Particle production from symmetry breaking after inflation

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Recent studies suggest that the process of symmetry breaking after inflation typically occurs very fast, within a single oscillation of the symmetry-breaking field, due to the spinodal growth of its long-wave modes, otherwise known as 'tachyonic preheating'. In this letter we show how this sudden transition from the false to the true vacuum can induce a significant production of particles, bosons and fermions, coupled to the symmetry-breaking field. We find that this new mechanism of particle production in the early Universe may have interesting consequences for the origin of dark matter and the generation of the observed baryon asymmetry through leptogenesis.

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Spontaneous symmetry breaking (SSB) is one of the basic ingredients of modern theories of elementary particles. It is usually assumed that SSB in Grand Unified and Electroweak theories took place in the early Universe through a thermal phase transition. However, it is also possible that some of these symmetries were broken at the end of a period of inflation [1], when the Universe had zero temperature and the negative mass term for the Higgs field appeared suddenly, i.e. in a time scale much shorter than the time required for SSB to occur. In this case, as was recently shown in Refs. [2,3], the process of symmetry breaking is extremely fast. The exponential growth of the Higgs quantum fluctuations is so efficient that SSB is typically completed within a single oscillation, while the field rolls down towards the minimum of its effective potential. This process, known as tachyonic preheating, leads to an almost instant conversion of the initial vacuum energy into classical waves of the scalar fields, in contrast with the process of 'parametric preheating', in which the inflaton field performs many oscillations before reheating the Universe [4].

In this letter we describe how this sudden transition from the false to the true vacuum can induce the non-adiabatic production of particles coupled to the Higgs. We also studied the consequences that this new process may have on the generation of the dark matter and the baryon asymmetry via leptogenesis. The phenomenon of particle production from symmetry breaking is analogous to the Schwinger mechanism [5], where the role of the external electric field pulse is played here by the time-dependent Higgs expectation value. It is also similar to the well known process of particle production by a time-dependent gravitational background [6], responsible for the observed anisotropies of the microwave background, as well as for Hawking radiation [7].

We will consider here a simplified model of SSB in which the Higgs instantly acquires a negative mass-squared term [2,3]. This 'quench' approximation corresponds to the limiting case of a hybrid inflation model [8] satisfying the so-called 'waterfall' condition. We therefore assume that the complex symmetry breaking field ϕ starts in the false vacuum at the top of its potential $V(\phi) = \lambda(|\phi|^2 - v^2)^2/4$, with zero mean, $\langle\phi\rangle = 0$, and initial conditions given by vacuum quantum fluctuations in de Sitter space. Other fields, scalars χ and fermions ψ, couple to the Higgs field with the usual scalar $g^2|\phi|^2\chi^2$ and Yukawa $h\phi\bar{\psi}\psi$ interactions. As we will show later, the backreaction of these fields on the Higgs evolution is negligible. Therefore, we can first solve the process of SSB and then use the evolution of the Higgs vacuum expectation value (vev) to study particle production.

The dynamics of symmetry breaking for different potentials in Minkowski space has been studied in detail in Refs. [2,3]. Here we only summarize the main results needed for our analysis. At the initial stages of SSB, when re-scattering effects are still unimportant, the Higgs modes follow the linear equation $\ddot{\phi}_k + (k^2 - m^2)\,\phi_k = 0$, where $m^2 = \lambda\, v^2$. With de Sitter initial conditions, all long-wavelength modes within the horizon $(H < k < m)$ grow exponentially $\phi_k(t) = \phi_k(0)\exp(t\sqrt{m^2-k^2})$, driving the fast growth of their occupation numbers [2], while modes with $k > m$ oscillate with constant amplitude. The exponential growth continues until the long-wave modes reach a value for which the effective Higgs mass becomes positive, i.e. when $\langle|\phi|^2\rangle \geq v^2/\sqrt{3}$, and the symmetry is broken soon after.

We have chosen $n_k + \frac{1}{2} = |\phi_k(t)\dot{\phi}_k(t)|$ as a proper definition for the occupation numbers of the Higgs tachyonic modes. This expression does not require an *a priori* definition of a mode frequency ω_k, and matches smoothly the one used in [3] for positive frequencies. For the growing modes we have

$$n_k + \frac{1}{2} = \frac{1}{2}\Big|e^{2t\sqrt{m^2-k^2}}\Big| \approx \frac{1}{2}\,e^{2mt}\,e^{-\frac{k^2}{2k_*^2}}\,. \qquad (1)$$

The occupation numbers of long-wavelength modes become exponentially large very quickly, although n_k drops abruptly for $k > k_* = m\,(2mt)^{-1/2}$, which gives a natural cutoff for the problem. Note that, while the dynamics conserve $\langle\phi\rangle = 0$, the Higgs dispersion grows in time as

$$\langle|\phi|^2\rangle = \frac{H^2}{8\pi^2}\int_H^m \frac{dk^2}{m^2}\, e^{2t\sqrt{m^2-k^2}} \sim \frac{H^2}{8\pi^2}\, e^{2mt}\,, \qquad (2)$$

where this expression has been regularized, as described below. The time it takes for the system to break the symmetry, i.e. when $\langle|\phi|^2(t_*)\rangle \simeq v^2$, can then be estimated as $mt_* \simeq \log(32\pi^2 m^2/\lambda H^2)^{1/2}$, which depends only logarithmically on the initial conditions and the Higgs self-coupling λ. For typical values, $\lambda = 10^{-2}$ and $v = 10^{-2} M_{\rm P}$, the symmetry is broken within $mt_* \sim 8$, and the typical cut-off frequency becomes $k_* \sim m/2$. This means that, by that time, the occupation numbers of modes with $k < k_*$ is exponentially large, $n_k \approx \frac{1}{2} e^{2mt_*} \simeq 16\pi^2 m^2/\lambda H^2 \sim 10^8$. These large occupation numbers allow us to treat these modes as semiclassical waves and match the solutions of the linear equations with the fully non-linear numerical lattice simulations [9]. The non-linear dynamics is studied by solving the real time evolution equations of classical fields, using a modified version of the lattice simulation program LATTICEEASY of Felder and Tkachev [10]. We start with initial fluctuations described by a Gaussian random field with zero mean, $\langle\phi_k\rangle = 0$, and regularized dispersion $\langle\phi_k^2\rangle_{\rm reg} \equiv \langle\phi_k^2\rangle - 1/2\omega_k = H^2/2\omega_k^3$. This prescription amounts to substituting quantum averages by ensemble averages. Note that in regularizing the dispersion we ensure that the physical masses and energies are not ultraviolet divergent.

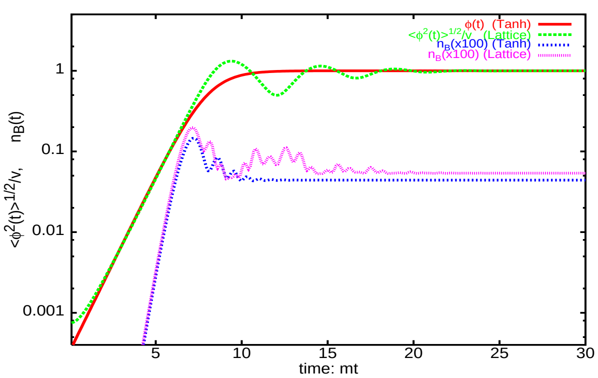

FIG. 1. The time evolution of the vacuum expectation value $\langle\phi^2(t)\rangle^{1/2}/v$, as compared with the approximate solution (3) with $mt_* = 8$. We have used a lattice of size N=128 and length L=100π, which gives $k_{\rm min} = 2H = 0.02m$ and $k_{\rm max} = 2.22m$. We also show the evolution of the number density, $n_{\rm B}(t)$, in units of $100\, m^{-3}$, for bosonic particles coupled to the Higgs with $g = 0.5$.

We can estimate the gradient energy density of the Higgs field at the time of symmetry breaking as $\langle(\nabla\phi)^2\rangle \sim k_*^2 v^2 \sim m^2 v^2/4 = V_0$, implying that a large fraction of the potential vacuum energy density at the phase transition has gone into the gradient energy of the Higgs field, not into its kinetic energy, therefore damping the subsequent oscillations around the true vacuum. Numerical lattice simulations show that, indeed, the picture described above is correct. The exponential growth of the Higgs vev can be written to very good approximation by

$$\phi(t) \equiv \langle|\phi|^2(t)\rangle^{1/2} = \frac{v}{2}\Big(1 + \tanh\frac{m(t-t_*)}{2}\Big)\,, \qquad (3)$$

which ignores the strongly damped oscillations after symmetry breaking [2]. We have checked that these low-amplitude oscillations do not contribute to the non-adiabatic production of particles, see Fig. 1. That is, parametric preheating is inefficient after symmetry breaking, a result anticipated in Ref. [11] for the case of hybrid inflation. Note that we will use the Higgs vev (3) as a homogeneous background field, while it is actually a coherent sum of tachyonic modes with different frequencies. Fortunately, the fact that all modes with $k < k_*$ grow essentially at the same speed implies that the correlation length of the phase transition is of order $\xi \sim k_*^{-1}$ during SSB and even larger after it. As a consequence, the Higgs vev will appear homogeneous on scales $m^{-1} < l < H^{-1}$. Therefore, those particles that couple to the Higgs will feel a homogeneous background field that grows exponentially during SSB and drives their effective mass towards their true vacuum value, creating particles in the process.

Let us now compute the production of bosons and fermions coupled to the Higgs using the formalism of quantum fields in strong backgrounds [12,13]. We can write the mode equations for these fields in terms of rescaled ones, $X_k(t) = a^{3/2}\chi_k$ and $\Psi(t) = a^{3/2}\psi$, as

$$\partial_t^2 X_k + \Big(k^2 + m_{\rm B}^2(t) a^2(t)\Big) X_k = 0\,, \qquad (4)$$

$$\Big(i\gamma^\mu\partial_\mu - m_{\rm F}(t) a(t)\Big)\Psi = 0\,, \qquad (5)$$

where both the mass $m(t)$ and the scale factor $a(t)$ depend on time. In practice, for most hybrid inflation models (for which the quench approximation described here is valid), the rate of expansion is typically much smaller than the mass scales of both the particles and the symmetry-breaking field, and therefore we can take the scale factor to be constant ($a = 1$) during SSB. We will only consider here the non-adiabatic production of particles due to the change of vacuum as it induces a sudden change in the inertia (masses) of bosons $m_{\rm B}^2(t) = g^2\langle|\phi|^2\rangle$ and fermions $m_{\rm F}(t) = h\langle|\phi|^2\rangle^{1/2}$, through the Higgs mechanism.

We have solved the mode equations (4) and (5) both numerically, with the Higgs vev computed by LATTICEEASY, and analytically, within the approximation (3), in terms of hypergeometric functions. We can thus obtain the number density of created particles as seen by a future observer in the true vacuum, $n_X = (g_s/2\pi^2 a^3)\int dk\, k^2\,|\beta_k|^2$, where β_k are the Bogoliubov coefficients that relate the *in* ($t \to -\infty$) and *out* ($t \to +\infty$) mode functions; k is the comoving momentum, and we have summed over spin indices ($g_s = 1$ for scalars, 2 for spinors). In the case of charged fields, n_X gives the number of particles, equal to that of antiparticles (i.e. $g_s = 2$ for complex scalars, 2 for Majorana fermions, 4 for Dirac fermions).

Let us consider first the production of bosons. The mode functions $X_k(t)$ of the scalar field are solutions of the oscillator equation (4) with time-dependent frequency $\omega_k^2 = k^2 + m_X^2$, and initial vacuum conditions, $X_k(0) = 1/\sqrt{2\omega_k}$ and $X_k'(0) = -i\omega_k X_k$. In this case, the Bogoliubov coefficient at $t \to \infty$ can be written as $n_k^{\rm B} = |\beta_k|^2 = (|X_k'|^2 + \omega_k^2 |X_k|^2 - \omega_k)/2\,\omega_k$. Substituting the exact solutions of (4) in terms of hypergeometric functions [12], we find

$$n_k^{\rm B} = \frac{\cos[\pi\sqrt{1-4\alpha^2}] + \cosh[2\pi(\omega_+ - \omega_-)/m]}{\sinh[2\pi\omega_-/m]\,\sinh[2\pi\omega_+/m]}\,, \qquad (6)$$

where $\omega_\pm$ are the *in/out* asymptotic frequencies, $\omega_-(k) = k$ and $\omega_+(k) = \sqrt{k^2 + \alpha^2 m^2}$, while $\alpha^2 \equiv g^2/\lambda$.

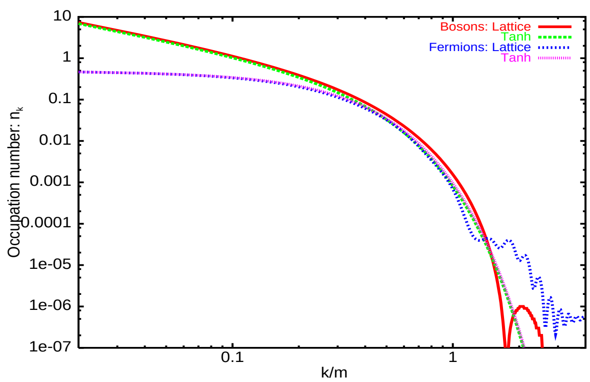

FIG. 2. The spectrum of occupation numbers for both bosons and fermions in the (asymptotic) true vacuum, using the lattice results for the Higgs vev, as compared with the analytical formulae (6) and (7). The parameters chosen here are $\lambda = 0.01$, $g = 0.5$ and $h = 0.5$. The tiny peaks at large momenta correspond to small resonances due to the strongly damped Higgs oscillations after SSB.

A similar analysis can be done in the case of fermions, where the first-order Dirac equation (5) for the two spinor components can be written as an oscillator equation with complex frequency [12], $X_k'' + (k^2 + m_X^2 - i m_X')\,X_k = 0$. Given the initial vacuum conditions, $X_k(0) = (1 - m_X/\omega_k)^{1/2}$ and $X_k'(0) = -i\omega_k X_k$, we can write the occupation number as $n_k^{\rm F} = |\beta_k|^2 = (\omega_k - m_X - {\rm Im} X_k X_k'^*)/2\,\omega_k$. The exact solutions of the fermionic equations, in terms of Hypergeometric functions [12], with $\alpha = h/\sqrt{\lambda}$, give

$$n_k^{\rm F} = \frac{\cosh[2\pi\alpha] - \cosh[2\pi(\omega_+ - \omega_-)/m]}{2\sinh[2\pi\omega_-/m]\,\sinh[2\pi\omega_+/m]}\,. \qquad (7)$$

The final occupation numbers obatined with the numerical solutions of Eqs. (4) and (5), using the full non-linear lattice solution of the Higgs background field, agree very well with the analytical formulae (6) and (7), see Fig. 2, except at large momenta, where resonances due to the strongly damped Higgs oscillations may appear, but are not significant for the total particle production.

The occupation numbers (6) and (7) are not thermally distributed. However, writing an effective temperature as a function of momentum as $T(k) = \omega_k/\log(n_k^{-1} \pm 1)$, one finds that, in the limit of large physical momenta k/a, the effective temperature observed by an asymptotic (future) observer is given by $T_{\rm eff} = m/4\pi$, for both bosons and fermions. This result is not surprising, since it corresponds to the gravitational analogue of a Rindler Universe with an acceleration $\kappa = m/2$, see Ref. [13], as would be expected from the exponential growth of the Higgs (3).

We will now compute the ratio of energy densities of the particles produced at the end of symmetry breaking to the initial false vacuum energy density, $\rho_0 = m^4/4\lambda$:

$$\frac{\rho_X}{\rho_0} = \frac{2g_s\lambda}{\pi^2}\int d\bar{k}\,\bar{k}^2\,n_{\bar{k}}(\alpha)\,\omega_{\bar{k}}(\alpha)\,, \qquad (8)$$

where $\bar{k} = k/m$. We can obtain a fit to the final energy density of bosons and fermions produced during symmetry breaking as

$$\frac{\rho_{\rm B}}{\rho_0} \simeq 2\cdot 10^{-3}\,g_s\lambda\,f(\alpha, 1.3)\,, \qquad (9)$$

$$\frac{\rho_{\rm F}}{\rho_0} \simeq 1.5\cdot 10^{-3}\,g_s\lambda\,f(\alpha, 0.8)\,, \qquad (10)$$

where $f(\alpha,\gamma) \equiv \sqrt{\alpha^2 + \gamma^2} - \gamma$. In the case of a very large coupling, $g^2, h^2 \gg \lambda$, the produced particles are non-relativistic, $\omega_k \approx m_X$, and their energy density is given by $\rho_{\rm X} = m_{\rm X} n_{\rm X}$, where X denotes either bosons or fermions. Note that, unless the couplings are unnaturally large, the fractional energy density in bosons (9) and fermions (10) is always small, so we do not expect an important backreaction on the evolution of the Higgs condensate as the symmetry is broken. Moreover, contrary to the case of Ref. [14], in which particles are produced long after symmetry breaking from non-linear rescattering, our mechanism of particle production from symmetry breaking gives an upper limit to the occupation numbers of bosons produced in the range $H < k < m$, even for arbitrarily large coupling g, which is of order $n_k \lesssim 10$. This prevents us from using LATTICEEASY to compute their energy density and backreaction.

We would now like to explore the cosmological consequences that this production of particles may have for the evolution of the Universe. First of all, these particles, either bosons of fermions, can be copiously produced if the self-coupling of the Higgs, and thus its mass, is large, driving a very sharp growth of the vev towards the true vacuum. The production can also be significant if the scalar and Yukawa couplings are large, see (9) and (10). In that case, the particles are non-relativistic and out of equilibrium at the end of symmetry breaking. Their energy density decays like matter, $\rho \sim a^{-3}$, while the energy density in Higgs particles is dominated by the gradient term, with a radiation equation of state. Assuming that the produced particles decay into stable relics, with branching ratio r, these relics could contribute to the present dark matter of the Universe. The ratio of their energy density today to that of radiation can be estimated as

$$\frac{\Omega_{\rm DM} h^2}{\Omega_\gamma h^2} = \frac{\rho_{\rm DM}}{\rho_0}\Big(\frac{H_{\rm end}}{\Gamma_{N_1}}\Big)^{1/2}\frac{T_{\rm rh}}{T_0} \approx 2\cdot 10^{11}\,\frac{r\rho_X}{\rho_0}\,\frac{v}{\rm GeV}\,, \tag{11}$$

where $H_{\rm end}$ is the rate of expansion at the end of inflation; $T_{\rm rh}$ is the reheating temperature, computed in terms of the perturbative Higgs decay rate [15], $T_{\rm rh} \simeq 0.12\sqrt{\Gamma_\phi\, M_{\rm P}}$, and $T_0 \simeq 3$ K is the temperature of the Universe today. Taking the present bound on non-relativistic dark matter as $\Omega_{\rm DM} = 0.3 \pm 0.1$ at 90% c.l., and the rate of expansion $H_0 = 72 \pm 7$ km/s/Mpc, we constrain the density of stable X particles produced at symmetry breaking as

$$\frac{\rho_X}{\rho_0} \approx 3\cdot 10^{-3} g_s h\sqrt{\lambda} < \frac{2.4\cdot 10^{-22}}{r}\,\Big(\frac{10^{14}\,{\rm GeV}}{v}\Big)\,. \tag{12}$$

One intriguing possibility is that heavy supersymmetric particles with $h > 1$ may have been produced at the EW phase transition ($v = 246$ GeV) through this non-thermal mechanism, and later decayed with very small branching ratios, $r < 4\cdot 10^{-8}$, into the neutralino (LSP). These could be responsible today for a large fraction of the observed dark matter; see Ref. [16] for an alternative scenario within preheating. Another possibility is that those particles that are produced very far from equilibrium eventually decay into relativistic particles that thermalize with the rest of the Higgs decays, then redshift like radiation and do not contribute to the present dark matter; or they decay much later, providing a source for ultra high energy cosmic rays [17].

Furthermore, we can take advantage of such a population of non-equilibrium fermions for proposing a new scenario of leptogenesis. Suppose that the symmetry breaking occurs in a GUT theory with $(B-L)$-violating interactions, like SO(10) or $\mathrm{SU}(2)_L\times\mathrm{SU}(2)_R$, and that the corresponding Higgs copiously produces out of equilibrium right-handed (RH) neutrinos. These massive neutrinos decay into leptons and SM Higgses with a leptonic asymmetry [18] that later gets converted into the baryon asymmetry of the Universe via sphaleron transitions [19]. We will briefly describe here the outline of constraints and leave for a future publication the details of the model of leptogenesis.

We will assume, as usual [20], that the lightest RH neutrino N_1, with mass M_1, is responsible for leptogenesis. The other two are too heavy to be produced at symmetry breaking, or they have already decayed by the time N_1 decays. Assuming that the masses of the SM light neutrinos are generated by the see-saw mechanism, $m_{\nu_i} = (m_D m_D^\dagger)_{ii}/M_i$, the decay rate of the RH neutrino can be written as

$$\Gamma_{N_1} = 6.6\cdot 10^5\,{\rm GeV}\,\Big(\frac{M_1}{10^{13}\,{\rm GeV}}\Big)^2\Big(\frac{m_{\nu_1}}{10^{-5}\,{\rm eV}}\Big)\,. \tag{13}$$

For definiteness, we will consider a GUT SB with a vev $v = 10^{14}$ GeV and a Higgs self-coupling $\lambda = 10^{-2}$, while the RH neutrino acquires its mass through the GUT Higgs mechanism, $M_1 = hv = 10^{13}$ GeV. Using the above formulae, we can estimate its number density and fractional energy density at symmetry breaking to be $n_{N_1} \simeq 2\cdot 10^{-4} m^3$ and $\rho_{N_1}/\rho_0 \simeq 1.44\cdot 10^{-5}$.

In order for leptogenesis to occur, we will have to satisfy a series of constraints [20,21], and check whether the model is consistent. First of all, the N_1 lifetime should be greater than the time of symmetry breaking t_*. In the model we are considering we find $mt_* \sim 16$, so we should satisfy $\Gamma_{N_1} < m/16 = 6.5\cdot 10^{11}$ GeV, or

$$m_{\nu_1} < 10\ {\rm eV}\Big(\frac{10^{13}\,{\rm GeV}}{M_1}\Big)^2\,. \tag{14}$$

We also have to ensure that the backreaction of the produced RH neutrinos on the Higgs field is negligible during symmetry breaking; that is, we should impose the constraint that the annihilation rate of N_1 into Higgses, $\Gamma_{\rm ann} = n_{N_1}\sigma_{\rm ann}$ be smaller than t_*^{-1}, where $\sigma_{\rm ann} = h^4/16\pi^2 M_1^2$. This gives a constraint on the couplings, $h\sqrt{\lambda} < 200$, or $\rho_{N_1} < 0.3\rho_0$, which is indeed satisfied. Moreover, this annihilation rate should be smaller than the decay rate, $\Gamma_{\rm ann} < \Gamma_{N_1}$, otherwise no RH neutrinos would be left to produce the lepton asymmetry. This imposes the constraint $m_{\nu_1} > 2\cdot 10^{-10}$ eV, which can be easily accommodated.

We can compute the time it takes for the N_1 neutrinos to decay and estimate the effective temperature that its products have at that time, denoting it by T_1. Since the equation of state after symmetry breaking is essentially that of radiation (see above) while the RH neutrinos behave like matter, the ratio of densities has an extra factor $a(t_1)/a(t_{\rm SB})$, coming from the expansion of the Universe. The energy density of RH neutrinos at t_1 can then be estimated as

$$\rho_{N_1}(t_1) = \Big(\frac{\rho_{N_1}}{\rho_\phi}\Big)_{\rm SB}\Big(\frac{H_{\rm end}}{\Gamma_{N_1}}\Big)^{1/2}\rho_\phi(t_1)\,, \tag{15}$$

where we have used the fact that at t_1, $H(t_1) = \Gamma_{N_1}$, giving $T_1 = 10^{11}\,{\rm GeV}\,(m_{\nu_1}/10^{-5}\,{\rm eV})^{5/8}$.

We also have to be careful that lepton-number-violating processes do not wash out the lepton asymmetry by the time the RH neutrinos decay, at t_1. This imposes the constraint $\Gamma_{\Delta L=2} = T_1^3(\sum m_{\nu_i}^2)/\pi^2 v_{\rm EW}^4 < \Gamma_{N_1}$, or

$$m_{\nu_1} < 1.2\ {\rm eV}\Big(\frac{3\cdot 10^{-3}{\rm eV}^2}{\sum m_{\nu_i}^2}\Big)^{8/7}\,. \tag{16}$$

Furthermore, the right-handed neutrinos should not decay before the final reheating of the Universe, $\Gamma_{N_1} > \Gamma_\phi = H(t_{\rm rh})$, giving the constraint,

$$m_{\nu_1} > 7 \cdot 10^{-9}\ \mathrm{eV} \left(\frac{T_{\rm rh}}{10^{10}\,\mathrm{GeV}}\right)^2 . \qquad (17)$$

Finally, we may ask how efficient is this mechanism for producing the required amount of baryons in the Universe, $n_{\rm B}/s = (4-7)\cdot 10^{-11}$, where s is the entropy density, and we use the fact that this ratio remains constant since reheating.

The baryon asymmetry is produced via sphaleron transitions that violate $(B+L)$ and convert a lepton asymmetry into a baryon asymmetry, $n_{\rm B} = (28/79) n_{\rm L}$ [19]. The entropy density at reheating can be computed in terms of the energy density, $s = (4/3)\rho_\phi/T_{\rm rh}$, and can be estimated with the help of (15) as

$$\begin{aligned}\frac{n_{\rm B}}{s} &= \frac{28}{79}\frac{3n_{\rm L}}{4n_{N_1}}\frac{T_{\rm rh}}{M_1}\left(\frac{\rho_{N_1}}{\rho_\phi}\right)_{\rm SB}\left(\frac{H_{\rm end}}{\Gamma_{N_1}}\right)^{1/2} \\ &= 2\cdot 10^{-6}\,\varepsilon\left(\frac{T_{\rm rh}}{10^{10}\,\mathrm{GeV}}\right)\left(\frac{m_{\nu_1}}{10^{-5}\,\mathrm{eV}}\right)^{1/2}, \qquad (18)\end{aligned}$$

where we have assumed that the lepton number density is directly related to the number density of RH neutrinos via $n_{\rm L} = \varepsilon\, n_{N_1}$. The baryon asymmetry of the Universe can thus be accommodated rather naturally with the lepton asymmetry parametrized by $\varepsilon \sim 3\epsilon^4/16\pi \sim 2\cdot 10^{-5}$, with $\epsilon \sim 1/13$ in a hierarchical model [20].

In conclusion, particle production at SSB from the exponential growth of the Higgs vev towards the true vacuum could be responsible for the present dark matter of the Universe, as well as the observed baryon asymmetry via leptogenesis. It would be very interesting to explore the possibility to observe these particle production effects in the 'little Big Bang' thought to occur in heavy ion collisions, where a supercooled chiral phase transition could have taken place [22].

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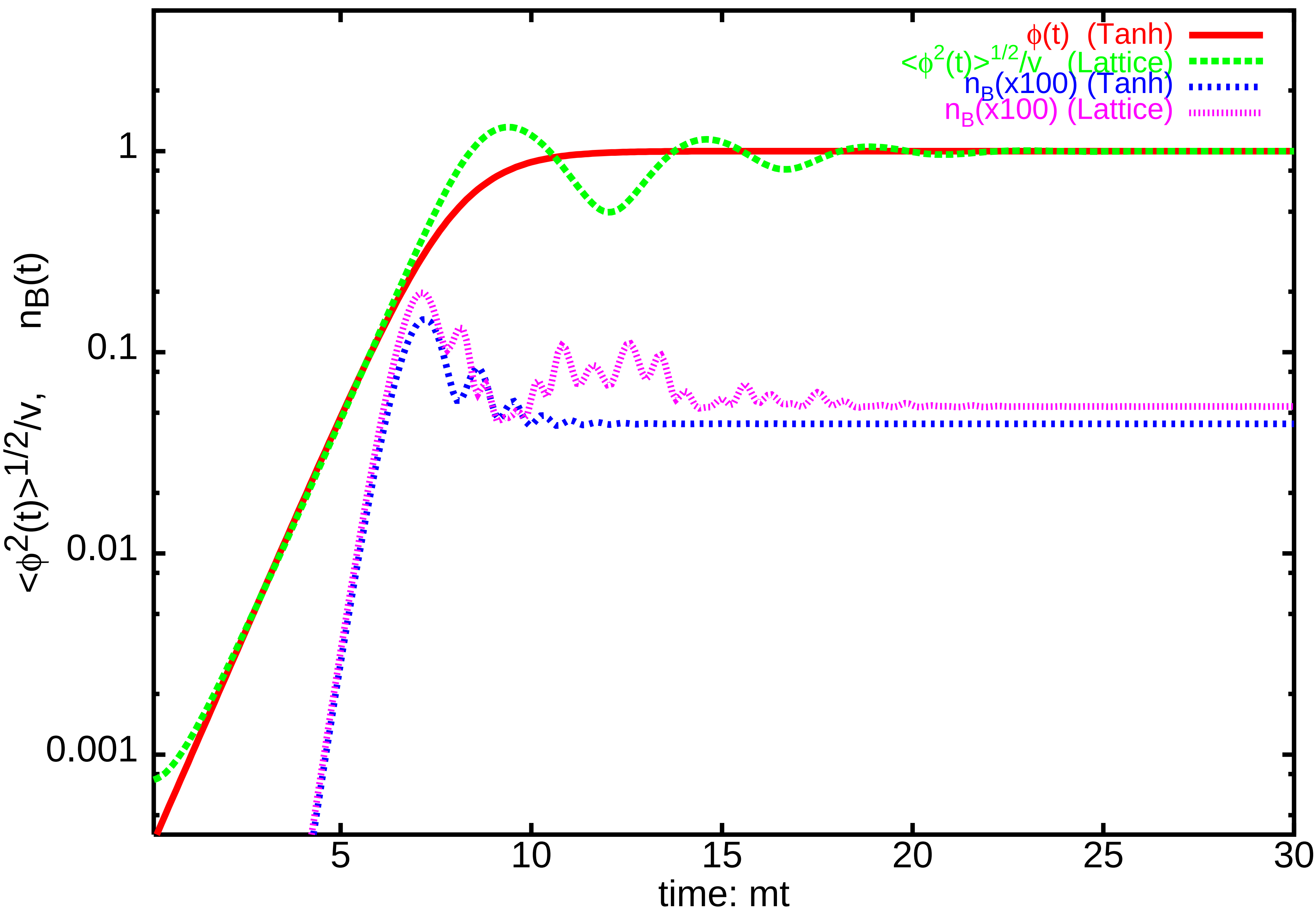
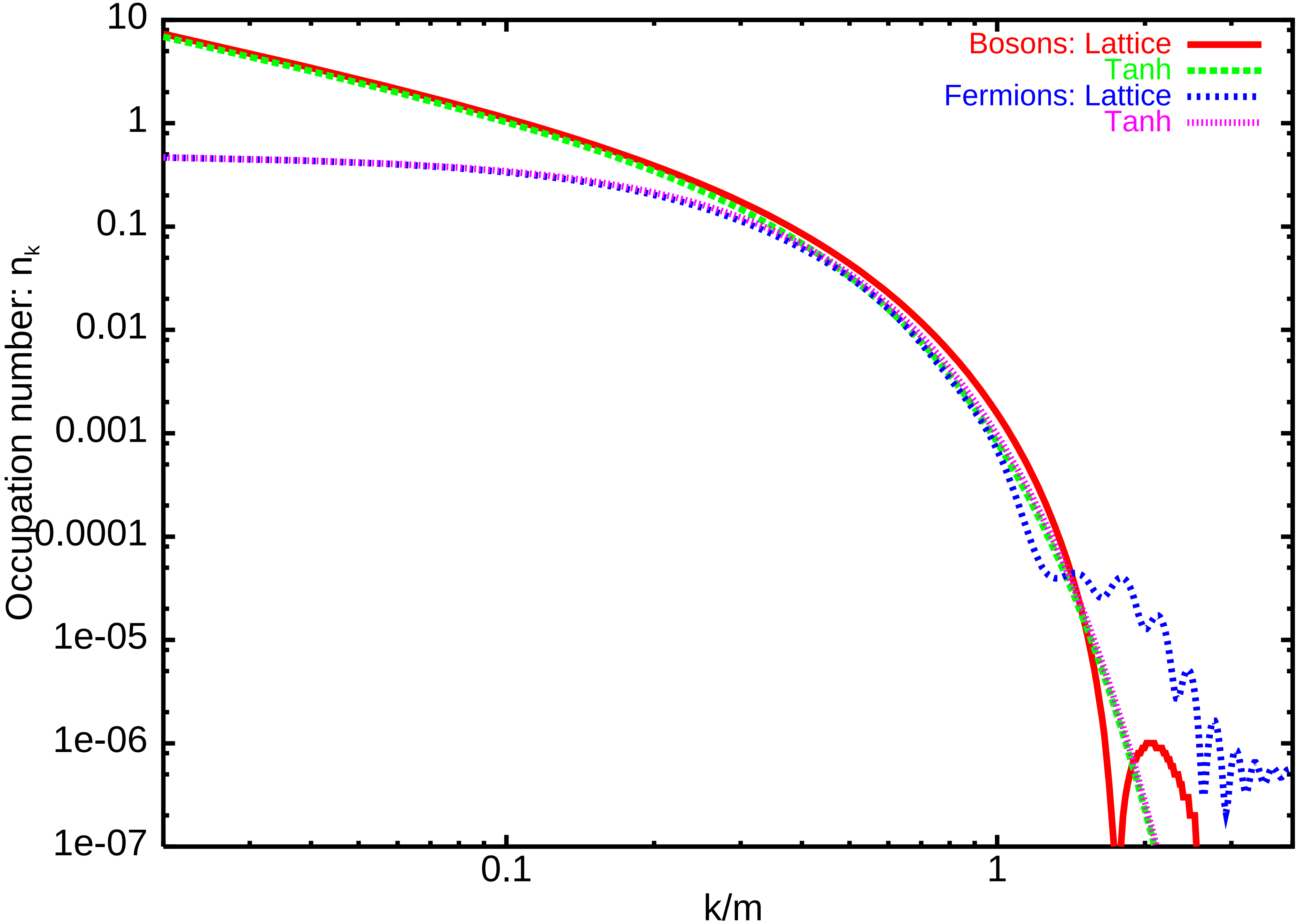


ϕ(t) (Tanh)
<ϕ^2(t)>^1/2/v (Lattice)
n_B(x100) (Tanh)
n_B(x100) (Lattice)
1
0.1
0.01
0.001
5
10
15
20
25
30
time: mt
<ϕ^2(t)>^1/2/v, n_B(t)



Bosons: Lattice
Tanh
Fermions: Lattice
Tanh
Occupation number: n_k
k/m
10
1
0.1
0.01
0.001
0.0001
1e-05
1e-06
1e-07
0.1
1