Eur. Phys. J. C 14, 547–552 (2000)
Digital Object Identifier (DOI) 10.1007/s100520000236

THE EUROPEAN PHYSICAL JOURNAL C
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Bosonic preheating in left-right-symmetric SUSY GUTs

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Received: 4 August 1999 / Published online: 14 April 2000 – © Springer-Verlag 2000

Abstract. We investigate the possibility of a bosonic preheating in the simplest model of supersymmetric Hybrid Inflation (F-term inflation), where the inflationary superpotential is of the O'Raifertaigh-Witten type. The end of inflation is related to a spontaneous symmetry breaking, which in the context of left-right symmetric models lowers the rank of the gauge group. Our results indicate that the possibility for bosonic preheating in this model is limited.

1 Introduction

The recent success in detecting neutrino masses has caused a renaissance of the idea of grand unification. In left-right-symmetric models, such as $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{(B-L)}$ or $SO(10)$, massive Majorana neutrinos arise naturally. As the difference of baryon- and lepton number is a gauged symmetry in these models, there could exist interesting mechanisms for the creation of the corresponding asymmetries. On the other hand, supersymmetry not only provides the right mass scale for the heavy neutrinos, but also gives the possibility of a inflationary potential in the GUT-Higgs-Sektor. Then it is possible to construct a cosmological model based on a consistent supersymmetric grand unified theory, which can be judged by both, its cosmological and its particle theoretical features. In this paper we will concentrate on an O'Raifertaigh-Witten model, which by the autors of [1] was shown to be a realisation of Linde's Hybrid Inflation scenario [2]. The superpotential can serve as a part of the Higgs sector of a left-right-symmetric model with global or local supersymmetry. The phase of inflation and the formation of density fluctuation in this model have been studied in [1,3]. Since a successful cosmological model needs a very effective mechanism for particle production after inflation, here we investigate the possibility of a bosonic preheating in the inflationary potential under consideration. As preheating possibly allows for the creation of superheavy particles, it is a very interesting scenario not only in the context of lepto- and baryogenesis but also for the creation of dark matter. Focusing on qualitative insights rather than quantitative accuracy, we tried to work with a minimum of numerical expense. Our investigation was inspired by those in [4,5], where preheating in the original version of Hybrid Inflation was studied. Our numerical results, which originally were presented in the context of an investigation of inflationary supersymmetric $SO(10)$-models in [6], are very similiar to those in a recent paper, based on a version of a NMSSM [7]. We will see, that the fourier methods used in their case as in ours have to be handled with care since the non-linear regime is reached early.

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2 Supersymmetric hybridinflation: the basic scenario

A common way to achieve an inflationary scenario from a supersymmetric theory, is to use the simplest O'Raifertaigh-Witten model for spontaneous symmetry breaking (SSB) [1]. It is given by the superpotential

$$W = X(\kappa \bar{C} C - \mu^2), \tag{2.1}$$

where the couplings are protected by either a continouus or a discrete R-symmetry, which could descend from a string theory. The superfield $C = (\hat{a}_C, \hat{\psi}_C)$ is a representation of the Lie-algebra A_G of the Lie-group G, and $\bar{C} = (\hat{a}_{\bar{C}}, \hat{\psi}_{\bar{C}})$ is the conjugate representation. Here the $\hat{a}$'s represent the scalar component while the $\hat{\psi}$'s are the spinors. The Superfield $X = (\hat{a}_X, \hat{\psi}_X)$ is an A_G-Singlet. The mass-scale μ could be caused by a string compactification. In the case, that G represents a gauge symmetry we have the tree-level scalar potential

$$\begin{aligned} V^{(0)} &= |F_C|^2 + |F_{\bar{C}}|^2 + |F_X|^2 + \frac{1}{2}\sum_{r=1}^{dim(G)} |D_r|^2 \\ &= \kappa^2 |a_X a_{\bar{C}}|^2 + \kappa^2 |a_X a_C|^2 + |\kappa a_{\bar{C}} a_C - \mu^2|^2 \\ &\quad + \frac{1}{2}\sum_{r=1}^{dim(G)} |D_r|^2. \end{aligned} \tag{2.2}$$

Here a_i represents the vacuum expectation value (vev) of the quantum field $\hat{a}_i$. Then $V^{(0)}$ is minimized for $\arg a_{\bar{C}} +$

$\arg a_C = 0$ and it is independent of $\arg a_{\bar{C}} + \arg a_X$ and $\arg a_C + \arg a_X$. Supersymmetry requires the D-terms to vanish identically for each single generator of the group G. If G represents a product group including an $U(1)$-factor, this means that $|a_C| = |a_{\bar{C}}|$.

In the following we will be interested only in the direction of the A_G-multiplets a_C and $a_{\bar{C}}$, which aquire a G breaking vev. Therefore they will be denoted by the same names. Now, concentrating on the $D = 0$-direction and using certain R-and gauge transformations [8] we can bring the scalar components a_i to the real axis. The corresponding canonically normalized scalarfields ϕ, σ and couplings λ, g are given by [8]:

$$\sigma := 2a_C = 2a_{\bar{C}} \tag{2.3}$$

$$\phi := \sqrt{2}a_X \tag{2.4}$$

$$\lambda := \frac{\kappa^2}{4} \tag{2.5}$$

$$g := \frac{\kappa}{\sqrt{2}} \tag{2.6}$$

$$M := \sqrt{\kappa}\mu, \tag{2.7}$$

where typically $g^2 = 2\lambda$. Using this conventions the scalar effective potential in terms of the fields ϕ and σ reads

$$V(\phi, \sigma) = \frac{1}{4\lambda}\left(M^2 - \lambda\sigma^2\right)^2 + \frac{1}{2}g^2\phi^2\sigma^2 + V^{(1)}(\phi), \tag{2.8}$$

which resembles very much Linde's original model of Hybrid Inflation [2]. Here $V^{(1)}(\phi, \sigma)$ represents the loop corrections to the tree level potential, which vanish for the supersymmetric case. For σ=0, $\phi > M/g =: \phi_c$ the gauge symmetry remains conserved while supersymmetry is broken in the sector of the gauge singlet X inducing loop-corrections $V^{(1)}(\phi, \sigma)$. Including 1-loop-corrections the effective potential in this regime is given by [1]

$$\begin{aligned} V(\phi > \phi_c) = \frac{M^4}{4\lambda}\Bigg(1 + \frac{\lambda}{4\pi^2}\Bigg[\ln\frac{2\lambda\phi^2}{M^2} \\ + \left(\frac{2\lambda\phi^2}{M^2} - 1\right)^2 \ln\left(1 - \frac{M^2}{2\lambda\phi^2}\right) \\ + \left(\frac{2\lambda\phi^2}{M^2} + 1\right)^2 \ln\left(1 + \frac{M^2}{2\lambda\phi^2}\right) \\ + \ln\frac{M^2}{\Lambda^2}\Bigg]\Bigg), \end{aligned} \tag{2.9}$$

where Λ is a renormalization scale. Since soft susy breaking terms lead to scalar masses of $\mathcal{O}(\text{TeV})$, such terms are neglegible compared to the GUT-scale mass parameter M. For a large range of the parameters κ and μ the potential above satisfies the Slow Roll conditions for inflation. Here inflation is caused by the "cosmological constant" $\frac{M^4}{4\lambda}$. The scalar gauge singlet is the only degree of freedom, which then has a nonvanishing vev ϕ. This vev very soon is dominated by the zero momentum Fourier mode and it forms a condensate. Referring to [9] it can be described by a homogeneous classical scalar field. For this reason we will call ϕ the inflaton and identify it with the zero momentum mode. The slow roll regime of the singlet is a typical aspect of the superpotential (2.1) and has become popular as effect of "a sliding field" outside the context of inflation [10].

The end of inflation is connected to the end of the gauge symmetric phase: as soon as ϕ reaches the "critical point" ϕ_c the effective mass $m_\sigma^2 = -M^2 + g^2\phi^2$ of the Higgs-field vanishes. The classical equations of motions alone cannot tell us something about the dynamics of the spontaneous symmetry breaking, since the derivative $V(\phi_c, \sigma = 0)_{,\sigma}$ vanishes identically and the vev of $\hat{\sigma}$ remains zero. But, following the argumentation of Garcia-Bellido and Linde [4], as ϕ slides towards zero, quantum fluctuations around this vev will get tachyonic masses stimulating an exponential growth of modes, whose momenta are smaller than the effective mass, $k < |m_\sigma|$, where k means the comoving momentum. The modes with $k > |m_\sigma|$ will not grow at all. The result is a spontaneous breakdown of the gauge symmetry caused by the inhomogenous distribution of the field $\hat{\sigma}$ with $\langle\hat{\sigma}\rangle = 0$. Inflation ends typically near to this point.

As pointed out in [4] the distribution of the Higgs field would be homogeneous on scales $l \sim |m_\sigma|^{-1}$ or even somewhat greater, if m_σ was constant in time. However, we will see soon that this mass is oscillatory. This changes the range and the mechanism of amplification of the "tachyonic modes". Due to this fact, the distribution $\sigma_k(t)$ in momentum space will be even more localized around the zero momentum mode, such that the spatial distribution look homogeneously even on a scale l which is somewhat greater than $|m_\sigma|^{-1}$. Assuming the zero momentum mode is dominating, this regime effectively may be described by a classical homogenous scalar field $\sigma(t)$, i.e. as the zero momentum mode of the Higgs field, which rolls down from the critical point. We will call $\sigma(t)$ the Higgs-condensate.

As pointed out in [3] local supersymmetry modifies the whole scenario because of the non-renormalizable terms, whose influence on the inflation depends on the value of the coupling constant κ. But in most of these cases the end of inflation is dominated by the renormalizable terms in the potential. Thus, non-renormalizable couplings will not contribute after the phase transition and will be ignored in the following investigation.

3 Evolution of the background

After SSB both of the fields $\hat{\phi}(t, \mathbf{x})$ and $\hat{\sigma}(t, \mathbf{x})$ form condensates, which essentially can be described by the two homogeneous classical fields $\phi(t)$ and $\sigma(t)$. Higgs or singlet particles will be identified with quantum fluctuations around these condensates. In this section we concentrate on the evolution of the background and we neglect the effect of the quantum fluctuations. For reasons of simplification we use a "M-rescaling" to natural variables:

$$t \to y := Mt$$

$$\mathbf{x} \to \xi := M\mathbf{x}$$

$$\mathbf{k} \to \mathbf{K} := \frac{\mathbf{k}}{M}$$

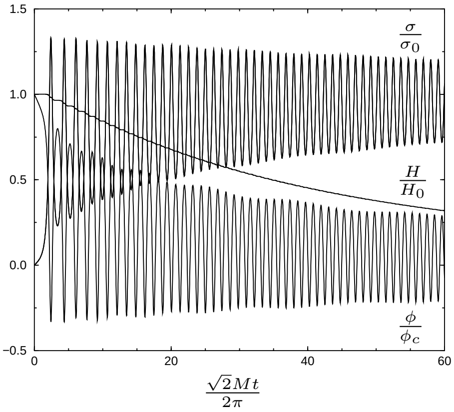

Fig. 3.1. Evolution of the background

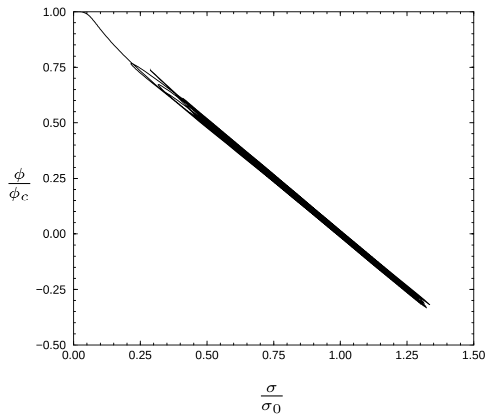

Fig. 3.2. Oscillations in the ϕ-σ-plane

$$\begin{aligned} \phi \to f &:= \frac{\phi g}{M} = \frac{\phi}{\phi_c} \\ \sigma \to s &:= \frac{\sigma\sqrt{\lambda}}{M} = \frac{\sigma}{\sigma_0}. \end{aligned} \tag{3.1}$$

Here the comoving momenta $\mathbf{K}$ and the scale factor $a(y)$ are normalized in such a way that $a(y_c) = 1$ is fullfilled, where y_c means the M-rescaled time at the phase transition.

Neglecting the sub-dominant loop corrections, which certainly will appear, as long as the supersymmetric minimum is not reached, the equations of motion after the phase transition read:

$$\left(\frac{d^2}{dy^2} + 3h(y)\frac{d}{dy} + 2s^2(y)\right) f(y) = 0 \tag{3.2}$$

$$\left(\frac{d^2}{dy^2} + 3h(y)\frac{d}{dy} + \left(-1 + f^2(y) + s^2(y)\right)\right) s(y) = 0 \tag{3.3}$$

where h is the M-rescaled Hubble Parameter, $h := (\frac{d}{dy}a)/a$, with

$$\begin{aligned} h^2(y) = \frac{2\pi}{3}\frac{M^2}{\lambda M_{pl}^2} &\left[\dot{f}^2(y) + 2\dot{s}^2(y)\right. \\ &\left.+(1 - s^2(y))^2 + 2f^2(y)s^2(y)\right]. \end{aligned} \tag{3.4}$$

Important qualitative insights then are possible without any detailed calculation: Both of the classical fields, $f = \frac{\phi}{\phi_c}$ and $s = \frac{\sigma}{\sigma_0}$, effectively oscillate with the same frequency, approximately given by $\omega = \sqrt{2}$, around the supersymmetric minimum, $\phi = 0, \sigma = M/\sqrt{\lambda}$. For energetic reasons the phase difference amounts to $\pi/2$. Coming from the critical point ($\phi = M/g, \sigma = 0$), depending on the initial values the dynamics in the $\phi - \sigma$-plane should tend to be nearly one dimensional. Since in our case the initial time derivates of the fields should be very small slow roll values the trajectory should be near to a straight line. Another interesting point, obvious without any calculation, is that the dynamics of the post-inflationary system, measured in its natural time and lengthscale M^{-1}, only depends on the scale of the phase transition, $M/\sqrt{\lambda}$. As the M-rescaled classical fields f and s and their derivatives, take maximal values of $\mathcal{O}(1)$ effectively, this scale has influence on the damping but not on other aspects of the background oscillations: the smaller the scale is, the longer the system swings correponding to its natural timescale M^{-1}. Thus, qualitative results of one breaking scale will also appear at another breaking scale.

This discussion is confirmed by our numerical integration of the system of equations, as shown in Figs. 3.1 and 3.2. The parameters used in this and all of the following investigations are given by

$$\begin{aligned} \lambda &= 0.625 \times 10^{-3} \\ g^2 &= 0.125 \times 10^{-2} \\ M &= 0.350 \times 10^{15} \\ \frac{M}{\sqrt{\lambda}} &= 1.40 \times 10^{16}, \end{aligned} \tag{3.5}$$

which corresponds to a SSB near to the SUSY-GUT scale in a supersymmetric $SO(10)$ model in [6]. The solutions are essentially damped oscillations around the sypersymmetric minimum, with the predicted phase. The dynamics in the f-s-plane is practically one dimensonal and can be restricted to the straight line for very small initial values of the time derivatives of the fields. Anharmonic behaviour is caused by the interaction of the classical fields and by the fact, that the Higgs field has a negative mass contribution from the tachyonic mass.

Our result for the supersymmetric case ($g^2 = 2\lambda$) is quite different to the case with $g^2 = \lambda$, which was considered in [4], where the classical fields oscillate in a rather chaotic way. But it essentially equals the result in [7].

Now we consider the comoving energy densities after the phase transition:

$$R := \rho a^3, \tag{3.6}$$

$$R_\phi := \rho_\phi a^3, \tag{3.7}$$

$$R_\sigma := \rho_\sigma a^3. \tag{3.8}$$

Here approximately $\rho = M^4/4\lambda$ is the system's total energy density at the phase transition. As the Universe passes a phase of matter domination after inflation, such that $\rho \sim a^{-3}$, the product ρa^3 will be constant. This is what

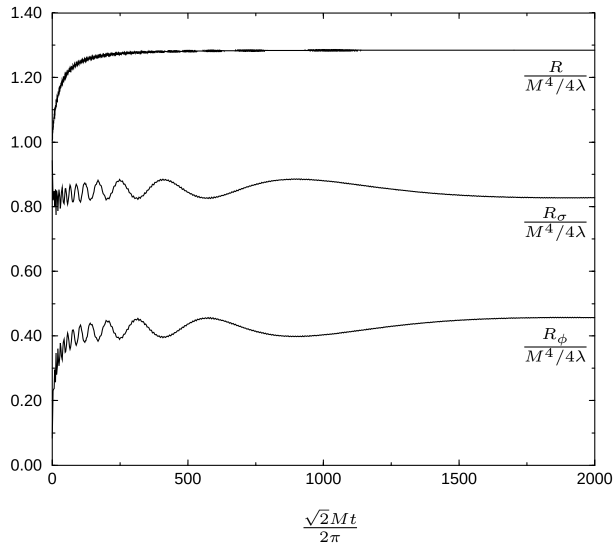

Fig. 3.3. Energy densities of the background

our numerical investigations show in Fig. 3.3. After the phase transition R is growing for a short time, up to the value of $1.24M^4/4\lambda$ and then remains practically constant. It takes the background about 120 oscillations until it swings in such a ,,harmonic" way, that the universe behaves matter dominated. Since we were interested in the effective distribution of energy between both of the swinging classical fields, in the calculation of R_ϕ and R_σ we integrated out the short time dynamics ($t \sim \mathcal{O}(M^{-1})$) by averaging over the the background oscillations. Here we respected the fact, that the frequency of the oscillations is slightly varying with time. The plot shows, how energy is exchanged between the classical fields - one reason for anharmonic behaviour in Fig. 3.1. After about 2000 oscillations the energy distribution between the oscillation of the inflaton and Higgs-condensate is approximately 1/3 : 2/3. Then the classical fields essentially are decoupled oscillators.

4 Parametric excitation of quantum fluctuations

Our investigation of the quantum fluctuations uses the M-rescaling introduced in Sect. 3. We studied the mode equations following from the Heisenberg expansion of a quantum field $\hat{\chi}$, which stands for the M-rescaled quantum fluctuations $\delta\hat{s}(y,\xi)$ and $\delta\hat{f}(y,\xi)$ around the classical fields $s(y)$ and $f(y)$:

$$\hat{\chi}(y,\mathbf{x}) = \frac{1}{(2\pi)^3 a^{3/2}(y)} \int d^3K (\hat{b}_K X_K(y) e^{-i\xi\mathbf{K}} + \hat{b}_K^+ X_K^*(y) e^{i\xi\mathbf{K}}), \tag{4.1}$$

where $\mathbf{K}$ is the M-rescaled commoving momentum. The operators $\hat{b}_K^{(+)}$ satisfy the commutation relation:

$$\left[\hat{b}_K, \hat{b}_{K'}^+\right] = (2\pi)^3 \, \delta(\mathbf{K} - \mathbf{K}'). \tag{4.2}$$

Then the occupation number is given by

$$n_K = \frac{\omega_K}{2}\Big(\frac{|\dot{X}_K|^2}{\omega_K^2} + |X_K|^2\Big) - \frac{1}{2}, \tag{4.3}$$

and the particle density reads:

$$n_\chi(y) = \int \frac{d^3K}{(2\pi a)^3} n_K. \tag{4.4}$$

Neglecting higher order terms in the quantum fluctuations, the mode equations read:

$$\left(\frac{d^2}{dy^2} + \left(\frac{K^2}{a^2} + 2s^2 - \frac{3}{2}\frac{\ddot{a}a + \dot{a}^2/2}{a^2}\right)\right)\delta f_K(y) = 0 \tag{4.5}$$

$$\left(\frac{d^2}{dy^2} + \left(\frac{K^2}{a^2} + 3s^2 + f^2 - 1 - \frac{3}{2}\frac{\ddot{a}a + \dot{a}^2/2}{a^2}\right)\right)\delta s_K(y) = 0, \tag{4.6}$$

As numerical calculations show, very soon $(3/2)(\ddot{a}a+\dot{a}^2/2)/(a^2) \sim 0$ is fullfilled within the numerical accuracy and this term can be neglected. This corresponds to a matter dominated universe, where typically $\ddot{a}a = -\dot{a}^2/2$. We performed a numerical integration of the full set of differential equations. The values of K, for which these investigations were done, were selected by a pre-investigation using the well known approximation by the Mathieu equation [11].

Before turning over to our numerical results, it is sensible first to recall the limitations of our ansatz. Clearly, it breaks down, when the system enters the non-linear regime. This latter stage can be investigate in the following way. As clearly pointed out in [12] the exponential growth of quantum fluctuation during preheating causes the system to be essentially semi-classical. Quantum fluctuations then behave essentially like random classical fields and can be treated e.g. on a lattice.

However, in the beginning the evolution of the system typically is in a linear regime, such that one can try to work with fourier methods.

As in the last section, we performed numerical calculations of the full system of differential equations mentioned above. We did not find any parametric amplification of the Higgs-fluctuations with $K \gtrsim 1/3$. For $K \lesssim 1/3$ there exists an amplification which can be described as a mixing of both effects: mainly parametric resonance and partly tachyonic mass. In the well known picture of the Mathieu equation,

$$\left(\frac{d^2}{dz^2} + A_k - 2q\cos(2z)\right) X_K = 0, \tag{4.7}$$

which is applicable restricting only to some few oscillations, this resonance starts inside the first resonance band ($1 - q \lesssim A_k \lesssim 1 + q$) with a resonance parameter $q \sim 0.9$, which corresponds to a strong narrow resonance regime. The amplification of these modes is essentially independent of the momentum and is given by the solution for

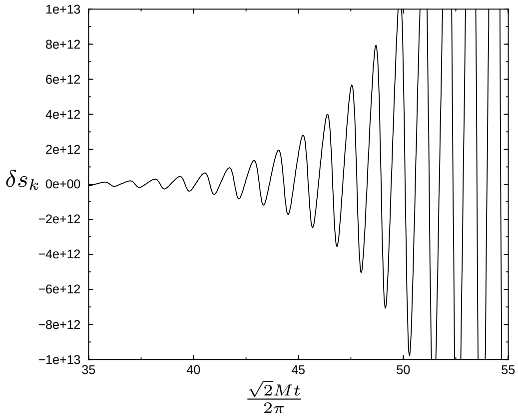

Fig. 4.1. mode function δs_k for $k = M/10$

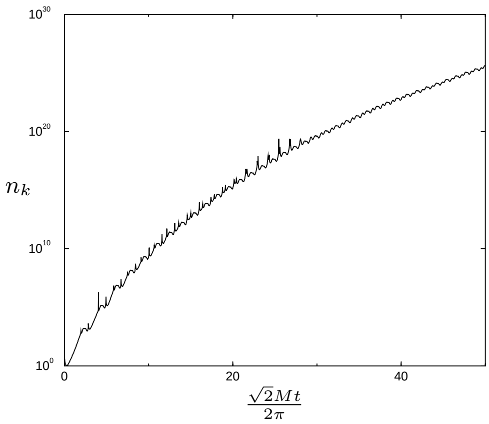

Fig. 4.2. resonance of the mode with $k = M/10$

$K = 0$. Their amplitude is not supressed compared to the zero mode and therefore the non-linear regime will be reached for each of these modes when $\sigma \sim \sigma_0$. For that reason our result only can tell us something about the region in momentum space, where particle production could occur, but not if it really does and how strong it is.

In the case of inflaton fluctuations the tachyonic mass is missing and the effect of the resonance is much weaker. There is no amplification of modes with $K \gtrsim 1/10$ and similiar to the above case and we observed a weak parametric amplification for modes with $K \lesssim 1/10$. But since the linearization of the equations of motion breaks down early, we also should not trust this result too much.

In order to learn more about the space/momentum dependence of the fields one could calculate the full (i.e. non simplified) mode equation for the fourier modes. Since the original equations of motions are non-linear, the mode equations then will be rather complicated integro-differential equations [13], that do not seem to be particuarly helpful for a numerical solution of the problem.

Our results show that after the SSB the inflaton and the Higgs condensate probably decay into Higgs and singlet particles, whose momenta are smaller than their mass.

Until now our investigations did not consider the effect of backreaction of the quantum fluctuations on the classical fields, which is very important in the case of a preheating with a very efficient particle production. As pointed out in [11], the main effect of the backreaction is to lower the resonance parameters in the picture of the Mathieu equation. This means, that appearing resonances will become weaker and will last for a shorter time.

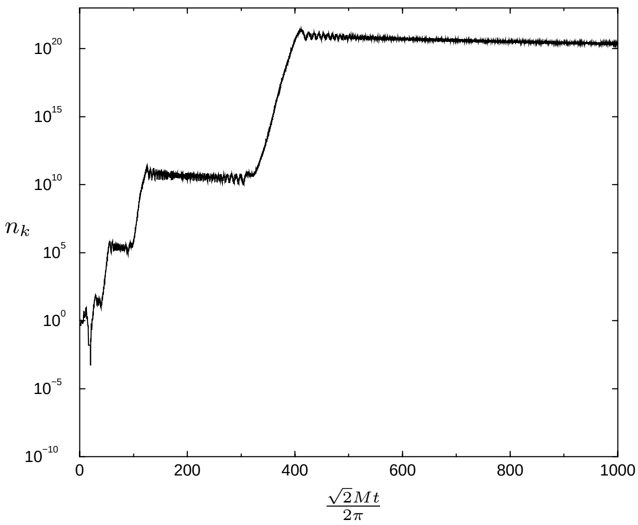

Fig. 5.1. parametric excitation of the $k = M$-mode of a scalar field $\hat{\chi}$

In our case soon the fluctuations become large and should terminate the resonance early. One could investigate this effect using the Hartree approximation, but a definite way to clarify this point are lattice calculations [12,13].

5 Parametric production of external fields

Now we try to estimate, if an external scalar quantum field $\hat{\chi}$, could be produced, coupling to the Higgs sector via

$$\frac{1}{2}h_1^2\phi^2\hat{\chi}^2 + \frac{1}{2}h_2^2\sigma^2\hat{\chi}^2. \tag{5.8}$$

Then the equations of motions will be modified in an obvious way. The mode equation for the modes $X_K(t)$ reads

$$\left(\frac{d^2}{dy^2} + \left(\frac{K^2}{a^2(t)} + \frac{h_1^2}{g^2}f^2 + \frac{h_2^2}{\lambda}s^2\right)\right) X_K(y,\xi) = 0, \tag{5.9}$$

where again $\ddot{a}a = -\dot{a}^2/2$ was assumed.

We used the same methods as before in our numerical investigation. We think, that they are able to tell us something about the possibility of particle production (even if we do not trust them quantitatively).

Our result is very similiar to the result in [4] for $g^2 = \lambda$. We found a strong resonance only for the case without coupling to the Higgs-condensate, $h_2 \simeq 0$ (see Fig. 5.1). But this situation is impossible, if we use the superpotential (2.1). The F-terms always lead to a coupling of χ to the higgs field, even if in the superpotential it only couples to the singlet. Thus, it seems to be very difficult, if not impossible, to construct a realistic coupling allowed by supersymmetry and leading to a bosonic preheating.

6 Consequences for left-right symmetric models

In left-right-symmetric models the phase transition that leads to the gauge symmetry of the Standard Model, necessarily lowers the rank of the Lie group. The only renormalizable superpotential with this feature has the structure

$$W = X\bar{C}C + \text{polynomial in } X, \qquad (6.1)$$

where X is a gauge singlet and $\bar{C}, C$ are appropriate spinor representations of the left-right-symmetric gauge group. It is an interesting question, if, using the superpotential (2.1), there could be a natural embedding of inflation into the context of a right-left-symmetric model such as $SO(10)$ [6] or $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{(B-L)}$, which leads to a senseful cosmological model.

In minimal $SO(10)$-models [14,15] for example C should be a 16-dimensional spinor representation, the scalar component of which aquires a GUT-scale vev in the $SU(5)$-singlet direction during the phase transition. Since the $SO(10)$ (more precisely: $Spin(10)$) is a simple and simply connected Lie group, breaking it down to the SM-gauge group will lead to unwanted monopoles. They appear in a first phase transition, when a 45 dimensional Higgs representation aquires a suitable vev to break $Spin(10)$ down to the left-right-symmetric group $G_{LR} := SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{(B-L)}$. The inflaton then dilutes the unwanted remnants and ends in a phase transition which breakes the remaining gauge symmetry tothe standard model group. This picture remains the same in a pure G_{LR}-Model, C being a $SU(2)_R \times U(1)_{(B-L)}$-Higgs-doublet. Our investigations indicate, that fourier methods are not applicable to such models, if one wants to do a quantitative investigation of the particle production after inflation. The reason is, that unlike to the case of chaotic inflation here the non-linear regime is reached early. But there are still some information we can extract by those methods. One is, that the creation of Higgs and singlet particles will be possible only within a “low momentum zone” with $k \ll M$. Probably then the decay of the two condensates will be incomplete during the linear stage. The other is, that other bosonic degrees of freedom cannot be created efficiently because of the structure of the supersymmetric lagrangian. We think the only approach, which will enable us to get quantitative result will be lattice calculations [12]. After having resolved this the next step would be to study the possibility of fermionic preheating for this superpotential, which recently was shown to be possibly very efficient [16]. Since we deal with a supersymmetric theory, there is no reason for neglecting the fermionic superpartners. Also a parametric creation of Majorana neutrinos could appear, which should be very interesting in the context of leptogenesis. In this context it also would be nessecary to rule out the creation of the helicity-1/2-gravitinos by non-perturbative effects, which recently was shown to be of possible danger for cosmological models which involve supersymmetry [17]. Particle production in left-right-symmetric models could also be caused by perturbative effects. Following [18] then still the production of super heavy matter could be possible. But, as we work with two condensates, both of them would have to decay very efficiently, in order to be not in conflict with standard cosmology. Considering the gravitino constraint on the reheating temperature of the Universe, this would mean strong limitations to possible couplings to other fields. Concluding, there is still some work to do, until the particle creation after inflation in this very simple supersymmetric model will be fully investigated.

Acknowledgements. We would like to thank S. J. Huber and P. John for useful discussions. We especially thank Igor Tkachev for his helpful coments on our work. This work was supported in part by the TMR network *Finite Temperature Phase Transitions in Particle Physics*, EU contract no. ERBFMRXCT97-0122.

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