Thermal leptogenesis with triplet Higgs boson and mass varying neutrinos

Peihong Gu* and Xiao-Jun Bi

Institute of High Energy Physics, Chinese Academy of Sciences,
P.O. Box 918-4, Beijing 100039, People’s Republic of China

Abstract

We study the thermal leptogenesis in the scenario where the standard model is extended to include one $SU(2)_L$ triplet Higgs boson, in addition to three generations of the right-handed neutrinos. And in the model we introduce the coupling between the Quintessence and the right-handed neutrinos, the triplet Higgs boson, so that the light neutrino masses vary during the evolution of the universe. Assuming that the lepton number asymmetry is generated by the decays of the lightest right-handed neutrino $N_1$, we find the thermal leptogenesis can be characterized by four model independent parameters. In the case where the contribution of the triplet Higgs to the lepton asymmetry is dominant, we give the relation between the minimal $M_1$ and the absolute mass scale $\bar{m}$ of the light neutrinos, by solving the Boltzmann equations numerically. We will also show that with the varying neutrino masses, the reheating temperature can be lowered in comparison with the traditional thermal leptogenesis.

*Electronic address: guph@mail.ihep.ac.cn
The baryon number asymmetry of the Universe has been determined precisely\(^1\):

\[\eta_B \equiv \frac{n_B}{n_\gamma} = (6.3 \pm 0.3) \times 10^{-10},\]  

where \(n_B = n_b - n_\bar{b}\) and \(n_\gamma\) are the baryon and photon number densities, respectively. As the confirmation of the neutrino oscillations by several experiments\(^2, 3\), leptogenesis\(^4\) is now an attractive scenario to explain the observed baryon number asymmetry, where the lepton number asymmetry is first produced and then converted to the baryon number asymmetry via the \((B + L)\)-violating sphaleron interactions\(^5\).

The minimal thermal leptogenesis is quite economic and only requires three generations of the right-handed Majorana neutrinos beyond the standard model, which are also necessary to explain the small neutrino masses through the seesaw mechanism\(^6\). However, this scenario seems require too high reheating temperature which may conflict with the upper bound of the reheating temperature set by the gravitino problem\(^7\), and hierarchical neutrino spectrum with \(m_i \lesssim 0.12\) eV\(^8\). Furthermore, if the light neutrinos are degenerate as indicated by the experimental signal of neutrinoless double beta decay\(^9\), it is hard to imagine that the Dirac and the Majorana neutrino mass matrices, both naturally having hierarchical eigenvalues, conspire to produce the degenerate neutrino spectrum via the seesaw mechanism. The degenerate neutrino spectrum is naturally produced in the type II seesaw\(^10\) model, where a triplet Higgs boson is introduced, whose vacuum expectation value \((vev)\) gives the common neutrino mass scale and the type I seesaw produces the mass square differences required by the oscillation experiments.

A way to reconcile the minimal thermal leptogenesis and the gravitino problem is to consider that the neutrino masses are cosmological variable\(^11, 12\). In a recent work we studied the scenario\(^13\) where the interaction between the right-handed neutrinos and the Quintessence\(^14, 15, 16, 17\), a dynamical scalar field as a candidate for the dark energy\(^18\), which drives the accelerating of the Universe at the present time\(^19\), makes the masses of the right-handed neutrinos vary during the evolution of the universe. In this scenario, the reheating temperature is lowered and compatible with the limits set by the gravitino problem, and degenerate light neutrino spectrum is also permitted. However, the two conditions, the low reheating temperature and degenerate neutrino spectrum, can not be satisfied simultaneously yet.

In this paper, we study a scenario of the thermal leptogenesis where the light neutrinos
are degenerate, to explain the neutrinoless double beta decay, and at the same time the reheating temperature is low. This scenario is realized in the type II seesaw model with variable neutrino masses. The type II seesaw model including one $SU(2)_L$ triplet Higgs boson, in addition to three generations of the right-handed neutrinos, is a general scenario derived from the left-right symmetric models.

In this scenario, the lepton number asymmetry can be generated by the decays of the right-handed neutrinos and/or the $SU(2)_L$ triplet Higgs. Assuming a hierarchical right-handed neutrino spectrum, $M_1 \ll M_2, M_3$, and $M_1$ is also much lighter than the mass of the triplet Higgs $M_1 \ll M_\Delta$, the lepton number asymmetry comes mainly from the decays of the lightest right-handed neutrino $N_1$. We find the thermal leptogenesis can be characterized by four model independent parameters: the $CP$ asymmetry $\varepsilon_1$ of $N_1$ decays, the heavy neutrino mass $M_1$, the absolute mass scale $\tilde{m}$ of the light neutrinos, and the effective light neutrino mass $\tilde{m}_1$, which is a similar result as in the minimal thermal leptogenesis.

The Lagrangian relevant to leptogenesis reads:

$$-\mathcal{L} = \frac{1}{2} M_i \bar{N}_R^i N_R^i + M_\Delta^2 \Delta_L^\dagger \Delta_L + g_{ij}^\nu \bar{\psi}_L^i N_R^j \phi + g_{ij}^{\Delta} \bar{\psi}_L^i i\tau_2 \Delta_L^\dagger \psi_L^j - \mu \phi^T i\tau_2 \Delta_L \phi + h.c. \quad (2)$$

where $\psi_L = (\nu, l)^T$, $\phi = (\phi^0, \phi^-)^T$ are the lepton and the Higgs doublets, and

$$\Delta_L = \begin{pmatrix} \frac{1}{\sqrt{2}} \delta^{\dagger} & \delta^{\dagger \dagger} \\ \delta^0 & -\frac{1}{\sqrt{2}} \delta^{\dagger} \end{pmatrix}$$

is the Higgs triplet.

After the electroweak phase transition, the left-handed neutrino mass matrix can be written as

$$m_\nu = -g^{\nu *}_\nu \frac{1}{M} g^{\nu \dagger} v^2 + 2g^A_v v_L = m^I_\nu + m^{II}_\nu \quad (3)$$

where $m^I_\nu$ is the type I seesaw mass term, $m^{II}_\nu$ is the type II seesaw mass term, and $v = 174 GeV$, $v_L \simeq \mu^* v^2 / M_\Delta^2$ are the vevs of $\phi$ and $\Delta_L$, respectively.

The $CP$ asymmetry $\varepsilon_1$ is generated by the interference of the loop diagrams, shown in Fig. 1, with the tree diagram of $N_1$ decay. Besides the two same diagrams as in the minimal seesaw scenario, diagrams (a) and (b) in Fig. 1, there is an additional diagram (c) with the exchange of the Higgs triplet. We then have

$$\varepsilon_1 = \varepsilon^N_1 + \varepsilon^\Delta_1 \quad (4)$$
with $\varepsilon_1^N$ and $\varepsilon_1^\Delta$ the CP asymmetry of $N_1$ decays due to the exchange of the right-handed neutrinos and the Higgs triplet, respectively. For $M_1 \ll M_2, M_3, M_\Delta$, we have

$$\varepsilon_1^N \simeq \frac{3}{16\pi} \frac{M_1 \sum_{i,j} \text{Im}[g_{ij}^\dagger g_{ij}^\dagger (m_{\nu}^*)_{ij}]}{(g^{\nu} g^{\nu})_{11}},$$

(5)

$$\varepsilon_1^\Delta \simeq \frac{3}{16\pi} \frac{M_1 \sum_{i,j} \text{Im}[g_{ij}^\dagger g_{ij}^\dagger (m_{\nu}^{II*})_{ij}]}{(g^{\nu} g^{\nu})_{11}}.$$  

(6)

It is interesting to see that the triplet contribution will dominate the CP asymmetry $\varepsilon_1$ when it dominates the neutrino mass matrix $m_{\nu}$[20]. In this case, there is an upper bound on the asymmetry $\varepsilon_1$,

$$|\varepsilon_1| \simeq |\varepsilon_1^\Delta| \leq \frac{3M_1 m_3}{16\pi v^2} \simeq \varepsilon_1^{\text{max}}.$$  

(7)

In order to calculate the final baryon number asymmetry, we have calculated the washout effect by solving the Boltzmann equations numerically. All the relevant processes should be taken into account, which include $N_1$ decays and inverse-decays; the $\Delta L = 1$ scatterings mediated by exchanging doublet Higgs; and $\Delta L = 2$ scatterings mediated by exchanging the right-handed neutrinos and the triplet Higgs. By solving the Boltzmann equations, we can get the baryon-to photon ratio $\eta_B$.

In comparison with the minimal seesaw scenario[23, 27, 28, 29], there new $\Delta L = 2$ scattering processes with the exchanges of the triplet Higgs should be considered. We give the reduced cross sections $\tilde{\sigma}_N$ for the process $l \bar{\phi} \leftrightarrow \bar{l}\phi$, and $\tilde{\sigma}_{N,t}$ for the process $ll(\bar{l}\bar{l}) \leftrightarrow \phi\phi(\bar{\phi}\bar{\phi})$ as:

$$\tilde{\sigma}_{N(N,t)}(s) = \frac{1}{2\pi} \sum_i (g^{\nu} g^{\nu})_{ii}^2 f_{ii}^{N(N,t)}(x) + \sum_{i,j} \text{Re} [(g^{\nu} g^{\nu})_{ij}] f_{ij}^{N(N,t)}(x)$$

$$+ \sum_{i,j,k} \text{Re} [g_{ij}^\dagger g_{kj}^\dagger g_{kj}^\dagger g_{ij}^\dagger M_1] f_{ijk}^{N(N,t)}(x) + \sum_i (g^{\Delta} g^{\Delta})_{ii} \frac{\mu^2}{M_1^2} f_{ii}^{\Delta(\Delta,t)}(x),$$

(8)
with

\[ f_{ii}^N(x) = 1 + \frac{a_i}{D_i(x)} + \frac{xa_i}{2D_i^2(x)} - \frac{a_i}{x}[1 + \frac{x + a_i}{D_i(x)}]\log(1 + \frac{x}{a_i}), \]

\[ f_{ij}^N(x) = \frac{1}{2}\sqrt{a_ia_j}\left(\frac{D_i(x)}{D_j(x)} + \frac{D_j(x)}{D_i(x)}\right) + \left(1 + \frac{a_i}{a_j}\right)\log(1 + \frac{x}{a_i}) + \left(1 + \frac{a_j}{a_i}\right)\log(1 + \frac{x}{a_j})\right), \]

\[ f_{ij}^\Delta(x) = 12\left[\frac{1}{x}\log(1 + \frac{x}{y}) - \frac{1}{x + y}\right], \]

\[ f_{ijk}^N(x) = \frac{8\sqrt{a_k(x-a_k)}[(1-\frac{y}{x}\log(1+x)] + 4\sqrt{a_k}\frac{y}{x + a_k} + y\frac{x}{y}\log(1 + \frac{x}{y}) - (1 + \frac{a_k}{x})\log(1 + \frac{x}{a_k})\right), \]

\[ f_{ij}^{N,t}(x) = \frac{x}{x + a_i} + \frac{a_i}{x + 2a_i}\log(1 + \frac{x}{a_i}), \]

\[ f_{ij}^{\Delta,t}(x) = 6\left[\frac{x(x-y)}{(x-y)^2 + y\Delta}\log(1 + \frac{x}{a_i})\right], \]

\[ f_{ijk}^{\Delta,t}(x) = 6\sqrt{a_k}\frac{x-y}{(x-y)^2 + y\Delta}\log(1 + \frac{x}{a_k})\right). \]

Here \(x = \frac{s}{M_1^2}, a_i \equiv M_i^2/M_1^2\) and \(1/D_i(x) \equiv (x - a_i)/[(x - a_i)^2 + a_i C_i]\) is the off-shell part of the \(N_i\) propagator with \(C_i \equiv a_i(g^{\rho\nu}\gamma^\nu_i/8\pi)^2\), and \(C_D \equiv \Gamma^2/M_1^2\). Similar to Ref. [23], the reaction density \(\gamma_N + \gamma_{N,t}\) can be separated into two parts: the resonance contribution which is highly peaked around \(x = 1\) and the contribution comes from the region \(x \ll 1\) which corresponds to \(z \gg 1\)

\[ \gamma_N^{res} = \frac{M_1^5}{64\pi^3 v^2} \tilde{m}_1 \frac{1}{z} K_1(z), \]

\[ \gamma_N(z \gg 1) \simeq \gamma_{N,t}(z \gg 1) \simeq \frac{3M_1^6}{8\pi^5 v^4} \tilde{m}_2 \frac{1}{z^6}, \]

where \(\tilde{m}_1 \equiv (g^{\rho\nu}\gamma^\nu_i/8\pi)^2\) and \(\tilde{m}_2 \equiv m_1^2 + m_2^2 + m_3^2 = tr(m^I_\nu m_\nu) = tr((m^I_\nu + m^I_\nu)^I(m^I_\nu + m^I_\nu))\). Analysis for \(z < 1\) is also similar to Ref. [23]. The reaction densities \(\gamma_i\) is defined as:

\[ \gamma_i(z) = \frac{M_i^4}{64\pi^2 z} \int_{(m_2 + m_3^2)/m_1^2}^\infty dx \bar{\sigma}_i(x) \sqrt{x} K_1(z \sqrt{x}), \]

where \(m_a\) and \(m_b\) are the masses of the two particles in the initial state. Since \(\gamma_N + \gamma_{N,t}\) is not changed compared with the results of Ref. [23], the thermal leptogenesis can be still characterized by four parameters: \(\varepsilon_1, M_1, \tilde{m}, \tilde{m}_1\), even in the presence of the \(SU(2)_L\) triplet Higgs.
The washout effect mainly depends on the effective neutrino mass $\tilde{m}_1$. Since $m_1 \leq \tilde{m}_1 \lesssim m_3$ should not be satisfied when the triplet contribution is dominant, we can always adjust $\tilde{m}_1$ to avoid the large washout effect. Therefore the neutrino mass spectrum can be degenerate, even above the cosmological bound. The numerical result is shown in Fig. 2. For $\tilde{m} \simeq 0.051 \text{eV}$, the low limit of the neutrino mass scale from the oscillation experiments constraint, we get $M_1 \simeq 3.4 \times 10^8 \text{GeV}$. We can get $M_1 \simeq 2.7 \times 10^8 \text{GeV}$ for $\tilde{m} \simeq 1.0 \text{eV}$, which is the upper bound from the cosmological constraint $\sum_i m_i < 1.7 \text{eV}$.

We notice that these values of $M_1$ are only marginally consistent with the bound set by the gravitino problem. In order to solve the problem, we consider the light neutrino masses are varying during the evolution of the universe. We introduce a parameter $k$ which indicates the ratio of the light neutrino masses at the leptogenesis epoch and the present epoch. When solving the Boltzmann equations, the $M_1$, $\tilde{m}$, and $\tilde{m}_1$ should all take the values at the leptogenesis epoch. If $\tilde{m}$ takes the value at the present epoch, we should replace $\tilde{m}$ by $k\tilde{m}$ in the Boltzmann equations. By solving the Boltzmann equations numerically, we can see the reheating temperature are lowered with the increasing $k$. For $k = 10$, we can get $M_1 \simeq 3.1 \times 10^8 \text{GeV}$ for $\tilde{m} \simeq 0.051 \text{eV}$, and $M_1 \simeq 2.7 \times 10^7 \text{GeV}$ for $\tilde{m} \simeq 1.0 \text{eV}$. $M_1$ is lowered to $M_1 \simeq 3.1 \times 10^7 \text{GeV}$ with $\tilde{m} \simeq 0.051 \text{eV}$ and $M_1 \simeq 2.5 \times 10^6 \text{GeV}$ with $\tilde{m} \simeq 1.0 \text{eV}$ for $k = 100$. In this paper, the vales of $M_1$ are all at the leptogenesis epoch, and $\tilde{m}$ takes the value at the present epoch.

In this paper, we get the varying neutrino masses by introducing the interaction between the Quintessence and the right-handed neutrinos, the triplet Higgs in Eq.(2). Assume these interactions take simple forms as:

\[
M_i \rightarrow M_i(Q) = \tilde{M}_i e^{\frac{\beta Q}{M_{\text{pl}}}}, \tag{20}
\]

\[
M_\Delta \rightarrow M_\Delta(Q) = \tilde{M}_\Delta e^{\frac{\beta Q}{M_{\text{pl}}}}, \tag{21}
\]

where $\beta$ is a $O(1)$ coefficient. Then we get

\[
m_\nu \propto e^{-\frac{\beta Q_0}{M_{\text{pl}}}}, \tag{22}
\]

and

\[
k = e^{\frac{\beta Q_0 - Q_{\text{D}}}{M_{\text{pl}}}}. \tag{23}
\]

$Q_0$, $Q_{\text{D}}$ are the values of the Quintessence field at the present epoch and the leptogenesis epoch, respectively.
FIG. 2: The minimal $M_1$ as a function of $\bar{m}$ for $\eta_B = 5.4 \times 10^{-10}$. $M_1$ is the value of the right-handed neutrino mass at the leptogenesis epoch, and $\bar{m}$ is the absolute mass scale of the light neutrino at the present epoch. The curve with $k = 1$ stand for the case without the variation of the light neutrino masses.

For a numerical calculation of $k$, we consider a model of the Quintessence with the double exponential potential $^{33}$

$$V = V_0 (e^{\lambda Q} + e^{\alpha Q}) .$$

(24)

The equations of motion of the Quintessence, which for a flat universe, are given by,

$$H^2 = \frac{8\pi G}{3} (\rho_B + \frac{\dot{Q}^2}{2} + V(Q)),$$

(25)

$$\ddot{Q} + 3H \dot{Q} + \frac{dV(Q)}{dQ} = 0 ,$$

(26)

where $\rho_B$ represent the energy densities of the background fluid. This model has the tracking
FIG. 3: The evolution of $w_Q$ and $Q$ as a function of the temperature $T$ for the double exponential Quintessence model.

property for suitable parameters. Here, we choose $\lambda = 100 M_{pl}^{-1}$, $\alpha = -100 M_{pl}^{-1}$, the initial value of Quintessence field $Q_i = 1.374 M_{pl}$ and for the equation-of-state, which is defined as

$$w_Q = \frac{\dot{Q}^2 / 2 - V(Q)}{Q^2 / 2 + V(Q)} ,$$

the initial value is $w_{Q_i} = -1$. We obtain that $\Omega_{Q_0} \simeq 0.72$ and the present equation-of-state of the Quintessence is $w_{Q_0} \simeq -1$ which are consistent with the observational data. In Fig. 3, we show the evolution of $w_Q$ and $Q$ with the temperature $T$.

Taking into account the interaction with the right-handed neutrinos and the triplet Higgs, we get the equation of motion of the Quintessence as

$$\ddot{Q} + 3H \dot{Q} + \frac{dV(Q)}{dQ} + \frac{dV_I(Q)}{dQ} = 0 .$$

The source term in the equation above is given by

$$\frac{dV_I(Q)}{dQ} = \sum_i n_i \frac{dM_i}{dQ} \left< \frac{M_i}{E_i} \right> + n_\Delta \frac{dM_\Delta}{dQ} \left< \frac{M_\Delta}{E_\Delta} \right>$$

$$= \frac{\beta}{M_{pl} \pi^2} T \sum_i M_i^3 K_1(M_i/T) + \frac{3}{4} \frac{\beta}{M_{pl} \pi^2} T M_\Delta^3 K_1(M_\Delta/T) ,$$

8
where \( n_i \) and \( E_i \) are the number density and the energy of the right-handed neutrinos respectively, \( n_\Delta \) and \( E_\Delta \) belong to the triplet Higgs, \( \langle \rangle \) indicates thermal average, and \( K_1 \) is the modified Bessel function. For simplicity, we have taken the Maxwell-Boltzmann distribution of the right-handed neutrinos and the triplet Higgs in the last step of the equation.

We then solve the equation (28) numerically, assuming \( \overline{M}_3 = 10 \overline{M}_2 = 10^{4} \overline{M}_1 \) and \( \overline{M}_\Delta = 10^{8} \overline{M}_1 \). The numerical results are shown in Figs. 4, 5, 6, and 7, where we have taken the same definition of \( w_Q \) as in Eq.(27). In Fig. 4 and 5, we take \( \beta = -1.68, \overline{M}_1 = 3.1 \times 10^{9} \) GeV, and \( \beta = -3.35, \overline{M}_1 = 3.1 \times 10^{9} \) GeV, respectively, which give rise to \( Q_0 \simeq 0 \) and \( Q_D \simeq 1.374M_{\text{pl}} \). We then have \( M_1 \simeq 3.1 \times 10^{8} \text{GeV} \) for \( k = 10 \), and \( M_1 \simeq 3.1 \times 10^{7} \text{GeV} \) for \( k = 100 \), corresponding to the case we considered in the hierarchical neutrino spectrum with \( \bar{m} \simeq 0.051 \text{eV} \). In Fig. 6 and 7, we choose the parameters \( \beta = -1.68, \overline{M}_1 = 2.7 \times 10^{8} \text{GeV} \) and \( \beta = -3.35, \overline{M}_1 = 2.5 \times 10^{8} \text{GeV} \), respectively. We find the values of \( Q_0 \) and \( Q_D \) are almost the same as the above case. We then have \( M_1 \simeq 2.7 \times 10^{7} \text{GeV} \) for \( k = 10 \) and \( M_1 \simeq 2.5 \times 10^{6} \text{GeV} \) for \( k = 100 \), corresponding to the case that satisfies the cosmic limit \( \bar{m} \simeq 1.0 \text{eV} \) for the degenerate neutrinos.

Comparing Figs. 3 with Figs. 4, 5, 6 and 7, one can see that the interaction of the Quintessence with the right-handed neutrinos and the triplet Higgs, does change the equation-of-state of the Quintessence field, however, does not change the tracking properties of this model. Furthermore, the value of the Quintessence field \( Q \) changes very little in this model until \( T \sim 10^{4} \text{GeV} \) which satisfies our assumption for a constant \( k \) during the period of leptogenesis.

In summary, we study the thermal leptogenesis in the scenario where the standard model is extended to include one \( SU(2)_L \) triplet Higgs boson, in addition to three generations of the right-handed neutrinos. And in the model we introduce the coupling between the Quintessence and the right-handed neutrinos, the triplet Higgs boson, so that the light neutrino masses vary during the evolution of the universe. Assuming that the lepton number asymmetry is generated by the decays of the lightest right-handed neutrino \( N_1 \), and find the thermal leptogenesis can be characterized by four model independent parameters: \( \varepsilon_1, M_1, \bar{m}, \bar{m} \). With the dominant contribution of the triplet Higgs to the lepton asymmetry and the varying neutrino masses, we find the degenerate spectrum of the light neutrino masses and the lower reheating temperature can be get simultaneously, by solving the Boltz-
FIG. 4: The evolution of $w_Q$, $Q$ and $k$ as a function of the temperature $T$ for the double exponential Quintessence model including the coupling with the right-handed neutrinos and the triplet Higgs. We take $\beta = -1.68$ and $M_1 = 3.1 \times 10^9 GeV$.

mann equations numerically.

Acknowledgment: We thank Prof. Xinmin Zhang and Dr. Bo Feng for discussions. This work is supported in part by the National Natural Science Foundation of China under the Grand No. 90303004, 10105004.

[1] M. Tegmark et. al., astro-ph/0310723
[2] G. L. Fogli, E. Lisi, A. Marrone, D. Montanino, A. Palazzo, A. M. Rotunno, hep-ph/0310012
M. H. Ahn et. al., K2K Collaboration, Phys. Rev. Lett. 90 (2003) 041801; M. Shiozawa et. al.
FIG. 5: The evolution of $w_Q$, $Q$ and $k$ as a function of the temperature $T$ for the double exponential Quintessence model including the coupling with the right-handed neutrinos and the triplet Higgs. We take $\beta = -3.35$ and $\overline{M}_1 = 3.1 \times 10^9 GeV$.

[3] Q. R. Ahmad et al., SNO Collaboration, nucl-ex/0309004; K. Eguchi et al., KamLAND Collaboration, Phys. Rev. Lett. 90 (2003) 021802.

[4] M. Fukugita and T. Yanagida, Phys. Lett. B174, 45 (1986).

[5] N. S. Manton, Phys. Rev. D28, 2019 (1983); F. R. Klinkhamer and N. S. Manton, Phys. Rev. D30, 2212 (1984); V. A. Kuzmin, V. A. Rubakov and M. E. Shaposhnikov, Phys. Lett. B155, 36 (1985).

[6] M. Gell-Mann, P. Ramond and R. Slansky, Proceedings of the Supergravity Stony Brook Workshop, New York, 1979, eds. P. Van Nieuwenhuizen and D. Freedman (North-Holland,
FIG. 6: The evolution of $w_Q$, $Q$ and $k$ as a function of the temperature $T$ for the double exponential Quintessence model including the coupling with the right-handed neutrinos and the triplet Higgs. We take $\beta = -1.68$ and $M_1 = 2.7 \times 10^8 GeV$.

Amsterdam); T. Yanagida, Proceedings of the Workshop on Unified Theories and Baryon Number in the Universe, Tsukuba, Japan 1979 (eds. A. Sawada and A. Sugamoto, KEK Report No. 79-18, Tsukuba); R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. 44, 912 (1980); S. L. Glashow, *Caranse lectures*, (1979).

[7] S. Weinberg, Phys. Rev. Lett. 48 (1982) 1303; H. Pagels and J. R. Primack, Phys. Rev. Lett. 48 (1982) 223.

[8] W. Buchmüller, P. Di Bari and M. Plümacher, Nucl. Phys. B665 (2003) 445.

[9] H.V. Klapdor-Kleingrothaus, I.V. Krivosheina, A. Dietz, O. Chkvorets, Phys.Lett. B586 (2004) 198.
FIG. 7: The evolution of $w_Q$, $Q$ and $k$ as a function of the temperature $T$ for the double exponential Quintessence model including the coupling with the right-handed neutrinos and the triplet Higgs. We take $\beta = -3.35$ and $\overline{M}_1 = 2.5 \times 10^8 GeV$.

[10] G. Lazarides, Q. Shafi and C. Wetterich, Nucl. Phys. B181 (1981) 287; R. N. Mohapatra and G. Senjanović, Phys. Rev. D23 (1981) 165; C. Wetterich, Nucl. Phys. B187 (1981) 343.

[11] P. Gu, X. Wang and X. Zhang, Phys. Rev. D68 (2003) 087301.

[12] R. Fardon, A. E. Nelson, N. Weiner, astro-ph/0309800

[13] X. J. Bi, P. Gu, X. Wang and X. Zhang, Phys. Rev. D69 (2004) 113007.

[14] B. Ratra and P.J.E. Peebles, Phys. Rev. D 37, (1988) 3406.

[15] C. Wetterich, Nucl. Phys. B 302, (1988) 302.

[16] J.A. Frieman, C.T. Hill, A. Stebbins, and I. Waga, Phys. Rev. Lett. 75, (1995) 2077.

[17] I. Zlatev, L. Wang, and P.J. Steinhardt, Phys. Rev. Lett. 82, (1999) 896; P.J. Steinhardt, L.
Wang and I. Zlatev, Phys. Rev. D 59, (1999) 123504.

[18] M. S. Turner, astro-ph/0108103.

[19] S. Pelmutter et al., Astrophys. J. 483 (1997) 565.

[20] T. Hambye and G. Senjanović, Phys. Lett. B582 (2004) 73.

[21] J. C. Pati and A. Salam, Phys. Rev. D10 (1974) 275; R. N. Mohapatra and J. C. Pati, Phys. Rev. D11 (1975) 566, Phys. Rev. D11 (1975) 1502; G. Senjanović and R. N. Mohapatra, Phys. Rev. D12 (1975) 1502; G. Senjanović, Nucl. Phys. B153 (1979) 334.

[22] R. N. Mohapatra and X. Zhang, Phys. Rev. D46 (1992) 5331; P. J. O’Donnell and U. Sarkar, Phys. Rev. D49 (1994) 2118.

[23] W. Buchmüller, P. Di Bari and M. Plümeracher, Nucl. Phys. B643 (2002) 367.

[24] L. Covi, E. Roulet and F. Vissani, Phys. Lett. B384 (1996) 169; W. Buchmüller and M. Plümeracher, Phys. Lett. B431 (1998) 354; A. Pilaftisis, Int. J. Mod. A14 (1999) 1811.

[25] S. Antusch, S. F. King, hep-ph/0405093.

[26] S. Davidson and A. Ibarra, Phys. Lett. B535 (2002) 25.

[27] M. A. Luty, Phys. Rev. D45 (1992) 455.

[28] M. Plümeracher, Z. Phys. C74 (1997) 549.

[29] G. F. Giudice, A. Notari, M. Raidal, A. Riotto, A. Strumia, Nucl. Phys. B685 (2004) 89.

[30] E. Chun and S. Kang, Phys. Rev. D63 (2001) 097902

[31] W. Rodejohann, hep-ph/0403236.

[32] W. Buchmüller, P. Di Bari and M. Plümeracher, hep-ph/0401240.

[33] T. Barreiro, E. J. Copeland, N. J. Nunes, Phys. Rev. D 61 (2000) 127301.

[34] G. R. Farrar and P. J. E. Peebles, Astrophys.J. 604 (2004) 1.