INFLATION, STRINGS, CMB ANISOTROPIES AND BARYOGENESIS

R. JEANNEROT
Instituut-Lorentz for Theoretical Physics,
Niels Bohrweg 2, 2333 CA Leiden, The Netherlands

In this talk, I shall focus on theories beyond the Standard Model which predict massive neutrinos. Hybrid inflation emerges naturally in these theories: the slow-rolling inflaton field is a gauge singlet which couples with a GUT Higgs field which triggers the end of inflation. In the standard scenario, spontaneous symmetry breaking takes place at the end of inflation at a scale M; $M_{GUT} > M > M_Z$ for inflation to solve the GUT monopole problem and cosmic strings always form at this intermediate scale. WMAP data constrain $M \in [10^{14.5} - 10^{15.5}]$ GeV and the singlet-Higgs coupling $\kappa \in [10^{-7} - 10^{-2}]$. The spectral index $n_s \gtrsim 0.98$ in slight conflict with WMAP3. When the symmetry which is broken at the end of inflation is gauged $B - L$, both the inflaton and the strings decay into right-handed neutrinos. There are then two competing non-thermal scenarios for baryogenesis via leptogenesis which take place at the end of inflation, during reheating and from cosmic strings decay. Which of the two scenarios dominates depends on the inflaton-neutrino sector parameters.

1 Introduction

Up to the discovery of the 'acoustic' peaks in the CMB power spectrum [1], there were two compelling mechanisms for explaining cosmological perturbations: inflation and cosmic strings [2]. Since cosmic strings predict a single peak, they are now excluded as main source of the cosmological perturbations. However, a mixed scenario with both inflation and cosmic strings with a string contribution less than about 10% is still allowed by the data[3,4]. In many models with both inflation and strings, the scalar perturbations are dominated by the scalar perturbations from inflation, and the string contribution may be too low for detection via the CMB temperature anisotropies. However they could be detected via the B-type polarization of the CMB [5].

From a theoretical point of view, inflation is often associated with the formation of cosmic strings. Perhaps the best particle physics motivated model of inflation is hybrid inflation [6]. It arises naturally in Supersymmetric (SUSY) Grand unified Theories (GUTs) [7,8], in effective strings theories and in brane worlds. Naturally meaning that the fields and the potential leading to hybrid inflation are needed to build the theory itself (I now focus on the case of SUSY GUTs) and the coupling constant which enters is the order unity. In either cases, spontaneous symmetry breaking takes place at the end of inflation [a] and cosmic strings form [8,12,13,14]. In this talk I will consider Standard hybrid inflation in the context of SUSY GUTs. And I shall be mainly concerned about models which contain $B-L$ as a gauge symmetry and predict massive neutrinos via the See-saw mechanism [15].

[a] In non minimal models of hybrid inflation such as shifted inflation [10] or smooth inflation [11], spontaneous symmetry breaking takes place before or during inflation, and no defect form at the of inflation.

In section 2, I show that cosmic strings always form at the end of standard [b] hybrid inflation when inflation solves the GUT monopole problem. In section 3, I study the CMB anisotropies which are predicted by these models. Matching theoretical predictions with the data gives constraints on two of the GUT parameters, the Spontaneous Symmetry Breaking (SSB) scale at the end of inflation and the relevant coupling constant [16,17]. In section 4, I show that when the symmetry which is broken at the end of inflation is gauged $B-L$, there are two competing non-thermal baryogenesis scenarios which take place after inflation: from reheating, and from cosmic strings decay [18].

2 Inflation and cosmic strings

2.1 Inflation from particle physics

Inflation must come from the particle physics model describing fundamental interactions at high energies. As a particle physicist, the first question i will ask is 'Can we get inflation from the Standard Model?' On general grounds, the answer is 'No', because the inflationary energy scale would be the order of 100 GeV which is far too low to produced the required amount of primordial perturbations [c]. The next question i will ask is 'Can we get inflation from the simplest extensions of the Standard Model?' As an aparté, we know since the discovery of neutrinos oscillations that neutrinos are massive and hence that the Standard Model must be extended. In order to explain the smallness of the observed neutrinos masses, one could just add a gauge singlet and a tiny coupling constant. However, by adding a $U(1)_{B-L}$ gauge symmetry to the Standard Model gauge group $G_{SM} = SU(3)_c \times SU(2)_L \times U(1)_Y$, massive neutrinos become a prediction [15]. Adding the idea of unification of the gauge coupling constants, one is lead to grand unified theories. So I shall rephrase the question as 'Can we get inflation from a grand unified theory?' At first sight, 'the unification scale $M_{GUT} \sim 10^{16}$ GeV is just the energy scale needed for inflation to explain the cosmological perturbations'.

It turns out that when building a model of slow-roll inflation in a theory beyond the Standard Model three ingredients are usually needed: SUSY, which provides the required flatness of the potential, a Standard Model singlet, the slow-rolling field, and GUT Higgs fields transforming under a gauge group G whose rank is larger than the rank of the Standard Model gauge group, i.e. rank$(G) > 4$ [8,13].

2.2 Standard hybrid inflation

Hybrid inflation [6] uses two fields instead of one, a gauge singlet S and a Higgs field Φ. Hybrid inflation is arguably the best particle physics motivated model of inflation. In the context of SUSY GUTs, there are two Higgs superfields Φ and $\bar{\Phi}$ in complex conjugate representations of the GUT gauge group G_{GUT} which lower the rank of the group by one unit when acquiring vacuum expectation values (VEV) at the end of inflation. The superpotential is given by

$$W_{\rm inf} = \kappa S(\bar{\Phi}\Phi - M^2), \tag{1}$$

where a suitable U(1) R-symmetry under which W and S transform in the same way ensures the uniqueness of this superpotential at the renormalizable level. The scalar potential has an inflationary valley, which is a valley of local minima, at $S > M$ and $|\Phi| = |\bar{\Phi}| = 0$. At tree level, the potential along this valley is $V_{infl} = \kappa^2 M^2$. S is the slowing rolling field and slow-roll conditions thus apply to S. Since $|\Phi| = |\bar{\Phi}| = 0$ during inflation, there is no symmetry

[b] *Standard* refers to the standard model of SUSY hybrid inflation [7] where SSB takes place at the end of inflation

[c] However it has been recently suggested that an MSSM flat direction might be suitable for inflation [19]. Even though this proposal requires strong fine-tuning, it is interesting in two ways: first of all it uses standard model physics, and second there is no need of standard model singlet.

breaking induced by these Higgs fields VEV during inflation. Inflation terminates as S falls below its critical value $S_c = M$ and inflation ends in a phase transition during which the Higgs fields acquire non-zero VEV equal to M: Spontaneous Symmetry Breaking (SSB) takes place at the end of inflation. The SSB scale M and is proportional to the inflationary scale $V_{infl}^{1/4}$, the proportionality coefficient being the squared root of the singlet-Higgs coupling κ.

The Higgs fields representations Φ and $\bar{\Phi}$ are conjugate N-dimensional representations of the GUT gauge group. We are now focusing on GUT which contain gauged $U(1)_{B-L}$ and predict massive neutrinos via See-saw. The component of of Φ (and $\bar{\Phi}$) which gets a VEV at the end of inflation transforms as an Standard Model singlet, and it also transforms either as an $SU(2)_R$ doublet or as an $SU(2)_R$ triplet. In a realistic model where there are no unwanted light fields between the scale M and the GUT scale, it is the only component which remains light below M_{GUT} [17] ($M < M_{GUT}$, see section 2.3). The scalar potential along the inflationary valley is flat at tree level. It is lifted by loop corrections, which are non-zero during inflation because SUSY is spontaneously broken, and by SUGRA corrections. Assuming minimum Khäler potential it is given by [16,17]

$$\begin{aligned} \frac{V}{\kappa^2 M^4} &= 1 + \frac{\kappa^2 \mathcal{N}}{32\pi^2}\Big[2\ln\Big(\frac{\kappa^2 M^2 x^2}{\Lambda^2}\Big) + (x^2+1)^2\ln(1+x^{-2}) + (x^2-1)^2\ln(1-x^{-2})\Big] \\ &+ 2x^4\Big(\frac{M}{m_{\rm p}}\Big)^4 + |a|^2 x^2\Big(\frac{M}{m_{\rm p}}\Big)^2 + A\frac{m_{3/2}}{M}x, \end{aligned} \tag{2}$$

where $m_{\rm p}$ is the reduced Planck mass and Λ a cutoff scale; $x = |S|/M$ so that $x \rightarrow 1$ at the critical point; $A = 4\cos(\arg m_{3/2} - \arg S)$, we assume that $\arg S$ is constant during inflation; $\mathcal{N} = 1 - 3$ depending on wether the components of Φ and $\bar{\Phi}$ which get a VEV at the end of inflation transform as an $SU(2)_R$ doublet or triplet and wether the symmetry group which breaks at the end of inflation contains an $SU(2)_R$ or an $U(1)_R$ symmetry. Hidden sector VEV which lead to low energy SUSY breaking are $\langle z\rangle = a m_{\rm p}$ and $\langle W_{\rm hid}(z)\rangle = m_{3/2}\exp^{-|a|^2/2} m_{\rm p}^2$, with $m_{3/2}$ the gravitino mass; the cosmological constant in the global minimum is set to zero by hand. All subdominant terms are dropped.

2.3 Cosmic strings form at the end of standard hybrid inflation

Since SSB takes place at the end of inflation, cosmic strings always form if the later solves the GUT monopole problem [13]. The underlying reason being that the rank of the gauge group is lowered by one unit at the end of inflation [8,13]. This is illustrated in reference [12] where an exhaustive study of all SSB breaking patterns for all GUT gauge groups with rank less than height and phenomenologically acceptable has been performed. The aim of this section is to understand why indeed cosmic strings form. Further details can fe found in reference [13].

Suppose that the Standard Model gauge group G_{SM} is embedded in a GUT gauge group G_{GUT}. This must be broken down to G_{SM} at around $M_{GUT} \sim 10^{16}$ GeV, which is the scale at which the gauge couplings unify

$$G_{GUT} \stackrel{M_{GUT}}{\rightarrow} \ldots \rightarrow G_{SM} \stackrel{10^2\,\mathrm{GeV}}{\rightarrow} SU(3)_c \times U(1)_Q. \tag{3}$$

In SUSY, the breaking of G_{GUT} down to G_{SM} can be direct of via intermediate symmetry groups, whereas in the non SUSY case there must be at least one intermediate step. If (some of) the Higgs fields used to break G_{GUT} have a superpotential given by Eq.(1) [d], inflation takes place and the spontaneous symmetry breaking of G_{GUT} takes place at the end of inflation. But in this scenario, cosmologically catastrophic monopoles which ought to form in all GUTs, form after inflation. In order to cure the monopole problem, one must introduce an intermediate

[d] The rank of G_{GUT} has to be strictly greater than the rank of G_{SM} [13]

symmetry group H, a subgroup of G_{GUT}, and use $\Phi, \bar{\Phi}$ not to break the GUT itself but this intermediate symmetry group H; the symmetry breaking scale M of H is $< M_{GUT}$. H must be chosen in such a way that the monopoles form between G_{GUT} and H and no unwanted defect form when H breaks down to G_{SM}

$$G_{GUT} \stackrel{\text{Monopoles}}{\rightarrow} ...H \stackrel{\Phi.\bar{\Phi},\text{ No unwanted defect}}{\rightarrow} ... \rightarrow G_{SM} \stackrel{10^2\text{ GeV}}{\rightarrow} SU(3)_c \times U(1)_Q. \quad (4)$$

It can be shown that the rank of H must be greater than five, that it must contain a U(1) factor [8,13], and that cosmic strings always form when H breaks down to G_{SM} [8,12,13].

$$G_{GUT} \stackrel{\text{Monopoles}}{\rightarrow} ...H \stackrel{\text{Inflation, Cosmic Strings}}{\rightarrow} ... \rightarrow G_{SM} \stackrel{10^2\text{ GeV}}{\rightarrow} SU(3)_c \times U(1)_Q. \quad (5)$$

3 CMB constraints and predictions

If the strings which form at the end of inflation are stable down to low energy, they will contribute to the CMB temperature anisotropies. The perturbations from inflation and cosmic strings are uncorrelated and they add up independently [8]

$$\left(\frac{\delta T}{T}\right)_{\rm tot} = \sqrt{\left(\frac{\delta T}{T}\right)^2_{\rm inf} + \left(\frac{\delta T}{T}\right)^2_{\rm cs}}. \quad (6)$$

The inflation contribution to the quadrapole is

$$\left(\frac{\delta T}{T}\right)_{\rm inf} = \frac{1}{12\sqrt{5}\pi m_{\rm p}^3} \frac{V^{3/2}}{V'}\bigg|_{\sigma=\sigma_Q}, \quad (7)$$

with a prime denoting derivative w.r.t. the real normalized inflaton field $\sigma = \sqrt{2}|S|$, and the subscript Q denoting the time observable scales leave the horizon. V is the scalar potential along the inflationary trajectory given by Eq.(2). The tensor perturbations from inflation $(\delta T/T)_{\rm tens} \sim 10^{-2} H/m_{\rm p}$ are very small.

The string induced perturbations are proportional to the string tension $(\delta T/T)_{\rm cs} = yG\mu$, with μ the tension and y parameterizing the density of the string network. Recent simulations predicts $y = 9 \pm 2.5$ [20], but values in the range $y = 3 - 12$ can be found in the literature [2]. The strings are formed by the Higgs fields Φ and $\bar{\Phi}$ which wind around the string in opposite directions. They are not BPS and do not satisfy the Bogomolnyi bound and hence they are lighter than BPS strings forming at the same energy scale [16]. The string tension is

$$\mu = 2\pi M^2 \theta(\beta), \quad \text{with} \quad \theta(\beta) =\sim 2.4 \ln(2/\beta)^{-1} \quad (8)$$

where the function θ encodes the correction away from the BPS limit and $\beta = (m_\phi/m_A)^2 \simeq (\kappa/g_{\rm GUT})^2$ with $g^2_{\rm GUT} \approx 4\pi/25$. Requiring the non-adiabatic string contribution to the quadrupole to be less than 10% gives the bound [16]

$$G\mu < 2.3 \times 10^{-7}\left(\frac{9}{y}\right) \quad \Rightarrow \quad M < 2.3 \times 10^{15}\sqrt{\frac{(9/y)}{\theta(\beta)}}. \quad (9)$$

The bound which comes from pulsar timing (the stochastic gravitational wave background produced by cosmic strings can disrupt pulsar timing and this has not been observed) is $G\mu < 1.0 \times 10^{-7}$ [21]; it is more stringent, but it has also more uncertainties. It corresponds to the 10% bound with $y = 20.7$.

Temperature anisotropies from both inflation and cosmic strings depend on two parameters, the SSB scale M of the intermediate symmetry group H, see Sec 2., and the singlet-Higgs coupling constant κ, see Eq.(1). Matching the theoretical predictions with the observed value $(\delta T/T) = 6.6 \times 10^{-6}$ [4] gives a constrain on M versus κ, see figure 1 [16,17]. The intermediate symmetry breaking scale must be very close to the GUT scale, $M \in [10^{14.5} - 10^{15.5}]$ GeV and the coupling constant $\kappa \in [10^{-7} - 10^{-2}]$. If the strings are unstable [17], larger values of κ are allowed.

The spectral index n_s is calculated using the slow-roll parameters and also depends on the singlet-Higgs coupling constant κ; it is also shown on figure 1 [16,17]. n_s is undistinguishable from unity for small values of κ, smaller than unity for intermediate values of κ and bigger than unity for large values of κ. It is extremely difficult to get a spectral index smaller than 0.98 with hybrid inflation except maybe with non-minimal models [27]. This is in slight conflict with WMAP 3-years data, and if these were to be confirmed, it would be excluded.

The string contribution $B = \left(\frac{\delta T}{T}\right)_{\rm tot} / \left(\frac{\delta T}{T}\right)_{\rm cs}$ is also a function of the coupling κ. It is negligibly small for most of the parameter space and saturates the 10% bound for large values of κ, which is the best interesting region for both n_s and baryogenesis (see Sec. 4). It is shown in figure 2 [16].

Further details can be found in references [16,17].

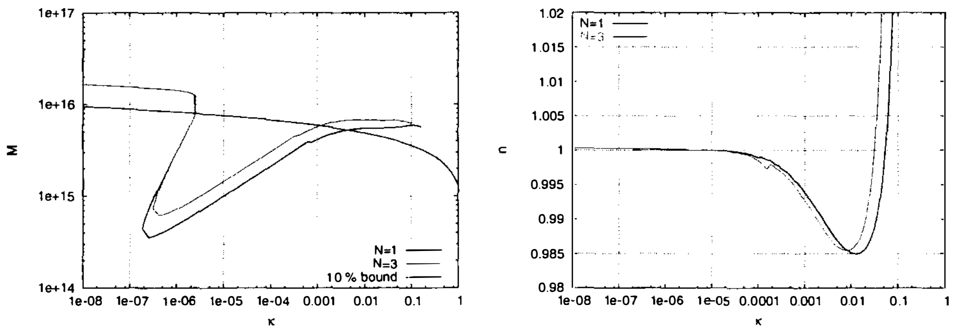

Figure 1: Left: CMB constraints on M as a function κ for $\mathcal{N} = 1, 3$ (blue curves) and the 10% bound (pink curve). Right: predictions for the spectral index n_s as function of κ for $\mathcal{N} = 1, 3$.

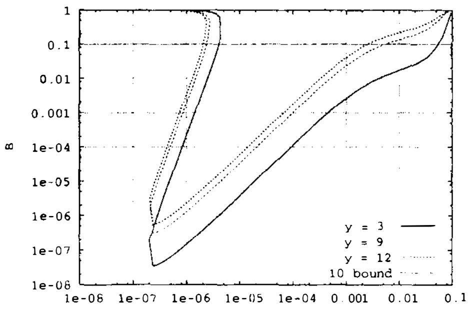

Figure 2: The string contribution $\left(\frac{\delta T}{T}\right)_{tot} / \left(\frac{\delta T}{T}\right)_{cs}$ as a function of κ for $y = 3, 6, 9$.

4 Baryogenesis via leptogenesis at the end of inflation

Baryogenesis aims to explain the observed matter-antimatter asymmetry of the Universe. It must take place after inflation, since any previously produced baryon asymmetry is washed-out. Standard GUT baryogenesis is ruled out because any GUT scale produced baryon asymmetry is erased by sphalerons transitions unless the universe possesses a $B-L$ asymmetry [e] [22]. A primordial $B-L$ asymmetry is naturally obtained in theories beyond the standard model which contain gauged $B-L$ via the out-of-equilibrium decay of heavy Majorana right-handed neutrinos [23]. This scenario known has leptogenesis is perhaps the best particle physics motivated model of baryogenesis. Thermal leptogenesis requires a symmetry breaking scale $\sim 10^{15}$ GeV and a reheating temperature $T_R \sim 10^{10}$ GeV [24]. Such high reheating temperature leads to an overproduction of gravitinos which decay lately and disrupt the predictions of nucleosynthesis.

When $G_{GUT}, H \supset U(1)_{B-L}$ and the Φ and $\bar{\Phi}$ fields entering the inflation superpotential given by equation (1) are the $B-L$ breaking Higgs fields, gauged $B-L$ is broken at the end of inflation and the strings which form at the end of inflation are the so-called $B-L$ cosmic strings [25]. There are then two competing non-thermal scenarios for leptogenesis which take place after inflation: from reheating during inflation [26] and from cosmic strings decay [18,25].

- Non-thermal leptogenesis during reheating

  The $B-L$ breaking Higgs field Φ which enters the inflationary superpotential Eq.(1) gives a superheavy Majorana mass to the right-handed neutrinos ($W \supset \Phi NN$ or $W \supset \Phi^2 NN/m_{\rm p}$) and reheating proceeds via production of heavy right-handed neutrinos and sneutrinos. Right-handed (s)neutrinos decay into electroweak Higgs(ino) and (s)leptons ($W \supset H_u LN$), CP is violated through the one-loop radiative correction involving a Higgs particle and by the self-energy correction, and lepton asymmetry is non-thermally produced when the right-handed neutrinos are out-of-equilibrium, i.e. when $T_R < M_{N_i}$. If $T_R > M_{N_1}$, the lepton asymmetry produced is wash-out by L-violating processes involving right-handed neutrinos until $T < M_{N_1}$, where M_{N_1} is the mass of the lightest right-handed neutrino. If $M_{N_1} > m_\Phi/2$, where m_Φ is the mass of the Higgs field in the true vacuum, the inflaton cannot decay into right-handed neutrinos and reheating must be gravitational. The resulting baryon asymmetry depends on the reheating temperature at the end of inflation, which depends on the mass M_{N_i} of the heaviest right-handed neutrinos the inflaton can decay into, on the symmetry breaking scale M which is constrained by CMB data as a function of the coupling κ (see Sec. 3) and on the CP violating parameter [18].

- Non-thermal leptogenesis from cosmic strings decay

  The strings which form at the end of inflation are the so-called $B-L$ cosmic strings [25]. The main decay channel of $B-L$ strings is into right-handed neutrinos and they also lead to non-thermal leptogenesis [25,18]. The resulting baryon asymmetry depends upon the amount of energy loss by the network into right-handed neutrinos and on the density of strings at the end of inflation; it also depends on the symmetry breaking scale M at the end of inflation which is constrain by CMB as a function of κ [18].

Which of these two scenarios dominates depends on wether the inflaton decay into right-handed neutrino is kinematically allowed and wether the lightest right-handed neutrino N_1 is in thermal equilibrium at reheating. Results, which take into account the CMB constraints derived in the previous section, are shown in figures 3 and 4. Further details can be found in reference [18].

[e] Sphalerons transition violate $B+L$ and conserve $B-L$, where B and L are respectively number and lepton number.

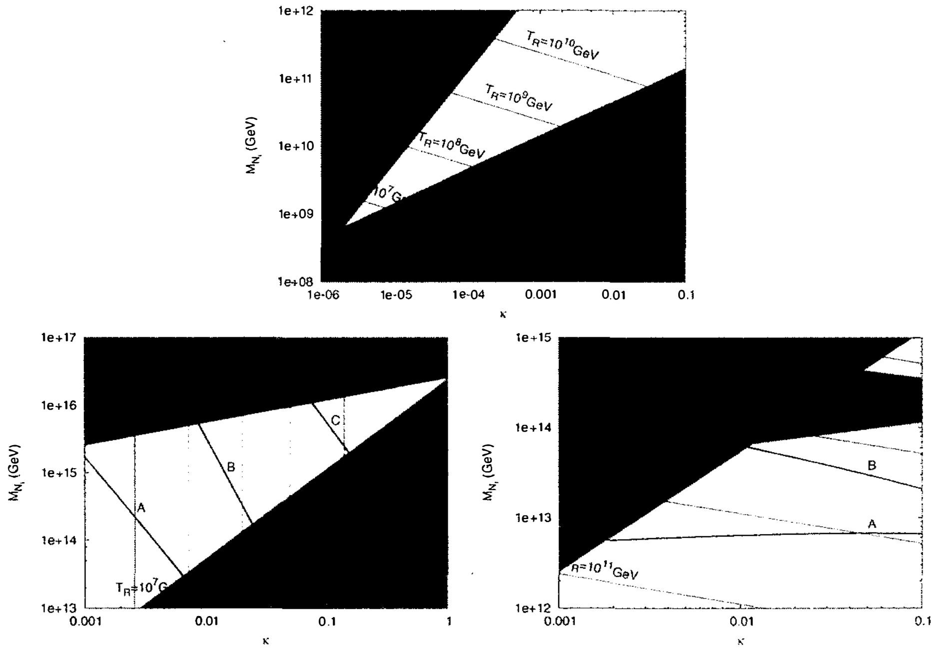

Figure 3: The mass M_{N_i} of the heaviest right-handed neutrino the inflaton can decay into as a function of the coupling κ. The white regions give the measured value of n_B/s. The colored regions are excluded. Top: $M_{N_1} < m_\phi/2$ and $M_{N_1} > T_R$, both strings and inflation contribute non-thermally to ΔB. Bottom: $M_{N_1} > m_\phi/2$, reheating is gravitational and only strings contribute non-thermally to ΔB. Bottom Left: $M_{N_1} > T_R$, there is no wash out. Bottom Right: $M_{N_1} < T_R$, there is also a thermal contribution.

5 Conclusions

GUT which predict massive neutrinos are good candidates for hybrid inflation. Cosmic strings form at the end of inflation and, if they stable down to low energy, they contribute to CMB anisotropies together with inflation. The symmetry breaking scale is constrained by CMB data to the range $M \in [10^{14.5} - 10^{15.5}]$ GeV and the relevant coupling $\kappa \in [10^{-7} - 10^{-2}]$. Scalar perturbations from inflation dominate for a large part of the parameter space and it might be impossible to detect the strings using the temperature anisotropies of the CMB. They could however be detected via the B-type polarization of the CMB [5].

Hybrid inflation predicts a spectral index $n_s \gtrsim 0.98$. It is very difficult to get smaller values except maybe by going to non minimal models [27]; hence if the three year WMAP central value $n_s = 0.951^{+0.015}_{-0.019}$ were to be confirmed, hybrid inflation with minimal SUGRA could be excluded. But even if scalar perturbations are dominated by scalar perturbations from inflation, tensor perturbations (which are negligible for hybrid inflation) can nonetheless be dominated by tensor perturbations from cosmic strings; this can allow a larger value of n_s [4].

Finally, when $B - L$ is broken at the end of inflation, baryogenesis via leptogenesis takes place after inflation during reheating and/or via cosmic strings decay; which of the two scenarios dominates depends upon the various parameters in the inflaton-neutrino sector.

Acknowledgments

I wish to acknowledge Marieke Postma for enjoyable collaboration. I also wish to thank The Dutch Organization for Scientific Research [NWO] for financial support.

References

1. C. B. Netterfield *et al.* [Boomerang Collaboration], Astrophys. J. **571**, 604 (2002) [arXiv:astro-ph/0104460].
2. A. Vilenkin and E. P. S. Shellard, *"Cosmic strings and other topological defects"*, Cambridge monographs on mathematical physics, Cambridge University Press, England, 1994; M. B. Hindmarsh and T. W. B. Kibble, Rept. Prog. Phys. **58** (1995) 477 [arXiv:hep-ph/9411342].
3. L. Pogosian, I. Wasserman and M. Wyman, arXiv:astro-ph/0604141.
4. D. N. Spergel *et al.*, arXiv:astro-ph/0603449.
5. U. Seljak and A. Slosar, arXiv:astro-ph/0604143.
6. A. Linde, Phys. Rev. D **49**, 748 (1994) [arXiv:astro-ph/9307002].
7. G. R. Dvali, Q. Shafi and R. K. Schaefer, Phys. Rev. Lett. **73**, 1886 (1994) [arXiv:hep-ph/9406319].
8. R. Jeannerot, Phys. Rev. D **56** (1997) 6205 [arXiv:hep-ph/9706391].
9. R. Kallosh, arXiv:hep-th/0109168, R. Kallosh and A. Linde, JCAP **0310**, 008 (2003) [arXiv:hep-th/0306058].
10. R. Jeannerot, S. Khalil, G. Lazarides and Q. Shafi, JHEP **0010**, 012 (2000) [arXiv:hep-ph/0002151].
11. G. Lazarides and C. Panagiotakopoulos, Phys. Rev. D **52**, 559 (1995) [arXiv:hep-ph/9506325].
12. R. Jeannerot, J. Rocher and M. Sakellariadou, Phys. Rev. D **68**, 103514 (2003) [arXiv:hep-ph/0308134].
13. R. Jeannerot, arXiv:hep-th/0604214.
14. N. T. Jones, H. Stoica and S. H. H. Tye, JHEP **0207**, 051 (2002) [arXiv:hep-th/0203163].
15. M. Gell-Mann, P. Ramond and R. Slansky, in *Supergravity*, P. van Nieuwenhuizen and D. Freeman eds., North Holland, Amsterdam 1979, p. 315; T. Yanagida, Prog. Theor. Phys. **64** (1980) 1103; R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. **44** (1980) 912; J. Schechter and J. W. F. Valle, Phys. Rev. D **22** (1980) 2227.
16. R. Jeannerot and M. Postma, JHEP **0505** (2005) 071 [arXiv:hep-ph/0503146].
17. R. Jeannerot and M. Postma, arXiv:hep-th/0604216.
18. R. Jeannerot and M. Postma, JCAP **0512** (2005) 006 [arXiv:hep-ph/0507162].
19. R. Allahverdi, K. Enqvist, J. Garcia-Bellido and A. Mazumdar, arXiv:hep-ph/0605035, D. H. Lyth, arXiv:hep-ph/0605283.
20. M. Landriau and E. P. S. Shellard, Phys. Rev. D **69** (2004) 023003 [arXiv:astro-ph/0302166].
21. A. N. Lommen, arXiv:astro-ph/0208572.
22. V. A. Kuzmin, V. A. Rubakov and M. E. Shaposhnikov, Phys. Lett. B **191**, 171 (1987).
23. M. Fukugita and T. Yanagida, Phys. Lett. B **174**, 45 (1986).
24. W. Buchmuller, R. D. Peccei and T. Yanagida, Ann. Rev. Nucl. Part. Sci. **55**, 311 (2005) [arXiv:hep-ph/0502169].
25. R. Jeannerot, Phys. Rev. Lett. **77**, 3292 (1996) [arXiv:hep-ph/9609442].
26. G. Lazarides and Q. Shafi, Phys. Lett. B **258**, 305 (1991). G. Lazarides, Springer Tracts Mod. Phys. **163**, 227 (2000) [arXiv:hep-ph/9904428].
27. M. Bastero-Gil, S. F. King and Q. Shafi, arXiv:hep-ph/0604198.