The Revival of Cosmic Strings

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Cosmic strings are one-dimensional topological defects which could have been formed in the early stages of our Universe. They triggered a lot of interest, mainly for their cosmological implications: they could offer an alternative to inflation for the generation of density perturbations. It was shown however that cosmic strings lead to inconsistencies with the measurements of the cosmic microwave background temperature anisotropies. The picture is changed recently. It was shown that, on the one hand, cosmic strings can be generically formed in the framework of supersymmetric grand unified theories and that, on the other hand, cosmic superstrings could play the rôle of cosmic strings. There is also some possible observational support. All this lead to a revival of cosmic strings research and this is the topic of my lecture.

1 Introduction

Cosmic strings attracted a lot of interest around the eighties and nineties. They offered an alternative mechanism to cosmological inflation for the generation of the primordial density perturbations leading to the large-scale structure formation one observes. However, towards the turn of the century cosmic strings lost their appeal, since it was shown that they lead to inconsistencies with the Cosmic Microwave Background (CMB) measurements. Nevertheless, the story of cosmic strings does not end here. In the last few years there has been a remarkable revival of the theoretical and observational activity.

In this lecture, I will discuss the present view on the cosmological rôle of cosmic strings. In Section 2, I will discuss aspects of cosmic strings in the framework of Grand Unified Theories (GUTs). I will first analyse the formation and classification of topological as well as embedded defects. I will then briefly discuss the CMB temperature anisotropies and I will compare the predictions of topological defects models with current measurements. I will then conclude that topological defects in general, and cosmic strings in particular, are ruled out as the unique source of density perturbations leading to the observed structure formation. At this point I do not conclude that cosmic strings are ruled out, but I ask instead which are the implications for the models of high energy physics which we employed to construct our cosmological scenario. The first question is whether cosmic strings are expected to be generically formed. I will address this question in the framework of Supersymmetric Grand Unified Theories (SUSY GUTs). I will show that cosmic strings are indeed expected to be generically formed within a large class of models within SUSY GUTs and therefore one has to use mixed models, consisting in inflation with a sub-dominant partner of cosmic strings. I will then examine whether such mixed models are indeed compatible with the CMB data. I will present two well-studied inflationary models within supersymmetric theories, namely F/D-term hybrid inflation. I will impose constraints on the free parameters of the models (masses and couplings) so that there is agreement between theory and measurements. In Section 3, I will address the issue of cosmic superstrings as cosmic strings candidates, in the context of braneworld cosmologies. In Section 4, I will
discuss a candidate of a gravitational lensing by a cosmic string. I will round up with the conclusions in Section 5.

2 Topological Defects

2.1 Topological Defects in GUTs

The Universe has steadily cooled down since the Planck time, leading to a series of Spontaneously Broken Symmetries (SSB). SSB may lead to the creation of topological defects [1, 2], which are false vacuum remnants, such as domain walls, cosmic strings, monopoles, or textures, via the Kibble mechanism [3].

The formation or not of topological defects during phase transitions, followed by SSB, and the determination of the type of the defects, depend on the topology of the vacuum manifold \( \mathcal{M}_n \). The properties of \( \mathcal{M}_n \) are usually described by the \( k \)-th homotopy group \( \pi_k(\mathcal{M}_n) \), which classifies distinct mappings from the \( k \)-dimensional sphere \( S^k \) into the manifold \( \mathcal{M}_n \). To illustrate that, let me consider the symmetry breaking of a group \( G \) down to a subgroup \( H \) of \( G \). If \( \mathcal{M}_n = G/H \) has disconnected components, or equivalently if the order \( k \) of the nontrivial homotopy group is \( k = 0 \), then two-dimensional defects, called domain walls, form. The spacetime dimension \( d \) of the defects is given in terms of the order of the non-trivial homotopy group by \( d = 4 - 1 - k \). If \( \mathcal{M}_n \) is not simply connected, in other words if \( \mathcal{M}_n \) contains loops which cannot be continuously shrunk into a point, then cosmic strings form. A necessary, but not sufficient, condition for the existence of stable strings is that the first homotopy group (the fundamental group) \( \pi_1(\mathcal{M}_n) \) of \( \mathcal{M}_n \), is nontrivial, or multiply connected. Cosmic strings are line-like defects, \( d = 2 \). If \( \mathcal{M}_n \) contains unshrinkable surfaces, then monopoles form, for which \( k = 1, d = 1 \). If \( \mathcal{M}_n \) contains noncontractible three-spheres, then event-like defects, textures, form for which \( k = 3, d = 0 \).

Depending on whether the symmetry is local (gauged) or global (rigid), topological defects are called local or global. The energy of local defects is strongly confined, while the gradient energy of global defects is spread out over the causal horizon at defect formation. Patterns of symmetry breaking which lead to the formation of local monopoles or local domain walls are ruled out, since they should soon dominate the energy density of the Universe and close it, unless an inflationary era took place after their formation. Local textures are insignificant in cosmology since their relative contribution to the energy density of the Universe decreases rapidly with time [4].

Even if the nontrivial topology required for the existence of a defect is absent in a field theory, it may still be possible to have defect-like solutions. Defects may be embedded in such topologically trivial field theories [5]. While stability of topological defects is guaranteed by topology, embedded defects are in general unstable under small perturbations.

2.2 Cosmic Microwave Background Temperature Anisotropies

The CMB temperature anisotropies offer a powerful test for theoretical models aiming at describing the early Universe. The characteristics of the CMB multipole moments can be used to discriminate among theoretical models and to constrain the parameters space.

The spherical harmonic expansion of the CMB temperature anisotropies, as a function of angular position, is given by

\[
\frac{\delta T}{T}(\mathbf{n}) = \sum_{\ell m} a_{\ell m} W_\ell Y_{\ell m}(\mathbf{n}) \quad \text{with} \quad a_{\ell m} = \int d\Omega_n \frac{\delta T}{T}(\mathbf{n}) Y^{*}_{\ell m}(\mathbf{n}) ;
\]

(1)

\( W_\ell \) stands for the \( \ell \)-dependent window function of the particular experiment. The angular power spectrum of CMB temperature anisotropies is expressed in terms of the dimensionless coefficients \( C_\ell \), which appear in the expansion of the angular correlation function in terms of the Legendre polynomials \( P_\ell \):

\[
\left\langle \frac{\delta T}{T}(\mathbf{n}) \frac{\delta T}{T}(\mathbf{n}') \right\rangle |_{(\mathbf{n}, \mathbf{n}' \equiv \cos \vartheta)} = \frac{1}{4\pi} \sum_\ell (2\ell + 1) C_\ell P_\ell(\cos \vartheta) W_\ell^2 .
\]

(2)
It compares points in the sky separated by an angle \( \vartheta \). Here, the brackets denote spatial average, or expectation values if perturbations are quantised. Equation (2) holds only if the initial state for cosmological perturbations of quantum-mechanical origin is the vacuum \([6, 7]\). The value of \( C_\ell \) is determined by fluctuations on angular scales of the order of \( \pi/\ell \). The angular power spectrum of anisotropies observed today is usually given by the power per logarithmic interval in \( \ell \), plotting \( \ell(\ell + 1)C_\ell \) versus \( \ell \).

The predictions of the defects models regarding the characteristics of the CMB spectrum are:

- Global \( \mathcal{O}(4) \) textures lead to a position of the first acoustic peak at \( \ell \approx 350 \) with an amplitude \( \sim 1.5 \) times higher than the Sachs-Wolfe plateau \([8]\).

- Global \( \mathcal{O}(N) \) textures in the large \( N \) limit lead to a quite flat spectrum, with a slow decay after \( \ell \sim 100 \) \([9]\). Similar are the predictions of other global \( \mathcal{O}(N) \) defects \([10, 11]\).

- Local cosmic strings predictions are not very well established and range from an almost flat spectrum \([12]\) to a single wide bump at \( \ell \sim 500 \) \([13]\) with an extremely rapidly decaying tail.

The position and amplitude of the acoustic peaks, as found by the CMB measurements \([14, 15, 16, 17]\), are in disagreement with the predictions of topological defects models. Thus, CMB measurements rule out pure topological defects models as the origin of initial density perturbations leading to the observed structure formation. At this point one has to ask which are the implications for the high energy physics models upon which our cosmological model was built. I will thus first ask whether cosmic strings formation is indeed generic. I will address this question in the framework of SUSY GUTs. I am only interested in cosmic strings, since I consider gauge theories, for which domain walls and monopoles are dangerous, while textures are cosmologically uninteresting \([4]\).

### 2.3 Genericity of Cosmic Strings Formation within SUSY GUTs

I will address the question of whether cosmic strings formation is generic, in the context of SUSY GUTs. Even though the Standard Model (SM) has been tested to a very high precision, it is incapable of explaining neutrino masses \([18, 19, 20]\). An extension of the SM gauge group can be realised within Supersymmetry (SUSY). SUSY offers a solution to the gauge hierarchy problem, while in the supersymmetric standard model the gauge coupling constants of the strong, weak and electromagnetic interactions meet at a single point \( M_{\text{GUT}} \approx (2-3) \times 10^{16} \text{ GeV} \). In addition, SUSY GUTs can provide the scalar field which could drive inflation, explain the matter-antimatter asymmetry of the Universe, and propose a candidate, the lightest superparticle, for cold dark matter.

Within SUSY GUTs there is a large number of SSB patterns leading from a large gauge group \( G \) to the SM gauge group \( G_{\text{SM}} \equiv \text{SU}(3)_C \times \text{SU}(2)_L \times \text{U}(1)_Y \). The study of the homotopy group of the false vacuum for each SSB scheme will determine whether there is defect formation and it will identify the type of the formed defect. Clearly, if there is formation of domain walls or monopoles, one will have to place an era of supersymmetric hybrid inflation to dilute them. To consider a SSB scheme as a successful one, it should be able to explain the matter/anti-matter asymmetry of the Universe and to account for the proton lifetime measurements \([18]\). In what follows, I consider a mechanism of baryogenesis via leptogenesis, which can be thermal or nonthermal one. In the case of nonthermal leptogenesis, \( \text{U}(1)_{B-L} \) (\( B \) and \( L \), are the baryon and lepton numbers, respectively) is a sub-group of the GUT gauge group, \( G_{\text{GUT}} \), and \( B-L \) is broken at the end or after inflation. In the case of thermal leptogenesis, \( B-L \) is broken independently of inflation. If leptogenesis is thermal and \( B-L \) is broken before the inflationary era, then one should check whether the temperature at which \( B-L \) is broken, which will define the mass of the right-handed neutrinos, is smaller than the reheating temperature which should be lower than the limit imposed by the gravitino. To ensure the stability of proton, the discrete symmetry \( \mathbb{Z}_2 \), which is contained in \( \text{U}(1)_{B-L} \), must be kept unbroken down to low energies. This implies that the successful SSB schemes should end at \( G_{\text{SM}} \times \mathbb{Z}_2 \). I will then examine how often cosmic strings have survived at the inflationary era, within all acceptable SSB patterns.
To accomplish this task one has to choose the large gauge group $G_{\text{GUT}}$. In Ref. [21] this study has been done explicitly for a large number of simple Lie groups. Since I consider GUTs based on simple gauge groups, the type of supersymmetric hybrid inflation will be of the F-type. The minimum rank of $G_{\text{GUT}}$ has to be at least equal to 4, to contain the $G_{\text{SM}}$ as a subgroup. Then one has to study the possible embeddings of $G_{\text{SM}}$ in $G_{\text{GUT}}$ to be in agreement with the SM phenomenology and especially with the hypercharges of the known particles. Moreover, the group must include a complex representation, needed to describe the SM fermions, and it must be anomaly free. Since, in principle, $SU(n)$ may not be anomaly free, I assume that the $SU(n)$ groups which I use have indeed a fermionic representation that certifies that the model is anomaly free. I set as the upper bound on the rank $r$ of the group, $r \leq 8$. Clearly, the choice of the maximum rank is in principle arbitrary. This choice could, in a sense, be motivated by the Horava-Witten [22] model, based on $E_8 \times E_8$. Thus, the large gauge group $G_{\text{GUT}}$ could be one of the following: SO(10), $E_6$, SO(14), SU(8), SU(9); flipped SU(5) and $[SU(3)]^3$ are included within this list as subgroups of SO(10) and $E_6$, respectively.

A detailed study of all the SSB schemes which bring us from $G_{\text{GUT}}$ down to the SM gauge group $G_{\text{SM}}$, by one or more intermediate steps, shows that cosmic strings are generically formed at the end of hybrid inflation. If the large gauge group $G_{\text{GUT}}$ is SO(10) then cosmic strings formation is unavoidable [21]. For $E_6$ it depends whether one considers thermal or nonthermal leptogenesis. More precisely, under the assumption of nonthermal leptogenesis then cosmic strings formation is unavoidable. If I consider thermal leptogenesis then cosmic strings formation at the end of hybrid inflation arises in 98% of the acceptable SSB schemes [23]. If the requirement of having $Z_2$ unbroken down to low energies is relaxed and thermal leptogenesis is considered as being the mechanism for baryogenesis, then cosmic strings formation accompanies hybrid inflation in 80% of the SSB schemes [23]. The SSB schemes of SU(6) and SU(7) down to the $G_{\text{SM}}$ which could accommodate an inflationary era with no defect (of any kind) at later times are inconsistent with proton lifetime measurements and minimal SU(6) and SU(7) do not predict neutrino masses [21], implying that these models are incompatible with high energy physics phenomenology. Higher rank groups, namely SO(14), SU(8) and SU(9), should in general lead to cosmic strings formation at the end of hybrid inflation. In all these schemes, cosmic strings formation is sometimes accompanied by the formation of embedded strings. The strings which form at the end of hybrid inflation have a mass which is proportional to the inflationary scale.

### 2.4 Mixed Models

Since cosmic strings are expected to be generically formed in the context of SUSY GUTs, one should consider mixed perturbation models where the dominant rôle is played by the inflaton field but cosmic strings have also a contribution, small but not negligible. Restricting ourselves to the angular power spectrum, we can remain in the linear regime. In this case,

$$C_\ell = \alpha C_{\ell}^I + (1 - \alpha)C_{\ell}^S,$$  

(3)

where $C_{\ell}^I$ 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. (3) is a free parameter giving the relative amplitude for the two contributions. Comparing the $C_{\ell}$, given by Eq. (3), with data obtained from the most recent CMB measurements, one gets that a cosmic strings contribution to the primordial fluctuations higher than 14% is excluded up to 95% confidence level [23, 25, 26]. In what follows, I will be on the conservative side and I will not allow cosmic strings to contribute more than 10% to the CMB temperature anisotropies.

### 2.5 Supersymmetric Hybrid Inflation

Inflation remains the most appealing scenario for describing the early Universe. Inflation essentially consists of a phase of accelerated expansion which took place at a very high energy scale. However, despite its
success, it faces a number of questions, as for example how generic is the onset of inflation \cite{27} and how one can guarantee a natural and successful inflationary model. Unfortunately, inflation is still a paradigm in search of a model. I will discuss two well-studied inflationary models in the framework of supersymmetry, namely F/D-term inflation.

2.5.1 F-term Inflation

F-term inflation can be naturally accommodated in the framework of GUTs when a GUT gauge group $G_{\text{GUT}}$ is broken down to the GSM at an energy $M_{\text{GUT}}$ according to the scheme

$$G_{\text{GUT}} \xrightarrow{M_{\text{GUT}}} H_1 \xrightarrow{M_{\text{infl}}} H_2 \rightarrow \text{GSM} ;$$

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

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

$$W^F_{\text{infl}} = \kappa S(\Phi_+ \Phi_- - M^2) ,$$

where $S$ is a GUT gauge singlet left handed superfield, $\Phi_+$ and $\Phi_-$ are defined above; $\kappa$ and $M$ are two constants ($M$ has dimensions of mass) which can be taken positive with field redefinition. The chiral superfields $S, \Phi_+, \Phi_-$ are taken to have canonical kinetic terms. This superpotential is the most general one consistent with an R-symmetry under which $W \rightarrow e^{i\beta} W, \Phi_+ \rightarrow e^{-i\beta} \Phi_+, \Phi_- \rightarrow e^{i\beta} \Phi_-$, and $S \rightarrow e^{i\beta} S$. An R-symmetry can ensure that the rest of the renormalisable terms are either absent or irrelevant.

The scalar potential reads

$$V(\phi_+, \phi_-, S) = |F_{\phi_+}|^2 + |F_{\phi_-}|^2 + |F_S|^2 + \frac{1}{2} \sum_a g_a^2 D_a^2 .$$

The F-term is such that $F_{\phi_i} = |\partial W/\partial \Phi_i|_{\theta = 0}$, where we take the scalar component of the superfields once we differentiate with respect to $\Phi_i = \Phi_+, \Phi_-, S$. The D-terms are

$$D_a = \tilde{\phi}_i (T_a)^i_j \phi^j + \xi_a ,$$

with $a$ the label of the gauge group generators $T_a, g_a$ the gauge coupling, and $\xi_a$ the Fayet-Iliopoulos term. By definition, in the F-term inflation the real constant $\xi_a$ is zero; it can only be nonzero if $T_a$ generates an extra U(1) group. In the context of F-term hybrid inflation, the F-terms give rise to the inflationary potential energy density, while the D-terms are flat along the inflationary trajectory, thus one may neglect them during inflation.

The potential has one valley of local minima, $V = \kappa^2 M^4$, for $S > M$ with $\phi_+ = \phi_- = 0$, and one global supersymmetric minimum, $V = 0$, at $S = 0$ and $\phi_+ = \phi_-_ = M$. Imposing 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 nonzero, SUSY 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 an effective potential $\cite{28, 29, 30, 31}$

$$V^F_{\text{eff}}(|S|) = \kappa^2 M^4 \left\{ 1 + \frac{\kappa^2 N}{32 \pi^2} \left[ 2 \ln \frac{|S|^2}{M^2} + \left( \frac{|S|^2}{M^2} + 1 \right)^2 \ln \left( 1 + \frac{M^2}{|S|^2} \right) + \left( \frac{|S|^2}{M^2} - 1 \right)^2 \ln \left( 1 - \frac{M^2}{|S|^2} \right) \right] \right\} ;$$
Λ 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.

Considering only large angular scales, one can get the contributions to the CMB temperature anisotropies analytically. In Ref. [31], the Sachs-Wolfe effect has been explicitly calculated. The quadrupole anisotropy has one contribution coming from the inflaton field, calculated using Eq. (8), and one contribution coming from the cosmic strings network, given by numerical simulations [32]. Fixing the number of e-foldings to 60, then for a given gauge group \( G_{\text{GUT}} \), 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. In Ref. [31] we have found that the cosmic strings contribution is consistent with the CMB measurements, provided

\[
M \lesssim 2 \times 10^{15} \text{GeV} \Leftrightarrow \kappa \lesssim 7 \times 10^{-7}. \tag{9}
\]

This constraint on \( \kappa \) is in agreement with the one found in Ref. [33]. Strictly speaking the above condition was found in the context of SO(10) gauge group, but the conditions imposed in the context of other gauge groups are of the same order of magnitude since \( M \) is a slowly varying function of the dimensionality \( N \) of the representations to which the scalar components of the chiral Higgs superfields belong.

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 on the reheating temperature can be converted to a constraint on the parameter \( \kappa \). The gravitino constraint on \( \kappa \) reads [31] \( \kappa \lesssim 8 \times 10^{-3} \), which is a weaker constraint.

Concluding, F-term inflation leads generically to cosmic strings formation at the end of the inflationary era. The cosmic strings formed are of the GUT scale. This class of models can be compatible with CMB measurements, provided the superpotential coupling is smaller than \( 10^{-6} \). This tuning of the free parameter \( \kappa \) can be softened if one allows for the curvaton mechanism. According to the curvaton mechanism [34, 35], another scalar field, called the curvaton, could generate the initial density perturbations whereas the inflaton field is only responsible for the dynamics of the Universe. 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 [31] the following limit on the initial value of the curvaton field

\[
\psi_{\text{init}} \lesssim 5 \times 10^{13} \left( \frac{\kappa}{10^{-2}} \right) \text{GeV}, \tag{10}
\]

provided the parameter \( \kappa \) is in the range \([10^{-6}, 1]\).

2.5.2 D-term Inflation

D-term inflation received a lot of interest, mainly because it is not plagued by the Hubble-induced mass problem, and in addition it can be easily implemented in string theory. D-term inflation is derived from the superpotential

\[
W_{\text{init}}^D = \lambda S \Phi_+ \Phi_-; \tag{11}
\]
$S$, $\Phi_+$, $\Phi_-$ are three chiral superfields and $\lambda$ is the superpotential coupling. D-term inflation requires the existence of a nonzero Fayet-Iliopoulos term $\xi$, which can be added to the Lagrangian only in the presence of an extra U(1) gauge symmetry, under which the three chiral superfields have charges $Q_S = 0$, $Q_{\Phi_+} = +1$ and $Q_{\Phi_-} = -1$, respectively. This extra U(1) gauge symmetry symmetry can be of a different origin; hereafter we consider a nonanomalous U(1) gauge symmetry. Thus, D-term inflation requires a scheme, like

$$G_{\text{GUT}} \times U(1) \xrightarrow{M_{\text{GUT}}/\Phi_+ \Phi_-} H \xrightarrow{M_{\text{SM}}} H \rightarrow \text{GSM}.$$  \hfill (12)

The symmetry breaking at the end of the inflationary phase implies that cosmic strings are always formed at the end of D-term hybrid inflation. To avoid cosmic strings, several mechanisms have been proposed which either consider more complicated models or require additional ingredients. For example, one can add a nonrenormalisable term in the potential [36], or add an additional discrete symmetry [37], or consider GUT models based on nonsimple groups [38], or introduce a new pair of charged superfields [39] so that cosmic strings formation is avoided within D-term inflation. In what follows, I will show that standard D-term inflation leading to cosmic strings production is still compatible with CMB data since cosmic strings contribution to the CMB data is not constant nor dominant. This implies that one does not have to invoke some new physics. The reader can find a detailed study in Ref. [31, 40].

In the global supersymmetric limit, Eqs. (6), (11) lead to the following expression for the scalar potential

$$V^D(\phi_+, \phi_-, S) = \lambda^2 \left[ |S|^2 (|\phi_+|^2 + |\phi_-|^2) + |\phi_+\phi_-|^2 \right] + \frac{g^2}{2} (|\phi_+|^2 - |\phi_-|^2 + \xi)^2,$$  \hfill (13)

where $g$ is the gauge coupling of the U(1) symmetry and $\xi$ is a Fayet-Iliopoulos term, chosen to be positive.

In D-term inflation, as opposed to F-term inflation, the inflaton mass acquires values of the order of Planck mass, and therefore, the correct analysis must be done in the framework of SUGRA. The SSB of SUSY in the inflationary valley introduces a splitting in the masses of the components of the chiral superfields $\Phi_\pm$. As a result, we obtain [40] two scalars with squared masses $m_+^2 = \lambda^2 |S|^2 \exp \left( (|S|^2/M_{\text{Pl}}^2) \pm g^2 \xi \right)$ and a Dirac fermion with squared mass $m_f^2 = \lambda^2 |S|^2 \exp \left( (|S|^2/M_{\text{Pl}}^2) \right)$. Thus, calculating the radiative corrections, the effective scalar potential for minimal supergravity reads [31, 40]

$$V_{\text{eff}} = \frac{g^2 \xi^2}{2} \left[ 1 + \frac{g^2}{16 \pi^2} \right] \left[ 2 \ln \frac{|S|^2 \lambda^2}{\lambda^2} e^{\frac{|S|^2}{\lambda^2}} + \left( \frac{\lambda^2 |S|^2}{g^2 \xi} e^{\frac{|S|^2}{g^2 \xi}} + 1 \right)^2 \ln \left( 1 + \frac{g^2 \xi e^{-\frac{|S|^2}{g^2 \xi}}}{\lambda^2 |S|^2} \right) \right] \left( \frac{\lambda^2 |S|^2}{g^2 \xi} e^{\frac{|S|^2}{g^2 \xi}} - 1 \right) \ln \left( 1 - \frac{g^2 \xi e^{-\frac{|S|^2}{g^2 \xi}}}{\lambda^2 |S|^2} \right) \left( \frac{\lambda^2 |S|^2}{g^2 \xi} e^{\frac{|S|^2}{g^2 \xi}} - 1 \right) \ln \left( 1 - \frac{g^2 \xi e^{-\frac{|S|^2}{g^2 \xi}}}{\lambda^2 |S|^2} \right) \right] \right] \right] \right].$$  \hfill (14)

In Refs. [31, 40], we have properly addressed the question of cosmic strings contribution to the CMB data and we got that standard D-term inflation can be compatible with measurements; the cosmic strings contribution to the CMB is actually model-dependent. Our most important finding was that cosmic strings contribution is not constant, nor is it always dominant.

More precisely, we obtained [31, 40] that $g \gtrsim 2 \times 10^{-2}$ is incompatible with the allowed cosmic strings contribution to the WMAP measurements. For $g \lesssim 2 \times 10^{-2}$, the constraint on the superpotential coupling $\lambda$ reads $\lambda \lesssim 3 \times 10^{-5}$. SUGRA corrections impose in addition a lower limit to $\lambda$. The constraints induced on the couplings by the CMB measurements can be expressed [31, 40] as a single constraint on the Fayet-Iliopoulos term $\xi$, namely $\sqrt{\xi} \lesssim 2 \times 10^{15} \text{ GeV}$.

Concluding, standard D-term inflation always leads to cosmic strings formation at the end of the inflationary era. The cosmic strings formed are of the GUT scale. This class of models is still compatible with CMB measurements, provided the couplings are small enough. As in the case of F-term inflation the fine
tuning of the couplings can be softened provided one considers the curvaton mechanism. In this case, the imposed CMB constraint on the initial value of the curvaton field reads [31, 40]

\[ \psi_{\text{init}} \lesssim 3 \times 10^{14} \left( \frac{g}{10^{-2}} \right) \text{GeV}, \]  

for \( \lambda \in [10^{-1}, 10^{-4}] \).

Our conclusions are still valid in the revised version of D-term inflation, in the framework of SUGRA with constant Fayet-Iliopoulos terms. In the context of N=1, 3+1 SUGRA, the presence of constant Fayet-Iliopoulos terms shows up in covariant derivatives of all fermions. In addition, since the relevant local U(1) symmetry is a gauged R-symmetry [41], the constant Fayet-Iliopoulos terms also show up in the supersymmetry transformation laws. In Ref. [42] there were presented all corrections of order \( g \xi / M_{\text{Pl}}^2 \) to the classical SUGRA action required by local supersymmetry. Under U(1) gauge transformations in the directions in which there are constant Fayet-Iliopoulos terms \( \xi \), the superpotential must transform as [41]

\[ \delta W = -i \frac{g \xi}{M_{\text{Pl}}^2} W, \]  

otherwise the constant Fayet-Iliopoulos term \( \xi \) vanishes. This requirement is consistent with the fact that in the gauge theory at \( M_{\text{Pl}} \to \infty \) the potential is U(1) invariant. To promote the simple SUSY D-term inflation model, Eq. (11), to SUGRA with constant Fayet-Iliopoulos terms, one has to change the charge assignments for the chiral superfields, so that the superpotential transforms under local R-symmetry [42]. In SUSY, the D-term potential is neutral under U(1) symmetry, while in SUGRA the total charge of \( \Phi^{\pm} \) fields does not vanish but is equal to \( -\xi / M_{\text{Pl}}^2 \). More precisely, the D-term contribution to the scalar potential \( V \) [see Eq. (13)], should be replaced by \( (g^2/2)(q_+|\phi_+|^2 + q_-|\phi_-|^2 + \xi)^2 \) where

\[ q_\pm = \pm 1 - \rho_\pm \frac{\xi}{M_{\text{Pl}}^2} \quad \text{with} \quad \rho_+ + \rho_- = 1. \]  

In addition, the squared masses of the scalar components \( \phi_\pm \) become

\[ m_\pm^2 = \lambda^2 |S|^2 \exp \left( |S|^2 / M_{\text{Pl}}^2 \right) \pm g^2 \xi q_\pm ; \]  

the Dirac fermion mass remains unchanged. However, since for the limits we imposed on the Fayet-Iliopoulos term \( \xi \), the correction \( \xi / M_{\text{Pl}}^2 \) is \( \sim 10^{-6} \), I conclude that our results also also valid in the revised version of D-term inflation within SUGRA.

### 3 Superstrings as Cosmic Strings Candidates

In the context of perturbative string theory, superstrings of cosmic size were excluded, mainly because they should have too large a tension. More precisely, perturbative strings have a tension close to the Planck scale, producing CMB inhomogeneities far larger than observed. Moreover, since the scale of their tension also exceeds the upper bound on the energy scale of the inflationary vacuum, such strings could have only been produced before inflation, and therefore diluted. In addition, there are instabilities that would prevent such long strings from surviving on cosmic time scales [43]. Thus, for years was a clear distinction between fundamental strings and cosmic strings.

Recently, this whole picture has changed. In addition to the fundamental F-strings, there are also D-strings as a particular case of higher-dimensional Dp-branes (D stands for Dirichlet and p denotes the dimensionality of the brane), partially wrapped on compact cycles resulting to only one noncompact dimension. In the braneworld approach, our Universe represents a D3-brane on which open strings can end [44], embedded in a higher dimensional space, called the bulk. Brane interactions can unwind and evaporate higher dimensional Dp-branes so that we are left with D3-branes embedded in a higher dimensional...
bulk; one of these D3-branes plays the rôle of our Universe. Since gauge charges are attached to the ends of strings, gauge particles and fermions can propagate only along the D3-branes while gravitons (and dilatons, ...) which are closed string modes can move in the bulk. Since gravity has been proved only down to scales of about 0.1 mm, the dimensions of the bulk can be much larger than the string scale. In the braneworld context, the extra dimensions can even be infinite, if the geometry is nontrivial. Large extra dimensions can be employed to address the hierarchy problem, a result which lead to an increasing interest in braneworld scenarios. Apart from the Dp-branes, there are also antibranes, which differ from the Dp-branes by having an equal and opposite conserved Ramond-Ramond charge, which implies an attractive force between them.

Braneworld cosmology can also offer a natural inflationary scenario. Assuming the early Universe contained an extra brane and antibrane, then an inflationary era could be driven by the potential between the two branes, while the separation between the branes will play the rôle of the inflaton. The inflaton potential is rather flat when the branes are separated and steepens as they approach, until at some point a field becomes tachyonic, which indicates an instability, leading to a rapid brane-antibrane annihilation.

D-brane-antibrane inflation leads to the abundant production of lower dimensional D-branes that are one-dimensional in the noncompact directions. Luckily, zero-dimensional defects (monopoles) and two-dimensional ones (domain walls), which are cosmologically undesirable, are not produced. In these models, the large compact dimensions and the large warp factors allow cosmic superstring tensions to be in the range between \(10^{-11} < G_\mu < 10^{-9}\), depending on the model.

Cosmic superstrings share a number of properties with the cosmic strings, but there are also differences which may lead to distinctive observational signatures. 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 collisions between all possible pairs of superstrings have been studied in string perturbation theory. For F-strings, the reconnection probability is of the order of \(g^2\), where \(g\) stands for the string coupling. For F-F string collisions, it was found that the reconnection probability \(P\) is \(10^{-3} \lesssim P \lesssim 1\). For D-D string collisions, one has \(10^{-1} \lesssim P \lesssim 1\). Finally, for F-D string collisions, the reconnection probability can take any value between 0 and 1. These results have been confirmed by a quantum calculation of the reconnection probability for colliding D-strings. Similarly, the string self-intersection probability is reduced. Moreover, when D- and F-strings meet they can form a three-string junction, with a composite DF-string. It is also possible in IIB string theory to have bound states of p F-strings and q D-strings, where p and q are coprime. This leads to the question of whether there are frozen networks dominating the matter content of the Universe, or whether scaling solutions can be achieved.

To study the evolution of cosmic superstrings, I have performed numerical simulations of independent stochastic networks of D- and F-strings. I found that the characteristic length scale \(\xi\), which gives the typical distance between the nearest string segments and the typical curvature of strings, grows linearly with time

\[
\xi(t) \propto \zeta t ,
\]

where the slope \(\zeta\) depends on the reconnection probability \(P\), and on the energy of the smallest allowed loops (i.e., the energy cutoff). For reconnection (or intercommuting) probability in the range \(10^{-3} \lesssim P \lesssim 0.3\), I found that

\[
\zeta \propto \sqrt{P} \Rightarrow \xi(t) \propto \sqrt{P} t ,
\]

in agreement with my old results. I thus disagree with the statement that \(\xi(t) \propto P t\). In Ref. it is claimed that the energy density of long strings \(\rho_t\) evolves as \(\dot{\rho}_t = 2(\dot{a}/a)\rho_t - P(\rho_t/\xi)\), where \(H = \dot{a}/a\) is the Hubble constant. Then substituting the ansatz \(\xi(t) = \gamma(t)t\), the authors of Ref. obtain
\[
\dot{\gamma} = -\frac{1}{2t}(\gamma - \mathcal{P}),
\]
during the radiation-dominated era. Since this equation has a stable fixed point at \(\gamma(t) = \mathcal{P}\), the authors state [53] that \(\xi \simeq \mathcal{P}t\). My disagreement with Ref. [53] is based on the fact that intersections between two long strings is not the most efficient mechanism for energy loss of the string network. The possible string intersections can be divided into three possible cases: (i) two long strings collide in one point and exchange partners with intercommuting probability \(P_1\); (ii) two strings collide in two points and exchange partners chopping off a small loop with intercommuting probability \(P_2^1\); and (iii) one long string self-intersects in one point and chops off a loop with intercommuting probability \(P_2^2\), which in general is different than \(P_1\). Clearly, only cases (ii) and (iii) lead to a closed loop formation and therefore remove energy from the long string network. Between cases (ii) and (iii), only case (iii) is an efficient way of forming loops and therefore dissipating energy. I have checked numerically [51] that case (iii) appears more often than case (ii), and besides, case (ii) has in general a smaller probability, since one expects that \(P_1 \sim P_2\). However, the heuristic argument employed in Ref. [53] does not refer to self-string intersections (i.e., case (iii)); it only applies to intersections between two long strings. This is clear since intersections between two long strings depend on the string velocity, however self-string intersections should not depend on how fast the string moves. In other words, a string can intersect itself even if it does not move but it just oscillates locally.

Studying the time evolution of the slope \(\zeta\), I found [51] that it reaches a constant value at relatively the same time \(t\) for various values of \(\mathcal{P}\), which implies that the long strings reach scaling. This result has been confirmed by studying numerically the behavior of a network of interacting Dirichlet-fundamental strings \((p, q)\) in Ref. [54]. To model \((p, q)\) strings arising from compactifications of type IIB string theory, the authors studied [54] the evolution of nonabelian string networks. The positive element of such nonabelian networks is that they contain multiple vertices where many different types of string join together. Such networks have the potential of leading to a string-dominated Universe due to tangled networks of interacting \((p, q)\) strings that freeze. It was shown [54] that such freezing does not take place and the network reaches a scaling limit. In this field theory approach however strings are not allowed to have different tensions, which is a characteristic property of cosmic superstrings. Recently, this has been done in the context of modelling \((p, q)\) cosmic superstrings [55]. It was found that such networks rapidly approach a stable scaling solution, where once scaling is reached, only a small number of the lowest tension states is populated substantially. An interesting question is to find out whether the field theory approach of Ref. [54] mimics the results of the modelling approach of Ref. [55].

The cosmic superstring network is characterised [51] by two components: there are a few long strings with a scale-invariant evolution; the characteristic curvature radius of long strings, as well as the typical separation between two long strings are both comparable to the horizon size, \(\xi(t) \simeq \sqrt{\mathcal{P}}t\), and there is a large number of small closed loops having sizes \(\ll t\). Assuming there are string interactions, the long strings network will reach an asymptotic energy density where

\[
\rho_l = \frac{\mu}{\mathcal{P}t^2}.
\]

Thus, for fixed linear mass density, the cosmic superstring energy density may be higher than the field theory case, but at most only by one order of magnitude. More precisely, the fraction of the total density in the form of strings in the radiation-dominated era reads

\[
\frac{\rho_{\text{str}}}{\rho_{\text{total}}} = \frac{32\pi G\mu}{3\mathcal{P}}.
\]

Oscillating string loops loose energy by emitting graviton, dilaton and Ramond-Ramond (RR) fields. Accelerated cosmic strings are sources of gravitational radiation, in particular from the vicinity of the cusps where the string velocity approaches the speed of light. Similarly, cosmic superstrings emit gravity waves but since the intercommutation probability is less than unity, their network is denser with more cusps, resulting in an enhancement of the emitted gravitational radiation. As it was pointed out [56], the
gravitational wave bursts emitted from cusps of oscillating string or superstring loops could be detectable with the gravitational-wave interferometers LIGO/VIRGO and LISA.

One can place constraints on the energy scale of cosmic strings from the observational bounds on dilaton decays. Considering that the dilaton lifetime is in the range $10^7 s \lesssim \tau \lesssim \tau_{\text{dec}}$, I obtained an upper bound $\eta \lesssim \mathcal{P}^{-1/3} \lesssim 10^{11} \text{GeV}$ for the energy scale of cosmic superstrings, which determines the critical temperature for the transition leading to string formation. A lower reconnection probability allows a higher energy scale of strings, at most by one order of magnitude.

4 Cosmic Strings in the Sky

As a theoretician, I believe that it is of great importance to get observational support for the existence of cosmic strings. Unfortunately, up to recently the attempts to find cosmic strings in the sky were unsuccessful.

A Russian-Italian collaboration claims to have found the first signature of a cosmic string in the sky. More precisely, the authors of Refs. [58, 59, 60] point out that the peculiar properties of the gravitational lens CSL-1 (Capodimonte - Sternberg Lens Candidate no.1) could be only explained as the first case of lensing by a cosmic string. CSL-1, found in the OACDF (Osservatorio Astronomico di Capodimonte - Deep Field) consists of two identical images, separated by $1.9''$. The two sources have very similar morphology, namely they consist or a bright nucleus surrounded by a faint halo with undistorted and quite circular isophotes. The most relevant feature of these images is indeed that their isophotes appear to be undistorted. The performed photometric and spectroscopic analysis [58, 60] revealed that both the two components of CSL-1 are giant elliptical galaxies at redshift $z = 0.46$. The possibility that CSL-1 could be interpreted as the projection of two giant elliptical galaxies, identical at a $99\%$ confidence level, has been disregarded as unlikely. Moreover, the peculiar properties of CSL-1 cannot be explained in terms of lensing by a compact lens model, since a usual gravitational lens created by a bound clump of matter lead to inhomogeneous gravitational fields which always distort background extended images. Thus, the most favorite explanation of CSL-1 in the framework of gravitational lensing is, according to the authors of Refs. [58, 59], that of lensing by a cosmic string. Assuming that CSL-1 is indeed the first lensing by a cosmic string then the observed separation of the two images corresponds to a particular value for the deficit angle which implies that $G\mu > 4 \times 10^{-7}$; for kinky strings $G\mu$ could be less. As it was recently pointed out [61], high string velocities enhance lensing effects by a factor $1/\sqrt{1-v^2}$, where $v$ stands for the string velocity. This decreases the lower bound on $G\mu$ placed by CSL-1.

In Ref. [59], the authors study the statistics of lens candidates in the vicinity of CSL-1. They claim that they expect 7-9 lens candidates, which is a relatively high number with respect to the one expected from normal gravitational lens statistics. This excess of gravitational lens candidates in the neighborhood of CSL-1 is claimed to be compatible with the proposed in Ref. [58] cosmic string scenario. As the authors however state only once we have spectroscopic studies of these candidates, we will be able to extract robust conclusions.

It is crucial to confirm or infirm this finding by further and independent studies.

5 Conclusions

In this lecture I presented the story of cosmic strings, as we know it at present. Cosmic strings are expected to be generically formed in a large class of models based on SUSY GUTs. If the predictions of cosmic strings are inconsistent with the various measurements, then either the theories which predict the formation of cosmic strings are altogether wrong, or the models have to be more complicated to avoid strings formation. Luckily, neither is needed. The free parameters of the models can be constrained so that there is agreement between predictions and measurements. Cosmological inflation is an attractive model with however too many possible choices. It is crucial to find out which are the natural inflationary models and
to constrain their free parameters. Therefore, the constraints imposed by the cosmological implications of cosmic strings are indeed important. The recent proposal that cosmic superstrings can be considered as cosmic strings candidates opens new perspectives on the theoretical point of view. Therefore, even though cosmic strings cannot play a dominant rôle in structure formation, one has to consider them as a sub-dominant partner of inflation. The possible observational support which was announced recently is of course a major issue and cosmic strings have just entered a new flourishing era.

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