A general overview is given of the warm inflation scenario.

Talk presented at PASCOS-98, Northeastern University, March 1998

The thermodynamics of inflationary expansion can be either isentropic or non-isentropic. Representing the early universe by a two fluid mixture of radiation energy density $\rho_r$ and vacuum energy density $\rho_v$, the inflationary regime, when the scale factor accelerates $\ddot{R} > 0$, is for $\rho_v > \rho_r$. From the point of view of two fluid Friedmann cosmology, isentropic inflationary expansion appears as a limiting case within the general regime of non-isentropic inflation.

In particle physics the vacuum equation of state $\rho_v = -p_v$ is realized by a scalar field with energy density $\rho(\phi) = \dot{\phi}^2/2 + (\nabla \phi)^2/2 + V(\phi)$, in which the potential energy density dominates $V(\phi) \gg \frac{1}{2} \dot{\phi}^2, \frac{1}{2} (\nabla \phi)^2$. (1)

Most field theory descriptions of inflation represent the vacuum energy through a scalar field satisfying eq. (1), with $\phi$ referred to as the inflaton. The goal of inflationary scalar field dynamics is to sustain the vacuum energy sufficiently long for expansion of the scale factor to exceed observational lower bounds and then end the inflationary epoch by entering the radiation dominated epoch. In the context of dynamical models, isentropic inflation is referred to as super-cooled inflation and non-isentropic inflation is referred to as warm inflation.

Scalar field models that represent the vacuum energy have no fundamental motivation. They provide a convenient mode for studying the complex dynamics of inflation. In addition, phase transitions can be represented by such models, and that is another key element to the particle physics picture of the early universe.
From the point of view of particle physics, a system of interacting fields will exchange energy at all times, with the inflationary epoch having no a priori reason for being different. To insist otherwise imposes a sharp division between an expansion regime with negligible energy exchange and then a subsequent reheating period, which is essentially a second Big-Bang.

On the one hand, for a first-order phase transition with bubble nucleation kinetics, this is a natural scenario dictated by the dynamics. This conception of Guth’s proved unsuccessful due to the conflicting requirements of having a slow nucleation rate to the true vacuum for obtaining adequate expansion versus a fast rate to allow bubble collisions after expansion, which would reheat the universe.

On the other hand, for continuous transitions, no qualitative feature of the dynamics requires a sharp division between the expansion and heating periods. Such a division can be constructed by requiring the scalar field potential to be ultra-flat during the inflationary expansion period and then sharply cusped to permit a subsequent reheating period that ends inflation and begins the radiation dominated regime. This is the new inflation picture. While the scalar field in on the ultra-flat portion of the potential, it will have weak self-interaction, since such a potential requires this. Also, it will interact weakly with other fields, since inflationary expansion will rapidly dilute any pre-existing field energy apart from the vacuum energy. An ultra-flat potential is sufficient for satisfying the requirements of inflation, but it is not necessary.

The general case is to allow the scalar field to interact with other fields during the entire evolution down the potential well. With this follows reactive forces on the inflaton, and they can slow the motion of the inflaton. The generic kinetics of continuous phase transitions for terrestrial systems, Ginzburg-Landau kinetics, is emphatic on the dissipational properties of the order parameter. Cosmological theories of phase transitions assume that the statistical mechanical principles on such large scales are no different from those on terrestrial scales. Thus Ginzburg-Landau kinetics is a viable possibility for the inflaton, provided that appropriate conditions can be realized. In particular the dissipational dynamics of Ginzburg-Landau kinetics assumes the presence of a large heat bath which interacts with the order parameter. For inflation, this requirement imposes the following self-consistency condition. The inflaton must release adequate vacuum energy into the heat bath to compensate for dilution of the heat bath due to inflationary expansion. Simultaneously the reaction of the heat bath on the inflaton must be sufficient to slow the inflaton’s roll down the potential. The questions are first can these requirements be satisfied by a sensible phenomenological dynamics and second can this dynamics be derived from quantum field theory?
Consider stochastic evolution for the inflaton governed by the Langevin-like equation
\[ \ddot{\phi}(t) + \left[ \Gamma + 3 \frac{\dot{R}(t)}{R(t)} \right] \dot{\phi} + V'(\phi) = \eta(t) \]
(2)
where \( \eta(t) \) is a random force function with vanishing ensemble averaged expectation value \( \langle \eta(t) \rangle = 0 \). Consider the limit of strong dissipation \( \Gamma \gg \dot{R}(t)/R(t) \) and the overdamped regime
\[ \Gamma |\dot{\phi}| \gg |\ddot{\phi}|. \]
(3)
Then eq. (2) has the Ginzburg-Landau form
\[ \frac{d\phi}{dt} = -\frac{1}{\Gamma} \frac{dV(\phi)}{d\phi}. \]
(4)
In this limit, the inflaton has a vacuum equation of state since \( \rho_\phi \approx V(\phi) \) which for a scalar field implies \( p_\phi = -\rho_\phi \). In addition, in a two fluid model composed of the scalar field and radiation, by energy conservation
\[ \dot{\rho}_r = -4 \frac{\dot{R}(t)}{R(t)} \rho_r - \dot{\rho}_\phi. \]
(5)
The Friedmann cosmology determined by eqs. (4) and (5) has been studied in [7] for a variety of vacuum functions \( V(\phi) = \lambda M^{4-n}(M - \phi)^n \). It was found there that for \( 2 \leq n < 4 \) the scale factor will go from a radiation dominated behavior into an inflationary behavior and then smoothly back to a radiation dominated behavior, and the latter occurring without reheating. The expansion e-folds, \( N_e \), during inflation and the drop in \( \rho_r(t) \) during the inflationary period are determined by the index \( n \) of the potential and the dissipative coefficient \( \Gamma \). For example for \( n = 2 \), the quadratic limit, \( N_e = \sqrt{2\pi/(3\lambda)}(\Gamma/m_p) \) and \( \rho_r(\tau_{EI})/\rho_r(\tau_{BI}) \approx 1/(4N_e^2) \), where the subscripts \( BI \) and \( EI \) signify begin and end inflation respectively.

The solutions in [7] are an existence proof of the warm inflation regime. A fundamental justification for such a dynamics requires deriving equation (2) from first principles and demonstrating the consistency of the limit eq. (3). Eq. (2) should not be confused with similar looking equations in earlier reheating models [9], since in the warm inflation case, the dissipative term represents frictional forces that arise from interaction of \( \phi \) with the heat bath. Furthermore, the overdamped limit eq. (3) is equivalent to an adiabatic limit. Under such conditions, eqs. (2) and (3) are an outcome of quantum mechanics both in flat spacetime [10] and for the cosmological setting of warm inflation [11]. These
equations have also been derived from a particular quantum field theory model in \[14\].

The fundamental origin of dissipation arises from the coupling of the inflaton to other fields which comprise the heat bath. In quantum field theory this implies the interactions $\phi^2 \chi^2$, $\phi \psi \bar{\psi}$ and $\phi^2 A^\mu_i A_{i\mu}$, for coupling to bosons $\chi$, fermions $\psi$ and gauge fields $A^\mu_i$. For such couplings, $\phi$ acts as a mass to the respective heat bath field and dissipation effects are only relevant when the mass it induces is less than or of order the temperature scale $T$. To enhance dissipative effects for large displacements of $\phi$, the couplings can also be modified by shifting $\phi$ only in the relevant interaction term. For example, for the bosonic case the shift $\phi^2 \chi^2 \to (\phi - M)^2 \chi^2$. Thus for several heat bath fields a distributed mass model suggests itself, which for the scalar case is

$$\sum_i g_i (\phi - M_i)^2 \chi_i^2$$  \hspace{1cm} (6)

or similarly a continuous mass model

$$\int d\mu g(\mu) (\phi - \mu)^2 \chi_\mu^2,$$  \hspace{1cm} (7)

with the mass spectrum given by $g_i, g(\mu)$ respectively. Such models could be motivated if the high energy world below the Planck scale and well above the electroweak scale tended away from an organized group theoretical structure towards a random one.

A second requirement of inflationary models is producing observationally consistent density perturbations $\delta \rho/\rho \sim 10^{-5}$, which is\[14\]. The general formula for $\delta \rho(t)/\rho(t)$ at horizon entry $t_f$ for perturbations produced by a scalar field, that exited the horizon at $t_i$ during inflation, is\[13\]

$$\frac{\delta \rho}{\rho} = \frac{\delta \rho}{\rho}(t_f) = \frac{V'(\phi(t_i)) \delta \phi(t_i)}{V(\phi(t_i))} \frac{1}{1 + w}$$  \hspace{1cm} (8)

where $w \equiv p/\rho$. In the warm inflation case\[13\]

$$\delta \phi^2 (H_i \hat{e}, t_i) = \int_{V=1/H_i^3} d^3\chi e^{iH_i \cdot \hat{e} \cdot \chi} (\phi(\chi, t_i) \phi(0, t_i))_\beta = O(1) \frac{H_i^3 T}{V'(\phi(t_i))},$$  \hspace{1cm} (9)

where $H_i$ is the (slowly varying) Hubble parameter at time $t_i$ during warm inflation. In warm inflation $\rho_r \gg \dot{\phi}^2/2$ so that

$$\frac{\delta \rho}{\rho}(t_f) = \frac{3V'(\phi(t_i)) \delta \phi(t_i)}{4\rho_r(t_i)}.$$  \hspace{1cm} (10)
In general the presence of non-negligible $\rho_r$ tends to suppress $\delta\rho/\rho$. In models that have been examined, regions with $\delta\rho/\rho < \ll 10^{-5}$ are much simpler to construct than the opposite region of large amplitude. In a phenomenological warm inflation model with a SU(5) Coleman-Weinberg potential, observational consistency has been obtained for both expansion e-folds $N_e > 60$ and density perturbations $\delta\rho/\rho \sim 10^{-5}$.

In conclusion, the warm inflation scenario has been shown to be consistent with observation for certain models and has nice features for treatment by quantum field theory. The warm inflation regime also has interesting possibilities for production of large scale cosmic magnetic fields and baryogenesis, since in a sense, the warm inflation regime is like a radiation dominated regime except with inflationary expansion.

References

1. A. H. Guth, *Phys. Rev.* D 23, 347 (1981).
2. A. Albrecht and P. J. Steinhardt, *Phys. Rev. Lett.* 48, 1220 (1982).
3. A. Linde, *Phys. Lett.* B 108, 389 (1982).
4. P. Spindel and R. Brout, *Phys. Lett.* B 320, 241 (1994).
5. H. P. de Oliveria and R. O. Ramos, *Phys. Rev.* D 57, 741 (1998); A. V. Nesteruk, R. Maartens, and E. Gunzig, astro-ph/9703137.
6. A. Berera, *Phys. Rev. Lett.* 75, 3218 (1995).
7. A. Berera, *Phys. Rev.* D 55, 3346 (1997).
8. A. Guth and E. Weinberg, *Nucl. Phys.* B 212, 321 (1983).
9. A. Albrecht, P. J. Steinhardt, M. S. Turner and F. Wilczek, *Phys. Rev. Lett.* 48, 1437 (1982); A. D. Dolgov and A. D. Linde, *Phys. Lett.* B 116, 329 (1982); L. F. Abbott, E. Farhi, and M. B. Wise, *Phys. Lett.* B 117, 29 (1982).
10. A. O. Caldeira and A. J. Leggett, *Ann. Phys.* 149, 374 (1983).
11. A. Berera, *Phys. Rev.* D 54, 2519 (1996).
12. A. Berera, M. Gleiser and R. O. Ramos, hep-ph/9803394.
13. J. M. Bardeen, *Phys. Rev.* D 22, 1882 (1980).
14. A. Berera and L. Z. Fang, *Phys. Rev. Lett.* 74, 1912 (1995).
15. R. Maartens and D. Tilley, *Gen. Rel. Grav.* 30 289, 1998.
16. P. L. Biermann, H. Falcke, Proceeding of Frontiers in Contemporary Physics International Lecture and Workshop, Vanderbilt University 1997; A. Berera, talk at same conference.