Towards a reliable effective field theory of inflation

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We present the first renormalizable quantum field theory model for inflation with a super-Hubble inflaton mass and sub-Planckian field excursions, which is thus technically natural and consistent with a high-energy completion within a theory of quantum gravity. This is done in the framework of warm inflation, where we show, for the first time, that strong dissipation can fully sustain a slow-roll trajectory with slow-roll parameters larger than unity in a way that is both theoretically and observationally consistent. The inflaton field corresponds to the relative phase between two complex scalar fields that collectively break a $U(1)$ gauge symmetry, and dissipates its energy into scalar degrees of freedom in the warm cosmic heat bath. A discrete interchange symmetry protects the inflaton mass from large thermal corrections. We further show that the dissipation coefficient decreases with temperature in certain parametric regimes, which prevents a large growth of thermal inflaton fluctuations. We find, in particular, a very good agreement with the Planck legacy data for a simple quadratic inflaton potential, predicting a low tensor-to-scalar ratio $r \lesssim 10^{-5}$.

The observational success of the inflationary paradigm is undeniable. Not only does it explain the spatial flatness, homogeneity and isotropy of our Universe on large scales, but it also generates the nearly-Gaussian and nearly-scale invariant spectrum of primordial density fluctuations required by measurements of the Cosmic Microwave Background (CMB) anisotropies and Large-Scale Structure.

From the theoretical perspective, however, inflation is far from being on solid grounds. Canonical models are based on an effective quantum field theory for a scalar field, $\phi$, that mimics a cosmological constant while on a slow-roll trajectory sustained by the Hubble friction. This imposes the well-known slow-roll conditions on the form of the scalar potential, $V(\phi)$,

$$\epsilon_\phi = M_P^2 \left( \frac{V'}{V} \right)^2 \ll 1, \quad |\eta_\phi| = M_P^2 \left| \frac{V''}{V} \right| \ll 1, \tag{1}$$

where $M_P \equiv (8\pi G)^{-1/2} \approx 2.4 \times 10^{18}$ GeV is the reduced Planck mass. The $\eta_\phi$-condition, in particular, implies a very light inflaton field, with effective mass well bellow the Hubble scale, $m_\phi \ll H$.

Light scalars are, however, extremely unnatural in any effective quantum field theory, since quantum corrections to their mass are quadratically divergent - a symptom of their sensitivity to new physics above the cut-off energy scale below which the theory can describe physical phenomena. The Higgs boson is the paramount example of this technical naturalness problem, better known as electroweak hierarchy problem, as its mass should naturally lie close to the Planck scale if no new particle states exist below the latter (see, e.g., Refs. 3, 4). According to ’t Hooft, this is related to no new symmetries emerging for vanishing scalar masses, as opposed to, e.g., fermions, for which a chiral symmetry is gained in this limit.

An effective field theory of inflation based on general relativity or any classical gravity theory necessarily fails above the Planck scale, making the scalar potential sensitive to Planck-suppressed non-renormalizable operators that generically drive the inflaton mass towards values above $H$. This is the well-known “eta-problem” in supergravity and string theory (see, e.g., Ref. 6), which are but examples of a more general inflationary naturalness problem, given that the cut-off scale for the effective field theory of inflation must necessarily lie above the Hubble scale, thus driving the inflaton mass to super-Hubble values in tension with the slow-roll conditions.

Many have tried to overcome this issue by employing symmetries that could enforce cancellations between different quantum corrections to the inflaton’s mass. The best-known example is supersymmetry, where bosonic and fermionic quantum corrections cancel out, but supersymmetry is broken by the finite energy density during inflation, always leaving $\mathcal{O}(H)$ corrections to the inflaton mass except, e.g., for some fine-tuned choices of the Kähler potential. Global symmetries are also expected to be broken by quantum gravity effects, such that shift symmetries also cannot guarantee the flatness of the potential. This is inherently assumed in models with axion-like fields, as well as models with asymptotic plateaux in the scalar potential, like in Higgs inflation and attractor models, hence necessarily involving some degree of fine-tuning, at least without a deeper understanding of the fundamental high-energy theory. Fine-tuning is, of course, undesirable in a theory that is supposed to dynamically generate the otherwise extremely fine-tuned initial conditions of the standard cosmological model.

There has also been an increased interest in the conditions leading to effective field theories admitting a consistent ultraviolet completion in quantum gravity and derived by explicit string theory constructions, known
as the “swampland” conjectures \cite{11,13}. In particular, these conjectures require, e.g., that \(M_P|V'|/V \gtrsim \mathcal{O}(1)\), alongside the criterion \(\Delta \phi \lesssim M_P\). While some inflationary scenarios can satisfy the second criterion, the first is in tension with the first slow-roll condition.

In addition, inflationary predictions for CMB observables rely on the consistency of quantum field theory in curved space-time, which is also not free from ambiguities, and for which we have no empirical guidance (see, e.g., Ref. \cite{14,19}). In particular, the form of the primordial spectrum of density perturbations on super-horizon scales depends on the particular choice of the Bunch-Davies vacuum when these scales are deep inside the horizon at the start of inflation. This choice is not unique and is, moreover, sensitive to transplanckian physics.

It had been understood early on in the development of warm inflation that, in the strongly dissipative regime, all these crucial issues can be overcome \cite{10,11} (see also \cite{22,23}), but realizing this in a quantum field theory model has proven very challenging. This Letter develops the first such model, which is renormalizable and both theoretically and observationally consistent.

Dissipation is an inherent physical process to any system interacting with its environment. While conventional inflationary models assume that matter and radiation are diluted away by accelerated expansion, being only regenerated during reheating after inflation, this need not be the case in general. In fact, if at the onset of inflation the universe is filled with a radiation bath, interactions between the inflaton and particles in the bath necessarily lead to dissipative effects and associated particle production. Part of the inflaton’s energy is thus continuously transferred to the radiation bath, acting as a heat source and thus keeping it warm despite the supercooling effect of accelerated expansion. This is the main premise of the alternative warm inflation paradigm, which was proposed more than two decades ago \cite{24,25}.

Dissipative effects are encoded in the effective equation for the inflaton field,

\[
\ddot{\phi} + (3H + \Upsilon)\dot{\phi} + V'(\phi) = 0, \tag{2}
\]

alongside the sourced equation for the radiation energy density, \(\rho_R\),

\[
\dot{\rho}_R + 4H\rho_R = \Upsilon\dot{\phi}^2, \tag{3}
\]

which can be derived from energy-momentum conservation or from an explicit computation of particle production rates. The dissipation coefficient \(\Upsilon = \Upsilon(\phi, T)\) can be computed from first principles using standard thermal field theory techniques, at least close to thermal equilibrium and in the regime \(T \gtrsim H\) where space-time curvature effects can be neglected (see Ref. \cite{27} for a review).

In the slow-roll regime we then have \(\dot{\phi} \approx V'/3H(1+Q)\), where \(Q = T/3H\) is the dissipative ratio, and \(\rho_R/V \approx \epsilon_\phi Q/\sqrt{2(1+Q)^2}\), provided that the slow-roll conditions,

\[
\epsilon_\phi, |\eta_\phi| < 1 + Q \tag{4}
\]

are satisfied. Hence, if strong dissipation can be achieved throughout inflation, \(Q \gg 1\), the eta-problem is avoided and inflation can occur for inflaton masses above the Hubble scale \cite{20,21,28}. It is also manifest that radiation is a sub-dominant component in the slow-roll regime, while it may come to dominate if \(Q\) becomes large at the end of inflation. In warm inflation there can thus be a smooth transition between inflation and the radiation era with no need for a separate reheating period \cite{24} where the inflaton decays away, allowing e.g. for cosmic magnetic field generation \cite{29}, baryogenesis \cite{30,31} and inflaton dark matter \cite{32} or dark energy \cite{33,34}.

Dissipative effects also lead to thermal inflaton fluctuations, as a result of the fluctuation-dissipation theorem \cite{20,35}. Inflaton perturbations are, in particular, sourced by a Gaussian white noise term with variance \(\langle \xi_k(\epsilon_k) \rangle = 2(\Upsilon + H)/\alpha^3 \times (2\pi)^3 \delta(\mathbf{k} + \mathbf{k}')\), valid when \(T \gg H\). Thus, thermal fluctuations are generically larger than their quantum counterparts, such that primordial density fluctuations are generated by classical field perturbations. Moreover, in the strong dissipation regime, when \(Q > 1\), their amplitude freezes out before they become superhorizon \cite{20}, nevertheless producing a nearly scale-invariant spectrum of classical curvature perturbations \cite{36,37,38}. Finally, such an enhancement of the primordial scalar perturbations typically leads to a lower inflationary energy scale and, therefore, to a lower tensor-to-scalar ratio, since gravitational wave production is unaffected by thermal effects below the Planck scale, as observed, e.g., in Refs. \cite{28,49,50}.

Despite all these appealing features, consistently realizing warm inflation in quantum field theory models has proved to be an enormous challenge \cite{12,11}. The inflaton typically gives a large mass to the particles it interacts with, similarly to the Higgs mechanism, and it is extremely hard to sustain sufficiently strong dissipative effects if their mass exceeds the ambient temperature due to Boltzmann suppression. Even if this can be avoided, thermal backreaction generically reintroduces the eta-problem through thermal inflaton mass corrections \(\Delta m_\phi \sim T \gtrsim H\) up to dimensionless couplings \cite{43}.

These problems were recently overcome by employing symmetry arguments, in a model akin to “Little Higgs” models for electroweak symmetry breaking and dubbed the “Warm Little Inflaton” (WLI) model \cite{50}. In this model, the inflaton corresponds to the relative phase between two complex scalar fields, \(\phi_1\) and \(\phi_2\), equally charged under a \(U(1)\) gauge symmetry, and which spontaneously break the latter. In the unitary gauge, the resulting vacuum manifold can be parametrized as

\[
\langle \phi_1 \rangle = \frac{M}{\sqrt{2}} e^{i\phi/M}, \quad \langle \phi_2 \rangle = \frac{M}{\sqrt{2}} e^{-i\phi/M}, \tag{5}
\]

where \(M\) is the symmetry breaking scale and \(\phi\) is the gauge-invariant inflaton field. In the original model, these complex scalar fields were coupled to fermion fields
\( \psi_{1,2} \) through Yukawa interactions satisfying a discrete interchange symmetry \( \phi_1 \leftrightarrow i \phi_2, \psi_1 \leftrightarrow \psi_2 \). This symmetry then ensures that if, e.g., the \( \psi_1 \) fermion couples to the linear combination \( \phi_1 + \phi_2 \), then \( \psi_2 \) couples to \( \phi_1 - \phi_2 \). As a result, the fermions acquire masses which are trigonometric functions of the inflaton field,

\[
m_1 = gM \cos(\phi/M), \quad m_2 = gM \sin(\phi/M),
\]

where \( g \) denotes the Yukawa coupling. These masses are thus bounded even if \( \phi \gg M \), and can be below the temperature during inflation. Moreover, the leading contributions in the finite-temperature effective potential for \( m_{1,2} \lesssim T \),

\[
\Delta V_T = -\frac{7\pi^2}{180} T^4 + \frac{1}{12}(m_1^2 + m_2^2) T^2 + \ldots ,
\]

are independent of the inflaton field, thus eliminating the troublesome thermal corrections to the inflaton’s mass.

This model thus yields a consistent realization of warm inflation, and we have moreover shown that its observational predictions are in agreement with the Planck data for a quartic inflaton potential, \( V(\phi) = \lambda \phi^4 \) [50–54]. However, this agreement requires weak dissipation, \( Q_1 \lesssim 1 \), at the time the relevant CMB scales become super-horizon, about 60 e-folds before the end of inflation. Even though \( Q_1 \) becomes large towards the end of the slow-roll regime, the eta-problem remains in this case.

This is an inherent consequence of the form of the dissipation coefficient in this scenario, \( \Upsilon \propto T \). As originally shown in Ref. [37], and further analyzed numerically in Ref. [65], the temperature dependence of the dissipation coefficient necessarily leads to a coupling between inflaton fluctuations and perturbations in the radiation fluid that modifies the evolution of the former. In particular, for \( dT/dT > 0 \), this results in a substantial enhancement of inflaton fluctuations if \( Q_1 \gtrsim 1 \). Physically, this is a consequence of dissipation increasing the temperature more in regions where it is already than average. If \( Q_1 \) grows during inflation, as for the quartic potential, the primordial perturbation spectrum then becomes blue-tilted, which is ruled out by Planck [2].

If, however, \( dT/dT < 0 \), inflaton perturbations are damped. More concretely, the dimensionless curvature power spectrum in the strong dissipation regime, \( Q_1 \gg 1 \), is generically well approximated by:

\[
\Delta^2_R \approx \frac{\sqrt{3} \pi}{24 \pi^2} e^{x_1} \frac{V(\phi_1)}{M^4} \left( \frac{T_v}{H_*} \right) Q_1^{5/2} Q_2^{3/2} e^{\beta_1},
\]

where the last factor results from the interplay between the inflaton and radiation fluctuations, with \( 1 \lesssim Q_2 \lesssim 10 \) and \( \beta_1 \) depending on \( c = d \log T / d \log T \) at horizon-crossing. For instance, numerically solving the coupled system of perturbations as detailed in, e.g., Ref. [65], we find \( \beta_1 \approx 2.5 \) for \( \Upsilon \propto T \), while for \( \Upsilon \propto T^{-1} \) yields \( \beta_1 \approx -1.6 \). Thus, it is tantamount to find a model with a dissipation coefficient decreasing with temperature, such as too allow for strong dissipation in an observationally consistent way.

Here, the complex scalar fields \( \phi_1 \) and \( \phi_2 \) are coupled to two other complex scalars \( \chi_1 \) and \( \chi_2 \) in the thermal bath. The interactions have a renormalizable bi-quadratic form and, as in the original WLI model, satisfy the discrete interchange symmetry \( \phi_1 \leftrightarrow i \phi_2, \chi_1 \leftrightarrow \chi_2 \). Without loss of generality, we may write the relevant interaction Lagrangian in the form:

\[
\mathcal{L}_{\phi \chi} = \frac{1}{2} g^2 |\phi_1 + \phi_2|^2 |\chi_1|^2 + \frac{1}{2} g^2 |\phi_1 - \phi_2|^2 |\chi_2|^2,
\]

which is a straightforward generalization of the fermionic WLI model. With the vacuum parametrization in Eq. (5), the zero-temperature masses of the \( \chi_1 \) and \( \chi_2 \) fields are also given by Eq. (6), being bounded functions of the inflaton field \( \phi \). Most importantly, their leading contributions to the inflaton thermal mass also cancel out as in the fermionic case.

The main difference between coupling fermions or scalars to the inflaton field lies in the form of the dissipation coefficient, due to their different statistics at non-zero temperature, which will be fundamental to set the present model apart from previous model building realizations of warm inflation. For on-shell particle production, which is the dominant process for \( T \gg m_{1,2} \), the dissipation coefficient is given by

\[
\Upsilon = \sum_{i,j=1,2} g^4 \sin^2(2\phi/M) \int \frac{d^3p}{(2\pi)^3} \frac{n_B(1+n_B)}{T} i\omega_{p,i}^2,
\]

where \( n_B(\omega_{p,i}) \) is the Bose-Einstein distribution, \( \omega_{p,i}^2 = p^2 + m_i^2, m_i^2 = m_i^2 + \alpha^2 T^2 \) is the thermally corrected mass of the \( \chi_1, \chi_2 \) fields and \( \Gamma_i \) their thermal decay width. These depend on interactions within the thermal bath, which we model as Yukawa interactions with light fermions \( \psi_{L,R} \) (with appropriate charges) and scalar self-interactions,

\[
\mathcal{L}_{\chi^3} = \sum_{i,j=1,2} \left( h^2 \chi_i \psi_L \psi_R + h.c. + \frac{\lambda}{2} |\chi_i|^4 + \lambda' |\chi_i|^2 |\chi_j|^2 \right).
\]

Scalar self-interactions contribute only at two-loop order to the thermal decay width, and we focus on parametric regimes where the decay into light fermions is dominant. Both types of interactions contribute nevertheless to the thermal mass at the same order, yielding \( \alpha^2 \approx [h^2 + \lambda(N+1) + \lambda'N]/12 \) if the \( \chi_i \) fields are in an \( N \)-dimensional representation of some gauge group. The resulting dissipation coefficient is then given by:

\[
\Upsilon \approx 4N g^2 M^2 T^2 \left[ 1 + \frac{1}{\sqrt{2\pi}} \left( \frac{m_\chi}{T} \right)^{3/2} \right] e^{-m_\chi/T},
\]

where we have taken the average of the oscillatory terms for field excursions \( \Delta \phi \gg M \), yielding an average mass \( m_\chi^2 \approx g^2 M^2 / 2 + \alpha^2 T^2 \) for both \( \chi_i \) scalar fields.
Although the dissipation coefficient has, in general, a non-trivial temperature dependence, when $m_\chi$ is dominated by thermal effects we have $m_\chi \simeq \alpha T$ and $\Upsilon \propto T^{-1}$. This will then yield the required damping of inflaton fluctuations as discussed above.

Let us consider the simplest scenario with a quadratic inflaton potential, \( V(\phi) = \frac{1}{2} m_\phi^2 \phi^2 \), since as we discussed above there are no symmetries protecting the inflaton from acquiring at least a Hubble-scale mass. The slow-roll equations can be integrated analytically for this potential when $\Upsilon \propto T^{-1}$ and from the form of the primordial perturbation spectrum at strong dissipation in Eq. (8) we obtain for the scalar spectral index $n_s$ a value of $0.965 \pm 0.001$, since as we discussed above there are no symmetries protecting the inflaton from acquiring at least a Hubble-scale mass.

The agreement with observational data for $Q_\ast \sim 100$ is, however, our most significant result, in this regime the slow-roll parameters $\epsilon_\phi = \eta_\phi \gtrsim 1$ throughout inflation. The slow-roll trajectory is thus fully sustained by the dissipative friction, yielding a consistent effective field theory for inflation with $m_\phi \gtrsim H$, $\Delta \phi \lesssim M_\phi$ and $M_\phi \phi / V \gtrsim 1$. We illustrate the dynamical evolution of the inflaton-radiation system in a representative case with $Q_\ast = 100$ in Fig. 2 where it is manifest that inflation can consistently occur with a super-Hubble inflaton mass and a sub-Planckian field excursion.

In this strong dissipative regime, primordial non-Gaussianity should generically be at the level $f_{NL}^{\text{warm}} \lesssim 10$ [57], and the dedicated searches by the Planck collaboration for the warm shape of the bispectrum [58, 59] allow for $Q_\ast < (3.2 - 4) \times 10^3$ (95% C.L.). Our scenario thus lies comfortably within these limits. We have, moreover, made a preliminary analysis of non-Gaussianity for the particular form of the dissipation coefficient in the present scenario, using the numerical codes developed in [60], obtaining $f_{NL}^{\text{warm}} \approx 3$ for $Q_\ast = 100$.

Scenarios with $Q_\ast \gg 100$, for the present model, typically require larger values for the coupling $g$ and hence larger zero-temperature masses for the $\chi_i$ fields. This results in larger deviations from $\Upsilon \propto T^{-1}$ and thus to a less efficient damping of inflaton fluctuations and a more blue-tilted spectrum.

In the example shown in Fig. 2, it is also manifest that strong dissipation can be sustained for a whole of 60 e-
folds of inflation, maintaining a slowly decreasing temperature $T \gg H$ \cite{61}. The temperature always satisfies $T \geq g M$, ensuring that the $\chi_1$ fields are relativistic, yet keeping the $U(1)$ gauge symmetry broken throughout inflation \cite{62}. The radiation abundance is $\rho_R/V \simeq 0.5 \epsilon_\phi/Q$ in the slow-roll regime for $Q \gg 1$, leading to a smooth transition to the radiation-dominated era.

Hence, this Letter shows, for the first time, that slow-roll inflation does not require an unnaturally light inflaton scalar field, within a renormalizable quantum field theory, with a simple quadratic scalar potential, that is robust against corrections from unknown new physics, particularly Planck-suppressed non-renormalizable operators. Moreover, the spectrum of primordial density fluctuations is fully described by classical thermal fluctuations is fully described by classical thermal fluctuations.

It had been pointed out early in the development of warm inflation \cite{20, 21} that in the strong dissipation regime it is possible for $\phi > H$ to solve the eta-problem, with sub-Planckian field excursions, thus solving the swampland criteria well before they were stated. It had also been understood that increasing dissipation lowers the energy scale of inflation and hence the tensor-to-scalar ratio \cite{20, 28} in line with subsequent CMB observations. However, the challenge has been to find a theoretically and observationally consistent model displaying all these appealing features. This Letter has achieved this by obtaining a quantum field theory model of inflation that is reliable in this respect, involving only a few fields, being free of the fine-tuning and ambiguities that generically plague the more conventional cold scenario.

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Dissipative effects mediated by heavy virtual particles may be sufficient to sustain a thermal bath during inflation, but only at the expense of very large field multiplicities in supersymmetric models [33–37, which may only be found in specific constructions in string theory [38] or possibly other scenarios with extra-dimensions [39].

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The exact critical temperature of the phase transition depends on the self-couplings of the $\phi_1$ and $\phi_2$ scalar fields which we have not specified, but we can assume $T_c \sim M$ within order unity factors.