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Gauge-Mediated Curvature of the Flat Directions During Preheating *

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Abstract

Usually one expects the inflaton field to be coupled to some gauge-charged particles allowing for its decay during reheating. Such particles then play a role of the messengers for the gauge-mediated supersymmetry breaking during and (shortly) after the inflation and radiatively induce soft masses to all other D -flat directions. We show that during the preheating stage this gauge-mediated soft masses are typically much greater than the Hubble parameter during inflation. The dramatic role is played by the supersymmetry (SUSY) breaking due to the parametric resonance effect, which ensures that the inflaton predominantly decays into the bosons and not the fermions. Difference in the Fermi-Bose occupation numbers results in the large gauge-mediated soft masses, which determine the post-inflationary evolution of the flat directions, suggesting that nonthermal phase transitions mediated by gauge messengers may play a crucial role in the Affleck-Dine mechanism for the generation of the baryon asymmetry.

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In the low-energy minimal supersymmetric standard model there exist a large number of D -flat directions along which squark, slepton and Higgs fields get expectation values. In flat space at zero temperature exact supersymmetry guarantees that the effective potential along these D -flat directions vanishes to all orders in perturbation theory (besides the possible presence of nonrenormalizable terms in the superpotential [1]). In the commonly studied supergravity scenario, supersymmetry breaking may take place in isolated hidden sectors [2] and then gets transferred to the other sectors by gravity. The typical curvature of D -flat directions resulting from this mechanism is

$$\widetilde{m}^2 \sim \frac{|F|^2}{M_{\text{Pl}}^2}, \quad (1)$$

where F is the vacuum expectation value (VEV) of the F -term breaking supersymmetry in the hidden sector. In order to generate soft masses of order of M_W in the matter sector, F is to be of order of $(M_W M_{\text{Pl}})$ and sfermion and Higgs masses along flat directions turn out to be in the TeV range. D -flat directions do not cause any cosmological problems. Some of them are extremely important in the Affleck-Dine (AD) scenario for baryogenesis [3] if large expectation values along flat vacua are present during the early stages of the evolving Universe. This is a necessary condition for the AD mechanism to be operative. In generic supergravity theories soft supersymmetric breaking masses are of order of the Hubble parameter H_I (typically $\sim 10^{13}$ or so) GeV during inflation [4]. This is due to the fact that inflation provides a nonzero energy density $V \sim |F|^2$ which breaks supersymmetry. Since in the inflationary phase the vacuum energy dominates, the Hubble parameter is given by $H^2 = (8\pi V/3M_{\text{Pl}}^2)$ and therefore the curvature along the D -flat directions becomes $\widetilde{m}^2 = cH_I^2$, where c may be either positive or negative. This fact has dramatic effects on what discussed so far. For the AD mechanism, large squark and slepton VEV's do not result if the induced soft mass squared is positive, but they do occur if it is negative and an acceptable baryon asymmetry can be obtained without subsequent entropy releases.

During inflation, D -flat directions, however, can get larger gauge-mediated soft masses [5]. It is well known that in the present vacuum (with zero energy) the gauge interaction can be of more efficient messenger of the SUSY breaking than the gravity, provided the messenger scale is below M_{Pl} . This is what usually happens in models with Gauge-

mediated Supersymmetry Breaking (GMSB) [6] where the message about the supersymmetry breaking (in a gauge-invariant direction) from the hidden sector is transferred to the observable sector through gauge interactions by the messenger sector. The latter is formed by some heavy superfields, transforming under the gauge group G as a real or conjugate representation, which suffer from a tree level supersymmetry breaking. The crucial point, however, is that even if gauge-mediated corrections are zero in the present vacuum they had to be important during inflation if the superfield X , which is dominating inflation, is coupled to some of the gauge nonsinglet superfields ϕ

$$W = g X \phi^2. \quad (2)$$

In most of the inflationary scenarios such couplings are expected to be there, in order to allow for the efficient reheating through the final decay of the inflaton field. In such a case ϕ -fields would play a role of the messengers of the gauge-mediated SUSY breaking during inflation and all other D -flat directions would obtain a radiative two-loop mass given by [5]

$$\widetilde{m}^2 \sim \left(\frac{\alpha}{4\pi}\right)^2 g^2 \frac{|F_X|^2}{M_\phi^2}, \quad (3)$$

where $M_\phi \sim g |X|$ is the mass term of the ϕ , α is the gauge coupling of the gauge group G and g is the coupling constant relating the superfield ϕ to the inflaton. Using again the relation between the vacuum energy and the Hubble constant during inflation $H_I^2 \sim |F_X|^2 / M_{\text{Pl}}^2$, one can rewrite the above relation as

$$\widetilde{m}^2 \sim H_I^2 \left(\frac{\alpha}{4\pi}\right)^2 \left(\frac{M_{\text{Pl}}}{|X|}\right)^2, \quad (4)$$

which shows that, in general, \widetilde{m}^2 may be larger than H_I^2 (typical magnitude of the soft masses induced by gravitational sources) if $|X|$ is somewhat below M_{Pl} [5]. Notice that these corrections are independent of coupling constant g . Although there is no generic proof, usually, e.g. for the low representations of the simple Grand Unified Groups (GUT's), these two-loop radiative corrections to \widetilde{m}^2 are positive and the D -flat directions are expected to be stabilised at the origin during inflation. This would be a disaster for the AD mechanism of baryogenesis which requires large expectation values along the flat

vacua after inflation to be operative. However, it has been recently point out [7] that if the inflaton couples to the superfields of the messenger sector and the latter are in the complex representation, supersymmetry breaking during inflation can generate one-loop Fayet-Iliopoulos D -terms. The corresponding soft masses are proportional to the abelian generators of G (*e.g.* hypercharge in the GUT's) and, therefore, can have either sign. They can dominate the gauge-mediated two-loop soft breaking terms, being of the order of

$$\widetilde{m}^2 \sim H_1^2 \left(\frac{\alpha}{4\pi} \right) \frac{M_{\text{pl}}}{|X|}. \quad (5)$$

Such (negative) masses can destabilize the sfermion flat directions during inflation, playing a crucial role for the AD mechanism of baryogenesis. Previously the induced gauge-mediated soft masses were analysed only during inflation. However, the crucial role for the post-inflationary evolution of the flat directions is played by their soft masses just before the reheating process. The aim of this letter is to analyse the issue of gauge-mediated supersymmetry breaking during a particular stage of the evolution of the early universe. The epoch we are referring to is called preheating [8] and is expected to occur after the end of chaotic inflation. The crucial relevance of supersymmetry breaking at preheating has been first pointed out in [9]. At the very beginning of this period, which is dominated by the coherent oscillations of the inflaton field, one can distinguish two possible cases depending whether the classical expectation value of the messenger¹ is fixed at its minimum or undergoes coherent oscillations together with the inflaton. The latter will be the case if the VEV of some ϕ component is nonzero in the minimum about which the inflaton oscillates. We will start considerations from the former case assuming no coherent oscillations of ϕ . Kofman, Linde and Starobinsky have recently pointed out that the explosive decay of the inflaton occurs at the first stage of reheating through the phenomenon of parametric resonance [8]. The inflaton energy is released

¹Below we will refer to the gauge-charged superfields coupled to the inflaton as ‘messengers’, although they are not assumed to be necessarily a messengers in the present vacuum, but only in the early universe.

in the form of inflaton decay products, whose occupation number is extremely large, and have energies much smaller than the temperature that would have been obtained by an instantaneous conversion of the inflaton energy density into radiation. Since it requires several scattering times for the low-energy decay products to form a thermal distribution, it is rather reasonable to consider the period in which most of the energy density of the Universe was in the form of the nonthermal quanta produced by inflaton decay as a separate cosmological era, dubbed as preheating to distinguish it from the subsequent stages of particle decay and thermalization which can be described by the techniques developed in [10]. Several aspects of the theory of explosive reheating have been studied in the case of slow-roll inflation [11] and first-order inflation [12]. One of the most important consequences of the stage of preheating is the possibility of nonthermal phase transitions with symmetry restoration [13–15]. These phase transitions appear due to extremely strong quantum corrections induced by particles produced at the stage of preheating. What is crucial for our considerations is that parametric resonance is a phenomenon peculiar of particles obeying Bose-Einstein statistics. Parametric resonant decay into fermions is very inefficient because of Pauli’s exclusion principle. This means that during the preheating period the Universe is only populated by a huge number of soft bosons and the occupation numbers of bosons and fermions belonging to the supermultiplet coupled to the inflaton superfield are completely unbalanced. Supersymmetry is then strongly broken during the preheating era [9] and large loop corrections may arise since the usual cancellation between diagrams involving bosons and fermions within the same supermultiplet is no longer operative. We shall see that the curvature along D -flat directions during the preheating era is much larger than the effective mass that they acquire in the inflationary stage. This makes the details of the effective potential along D -flat directions during inflation almost irrelevant as far the initial conditions of the condensates along the D -flat directions is concerned. Let us first assume that the F_X -term corresponding to the superfield X is dominating inflation and that the gauge-charged superfield ϕ is in the real (say adjoint) representation of G . The simplest superpotential (leading to chaotic inflation [16] and to the subsequent resonance decay of the inflaton) one can envisage relating X to the supermultiplet ϕ is

$$W = M_X X Z + g X \phi^2 \tag{6}$$

where Z is another gauge singlet superfield² and $\sim 10^{13}$ GeV for the density perturbations generated during the inflationary era to be consistent with COBE data [17]. G -invariant contraction of the indices is assumed. There are several possible choices of the discrete or continuous symmetries under which the above form is the most general renormalizable one. One example is a phase symmetry under which $X \rightarrow e^{i\theta} X$ and $Z \rightarrow e^{-i\theta} Z$ and $\phi \rightarrow e^{-i\theta/2} \phi$. The global minimum of the theory is at $X = 0$ and $\phi = \sqrt{-MZ/g}$ = arbitrary. For any non-zero value of X the minimum in all other fields is at $Z = \phi = 0$ and their masses are $M_Z^2 = M_X^2$ and $M_\phi^2 = g^2|X|^2$ respectively. Therefore, assuming the chaotic initial conditions $|X| \gg M_X$, we expect that ϕ -field will quickly settle at the origin due to very large curvature in its direction. Contrastly, the curvature in the X -direction is small and inflation occurs during the slow rolling of the scalar field X from its very large value³. Then inflaton oscillates with an initial amplitude $X_0 \sim 10^{-1} M_{\text{Pl}}$. Within few dozen oscillations the initial energy density $\rho_X \sim M_X^2 X_0^2$ is transferred through the interaction $g^2 X^2 \phi^2$ to *bosonic* ϕ -quanta in the regime of parametric resonance [8]. At the end of the broad parametric resonance the field X drops down to $X_e \sim 10^{-2} M_{\text{Pl}}$ and parametric resonance only occurs if $gX_e > M_X$. This implies $g > 10^{-4}$. Notice that the flatness of the inflaton potential during inflation is preserved for such large values of couplings g by supersymmetric cancellations. In the above example the one-loop corrections to the inflaton potential are simply zero, because no Fermi-Bose mass splitting occurs along the inflationary trajectory (the only non-zero F -component is the one of the Z -superfield $F_Z = M_X X$, which does not couple to the other fields). Below

²Without the Z superfield the preheating is necessarily marked by the coherent oscillations of ϕ (see the text below), in which case the parametric resonance requires further investigation. Here we want to make situation maximally adequate to the one studied in [8].

³The curvature in the Z -direction is also small, so that in principle both singlets can roll slowly, however for simplicity we assume that inflation in the X direction lasts longer, so that when it starts oscillations about the minimum, Z -field is already fixed there.

we will consider the case when the inflation is dominated by F_X and one-loop corrections are present, but, in any case, they only modify the inflaton potential by a logarithmic factor with a small coefficient). At the end of the preheating era the Universe is expected to be filled up with noninteracting ϕ -bosons with relatively small energy per particle, $E_\phi \sim 10^{-1} \sqrt{g M_X M_{\text{Pl}}}$ and with very large occupation numbers $n_\phi/E_\phi^3 \sim g^{-2}$. Here we are assuming that the energy E_ϕ is larger than any bare mass of the superfield ϕ (which is automatically the case in the above model). Our results do not crucially depend upon this assumption. The leading contribution to the curvature of D -flat directions comes from the two-loop exchange of the ϕ -bosons which are produced during the parametric decay of the inflaton and form the noninteracting gas of particles out of equilibrium during the preheating stage. Unfortunately, one cannot use the standard imaginary-time formalism since in the nonequilibrium case there is no relation between the density matrix of the system and the time evolution operator, which is of essential importance in the formalism. There is, however, the real-time formalism of Thermo Field Dynamics (TFD), which suites our purposes [18]. This approach leads to a 2×2 matrix structure for the free propagator for the ϕ -boson (only the (11)-component is physical)

$$\begin{pmatrix} D_{11}(K) & D_{12}(K) \\ D_{21}(K) & D_{22}(K) \end{pmatrix} = \begin{pmatrix} \Delta(K) & 0 \\ 0 & \Delta^*(K) \end{pmatrix} + \begin{pmatrix} f_\phi(k) & \theta(k_0) + f_\phi(k) \\ \theta(-k_0) + f_\phi(k) & f_\phi(k) \end{pmatrix} \times 2\pi\delta[K^2 - m_\phi^2], \quad (7)$$

with the usual vacuum Feynman propagator

$$\Delta(K) = \frac{i}{K^2 - m_\phi^2 + i\epsilon}. \quad (8)$$

The distribution function f_ϕ is chosen such that the number density of particles, $n_\phi = (2\pi)^{-3} \int d^3p f_\phi(p)$ and setting it equal to $\sim \rho_\phi/E_\phi$. Notice that at the preheating stage the occupation number $f_{\tilde{\phi}}$ of the fermionic partner $\tilde{\phi}$ of the ϕ -boson is much smaller than f_ϕ : even though supersymmetric cancellation may occur when only vacuum propagators are inserted, such a cancellation is no longer operative in the gas of ϕ -bosons where $f_\phi \gg f_{\tilde{\phi}}$. Making use of the standard TFD Feynman rules one can show that during the preheating era D -flat directions acquire a correction to the mass squared

$$\widetilde{m}^2 \sim \alpha^2 \frac{n_\phi}{E_\phi} \sim \frac{10^{-2}}{g} \alpha^2 M_X M_{\text{Pl}}, \quad (9)$$

which is much larger than the two-loop contribution $\sim (\alpha/4\pi)M_X^2$ that soft breaking terms may receive during inflation. Now let us consider the case when the messenger field undergoes the coherent oscillations driven by the oscillations of the inflaton. This will happen when the instant VEV of ϕ is a nontrivial function of the inflaton VEV. Such a behaviour is exhibited already by a simplest system: single inflaton superfield coupled to the messengers

$$W = \frac{1}{2} (M_X X^2 + g X \phi^2). \quad (10)$$

The global minimum is at $X = \phi = 0$, but for $0 < |X| < X_c = \frac{M_X}{g}$ the instant minimum of ϕ is at $|\phi| = \sqrt{\frac{2}{g}|X|(M_X - g|X|)}$. Thus, whenever X drops below X_c , ϕ will undergo the driven coherent oscillations. For $X > X_c$, ϕ vanishes and the tree level potential is dominated by the inflaton F -term $F_X = M_X X$, which splits masses of the Fermi-Bose components in the ϕ superfield. This splitting result in two things: 1) the one-loop corrections to the inflaton slope, which for large $|X|$ behave as

$$(\Delta V_{\text{eff}})_{|X| \rightarrow \infty} \sim \frac{g^2}{16\pi^2} M_X^2 |X|^2 \ln|X|^2; \quad (11)$$

and 2) the two-loop universal (up to charges) gauge-mediated soft masses for the D -flat directions

$$\widetilde{m}^2 \sim \left(\frac{\alpha}{\pi}\right)^2 M_X^2 \quad (12)$$

After the inflaton VEV drops below the critical value X_c both fields start to oscillate about the global minimum. Parametric resonance in such a case needs a special investigation, which will not be attempted here. Instead we will argue that there is an independent source of the supersymmetry breaking due to a coherent oscillations of the ϕ VEV. This condensate can be regarded as a gas of cold bosons with energies $\sim M_X$ and occupation numbers $n_\phi \sim \frac{M_X^3}{g^2}$. Again, since there are no fermions the two-loop gauge diagrams do not cancel out and the resulting soft masses can be estimated as

$$\widetilde{m}^2 \sim \frac{\alpha^2}{g^2} M_X^2 \quad (13)$$

Although these masses are smaller than (9), they are greater than the gravity-mediated contribution and thus, will play a dominant role in the cases in which parametric resonance is suppressed. Let us now assume another case, as suggested in [7], that the inflaton couples to some superfields of the messenger sector belonging to the complex representations. We introduce a pair of messengers ϕ and $\bar{\phi}$ with an opposite charges under a certain $U(1)$ -group. We will think of this $U(1)$ as being an abelian subgroup of some Grand Unified Theory symmetry under which ϕ and $\bar{\phi}$ transform in the complex representations. The simplest superpotential which leads to the messenger VEVs being fixed in their (minimum all the way until the inflaton settles in the global vacuum) has the form

$$W = W_0 + g X \bar{\phi} \phi \tag{14}$$

where W_X is a part of the superpotential responsible for the slope of inflaton potential, which can be taken to be $W_0 = M_X X Z$ as in (6). The inflation in this model will proceed in the same way as discussed above, except the preheating stage. The crucial difference is that now inflaton through the parametric resonance will decay into two different bosons ϕ and $\bar{\phi}$ which in general can have different occupation numbers. This can be the case if, for instance, one of these particles has a nonzero bare mass because of mixing with some other superfield A in the superpotential⁴

$$W' = M_A \phi A \tag{15}$$

and the bare mass M_A is so high to stop the production of ϕ -quanta during the parametric resonance. However, if the scale M_A is very high (much greater than $\sqrt{M_X X_e}$), it will suppress the production of $\bar{\phi}$ quanta as well. This is because of the superdecoupling arguments: below the energies $\sim M_A$ the ϕ and A fields decouple and the low-energy superpotential can not include any gauge-invariant coupling of X and $\bar{\phi}$ superfields. However, these arguments are not applicable if the SUSY-breaking scale during oscillations

⁴The mass of the ϕ -quanta induced by D -terms in presence of a nonvanishing VEV during inflation along AD flat vacua is zero.

(in our case $\sim \sqrt{M_X X_e}$) is comparable to M_A . So with $M_X \sim 10^{13}$ GeV the right order of magnitude for M_A would be somewhere around the GUT scale. Another possible source of the asymmetry between ϕ and $\bar{\phi}$ states can be their different cross couplings with an inflaton field in the Kähler potential

$$C^2 \int d^4\theta \frac{1}{4 M^2} X X^+ \bar{\phi} \bar{\phi}^+ = C^2 \frac{|F_X|^2}{M^2} |\bar{\phi}|^2 + \dots \quad (16)$$

Such couplings with $M \sim \frac{M_{\text{Pl}}}{\sqrt{8\pi}}$ will generically be presented in supergravity theories. Assuming $g = 0$ and $W_0 = M_X X^2$ in (14), this interaction induces an effective cross coupling $|X|^2 |\bar{\phi}|^2$ in the potential with the coefficient $g_{\text{eff}}^2 \sim C^2 \left(\frac{M_X}{M}\right)^2$ ⁵. Then, the initial energy density of the inflaton $\rho_X \sim M_X^2 X_e^2$ may be transferred through the interaction $g_{\text{eff}}^2 X^2 |\bar{\phi}|^2$ to *bosonic* $\bar{\phi}$ -quanta in the regime of parametric resonance [8]. Because of the mass difference, at the preheating stage the messengers ϕ and $\bar{\phi}$ have different number densities, $n_{\bar{\phi}} \gg n_{\phi}$. Their contributions to the curvature along D -flat directions do not cancel in the one-loop diagrams giving rise to a Fayet-Iliopoulos D -terms. The corresponding soft masses for sfermion fields are proportional to

$$\widetilde{m}^2 \sim (4 \pi \alpha) \frac{n_{\bar{\phi}}}{E_{\bar{\phi}}} \sim (4 \pi \alpha) \frac{10^{-2}}{g_{\text{eff}}} M_X M_{\text{Pl}}, \quad (17)$$

which is larger than the one-loop correction (5) obtained by soft breaking masses during inflation. Our result implies that during the explosive stage of preheating gauge-mediated supersymmetry breaking is stronger than at the stage of inflation, suggesting that non-thermal phase transitions mediated by gauge messengers may play a crucial role in the AD mechanism for the generation of the baryon asymmetry [9].

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⁵In the previous example, the coupling (16) with $M = M_A$ can be generated for very large values of M_A after integrating out the heavy ϕ and A states. However, in this case, as a matter of the potential structure, one can not use simply $W_0 = M_X X^2$.

REFERENCES

- [1] J. Ellis, K. Enqvist, D.V. Nanopoulos and K. Olive, Phys. Lett. **B191** (1987) 343; K.W. Wong, Nucl. Phys. **B321** (1989) 528.
- [2] R. Barbieri, S. Ferrara and C.A. Savoy, Phys. Lett. **B119** (1982) 343; A.C. Chamseddine, R. Arnowitt and P. Nath, Phys. Rev. Lett. **49** (1982) 970; L. Hall, J. Lykken and S. Weinberg, Phys. Rev. **D27** (1983) 2359.
- [3] I. Affleck and M. Dine, Nucl. Phys. **B249** (1985) 361.
- [4] G. Dvali, IFUP-TH-09-95, hep-ph/9503259; M. Dine, L. Randall and S. Thomas, Phys. Rev. Lett. **75** (1995) 398; *ibid* Nucl. Phys. **458** (1996) 291. The fact that during inflation fields generically have masses of order of H has been noted also by, M. Dine, W. Fischler and D. Nemeschinsky, Phys. Lett. **B136** (1984) 169; G.D. Coughlan, R. Holman, P. Ramond and G.G. Ross, Phys. Lett. **B140** (1984) 44; A.S. Goncharov, A.D. Linde, and M.I. Vysotsky, Phys. Lett. **147B**, 279 (1984); O. Bertolami and G.G. Ross, Phys. Lett. **B183** (1987) 163; E.J. Copeland et *al.*, Phys. Rev. **D49** (1994) 6410.
- [5] G. Dvali, Phys. Lett. **B355** (1995) 78.
- [6] M. Dine, W. Fischler, and M. Srednicki, Nucl. Phys. **B189** (1981) 575; S. Dimopoulos and S. Raby, Nucl. Phys. **B192** (1981) 353; M. Dine and W. Fischler, Phys. Lett. **B110** (1982) 227; M. Dine and M. Srednicki, Nucl. Phys. **B202** (1982) 238; M. Dine and W. Fischler, Nucl. Phys. **B204** (1982) 346; L. Alvarez-Gaumé, M. Claudson, and M. Wise, Nucl. Phys. **B207** (1982) 96; C.R. Nappi and B.A. Ovrut, Phys. Lett. **B113** (1982) 175; S. Dimopoulos and S. Raby, Nucl. Phys. **B219** (1983) 479.
- [7] G. Dvali, CERN-TH/96-129, hep-ph/9605445.
- [8] L. Kofman, A.D. Linde and A.A. Starobinsky, Phys. Rev. Lett. **73**, 3195 (1994).
- [9] G. Anderson, A. Linde and A. Riotto, FERMILAB-Pub-96/078-A, hep-ph/9606416, submitted to Phys. Rev. Lett. .

- [10] A.D. Dolgov and A.D. Linde, Phys. Lett. **B**, L.F. Abbott, E. Fahri and M. Wise, Phys. Lett. **B117** (1982) 29.
- [11] J. Traschen and R. Brandenberger, Phys. Rev. D **42**, 2491 (1990); Y. Shtanov, J. Traschen, and R. Brandenberger, Phys.Rev. D **51**, 5438 (1995); D. Boyanovsky, H.J. de Vega, R. Holman, D.S. Lee, and A. Singh, Phys. Rev. D **51**, 4419 (1995). M. Yoshimura, Prog. Theor. Phys. **94** (1995) 873; H. Fujisaki et *al.*, hep-ph/9508378 and hep-ph/9511381; D. Boyanovsky, M. D’Attanasio, H.J. de Vega, R. Holman and D.-S. Lee, Phys. Rev. **D52**, 6805 (1995); D. Kaiser, Phys. Rev. **D53** (1996) 1776; S. Kasuya and M. Kawasaki, hep-ph/9603317; S.Yu. Khlebnikov and I. Tkachev, OSU-TA-08/96, hep-ph/9603378.
- [12] E.W. Kolb and A. Riotto, FERMILAB-Pub-96-036-A, astro-ph 9602095.
- [13] L. Kofman, A.D. Linde and A.A. Starobinsky, Phys. Lett. **76** (1996) 1011.
- [14] I. Tkachev, Phys. Lett. **B376** (1996) 35.
- [15] A. Riotto and I. Tkachev, Fermilab-Pub-96/88-A, astroph/9604444.
- [16] A. Linde, Phys. Lett. **B129**, 177 (1983).
- [17] G.F. Smoot et *al.*, Astrophys. J. Lett. **396** (1992) L1.
- [18] H. Umezawa, H. Matsumoto and M. Tachiki, *Thermo Field Dynamics and Condensed States*, North Holland, 1982; P.A. Henning, Phys. Rep. **253** (1995) 235.