Cosmic strings: progress and problems

Alexander Vilenkin
Institute of Cosmology, Department of Physics and Astronomy,
Tufts University, Medford, MA 02155, USA

Recent developments in cosmic strings are reviewed, with emphasis on unresolved problems.

I. INTRODUCTION

It gives me great pleasure to contribute to the volume honoring Katsuhiro Sato at his 60th birthday. It is a bit embarrassing that the subject of my contribution is one of the very few topics in modern cosmology (and probably the only one in these Proceedings) on which Katsu has not worked at one time or another during his career. But then this only serves to illustrate the amazingly wide range of his research.

The idea of cosmic strings was popular in the 1980’s and early 90’s, mainly in the context of structure formation theory. Density perturbations that would be produced by strings of the grand unification energy scale would have the right magnitude to serve as seeds for galaxies and clusters. This scenario was then disfavored by CMB observations, and cosmic string research was all but abandoned for nearly a decade. Recently, however, there has been a revival of interest in the subject. This string Renaissance has been mostly fueled by the developments in superstring theory. It now appears that fundamental strings may have astronomical dimensions and play the role of cosmic strings. Also, it has been realized that cosmic strings can have observable effects even if their energy scale is well below grand unification. Furthermore, there are some intriguing indications that a cosmic string may have already been observed. Here, I will briefly review the present state of the art, emphasizing recent developments and unresolved problems.

II. “ORDINARY” STRINGS

Cosmic strings can be formed as linear defects at symmetry breaking phase transitions in the early universe. Their formation and evolution have been extensively discussed in the literature; for a review and references see [1,2].

Topologically stable strings do not have ends. They can either form closed loops or extend to infinity. The mass per unit length of string, $\mu$, which is also equal to the string tension, is of the order $\mu \sim \eta^2$, where $\eta$ is the energy scale of symmetry breaking. The dimensionless parameter $G\mu \sim \eta^2/M_p^2$, where $G$ is Newton’s constant and $M_p$ is the Planck mass, characterizes the strength of gravitational interactions of strings. For grand-unification-scale strings, $G\mu \sim 10^{-6}$.

At the time of formation, strings have the form of a tangled network, consisting of Brownian infinite strings and a distribution of closed loops. Curved strings move under the action of their tension, developing speeds close to the speed of light. When strings cross, they reconnect, or “exchange partners”. Self-intersections result in the formation of closed
loops. Numerical simulations of string evolution indicate that the network of long strings evolves in a scale-invariant fashion. The typical distance between the strings $d(t)$, and the coherence length $\xi(t)$, defined as the distance beyond which the directions along the string are uncorrelated, both scale with the cosmic horizon:

$$d(t) \sim \xi(t) \sim t.$$  \hfill (1)

Long strings exhibit significant small-scale structure, which is partly a remnant of the initial wiggliness of the strings and partly due to multiple kinks formed at string intersections. As strings move at relativistic speeds, each horizon-size segment intersect itself or another long string about once in a Hubble time $t$. As a result, one or few loops of size $\sim t$ are produced per Hubble volume per Hubble time. These large loops then shatter, through multiple self-intersections, into a large number of small loops. According to the standard lore, the characteristic size of loops, $l_{\text{loop}}(t)$ is set by the typical wavelength of the smallest wiggles on long strings, $l_{\text{wiggles}}(t)$. There is, however, no hard evidence for that. In string simulations performed in the 1990’s, both $l_{\text{wiggles}}$ and $l_{\text{loop}}$ remained at the resolution of the simulations, so neither of these scales could be reliably determined.

There have been claims that long strings lose a substantial amount of energy by direct emission of tiny loops of size not much greater than the string thickness [3] and by radiation of massive particles [4]. Personally, I find this hard to believe. It is not clear how the motion of astronomically large strings can excite the extremely short-wavelength perturbations on the scale of the string thickness. This view is supported by the numerical results in [5,6], which show no evidence for massive radiation.

The standard scenario of string evolution assumes that the main mechanism of string smoothing on the smallest scales is the gravitational damping of small-scale wiggles. A simple estimate then gives

$$l_{\text{wiggles}}(t) \sim l_{\text{loop}}(t) \sim \alpha t,$$  \hfill (2)

where

$$\alpha \sim 50G\mu.$$  \hfill (3)

Recent developments, however, put into question some of the key assumptions of the standard scenario. First, it has been realized that the gravitational radiation from counterstreaming wiggles on long strings is far less efficient at damping the wiggles than originally thought [7]. If indeed $\alpha$ is determined by the gravitational back-reaction, then the new analysis shows [8] that its value is sensitive to the spectrum of small-scale wiggles and is generally much smaller than (3). Moreover, high-resolution numerical simulations of strings in flat spacetime indicate that small wiggles evolve in a scale-invariant manner, even in the absence of gravitational damping [9]. This suggests that the parameter $\alpha$ may be unrelated to $G\mu$, at least for strings of very low energy scale. The time is now ripe for a new generation of string simulations, which may help to resolve these important issues.

III. COSMIC SUPERSTRINGS

Witten [10] was the first to consider the possibility of cosmic superstrings, but only to promptly rule it out. For fundamental strings, the mass per unit length is $\mu \sim M_s^2$, where
Ms is the string energy scale. If one assumes that $M_s \sim M_p$, then $G\mu \sim (M_s/M_p)^2 \sim 1$. This is far above the observational bounds, which require that $G\mu \lesssim 10^{-7}$.

Recently, however, it has been shown [11] that in models with large extra dimensions $M_s$ can be $\ll M_p$, yielding $G\mu \ll 1$. Another possibility is that the effective string tension may be small due to the bulk gravitational potentials, or “warp factors”, as in the Randall-Sundrum model [12]. Suppose the metric is of the form

$$ds^2 = F(y)\eta_{\mu\nu}dx^\mu dx^\nu + ds_y^2,$$

(4)

where $x^\mu$ are the 4D spacetime coordinates and $y$ labels the coordinates in the extra dimensions. If the strings are localized in a gravitational potential well at $y = y_0$, then their effective tension is $\mu_{\text{eff}} = F(y_0)\mu$. With $F(y_0) \ll 1$, we can have $G\mu_{\text{eff}} \ll 1$ even for $M_s \sim M_p$.

Apart from the fundamental, or $F$-strings, superstring theory provides another candidate for the role of cosmic strings. It is a $D$-string, which can either be a 1-dimensional $D$-brane, or a higher-dimensional brane with all but one dimensions compactified.

Both $F$ and $D$-strings are naturally formed in models of brane inflation. In such models, the inflationary expansion is driven by the attractive interaction potential between parallel $D$-brane and anti-$D$-brane, which are separated in extra dimensions. The two branes are slowly pulled toward one another, until they collide and annihilate [13]. The role of the inflaton field in this model is played by the separation between the branes. This field becomes tachyonic when the branes get close together, and the branes quickly annihilate.

Each brane has a $U(1)$ gauge field living on its worldsheet, so there is a $U(1) \times U(1)$ symmetry prior to brane annihilation. The tachyon field is coupled to the combination $(A_1 - A_2)$ of the two gauge fields, and Nielsen-Olesen vortices form as in the usual Higgs model when the tachyon develops an expectation value. These vortices are identified with $D$-strings. The orthogonal combination $(A_1 + A_2)$ is not higgsed. When the branes annihilate, the electric component of this field combination is squeezed into electric flux lines, which are identified with $F$-strings. A substantial fraction of energy of the annihilating branes can go into a stochastic string network. However, the details of the string formation mechanism are not yet fully understood and some points remain controversial [14–16]. For a more detailed review of cosmic superstrings, see [17].

Cosmic $F$ and $D$-strings are exciting new alternatives, but it should be emphasized that we have no reason to dismiss the good old field-theory strings resulting from symmetry breaking. It has been argued in [18] that such strings inevitably arise in a wide class of supersymmetric grand unified models which exhibit inflation.

### IV. STRING PROPERTIES

The values of $G\mu$ for $F$- or $D$-strings resulting from brane inflation are expected to be in the range [19,17]

$$10^{-11} \lesssim G\mu \lesssim 10^{-6}.$$  

(5)

This estimate was obtained assuming that the same period of brane inflation accounts also for the observed density inhomogeneities. $G\mu$ can be much smaller if the perturbations were generated by some other mechanism.
An important difference between fundamental and “ordinary” strings is in their reconnection properties. As I already mentioned, when ordinary strings cross, they always reconnect. For F-strings, on the other hand, reconnection is a quantum-mechanical process, whose probability is governed by the string tension $g_s$, $P \sim g_s^2$. Another difference is due to the higher-dimensional habitat of F-strings. Moving linear objects generically intersect in 3 dimensions, but can easily miss one another in higher dimensions. This leads to further suppression of reconnection probability [15,19]. A detailed analysis in [20] suggests that $10^{-3} \lesssim P \lesssim 1$ for F-strings and $0.1 \lesssim P \lesssim 1$ for D-strings.

An intersection of an F-string with a D-string cannot result in reconnection. Instead, segments of the two strings can join, forming a bound state - an FD-string. Such interactions between F and D networks can result in the formation of an interconnected web - an FD-network - with F, D, and FD-strings joining at three-way vertices [16,15].

In braneworld models, it is also important to consider the interaction of strings with D-branes. F-strings can end on a D-brane, with their endpoints playing the role of particles on the brane worldsheet. Even though brane-antibrane pairs may have annihilated at the end of inflation, some branes must have survived, since the Standard Model particles presumably live on a stack of branes. If Brownian F-strings are allowed to wiggle around in the bulk, they will have multiple intersections with the surviving branes and will quickly break up into pieces with their ends attached to the branes. The segments will oscillate under the action of the string tension and will dissipate their energy by radiation of gauge quanta, gravitational waves, and other light fields.

This quick demise of an F-string network is not inevitable. The network may be localized away from the surviving branes by bulk gravitational fields. For example, the brane annihilation, which is followed by string formation, may occur in one potential well, while the surviving branes may reside in another well. This is precisely what happens in the KKLMMT model of brane inflation [21]. In this case, F-strings can break up only if a piece of string tunnels through the potential barrier separating the two wells. The tunneling may be strongly suppressed, in which case the strings are practically stable [16].

Depending on the model, D-strings may or may not be able to break on surviving branes. If they can break, then the situation is identical to that of F-strings, and a stable network can exist only if the strings are separated from the branes by a potential barrier. An interesting alternative is the case when the bulk potentials localize D-strings on top of the branes. If the strings cannot break (as, e.g., in the case of D1-strings and D7-branes [17]), they will be superconducting, with open F-strings connecting D1 and D7 playing the role of massless charge carriers [15].

**V. STRING EVOLUTION**

The evolution of cosmic superstrings is rather model-dependent. One possibility is that the strings break up on the surviving branes and completely disappear before present.

A more interesting alternative is when strings of only one type -either F or D - survive until present. The strings would then evolve as “ordinary” cosmic strings, except the reconnection probability will generally be less than one and can even be $P \ll 1$. For a low reconnection probability, the number of long strings per horizon is expected to increase. The scaling regime of evolution can be sustained only if the long strings have one or few
reconnections per Hubble distance per Hubble time. With $P \ll 1$, it will take many string crossings to get a reconnection. Hence, the number of strings per horizon should be large, $N_s \gg 1$. Simple estimates suggest [22,23] $N_s \sim P^{-1/2}$, in agreement with earlier numerical simulations [24]. This feature of superstrings, that there may be many of them in our Hubble volume, may help to distinguish them observationally from ordinary strings.

If both $F$- and $D$-strings are formed and are confined to the same potential well in extra dimensions, the strings will combine to form an $FD$-network. Analytic and numerical models indicate that such interconnected networks evolve towards a scaling regime, where the typical inter-string distance remains a fixed fraction of the horizon [25–28],

$$d(t) \sim \zeta t.$$  \hspace{1cm} (6)

The energy density of the network is $\rho_{FD} \sim \mu/d^2(t)$. Requiring that it is much smaller than the total energy density of the universe yields the condition

$$\frac{\rho_{FD}}{\rho} \sim \frac{G\mu/\zeta^2}{\rho} \ll 1.$$  \hspace{1cm} (7)

The parameter $\zeta$ in (6) is determined by the rate of energy dissipation in the network [25]. The simulations of Refs. [26,27] were performed for global strings, which dissipate energy very efficiently, through the emission of Goldstone boson radiation. The corresponding values of $\zeta$ are $\zeta \sim 0.1 – 0.01$, and the condition (7) is satisfied for reasonable values of $G\mu$. For superstring networks, on the other hand, the dominant radiation mechanism is gravitational radiation. It has been argued in [25] that in this case $\zeta \sim G\mu$, and Eq. (7) gives $\rho_{FD}/\rho \sim 1/G\mu \gg 1$, indicating that the string network becomes so dense that it dominates the universe.

The conclusion appears to be that models predicting the formation of $FD$-networks are ruled out. The only caveat is that the networks may lose energy more efficiently through chopping off small nets, in a way similar to chopping off closed loops by ordinary strings. This issue needs to be further investigated in high-resolution simulations of evolving string networks.

VI. LOOKING FOR COSMIC STRINGS

Apart from the new insights from superstring theory, there is another, independent reason for the renewed interest in cosmic strings. As I already mentioned, it has been recently realized that gravitational waves from strings may be detectable for a very wide range of $G\mu$, extending to values well below the grand unification scale. This is particularly interesting now, when LIGO and other detectors are beginning their search for gravitational waves (GW).

GW from string loops oscillating at different cosmic epochs add up to a stochastic background, which spans many orders of magnitude in frequency. Such a background would introduce noise into the arrival times of signals coming from the millisecond pulsar. The bound on $G\mu$ from the pulsar data is [29]

$$G\mu \lesssim 10^{-7}.$$  \hspace{1cm} (8)
(For a detailed discussion of the current pulsar bounds, see [22].) A similar bound follows from the microwave background observations [30,31].

The key new development is the realization that, in addition to a featureless Gaussian component, the GW background from strings includes sharp GW bursts. Oscillating loops of string typically develop “cusps” once or few times during the period of oscillation, with the string velocity momentarily reaching the speed of light at the tip of the cusp. Short GW bursts emanating from cusps carry away about the same power as the radiation at low frequencies, comparable to the oscillation frequency of the loop. The analysis in [32] suggests that the bursts should be detectable by LIGO and LISA for values of $G\mu$ as low as $10^{-12} - 10^{-14}$. The effects of the low reconnection probability, $P \ll 1$, and of the uncertainty in the loop size parameter $\alpha$ have been investigated in [22], with the conclusion that the predictions of [32] are quite robust, at least when the loop sizes are not suppressed by many orders of magnitude relative to the standard scenario.

Another intriguing new development is the observation of two nearly identical galaxies at redshift $z = 0.46$ with angular separation of 1.9 arc seconds [33]. The spectra of the two galaxies coincide at 99.9% confidence level [34]. The most plausible interpretation of the data appears to be lensing by a cosmic string with $G\mu \sim 4 \times 10^{-7}$. This estimate assumes a slowly moving string orthogonal to the line of sight at a relatively low redshift ($z \lesssim 0.1$). Increasing the string redshift or changing its orientation would give a higher estimate for $G\mu$, which might be in conflict with the CMB and pulsar observations. On the other hand, the estimate for $G\mu$ could be decreased due to relativistic motion [35,36] or wiggliness [37] of the string. The only alternative to the string interpretation is that two very similar galaxies just happen to be next to one another. The issue is likely to be resolved by Space Telescope observations later this year.

If strings are superconducting, they can produce a rich variety of observational effects. Among them are gamma-ray bursts [38] and high-energy cosmic rays (for a discussion and references, see [39]).

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[1] A. Vilenkin and E.P.S. Shellard, *Cosmic Strings and Other Topological Defects*, Cambridge University Press (Cambridge, 2000).
[2] M.B. Hindmarsh and T.W.B. Kibble, Rep. Prog. Phys. 58, 477 (1995).
[3] G.R. Vincent, M. Hindmarsh and M. Sakellariadou, Phys. Rev. D56, 637 (1997).
[4] G.R. Vincent, N.D. Antunes and M. Hindmarsh, Phys. Rev. Lett. 80, 2277 (1998).
[5] J.N. Moore and E.P.S. Shellard, hep-ph/9808336.
[6] K.D. Olum and J.J. Blanco-Pillado, Phys. Rev. Lett. 84, 4288 (2000).
[7] K.D. Olum and X. Siemens, Nucl. Phys. B611, 125 (2001).
[8] K.D. Olum, X. Siemens and A. Vilenkin, Phys. Rev. D66, 043501 (2002).
[9] V. Vanchurin, K.D. Olum and A. Vilenkin, Cosmic string scaling in flat space, gr-qc/0501040.
[10] E. Witten, Nucl. Phys. B249, 557 (1985).
[11] N. Arkani-Hamed, S. Dimopoulos and G. Dvali, Phys. Rev. D59, 086004 (1999).
[12] L. Randall and R. Sundrum, Phys. Rev. Lett. **83**, 3370 (1999).
[13] G.R. Dvali and S.-H. Tye, Phys. Lett. **B450**, 72 (1999).
[14] S. Sarangi and S.-H. Tye, Phys. Lett. **B536**, 185 (2002).
[15] G. Dvali and A. Vilenkin, JCAP **0403**, 010 (2004).
[16] E.J.Copeland, R.C. Myers and J. Polchinski, JHEP **0406**, 013 (2004).
[17] J. Polchinski, Introduction to cosmic F- and D-strings, hep-th/0412244.
[18] R. Jeannerot, J. Rocher and M. sakellariadou, Phys. Rev. **D68**, 103514 (2003).
[19] N.T. Jones, H. Stoica and S.H. Tye, Phys. Lett. **B563**, 6 (2003).
[20] M.G. Jackson, N.T. Jones and J. Polchinski, Collisions of F- and D-strings, hep-th/0405229.
[21] S. Kachru, R. Kallosh, A. Linde, J. Maldacena, L. McAllister and S.P. Trivedi, JCAP **0310**, 013 (2003).
[22] T. Damour and A. Vilenkin, Phys. Rev. **D71**, 063510 (2005).
[23] M. Sakellariadou, JCAP **0504**, 003 (2005).
[24] M. Sakellariadou and A. Vilenkin, Phys. Rev. **D42**, 349 (1990).
[25] T. Vachaspati and A. Vilenkin, Phys. Rev. **D35**, 1131 (1987).
[26] U.-L. Pen and D.N. Spergel, Phys. Rev. **D51**, 4099 (1995).
[27] E.J. Copeland and P.M. Saffin, hep-th/0505110.
[28] S.H. Tye, I. Wasserman and M. Wyman, Phys. Rev. **D71**, 103508 (2005).
[29] V.M. Kaspi, J.H. Taylor and M.F. Ryba, Astrophys. J. **428**, 713 (1994).
[30] M. Wyman, L. Pogosian and I. Wasserman, Phys. Rev. **D72**, 032513 (2005).
[31] A. Fraisse, astro-ph/0503402.
[32] T. Damour and A. Vilenkin, Phys. Rev. **D64**, 064008 (2001).
[33] M. Sazhin et. al., Mon. Not. Roy. Astro. Soc. **343**, 353 (2003).
[34] M. Sazhin et. al., astro-ph/0506400.
[35] A. Vilenkin, Nature **322**, 613 (1986).
[36] B. Shlaer and S.H. Tye, hep-th/0502242.
[37] T. Vachaspati and A. Vilenkin, Phys. Rev. Lett. **67**, 1057 (1951).
[38] V. Berezinsky, B. Hnatyk and A. Vilenkin, Phys. Rev. **D64**, 043004 (2001).
[39] V. Berezinsky, P. Blasi and A. Vilenkin, Phys. Rev. **D58**, 103515 (1998).