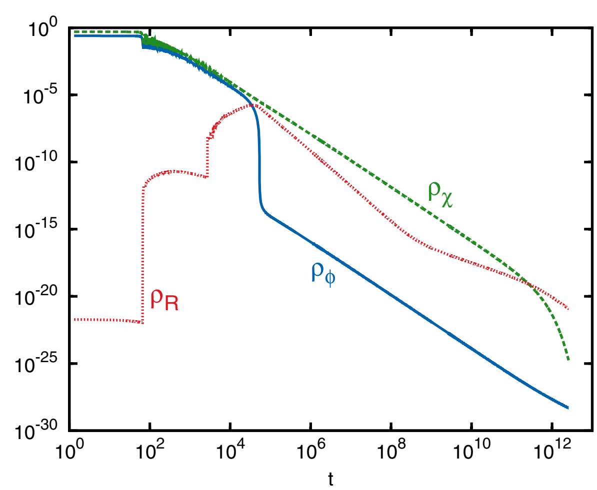

FIG. 3. Results of the numerical integration of the inflaton-waterfall-radiation dynamical equations for $M = 10^{-2}M_P$, $g = 10^{-5}$, $h = 1$, $h_\chi = 10^{-5}$ and $m_f = 0.51m_\phi$, showing the time evolution of the inflaton (solid blue curve), the waterfall (dashed green curve) and radiation (dotted red curve) energy densities. All quantities are given in Planck units such that $M_P = 1$.

Unlike in the minimal model, the reheating temperature is not controlled by the inflaton but rather by the waterfall decay, corresponding to the temperature for which $\Gamma_\chi = H$:

$$T_R \simeq 0.23 g_{*R}^{-1/4} h_\chi (m_\phi M_P)^{1/2}, \tag{4.8}$$

with g_{*R} being the effective number of light degrees of freedom at reheating and where we used $m_\phi = m_\chi$ since the fields are close to the global minimum at this stage.

B. Condensate evaporation: The WIMPlaton scenario

Condensate evaporation will proceed in the hybrid case as in the minimal model with a single inflaton field, at a rate given by Eq. (2.16), leading to a WIMPlaton scenario. In the hybrid case, however, freeze-out of the inflaton particles resulting from evaporation may occur either before or after the decay of the waterfall field. If it occurs after, in the radiation era, one obtains the same value for the inflaton mass yielding the observed dark matter relic abundance as in Eq. (2.20). For small values of the coupling h_χ, in particular

$$h_\chi \lesssim h_\chi^* \equiv 4\times 10^{-9} h \left(\frac{g_{*F}}{10}\right)^{3/8} \left(\frac{x_F}{25}\right)^{-11/8} \left(\frac{\Omega_{\phi 0} h_0^2}{0.1}\right)^{1/4}, \tag{4.9}$$

freeze-out will occur before the decay of the waterfall field, in a matter-dominated era. In this case the inflaton mass yielding the correct relic abundance for dark matter is given by

$$m_\phi = 756 h^2 \left(\frac{h_\chi}{h_\chi^*}\right)^{2/3} \left(\frac{g_{*F}}{10}\right)^{1/4} \left(\frac{x_F}{25}\right)^{-3/4} \times \left(\frac{\Omega_{\phi 0} h_0^2}{0.1}\right)^{1/2} \text{GeV}. \tag{4.10}$$

Note, however, that the coupling h_χ cannot be too small, or otherwise the reheating temperature will be below 100 MeV, spoiling the successful predictions of BBN.

In Fig. 4 we summarize on the plane (h, h_χ) the different possibilities for inflaton dark matter in the hybrid model. In both regions where the freeze-out of the inflaton abundance occurs either in the radiation or waterfall eras, the inflaton can account for the present dark matter abundance for masses in the GeV–TeV range with the reheating and freeze-out temperatures being well above the limit imposed by BBN. This shows that the WIMPlaton scenario introduced earlier is not an exclusive feature of the minimal model but also occurs in other models of inflation with additional dynamical fields.

While the inflaton mass values corresponding to the observed dark matter abundance are not very different in the two realizations that we have analyzed, in the hybrid scenario there are novel phenomenological possibilities. In particular, the inflaton decay products need not interact with the SM degrees of freedom, since these may be excited only after the decay of the waterfall field. Either the waterfall sector decays directly into SM particles or its decay products interact with some of the latter. We note that the waterfall fields may be charged under gauge symmetries, in which case the relevant terms in the superpotential are of the form $\Phi X_\pm \bar{X}_\pm$, etc., where $X_\pm$ and $\bar{X}_\pm$ transform in conjugate representations of the gauge group. This will then open up new avenues for model building in inflaton dark matter scenarios besides those described in the minimal model, which lie, however, outside the scope of this work.

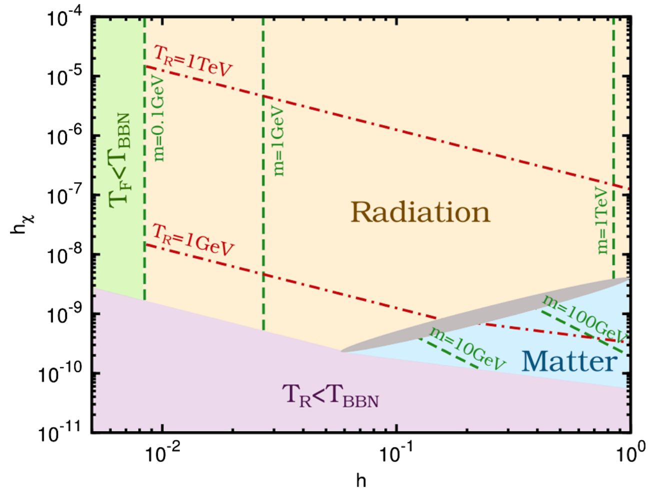

FIG. 4. Parameter space of the hybrid model of inflaton dark matter with $m_f = m_\phi$. In the orange (blue) region the abundance of the inflaton accounts for the present dark matter energy density and the freeze-out occurs in the radiation (waterfall matter) era. The purple (green) region is excluded because the reheating (freeze-out) temperature is below 100 MeV. Dashed (dash-dotted) lines are curves of constant inflaton mass (reheating temperature). The grey area represents approximately the transition between the regions where the freeze-out takes place in the radiation and waterfall era.

V. EMBEDDING IN A CONSISTENT INFLATIONARY MODEL

As we have concluded from the analysis above, the inflaton field can account for dark matter in the Universe at late times for masses below or around the TeV scale in the WIMPlaton scenario. This implies that the inflaton potential cannot be given solely by the terms that we considered in the previous sections. For a quadratic potential, the amplitude of CMB temperature anisotropies would yield

$$m_\phi \simeq \frac{\sqrt{6\pi^2 \Delta_\mathcal{R}^2}}{N_e} M_P \simeq 1.4\times 10^{13} \left(\frac{60}{N_e}\right) \text{GeV}, \tag{5.1}$$

which for 50–60 e-folds of slow-roll inflation largely exceeds the TeV scale. Similarly, in the standard SUSY hybrid models with minimal Kähler potential, inflation is driven essentially by the constant vacuum energy $V_0 = g^2M^4/8$ while the waterfall fields are stabilized at the origin. We may then use the normalization of the scalar curvature power spectrum to obtain the relation

$$m_\phi \gtrsim 2.5 \times 10^{15}|\eta| \text{ GeV}, \tag{5.2}$$

for $\phi \gtrsim M$, where $m_\phi = gM$ is the inflaton mass at the minimum. Since the scalar spectral index $n_s \simeq 1 + 2\eta \simeq 0.9603 \pm 0.0073$ at 68 C.L. [4] in these scenarios, we conclude that $m_\phi \lesssim$ TeV cannot yield an observationally consistent model. Typically we have from the normalization of the spectrum $M \sim 10^{13}$–10^{16} GeV, and then the WIMPlaton scenario requires small couplings $g \sim 10^{-13}$–10^{-10}, which are responsible for the very flat potential during inflation and the scale invariant spectrum. A nonminimal Kähler potential can yield the observed spectral index for lower values of the coupling, although at the expense of a slight increase in M [36,37], which again makes it difficult to achieve the required WIMPlaton mass values.

However, our analysis assumed only that m_ϕ is the inflaton mass as it oscillates about the minimum of the potential at the origin, while the effective inflaton mass can be much larger if slow roll occurs for significantly larger field values where self-interactions become important. For example, the discrete $\mathbb{Z}_2$ symmetry allows for quartic self-interactions such that

$$V(\phi) = \frac{\lambda}{4!}\phi^4 + \frac{1}{2}m_\phi\phi^2 + \cdots. \tag{5.3}$$

In the minimal model the quartic term dominates for $|\phi| > \sqrt{12/\lambda}m_\phi$, while our analysis is valid if the quartic term is subdominant for the field values $|\phi| \sim m_f/h$ at which decay into fermions occurs, which requires

$$h \gtrsim \sqrt{\frac{\lambda}{12}}\left(\frac{m_f}{m_\phi}\right) \simeq 3 \times 10^{-8}\left(\frac{m_f}{m_\phi}\right), \tag{5.4}$$

where we have used $\lambda \simeq 10^{-14}$ as imposed by the COBE normalization for inflation with a quartic potential. This is easily satisfied if the fermions are not much heavier than the inflaton given the more stringent bounds on the Yukawa coupling discussed earlier.

We note that, in supergravity models, such chaotic inflation scenarios can be obtained by considering a non-minimal Kähler potential for the inflaton, while taking the canonical one for the other superfields in the model, in particular the Z field. One possibility is to consider a Kähler potential with a shift symmetry [38], e.g. $K(\Phi, Z, \ldots) = (\Phi + \Phi^\dagger)^2/2 + ZZ^\dagger + \cdots$, with inflation taking place along the imaginary component of the scalar inflaton.

A quartic term is, however, not sufficient to produce a consistent model of inflation, since it predicts a too red-tilted spectrum for curvature perturbations and a tensor-to-scalar ratio already outside the bounds obtained by Planck [39] and BICEP2 [40,41]. A consistent spectrum may, for example, be achieved in warm inflation scenarios [42]. However, as warm inflation naturally leads to radiation becoming the dominant component at the end of the slow-roll regime, the post-inflationary evolution will necessarily differ from the dynamical picture discussed here.

Another interesting possibility is the inclusion of a nonminimal coupling to the gravitational sector, in particular a coupling between the inflaton and Ricci scalar of the form $\xi\phi^2R$, which is compatible with the discrete $\mathbb{Z}_2$ symmetry. The resulting inflationary scenario yields a perturbation spectrum that smoothly interpolates between the minimal quartic model and the Starobinsky model as the nonminimal coupling constant increases. On the one hand, the latter is characterized by a low tensor-to-scalar ratio and a spectral index $n_s = 0.96$–0.97 in agreement with the Planck results; on the other hand, a small non-minimal coupling constant is preferred to obtain a non-negligible tensor-to-scalar ratio. We refer the reader to [43,44] and the references therein for a more detailed discussion of these scenarios, since here we are mainly interested in the post-inflationary dynamics. The effect of the nonminimal coupling on the effective scalar potential in the Einstein frame becomes negligible for $\xi\phi^2/M_P^2 \lesssim 1$, such that consistency of our analysis implies

$$h \gtrsim \sqrt{\xi}\frac{m_f}{M_P}, \tag{5.5}$$

which is generically less stringent than Eq. (5.4) for masses in the TeV range and $\xi < 10^{16}$.

A quartic self-coupling can also be easily introduced in SUSY hybrid inflation by a superpotential coupling between the inflaton and the auxiliary Z field, $\sqrt{\lambda}\Phi^2Z/2$. We then have

$$\frac{\lambda\phi^4}{V_0} \sim 6\times 10^{16}\left(\frac{\lambda}{10^{-14}}\right)\left(\frac{1\,\text{TeV}}{m_\phi}\right)^2\left(\frac{M_P}{M}\right)^2\left(\frac{\phi}{M_P}\right)^4, \tag{5.6}$$

so that the quartic term will easily dominate over the vacuum term for super-Planckian values and the typical parameters required by the normalization of the spectrum and the WIMPlaton scenario. We note that the vacuum term will come to dominate the energy at small field values, but for $\phi \gtrsim \phi_c$ we have

$$\eta \sim \lambda\frac{\phi^2}{M^2}\frac{M_P^2}{m_\phi^2} \gtrsim \lambda\frac{M_P^2}{m_\phi^2} \gg 1, \tag{5.7}$$

such that slow-roll inflation never takes place in the small field regime, and inflation may then occur entirely in a chaotic regime [45–47]. An alternative interesting possibility that can realize inflation at low energy scales is to inflate along the waterfall trajectory, which has been shown to yield a sufficiently red-tilted spectrum for small couplings g within the context of both SUSY and non-SUSY hybrid models [48,49]. Although it has been recently realized that isocurvature contributions yield a too large enhancement of the amplitude of the spectrum unless $N_e \lesssim 60$ [50], the WIMPlaton scenario typically favors a long reheating period, with e.g. $N_e \sim 30$–40 being allowed by the data [50].

Finally, the discrete $\mathbb{Z}_2 \times S_2$ symmetry does not protect the minimal model from radiative corrections. In particular, the Yukawa interactions induce loop corrections of the Coleman-Weinberg form, which for large inflaton field values take the leading form:

$$\Delta V_f \approx -\frac{h^4\phi^4}{16\pi^2}\left(\log\left(\frac{h^2\phi^2}{\mu^2}\right) - \frac{3}{2}\right) \tag{5.8}$$

and therefore induce an effective quartic term in the potential. The effect of this term does not necessarily spoil the predictions of the nonminimally coupled quartic model, as discussed in [51], although one must ensure that the observed normalization of the perturbation spectrum is obtained. While for $\xi \ll 1$ the effective quartic coupling must have approximately the same value as in the minimally coupled case, which requires $h \lesssim 10^{-3}$, significantly larger values can be accommodated for large nonminimal couplings.

Radiative corrections can, however, be significantly reduced in supersymmetric scenarios, as we considered for the hybrid case. A supersymmetric version of the minimal model consistent with the discrete symmetry corresponds to a superpotential of the form

$$W = \frac{h}{2}\Phi(Y_+^2 - Y_-^2) + \frac{m_f}{2}(Y_+^2 + Y_-^2) + \frac{m_\phi}{2}\Phi^2 + \frac{\sqrt{\lambda}}{2}\Phi^2 Z, \tag{5.9}$$

where the auxiliary superfield Z induces the quartic term in the inflaton potential (noting that the discrete symmetry forbids cubic inflaton terms in the superpotential). Supersymmetry then cancels the leading contributions of scalars and bosons to the one-loop Coleman-Weinberg potential, which becomes

$$\Delta V_{\rm SUSY} \approx \frac{h^2}{16\pi^2}\log\left(\frac{h^2\phi^2}{\mu^2}\right)V(\phi). \tag{5.10}$$

This contribution is thus necessarily smaller than the tree-level potential $V(\phi) \simeq \lambda|\phi|^4/4$ for $h \lesssim 1$, therefore avoiding the generation of large effective self-interactions during inflation.

Besides the Yukawa terms considered so far, the supersymmetric model also yields scalar interactions between the inflaton and the scalar partners $y_\pm$ of the fermions $\psi_\pm$ which, apart from SUSY splittings that vanish at the origin, have the same mass $m_\pm = |m_f + h\phi|$. Trilinear terms in the scalar potential also lead to the decay $\phi \to y_\pm y_\pm$, with analogous kinematics and comparable widths to the fermionic decay channels, therefore yielding a similar incomplete decay of the inflaton as analyzed above. We note that the incomplete decay dynamics can be fully described in terms of scalar fields and is therefore not exclusive of fermion Yukawa couplings, although the required form of the scalar mass terms is more naturally motivated within a supersymmetric context. We also note the existence of quartic terms in the scalar potential which induce the three-body decay $\phi \to z y_\pm y_\pm$, where z is the scalar component of the Z chiral multiplet. The associated couplings have opposite signs for $y_\pm$ and are proportional to λ, so that they are typically subdominant with respect to the two-body decays and they are also kinematically forbidden at late times, so that they do not affect our earlier conclusions.

VI. CONCLUSION

In this work we have shown that the decay of the inflaton following the inflationary slow-roll regime can be incomplete, such that successful reheating is achieved while leaving a stable remnant that can account for the observed dark matter in the Universe. This is achieved by coupling the inflaton field ϕ to a pair of fermions $\psi_\pm$ (and/or similarly scalar fields) while imposing a discrete symmetry that simultaneously changes $\phi \to -\phi$ and interchanges the two fermions, corresponding to the discrete subgroup $C_2 \subset \mathbb{Z}_2 \times S_2$. This symmetry forbids all inflaton decay channels except for $\phi \to \psi_\pm\psi_\pm$ if it is preserved in the vacuum state, including decays mediated by off-shell fermion modes. The inflaton is thus stable at late times if the bare fermion mass $m_f > m_\phi/2$, where m_ϕ denotes the inflaton mass at the origin. However, since the physical fermion mass is field dependent, inflaton decay may occur during the initial oscillations while the field amplitude is sufficiently large.

The incomplete decay of the inflaton produces a bath of fermions and possibly other particles, including the Standard Model fields. These fermions will in turn scatter off individual bosons in the scalar field condensate, likely leading to its evaporation. The combination of decays and evaporation processes will then reheat the Universe after inflation and we have found that the BBN constraint on the reheating temperature points towards a complete evaporation of the scalar field condensate.

In the radiation era, the stable inflaton particles can reach a thermal equilibrium state. This is maintained through annihilation into fermions until it becomes inefficient compared to Hubble expansion and the inflaton particles fall out of equilibrium. Their frozen relic abundance can then account for the observed dark matter in the Universe for inflaton masses in the GeV–TeV range. Our mechanism thus conciliates reheating after inflation with the standard WIMP mechanism, so that we denote this generically as the WIMPlaton scenario.

In the simplest models with a single dynamical field, the inflaton decay products must interact with the Standard Model (SM) degrees of freedom in order to excite them in the thermal bath. We have explored different possibilities for such interactions, including the case where the inflaton decays into right-handed neutrinos or mili-charged particles in a hidden sector. We have also explored the alternative possibility of hybrid inflation models, where a dynamical waterfall sector, which is also charged under the discrete symmetry, is responsible for reheating the Universe. While the incomplete decay of the inflaton is still required to ensure the transition into a radiation-dominated epoch, in such scenarios the inflaton decay products need not couple directly to the SM fields, which opens up new phenomenological possibilities.

While inflaton mass values in the GeV–TeV range may *a priori* seem too low to yield the correct amplitude for the primordial spectrum of curvature perturbations, we have shown that the inflaton mass can be much larger during the slow-roll period than at the minimum of the potential. This is for example the case of large-field chaotic models, where the inflaton mass is dominantly set by a quartic self-interaction at super-Planckian values during inflation and by a bare quadratic term close to the minimum. Such models are consistent with the latest CMB anisotropy measurements by Planck if e.g. nonminimal couplings to gravity or dissipative effects are included. In the case of hybrid models inflation can also occur in a chaotic regime, but there are also other viable scenarios for inflation with low inflaton mass, such as along the waterfall direction. A possible alternative that may generically allow for small inflaton masses is the curvaton scenario [52,53], where the fluctuations of an additional field that is light during inflation dominate the total curvature perturbations when decaying.

Our results show that there exist concrete particle physics models where a single dynamical field can drive inflation in the early Universe and account for cold dark matter at late times. This is an appealing feature from the theoretical perspective, since it allows one to address two of the most important problems in modern cosmology within the same simple model. Moreover, we have found that consistent cosmological scenarios typically require the masses of both the inflaton and the particles it interacts directly with to lie in the GeV–TeV range, opening up the possibility of testing both inflation and dark matter at present particle colliders such as the LHC, as well as direct or indirect dark matter searches. Finally, this scenario also singles out particular classes of inflationary models that can be tested with CMB experiments such as Planck, which are now reaching unprecedented levels of precision. We thus hope that our work motivates further exploration of the phenomenological and observational consequences of the proposed generic mechanism in its several possible implementations.

ACKNOWLEDGMENTS

M. B. G. and R. C. are supported by the "Junta de Andalucía" group FQM101 and Project No. P10-FQM-6552. J. G. R. is supported by the FCT Grant No. SFRH/BPD/85969/2012 and partially by Grant No. PTDC/FIS/116625/2010, the CIDMA strategic Project No. UID/MAT/04106/2013 and the Marie Curie action NRHEP-295189-FP7-PEOPLE-2011-IRSES. The authors would like to thank the hospitality of the Higgs Centre for Theoretical Physics during the completion of this work. R. C. would like to thank the hospitality of the University of Aveiro and the "Junta de Andalucía" for supporting this visit.

[1] A. H. Guth, The inflationary universe: A possible solution to the horizon and flatness problems, Phys. Rev. D **23**, 347 (1981); A. Albrecht and P. J. Steinhardt, Cosmology for Grand Unified Theories with Radiatively Induced Symmetry Breaking, Phys. Rev. Lett. **48**, 1220 (1982); A. D. Linde, A new inflationary universe scenario: A possible solution of the horizon, flatness, homogeneity, isotropy and primordial monopole problems, Phys. Lett. B **108**, 389 (1982).

[2] A. Borriello and P. Salucci, The dark matter distribution in disk galaxies, Mon. Not. R. Astron. Soc. **323**, 285 (2001).

[3] E. Komatsu *et al.* (WMAP Collaboration), Five-year Wilkinson Microwave Anisotropy Probe (WMAP 1) observations: Cosmological interpretation, Astrophys. J. Suppl. Ser. **180**, 330 (2009).

[4] P. A. R. Ade *et al.* (Planck Collaboration), Planck 2013 results. I. Overview of products and scientific results, Astron. Astrophys. **571**, A1 (2014).

[5] M. Tegmark *et al.* (SDSS Collaboration), The 3D power spectrum of galaxies from the SDSS, Astrophys. J. **606**, 702 (2004).

[6] H. Hoekstra, H. Yee, and M. Gladders, Current status of weak gravitational lensing, New Astron. Rev. **46**, 767 (2002).

[7] M. Taoso, G. Bertone, and A. Masiero, Dark matter candidates: A ten-point test, J. Cosmol. Astropart. Phys. 03 (2008) 022.

[8] M. S. Turner, Coherent scalar field oscillations in an expanding universe, Phys. Rev. D **28**, 1243 (1983).
[9] A. R. Liddle and L. A. Urena-Lopez, Inflation, Dark Matter and Dark Energy in the String Landscape, Phys. Rev. Lett. **97**, 161301 (2006).
[10] A. Albrecht, P. J. Steinhardt, M. S. Turner, and F. Wilczek, Reheating an Inflationary Universe, Phys. Rev. Lett. **48**, 1437 (1982); A. D. Dolgov and A. D. Linde, Baryon asymmetry in the inflationary universe, Phys. Lett. B **116**, 329 (1982); L. F. Abbott, E. Fahri, and M. Wise, Particle production in the new inflationary cosmology, Phys. Lett. B **117**, 29 (1982); L. Kofman, A. D. Linde, and A. A. Starobinsky, Reheating After Inflation, Phys. Rev. Lett. **73**, 3195 (1994).
[11] L. Kofman, A. D. Linde, and A. A. Starobinsky, Towards the theory of reheating after inflation, Phys. Rev. D **56**, 3258 (1997).
[12] A. Berera and L.-Z. Fang, Thermally Induced Density Perturbations in the Inflation Era, Phys. Rev. Lett. **74**, 1912 (1995); A. Berera, Warm Inflation, Phys. Rev. Lett. **75**, 3218 (1995).
[13] A. Berera, I. G. Moss, and R. O. Ramos, Warm inflation and its microphysical basis, Rep. Prog. Phys. **72**, 026901 (2009).
[14] V. H. Cardenas, Inflation, reheating and dark matter, Phys. Rev. D **75**, 083512 (2007); G. Panotopoulos, A brief note on how to unify dark matter, dark energy, and inflation, Phys. Rev. D **75**, 127301 (2007); K. Mukaida and K. Nakayama, Dark matter chaotic inflation in light of BICEP2, J. Cosmol. Astropart. Phys. 08 (2014) 062.
[15] A. R. Liddle, C. Pahud, and L. A. Urena-Lopez, Triple unification of inflation, dark matter, and dark energy using a single field, Phys. Rev. D **77**, 121301 (2008).
[16] N. Okada and Q. Shafi, WIMP dark matter inflation with observable gravity waves, Phys. Rev. D **84**, 043533 (2011); A. de la Macorra, Dark matter from the inflaton field, Astropart. Phys. **35**, 478 (2012).
[17] R. N. Lerner and J. McDonald, Gauge singlet scalar as inflaton and thermal relic dark matter, Phys. Rev. D **80**, 123507 (2009); V. V. Khoze, Inflation and dark matter in the Higgs portal of classically scale invariant Standard Model, J. High Energy Phys. 11 (2013) 215.
[18] N. Bose and A. S. Majumdar, Unified model of k-inflation, dark matter and dark energy, Phys. Rev. D **80**, 103508 (2009); J. De-Santiago and J. L. Cervantes-Cota, Generalizing a unified model of dark matter, dark energy, and inflation with noncanonical kinetic term, Phys. Rev. D **83**, 063502 (2011).
[19] G. Steigman and M. S. Turner, Cosmological constraints on the properties of weakly interacting massive particles, Nucl. Phys. **B253**, 375 (1985).
[20] W.-K. Tung, *Group Theory in Physics* (World Scientific Publishing Co. Pte. Ltd., Singapore, 1985).
[21] E. W. Kolb and M. S. Turner, *The Early Universe (Frontiers in Physics)* (Addison-Wesley Publising Co., Redwood City, 1988), Vol. 70, p. 719.
[22] K. A. Olive, G. Steigman, and T. P. Walker, Primordial nucleosynthesis: Theory and observations, Phys. Rep. **333–334**, 389 (2000).
[23] P. B. Greene and L. Kofman, Preheating of fermions, Phys. Lett. B **448**, 6 (1999).
[24] G. F. Giudice, M. Peloso, A. Riotto, and I. Tkachev, Production of massive fermions at preheating and leptogenesis, J. High Energy Phys. 08 (1999) 014.
[25] D. J. H. Chung, E. W. Kolb, A. Riotto, and I. I. Tkachev, Probing Planckian physics: Resonant production of particles during inflation and features in the primordial power spectrum, Phys. Rev. D **62**, 043508 (2000).
[26] I. Rudenok, Y. Shtanov, and S. Vilchinskii, Post-inflationary preheating with weak coupling, Phys. Lett. B **733**, 193 (2014).
[27] G. Aad *et al.* (ATLAS Collaboration), Observation of a new particle in the search for the Standard Model Higgs boson with the ATLAS detector at the LHC, Phys. Lett. B **716**, 1 (2012); S. Chatrchyan *et al.* (CMS Collaboration), Observation of a new boson at a mass of 125 GeV with the CMS experiment at the LHC, Phys. Lett. B **716**, 30 (2012).
[28] H. Vogel and J. Redondo, Dark radiation constraints on minicharged particles in models with a hidden photon, J. Cosmol. Astropart. Phys. 02 (2014) 029.
[29] R. Foot and S. Vagnozzi, Dissipative hidden sector dark matter, Phys. Rev. D **91**, 023512 (2015); Diurnal modulation signal from dissipative hidden sector dark matter, Phys. Lett. B **748**, 61 (2015).
[30] G. R. Dvali, Q. Shafi, and R. K. Schaefer, Large Scale Structure and Supersymmetric Inflation without Fine Tuning, Phys. Rev. Lett. **73**, 1886 (1994).
[31] A. D. Linde, Eternal extended inflation and graceful exit from old inflation without Jordan-Brans-Dicke, Phys. Lett. B **249**, 18 (1990); Axions in inflationary cosmology, Phys. Lett. B **259**, 38 (1991); Hybrid inflation, Phys. Rev. D **49**, 748 (1994).
[32] E. J. Copeland, A. R. Liddle, D. H. Lyth, E. D. Stewart, and D. Wands, False vacuum inflation with Einstein gravity, Phys. Rev. D **49**, 6410 (1994).
[33] P. Binetruy and G. Dvali, D-Term Inflation, Phys. Lett. B **388**, 241 (1996); E. Halyo, Hybrid Inflation from Supergravity D Terms, Phys. Lett. B **387**, 43 (1996).
[34] M. Bastero-Gil and S. F. King, A next-to-minimal supersymmetric model of hybrid inflation, Phys. Lett. B **423**, 27 (1998); F term hybrid inflation in effective supergravity theories, Nucl. Phys. **B549**, 391 (1999).
[35] S. Antusch, K. Dutta, and P. M. Kostka, SUGRA hybrid inflation with shift symmetry, Phys. Lett. B **677**, 221 (2009); Tribrid inflation in supergravity, AIP Conf. Proc. **1200**, 1007 (2010).
[36] M. Bastero-Gil, S. F. King, and Q. Shafi, Supersymmetric hybrid inflation with nonminimal Kahler potential, Phys. Lett. B **651**, 345 (2007).
[37] M. ur Rehman, V. N. Senoguz, and Q. Shafi, Supersymmetric and smooth hybrid inflation in the light of WMAP3, Phys. Rev. D **75**, 043522 (2007).
[38] M. Kawasaki, M. Yamaguchi, and T. Yanagida, Natural Chaotic Inflation in Supergravity, Phys. Rev. Lett. **85**, 3572 (2000); P. Brax and J. Martin, Shift symmetry and inflation in supergravity, Phys. Rev. D **72**, 023518 (2005); R. Kallosh and A. Linde, New models of chaotic inflation in supergravity, J. Cosmol. Astropart. Phys. 11 (2010) 011.
[39] P. A. R. Ade *et al.* (Planck Collaboration), Planck 2015 results. XIII. Cosmological parameters, arXiv:1502.01589.
[40] P. A. R. Ade *et al.* (BICEP2 Collaboration), Detection of *B*-Mode Polarization at Degree Angular Scales by BICEP2, Phys. Rev. Lett. **112**, 241101 (2014).

[41] P. A. R. Ade *et al.* (BICEP2 and Planck Collaborations), Joint Analysis of BICEP2/Keck Array and Planck Data, Phys. Rev. Lett. **114**, 101301 (2015).

[42] S. Bartrum, M. Bastero-Gil, A. Berera, R. Cerezo, R. O. Ramos, and J. G. Rosa, The importance of being warm (during inflation), Phys. Lett. B **732**, 116 (2014).

[43] D. S. Salopek, J. R. Bond, and J. M. Bardeen, Designing density fluctuation spectra in inflation, Phys. Rev. D **40**, 1753 (1989); B. L. Spokoiny, Inflation and generation of perturbations in broken symmetric theory of gravity, Phys. Lett. B **147**, 39 (1984); T. Futamase and K. Maeda, Chaotic inflationary scenario in models having nonminimal coupling with curvature, Phys. Rev. D **39**, 399 (1989); R. Fakir and W. G. Unruh, Improvement on cosmological chaotic inflation through nonminimal coupling, Phys. Rev. D **41**, 1783 (1990); D. I. Kaiser, Primordial spectral indices from generalized Einstein theories, Phys. Rev. D **52**, 4295 (1995); E. Komatsu and T. Futamase, Complete constraints on a nonminimally coupled chaotic inflationary scenario from the cosmic microwave background, Phys. Rev. D **59**, 064029 (1999); K. Nozari and S. D. Sadatian, Nonminimal inflation after WMAP3, Mod. Phys. Lett. A **23**, 2933 (2008).

[44] F. L. Bezrukov, Nonminimal coupling in inflation and inflating with the Higgs boson, arXiv:0810.3165; F. L. Bezrukov and M. Shaposhnikov, The Standard Model Higgs boson as the inflaton, Phys. Lett. B **659**, 703 (2008); F. L. Bezrukov, A. Magnin, and M. Shaposhnikov, Standard Model Higgs boson mass from inflation, Phys. Lett. B **675**, 88 (2009); F. Bezrukov, D. Gorbunov, and M. Shaposhnikov, On initial conditions for the hot big bang, J. Cosmol. Astropart. Phys. 06 (2009) 029; F. Bezrukov and M. Shaposhnikov, Standard Model Higgs boson mass from inflation: Two loop analysis, J. High Energy Phys. 07 (2009) 089.

[45] S. Clesse and J. Rocher, Avoiding the blue spectrum and the fine-tuning of initial conditions in hybrid inflation, Phys. Rev. D **79**, 103507 (2009).

[46] S. Clesse and J. Rekier, Updated constraints on large field hybrid inflation, Phys. Rev. D **90**, 083527 (2014).

[47] L. M. Carpenter and S. Raby, Chaotic hybrid inflation with a gauged BL, Phys. Lett. B **738**, 109 (2014).

[48] S. Clesse, Hybrid inflation along waterfall trajectories, Phys. Rev. D **83**, 063518 (2011); S. Clesse and B. Garbrecht, Slow roll during the waterfall regime: The small coupling window for SUSY hybrid inflation, Phys. Rev. D **86**, 023525 (2012).

[49] H. Kodama, K. Kohri, and K. Nakayama, On the waterfall behavior in hybrid inflation, Prog. Theor. Phys. **126**, 331 (2011).

[50] S. Clesse, B. Garbrecht, and Y. Zhu, Non-Gaussianities and curvature perturbations from hybrid inflation, Phys. Rev. D **89**, 063519 (2014).

[51] N. Okada, M. U. Rehman, and Q. Shafi, Tensor to scalar ratio in nonminimal ϕ^4 inflation, Phys. Rev. D **82**, 043502 (2010).

[52] D. H. Lyth and D. Wands, Generating the curvature perturbation without an inflaton, Phys. Lett. B **524**, 5 (2002); T. Moroi and T. Takahashi, Effects of Cosmological Moduli Fields on Cosmic Microwave Background, Phys. Lett. B **522**, 215 (2001); **539**, 303(E) (2002); K. Enqvist and M. S. Sloth, Adiabatic CMB perturbations in pre-big bang string cosmology, Nucl. Phys. **B626**, 395 (2002).

[53] For a review on curvaton and particle physics models see, for example, A. Mazumdar and J. Rocher, Particle physics models of inflation and curvaton scenarios, Phys. Rep. **497**, 85 (2011).