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Cosmic F- and D-strings

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Abstract: Macroscopic fundamental and Dirichlet strings have several potential instabilities: breakage, tachyon decays, and confinement by axion domain walls. We investigate the conditions under which metastable strings can exist, and we find that such strings are present in many models. There are various possibilities, the most notable being a network of (p, q) strings. Cosmic strings give a potentially large window into string physics.

Keywords: Superstrings and Heterotic Strings, D-branes, Cosmology of Theories beyond the SM.

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1. Introduction

Before the 'second superstring revolution' there appeared to be a clear distinction between fundamental strings and cosmic strings. Fundamental strings were believed to have tensions μ close to the Planck scale. In perturbative heterotic string theory, for example, $G\mu = \alpha_{\rm GUT}/16\pi \gtrsim 10^{-3}$, whereas the isotropy of the cosmic microwave background implied (even before COBE) that any string of cosmic size must have $G\mu \lesssim 10^{-5}$ [1]. Thus any cosmic strings would have had to arise in the low energy effective field theory, as magnetic or electric flux tubes.

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In principle, fundamental strings could have been produced in the early universe and then grown to macroscopic size with the expansion of the universe. Inflation provides a simple explanation for the absence of cosmic fundamental strings of such high tension. However, even aside from inflation it was noted in ref. [2] that there are effects that would prevent the appearance of cosmic fundamental strings. Macroscopic type-I strings break up on a stringy time scale into short open strings, and so would never form. Macroscopic heterotic strings always appear as boundaries of axion domain walls, whose tension would force the strings to collapse rather than grow to cosmic scales [3]. At the time of ref. [2] no instability of long type-II strings was known, but it is now clear that NS5-brane instantons [4] (in combination with supersymmetry breaking to lift the zero modes) will produce an axion potential and so lead to domain walls.

Today the situation in string theory is much richer. First, many new one-dimensional objects are known: in addition to the fundamental F-strings, there are D-strings, as well as higher dimensional D-, NS-, M-branes that are partially wrapped on compact cycles so that only one noncompact dimension remains. Second, the possibility of large compact dimensions [5] and large warp factors [6] allows much lower tensions for these strings. Third, the various string-string and string-field dualities relate these objects to each other, and to the field-theoretic flux tubes, so that they are actually the same object as it appears in different parts of parameter space. Thus it is important to revisit this subject, and ask whether some of these strings may be cosmologically interesting.

Indeed, it has been argued by Jones, Sarangi, and Tye and by Stoica and Tye [7] that D-brane-antibrane inflation [8] leads to the copious production of lower-dimensional D-branes that are one-dimensional in the noncompact directions. This is a special case of the production of strings in hybrid inflation [9]. Refs. [7] also make the important observation that zero-dimensional defects (monopoles) and two-dimensional defects (domain walls) are not produced; either of these would have led to severe difficulties.

It is necessary then to investigate the stability of possible cosmic strings against the processes noted above. A naive extrapolation of the results of ref. [2] would suggest that all BPS strings are confined by domain walls, and that all non-BPS strings are unstable to breakage or tachyon decay. We will find that the situation is more interesting, and that stable strings exist in certain classes of models but not in others.

In section 2 we investigate this subject, and identify conditions under which long strings can be at least metastable. Our focus is on type-I/II theories, though the same principles will hold in heterotic and M theory compactifications. In section 3 we apply our conditions in the string theory inflation model of Kachru, Kallosh, Linde, Maldacena, McAllister, and Trivedi ($\mathbb{KLMT}$) [10], which is based on the stabilization of all moduli in the warped IIB framework of ref. [11]. We find that the nature of the cosmic strings in this model depends on precisely how the Standard Model fields and the moduli stabilization are introduced. We identify three possibilities: (a) no strings; (b) D1-branes only (or fundamental strings only); (c) (p, q) strings — bound states of p fundamental strings and q D-strings for relatively prime (p, q) — with an upper bound on p. In section 4 we briefly discuss large compact dimension models without large warp factors, and find a similar range of possibilities. In section 5 we discuss the observational signatures of these cosmic strings. Although the

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various bounds currently are all in the area of $G\mu \lesssim 10^{-6}$, there are future observations that will reach many orders of magnitude further in $G\mu$. With (p, q) strings there are more complicated string networks than when only one type of string is present. This may enhance the signatures for these strings, and possibly place strong constraints on models of these types.

While we were completing this work we learned of papers on related subjects by Dvali, Kallosh and Van Proeyen [12] and by Dvali and Vilenkin [13].

2. Stability conditions

2.1 Breakage

2.1.1 Breakage on an M^4-filling brane

The prototype example of string breakage is the type-I string. On a long type-I string there is a constant rate per unit length for the string to break by formation of a pair of endpoints [14]. By this process the string will rapidly convert to ordinary quanta as in figure 1. The modern interpretation is that these endpoints are attached to a spacetime-filling D9-brane. Similarly, the fundamental type-II string (F-string) can end on any D-brane, and so in any type-II model with Dp-branes that fill the noncompact M^4 dimensions there will be an amplitude for the string to break. Of course, for $p < 9$ the D-brane does not fill the *compact* dimensions, and there is the possibility that the breaking is suppressed due to transverse separation between the string and the D-brane. This will allow metastable strings to exist in some models.

By S-duality from F1/D3, a D-string can end on a D3-brane [15, 16] and so will be unstable in compactifications with D3-branes. A D-string can also end on a D5-brane, but actually this does not represent an instability. When an F- or D-string ends on a D3-brane, the endpoints are electric or magnetic sources on the brane. The endpoint pair that is produced when a string breaks are still connected by flux lines in the brane — one can think of the string as dissolving into the brane and becoming flux. For a D-string oriented in the 1-direction, the flux is one unit of integrated F_{23}, which can lower its energy toward zero by spreading in the 23-plane as in figure 2. Similarly an F-string dissolving in any D-brane becomes one unit of F_{01} which can again spread. For a D-string dissolving in a D5-brane however, there is one unit of $F \wedge F$ in the four transverse directions, and the energy of this flux is not lowered by spreading out — this is the scale invariance of the instanton action. In other words, there is still a nonzero energy per unit length, which is not surprising given that the parallel D1/D5 system already saturates the BPS bound.[1] So

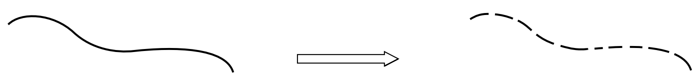

Figure 1: A long open string converts to short open strings on a stringy time scale.

[1] The exact scale invariance will be broken as a consequence of supersymmetry breaking, but there will still be a positive minimum to the string tension.

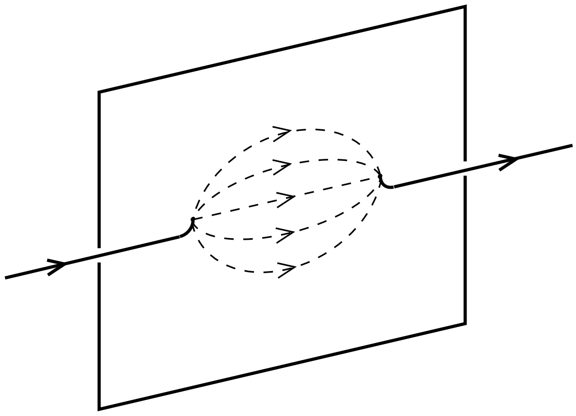

Figure 2: When a string ends on a brane, flux lines in the brane still connect the endpoints.

from the four-dimensional point of view the string does not break, and we might say that while a D-string can *attach* to a D5-brane it cannot *break* on one. Similarly it actually costs energy for a D-string to attach to a D7-brane, and so again the string cannot break.

As another example, a string which arises from a D3-brane wrapped on a 2-cycle can end on a D5-brane which wraps the same cycle and fills the noncompact dimensions. Other cases can be obtained from these by S- and T-duality.

In summary, stability requires that there be no M^4-filling branes on which the string can end, or that the decay be suppressed by transverse separation. We will also encounter a related instability, the tachyonic decay of an unstable D-brane. This is similar to breakage in that the D-brane turns into ordinary quanta on a stringy time scale by decays everywhere along its length. An unstable D-brane in superstring theory is constructed from a D-$\overline{\mathrm{D}}$ pair (for a review see ref. [17]), and the tachyon is an open string stretching between them. As with breakage, this decay can also be suppressed by transverse separation, in this case of the D and the $\overline{\mathrm{D}}$.

2.1.2 Breakage on a 'baryon'

The U(1) gauge field on the M^4-filling brane plays an essential role in allowing the strings to break. Strings couple to forms, for example

$$\int_{\mathrm{F}_1} B_{(2)}\,, \qquad \int_{\mathrm{D}_1} C_{(2)} \tag{2.1}$$

for an F-string or D-string respectively, where the integrals run over the string world-volumes.[2] For a world-volume without boundary, these actions are invariant under the form gauge transformations

$$\delta B_{(2)} = d\lambda_{(1)}\,, \qquad \delta C_{(2)} = d\chi_{(1)}\,, \tag{2.2}$$

[2]Our notation is that $C_{(q)}$, $B_{(2)}$, and $A_{(1)}$ are the R-R, NS-NS, and brane gauge potentials, and $F_{(q+1)}$, $H_{(3)}$, and $G_{(2)}$ the corresponding field strengths. A tilde on a field strength denotes the inclusion of Chern-Simons terms. A four-dimensional form obtained by integrating a higher-rank form over a compact submanifold is denoted by square brackets, e.g. $C_{[q]}$.

but when the string breaks there is a surface term in the transformation. The point is that the form field also couples to the U(1) flux, and one sees from figure 2 that this provides the necessary continuity. The relevant terms in the brane action are

$$S' = \int_{\rm D3} \frac{1}{2}|G_{(2)} + B_{(2)}|^2 + G_{(2)} \wedge C_{(2)} \,. \tag{2.3}$$

Here $G_{(2)} = dA_{(1)}$ is the gauge field strength on the brane, and we have included its Chern-Simons couplings to the forms. We take the example of the D3-brane, and normalizations are simplified for clarity. The variation of S' under the form transformations is

$$\delta S' = -\int_{\rm D3} \lambda_{(1)} d{*}(G_{(2)} + B_{(2)}) + \chi_{(1)} dG_{(2)} \,. \tag{2.4}$$

The F1-brane endpoint is an electric source for $G_{(2)} + B_{(2)}$, and the D-brane endpoint is a magnetic source for $G_{(2)}$, each giving a delta function in one of the terms in the variation (2.4). Thus the variation of the actions (2.1) and (2.3) cancel when a string breaks on a brane. The apparent asymmetry between the forms $B_{(2)}$ and $C_{(2)}$ in the action (2.3) can be reversed by an electric-magnetic duality transformation on the brane.

Type-II theories also have U(1) gauge fields arising from the bulk RR fields, and these can play the same role. Consider the IIB theory, for example. The four-form potential with one index along M^4 and three in the compact directions gives rise to a four-dimensional gauge field for each three-cycle $\mathcal{K}_3$

$$C_{[1]} = \int_{\mathcal{K}_3} C_{(4)} \,. \tag{2.5}$$

The five-form field strength includes a Chern-Simons term $\tilde{F}_{(5)} = dC_{(4)} + B_{(2)} \wedge F_{(3)}$. One is generally interested in compactifications in which some of the background three-form fluxes are nonvanishing. If, for example,

$$\int_{\mathcal{K}_3} F_{(3)} = M \,, \tag{2.6}$$

then the effective two-dimensional field strength is

$$\tilde{F}_{[2]} = \int_{\mathcal{K}_3} \tilde{F}_{(5)} = dC_{[1]} + MB_{(2)} \,. \tag{2.7}$$

Comparing with the first term in the action (2.3), we see that we have the necessary coupling to allow the F1-brane to break; the role of the integer M will be seen below.

To complete the picture we need the mechanism by which the string breaks. Consider a D3-brane wrapped on the same cycle $\mathcal{K}_3$. This is a localized particle in four dimensions, which we refer to loosely as a baryon, because this is the role it plays in gauge/string duality [18]. The gauge field equation on this brane's world-volume is $d * G_{(2)} = F_{(3)}$, again following from the action (2.3), and this is inconsistent with the flux (2.6) unless M F1-branes end on the baryon. We conclude that precisely if $M = 1$ the F1-brane can break by pair production of baryons. If $M \geq 2$ then instead the baryon is a vertex at which M F-strings meet.

More generally, if

$$\int_{\mathcal{K}_3} F_{(3)} = M\,, \qquad \int_{\mathcal{K}_3} H_{(3)} = M' \tag{2.8}$$

then the baryon is a vertex at which M F1-branes and M' D1-branes end.

2.2 Axion domain walls

Any BPS p-brane (that is, the infinite flat brane is a BPS state) must couple to a $(p+1)$-form field $C_{(p+1)}$ (we are using the notation for an R-R brane, but the same applies to NS-NS and M theory branes). This gives the repulsive force that offsets the gravitational and dilatonic attraction between a pair of parallel branes. Let the brane be wrapped on a $(p-1)$-cycle $\mathcal{K}_{p-1}$, thus leaving a string in the noncompact directions. In four dimensions there will be a two-form potential

$$C_{[2]} = \int_{\mathcal{K}_{p-1}} C_{(p+1)}\,, \tag{2.9}$$

which couples to the world-volume of the string. In the four-dimensional theory this is dual to an axion field ϕ,

$$dC_{[2]} = *_4 d\phi + \text{source terms}\,. \tag{2.10}$$

The string is an electric source for $C_{[2]}$ and so a magnetic, topological, source for ϕ. That is, the axion field (appropriately normalized) changes by 2π in going around the string:

$$\oint_C dx \cdot \partial\phi = 2\pi \tag{2.11}$$

on any contour C encircling the string.

Now consider a euclidean $(6-p)$-brane instanton, which couples magnetically to $C_{(p+1)}$, wrapping a $(7-p)$-cycle $\mathcal{K}_{7-p}$ that intersects $\mathcal{K}_{p-1}$ once. A magnetic source for $C_{(p+1)}$ is an electric source for ϕ, and so the instanton amplitude is proportional to $e^{i\phi}$. Since all supersymmetries are ultimately broken, the fermion zero modes in the instanton amplitude are lifted, and this produces a periodic potential for ϕ. From eq. (2.11) it follows that ϕ cannot sit in the minimum of this potential everywhere as we encircle the string — there is a kink where it changes by 2π and passes over the maximum of the potential. Since the kink in ϕ intersects any contour C that circles the string, it defines a domain wall ending on the string. Unless the domain wall tension is exceedingly small this will cause the strings to rapidly disappear [3].

2.3 Two puzzles

2.3.1 The first puzzle

A string cannot both break and bound a domain wall, because the boundary of a boundary is zero. What happens if both instabilities appear to be present? Consider the case of a D1-brane which can end on a D3-brane and which also couples to a massless four-dimensional field $C_{(2)}$. The effective four-dimensional action is

$$\frac{1}{2}\int |dC_{(2)}|^2 + |G_{(2)}|^2 + G_{(2)} \wedge C_{(2)}\,. \tag{2.12}$$

We are writing explicitly only the terms that govern the local dynamics, not the source terms. For a higher dimensional wrapped D-brane one just replaces $C_{(2)}$ by $C_{[2]}$, and for an F-string by $B_{(2)}$ (after a D3 world-volume electric/magnetic duality). The effective four-dimensional field equation for the massless field $C_{(2)}$ is

$$d{*_4}dC_{(2)} = G_{(2)} \;, \tag{2.13}$$

which implies that the axion field must be defined

$$dC_{(2)} = *_4(d\phi + A_{(1)})\,. \tag{2.14}$$

It follows that ϕ transforms nonlinearly under gauge transformations of the brane gauge field $A_{(1)}$,

$$\delta A_{(1)} = d\lambda\,, \qquad \delta\phi = -\lambda\,, \tag{2.15}$$

and so the U(1) on the brane is Higgsed. There is no domain wall because the axion is removed by the Higgs mechanism.[3]

Let us look in more detail at the external field of a D-string running in the 1-direction, again in the effective four-dimensional picture. The nonzero field components are C_{01} and G_{23}. The field equation for $A_{(1)}$,

$$d(*G_{(2)} + C_{(2)}) = 0 \tag{2.16}$$

implies that $C_{01} - G_{23}$ is constant in the transverse directions, and so vanishes. The field equation (2.13) with source term from the D-string then implies that

$$\partial_\perp^2 C_{01} = C_{01} + \delta^2(x_\perp)\,. \tag{2.17}$$

It follows that the R-R field of the string falls exponentially at infinity (the mass is 1 because the constants have been dropped). The R-R field of the D-string is screened, so there is no obstruction to the D-string breaking. The D-string dissolves in the D3-brane to become a tube of G_{23} flux; the total flux in this tube is equal and opposite to the screening flux, and so the G_{23} field dissolves into ordinary quanta.

2.3.2 The second puzzle

We have concluded that there is no stable D-string. However, there is a spontaneously broken U(1) gauge symmetry (2.15), and so one might have expected a topologically stable string associated with this broken symmetry [20]. Such a string would involve degrees of freedom beyond those that we are considering, namely the 'radial' Higgs field associated with the angular part ϕ, and so a higher energy scale; whether such strings exist depends on the topology of this higher-energy field space.

[3] The instanton amplitude $e^{i\phi}$ would seem to violate gauge invariance, but this is balanced by the production of charged strings. Consider for example a D1 string which can end on a D3-brane and whose axion couples to a euclidean D5 instanton. The instanton and D3-brane intersect in a spacetime point, and so the string creation process is essentially the same as for a D0-brane passing through a D8-brane [19].

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Unlike D-strings, such 'ϕ-strings' would not generically be produced in brane-antibrane inflation. Ref. [7] shows that the production of strings can be understood by applying the Kibble argument [21] to the vacuum manifold for the brane-antibrane tachyon, as in K-theory [17, 22]. Add in a D3-$\overline{\rm D3}$ pair, such as would be present during inflation. If we first ignore the Higgsing by ϕ, the vacuum manifold is

$$\frac{\mathrm{U}(2) \times \mathrm{U}(1)}{\mathrm{U}(1) \times \mathrm{U}(1)} = S^3 \,. \tag{2.18}$$

The breaking is the same as the Weinberg-Salam model, $\mathrm{SU}(2) \times \mathrm{U}(1) \to \mathrm{U}(1)$. The vacuum manifold is simply connected, so there is no stable D-string, reflecting the fact that it can break on the surviving D3-brane.

Now include the the nonlinear Higgsing (2.15). This breaks one of the denominator U(1)'s. However, the important point is that this U(1) is broken even before the brane annihilation, by the same nonlinear mechanism: one of the numerator U(1)'s is broken as well. Thus the vacuum manifold for D3-$\overline{\rm D3}$ annihilation is still $\mathrm{U}(2)/\,\mathrm{U}(1) = S^3$ and so the Kibble mechanism does not lead to strings. Any strings associated with this already-broken symmetry will have been diluted away during inflation, and the nonlinear Higgs field will be ordered over long distances. Of course, if the number of e-foldings is sufficiently small, or if additional phase transitions take place after inflation, then there may be interesting strings beyond those that we consider.

3. The $\mathbb{KLMT}$ model

3.1 Generalities

We now apply the above analysis to the $\mathbb{KLMT}$ model [10]. This is presently the most well-developed model of inflation in string theory, being based on a framework in which all moduli are stabilized [11]. Although many open issues remain, some of which will be encountered below, this provides an excellent test case for cosmic strings.

The $\mathbb{KLMT}$ model is based on IIB string theory on a Calabi-Yau manifold. This is orientifolded by a $\mathbb{Z}_2$ symmetry that has isolated fixed points, which become O3-planes.[4] The spacetime metric is warped,

$$ds^2 = e^{2A(x_\perp)} \eta_{\mu\nu} dx^\mu dx^\nu + ds_\perp^2 \,. \tag{3.1}$$

The inflaton is the separation between a D3-brane and an anti-D3-brane, whose annihilation leads to reheating. The annihilation occurs in a region (throat) of large gravitational redshift, $\min\{e^{A(x_\perp)}\} = e^{A_0} \ll 1$, where $e^{A(x_\perp)}$ is normalized to be $O(1)$ in the bulk of the Calabi-Yau. Note that in figure 3 the covering manifold has two throats which are identified under the $\mathbb{Z}_2$ rather than a single throat identified with itself; this is equivalent to the statement that there is no O3-plane in the throat.

[4]We focus on this special case of the more general F theory construction, but as we note below the results in the latter case are much the same.

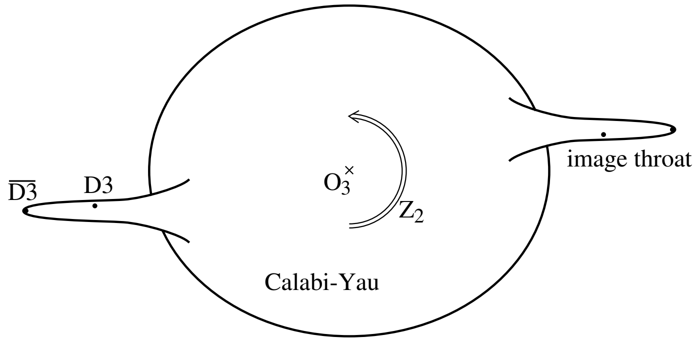

Figure 3: Schematic picture of the $\mathbb{KLMT}$ geometry: a warped Calabi-Yau manifold with throats, identified under a $\mathbb{Z}_2$ orientifold.

The redshift in the throat plays a key role: both the inflationary scale and the scale of string tension, as measured by a ten-dimensional inertial observer, are set by string physics and are close to the four-dimensional Planck scale. The corresponding energy scales as measured by a four-dimensional physicist are then suppressed by a factor of e^{A_0}.

As discussed in ref. [7], one expects copious production only for objects that are one-dimensional in the noncompact dimensions and lie entirely within the region of reheating. The obvious candidates are then the F1-brane (fundamental IIB string) and D1-brane, localized in the throat. There is also another possibility, which however can be quickly disposed of. In the $\mathbb{KLMT}$ model, the throat is a Klebanov-Strassler geometry [23], whose cross section is topologically $S^2 \times S^3$. A D3-brane wrapped on the S^2 also gives a string in four dimensions. However, the S^2 is topologically trivial, it collapses to a point at the end of the throat [23], and so this string can break rapidly.

The D1-branes can be regarded as topological defects in the tachyon field that describes D3-$\overline{\text{D3}}$ annihilation [17, 22], and so these will be produced by the Kibble mechanism [21]. The F1-branes do not have a classical description in these same variables, but in an S-dual description they are topological defects and so must be produced in the same way. Of course, only one of the S-dual descriptions can be quantitatively valid, and if the string coupling is of order one then neither is. However, the Kibble argument depends only on causality and so it is probably valid for both kinds of string in all regimes (it has also been argued in ref. [24] that fundamental strings will be created). All strings created at the end of inflation are at the bottom of the inflationary throat, and they remain there because they are in a deep potential well: their effective four-dimensional tension μ depends on the warp factor at their location and their ten-dimensional tension $\bar{\mu}$ as

$$\mu = e^{2A(x_\perp)}\bar{\mu}\,. \tag{3.2}$$

Given both F1-branes and D1-branes, there will also be (p, q) strings for relatively prime integers p and q. These were found in ref. [25] using the $\mathrm{SL}(2, \mathbb{Z})$ duality of the IIB superstring theory, and are now interpreted [26, 27] as bound states of p F1-branes and q D1-branes. Their tension in the ten-dimensional theory is

$$\bar{\mu}_{p,q} = \frac{1}{2\pi\alpha'}\sqrt{(p - Cq)^2 + e^{-2\Phi}q^2}\,, \tag{3.3}$$

where C is the Ramond-Ramond scalar (normalized to periodicity 1) and Φ is the dilaton ($e^{\Phi} = g_{\mathrm{s}}$), both evaluated at the location of the string. In the $\mathbb{KLMT}$ model the values of these fields are fixed in terms of discrete three-form fluxes.

It remains to discuss the stability of these strings. From the previous section, it is essential to know what branes remain in the theory after inflation, upon which the (p, q) strings might break. In the $\mathbb{KLMT}$ model there must be such branes for two reasons: (a) there must be the branes on which the Standard Model fields live, and (b) the moduli stabilization in this model involves one or more anti-D3-branes located in a throat (these could possibly be the same as the Standard Model branes).

3.2 Scenarios

As we will explain, the only branes that are relevant for the stability of the cosmic strings are those that are located in, or intersect, the inflationary throat. We thus consider the various possibilities.

3.2.1 No branes in the throat

Here we consider the case that the Standard Model and stabilizing branes are located outside the inflationary throat. Let us first ask whether the (p, q) strings are BPS. In fact they are not. The easiest way to see this is to consider the special case of the $T^6/\mathbb{Z}_2$ orientifold, which is T-dual to the type-I theory. The D1-branes become type-I D7-branes, which are non-BPS [22]. The type-I orientation reversal turns a IIB D7-brane into a $\overline{\mathrm{D7}}$-brane, so the type-I D7-brane is actually a bound state of a IIB D7-$\overline{\mathrm{D7}}$ pair. The F1-branes become type-I fundamental strings, which are again non-BPS. Another way to understand this is through the orientifold projection, under which the forms $B_{\mu\nu}$ and $C_{\mu\nu}$ that couple to the F1- and D1-branes are odd [28]. Thus the zero modes of these fields are removed, and do not appear in the massless spectrum.[5]

The orientifold projection in the $\mathbb{KLMT}$ model turns a D1-brane into an anti-D1-brane, which is to say it reverses the orientation. The four-dimensional D-string is thus in this construction a D1-$\overline{\mathrm{D1}}$ bound state, a potentially unstable configuration. We have noted above that the D1 is confined to the tip of the inflationary potential. The key point is then that the image $\overline{\mathrm{D1}}$ is in the image throat. The length of each throat is somewhat greater than unity in string units, so the D1-$\overline{\mathrm{D1}}$ strings are stretched and nontachyonic.

Even without a tachyon it is possible for the D1-string to fluctuate into the other well and annihilate with the $\overline{\mathrm{D1}}$. In order for the D-string to be an interesting cosmic string the decay rate must be suppressed from its natural string scale, of order α'^{-1} per unit string length and time, to a rate of order the cosmic scale H^2 or less, a suppression of roughly 10^{50} in time scale and so 10^{100} in rate per unit length. In fact the suppression is much greater than this, due to the warping of the compact space. The decay proceeds through the appearance of a hole in the euclidean world-sheet, in which the D1 and $\overline{\mathrm{D1}}$ have annihilated. At the edge of this hole the D1-brane crosses over the to the image throat

[5]This will also be true in the more general F theory constructions. The zero modes of $B_{\mu\nu}$ and $C_{\mu\nu}$ are inconsistent with the nontrivial $\mathrm{SL}(2, \mathbb{Z})$ holonomy of the 7-branes.

and annihilates with the $\overline{\mathrm{D1}}$. The rate is given by the usual Schwinger expression for this instanton,

$$e^{-B}, \qquad B = \frac{\pi\sigma^2}{\rho}. \tag{3.4}$$

Here ρ is the action per unit area, i.e. the tension $e^{2A_0}/2\pi\alpha' g_{\mathrm{s}}$. The parameter σ is the action per unit length for the boundary of the hole. Since the D1-brane must pass through the unwarped bulk of the Calabi-Yau, this is of order $R/2\pi\alpha' g_{\mathrm{s}}$ where R is the distance between the throat and its image. The dominant factor in B is from the warping,

$$B \sim e^{-2A_0} \sim 10^8, \tag{3.5}$$

where the numerical value is taken from ref. [10] and will be discussed further in section 5. The warping suppresses the decay by the impressive factor e^{-B}, even though the D1 and its image are only a few string units apart![6] The decay of the F-string and all the other (p, q) strings are similarly stabilized, because they involve world-sheets that stretch from one throat to its image, through the unwarped region.

We can now see why the branes outside the throat are irrelevant. In order for a string to break on one of these, it must fluctuate out of the throat, which again costs the suppression (3.5).

Finally, let us consider decay via baryon pair production. The only relevant baryons are the D3-branes that wrap the S^3 in the Klebanov-Strassler throat. All other three-cycles pass through the bulk of the Calabi-Yau, and so the masses of the corresponding baryon are at the string scale, unsuppressed by the warp factor. For these the pair production rate is then suppressed by the same factor (3.5).

The S^3 in the throat carries M units of $F_{(3)}$ [23, 28]. It follows that if $M = 1$ the F-strings are unstable to baryon production. More generally, a baryon allows a (p, q) string to decay to a $(p - M, q)$ string, and so the stable strings have $|p| \leq M/2$.

In the $\mathbb{KLMT}$ model there is a lower bound on M. During inflation there is at least one $\overline{\mathrm{D3}}$ in the inflationary throat. It is shown in ref. [29] that this is classically unstable unless

$$M \gtrsim 12. \tag{3.6}$$

Thus the baryon-mediated decay will not destabilize the (p, q) strings with small p, and there is still a rich spectrum.

The numerical value of the tunnelling rate is again given by eq. (3.4), where now σ is the baryon mass

$$\sigma = \overline{T}_{\mathrm{D3}} V_{S^3}, \qquad \overline{T}_{\mathrm{D3}} = \frac{1}{(2\pi)^3 \alpha'^2 g_{\mathrm{s}}}, \tag{3.7}$$

and

$$\rho = \bar{\mu}_{p,q} - \bar{\mu}_{p-M,q}. \tag{3.8}$$

[6] The radial metric in the throat goes as dr^2/r^2, so the length is only logarithmic in $e^{A_0} \sim r_{\min}/r_{\max}$.

Obtaining the volume of the S^3 from eq. (65) of ref. [30], one obtains for $(p, q) \to (p - M, q)$ with $p > M/2$ a decay rate of order

$$e^{-B}, \qquad B \sim 0.2 \times \frac{qM^2}{2p - M}; \tag{3.9}$$

we have taken $g_s \ll 1$ to simplify the expression. There is no suppression from the warp factor because the decay takes place entirely in the throat, and the decay will be rapid on a cosmological time scale ($B < 200$) for a wide range of parameters. For example, if $q = 1$ and $2p - M = O(M)$, the decay is rapid for $M < 10^3$. The baryon is a 'bead' at which a (p, q) string and a $(p - M, q)$ string join, and it will accelerate rapidly in the direction of the higher-tension string.

Incidentally, a process similar to the pair production of baryons is the decay of the inflationary $\overline{\mathrm{D3}}$ by an NS5-brane instanton wrapped on the S^3 [29, 31]. This is negligible on cosmological time scales unless M is very close to the bound (3.6).

In summary, in this case we obtain (p, q) strings for relatively prime p, q with $|p| \leq M/2$.

3.2.2 Stabilizing $\overline{\mathrm{D3}}$-branes in the throat

Now suppose that the stabilizing $\overline{\mathrm{D3}}$-brane(s) are in the same throat in which inflation occurs. In other words, inflation is driven by N D3-branes and $\Delta + N$ $\overline{\mathrm{D3}}$-branes for some $\Delta > 0$, so that after annihilation Δ $\overline{\mathrm{D3}}$-branes remain. In this case the F-strings and D-strings can both break, as can all the (p, q) strings, and there are no cosmic strings. In terms of K theory, the tachyon vacuum manifold is

$$\frac{\mathrm{U}(N) \times \mathrm{U}(\Delta + N)}{\mathrm{U}(N) \times \mathrm{U}(\Delta)}. \tag{3.10}$$

This contains a U(1) and supports D-string vortices only if the number Δ of residual $\overline{\mathrm{D3}}$-branes vanishes.

The various scenarios that we are exploring must satisfy a number of nontrivial constraints, including stability of the weak scale, stability of the moduli during inflation, and sufficient reheating [32]. We will not attempt to discuss these in detail here, but it is worth noting that the present scenario has a serious, probably fatal, problem with reheating. The point is that there is a U(1) gauge field on the stabilizing $\overline{\mathrm{D3}}$. This is the only massless degree of freedom in the inflationary throat, and so the only one that couples directly to the inflationary fields. Then almost all of the energy at reheating goes into these U(1) gauge bosons rather than into the Standard Model fields.

Of course, many of the specific features of the $\mathbb{KLMT}$ model may be absent in more general constructions. The role of the $\overline{\mathrm{D3}}$-branes is to break supersymmetry, raising the supersymmetric AdS vacuum to a state of approximately zero cosmological constant [11]. It seems likely that there are other dynamical mechanisms that would accomplish the same thing,[7] and that there need not be an associated U(1) field in general. This eliminates the immediate problem with reheating (though this is still a nontrivial issue to explore).

[7] We thank S. Kachru and E. Silverstein for comments on this point.

However, the alternate stabilizing mechanism no longer mediates breaking of (p, q) strings, as this also requires a U(1) gauge field as in figure 2. Thus we would obtain the same strings as in the no-brane case.

3.2.3 Standard model branes in the throat

Now let us consider the Standard Model branes. In the $\mathbb{KLMT}$ model it is natural to introduce D3-branes and/or $\overline{\text{D3}}$-branes, as well as D7-branes in the F theory constructions. Ref. [33] gives constructions of the Standard Model in terms of local configurations with these ingredients. These local constructions can then be incorporated into the $\mathbb{KLMT}$ model. In order to reduce the supersymmetry of the low energy theory from $\mathcal{N} = 4$ to $\mathcal{N} = 1$ or $\mathcal{N} = 0$ it is necessary that the D3-branes (or $\overline{\text{D3}}$-branes) be fixed at an appropriate singularity, for example an orbifold fixed point. The D7-branes are needed for anomaly cancellation, and they must also pass through this singularity. In some cases (certain non-abelian orbifolds) the D7-branes are not needed [33, 34].

Consider first the case that the 3-branes and the singularity are outside the inflationary throat but at least one of the D7-branes passes through the bottom of the throat. By the general discussion of breakage, the F1-branes will now be able to break but the D1-branes will not, and so the only cosmic string is the D-string. However, there is a caveat. The breaking of the F-strings requires the U(1) gauge field on the D7-brane. Since the D7-branes are extended, their dynamics depends on the physics not only in the throat but in the full Calabi-Yau. If the U(1) is confined in some way, then again a larger set of (p, q) strings will be stable. This may be the case, because in F theory compactifications there is not in general a U(1) gauge field locally associated to each 7-brane [35]; some of the U(1)'s are frozen out by the global dynamics.

Now consider the case that the Standard Model 3-branes and the associated singularity are located in the inflationary throat. For this discussion it will not matter whether these are D3-branes or $\overline{\text{D3}}$-branes, although of course this will strongly affect the dynamics and the breaking of supersymmetry. With 3-branes in the throat, all the (p, q) strings are unstable to breakage.

The presence of the singularity gives rise to a new possibility. We focus on the example of the $\mathbb{C}^3/\mathbb{Z}_3$ orbifold singularity for simplicity. Hidden in the singularity are a collapsed two-cycle and a collapsed four-cycle. A D3-brane wrapped on the two-cycle and a D5-brane wrapped on the four-cycle are both effectively one-dimensional, not only as seen in four dimensions but as seen in ten — they are fully located in the throat. Certain bound states of the D1 and wrapped D3 and D5 are referred to as fractional strings, from their perturbative orbifold description [36, 37].

If we consider the $\mathbb{Z}_3$ orbifold of the IIB theory before the $\mathbb{Z}_2$ orientifold, and ignore for now the D3-branes, then the fractional strings are BPS states. They couple to the global $C_{(2)}$ field, and also to two additional two-forms which are obtained from $C_{(4)}$ integrated over the collapsed two-cycle and from $C_{(6)}$ integrated over the collapsed four-cycle. In the perturbative picture the latter two two-forms are twisted sector states. The orientifold projection removes $C_{(2)}$ but not the other two — there are independent collapsed cycles and corresponding two-form fields in each image throat, and one linear combination of each survives. Thus after the orientifold these fractional strings are axionic.

To complete the story we must consider the M^4-filling branes. As with the D1-branes, there are fractional D3-branes, which arise from D5-branes wrapped on the collapsed two-cycle and from D7-branes wrapped on the collapsed four-cycle. The respective strings can break on these. In the models of ref. [33] there are M^4-filling branes of all three types — there are three U(1)'s in particular — so that all the strings are unstable to breakage.

We conclude that with the Standard Model 3-branes in the inflationary throat there are no interesting cosmic strings.

4. Large dimension models

In the $\mathbb{KLMT}$ model the string scale is lowered by a large warp factor. It can also be lowered in the context of large compact dimensions without such a large warping; this was the context for the analysis in ref. [7]. Although in these models there are not yet examples with all moduli stabilized, we can still investigate the stability of potential cosmic strings.

Consider first the prototype example, the $T^6/\mathbb{Z}_2$ orientifold, where the $\mathbb{Z}_2$ reflects k coordinates. This model is equivalent to the type I string under T-duality on k axes. It has Dp-branes and Op-planes for $p = 9 - k$. There are potential cosmic strings from wrapped Dq-branes, where $p - q$ must be even ($p - q$ odd gives non-BPS strings that immediately decay [17, 22]). The key property that characterizes the strings is the number ν of ND or DN directions, in which the Dp-branes are extended and the Dq-branes are not or vice versa. This number must be even (like $p - q$) and is at least two, from the noncompact dimensions transverse to the string. If $\nu = 2$ mod 4, the strings are non-BPS, like the D1-branes of the $\mathbb{KLMT}$ model. However, the strings here are free to fluctuate over the whole compact space and so can rapidly find their $\mathbb{Z}_2$ images and decay. If $\nu = 4$ mod 4, the strings are BPS and so axionic; only the $\nu = 2$ strings can break so these axions are not Higgsed. Because this theory has $\mathcal{N} = 4$ supersymmetry, instantons cannot generate an axion potential, and the BPS strings are stable global strings.

Reducing the supersymmetry to a realistic level will have several effects. First, it produces an axion potential so the BPS strings are now confined by domain walls. The instanton action is of order of the volume in string units of the cycle $\mathcal{K}_{7-p}$ on which it wraps. If this is of order 10^2 or more, the domain wall tension could be small enough not to confine the strings. One would expect that the supersymmetry breaking would also produce a potential for the positions of the Dq-branes in the compact space, so that these will be trapped in local minima. If a Dq-brane and its $\overline{\mathrm{D}q}$ image are in different minima then their annihilation may be suppressed as in the $\mathbb{KLMT}$ model. There the suppression was due to a deep warp factor, and here it would be due to a large distance l between the Dq and $\overline{\mathrm{D}q}$. Estimating the tunneling action (3.4) for the example of D1-branes, we have $\sigma \sim l/g_{\mathrm{s}}\alpha'$ and $\rho \sim 1/g_{\mathrm{s}}\alpha'$, and so

$$B \sim \frac{l^2}{g_{\mathrm{s}}\alpha'} . \tag{4.1}$$

This can be quite large for large l. Similarly there can be long-lived F-strings, and so (p, q) strings. There could be an even richer spectrum of strings in the large dimension models from wrapping on different cycles [7]. The Standard Model branes may allow some of these to break, but others — those wrapped on different cycles from the Standard Model branes, for example — will remain stable.

5. Observational consequences

5.1 String properties

The evolution of string networks, and their signatures, depend on certain physical properties of the strings: their tensions, their interaction with nongravitational fields, and their intercommutation properties.

5.1.1 Tension

For any given geometry the form of the brane-antibrane potential is known, and so one obtains a relation between the observed magnitude of density fluctuations δ_H and the parameters of the model. For the models of ref. [7], which are based on unwarped compactifications with the moduli fixed by hand, the authors find a range

$$10^{-11} \lesssim G\mu \lesssim 10^{-6}\,, \tag{5.1}$$

with a narrower range around 10^{-7} for their favored models based on branes at small angles. For the $\mathbb{KLMT}$ model, combining equations (3.7), (3.9), and (C.12) of ref. [10] gives

$$\frac{G^2 e^{4A_0}}{(2\pi\alpha')^2 g_{\rm s}} = \frac{\delta_H^3}{32\pi C_1^3 N_e^{5/2}}\,, \tag{5.2}$$

where C_1 is a model-dependent constant of order unity and N_e is the number of e-foldings for the observed fluctuations. Inserting the numerical values [10] $\delta_H = 1.9 \times 10^{-5}$, $C_1 = 0.39$, and $N_e = 60$ gives 4×10^{-20} for the right-hand side. The left-hand side is simply the product of the values of $G\mu$ for the F-string and the D-string, assuming for simplicity that the RR scalar C vanishes. Thus,

$$\sqrt{G\mu_{\rm F}\, G\mu_{\rm D}} \sim 2 \times 10^{-10}\,, \qquad \frac{\mu_{\rm D}}{\mu_{\rm F}} = \frac{1}{g_{\rm s}}\,. \tag{5.3}$$

The value of $g_{\rm s}$ is likely in the range 0.1 to 1 ($g_{\rm s} > 1$ can be converted to $g_{\rm s} < 1$ by S-duality), so $G\mu$ for both strings is in the range 10^{-10} to 10^{-9}.

Note that these string tensions are all deduced assuming that the cosmological perturbations arise from the fluctuations of the inflaton field. More general mechanisms would allow a wider range of scales.

5.1.2 Superconductivity

Superconducting strings carry massless charged degrees of freedom [38]. Of the strings that we have discussed, the (p, q) strings are not superconducting. They are at a potential minimum distant from any other branes in the compact space, and so their only massless fluctuations are their transverse fluctuations and the fermionic partners of these.

The D-strings, obtained when a D7-brane passes through the throat, will be superconducting with respect to the D7-brane gauge field if the D7-brane and the D1-brane are precisely coincident. In general there is no reason that this should be the case because the positions are determined by different dynamics — that of the D1-brane by the local warp factor, and that of the D7-brane by global considerations. Further, the virtual 1-7 strings contribute a short-distance repulsive term. Thus these are unlikely to be superconducting. If the D1-D7 separation is small there will be a strong coupling of the string to the Standard Model fields; such a coupling is absent for the (p, q) strings.

5.1.3 Intercommutation

Consider first the case that there is only one kind of string. When two strings cross they may pass through one another or intercommute (reconnect) as in figure 4. For gauge theory strings the reconnection probability is essentially unity [39]. For fundamental bosonic strings the probability was obtained in ref. [40]. The same result holds for the superstring [42]: the probability is of order

$$P \sim 0.5 g_{\rm s}^2 \frac{V_{\rm min}}{V} \tag{5.4}$$

for a typical collision. Here V is the six-volume within which the strings are confined, and $V_{\rm min} = (4\pi^2\alpha')^3$ is the minimum volume for toroidal compactification. For $g_{\rm s} > 1$ (and so also for D-strings) P should saturate, possibly with a slow rise, around $P \sim V_{\rm min}/V$.

For the large dimension models, the factor $V_{\rm min}/V$ will give a large suppression if the strings are free to fluctuate over the compact space [7]. This may be the case in highly supersymmetric models, but as we have noted in the previous section we expect that in realistic models any strings will be localized by a potential and so the effective V will be smaller than the total volume of the compact space. For a string of tension $\bar\mu$ whose transverse fluctuations feel an external potential characterized by mass-squared $\overline{m}^2$ (the bars indicate that we are working in the ten-dimensional metric), the quantum fluctuations of the string give

$$\langle(\delta X)^2\rangle \sim \frac{1}{4\pi\bar\mu}\ln\frac{\bar\mu}{\overline{m}^2} \sim V_{\rm min}^{1/3}\frac{\bar\mu_{\rm F}}{8\pi^2\bar\mu}\ln\frac{\bar\mu}{\overline{m}^2}\,. \tag{5.5}$$

We have used two-dimensional free field theory with a UV cutoff at $\bar\mu$. From this it appears that the fluctuations around the potential minimum will not give a large enhancement

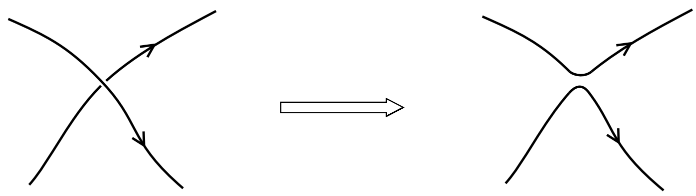

Figure 4: When two strings of the same type cross they can intercommute (reconnect).

of $V/V_{\rm min}$; in particular, the effective V grows only logarithmically with m^{-1}. Hence the probability (5.4) for strings to intercommute is not unusually suppressed. Thermal fluctuations at $T \gg m$ would give some suppression of P.

When strings of two different types cross they cannot intercommute in the same way. Rather they can produce a pair of trilinear vertices connected by a segment of string (figure 5). For example, the crossing of a (p, q) string and a (p', q') string can produce a

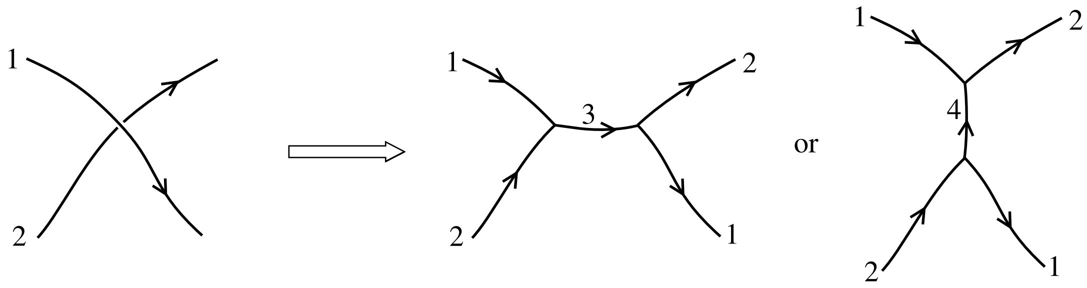

Figure 5: Crossing of strings of different types.

$(p + p', q + q')$ string or a $(p - p', q - q')$ string.[8] Only for $(p', q') = \pm(p, q)$ is the usual intercommutation possible.

A more complete treatment of this subject is in preparation [42].

5.1.4 Network properties

After their formation at the end of inflation, string networks decay by the combined processes of intercommutation, which breaks longer strings into smaller loops, and gravitational radiation. For gauge theory strings there is also the decay process due to the fields themselves, and there is a debate as to which of the decay processes dominate (gravitational versus Higgs) [43, 44]. Assuming gravitational radiation dominates, then if the decay proceeds at the maximum rate consistent with causality, the distribution of strings will scale with the horizon volume. This *scaling behavior* leads to an energy density that goes as t^{-2}, so that $w = 1/3$ in the radiation-dominated era and $w = 0$ in the matter-dominated era. That is, in each era the energy density in strings is a fixed fraction of the total energy density. Simulations indicate that networks composed of a single kind of string do scale, with $\rho_{\rm string}/\rho_{\rm rad} \sim 400G\mu$ in the radiation dominated era and $\rho_{\rm string}/\rho_{\rm mat} \sim 60G\mu$ in the matter dominated era. These values are for recombination probability $P = 1$ and increase as P decreases [45].

When there are several kinds of string, with trilinear vertices, then there is the possibility that the network evolves to a three-dimensional structure which freezes in a local minimum of the potential energy and then just grows with the expansion of the universe. In this case the energy density would evolve as $a(t)^{-2}$, or $w = -1/3$. Whether the network freezes or scales is a complicated dynamical problem. Such networks arise in field theory when a symmetry group G breaks to a discrete subgroup K. When K is non-abelian,

[8] There has been extensive study of two-dimensional supersymmetric networks formed in this way; see for example ref. [41].

intercommutation cannot occur, rather the network evolves as in figure 5. Simulations of $K = Z_3$ (with string vertices provided by monopoles) [46] and $K = S_3$ [47, 48] indicate that these systems scale rather than freeze, but with some enhanced density of strings. Simulations of $K = S_8$ [47] show an energy density that grows relative to the scaling solution and appears to indicate freezing behavior. This is consistent with the fact that the larger group allows networks of greater topological complexity, but it could also be a reflection that the simulations in [47] have not yet managed to reach the scaling regime. It is worth investigating this issue further.

To determine whether networks of (p, q) strings scale or freeze will ultimately require simulations. We conjecture that they scale, in that their topological complexity appears to be roughly that of the S_3 networks. If $g_{\rm s}$ is close to one, only the four lowlying strings with $|p|, |q| \leq 1$ are likely to be heavily populated. For any crossing between two lowlying strings, one of the two processes in figure 5 will again involve only lowlying strings, and for most angles of crossing this process will be energetically favored. Another argument for scaling behavior is to consider the limit $g_{\rm s} \ll 1$, where the D-strings are much heavier than the F-strings. The D-strings should then evolve largely independently of the F-strings, and so scale like a single-string network; after the D-strings decay to the scaling distribution on a given length scale, the F-strings in turn evolve like a single-string network.

5.2 Observational bounds

If a string network freezes into a $w = -1/3$ state, it quickly comes to dominate the energy density of the universe unless the initial energy scale is much lower than those considered above: it must be of order the weak scale or less. Thus if (p, q) string networks freeze, models with (p, q) strings are excluded with the parameters considered here. Again, our conjecture is that they do not freeze.

Assuming a scaling distribution of cosmic strings, the current upper bound on $G\mu$ comes from the power spectrum of the CMB, based on numerical evolution of the Nambu-Goto equations: $G\mu \lesssim 0.7 \times 10^{-6}$ [49] (see also ref. [50]). The tensions given in section 5.1.1 for the various models are below this bound. On the other hand numerical evolution of the underlying abelian-Higgs field theory has led Vincent et al to argue that the bound is closer to $G\mu \approx 10^{-8}$ [43] (however see also [44]).

Superconducting strings may act as sources for vortons [51], loops of cosmic string with charge and current stabilized by the angular momentum of the charge carriers. In this case they would be subject to bounds on their allowed tension, with $10^{-28} \lesssim G\mu \lesssim 10^{-10}$ being claimed to be a cosmologically unacceptable range of values [52]. If the energy scale associated with superconducting strings were close to the electroweak scale, then the vortons could become serious candidates for cold dark matter. In the context of the $\mathbb{KLMT}$ model, this would correspond to having inflation in the throat occurring at or around the electroweak scale.

Cosmic strings produce large quantities of gravitational waves, because they are relativistic and inhomogeneous. Pulsar timing measurements then place an upper bound on $G\mu$ which is roughly comparable to that from the CMB, depending on uncertainties from network properties [53].

Remarkably, future measurements of non-gaussian emission of gravitational waves from cusps on strings will be sensitive to cosmic strings with values of $G\mu$ seven orders of magnitude below the current bound, covering the entire range of tensions discussed in section 5.1.1. According to ref. [54], even LIGO 1 may be sensitive to a range around $G\mu \sim 10^{-10}$, while LIGO 2 will reach down to $G\mu \sim 10^{-11}$ and LISA to $G\mu \sim 10^{-13}$. In addition [54], pulsar timing measurements may reach a sensitivity of $G\mu \sim 10^{-11}$. Thus, gravitational waves provide a potentially large window into string physics, if we have a model in which strings are produced after inflation and are metastable.

If the string couples strongly to Standard Model fields then instead of primarily producing gravitational radiation the string network may decay through the production of high energy cosmic rays, photons and neutrinos from string cusps [55]. These authors have calculated the predicted flux of high energy gamma rays, neutrinos and cosmic ray antiprotons and protons as a function of the scale of symmetry breaking at which the strings are produced, and argued that in order to reproduce the (possibly) observed distribution of particles above the GZK cut-off, they require $G\mu \leq 10^{-9}$. Given the values we expect in the $\mathbb{KLMT}$ model this remains in the interesting regime for cosmic strings arising out of string theory. Note however refs. [56], which argue that the cosmic radiation from cusps is suppressed.

6. Conclusions

We have found that both fundamental and Dirichlet strings might be observed as cosmic strings. The issue is model-dependent — it depends on having brane inflation to produce the strings, and on having a scenario in which the strings are stable. Nevertheless, this is a potentially large and rather direct window onto string theory.

Of course, if cosmic strings are discovered, the problem will be to distinguish fundamental objects from gauge theory solitons. Indeed, this is not a completely sharp question, because these are dual descriptions of the same objects. If one can infer that the strings have intercommutation probabilities less than unity, this is a strong indication that they are weakly coupled F-strings. Discovery of a (p, q) spectrum of strings would be a promising signal for F- and D-strings. Note however that these throats have a dual gauge description [23] and therefore such strings can also be obtained in gauge theory; the spectrum is actually a signal of an $\mathrm{SL}(2, \mathbb{Z})$ duality and so might arise in other ways as well. If cosmic strings are found through the gravitational radiation from cusps, determining their tensions and intercommutation properties will require a spectrum of many events as well as precise simulations of the evolution of string networks.

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