Hybrid inflation within supersymmetric grand unified theories, as well as inflation through brane collisions within braneworld cosmological models, lead to the formation of one-dimensional defects. Observational data, mainly from the cosmic microwave background temperature anisotropies but also from the gravitational wave background, impose constraints on the free parameters of the models. I review these inflationary models and discuss the constraints from the currently available data.

1 Introduction

The inflationary scenario is with no doubt extremely successful in solving the shortcomings of the standard hot Big Bang model. Inflation consists of a phase of accelerated expansion which took place at a very high energy scale. Inflation is deeply rooted in the basic principles of general relativity and field theory, while when the principles of quantum mechanics are taken into account, it provides a successful explanation for the origin of the large scale structures and the measured temperature anisotropies in the Cosmic Microwave Background (CMB) spectrum.

However, despite its success inflation still remains a paradigm in search of model. One should search for an inflationary model inspired from some fundamental theory and subsequently test its predictions against current data. This offers a beautiful and fruitful example of the interplay between high energy physics and astrophysics/cosmology. Inflation should also prove itself generic, meaning that the onset of inflation should be independent of any particular initial conditions. This issue, already addressed in the past, has been recently re-investigated.

Hybrid inflation is based on Einstein’s gravity but driven by false vacuum. The inflaton field rolls down its potential while another scalar field is trapped in an unstable false vacuum. Once the inflaton field becomes much smaller than some critical value, a phase transition to the
true vacuum takes place signaling the end of inflation. Such phase transition may leave behind topological defects as false vacuum remnants.

In this lecture, I will first discuss hybrid inflation within Supersymmetric Grand Unified Theories (SUSY GUTs) and then within braneworld cosmological models. In both cases one-dimensional topological defects are generically produced at the end of inflation. I will discuss the predictions of these models, in particular regarding the CMB temperature anisotropies spectrum but also with respect to the gravitational wave background, which consequently induce constraints on the parameters space of the models.

2 Inflation within supersymmetric grand unified theories

Theoretically motivated inflationary models can be built in the context of Supersymmetry (SUSY) or Supergravity (SUGRA). $N = 1$ SUSY models contain complex scalar fields which often have flat directions in their potential, thus offering natural candidates for inflationary models. In this framework, hybrid inflation driven by F-terms or D-terms is the standard inflationary model, leading generically to cosmic string formation at the end of inflation. F-term inflation is potentially plagued with the $\eta$-problem, while D-term inflation avoids it.

2.1 F-term inflation

F-term inflation can be naturally accommodated in the framework of GUTs when a GUT gauge group, $G_{GUT}$, is broken down to the Standard Model (SM) gauge group, $G_{SM}$, at an energy scale $M_{GUT}$ according to the scheme

$$G_{GUT} \xrightarrow{M_{GUT}} H_1 \xrightarrow{M_{inf}} H_2 \xrightarrow{F} G_{SM},$$

where $\Phi^+, \Phi^-$ is a pair of GUT Higgs superfields in non-trivial complex conjugate representations, which lower the rank of the group by one unit when acquiring non-zero vacuum expectation value. The inflationary phase takes place at the beginning of the symmetry breaking $H_1 \xrightarrow{M_{inf}} H_2$.

The gauge symmetry is spontaneously broken by adding F-terms to the superpotential. The Higgs mechanism leads generically to Abrikosov-Nielsen-Olesen strings, called F-term strings.

F-term inflation is based on the globally supersymmetric renormalisable superpotential

$$W_{\text{inf}}^F = \kappa S (\Phi^+ \Phi^- - M^2),$$

where $S$ is a GUT gauge singlet left handed superfield and $\kappa$, $M$ are two constants ($M$ has dimensions of mass) which can be taken positive with field redefinition. The scalar potential as a function of the scalar complex component of the respective chiral superfields $\Phi^\pm, S$ reads

$$V(\phi^+, \phi^-, S) = |F_{\Phi^+}|^2 + |F_{\Phi^-}|^2 + |F_S|^2 + \frac{1}{2} \sum_a g_a^2 D_a^2.$$
initially $S \gg M$, the fields quickly settle down the valley of local minima. Since in the slow roll inflationary valley the ground state of the scalar potential is non-zero, supersymmetry is broken. In the tree level, along the inflationary valley the potential is constant, therefore perfectly flat. A slope along the potential can be generated by including the one-loop radiative corrections. Thus, the scalar potential gets a little tilt which helps the inflaton field $S$ to slowly roll down the valley of minima. The one-loop radiative corrections to the scalar potential along the inflationary valley lead to the effective potential

\[
V_{\text{eff}}(|S|) = \kappa^2 M^4 \left\{ 1 + \frac{\kappa^2 N}{32 \pi^2} \left[ 2 \ln \frac{|S|^2 \kappa^2}{\Lambda^2} + (z + 1)^2 \ln(1 + z^{-1}) \right] + (z - 1)^2 \ln(1 - z^{-1}) \right\}; \quad z = \frac{|S|^2}{M^2},
\]

$\Lambda$ is a renormalisation scale and $N$ stands for the dimensionality of the representation to which the complex scalar components $\phi_+, \phi_-$ of the chiral superfields $\Phi_+, \Phi_-$ belong.

### 2.2 D-term inflation

D-term inflation is one of the most interesting and versatile models of inflation. It is possible to implement naturally D-term inflation within high energy physics, as for example SUSY GUTs, SUGRA, or string theories. In addition, D-term inflation avoids the Hubble-induced mass problem. Here, the gauge symmetry is spontaneously broken by introducing Fayet-Iliopoulos (FI) D-terms. In standard D-term inflation, the constant FI term gets compensated by a single complex scalar field at the end of the inflationary era, which implies that standard D-term inflation ends with the formation of cosmic strings, called D-strings. A supersymmetric description of the standard D-term inflation is insufficient; the inflaton field reaches values of the order of the Planck mass, or above it, even if one concentrates only around the last 60 e-folds of inflation; the correct analysis is therefore in the context of supergravity.

Standard D-term inflation requires a scheme

\[
G_{\text{GUT}} \times U(1) \xrightarrow{M_{\text{GUT}}} H \times U(1) \xrightarrow{M_{\text{inf}}} H \rightarrow G_{\text{SM}},
\]

It is based on the superpotential

\[
W = \lambda S \Phi_+ \Phi_-,
\]

where $S, \Phi_+, \Phi_-$ are three chiral superfields and $\lambda$ is the superpotential coupling. It assumes an invariance under an Abelian gauge group $U(1)_{\xi}$, under which the superfields $S, \Phi_+, \Phi_-$ have charges 0, +1 and −1, respectively. It also assumes the existence of a constant FI term $\xi$. In the standard supergravity formulation the Lagrangian depends on the Kähler potential $K(\Phi_i, \bar{\Phi}_i)$ and the superpotential $W(\Phi_i)$ only through the combination given in

\[
G(\Phi_i, \bar{\Phi}_i) = \frac{K(\Phi_i, \bar{\Phi}_i)}{M_{\text{Pl}}^2} + \ln \frac{|W(\Phi_i)|^2}{M_{\text{Pl}}^6}.
\]

However, this standard supergravity formulation is inappropriate to describe D-term inflation. In D-term inflation the superpotential vanishes at the unstable de Sitter vacuum (anywhere else the superpotential is non-zero). Thus, standard supergravity is inappropriate, since the theory is ill-defined at $W = 0$. In conclusion, D-term inflation must be described with a non-singular formulation of supergravity when the superpotential vanishes.

Various formulations of effective supergravity can be constructed from the superconformal field theory. One must first build a Lagrangian with full superconformal theory, and then the
gauge symmetries that are absent in Poincaré supergravity must be gauge fixed. In this way, one can construct a non-singular theory at $W = 0$, where the action depends on all three functions: the Kähler potential $K(\Phi_i, \bar{\Phi}_i)$, the superpotential $W(\Phi_i)$ and the kinetic function $f_{ab}(\Phi_i)$ for the vector multiplets. To construct a formulation of supergravity with constant FI terms from superconformal theory, one finds that under U(1) gauge transformations in the directions in which there are constant FI terms $\xi_\alpha$, the superpotential $W$ must transform as $\delta_\alpha W = \eta_{\alpha i} \partial_i W = -i(g_\xi/M_{Pl}^2) W$; one cannot keep any longer the same charge assignments as in standard supergravity.

D-term inflationary models can be built with different choices of the Kähler geometry. Let us first consider D-term inflation within minimal supergravity. It is based on

$$K_{\text{min}} = \sum_i |\Phi_i|^2 = |\Phi_-|^2 + |\Phi_+|^2 + |S|^2 ,$$

with $f_{ab}(\Phi_i) = \delta_{ab}$. The tree level scalar potential is

$$V_{\text{min}} = \lambda^2 \exp \left( \frac{|\phi_-|^2 + |\phi_+|^2 + |S|^2}{M_{Pl}^2} \right) \left[ |\phi_+\phi_-|^2 \left( 1 + \frac{|S|^4}{M_{Pl}^4} \right) + |\phi_+ S|^2 \left( 1 + \frac{|\phi_-|^4}{M_{Pl}^4} \right) \right] + |\phi_- S|^2 \left( 1 + \frac{|\phi_+|^4}{M_{Pl}^4} \right) + 3 \frac{|\phi_+\phi_- S|^2}{M_{Pl}^2} + \frac{g^2}{2} \left( q_+|\phi_+|^2 + q_-|\phi_-|^2 + \xi \right)^2 ,$$

with $q_\pm = \pm 1 - \xi/(2M_{Pl}^2)$. The potential has two minima: One global minimum at zero and one local minimum equal to $V_0 = (g^2/2)\xi^2$. For arbitrary large $S$ the tree level value of the potential remains constant and equal to $V_0$; the $S$ plays the rôle of the inflaton field. Assuming chaotic initial conditions $|S| \gg S_\text{ini}$ inflation begins. Along the inflationary trajectory the D-term, which is the dominant one, splits the masses in the $\Phi_\pm$ superfields, leading to the one-loop effective potential for the inflaton field. Considering also the one-loop radiative corrections,

$$V_{\text{min}}(\pm S) = \frac{g^2 \xi^2}{2} \left( 1 + \frac{g^2}{16\pi^2} \left[ 2 \ln \left( 1 + \frac{\lambda^2 g^2 \xi |S|^2}{M_{Pl}^4} \right) + f_V(z) \right] \right) ,$$

where

$$f_V(z) = (z + 1)^2 \ln \left( 1 + \frac{1}{z} \right) + (z - 1)^2 \ln \left( 1 - \frac{1}{z} \right) , \quad z \equiv \frac{\lambda^2 g^2 \xi |S|^2}{M_{Pl}^4} \exp \left( \frac{|S|^2}{M_{Pl}^2} \right) .$$

As a second example, consider D-term inflation based on Kähler geometry with shift symmetry,

$$K_{\text{shift}} = \frac{1}{2} (S + \bar{S})^2 + |\phi_+|^2 + |\phi_-|^2 ,$$

and minimal structure for the kinetic function. The scalar potential reads

$$V_{\text{shift}} \simeq \frac{g^2}{2} \left( |\phi_+|^2 - |\phi_-|^2 + \xi \right)^2$$

$$+ \lambda^2 \exp \left( \frac{|\phi_-|^2 + |\phi_+|^2}{M_{Pl}^2} \right) \exp \left[ \frac{(S + \bar{S})^2}{2M_{Pl}^2} \right]$$

$$\times \left[ |\phi_+\phi_-|^2 \left( 1 + \frac{S^2 + \bar{S}^2}{M_{Pl}^4} + \frac{|S|^2|S + \bar{S}|^2}{M_{Pl}^4} \right) + |\phi_+ S|^2 \left( 1 + \frac{|\phi_-|^4}{M_{Pl}^4} \right) \right]$$

$$+ |\phi_- S|^2 \left( 1 + \frac{|\phi_+|^4}{M_{Pl}^4} \right) + 3 \frac{|\phi_+\phi_- S|^2}{M_{Pl}^2} .$$

(13)
As in D-term inflation within minimal SUGRA, the potential has a global minimum at zero for \( \langle \Phi_+ \rangle = 0 \) and \( \langle \Phi_- \rangle = \sqrt{\xi} \eta \) and a local minimum equal to \( V_0 = (g^2/2)\xi^2 \) for \( \langle S \rangle \gg S_c \) and \( \langle \Phi_\pm \rangle = 0 \). The exponential factor \( e^{S/2} \) in minimal SUGRA has been replaced by \( e^{(S+S_c)/2} \). Writing \( S = \eta + i\phi_0 \) one gets \( e^{(S+S_c)/2} = e^{i\eta} \). If \( \eta \) plays the role of the inflaton field, we obtain the same potential as for minimal D-term inflation. If instead \( \phi_0 \) is the inflaton field, the inflationary potential is identical to that of the usual D-term inflation within global SUSY. The latter case is better adapted with the choice \( K_{\text{shift}} \), since in this case the exponential term is constant during inflation, thus it cannot spoil the slow roll conditions.

As a last example, consider a Kähler potential with non-renormalisable terms:

\[
K_{\text{non-renorm}} = |S|^2 + |\Phi_+|^2 + |\Phi_-|^2 + f_+(|S|^2/M^2_{\text{Pl}})|\Phi_+|^2 + f_-|S|^2|\Phi_-|^2 + b|S|^4/M^2_{\text{Pl}}^2, \quad (14)
\]

where \( f_\pm \) are arbitrary functions of \( (|S|^2/M^2_{\text{Pl}}) \) and the superpotential is given in Eq. (6). The effective potential reads

\[
V_{\text{eff}}^{\text{non-renorm}}(|S|) = \frac{g^2\xi^2}{2} \left\{ 1 + \frac{g^2}{16\pi^2} \left[ 2 \ln \left( z^2 \xi/\Lambda^2 \right) + f_V(z) \right] \right\}, \quad (15)
\]

where

\[
f_V(z) = (z + 1)^2 \ln \left( 1 + \frac{1}{z} \right) + (z - 1)^2 \ln \left( 1 - \frac{1}{z} \right) \quad (16)
\]

with

\[
z \equiv \frac{\lambda^2|S|^2}{g^2\xi} \exp \left( \frac{|S|^2}{M^2_{\text{Pl}}} + b \frac{|S|^4}{M^4_{\text{Pl}}} \right) \frac{1}{(1 + f_+)(1 + f_-)}. \quad (17)
\]

3 **Inflation within braneworld cosmologies**

In the context of braneworld cosmology, brane inflation occurs in a similar way as hybrid inflation within supergravity, leading to string-like objects. In string theories, D-brane D-anti-brane annihilation leads generically to the production of lower dimensional D-branes, with D3- and D1-branes (D-strings) being predominant. To sketch brane inflation (for example see ), consider a Dp-/Dp system in the context of IIB string theory. Six of the spatial dimensions are compactified on a torus; all branes move relatively to each other in some directions. A simple and well-motivated inflationary model is brane inflation where the inflaton is simply the position of a Dp-brane moving in the bulk. As two branes approach, the open string modes between the branes develop a tachyon, indicating an instability. The relative Dp-/Dp-brane position is the inflaton field and the inflaton potential comes from their tensions and interactions. Brane inflation ends by a phase transition mediated by open string tachyons. The annihilation of the branes releases the brane tension energy that heats up the universe so that the hot big bang epoch can take place. Since the tachyonic vacuum has a non-trivial \( \pi_1 \) homotopy group, there exist stable tachyonic string solutions with \( (p-2) \) co-dimensions. These daughter branes have all dimensions compact; a four-dimensionnal observer perceives them as one-dimensional objects, the D-strings. Zero-dimensional defects (monopoles) and two-dimensional ones (domain walls), which are cosmologically undesirable, are not produced during brane intersections. Fundamental strings (F-strings) and D-strings that survive the cosmological evolution become cosmic superstrings. In these models, the large compact dimensions and the large warp factors allow cosmic superstring tensions to be in the range between \( 10^{-13} < G\mu < 10^{-6} \), depending on the model. These cosmic superstrings are stable, or at least their lifetime is comparable to the age of the universe, so they can survive to form a cosmic superstring network.
Cosmic superstrings share a number of properties with the cosmic strings, but there are also differences which may lead to distinctive observational signatures. When F- and D-strings meet they can form a three-string junction, with a composite FD-string. IIB string theory allows the existence of bound \((p,q)\) states of \(p\) F-strings and \(q\) D-strings, where \(p\) and \(q\) are coprime. String intersections lead to intercommutation and loop production. For cosmic strings the probability of intercommutation \(P\) is equal to 1, whereas this is not the case for F- and D-strings. Clearly, D-strings can miss each other in the compact dimension, leading to a smaller \(P\), while for F-strings the scattering has to be calculated quantum mechanically since these are quantum mechanical objects. The evolution of cosmic superstring networks has been studied numerically. All studies\(^\text{[13][14][15]}\) conclude that the network will reach a scaling regime.

Cosmic superstrings interact with the standard model particles only via gravity, implying that their detection involves gravitational interactions. Since the particular brane inflationary scenario remains unknown, the tensions of superstrings are only loosely constrained.

4 Compatibility between predictions and data

4.1 CMB temperature anisotropies

Cosmic strings and string-like objects are expected to be generically formed within a large class of models (SUSY, SUGRA and string theories), thus one should consider mixed perturbation models where the dominant rôle is played by the inflaton field and cosmic strings have a small, but not negligible, contribution. Restricting ourselves to the angular power spectrum we remain in the linear regime, where

\[
C_\ell = \alpha C_\ell^I + (1 - \alpha) C_\ell^S; \tag{18}
\]

\(C_\ell^I\) and \(C_\ell^S\) denote the (COBE normalized) Legendre coefficients due to adiabatic inflaton fluctuations and those stemming from the cosmic strings network, respectively. The coefficient \(\alpha\) in Eq. (18) is a free parameter giving the relative amplitude for the two contributions. Comparing the \(C_\ell\), given by Eq. (18), with data obtained from the most recent CMB measurements, one can impose constraints on the parameters space of the models. The upper limit imposed\(^\text{[16][17][18]}\) on the cosmic string contribution to the CMB data depends on the numerical simulation employed in order to calculate the cosmic string power spectrum. In what follows I will not allow cosmic strings to contribute more than 9% to the CMB temperature anisotropies.

F-term inflation

Considering only large angular scales one can get the contributions to the CMB temperature anisotropies analytically. The quadrupole anisotropy has one contribution coming from the inflaton field, calculated using Eq. (4), and one contribution coming from the cosmic strings network. Fixing the number of e-foldings to 60, the inflaton and cosmic strings contribution to the CMB depend on the superpotential coupling \(\kappa\), or equivalently on the symmetry breaking scale \(M\) associated with the inflaton mass scale, which coincides with the string mass scale. The total quadrupole anisotropy has to be normalised to the COBE data. The cosmic strings contribution is consistent with the CMB measurements provided\(^\text{[17]}\)

\[
M \lesssim 2 \times 10^{15} \text{GeV} \quad \Leftrightarrow \quad \kappa \lesssim 7 \times 10^{-7}. \tag{19}
\]

The superpotential coupling \(\kappa\) is also subject to the gravitino constraint which imposes an upper limit to the reheating temperature, to avoid gravitino overproduction. Within the framework of SUSY GUTs and assuming a see-saw mechanism to give rise to massive neutrinos, the inflaton field decays during reheating into pairs of right-handed neutrinos. This constraint
The tuning of the free parameter $\kappa$ can be softened if one allows for the curvaton mechanism. The curvaton is a scalar field that is sub-dominant during the inflationary era as well as at the beginning of the radiation dominated era which follows the inflationary phase. There is no correlation between the primordial fluctuations of the inflaton and curvaton fields. Clearly, within supersymmetric theories such scalar fields are expected to exist. In addition, embedded strings, if they accompany the formation of cosmic strings, they may offer a natural curvaton candidate, provided the decay product of embedded strings gives rise to a scalar field before the onset of inflation. Considering the curvaton scenario, the coupling $\kappa$ is only constrained by the gravitino limit. More precisely, assuming the existence of a curvaton field there is an additional contribution to the temperature anisotropies. The WMAP CMB measurements impose the following limit on the initial value of the curvaton field

$$\psi_{\text{init}} \lesssim 5 \times 10^{13} \left(\frac{\kappa}{10^{-2}}\right) \text{GeV} \quad \text{for} \quad \kappa \in [10^{-6}, 1]. \quad (20)$$

**D-term inflation**

D-term inflation leads to cosmic string formation at the end of the inflationary era. In the case of minimal SUGRA, consistency between CMB measurements and theoretical predictions impose\cite{7} that $g \lesssim 2 \times 10^{-2}$ and $\lambda \lesssim 3 \times 10^{-5}$, which can be expressed as a single constraint on the Fayet-Iliopoulos term $\xi$, namely $\sqrt{\xi} \lesssim 2 \times 10^{15} \text{GeV}$. Note that for minimal D-term inflation one can neglect the corrections introduced by the superconformal origin of supergravity. Within minimal SUGRA the couplings and masses must be fine tuned to achieve compatibility between measurements on the CMB anisotropies and theoretical predictions. The fine tuning on the couplings can be softened if one invokes the curvaton mechanism and provided the initial value of the curvaton field is\cite{9}

$$\psi_{\text{init}} \lesssim 3 \times 10^{14} \left(\frac{g}{10^{-2}}\right) \text{GeV} \quad \text{for} \quad \lambda \in [10^{-1}, 10^{-4}] . \quad (21)$$

The constraints on the couplings remain qualitatively valid in non-minimal SUGRA theories, with the superpotential $W$ given in Eq. (6) and a non-minimal Kähler potential. Let us first consider D-term inflation based on Kähler geometry with shift symmetry. The cosmic string contribution to the CMB anisotropies is dominant, in contradiction with the CMB measurements, unless the superpotential coupling is\cite{10} $\lambda \lesssim 3 \times 10^{-5}$. Also in the case of D-term inflation based on a Kähler potential with non-renormalisable terms, the contribution of cosmic strings dominates if the superpotential coupling $\lambda$ is close to unity. The constraints on $\lambda$ read\cite{10}

$$(0.1 - 5) \times 10^{-8} \leq \lambda \leq (2 - 5) \times 10^{-5} \text{ or, equivalently } \sqrt{\xi} \leq 2 \times 10^{15} \text{ GeV} , \quad (22)$$

implying $G\mu \leq 8.4 \times 10^{-7}$. In conclusion, higher order Kähler potentials do not suppress cosmic string contribution, as it was incorrectly claimed in the literature.

**Brane inflation**

The CMB temperature anisotropies originate from the amplification of quantum fluctuations during inflation, as well as from the cosmic superstring network. Provide the scaling regime of the superstring network is the unique source of the density perturbation, the COBE data imply $G\mu \simeq 10^{-6}$. The latest WMAP data allow at most a 9% contribution from strings, which imply a bound $G\mu \lesssim 7 \times 10^{-7}$. Thus, the cosmic superstrings produced towards the end of inflation in the context of braneworld cosmological models is in agreement with the present CMB data.
Strings/superstrings moving with velocity \( v \) across the sky lead to a shift \( \Delta T/T \approx 8\pi G\mu v \gamma \) in the CMB temperature. Current CMB data may probe \( G\mu \approx 10^{-10} \); this limit may be possible to go down to \( G\mu \approx 10^{-13} \).

4.2 Gravitational wave background

Oscillating cosmic string loops emit\(^{19}\) Gravitational Waves (GW). Long strings are not straight but they have a superimposed wiggly small-scale structure due to string intercommutations, thus they also emit\(^{20}\) GW. Cosmic superstrings can as well generate\(^{21}\) a stochastic GW background. Therefore, provided the emission of gravity waves is the efficient mechanism\(^{22}\) for the decay of string loops, cosmic strings/superstrings could provide a source for the stochastic GW spectrum in the low-frequency band. The stochastic GW spectrum has an almost flat region in the frequency range \( 10^{-8} - 10^{10} \) Hz. Within this window, both ADVANCED LIGO/VIRGO (sensitive at a frequency \( f \sim 10^{2} \) Hz) and LISA (sensitive at \( f \sim 10^{-2} \) Hz) interferometers may have a chance of detectability.

Strongly focused beams of relatively high-frequency GW are emitted by cusps and kinks in oscillating strings/superstrings. The distinctive waveform of the emitted bursts of GW may be the most sensitive test of strings/superstrings. ADVANCED LIGO/VIRGO may detect bursts of GW for values of \( G\mu \) as low as \( 10^{-13} \), and LISA for values down to \( G\mu \geq 10^{-15} \). At this point, I would like to remind to the reader that there is still a number of theoretical uncertainties for the evolution of a string/superstring network.

Recently, they have been imposed limits\(^{23}\) on an isotropic gravitational wave background using pulsar timing observations, which offer a chance of studying low-frequency (in the range between \( 10^{-9} - 10^{-7} \) Hz) gravitational waves. The imposed limit on the energy density of the background per unit logarithmic frequency interval reads \( \Omega_{gw}^{cs}(1/8\text{yr})h^2 \leq 1.9 \times 10^{-8} \) (where \( h \) stands for the dimensionless amplitude in GW bursts). If the source of this isotropic GW background is a cosmic string/superstring network, then it leads to an upper bound on the dimensionless tension of a cosmic string background. Under reasonable assumptions for the string network the upper bound on the string tension reads\(^{23}\) \( G\mu \leq 1.5 \times 10^{-8} \). This is a strongest limit than the one imposed from the CMB temperature anisotropies. Clearly, to achieve compatibility with this constraint, F- and D-term inflation become even more fine tuned, unless one invokes the curvaton mechanism. This limit does not affect cosmic superstrings. However, it has been argued\(^{23}\) that with the full Parkes Pulsar Timing Array (PPTA) project the upper bound will become \( G\mu \leq 5 \times 10^{-12} \), which is directly relevant for cosmic superstrings. In conclusion, the full PPTA will either detect gravity waves from strings or they will rule out a number of models.

5 Conclusions

Cosmic strings and string-like objects are generically formed at the end of inflation in the framework of SUSY GUTs as well as in the context of branewold models. These objects contribute to the CMB temperature anisotropies implying constraints on the parameters of the models. More precisely, by allowing a small but non-negligible contribution of strings to the angular power spectrum of CMB anisotropies, we constrain the couplings of the inflationary models, or equivalently the dimensionless string tension. These models remain compatible with the most current CMB measurements, even when we calculate\(^{24}\) the spectral index. Namely, the inclusion of a sub-dominant string contribution to the large scale power spectrum amplitude of the CMB increases the preferred value for the spectral index. In addition, cosmic strings/superstrings can contribute to the stochastic gravitational background, thus limits can be imposed to the dimensionless string tension; these limits are indeed strongest.
F-/D-term inflation within SUSY GUTs or branewold inflation leading to string-like defects are inflationary models which are definitely not ruled out. Stable strings/superstrings are indeed generically produced in the simplest versions of these models, which then become severely constrained from currently available data. However, one can include additional ingredients so that the one-dimensional defects formed at the end of inflation become unstable. This may be indeed the most realistic approach to model building.

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