

Viability of warm inflation with standard model interactions

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The minimal warm inflation scenario proposed in [K. V. Berghaus *et al.*, Warm inflation with the standard model, [arXiv:2503.18829](#).]—featuring an axionlike inflaton coupled to Standard Model (SM) gluons via the standard interaction $\phi G\tilde{G}$ —offers a compelling bridge between inflationary dynamics and SM particle content. While the model retains only the inflaton as a beyond-SM field, its original analysis relied on some approximate treatments of warm inflation’s (WI) dynamics. Here, we revisit this scenario using `wi2easy`, a precision computational tool for WI dynamics [G. S. Rodrigues and R. O. Ramos, `wi2easy`: Warm inflation dynamics made easy, [arXiv:2504.17760](#).], to rigorously evaluate the model’s viability and full range of model’s parameters compatible with the observational parameters. Overall, we find that the results of Berghaus *et al.* hold, but with significant differences in the weak and strong dissipative regimes of WI.

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I. INTRODUCTION

Warm inflation (WI) is a cosmological framework where the inflaton field interacts continuously with a thermal radiation bath during the accelerated expansion phase, generating entropy (particles) through dissipative dynamics rather than requiring a separate reheating epoch [1–4]. Unlike cold inflation (CI), which assumes a near-zero temperature universe until reheating begins, WI maintains a thermalized environment via particle interactions, with dissipation encoded in a coefficient Υ , which typically can be a function of the temperature and the inflaton amplitude, besides of microscopic parameters entailing the microscopic derivation of it. The early ground breaking work on WI has shown how thermal fluctuations—rather than quantum ones—could seed primordial density perturbations, while key developments included deriving consistent perturbation equations for dissipative systems [5–8] and calculating modified observational signatures. In order to preserve the flatness of the inflaton potential and prevent it from receiving large thermal corrections which could disrupt the inflationary attractor solution, early models of WI explored SUSY-inspired interactions (for a review, see, e.g., Ref. [9]), while later studies tried to better integrate WI

with particle physics frameworks making use of interactions where the inflaton would be a pseudo-Goldstone boson [10,11] or to have axionlike couplings to gauge fields [12], such as to protect the inflaton potential from large quantum and thermal corrections due to the models’ symmetries¹ (for a recent review on model building proposals for WI, see, e.g., Ref. [16]).

Recently, the authors of Ref. [17] have proposed a WI construction with a minimal extension of the standard model (SM) of particle physics, where the inflaton is coupled to gluons via the standard $\phi G\tilde{G}$ axionlike interaction. Some earlier attempts of building WI and WI motivated models with SM interactions include for example Refs. [18–21]. In particular, in the Ref. [20], the authors considered the minimal supersymmetric standard model, with the inflaton considered to be a flat direction in the model. The model was, however, unsuccessful for WI because of the presence of large thermal corrections to the inflaton potential.² The work of Ref. [21] implemented a SM WI extension with a neutrino portal and using the symmetries and construction of the warm little inflaton model proposed in the Refs. [10,11]. The possibility of having a WI model built with genuine SM interactions is a major result. Constructing a WI model within the SM framework offers compelling theoretical and phenomenological advantages. By leveraging known SM interactions

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¹For earlier WI models based on axion inflation, see also Refs. [13–15].

²Note, however, that one of the first working models for WI was based on supersymmetry and proposed in Ref. [22], with dissipation coefficients originally derived in Ref. [23].

and minimal extensions of them—such as the inflaton coupling to the gluons via the axionlike interaction—this model not only minimizes beyond-the-SM (BSM) particle content, but also enhances the testability of the WI general predictions and reduces possible fine-tuning in previous WI constructions. In this context, the thermal bath in WI naturally arises from SM particle production, sidestepping *ad hoc* reheating mechanisms, while providing a seamless transition to the radiation-dominated era. Crucially, the SM-derived dissipation coefficient comes from sphaleron decay, which is a most studied subject in the context of gauge field thermodynamics [24–28]. Furthermore, by having WI embed with SM interactions enables a formulation grounded in particle physics experimentally, enabling cross-validation through collider constraints or early universe thermodynamics.

We revisit the WI scenario with SM interactions proposed in Ref. [17] using the `WI2easy` code [29], which replaces the explicit stochastic perturbation methods commonly used previously in WI [5,6,8,30] with a deterministic Fokker-Planck formalism [31,32]. While mathematically equivalent to prior approaches, the Fokker-Planck framework offers two key advantages: simplified numerical implementation and enhanced precision. `WI2easy` is then a precision computational tool for WI dynamics. As such, we can rigorously evaluate the model’s viability, bypassing approximations used in the analysis done in Ref. [17]. In particular, we quantify deviations from Ref. [17] approximations, revealing some significant differences in both weak and strong dissipation regimes of WI, although in intermediate regimes the results of that reference holds approximately well.

The organization of the paper is as follows. In Sec. II, we briefly introduce the construction of Ref. [17]. We also take advantage of some of the well-known thermodynamics and kinetic theory for quantum chromodynamics (QCD) to explore advances in the study considered in here. In Sec. III, we apply the `WI2easy` to the problem under study and give the main results obtained from our analysis. In Sec. IV, we give our conclusions and final remarks.

II. WI WITH SM INTERACTIONS

WI is characterized by the evolution equations for the inflaton field ϕ and radiation energy density ρ_r ,

$$\ddot{\phi} + (3H + \Upsilon)\dot{\phi} + V_{,\phi} = 0, \quad (2.1)$$

$$\dot{\rho}_r + 4H\rho_r = \Upsilon\dot{\phi}^2, \quad (2.2)$$

where Υ is the dissipation coefficient, which is assumed to be derived from first principles quantum field theory (for earlier explicit derivations, see for example Refs. [3,10,11,33,34]). For a thermalized radiation bath, it also follows that

$$\rho_r = \frac{g_*\pi^2}{30}T^4, \quad (2.3)$$

where g_* is the radiation bath degrees of freedom. For the temperatures that we obtain here, $T \gtrsim 10^{11}$ GeV, we can assume that all SM particles are thermalized, $g_* = 106.75$, while when also considering the inflaton field thermalized, then $g_* = 106.75 + 1$.

In the WI construction of Ref. [17], it is made use of an axionlike field making the role of the inflaton and which is coupled to the QCD gluon through the dimension five interaction,

$$\mathcal{L}_{\text{int}} = -\frac{\alpha_g}{8\pi f_a} \phi G_{\mu\nu}^c \tilde{G}^{c\mu\nu}, \quad (2.4)$$

where $c = 1, \dots, N_c^2 - 1$, $N_c = 3$, $G_{\mu\nu}^c$ is the gluon field tensor, $\tilde{G}^{c\mu\nu}$ is its dual and $\alpha_g = g^2/(4\pi)$, with g the strong coupling gauge constant and f_a is the axion decay constant.

As first realized in Ref. [12], the dimension five axionlike interaction to $SU(N_c)$ gauge fields can lead to a successful WI realization where the dissipation coefficient Υ is a consequence of sphaleron decay in a thermal bath. In a $SU(N_c)$ gauge field theory, the sphaleron rate is [24]

$$\Gamma_{\text{sphal}} \simeq \kappa(\alpha_g N_c)^5 T^4, \quad (2.5)$$

where the constant κ here is given by

$$\kappa \simeq 0.21 \left(\frac{N_c g^2 T^2}{m_D^2} \right) \frac{N_c^2 - 1}{N_c^2} \left[\ln \left(\frac{m_D}{\gamma} \right) + 3.041 \right], \quad (2.6)$$

where m_D is color Debye mass,

$$m_D^2 = (2N_c^2 + N_f) \frac{g^2 T^2}{6}, \quad (2.7)$$

with N_f the number of quark flavors and γ is the rate of color randomization,

$$\gamma = \frac{g^2 N_c T}{4\pi} \left[\ln \left(\frac{m_D}{\gamma} \right) + 3.041 \right]. \quad (2.8)$$

Explicitly, κ is found to be given by

$$\kappa = 1.26 \frac{N_c^2 - 1}{(2N_c + N_f)N_c} W \left(\frac{e^{3.041}}{N_c} \sqrt{\frac{2\pi(2N_c + N_f)}{3\alpha_g}} \right), \quad (2.9)$$

where $W(x)$ is the Lambert function, given by the principal solution of $x = we^w$. Here, we also consider the running of the gluon gauge coupling as a function of the temperature, $\alpha_g \equiv \alpha_g(T)$, which at one-loop order is [35]

$$\alpha_g^{-1}(T) = \alpha_g^{-1}(M_Z) + \frac{11N_c - 2N_f}{6\pi} \ln \left(\frac{T}{M_Z} \right), \quad (2.10)$$

where $\alpha_g(M_Z) = 0.1172$ with $M_Z = 91.188$ GeV.

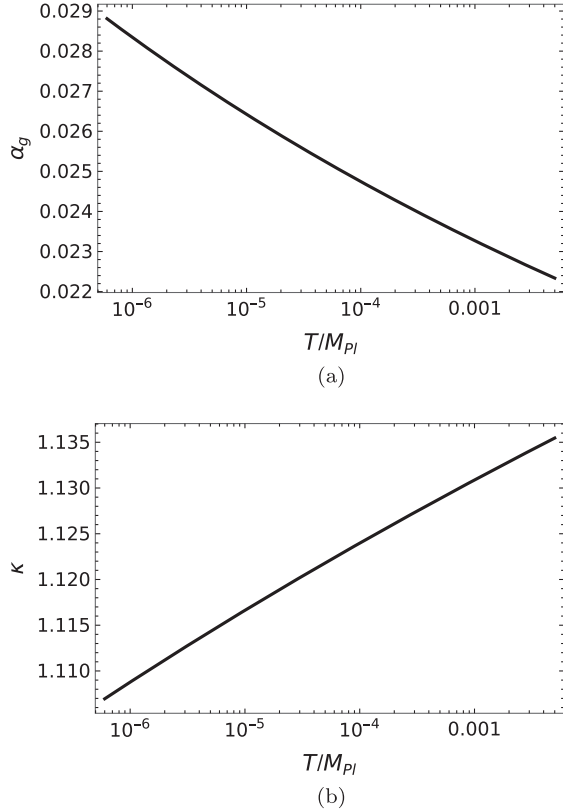


FIG. 1. The temperature dependence of α_g (panel a) from Eq. (2.10) and the coefficient κ (panel b) in the sphaleron rate Eq. (2.5).

As we are going to verify in the next section, in the observational region of interest for the SM WI model, the typical temperatures that we find have values that fall in the range $T \in [6 \times 10^{-7}, 5 \times 10^{-3}] M_{\text{Pl}}$, where M_{Pl} is the reduced Planck mass, $M_{\text{Pl}} \simeq 2.44 \times 10^{18}$ GeV. Using Eqs. (2.9) with (2.10), we show in Fig. 1 the overall temperature dependence of the strong fine structure coupling α_g and also the dimensionless factor κ in the temperature range of interest. In particular, the result for κ is consistent with the value $\kappa \sim 1$ assumed in Ref. [17], but κ and also α_g still display some relevant temperature dependence. Hence, we keep in our study the full $\kappa(T)$ result from Eq. (2.9), considering also Eq. (2.10).

Finally, the sphaleron rate Γ_{sphal} is related to the dissipation coefficient Υ in WI as [12,36]

$$\Upsilon_{\text{sph}} = \frac{\Gamma_{\text{sphal}}}{2Tf_a^2}. \quad (2.11)$$

The result Eq. (2.11) applies in the pure gauge model in the absence of fermions. The presence of fermions (quarks) tends to suppress the sphaleron dissipation in the case of chiral symmetry breaking ($m_q \neq 0$) and can even make it completely vanish for massless fermions [28,37,38]. The effect of the expansion also changes the sphaleron

dissipation expression, in which case the effective dissipation in the SM WI model becomes [17]

$$\Upsilon_{\text{eff}} = \frac{\Upsilon_{\text{sph}}}{1 + \frac{4\tilde{N}_f f_a^2 \Upsilon_{\text{sph}}}{HN_c T^2}}, \quad (2.12)$$

where $\tilde{N}_f = 5$ is the effective number of quarks contributing to Υ_{eff} , as considered and explained in Ref. [17].³ The Hubble parameter in Eq. (2.12) is well approximated by the slow-roll result,

$$H \simeq \sqrt{\frac{V(\phi)}{3M_{\text{Pl}}^2}}. \quad (2.13)$$

The expression Eq. (2.12) turns out to be a complicate function of both temperature and inflaton amplitude. As is known in WI [5,6], the functional dependence of the dissipation coefficient on both T and ϕ can strongly affect the perturbation spectrum and observable quantities derived from it. This effect is traditionally modeled through a function $G(Q)$ of the dissipation ratio $Q = \Upsilon/(3H)$ and that multiplies the scalar of curvature power spectrum (evaluated at Hubble exit) in WI [16],

$$P_{\mathcal{R}} = \left(\frac{H^2}{2\pi\dot{\phi}} \right)^2 \left(1 + 2n_* + \frac{2\sqrt{3}\pi Q}{\sqrt{3 + 4\pi Q} H} T \right) G(Q) \Big|_{k=aH}, \quad (2.14)$$

where n_* in the above expression denotes the possible statistical distribution for the inflaton due to the presence of the radiation bath. The function $G(Q)$ can only be derived numerically, by solving the complete set of perturbation equations in WI. Recent public codes have been dedicated to the precise numerical determination of $G(Q)$ depending on the WI model used [29,30]. In Ref. [17], the authors made use of a fitting form for the power spectrum in WI proposed in the Ref. [41], where

$$P_{\mathcal{R}} \simeq \frac{H^3 T}{4\pi^2 \dot{\phi}^2} F(Q) \Big|_{k=aH}, \quad (2.15)$$

and with the function $F(Q)$ given by

³Note that ultraviolet completions of axionlike models involve beyond SM heavy fields like, for example, in the Kim-Shifman-Vainshtein-Zakharov (KSVZ) axion model [39,40], which introduces heavy exotic quarks and generates the axion coupling Eq. (2.4) through the chiral anomaly. Even though the heavy exotic quarks decouple in the effective theory, they could still affect the dissipation through extra decay channels similar to the cases reported in Refs. [28,37,38]. However, we can always choose appropriate parameters (e.g., the Yukawa couplings in these models) such that the chiral decay rates from the interactions remain much smaller than the Hubble rate, $\Gamma_{\text{ch}} \ll H$ and the result Eq. (2.12) still remains valid.

$$F(Q) = 3.27 \times 10^{-4} Q^7 + 168Q \left[\frac{1}{3} \left(1 + \frac{9Q^2}{25} \right) + \frac{2}{3} \tanh \frac{1}{30Q} \right]. \quad (2.16)$$

Essentially, the authors of Ref. [41] have estimated the result Eq. (2.15) by considering two critical assumptions. First, that the dissipation coefficient was the one coming from the sphaleron decay for pure gauge fields, Eq. (2.11), i.e., $\Upsilon \propto T^3$. As already noticed from Eq. (2.12), the effective dissipation coefficient Υ_{eff} in the SM WI model is expected to have some quite different functional form. Second, they have neglected the statistical distribution term n_* in Eq. (2.14), which means in particular that the inflaton perturbations should remain in the vacuum state throughout the dynamics and never thermalize. Here we will test explicitly the validity of these two hypothesis.

Infer the thermalization of the inflaton perturbations in WI is a difficult problem. This is because we need to know precisely the interactions of the inflaton field with all other degrees of freedom to properly estimate the scattering rate of the inflaton field. Then, we need to solve the kinetic equation for the inflaton particle distribution in the expanding spacetime to check whether the statistical distribution for the inflaton field would indeed approach a thermal one. In the Ref. [42] it was shown that in WI, even for scattering rates much smaller than the Hubble rate, the inflaton distribution can acquire a shape close to a thermal one. Then when $\Gamma_{\text{scat}} \gtrsim H$, the distribution approaches a thermal equilibrium form, given by a Bose-Einstein distribution, $n_* \rightarrow n_{\text{BE}}$, and as indeed expected in general. Here, we can take advantage of the existent literature on the thermodynamics and kinetic theory of axions in QCD to then infer the possibility of thermalization of the inflaton field. Despite the fact that there is still a debate in the literature about the thermalization of axions in QCD [43,44], the temperatures we are interest in are sufficiently high that we can make use of the hard thermal loop estimates of Ref. [45]. The dominant process for axion thermalization comes from the scattering processes with the gluon, $\phi + g \rightleftharpoons g + g$, while the processes involving quarks are subdominant. Summing all processes $\Gamma_{\text{scat}} = \Gamma_{\phi+g \rightleftharpoons g+g} + \Gamma_{\phi+q \rightleftharpoons g+q} + \Gamma_{\phi+\bar{q} \rightleftharpoons g+\bar{q}} + \Gamma_{\phi+g \rightleftharpoons q+\bar{q}}$, the result is [45]

$$\Gamma_{\text{scat}} = \frac{\alpha_g^3 T^3}{f_a^2} \left\{ \frac{15\zeta(3)}{2\pi^4} \left[3 - \gamma_E + \frac{2\zeta'(3)}{\zeta(3)} - \ln(2\pi\alpha_g) - \frac{17}{12} \right] + \frac{N_f \zeta(3)}{6\pi^4} + \frac{N_f \zeta(3)}{2\pi^4} \left[\frac{1}{4} \left(9 + \ln 4 - 6\gamma_E + \frac{6\zeta'(3)}{\zeta(3)} \right) + \frac{3}{4} \left(3 - 2\gamma_E + \frac{2\zeta'(3)}{\zeta(3)} \right) - \frac{3}{2} \ln(2\pi\alpha_g) - \frac{15}{8} \right] \right\}, \quad (2.17)$$

where $\zeta(x)$ is the zeta-function and γ_E is the Euler-Mascheroni constant, $\gamma_E \simeq 0.577$.

The result given by Eq. (2.17) will be directly compared with the Hubble parameter in the next section. Finally, the suitability of the approximated power spectrum, Eq. (2.15), is compared to the one obtained from the full numerical solution of perturbations in WI using `WI2easy`.

III. BACKGROUND DYNAMICS AND PERTURBATION QUANTITIES IN THE SM WI MODEL

In our analysis of the SM WI model we will make use of the recently released code `WI2easy` for WI analysis of background and perturbation dynamics. The code can handle general forms of dissipation coefficients, $\Upsilon = C_\Upsilon f(T, \phi)$, which makes it well suited to treat to effective dissipation coefficient Eq. (2.12). Here, we identify the dimensionless constant C_Υ with the axion constant through

$$C_\Upsilon = \frac{M_{\text{Pl}}^2}{f_a^2}. \quad (3.1)$$

`WI2easy` automatically determines the numerical value for C_Υ once an initial value for the dissipation ratio Q is given, while also performing the normalization of the inflaton potential such that at the pivot scale the power spectrum has the amplitude according to the CMB normalization. Note that the dissipation ratio is here defined as

$$Q = \frac{\Upsilon_{\text{eff}}}{3H}. \quad (3.2)$$

The pivot scale has the value $k_p = 0.05 \text{ Mpc}^{-1}$, while the CMB amplitude normalization is $A_s \simeq 2.105 \times 10^{-9}$ (we follow Planck's default values [46]). For the inflaton potential, we consider the same used in Ref. [17], such that our results can be directly compared to,

$$V(\phi) = \lambda \phi^4 \equiv \frac{V_0}{M_{\text{Pl}}^4} \phi^4, \quad (3.3)$$

where the potential normalization V_0 is automatically computed in `WI2easy` such that the spectrum will have the proper normalization A_s at Hubble exit.

Under the assumptions used by the authors of Ref. [41] to obtain Eq. (2.15), the equivalent in the `WI2easy` options available in its built-in interface is to choose the option of including explicitly a radiation noise term in the equation for the radiation perturbation (see the discussion in Ref. [8] where this term was first identified and also the `WI2easy` paper [29] for additional discussions in connection to the code results). As already mentioned, Ref. [41] has also implicitly assumed $n_* = 0$ in Eq. (2.14). In fact they have dropped the entire $1 + 2n_*$ term in Eq. (2.14), which makes their expression for the power spectrum suitable only for $Q \gtrsim 1$, i.e., in the strong dissipative regime of WI. Hence, comparing Eq. (2.14) with (2.15), we can identify their

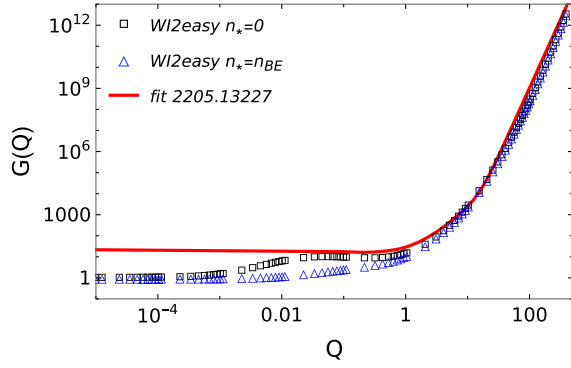


FIG. 2. The $G(Q)$ function from the fitting of Ref. [41] compared with the results from WI2easy in the absence and presence of inflaton thermalization.

approximated $G(Q)$ function as being

$$G_{\text{approx}}(Q) = \frac{\sqrt{3 + 4\pi Q}}{2\sqrt{3}\pi Q} F(Q). \quad (3.4)$$

We compare Eq. (3.4) with the results provided by WI2easy.

By fully implementing the SM WI model in WI2easy, we first make the comparison of Eq. (3.4) with our numerical result. This is shown in Fig. 2. Among the options available in WI2easy, we consider the cases of nonthermalized ($n_* = 0$) and fully thermalized ($n_* = n_{\text{BE}}$) inflaton perturbations. One notices that Eq. (3.4) only roughly fits the numerical results, for both nonthermalized and thermalized inflaton perturbations, for $Q \gtrsim 1$, while in the weak dissipative regime, $Q \lesssim 1$, it displays noticeable discrepancies, in particular in the case of thermalized inflaton perturbations.

To check for thermalization of the inflaton perturbations, we run WI2easy for a range of values of Q at Hubble exit time. The temperature is obtained at the different values of Q and used to check the ratio of the scattering rate Eq. (2.17) with the Hubble parameter, Γ_{scat}/H , evaluated at the Hubble exit point $k = aH$. The results are shown in Fig. 3. The result of Fig. 3 indicates that the inflaton

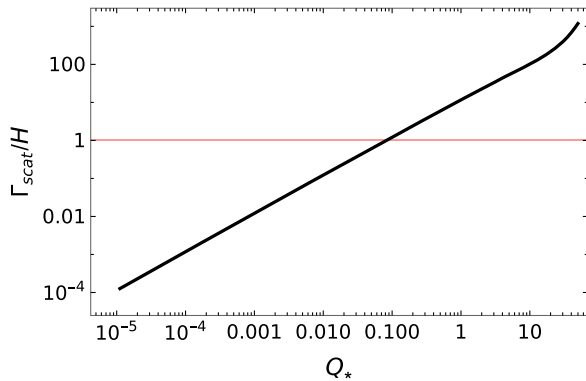


FIG. 3. The ratio of the scattering rate for the inflaton field, Eq. (2.17), with the Hubble parameter as a function of Q at Hubble exit.

perturbations are expected to thermalize for $Q_* \gtrsim 0.08$, hence, already in the weak dissipative regime of WI. This result, together with the ones shown in Fig. 2, indicates that the power spectrum can have significant differences in the

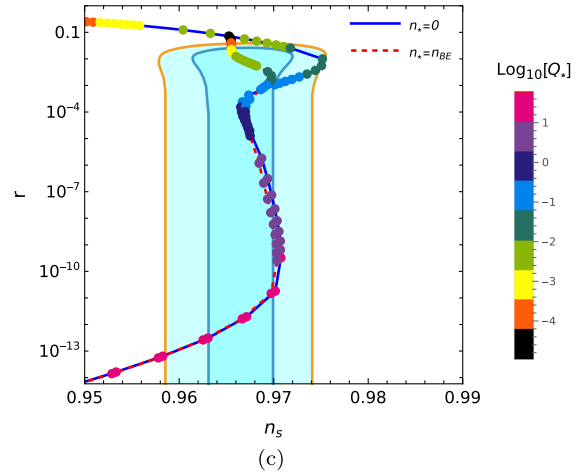
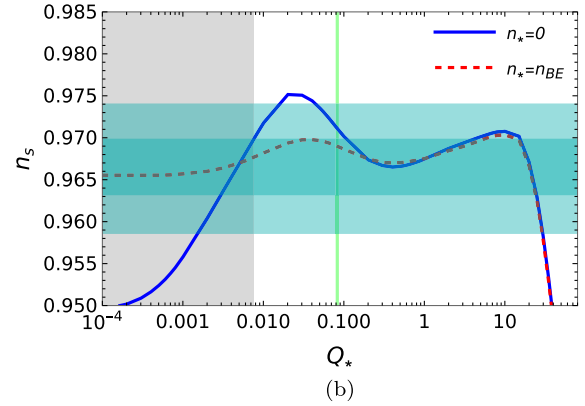
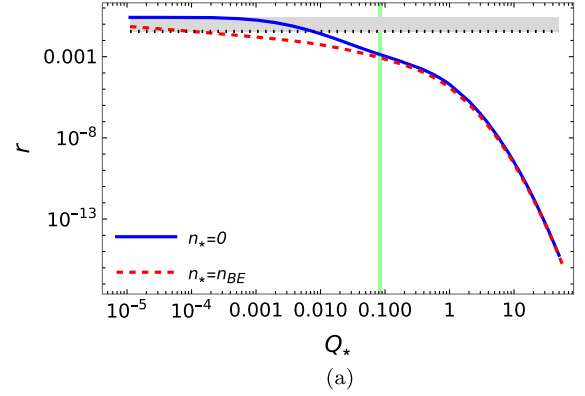


FIG. 4. The tensor-to-scalar ratio (panel (a)) and the spectral tilt (panel (b)) at different values of Q_* . The tensor-to-scalar as a function of n_s is presented in panel c. The top gray region in panel (a) and the left gray region in panel (b) indicate the region where $r > 0.036$, with the upper bound on r from the BICEP, Keck Array and Planck combined data [47]. The cyan regions in panels (b) and (c) indicate the one- and two-sigma constraints also from Ref. [47]. The vertical green line in panels a and b indicates the expected threshold for thermalization ($n_* = n_{\text{BE}}$).

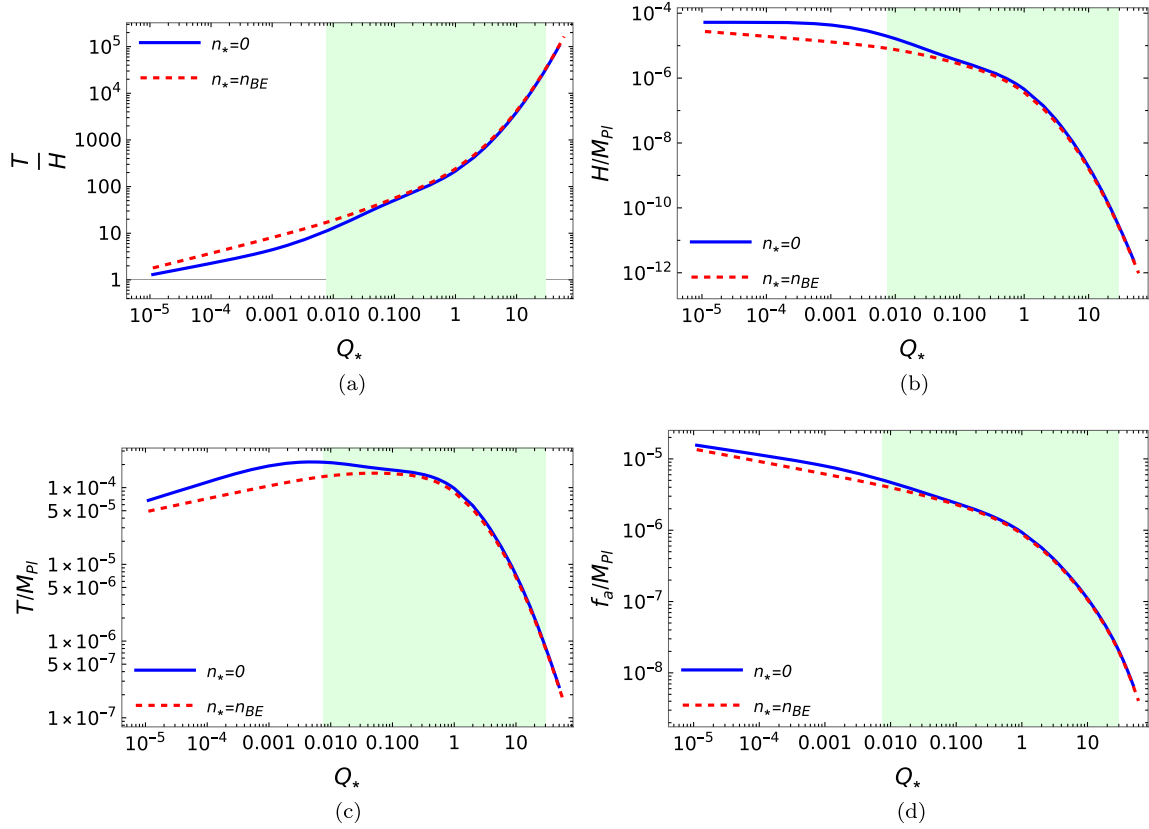


FIG. 5. Some of the key background quantities in WI, namely T/H (panel a), H (panel b) and T (panel c) and the axion decay constant f_a (panel d). The light green band in all the plots indicates the region of Q_* values for which r and n_s are within the observational window (see Fig. 4).

thermal case ($n_* = n_{BE}$) than in the nonthermalized situation, in special in the weak dissipative regime of WI.

Still using `wi2easy`, we obtain the dependence of the tensor-to-scalar ratio, r , and the spectral tilt, n_s , as a function of Q at Hubble exit. We again consider both cases of nonthermalized and thermalized inflaton perturbations. The corresponding results are presented in Fig. 4. One notices from Fig. 4(c) that our results agree with the ones shown in the Fig. 1 of Ref. [17] in the region around $Q_* \in [0.2, 15]$. But they deviate significantly for $Q \lesssim 0.2$ and $Q_* \gtrsim 15$. The difference in the results for the weak dissipative regime can be attributed due to the unsuitability of the fitting function for the power spectrum considered in Ref. [17] for that region of Q values, while the difference at larger values of Q_* can be attributed to the fact of the inflaton field amplitude dependence in the dissipation coefficient Eq. (2.12) becoming more prominent in the large dissipative regime of WI. This change of behavior in the spectrum at large Q_* turns out to happen here earlier than the estimate of $Q_* > 40$ quoted in Ref. [17]. In fact, as Q increases, the leading temperature dependence power in Υ_{eff} shifts to values lower than three, $\Upsilon_{\text{eff}} \propto T^c$, with $c < 3$. The power spectrum in this case is expected to have less power than in the case of $c = 3$ [5,6], forcing n_s to become redder.

The range of values for Q_* which falls inside the observational window are $Q_* \in [0.0076, 30]$. For $Q_* < 0.0076$ the tensor-to-scalar ratio becomes larger than the upper bound set by the BICEP, Keck Array and Planck combined data [47], $r < 0.036$, while for $Q_* \gtrsim 30$ the spectral tilt goes outside of the two-sigma region for the same combined data.

As a note, we point out that the recent released data results from the Atacama Cosmology Telescope (ACT) [48] push the results for n_s to slightly larger values compared to the earlier results from the Planck/BICEP and Keck ones. While in the WI models studied in Ref. [49] those update results tend to favor slight higher values of Q , in the present model they go in the opposite direction, as seen here in the Fig. 4(c).

Finally, in Fig. 5, we present the results for some of the background quantities of interest as a function of Q_* , namely, the temperature over the Hubble parameter, T/H , the Hubble parameter, the temperature and the axion decay constant f_a . The value of f_a is obtained from the coefficient of the effective dissipation coefficient and defined through Eq. (3.1). One notices in particular that in the observational window of valid r and n_s values, we have $f_a \in [4.98 \times 10^{10}, 1.23 \times 10^{13}]$ GeV, which is still consistent with present bounds on the QCD axion for instance [50,51].

IV. CONCLUSION

In this paper, we revisit the proposal of Ref. [17] to construct a WI scenario from SM gauge interactions. The embedding of WI within the SM framework represents a significant advancement, as it establishes, for the first time, a concrete connection between WI dynamics and the particle content of the SM. We reanalyze the results of Ref. [17] using the `WI2easy` code, a numerical tool designed for precision computation of background evolution and perturbations in WI. Furthermore, we critically reassess key assumptions in the original work. Specifically, we scrutinize the validity of the approximate power spectrum employed in that reference by comparing it to results obtained with `WI2easy`, incorporating the full temperature and inflaton amplitude dependencies of the dissipation coefficient. Additionally, we evaluate whether inflaton perturbations could undergo thermalization—a crucial consideration for ensuring the accuracy of WI power spectrum predictions and the reliability of derived observational parameters. These refinements address pivotal challenges in achieving robust, first-principles calculations in WI scenarios.

Our results indicate that the inflaton can achieve a thermal equilibrium distribution for values of the dissipation ratio at Hubble exit of $Q_* \gtrsim 0.08$, hence already in the weak dissipative regime of WI. While our more accurate results obtained with `WI2easy` show agreement with the results from Ref. [17] for values of Q_* between approximately 0.6 and 15, there are significant differences in the $Q_* \lesssim 0.6$ and $Q_* \gtrsim 15$ regions. We have identified the source of these differences, which can be attributed to the approximations considered in Ref. [17].

For future work, it would be of interest to also analyze possible non-Gaussianity signatures that this realization of WI might have, following for example the methods set forward in Ref. [52]. This would need a more dedicated analysis using the form of the dissipation coefficient given by Eq. (2.12), which differs from the most simpler form where $\Upsilon \propto T^c$. If one assumes the case with $c = 3$ analyzed

in Ref. [52], for the range of Q_* values of interest that we have obtained here, $Q_* \in [0.0076, 30]$, the estimated non-Gaussianity coefficient for the characteristic shape of WI obtained in Ref. [52] is $5 \lesssim |f_{\text{NL}}^{\text{warm}}| \lesssim 30$ (for a more recent analysis for this case, see also Ref. [41]). We expect that Eq. (2.12), which has an effective power in the temperature $c < 3$, to produce a smaller $|f_{\text{NL}}^{\text{warm}}|$. Observational constraints based on the warm shape of the bispectrum in WI are, unfortunately, very poorly known at the moment, with results derived only for the case of the bispectrum obtained in the very simple case of a constant dissipation coefficient (i.e., with no temperature dependence on Υ), with the result for $f_{\text{NL}}^{\text{warm}}$ estimated to be [53] $|f_{\text{NL}}^{\text{warm}}| = 48 \pm 27$.

In conclusion, our findings establish the viability of constructing successful WI scenarios directly from SM gauge interactions. Furthermore, they highlight the critical role of the `WI2easy` numerical toolkit in achieving reliable results: while approximate analytical methods remain valuable for conceptual studies, our work demonstrates that bypassing such approximations is essential for deriving observationally robust predictions. This precision is particularly crucial when computing observable quantities—such as the primordial power spectrum or spectral index—that must align with high-accuracy cosmological datasets.

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DATA AVAILABILITY

The data that support the findings of this article are openly available [54].

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