Physics Letters B 685 (2010) 235–238

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Inflation with the right-handed sneutrino revisited

O. Efthimiou [a,*], K. Tamvakis [a,b]

[a] *Physics Department, University of Ioannina, 45110 Ioannina, Greece*
[b] *Physics Department, CERN, CH-1211, Geneva 23, Switzerland*

ARTICLE INFO

Article history:
Received 18 December 2009
Accepted 1 February 2010
Available online 2 February 2010
Editor: L. Alvarez-Gaumé

ABSTRACT

We consider an extension of νMSSM with an extra $U(1)$ that realizes D-term inflation driven by the right-handed sneutrino. Non-renormalizable terms in the Kähler potential and the superpotential are considered, the latter controlled by a suitable discrete R-symmetry. We find that, for sub-Planckian inflaton values, the predictions of inflationary parameters are compatible with observations, establishing the right-handed sneutrino driven inflation as a viable scenario.

The most dramatic development of the past decade in particle physics, namely the discovery of neutrino mass [1], requires the presence of a right-handed neutrino field, which, in the supersymmetric version of the Standard Model, is accompanied by its scalar partner, a *right-handed sneutrino*. A straightforward way to explain the smallness of the neutrino mass is to invoke the *see-saw mechanism* [2], in which the right-handed neutrino possesses a large mass in the range of 10^{11} to 10^{15} GeV. What is remarkable is that right-handed sneutrino fields can be related to inflation, thus, providing a direct connection between cosmological inflation and particle physics. Models where the right-handed sneutrino is the inflaton have been proposed, either in the chaotic inflation framework [3,4] or in the framework of D-term inflation [5]. Although the minimal sneutrino inflationary scenarios, based on chaotic inflation, result in density perturbations compatible with existing data, from the point of view of particle physics, the model should be embeddable to supergravity. Then, one has to address the flatness of the scalar potential, which, in general, is problematic due to supergravity corrections of F-terms (η-*problem*). This can be circumvented if one adopts the D-term inflation [6,7] framework, where the required vacuum energy that drives inflation is supplied by Fayet–Iliopoulos D-term of a $U(1)$ gauge factor. An alternative approach is to employ a specific form of the Kähler potential in order to maintain the flatness of the scalar potential (F-term hybrid inflation) but extra fields are required [8,9].

In the present Letter we reconsider inflation driven by the right-handed sneutrino in the D-term inflation framework. We consider an extension of the standard νMSSM that realizes neutrino masses through the see saw mechanism, with an extra $U(1)$ gauge factor, under which all standard fields are neutral apart from a pair of MSSM singlets $\phi_\pm$. Suitable symmetries, such as R-parity and discrete R-symmetries, restrict the superpotential couplings of these fields. We consider leading non-renormalizable corrections to the superpotential and the Kähler potential, while restricting ourselves to sub-Planckian values of the inflaton field. We find that sneutrino driven inflation leads to inflationary parameters and, in particular, the spectral index, compatible with observations. Thus, modulo general inflationary issues, such as the gravitino problem, right-handed sneutrino driven inflation seems to be a viable scenario.

The extension of the MSSM with three right-handed neutrino superfields N_i^c realizes the seesaw mechanism through the renormalizable superpotential

$$\mathcal{W}_0 = \frac{M_R^{(i)}}{2} N_i^c N_i^c + Y_{ij} N_i^c L_j H^c. \tag{1}$$

This is the most general renormalizable superpotential, assuming R-parity conservation and taking the right-handed neutrino to be odd under it. We shall extend this model further by introducing an extra $U(1)$ gauge factor under which all fields are neutral except an oppositely charged pair ϕ_+ and ϕ_-. We are also going to assume that a non-zero Fayet–Iliopoulos D-term is present. A renormalizable superpotential term $N^c\phi_+\phi_-$ would lead to a neutrino state of electroweak mass but, here, it is forbidden by R-parity, if we take the product of $\phi_\pm$ to be even. Then, an extra symmetry would be necessary in order to forbid the direct $\phi_+\phi_-$ mass-term. A suitable such symmetry is *a discrete R-symmetry* and as a specific example we may take the following $\mathcal{Z}_3^{(R)}$ symmetry [1]

[1] All standard terms of the superpotential are allowed, provided $(H^c, H, L, E^c, Q, D^c, U^c) \to (1, \alpha^2, \alpha, \alpha^2, \alpha, \alpha^2, \alpha)$. This symmetry is not assumed to be a symmetry of the sector of the theory responsible for the spontaneous breaking of local supersymmetry. It will be explicitly broken by supersymmetry breaking.

* Corresponding author.
E-mail address: oefthimiou@grads.uoi.gr (O. Efthimiou).

0370-2693/$ – see front matter
doi:10.1016/j.physletb.2010.02.001

$$N_i^c \to \alpha N_i^c, \qquad \phi_\pm \to \phi_\pm, \qquad \mathcal{W} \to \alpha^2 \mathcal{W}. \tag{2}$$

Adopting this symmetry, we see that at the lowest non-renormalizable level the only allowed term is $\frac{\lambda_{ij}}{M} N_i^c N_j^c \phi_+\phi_-$. Also, there are no allowed terms of $O(M^{-2})$. Actually, the allowed superpotential form, involving only right-handed neutrino fields and $\phi_\pm$, to all orders can be written down as

$$\mathcal{W}_N = (N^c)^2 \mathcal{F}\big(\phi_+\phi_-, (N^c)^6\big) = \sum_{n,m=0} \frac{C_{n,m}}{M^{2m+6n-1}} (\phi_+\phi_-)^m (N^c)^{3n+2}. \tag{3}$$

The first few allowed terms, apart from (1), are

$$\frac{\lambda_{ij}}{M} N_i^c N_j^c \phi_+\phi_- + \frac{\lambda'_{ij}}{M^3} N_i^c N_j^c (\phi_+\phi_-)^2 + \frac{\lambda''_{ij}}{M^5} N_i^c N_j^c (\phi_+\phi_-)^3 + \frac{\lambda_{i_1\ldots i_8}}{M^5} N_{i_1}^c \cdots N_{i_8}^c + \cdots.$$

Throughout this Letter we assume that the defining scale of non-renormalizable terms will be of the order of the reduced Planck mass and we set $M = M_P \sim 2.4 \times 10^{18}$ GeV. The only other dimensionfull parameters appearing in the superpotential are the right-handed neutrino masses $M_R^{(i)}$. In order to obtain an acceptable neutrino mass through the seesaw mechanism, we must take the right-handed neutrino masses M_1, M_2, M_3 in the range 10^{10}–10^{14} GeV. We shall assume that one of these right-handed sneutrino fields, namely, the lightest, will play a role in inflation and we shall suppress family indices in what follows.

Since, we have considered non-renormalizable terms in the superpotential, we must do the same with the Kähler potential as well [10]. Restricting ourselves to the quadratic $\phi_\pm$ term, we have

$$\mathcal{K} = |N^c|^2 f(n) + g(n)\big(|\phi_+|^2 + |\phi_-|^2\big) + \cdots, \tag{4}$$

where $f(n)$ and $g(n)$ are arbitrary functions defined as

$$f(n) = \sum_{j=0} f_j n^j, \qquad g(n) = \sum_{j=0} g_j n^j, \quad n \equiv \frac{|N^c|^2}{M^2}. \tag{5}$$

The dots in (4) correspond to higher powers of $\phi_\pm$, which will be neglected, since we anticipate that $\phi_\pm$ will either stay at the origin or obtain values $\ll M$.

We may next proceed to calculate the scalar potential resulting from (4) and the superpotential

$$\mathcal{W} = \frac{M_R}{2} N^{c\,2} + \frac{\lambda}{2M} N^{c\,2}(\phi_+\phi_-) + \cdots, \tag{6}$$

where the dots denote terms of $O(M^{-3})$ or higher. The F-term part of the potential is

$$\begin{aligned} V_F &\approx e^{\mathcal{K}/M^2}\left\{\mathcal{F}_i\big(\mathcal{K}_j^i\big)^{-1}\mathcal{F}^j - 3\frac{|\mathcal{W}|^2}{M^2}\right\} \\ &\approx \frac{|N^c|^2}{A}\bigg[M_R^2 + \frac{\lambda^2}{M^2}|\phi_+\phi_-|^2 + \frac{\lambda M_R}{M}\big(\phi_+\phi_- + \phi_+^*\phi_-^*\big) \\ &\quad + \frac{M_R^2}{2M^2}\big(N^{c\,2} + N^{c*2}\big)\big(f + nf^{(1)}\big)\bigg] \\ &\quad - \frac{BM_R^2}{M^2A^2}|N^c|^2\big(|\phi_+|^2 + |\phi_-|^2\big) + \frac{\lambda^2|N^c|^4}{4g(n)M^2}\big(|\phi_+|^2 + |\phi_-|^2\big) \\ &\quad + \frac{M_R^2|N^c|^2}{AM^2}\big[f(n)|N^c|^2 + g(n)\big(|\phi_+|^2 + |\phi_-|^2\big)\big] \\ &\quad - \frac{3M_R^2}{4M^2}|N^c|^4 + \cdots. \end{aligned} \tag{7}$$

We have assumed that field values will be below M and kept terms up to $O(M^{-2})$, while keeping the functions f and g intact. The functions A and B stand for

$$\begin{aligned} A(n) &\equiv f(n) + 3nf^{(1)}(n) + n^2 f^{(2)}(n), \\ B(n) &\equiv g^{(1)} + ng^{(2)}(n). \end{aligned} \tag{8}$$

If we ignore the subleading terms proportional to M_R, we obtain

$$V_F \approx \frac{\lambda^2|N^c|^2}{AM^2}|\phi_+\phi_-|^2 + \frac{\lambda^2|N^c|^4}{4M^2 g(n)}\big(|\phi_+|^2 + |\phi_-|^2\big) + \cdots. \tag{9}$$

The D-term part of the scalar potential is

$$\begin{aligned} V_D &= \frac{\tilde{g}^2}{2}(\mathcal{K}_+\phi_+ - \mathcal{K}_-\phi_- + \xi)^2 \\ &= \frac{\tilde{g}^2}{2}\big(g(n)\big(|\phi_+|^2 - |\phi_-|^2\big) + \xi\big)^2, \end{aligned} \tag{10}$$

where we have introduced a Fayet–Iliopoulos D-term $\xi \gg M_R^2$.

The $\phi_\pm$-mass terms read off from the potential are

$$M_\pm^2 = \frac{\lambda^2}{4M^2 g(n)}|N^c|^4 \pm \tilde{g}^2\xi g(n). \tag{11}$$

Both masses are positive for

$$n^2 = \frac{|N^c|^4}{M^4} \geqslant \frac{4\tilde{g}^2}{\lambda^2}\frac{\xi}{M^2}g^2(n). \tag{12}$$

Anticipating sub-Planckian values for N^c, we may define the *critical field*

$$n_c \equiv \frac{2\tilde{g}}{\lambda} g_0 \frac{\sqrt{\xi}}{M}, \tag{13}$$

where $g(n) \approx g_0 + ng_1 + \cdots$. Thus, (12) becomes just $n \geqslant n_c$, or, in terms of a real field

$$N^c = \frac{\phi}{\sqrt{2}} \quad \Longrightarrow \quad \phi \geqslant \phi_c = 2\sqrt{\frac{g_0\tilde{g}}{\lambda}}\sqrt{\xi}M. \tag{14}$$

For $\phi > \phi_c$ both $\phi_\pm$ have positive masses and their vevs stay at the origin. In the unbroken phase the scalar potential is

$$V(\phi) \approx \frac{\tilde{g}^2\xi^2}{2} + O\big(M_R^2\big)\phi^2 + O\big(M_R^2/M^2\big)\phi^4 + \cdots,$$

which is approximately equal to $\tilde{g}^2\xi^2/2$ and very flat. Thus, when the Universe is in the above global vacuum the energy density is constant and inflation can occur. The amount of inflation depends on the initial and final value the sneutrino field which plays the role of the inflaton. When the critical value $\phi \sim \phi_c$ is reached the above local minimum ceases to exist and the Universe makes a transition to the global minimum

$$\phi_+ = 0, \qquad \phi_- \approx \sqrt{\xi/g_0} \tag{15}$$

and inflation stops.

At the local vacuum $\phi_\pm = 0$, the tree potential receives considerable radiative corrections given by the standard Coleman–Weinberg formula in the terms of the split masses of the $\phi_\pm$ superfields as[2]

[2] The kinetic terms, near the origin $\phi_\pm = 0$, are

$$A(n)\big|\partial_\mu N^c\big|^2 + g(n)\big(|D_\mu\phi_+|^2 + |D_\mu\phi_-|^2\big).$$

$$\Delta V = \frac{1}{32\pi^2}\sum_{\pm}\frac{M_{\pm}^4}{g^2(n)}\ln\left(\frac{M_{\pm}^2}{g(n)\Lambda^2}\right) - \frac{1}{16\pi^2}\frac{M^4(0)}{g^2(n)}\ln\left(\frac{M^2(0)}{g(n)\Lambda^2}\right),$$

where

$$M_{\pm}^2 = M^2(0) \pm \tilde{g}^2\xi g(n).$$

Thus, we finally obtain

$$V \approx \frac{\tilde{g}^2}{2}\xi^2 + \frac{\tilde{g}^4\xi^2}{16\pi^2}\ln\left(\frac{\phi^4}{g^2(n)\tilde{\Lambda}^4}\right), \tag{16}$$

where we have suitably chosen the cutoff to absorb the constant factors.

Let us conclude the presentation of the model by discussing the mass-scales involved. We have already assumed that the scale of the right-handed neutrino mass M_R is much smaller than the other scales involved. The assumed range of M_R is between 10^{10} and 10^{14} GeV. Apart from the reduced Planck scale $M \sim 2.4 \times 10^{18}$ GeV, the only other scale appearing is the Fayet–Iliopoulos scale $\sqrt{\xi}$. There is a well-known constraint for ξ coming from the formation of cosmic strings [11], namely

$$3.8 \times 10^{15}\ \text{GeV} \leqslant \sqrt{\xi} \leqslant 4.6 \times 10^{15}\ \text{GeV}. \tag{17}$$

In order to safely ignore contributions from the right-handed neutrino mass M_R, we should have

$$M_R^2\phi^2 \ll \tilde{g}^2\xi^2 \quad\Longrightarrow\quad M_R \ll \tilde{g}\xi/\phi_{\max}.$$

Taking $\sqrt{\xi} \sim 4.6 \times 10^{15}$ GeV and $\tilde{g} \sim O(0.1)$, for $\phi_{\max} \sim M$, we obtain

$$M_R \ll 10^{12}\ \text{GeV}.$$

Note that the non-minimality of the Kähler potential will not modify this significantly, since, for sub-Planckian values of n, $A(n) \approx f_0$ and the physical mass is just $M_R^* = M_R/\sqrt{f_0}$ with $f_0 \sim O(1)$.

In the local vacuum the scalar potential is well approximated by the constant $\tilde{g}^2\xi^2/2$ plus the radiative corrections part (16). Assuming that the slow-roll approximation is valid, namely $\ddot{\phi} \ll H\dot{\phi}$ and $(\dot{\phi})^2 \ll V(\phi)$, the classical evolution equations are

$$3H\dot{\phi} \approx -\frac{V'(\phi)}{A(\phi)},$$

$$H^2 \approx \frac{V(\phi)}{3M^2} \quad\Longrightarrow\quad \frac{d\phi}{d\ln a} \approx -\frac{M^2}{A(\phi)}\frac{V'(\phi)}{V(\phi)}. \tag{18}$$

The above expression can be integrated to give the *number of e-folds* $\mathcal{N}$ as

$$\mathcal{N} = \ln\left(\frac{a_f}{a_i}\right) \approx \frac{1}{M^2}\int_{\phi_f}^{\phi_i} d\phi\, A(\phi)\frac{V(\phi)}{V'(\phi)} = \frac{2\pi^2}{\tilde{g}^2}\int_{n_f}^{n_i} dn\left(\frac{f(n) + 3nf^{(1)}(n) + n^2 f^{(2)}(n)}{1 - n\frac{g^{(1)}}{g}}\right). \tag{19}$$

Anticipating $n \ll 1$, we approximate and obtain

$$\mathcal{N} \approx \frac{2\pi^2}{\tilde{g}^2}\int_{n_f}^{n_i} dn\left(f_0 + n\left(4f_1 + \frac{g_1}{g_0}f_1\right) + \cdots\right)$$

or

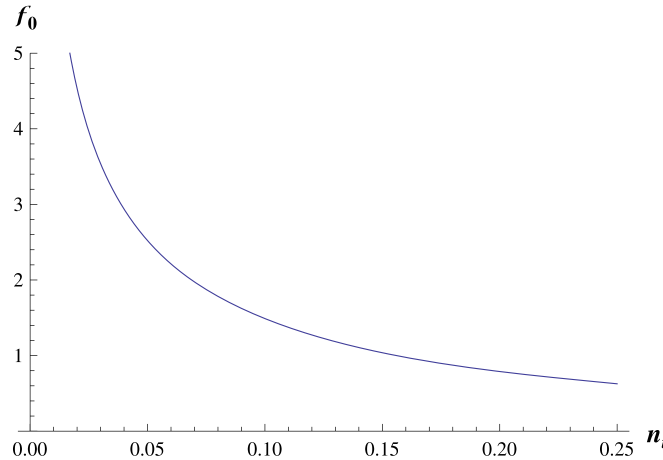

Fig. 1. The Kähler coefficient f_0 as a function of n_i.

$$\frac{\tilde{g}^2}{2\pi^2}\mathcal{N} \approx \left(f_0(n_i - n_f) + \frac{1}{2}(n_i^2 - n_f^2)\left(4f_1 + \frac{g_1}{g_0}f_0\right)\right). \tag{20}$$

The *comoving curvature perturbation*, in terms of the potential (16), is

$$\mathcal{R}_c = \frac{H^2}{2\pi|\dot{\phi}|} \approx \frac{V^{3/2}A}{2\pi\sqrt{3}M^3|V'|} = \frac{\pi}{\sqrt{2}\sqrt{3}\tilde{g}}\left(\frac{\xi}{M^2}\right)\left(\frac{\phi_i}{M}\right)\left(\frac{f(n_i) + 3n_i f^{(1)}(n_i) + n_i^2 f^{(2)}(n_i)}{1 - n_i\frac{g^{(1)}(n_i)}{g(n_i)}}\right) \approx \frac{\pi}{\sqrt{6}\tilde{g}}\left(\frac{\xi}{M^2}\right)\left(\frac{\phi_i}{M}\right)\left(f_0 + n_i\left(4f_1 + \frac{g_1}{g_0}f_0\right)\right). \tag{21}$$

Matching this to the observed value $\mathcal{R}_c \approx 4.7 \times 10^{-5}$ and choosing $\sqrt{\xi} \approx 4.6 \times 10^{15}$ GeV and $\tilde{g} \approx 0.1$, amounts to the constraint

$$\sqrt{n_i}\left(f_0 + n_i\left(4f_1 + \frac{g_1}{g_0}f_0\right)\right) \approx 0.705. \tag{22}$$

Similarly, (20), for $\mathcal{N} \sim 65$, can be thought off as an equation constraining n_f.

As a very rough approximation, we shall assume that $f_1 \sim f_0$ and $g_1 \sim g_0$. The range of values for f_0 is shown in Fig. 1 for a range of sub-Planckian values of n_i. Assuming that inflation ends when the value n_c is reached, we may identify

$$n_f \approx n_c = \frac{2\tilde{g}}{\lambda}g_0\frac{\sqrt{\xi}}{M}. \tag{23}$$

For the chosen values, of $\sqrt{\xi} \sim 4.6 \times 10^{15}$ GeV and $\tilde{g} \sim 0.1$, this corresponds to $n_c \sim 0.38 \times 10^{-3}(g_0/\lambda)$ and requires a small (but not unnatural) coupling $\lambda \sim g_0 O(0.01)$. In Fig. 2 n_f is plotted as a function of n_i.

Let us now consider the *slow-roll parameters* ϵ, η and ζ. They are given in terms of the potential and its derivatives. We have

$$\epsilon = M^2\left(\frac{V'}{V}\right)^2 \approx \frac{1}{8n_i}\left(\frac{\tilde{g}}{\pi}\right)^2\left(1 - 2n_i\frac{g_1}{g_0} + O(n_i^2)\right), \tag{24}$$

$$\eta = 2M^2\frac{V''}{V} \approx -\frac{1}{2n_i}\left(\frac{\tilde{g}}{\pi}\right)^2\left(1 + n_i\frac{g_1}{g_0} + O(n_i^2)\right), \tag{25}$$

$$\zeta = \frac{M^2}{2}\sqrt{\frac{V'''V'}{V^2}} \approx \frac{1}{4\sqrt{2n_i}}\left(\frac{\tilde{g}}{\pi}\right)^2\left(1 - n\frac{g_1}{g_0} + O(n_i^2)\right). \tag{26}$$

The corresponding *spectral index* n_s is

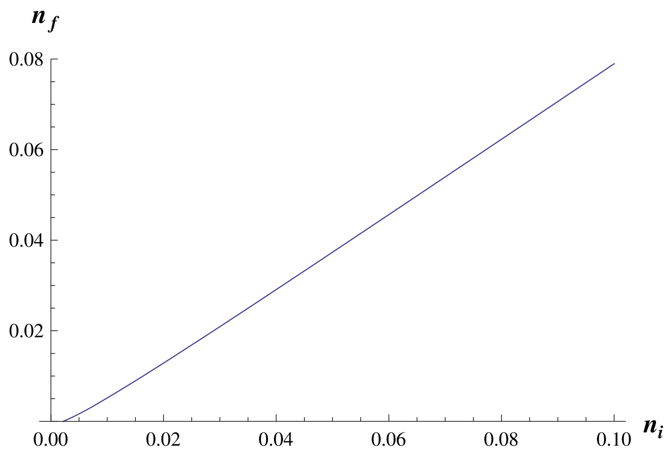

Fig. 2. n_f as a function of n_i.

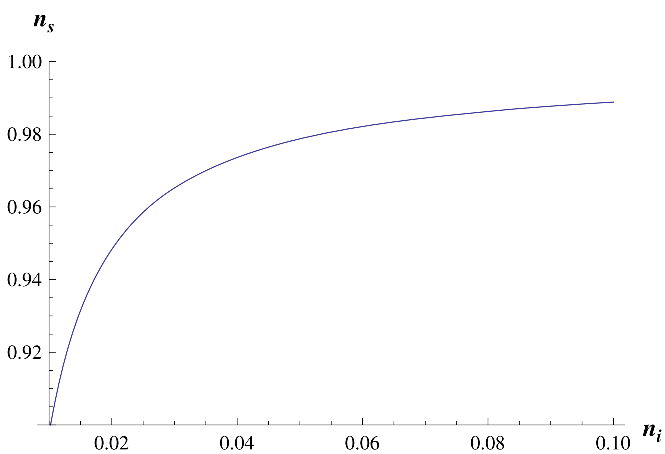

Fig. 3. Spectral index n_s as a function of n_i.

$$\begin{aligned} n_s &= 1 + 2\eta_i - 6\epsilon_i \approx 1 + 2\eta_i \\ &\approx 1 - \frac{1}{n_i}\left(\frac{\tilde{g}}{\pi}\right)^2 \left(1 + n_i \frac{g_1}{g_0} + O\left(n_i^2\right)\right) \end{aligned} \tag{27}$$

and it is plotted in Fig. 3 for the choice $\tilde{g} \sim 0.1$ and $g_0 \sim g_1$. From this plot one can immediately see that our values for the spectral index are compatible with the corresponding value from observational data, $n_s = 0.963^{+0.014}_{-0.015}$ [12].

We have assumed that the right-handed sneutrino that drives inflation is the lightest one. It is well known that this sneutrino through its lepton violating processes can produce the desired lepton number asymmetry. Its decays ($\tilde{N}^c \to \tilde{L}^* + H^{c*}$) will reheat the Universe to a temperature $T_R = 1.4 \times 10^{10} \sqrt{M_R^*(\mathrm{GeV})/10^{10}} \times \sqrt{\sum_j |Y_{1j}|^2/10^{-6}}$. In order to avoid overproduction of gravitinos this temperature must not exceed 10^6–10^7 GeV. This corresponds to the preferred values $Y \sim 10^{-6}$ and $M_N \sim 10^{10}$ GeV. Note however that it is possible that a late time entropy production will dilute the gravitino density without the need of small Yukawa couplings. Apart from the ubiquitous gravitino problem, there is also another issue, namely the problem associated with the decays of the heavy fields $\phi_\pm$, a typical problem of all models of D-term inflation. These decays can lead to a potentially high reheating temperature causing overproduction of gravitinos. Note however that the subsequent decay of the right-handed sneutrino at a lower temperature will produce additional entropy that can sufficiently dilute gravitinos. The precise way this can occur depends on the range of various parameters.

Let us now conclude summarizing the main points of this Letter. We have considered the right-handed neutrino extension of the MSSM that realizes neutrino masses through the see saw mechanism. We extended this model further with a $U(1)$ gauge group under which all standard fields are neutral apart from a pair of MSSM singlets $\phi_\pm$. Suitable symmetries, such as R-parity and discrete R-symmetries, restrict the superpotential couplings of these fields to a class of non-renormalizable operators. As it stands the model can realize the scenario of D-term inflation with the inflaton being identified with the right-handed sneutrino field. Considering also non-minimal corrections to the Kähler potential but staying in the sub-Planckian field space, we arrive at inflationary predictions compatible with existing data, thus, establishing the possibility of right-handed sneutrino driven inflation as a viable scenario.

Acknowledgements

The authors acknowledge the support of European Research and Training Networks MRTPN-CT-2006 035863-1 (UniverseNet) and “UNILHC” PITN-GA-2009-237920 (Unification in the LHC era). K.T. thanks the CERN Theory Group for its hospitality.

References

[1] Super-Kamiokande Collaborations, Y. Ashie, et al., Phys. Rev. D 71 (2005) 112005;
Y. Ashie, et al., Phys. Rev. Lett. 93 (2004) 101801;
KamLAND Collaboration, T. Araki, et al., Phys. Rev. Lett. 94 (2005) 081801;
SNO Collaboration, B. Aharmin, et al., Phys. Rev. C 72 (2005) 055502;
K2K Collaboration, E. Aliu, et al., Phys. Rev. Lett. 94 (2005) 081802;
CHOOZ Collaboration, M. Apollonio, et al., Nucl. Eur. Phys. J. C 27 (2003) 331;
Super-Kamiokande Collaborations, Y. Ashie, et al., Phys. Rev. D 71 (2005) 112005;
Y. Ashie, et al., Phys. Rev. Lett. 93 (2004) 101801.
[2] M. Gell-Mann, P. Ramond, R. Slansky, in: D. Freedman, P. van Nieuwenhuizen (Eds.), Supergravity, North-Holland, Amsterdam, 1979, p. 315;
T. Yanagida, in: O. Sawada, A. Sugamoto (Eds.), Proceedings of the Workshop on Unified Theory and Baryon Number in the Universe, KEK, Tsukuba, Japan, 1979;
R. Mohapatra, G. Senjanovic, Phys. Rev. Lett. 44 (1980) 912;
R. Mohapatra, G. Senjanovic, Phys. Rev. D 23 (1981) 165;
P. Minkowski, Phys. Lett. B 67 (1977) 421.
[3] H. Murayama, H. Suzuki, T. Yanagida, J. Yokoyama, Phys. Rev. Lett. 70 (1993) 1912;
H. Murayama, H. Suzuki, T. Yanagida, J. Yokoyama, Phys. Rev. D 50 (1994) 2356.
[4] J. Ellis, M. Raidal, T. Yanagida, Phys. Lett. B 581 (2004) 9.
[5] K. Kadota, J. Yokoyama, Phys. Rev. D 73 (2006) 043507.
[6] E. Halyo, Phys. Lett. B 387 (1996) 43.
[7] P. Binetruy, G. Dvali, Phys. Lett. B 388 (1996) 241.
[8] S. Antusch, M. Bastero-Gil, S. King, Q. Shafi, Phys. Rev. D 71 (2005) 083519.
[9] C.M. Lin, J. McDonald, Phys. Rev. D 77 (2008) 063529, arXiv:0710.4273 [hep-ph].
[10] O. Seto, J. Yokoyama, Phys. Rev. D 73 (2006) 023508.
[11] R. Jeannerot, Phys. Rev. D 56 (1997) 6205;
D. Lyth, A. Riotto, Phys. Lett. B 412 (1997) 28;
J. Rocher, M. Sakellariadou, JCAP 0503 (2005) 004;
M. Endo, M. Kawasaki, T. Moroi, Phys. Lett. B 569 (2003) 73.
[12] WMAP Collaboration, J. Dunkley, et al., Astrophys. J. Suppl. 180 (2009) 306, arXiv:0803.0586 [astro-ph].