



Constraining chameleon field driven warm inflation with Planck 2018 data

Haidar Sheikahmadi^{1,2,a}, Abolhassan Mohammadi^{3,4,b}, Ali Aghamohammadi^{5,c}, Tiberiu Harko^{6,7,d}, Ramón Herrera^{8,e}, Christian Corda^{9,10,f}, Amare Abebe^{2,g}, Khaled Saaidi^{3,h}

¹ School of Astronomy, Institute for Research in Fundamental Sciences (IPM), P. O. Box 19395-5531, Tehran, Iran

² Center for Space Research, North-West University, Mafikeng, South Africa

³ Department of Physics, Faculty of Science, University of Kurdistan, Sanandaj, Iran

⁴ Dipartimento di Fisica e Astronomia, Università di Bologna, via Iriero 46, 40126 Bologna, Italy

⁵ Sanandaj Branch, Islamic Azad University, Sanandaj, Iran

⁶ Department of Physics, Babes-Bolyai University, Kogalniceanu Street, 400084 Cluj-Napoca, Romania

⁷ School of Physics, Sun Yat-Sen University, Xingang Road, Guangzhou 510275, People's Republic of China

⁸ Instituto de Física, Pontificia Universidad de Católica de Valparaíso, Casilla 4950, Valparaíso, Chile

⁹ Research Institute for Astronomy and Astrophysics of Maragha (RIAAM), P.O. Box 55134-441 Maragha, Iran

¹⁰ International Institute for Applicable Mathematics and Information Sciences (IIAMIS), B.M. Birla Science Centre, Adarsh Nagar, Hyderabad 500 463, India

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Abstract We investigate warm inflationary scenario in which the accelerated expansion of the early Universe is driven by chameleon-like scalar fields. Due to the non-minimal coupling between the scalar field and the matter sector, the energy-momentum tensor of each fluid component is not conserved anymore, and the generalized balance equation is obtained. The new source term in the energy equation can be used to model warm inflation. On the other hand, if the coupling function varies slowly, the model reduces to the standard model used for the description of cold inflation. To test the validity of the warm chameleon inflation model, the results for warm inflationary scenarios are compared with the observational Planck2018 Cosmic Microwave Background data. In this regard, the perturbation parameters such as the amplitude of scalar perturbations, the scalar spectral index and the tensor-to-scalar ratio are derived at the horizon crossing in two approximations, corresponding to the weak and strong dissipative regimes. As a general result it turns out that the theoretical predictions of the chameleon

warm inflationary scenario are consistent with the Planck 2018 observations.

1 Introduction

Four decades after the introduction of the inflation model, it can now be considered as one of the cornerstones of modern cosmology [1–9]. The main success of this theory goes back to resolving the three main drawbacks of the Standard Big Bang (SBB) theory, namely the flatness, horizon and relic problems [10–15]. In inflation theory we deal with very high energy scales, between 200 GeV and 10^{12} TeV [16], and, by considering the quantum effects inherent to these energies, one can explain the origin of seeds for Large-Scale Structure (LSS) formation, besides the fluctuations generating the Cosmic Microwave Background (CMB) anisotropies [16–25]. In the present work, we will investigate in more detail the relation between quantum fluctuations and classical behaviour at the end of inflation.

Inflationary scenarios usually predict that the power spectrum originated from primary quantum fluctuations should have equal power on all scales, i.e., it should be scale invariant [8, 9, 26–30]. Although several interesting models have been proposed to explain the origin of the Universe [31–60], inflation, due to its potential for solving the major problems of the SBB theory, has become a crucial part of present day cosmology. To generate an inflationary phase for the very

^a e-mails: h.sh.ahmadi@gmail.com; h.sheikahmadi@ipm.ir

^b e-mails: abolhassanm@gmail.com; a.mohammadi@uok.ac.ir

^c e-mail: a.aqamohamadi@gmail.com

^d e-mail: tiberiu.harko@gmail.com

^e e-mail: ramon.herrera@pucv.cl

^f e-mail: cordac.galilei@gmail.com

^g e-mail: amare.abbebe@gmail.com

^h e-mails: khaledsaaidi@gmail.com; ksaaidi@uok.ac.ir

early evolution of the Universe, which could explain the symmetries of the cosmological principle, one of the important physical candidates is a scalar field. This type of exotic matter can be described as a fluid with negative pressure that inflation needs to have. In other words, the logarithm of the scalar field potential should move much slower than the kinetic part, leading to the rapid exponential expansion determining the inflationary expansion [1–6].

Despite all of its successes the nature and the origin of the inflationary theory remain a mystery, and it is still a matter of intense scientific investigation [10, 11, 16]. Nevertheless, this is not so surprising for us because inflation happened, as mentioned before, in a very high energy epoch. Up to now, there is no complete or convincing version of the Grand Unification Theory (GUT) that could describe the physics of the very early Universe, and of its beginning. Following [16], and from the observational point of view, we can classify the huge number of different models of inflation into three main categories, with their related sub-categories. These are the single-field inflation [61–74], the multiple-field inflation [75–82] and those models in which the fluid is not described by scalar fields [83–87].

One of the important requirements of the inflationary scenarios is to establish a physically realistic relation between the end of the early accelerated, de Sitter type evolution, and the beginning of the radiation-dominated epoch, which corresponds to the SBB model. Based on this requirement, one can consider two different types of inflation, namely the super-cold and the warm models of inflation [88–100].

Whereas each one of these scenarios has its advantages and drawbacks, in the present work we want to consider both of them in more detail for a single-field model of inflation, based on the chameleon mechanism. But before going through this model, let us explain the properties of super-cold and warm proposals of inflation and their successes and failures. In super-cold models, to establish a relation between the end of inflation and the beginning of the hot big bang model, usually the concept of quantum fluctuations, based on the scalar field oscillations, is considered. In other words, in this approach, one should assume a tachyonic pre-heating phase. Hence, based on this mechanism, the end of inflation is smoothly connected to the radiation-dominated epoch [75, 101–109]. The main problem of this model goes back to the consistency with observations. Unfortunately we know very little, from an observational point of view, about these eras and their evolutions, i.e., the pre-heating and the reheating, and we have no accurate criteria to compare the theoretical results with observations [16]. Recently, different proposals have been proposed to bind existing observational data together to obtain a better estimation of super-cold inflation [21, 22, 110–139]. Another important constraint originating from this setup, besides the adiabatic initial condition on the CMB, is that the temperature due to reheating has to be larger

than the BB Nucleosynthesis (BBN) scales. For more details and for different solutions to overcome this drawback, we refer the reader to Ref. [140], and references therein.

In the present paper we are going to investigate the properties of warm inflation in more detail. A necessary condition for a field theory to produce warm inflation is that the radiation must be in thermal equilibrium. A crucial advantage of this model of inflation is that it does not need any preheating and reheating mechanisms to realize the connection between the end of inflation and the beginning of the radiation era [88–92]. To see how this can be achieved, let us go back to the aforementioned crucial problem of initial quantum fluctuations, related to the super-cold inflation. As opposed to the super-cold model of inflation, in warm inflation one can assume the existence of an interaction between the different components of the initial stages of the Universe as an intrinsic physical property. This assumption plays the role of a master-key for warm inflation. In other words, when the scalar field acquires the zero-point energy of the inflaton, the responsible field in driving inflation, its reaction back on the inflaton can cause the damping of its motion [92, 93, 141]. We must stress that to solve the flatness and horizon problems this reaction should be strong [92, 93]. Therefore instead of slow-roll mechanisms in warm inflation the concept of over-damped motion can be introduced [92, 93].

The prediction of the cosmological quantum perturbations is one of the main achievements of the inflationary scenario. The perturbations are divided into three types: scalar, vector, and tensor perturbations, which evolve independently up to the linear order. The scalar perturbations are known to be the seeds of the LSS of the Universe, while the tensor perturbations are primordial gravitational waves. For the sake of completeness, we recall that the now-famous event GW150914, which is the first direct observation of gravitational waves from a binary black hole merger [142] occurred in the 100th anniversary of Albert Einstein's prediction of gravitational waves [143]. That event was a cornerstone for science in general and for gravitational physics in particular. It indeed gave definitive proof of the existence of gravitational waves, of the existence of black holes having mass greater than 25 solar masses, and of the existence of binary systems of black holes that merge in a time less than the age of the Universe [142].

Such a direct gravitational waves detection represented the starting point of the new era of gravitational waves astronomy. After the event GW150914, the LIGO Scientific Collaboration announced other six new gravitational waves detections, the events GW151226 [144], GW170104 [145], GW170814 [146], GW170817 [147], GW170608 [148] and, recently, GW151012 [149]. All the cited events again are the consequences of binary black hole coalescences, with the sole exception of the event GW170817, which represents the first gravitational waves detection from a neutron star merger

[147]. Gravitational waves detectors could be, in principle, decisive to confirm the physical consistency of Einstein's general theory of relativity, or, alternatively, to endorse the framework of extended theories of gravity [150, 151]. In fact, some differences between the GTR and alternative theories appear in the linearized theory of gravity, and they could be observed through different interferometer response functions [150–152].

As mentioned already, important candidates for the initial seeds of LSS formation are the quantum fluctuations. But the main question is how we can find a mechanism to link quantum fluctuations to these classical seeds, namely the quantum-to-classical transition problem [153–156]. In [156] it was shown that although by considering modes larger than the Hubble length one can treat the fluctuations classically, they are not covariant under canonical transformation [153, 154]. To solve this problem, a solution was proposed in [157–161], which is based on the neglect of the quantum interaction between macroscopically distinguished events, namely decoherence. Whereas in warm inflation the dissipation behaviour appears as a basic property, the decoherence can automatically generate it. In [92, 162–166] different conditions for warm inflation originating from quantum field theory were studied.

The present paper is organized as follows. We first briefly review the chameleon scalar field model in Sect. 2. In particular, we introduce the gravitational action, involving a non-minimal coupling between scalar field and the matter sector, and we write down the generalized Friedmann equations of the theory, as well as the energy balance equation, in which the coupling between scalar field and matter generates a new source term, given by the product of the derivative of the coupling function with respect to the scalar field, and the matter Lagrangian. The chameleon warm inflation model is introduced, and analyzed, in Sect. 3. First, in Sect. 3.1 we review the standard approach to warm inflation, in which a radiation component is present, together with the scalar, during inflation. Due to the interaction between these two components, there is an energy transfer from the scalar field to radiation, due to the presence of a new source term in the conservation equations. The slow-roll parameters and the number of e-folds are also obtained. The main theoretical model we investigate in the paper, the chameleon field driven warm inflation model, is presented in detail in Sect. 3.2. In the chameleon warm inflationary model the dissipative term in the conservation equations is automatically generated due to the scalar field-radiation coupling. This also fixes the functional forms of the Hubble function, of the scalar field potential, and of the radiation temperature as functions of the coupling function $f(\phi)$, and, consequently, of the scalar field. Under the assumption that the temperature dependence of the radiation is a power law function of the scalar field, the temperature evolution can be fully determined. Moreover, the scalar

spectral index, the amplitude of tensor perturbations, and the tensor spectral index, are also obtained. Usually, the warm inflationary models are investigated in two limiting regimes, corresponding to weak and strong dissipation, respectively. In Sect. 3.3 we investigate the weak dissipative regime, for which the dissipative ratio is much smaller than unity. We perform a detailed physical and numerical analysis of the model in this approximation, by obtaining the slow-roll parameters, the scalar spectral index n_s , and the tensor to scalar ratio r as functions of the coefficients n and m (characterizing the temperature evolution), and of the number of e-folds. Using the $r - n_s$ diagram from the Planck 2018 data, the comparison of the theoretical predictions of the model with the observational data is performed, and the allowable range of parameters is obtained. The evolution of the scalar field potential and of the coupling function $f(\phi)$ is also obtained. The strong dissipative regime is investigated in a similar way in Sect. 3.4, and the allowable ranges of the model parameters are obtained from the comparison with the observational data. The coupling function has an exponential behavior. Finally, in Sect. 4 we discuss and conclude our results.

2 A brief review of the chameleon model

As stated earlier, a scalar field-driven early epoch of inflationary expansion is nowadays taken to be as a standard ingredient to describe the expansionary history of the Universe. Less standard, however, is the nature of the inflationary scenario, and of the driving scalar field. One possible candidate suggested in the literature [167–179] is the so-called chameleon scalar field, a light scalar field whose mass depends on the ambient matter density. Originally put forward to overcome the quintessence mechanism drawbacks, and as a dark energy candidate for the late-time cosmic acceleration, recent studies [180–183] have shown that chameleonic inflation is also possible with the appropriate choice of the potential for the scalar field. This leads to a varying mass for the scalar field in which in a dense environment, the scalar field acquires a large mass, which results in short range effects.

In the original chameleon model, the scalar field is coupled to the matter through a conformal factor, $\tilde{g}_{\mu\nu} = e^{\frac{2\beta\phi}{M_P}} g_{\mu\nu}$, which is also the relation between the Jordan and the Einstein frames. However, our case is different in that the scalar field is non-minimally coupled to the matter sector, such that the Lagrangian of the model is given as

$$S = \int d^4x \sqrt{-g} \left(R/2 - \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - V(\phi) + f(\phi) L_m \right).$$

Here g is the determinant of the metric tensor $g_{\mu\nu}$, R the Ricci scalar constructed from the metric, L_m is the Lagrangian of the standard matter fields, $V(\phi)$ is the potential of the

chameleon-like scalar field ϕ , and $f(\phi)$ is some generic function of the scalar field to be described shortly.

Assuming a spatially flat Friedmann–Robertson–Walker (FRW) metric, the Friedmann equations of the model are given by

$$3H^2 = \rho_\phi + f(\phi)\rho, \quad 2\dot{H} + 3H^2 = -p_\phi - f(\phi)p, \quad (1)$$

where ρ_ϕ and p_ϕ are the energy density and pressure of the scalar field, respectively, given by

$$\rho_\phi = \frac{1}{2} \dot{\phi}^2 + V(\phi), \quad p_\phi = \frac{1}{2} \dot{\phi}^2 - V(\phi), \quad (2)$$

with an overdot representing the derivative with respect to the cosmic time t . Also, taking the variation of the action with respect to the scalar field ϕ , leads to the equation of motion of the scalar field (the generalized Klein–Gordon equation),

$$\ddot{\phi} + 3H\dot{\phi} + V'(\phi) = f'(\phi)L_m, \quad (3)$$

where the prime denotes the derivative with respect to the scalar field. The matter Lagrangian in the above equation must be specified in order to close the system of cosmological equation. There are many studies on this topic such as [184–192], where the authors explain that there are two definitions for the Lagrangian of a perfect fluid as $L_m^{(1)} = -\rho$ and $L_m^{(2)} = p$, respectively. However, in our case, this degeneracy is broken due to the non-minimal interaction of the scalar field and matter, and therefore we follow the approach proposed in [192], where it is shown that only for $L_m^{(2)} = p$ one has a geodesic motion for perfect fluids. Thus, in this work, we are going to adopt this Lagrangian for the perfect fluid. Because of the interaction between the scalar field and the matter sector, the conservation equations are generalized, and they are given by

$$\begin{aligned} \frac{d}{dt} [f(\phi)\rho] + 3Hf(\phi)(\rho + p) &= -\dot{f}(\phi)p, \\ \frac{d}{dt} \rho_\phi + 3H(\rho_\phi + p_\phi) &= \dot{f}(\phi)p. \end{aligned} \quad (4)$$

3 Chameleon warm inflation

In this section, we are going to explain the general features of warm inflation, in which the inflaton is described by a chameleon-like scalar field. In warm inflation a dissipation term appears on the right hand side of the conservation equation, which is usually included by hand. However, in the chameleon model of scalar field one could naturally obtain a dissipation term, as in Eq. (4), because of the presence of the interaction between the scalar field, and the matter term in the action. In the following, we first review the warm inflationary scenario, and then we combine the scenario with the chameleon model, where the chameleon scalar field plays the role of the inflaton.

3.1 Review of warm inflation

The warm inflationary scenario describes an accelerated expansion phase of the Universe at its earliest times, in which the scalar field is the dominant component. However, its main difference with the cold inflation is that beside a scalar field, there is another component of the cosmological fluid is present, usually taken as radiation. During inflation the scalar field and the matter (radiation) component interact. In such a case, the Friedmann equations are given by

$$3H^2 = \frac{1}{2} \dot{\phi}^2 + V(\phi) + \rho_r, \quad 2\dot{H} = -\dot{\phi}^2 - \frac{4}{3} \rho_r. \quad (5)$$

Due to this interaction, energy is transferred from the scalar field to the radiation fluid, and this process is described through the following conservation equations,

$$\dot{\rho}_r + 3H(\rho_r + p_r) = \Gamma \dot{\phi}^2, \quad (6)$$

$$\dot{\rho}_\phi + 3H(\rho_\phi + p_\phi) = -\Gamma \dot{\phi}^2. \quad (7)$$

Here Γ is a dissipation coefficient. Equation (7) is usually written in a different way, known as the equation of motion of the scalar field, as

$$\ddot{\phi} + 3H(1 + Q)\dot{\phi} + V'(\phi) = 0, \quad (8)$$

where the parameter $Q = \Gamma/3H$ is the ratio of the radiation production to expansion rate.

In warm inflation, the slow-roll approximations are still at work, i.e., in order to have a quasi-de Sitter expansion, the rate of the Hubble parameter during a Hubble time is assumed to be small, a condition that is imposed via the first slow-roll parameter

$$\varepsilon_1 = -\frac{\dot{H}}{H^2}. \quad (9)$$

The energy density of the scalar field dominates over the radiation energy density, and also the kinetic term of the scalar field is negligible w.r.t its potential, i.e. $\rho_\phi \gg \rho_r$ and $\rho_\phi \simeq V(\phi)$. In addition to the assumptions that we also have in cold inflation, in warm inflationary scenario it is supposed that the radiation production is quasi-stable during inflation, so that $\dot{\rho}_r \ll H\rho_r$ and $\dot{\rho}_r \ll \Gamma \dot{\phi}^2$. Then, from Eqs. (5), (7), and (8), it follows that

$$3H^2 \simeq V(\phi), \quad (10)$$

$$\dot{\phi} \simeq -\frac{V'(\phi)}{3H(1+Q)}, \quad (11)$$

$$\rho_r = C_\gamma T^4 = \frac{\Gamma}{4H} \dot{\phi}^2, \quad (12)$$

where T is the temperature of the fluid, $C_\gamma = \pi^2 g_\star/30$ is the Stefan–Boltzman constant, and g_\star is the number of degrees of freedom of the radiation field.

With the use of the above equations, it follows that the slow-roll parameters can be expressed as

$$\varepsilon_1 = \frac{1}{2(1+Q)} \frac{V'^2(\phi)}{V^2(\phi)}, \quad \varepsilon_2 = \frac{\dot{\varepsilon}_1}{H\varepsilon_1}. \tag{13}$$

It is common to use two other slow-roll parameters, which are stated in terms of the potential as,

$$\eta = \frac{1}{(1+Q)} \frac{V''}{V}, \quad \beta = \frac{1}{(1+Q)} \frac{V'\Gamma'}{V\Gamma}. \tag{14}$$

The slow-roll parameter ε_2 is related to the above slow-roll parameters through the relation

$$\varepsilon_2 = -2\eta + 4\varepsilon_1 + \frac{Q}{(1+Q)} (\beta - \varepsilon_1).$$

Due to the appearance of the term $(1+Q)$ in the dominator of the slow-roll parameters, it follows that the smallness of the parameters are guaranteed for a large range of potentials that could satisfy the slow-roll approximations.

Smallness of the slow-roll parameters ensures us that first the Universe had a quasi-exponential accelerated expansionary phase, and also it did stand in this phase for enough number of e-folds to solve the issues of the hot big bang model. The number of e-folds is given by the following relation,

$$\begin{aligned} N &= \int_{t_*}^{t_e} H dt = \int_{\phi_*}^{\phi_e} \frac{H}{\dot{\phi}} d\phi \\ &= - \int_{\phi_*}^{\phi_e} (1+Q) \frac{V(\phi)}{V'(\phi)} d\phi, \end{aligned} \tag{15}$$

where the last equality is obtained by using Eqs. (11) and (10). The subscripts “e” and “*” respectively indicate the quantity at the end of inflation, and at the horizon crossing time, respectively.

3.2 Warm inflation with a chameleon-like scalar field

In the chameleon scalar field model, there is a source term on the right hand side of the conservation Eq. (4), which automatically appears in the equation because of the interaction between the scalar field and the matter sector. Since the chameleon scalar field is taken as the inflaton, it follows that our source term in the conservation equation of the chameleon model plays the same role of the term $\Gamma\dot{\phi}^2$ in Eq. (7). Then, we are going to equate these two terms, i.e.

$$\Gamma\dot{\phi}^2 = -\frac{1}{3} \rho_{rs} \frac{\dot{f}(\phi)}{f(\phi)}, \tag{16}$$

where the second fluid component of the very early Universe is taken as radiation, with equation of state $p_r = \rho_r/3$. On the other hand, the sign of this interaction also shows up in the Friedmann equation, i.e., via the presence of the coupling function $f(\phi)$ in Eq. (1). To get back to the original evolution equation, we can define again an effective radiation energy

density and pressure as ρ_{rs} and p_{rs} . In this case, the evolution equations are reorganized as

$$\dot{\rho}_{rs} + 3H(\rho_{rs} + p_{rs}) = -p_{rs} \frac{\dot{f}(\phi)}{f(\phi)} = \Gamma\dot{\phi}^2, \tag{17}$$

$$\dot{\rho}_{\phi} + 3H(\rho_{\phi} + p_{\phi}) = p_{rs} \frac{\dot{f}(\phi)}{f(\phi)} = -\Gamma\dot{\phi}^2. \tag{18}$$

Assuming a quasi-stable production of the radiation, so that $\dot{\rho}_{rs} \ll H\rho_{rs}$ and $\dot{\rho}_{rs} \ll \Gamma\dot{\phi}^2$, which implies that $\dot{\rho}_{rs} \ll p_{rs}\dot{f}(\phi)/f(\phi)$, from Eq. (17) one can calculate the Hubble parameter as

$$H = -\frac{1}{12} \frac{\dot{f}(\phi)}{f(\phi)}. \tag{19}$$

Using Eq. (10), the potential of the scalar field is derived in terms of $f(\phi)$ and its derivative as

$$V(\phi) = \frac{1}{48} \left[\frac{\dot{f}(\phi)}{f(\phi)} \right]^2. \tag{20}$$

On the other hand, using the result of QFT in curved space, the dissipation coefficient is taken as its best known expression $\Gamma = \Gamma_0 T^m / \phi^{m-1}$. Then, by applying Eqs. (11), (12) and (16), the temperature of the radiation fluid is obtained as

$$T^{m-4} = \frac{C_{\gamma}}{\Gamma_0} (1+Q) H \frac{f'(\phi)}{f(\phi)} \frac{\phi^{m-1}}{V'(\phi)}, \tag{21}$$

where we have used the mathematical result $\dot{f}(\phi) = f'(\phi)\dot{\phi}$.

To go one step further, we assume the ansatz $\dot{f}(\phi)/f(\phi) = \alpha\phi^n$. With the use of this assumption in Eqs. (19) and (20), the Hubble parameter and the potential of the scalar field are obtained as

$$H = -\frac{\alpha}{12} \phi^n, \quad V(\phi) = \frac{\alpha^2}{48} \phi^{2n}. \tag{22}$$

Since the Hubble parameter is positive, α should be a negative constant.

On the other hand, with the use of Eqs. (11), (19) and (22) in Eq. (21), one arrives at the following non-linear equation for the temperature of the fluid,

$$\begin{aligned} T^m + \frac{\phi^{\frac{m+3n-3}{2}}}{\xi\Gamma_0} T^{\frac{m-4}{2}} - \frac{\alpha}{4\Gamma_0} \phi^{m+n-1} &= 0, \\ \xi &\equiv \left(-\frac{192C_{\gamma}}{n^2\alpha^3\Gamma_0} \right)^{1/2}. \end{aligned} \tag{23}$$

Due to the fact that the third term of the above equation is a power-law function of the scalar field, it follows that the temperature $T(\phi)$ should be a power-law function of the scalar field as well, in order to satisfy the equation. Hence, for the temperature we assume the functional form $T(\phi) = T_0\phi^q$. By substituting T in Eq. (21), we obtain

$$T_0^m \phi^{mq} + \frac{T_0^{\frac{m-4}{2}}}{\xi\Gamma_0} \phi^{\frac{m+3n-3}{2} + \frac{q(m-4)}{2}} - \frac{\alpha}{4\Gamma_0} \phi^{m+n-1} = 0.$$

Then, the power of the scalar field of the three terms should be equal, which gives the following relations:

- (a) equality of the first and second terms: $q = \frac{m+3n-3}{m+4}$;
- (b) equality of the first and third terms: $q = \frac{m+n-1}{m}$; and
- (c) equality of the second and third terms: $q = \frac{m-n+1}{m-4}$.

Hence, if the three conditions are satisfied, the proposed function for the temperature might be a proper solution. However, it can be shown that these conditions actually give one constraint that is extracted as follows:

- a = b: $m = \frac{2(n-1)}{n-3}$;
- a = c: $m = \frac{2(n-1)}{n-3}$; and
- b = c: $m = \frac{2(n-1)}{n-3}$.

Therefore, for the proposed functional form of the temperature, the following constraints must be satisfied:

$$m = \frac{2(n-1)}{n-3}, \quad q = \frac{n-1}{2}, \tag{24}$$

which show that by determining the power n , the power parameters m and q are also determined.

In the warm inflationary scenario, there are both quantum and thermal fluctuations. Thermal fluctuations are dependent on the fluid temperature T , and quantum fluctuations depend on the Hubble parameter H . A feature of warm inflation is that the fluid temperature is bigger than the Hubble parameter, $T > H$, stating that the thermal fluctuations overcome quantum fluctuations, and become the origin of the Universe’s LSS. To ensure that we stay in the warm inflation regime, the condition $T/H > 1$

$$\frac{T(\phi)}{H(\phi)} = -\frac{12T_0}{\alpha} \phi^{q-n} > 1, \tag{25}$$

must be satisfied during the cosmological expansion.

The amount of cosmic expansion during inflation is measured through the number of e-folds N defined as

$$\begin{aligned} N &= \int_{t_*}^{t_e} H dt = \int_{\phi_*}^{\phi_e} \frac{H}{\dot{\phi}} d\phi = -\frac{1}{12} \int_{\phi_*}^{\phi_e} \frac{f'(\phi)}{f(\phi)} d\phi \\ &= -\frac{1}{12} \ln \left[\frac{f(\phi_e)}{f(\phi_*)} \right]. \end{aligned} \tag{26}$$

Once the expression of the scalar field at the end of inflation is obtained from the relation $\varepsilon_1(\phi_e) = 1$, Eq. (26) is used to derive the scalar field at the time of the horizon crossing in terms of the number of e-folds. Then, all the perturbation parameters could be expressed in terms of the number of e-folds.

To test the validity of a theoretical model, one has to compare its predictions with the observational data. In this

regard, we will obtain some important perturbations parameters, such as the amplitude of scalar perturbations, the scalar spectral index and the tensor-to-scalar ratio, and they will be compared with the Planck-2018 data. Following [94,95], the amplitude of scalar perturbations is calculated as

$$\mathcal{P}_s = \left(\frac{H^2}{2\pi\dot{\phi}} \right)^2 \left(1 + 2n_{BE} + \frac{2\sqrt{3}\pi Q}{\sqrt{3+4\pi Q}} \frac{T}{H} \right) G(Q), \tag{27}$$

where n_{BE} is the Bose–Einstein distribution, given by $n_{BE} = [\exp(H/T_{\delta\phi}) - 1]^{-1}$, and $T_{\delta\phi}$ is the temperature of the inflaton fluctuations [95]. Here $G(Q)$, giving the growth of the fluctuations, is a function of Q and its presence is due to the coupling of the scalar field and radiation [94,95]. The scalar spectral index and its running are obtained from the amplitude of scalar perturbations, and it is defined as

$$n_s - 1 = \frac{d \ln(\mathcal{P}_s)}{d \ln(k)}, \quad \alpha_s = \frac{dn_s}{d \ln(k)}. \tag{28}$$

Tensor perturbations, known as gravitational waves, are measured indirectly through the tensor-to-scalar ratio parameter $r = \mathcal{P}_t/\mathcal{P}_s$. The amplitude of tensor perturbations is given by [94]

$$\mathcal{P}_t = \frac{2H^2}{\pi^2}. \tag{29}$$

The tensor spectral index is defined as

$$n_t = \frac{d \ln(\mathcal{P}_t)}{d \ln(k)}. \tag{30}$$

In the relevant literature, the scenario of warm inflation is usually considered in two regimes, known as the strong and the weak dissipative regimes, where the dissipative ratio is respectively $Q \gg 1$ and $Q \ll 1$. In the subsequent subsections, we are going to consider the discussed model in these two regimes.

3.3 Weak dissipative regime

In the weak dissipative regime, the dissipative ratio is much smaller than unity, i.e., $Q \ll 1$ ($\Gamma \ll 3H$), and thus we have $(1 + Q) \simeq 1$. This approximation makes the evolution equation easier to analyse, since the time derivative of the scalar field is expressed as

$$\dot{\phi} = \frac{n\alpha}{6} \phi^{n-1}. \tag{31}$$

Also, by taking this approximation into account, and following the same process as used previously, the temperature of the fluid is obtained easily as a function of the scalar field:

$$T^{m-4}(\phi) = \frac{-12C_\gamma}{n^2\alpha\Gamma_0} \frac{1}{\phi^{n-m-1}}. \tag{32}$$

In subsection (3.2), an ansatz was introduced for the model. Using this definition and applying Eq. (31), the coupling function $f(\phi)$ is obtained as an exponential function,

$$f(\phi) = f_0 \exp\left(\frac{3}{n} \phi^2\right), \tag{33}$$

where f_0 is a constant of integration.

The slow-roll parameters, introduced in Eqs. (13) and (14), are obtained as

$$\begin{aligned} \varepsilon_1(\phi) &= \frac{2n^2}{\phi^2}, \quad \eta(\phi) = \frac{(2n-1)}{n} \frac{1}{\phi^2}, \\ \beta(\phi) &= \frac{(6m-nm-4)}{n(m-4)} \varepsilon_1(\phi). \end{aligned} \tag{34}$$

Inflation ends for $\phi_e^2 = 2n^2$, and the scalar field at the horizon crossing is extracted from the number of e-folds. Hence from Eq. (26) one arrives at the expression,

$$N = \int_{\phi_*}^{\phi_e} \frac{H}{\dot{\phi}} d\phi = -\frac{1}{12} \int_{\phi_*}^{\phi_e} \frac{f'(\phi)}{f(\phi)} d\phi, \tag{35}$$

which subsequently gives

$$\phi_*^2 = 2n^2 \left(1 + \frac{2N}{n}\right). \tag{36}$$

In the weak dissipative regimes, the parameter $G(Q) \simeq 1$ at the time of horizon crossing. Then, from Eq. (27), the amplitude of scalar perturbations in these regimes is found as

$$\mathcal{P}_s = 2 \left(\frac{H^2}{2\pi\dot{\phi}}\right)^2 \frac{T}{H}. \tag{37}$$

From Eq. (28) the scalar spectral index and its running in the weak regimes are obtained as

$$n_s - 1 = -\frac{17}{8} \varepsilon_1 + \frac{3}{2} \eta - \frac{1}{4} \beta, \tag{38}$$

where η is defined through the slow-roll parameters ε_2 as $\eta = V''/(1+Q)V$, so that $\varepsilon_2 = \eta - \varepsilon_1$, and the slow-roll parameter β is given as $\beta = V'\Gamma'/(1+Q)V\Gamma$. The tensor-to-scalar ratio in this case is also obtained as

$$r = 8\varepsilon_1 \frac{H}{T}. \tag{39}$$

To compare the results of the model with observational data, the above perturbation parameters are computed at the horizon crossing. Using Eq. (36), the slow-roll parameters at horizon crossing are obtained as follows:

$$\varepsilon_1^* = \left(1 + \frac{2N}{n}\right)^{-1}, \tag{40}$$

$$\eta^* = \frac{2n-1}{n} \varepsilon_1^*, \tag{41}$$

$$\beta^* = -\frac{nm-6m+4}{n(m-4)} \varepsilon_1^*. \tag{42}$$

Substituting the above slow-roll parameters in Eq. (38), the scalar spectral index is given as a function of n, m and e-folding number N as

$$\begin{aligned} n_s(n, m, N) &= 1 - \left(\frac{17}{8} - \frac{3(2n-1)}{2n} - \frac{nm-6m+4}{4n(m-4)}\right) \\ &\times \left(1 + \frac{2N}{n}\right)^{-1}. \end{aligned} \tag{43}$$

From Eq. (37), and by using Eq. (36), the model constant parameter Γ_0 is found in terms of n, m and N as

$$\begin{aligned} \Gamma_0(n, m, N) &= -\frac{12C_\gamma}{\alpha n^2} \left(-\frac{\alpha}{96\pi^2 n^2 \mathcal{P}_s}\right)^{m-4} \\ &\times \left[2n^2 \left(1 + \frac{2N}{n}\right)\right]^{\frac{nm-5n+3m-7}{2}}, \end{aligned} \tag{44}$$

where \mathcal{P}_s is the amplitude of the scalar perturbations at horizon crossing. According to the latest observational data $\mathcal{P}_s = 2.17 \times 10^{-9}$. The tensor-to-scalar ratio for the same time is found from Eq. (39) as

$$\begin{aligned} r^*(n, m, N) &= -\frac{4\alpha n^2}{3} \left(-\frac{\alpha n^2}{12C_\gamma}\right)^{\frac{1}{m-4}} \Gamma_0^{\frac{1}{m-4}}(n, m, N) \\ &\times \left[2n^2 \left(1 + \frac{2N}{n}\right)\right]^{\frac{nm-3(n+m)+7}{2(m-4)}}. \end{aligned} \tag{45}$$

Using the $r - n_s$ diagram of Planck-2018, one could plot a $n - m$ diagram as shown in Fig. 1, where the dark blue color indicates an area of (n, m) where the results for n_s and r stand in 68% CL. The light blue color indicates an area of (n, m) in which the point (r, n_s) of the model stands in 95% CL.

Inserting the scalar field at horizon crossing in Eq. (20), the potential of the scalar field could be obtained in terms of

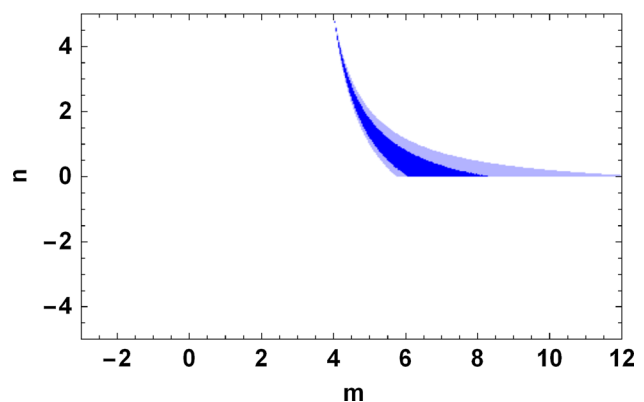


Fig. 1 Numerical values of the (n, m) parameters of the chameleon field driven warm inflation model in the weak dissipative regime for which the point $(r - n_s)$ is located in the observational region. The dark blue color shows the values of (n, m) in the 68% CL range of the Planck-2018 data, while the light blue color shows the parameter values in the range 95% CL

N for specific values of (n, m) that have been taken from Fig. 2. The results in the Figure indicate that the inflation did begin at an energy scale about 10^{15-16} GeV. At the onset of inflation, the scalar field stands on the top, and by passing time it rolls down slowly until the end of inflation.

Inserting the scalar field at horizon crossing in Eq. (33), the coupling function could be depicted versus N . Getting some point of (n, m) from Fig. 1, the behavior of the coupling function $f(\phi)$ in the WDR is illustrated in Fig. 3. The curves show the behavior of $\ln(f(N))$ versus $\ln(N)$ during the inflationary era; however, the scales on the axes show the actual ranges of $f(N)$ and N . The coupling function $f(N)$ has an exponential behavior, and it grows rapidly by enhancement of N .

An important feature of the warm inflationary scenario is that the thermal fluctuations overcome the quantum fluctuations, since the fluid temperature is bigger than the Hubble parameter. To have a healthy warm inflation, this condition should be satisfied during the cosmological evolution. Figure 4 describes the behavior of the ratio of the temperature to the Hubble parameter during this era. From the Figure one can see that the condition is satisfied during this phase of cosmological expansion.

3.4 Strong dissipative regime

In this regime, the dissipative ratio is much larger than unity, so that $(1 + Q) \simeq Q$. In this case, the analysis of the model becomes more complicated, and the cosmological dynamics is different from the weak dissipative regime. In this approximation the time derivative of the scalar field is given by

$$\dot{\phi} = -\frac{V'(\phi)}{\Gamma}. \tag{46}$$

Using Eq. (16), and by considering the definition of Γ , and Eq. (46), the temperature of the radiation fluid is obtained as

$$T^{m+4}(\phi) = -\frac{n^2\alpha^3}{192\Gamma_0 C_\gamma} \phi^{3n+m-3} = T_0^{m+4} \phi^{3n+m-3}. \tag{47}$$

Then, with the help of this relation, the time derivative of the scalar field can be found in terms of the scalar field, so that

$$\dot{\phi} = -\frac{n\alpha^2}{24\Gamma_0 T_0^m} \phi^{\frac{8n-nm+5m-8}{m+4}}. \tag{48}$$

The dissipative ratio Q is given by

$$Q(\phi) = \frac{\Gamma}{3H} = -\frac{4\Gamma_0 T_0^m}{\alpha} \phi^{\frac{2(nm-2n-3m+2)}{m+4}}. \tag{49}$$

From Eqs. (13) and (14), after substituting Eq. (49), we find out the slow-roll parameters as

$$\varepsilon_1(\phi) = -\frac{\alpha n^2}{2\Gamma_0 T_0^m} \phi^{\frac{-2(nm-2n-2m+6)}{m+4}}, \tag{50}$$

$$\eta(\phi) = \frac{(2n-1)}{n} \varepsilon_1(\phi), \tag{51}$$

$$\beta(\phi) = \frac{m(3n-6)+4}{n(m+4)} \varepsilon_1(\phi). \tag{52}$$

Then, the scalar field at the end of inflation is obtained as

$$\phi_e^{\frac{2(nm-2n-2m+6)}{m+4}} = -\frac{\alpha n^2}{2\Gamma_0 T_0^m}.$$

After obtaining $\dot{\phi}(\phi)$ and $Q(\phi)$, the coupling function $f(\phi)$ can also be determined through the definition introduced previously. It turns out that $f(\phi)$ is again an exponential function of the scalar field, and it is given by

$$f(\phi) = f_0 \exp \left(-\frac{12\Gamma_0 T_0^m}{n\alpha} \frac{m+4}{nm-2n-2m+6} \phi^{\frac{2(nm-2n-2m+6)}{m+4}} \right), \tag{53}$$

where f_0 is an arbitrary constant of integration. Using this result in the relation of N , after integration, the scalar field at the time of horizon crossing is obtained in terms of N as

$$\phi_\star^{\frac{2(nm-2n-2m+6)}{m+4}} = -\frac{\alpha n^2}{2\Gamma_0 T_0^m} \times \left(1 + \frac{2(nm-2n-2m+6)}{n(m+4)} N \right). \tag{54}$$

In the strong dissipative regime, the parameter $G(Q)$, depending on the different values of the parameter m , can be approximated as [94, 95]:

$$\begin{aligned} m = 1 &\longrightarrow G(Q) \simeq 1 + 0.127Q^{4.330} + 4.981Q^{1.946}, \\ m = 3 &\longrightarrow G(Q) \simeq 1 + 0.0185Q^{2.315} + 0.335Q^{1.364}, \\ m = -1 &\longrightarrow G(Q) \simeq \frac{1 + 0.4Q^{0.77}}{(1 + 0.15Q^{1.09})^2}. \end{aligned}$$

Therefore, in a more convenient way, we can write down the function $G(Q) = a_m Q^{b_m}$, where

$$\begin{aligned} m = 1 &\longrightarrow a_m = 0.127, \quad b_m = 4.330, \\ m = 3 &\longrightarrow a_m = 0.0185, \quad b_m = 2.315, \\ m = -1 &\longrightarrow a_m = 17.78, \quad b_m = -1.41. \end{aligned}$$

For the case $m = -1$, the power of the term Q is negative stating that the term $G(Q)$ tends to zero which in turn lead the amplitude of the scalar perturbation to zero. So, this case is not considered here.

Hence the amplitude of the scalar perturbations is obtained as

$$\mathcal{P}_s = \left(\frac{H^2}{2\pi\dot{\phi}} \right)^2 \sqrt{3\pi} \frac{T}{H} \times a_m Q^{b_m}. \tag{55}$$

By taking the time derivative of this equation according to Eq. (28) leads to the scalar spectral index and its running, given by

Fig. 2 The behavior of the scalar field potential $V(\phi)$ versus ϕ of the chameleon field driven warm inflation model in the weak dissipative regime for different values of n during inflation. The figure shows that for $n = 2$, the energy scale of inflation is about 10^{15} GeV. Once inflation is approaching to its end, the potential decreases in time

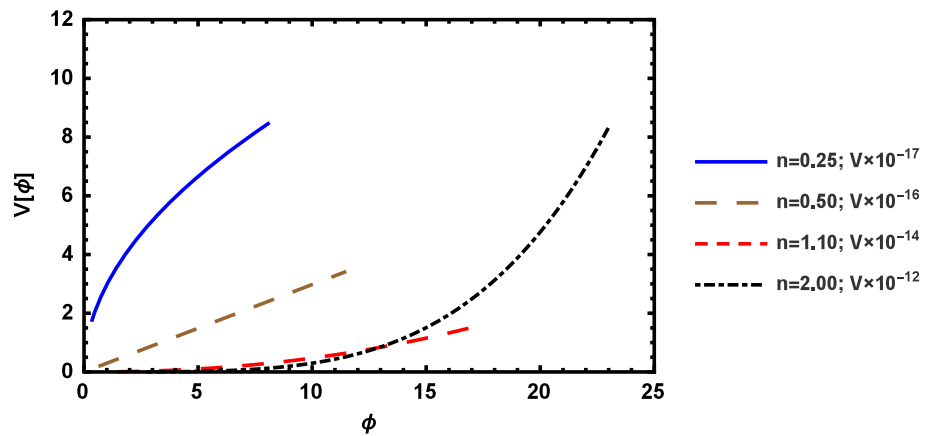


Fig. 3 The behavior of the coupling function $f(N)$ versus the number of e-fold for different values of n (in a logarithmic scale). The function has an exponential behavior and when approaching to the end of inflation it rapidly decreases

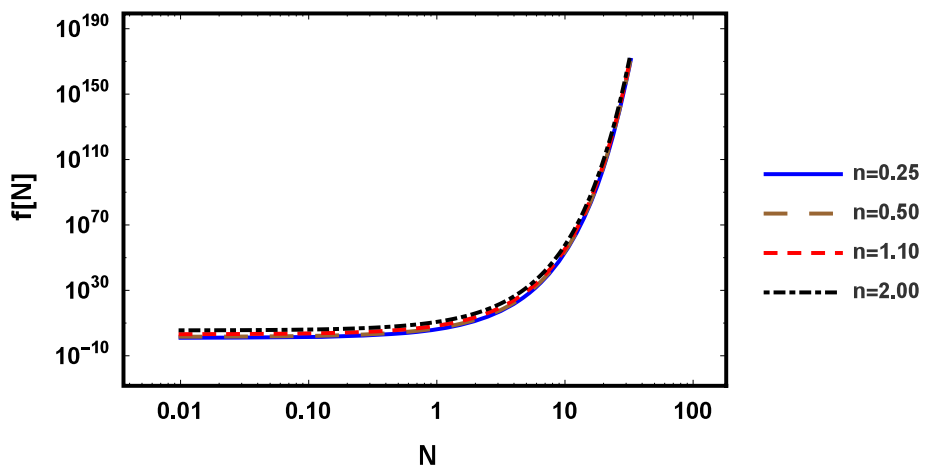
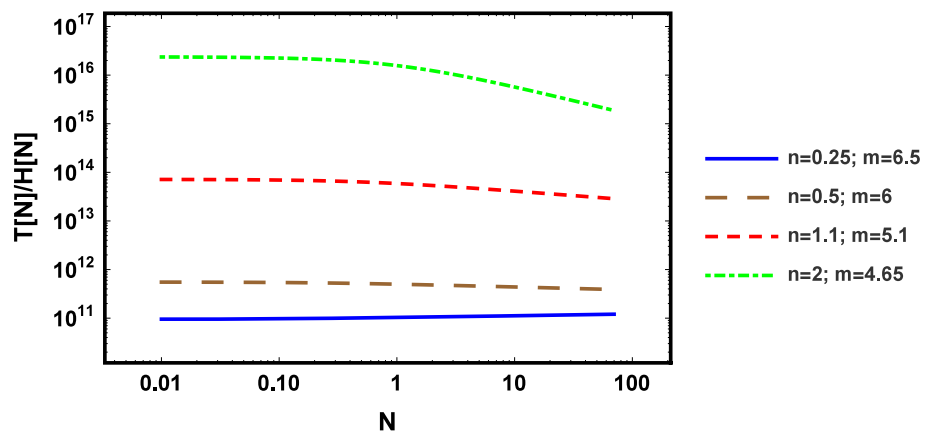


Fig. 4 The ratio of the temperature to the Hubble parameter during the inflationary period of the chameleon field driven warm inflation model in the weak dissipative regime versus the number of e-folds for different values of (n, m) . As one can see from the plots during inflation the temperature is larger than the Hubble parameter, and the condition $T > H$ is satisfied properly



$$n_s - 1 = -2q_1\varepsilon_1 - q_2\beta + \frac{3}{2}\eta, \tag{56}$$

where $q_1 = \frac{9}{8} - \frac{b_m}{2}$ and $q_2 = b_m + \frac{7}{4}$. We notice that the function $G(Q)$ for $m = 3$ has an acceptable behaviour for the value $m = -1$, or even for $m = 1$. The amplitude of tensor perturbations in strong dissipative regimes is given by [94,95]

$$\mathcal{P}_t = \frac{2H^2}{\pi^2} = \frac{\alpha^2}{72\pi^2}\phi^{2n}. \tag{57}$$

Then, the tensor-to-scalar ratio is obtained as follows:

$$r = \frac{24(\sqrt{3})^{b_m}}{\sqrt{3\pi} a_m} \left(\frac{4C_\gamma}{9}\right)^{\frac{1}{4}} \frac{V^{3/2}}{V^{(q_1-1)}\Gamma^{q_2}}. \tag{58}$$

Using Eq. (54), the above perturbation parameters can be obtained at the horizon crossing time. The slow-roll parameters in the SDR at this cosmological instance are given by

$$\varepsilon_1^* = \left(1 + \frac{2(nm - 2n - 2m + 6)N}{n(m + 4)}\right)^{-1}, \tag{59}$$

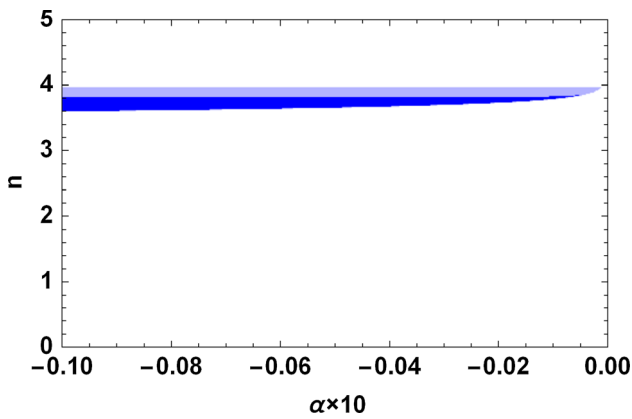


Fig. 5 The ranges of the parameters (n, α) of the chameleon field driven warm inflation model in the strong dissipative regime that are consistent with n_s and r located in the observational area. The light blue color indicates the areas with 95% CL, while the dark blue color indicates the areas with 68% CL

$$\eta^* = \frac{2n - 1}{n} \varepsilon_1^*, \tag{60}$$

$$\beta^* = \frac{3nm - 6m + 4}{n(m + 4)} \varepsilon_1^*. \tag{61}$$

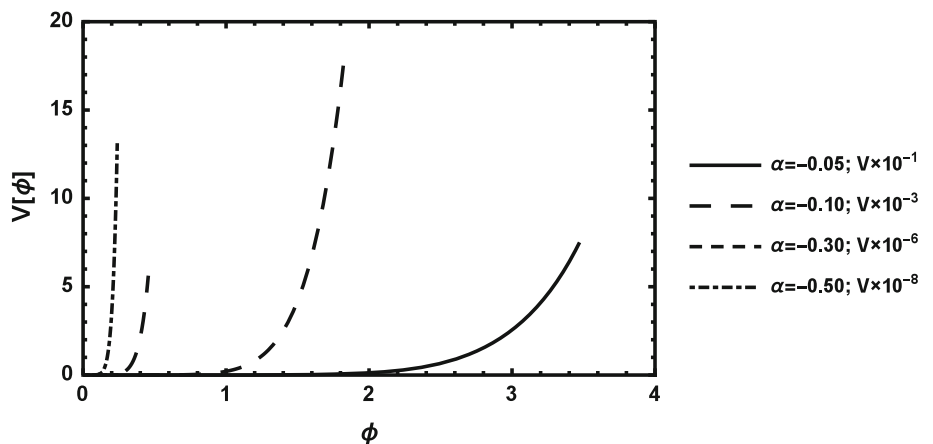
Hence we obtain the following scalar spectral index,

$$\begin{aligned} n_s(n, \alpha, N) &= 1 - \left(2q_1 + \frac{3nm - 6m + 4}{n(m + 4)} q_2 - \frac{3(2n - 1)}{2n} \right) \varepsilon_1^*. \end{aligned} \tag{62}$$

We notice that this is a function of n, m and N . The constant parameter Γ_0 can be determined by using the observational data for the amplitude of the scalar perturbations as

$$\Gamma_0(n, \alpha, N) = \frac{B}{\mathcal{P}_s^*} \left(-\frac{n^2 \alpha}{2} \left[-\frac{192C_\gamma}{n^2 \alpha^3} \right]^{\frac{m}{m+4}} \frac{1}{\varepsilon_1^*} \right)^{\frac{(m+4)\sigma}{2S(nm-2n-2m+6)}}, \tag{63}$$

Fig. 6 The behavior of the scalar field potential $V(\phi)$ versus the inflation scalar field in the chameleon field driven warm inflation model in the strong dissipative regime for different values of (n, α) , obtained from Fig. 5. The figure shows that the inflation could start from energy scales of the order of 10^{15} GeV, with the scalar field potential decreasing during the cosmological evolution



where

$$\begin{aligned} B(n, \alpha, N) &= \frac{1}{(2n)^{3/2}} \left(\frac{9}{4C_\gamma} \right)^{\frac{1}{4}} \frac{\sqrt{3\pi} a_m}{36(\sqrt{3})^{b_m} \pi^2} \left(\frac{\alpha^2}{48} \right)^{q_1 - \frac{3}{2}} \\ &\times \left[-\frac{192C_\gamma}{n^2 \alpha^3} \right]^{\frac{-mq_2}{m+4}}, \end{aligned}$$

$$S(n, \alpha, N) = \frac{4q_2}{m + 4} - \frac{4\sigma}{(nm - 2n - 2m + 6)},$$

and

$$\sigma(n, \alpha, N) = -\frac{3(2n - 1)}{2} + 2nq_1 + \frac{(3nm - 6m + 4)}{m + 4} q_2,$$

respectively. Using the above relations, the tensor-to-scalar ratio at the horizon crossing is obtained as

$$\begin{aligned} r^*(n, \alpha, N) &= \frac{\alpha^2}{72\pi^2} \frac{B^{-1}}{\mathcal{P}_s^*} \\ &\times \left(-\frac{n^2 \alpha}{2} \left[-\frac{192C_\gamma}{n^2 \alpha^3} \right]^{\frac{m}{m+4}} \frac{1}{\varepsilon_1^*} \right)^{\frac{(m+4)(2n-\sigma)}{2(nm-2n-2m+6)}} \Gamma_0^{\tilde{S}}, \end{aligned} \tag{64}$$

where we have defined

$$\tilde{S}(n, \alpha, N) = -\frac{4q_2}{m + 4} - \frac{2(2n - \sigma)}{(nm - 2n - 2m + 6)}.$$

Similarly as in the previous section, we use the $r - n_s$ diagram of Planck-2018 to find out the proper region of (n, α) , which leads for the model prediction to be agreement with observational data. Figure 5 shows this area, depicted with dark and light blue colors, respectively, displaying the values of (n, α) compatible with 68% and 95% CL, respectively. With the use of the value of the scalar field at horizon crossing, i.e. substituting Eq. (54) in Eq. (20), and by choosing the proper values of (n, α) from Fig. 5, the scalar field potential is expressed in terms of the number of e-folds. Figure 6 depicts the behavior of the scalar field potential during inflation, indicating that the inflation energy scale could be about 10^{15-16} GeV.

Fig. 7 The behavior of coupling function $f(N)$ versus the number of e-folds in the chameleon field driven warm inflation model in the strong dissipative regime for different values of (n, α) , which are selected from Fig. 5. The function $f(N)$ has an exponential behavior, and during inflation it decreases

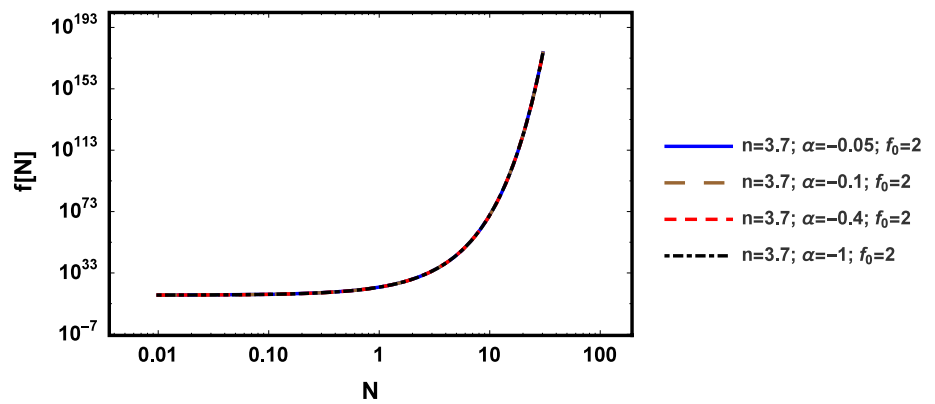
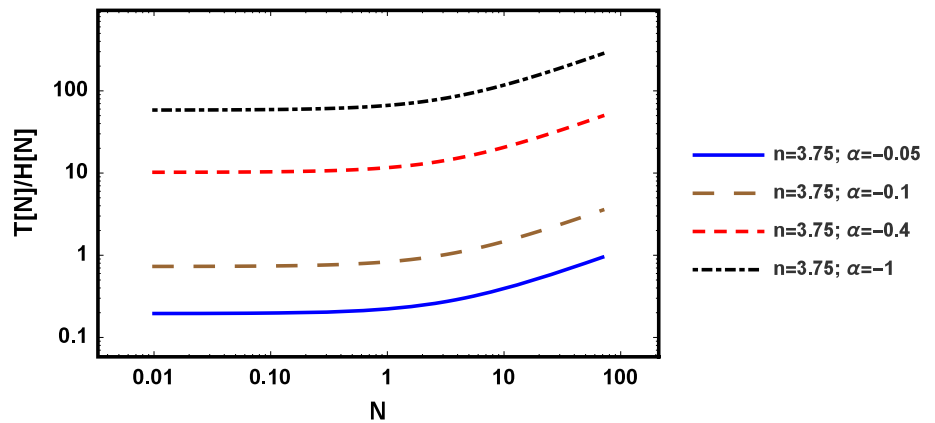


Fig. 8 The ratio of the temperature to the Hubble parameter during the inflationary evolution of the chameleon field driven warm inflation model in the strong dissipative regime versus the number of e-folds for different values of (n, α) . During inflation the temperature is bigger than the Hubble parameter, and their ratio increases when approaching the end of inflation



By using the value of the scalar field at horizon crossing, i.e. Eq. (54) in Eq. (53), the coupling function $f(\phi)$ can again be expressed in terms of the number of e-folds. Figure 7 portrays the behavior of the coupling function during inflation for different values of (n, α) that have been selected from Fig. 5. The curves depict the behavior of $\ln[f(N)]$ versus $\ln(N)$ during the inflationary times, however, the scales show the true values of $f(N)$ and N . The coupling function $f(N)$ has an exponential behavior and it grows rapidly by the enhancement of N .

As the final step, the condition $T/H > 1$ is considered by plotting the ratio of the temperature to the Hubble parameter versus the number of e-folds during inflation, as shown in Fig. 8. From the Figure it immediately follows that the condition is satisfied during the cosmological evolution.

4 Conclusion

In the present work, we have studied the behavior of chameleon-like scalar fields, non-minimally interacting with matter, as potential candidates for driving warm inflationary evolution. Since the scalar field interacts with the matter sector, in these models one has a generalized energy conservation law, where there is a source for the scalar field,

depending on the scalar field – matter coupling function, and on the matter Lagrangian. Besides, due to this interaction, the mass of the scalar field strongly depends on the environment energy density, so that on cosmological length scales, the scalar field mass is small, while on short scales, like, for example, around the Earth, the scalar field mass grows. This interesting feature of the chameleon scalar field model may be of major interest for cosmology.

The scenario of warm inflation is a different story compared to the cold one. Although this model has almost the same basic assumptions, here the scalar field interacts with the matter sector during inflation, and energy transfer from the scalar field to the matter part does occur during the accelerated expansion phase. Due to this feature, the chameleon scalar field model seems to be a suitable model for describing warm inflation, since it naturally predicts such an interaction. By using quantum field theory and particle physics, it can be shown that the dissipation coefficient, in general, depends on both scalar field and temperature. Applying this result, and imposing an appropriate ansatz, the Hubble parameter, the scalar field potential, the coupling function, and the temperature can be obtained in terms of the scalar field. Then, the chameleon field driven warm inflationary scenario was considered in detail in two approximations, corresponding to the weak and strong dissipative regimes, respectively. The main

perturbation parameters, such as the amplitude of scalar perturbations, the scalar spectral index, and tensor-to-scalar ratio were derived for each regime. After computing the scalar spectral index and tensor-to-scalar ratio, and by using the $r - n_s$ diagram of Planck-2018, we plotted the area of the model parameters (n, m) and (n, α) that places the model exactly in the 68% and 95% CL areas, which are depicted in Figs. 1 and 5, respectively, corresponding to the weak and strong dissipative regimes.

Using the results on the model parameters, the scalar field potentials were obtained for both regimes. Our results show that the inflation energy scale could be around 10^{15} GeV, and the scalar field stands on the top of its potential, and rolls down slowly during the cosmological evolution, when and approaching to the end of inflation. Also, the results of the investigation of the coupling function $f(\phi)$ shows that it has an exponential behavior during inflation in both regimes.

Another difference between warm inflation and cold inflation is about the type of perturbations. In cold inflation, the quantum perturbations are produced during inflation, and these perturbations are the seeds for LSS of the Universe. On the other hand, in warm inflation, there are both quantum and thermal perturbations, and it is the thermal perturbations that dominate, and are the basic seeds for the formation of the LSS of the Universe. This domination is expressed by the condition $T > H$, otherwise, the quantum fluctuations are overwhelmed and we do not have warm inflation anymore. So, this condition as an important feature of warm inflation was studied as well. Our investigations show that during inflation the temperature is larger than the Hubble parameter, and this result has been illustrated in Figs. 4 and 8, respectively.

The initial conditions of the inflation, as well as the matter content at the moment of the birth of the Universe are not fully known. Similarly, there are many theoretical and observational uncertainties in the physics of the decay and interactions of the scalar fields. Probably a full understanding of these processes may need physics beyond the Standard Model of particle physics to account for the radiation and other forms of matter creation process. And, to constrain the theoretical models, and find out what truly happened in the very early Universe, certainly more accurate observational data are needed. In the present investigation we have proposed some basic tools that may open the possibility of an in-depth comparison of the predictions of the theoretical models of inflation with the cosmological data.

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