Cosmic Rays From Cosmic Strings

A.J. Gill and T.W.B. Kibble

Blackett Laboratory
Imperial College
South Kensington
London SW7 2BZ
United Kingdom

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Abstract

It has been speculated that cosmic string networks could produce ultra-high energy cosmic rays as a by-product of their evolution. By making use of recent work on the evolution of such networks, it will be shown that the flux of cosmic rays from cosmologically useful, that is GUT scale strings, is too small to be used as a test for strings with any foreseeable technology.
1 Introduction

Cosmic strings produced during GUT scale phase transitions in the early universe could provide a mechanism for cosmological structure formation [1]. They also preserve regions of space-time in the symmetry unbroken phase due to boundary conditions which topologically restrict their decay. Under certain circumstances, these restrictions are removed and the energy stored in the unbroken vacuum phase is liberated in the form of the GUT scale quanta of the gauge and scalar fields which form the defects. It has been suggested that the subsequent decay of these massive GUT quanta could be a source of ultra-high energy cosmic rays [2] both at and above the highest energies currently observable. If enough of these cosmic rays are produced then they could provide a clean observational signal characteristic of cosmic strings.

Such a scheme for producing ultra-high energy cosmic rays obviates the need for acceleration mechanisms, such as those involving AGNs, which are currently thought to be responsible for most of the highest energy cosmic rays observed [3]. Thus, high energy neutral as well as charged particles could be produced, although the radiation from cosmic strings would not contain any nuclei heavier than hydrogen. Also, acceleration mechanisms in AGNs are thought to be incapable of producing cosmic rays which would reach us at energies above $10^{19}$ eV. Thus, if the flux of ultra-high-energy cosmic rays were large enough to be measurable, we would have a very clean, background free test for cosmic strings. Even knowing that strings are not capable of producing cosmic rays at such high energies is advantageous as anything which is then seen must be from some previously unknown process and therefore indicate the presence of entirely new physics. Experiments are now either running, being built or being designed to measure the ultra-high energy cosmic ray flux at appropriate energies [4] [5].

Prior to this, much work had already been done on the mechanisms by which strings decay to the particles commonly found in cosmic rays [6] and also on the propagation of high-energy cosmic rays [7]. Due to calculational and computational difficulties concerning the evolution of cosmic string networks, however, very little predictive work had been done on the likely flux of cosmic rays and no work at all had been done which included all possible mechanisms for removing topological stability. Thus the magnitude of the expected cosmic ray flux was a completely unknown quantity. In much previous work, the predictions were simply normalised to match the experimental cosmic ray data at around a joule. Despite the lack of a solid justification, no criticism was made of such techniques and the predictions were adopted uncritically by the astrophysics community [8]. In addition to this, our understanding of some of the assumptions made in the early work has changed significantly since this was done, to the extent that some of the earlier results must now be modified.

It is also possible that other topological defects could produce observable cosmic rays, but this is as yet a virgin field as our ideas of the evolution of other defects are not as advanced as those concerning strings.
The Mechanism Of Cosmic Ray Production

2 Cosmic strings, and topological defects in general, are soliton solutions of classical equations which, it is claimed, are the effective field equations of the, as yet unknown, GUT-scale physics [9] [10]. These solutions describe the vacuum expectation value of a scalar field $\langle \phi \rangle$ and its associated gauge field. The scalar field involved may thus be either a genuine scalar, such as the Higgs, or a composite structure somewhat like a Cooper pair. Such defects may be formed in gauge theories during phase transitions in which the full symmetry group is broken in some fashion [11] [12].

One of the advantages of the cosmic string model is the paucity of free parameters. All of the cosmological consequences of strings are relatively insensitive to the coupling constants of the field theory and depend almost entirely on a single parameter related to the temperature of the symmetry-breaking phase transition. This parameter is usually specified in the form of the dimensionless number $G\mu/c^4$ where $\mu$ is the energy per unit length of the string. For cosmologically useful strings, this would have to be of the order of $10^{-6}$. The cosmic ray flux also depends in principle not just on $\mu$ but also on the values of the coupling constants, although in practice this is not a problem as one can place a unitarity bound on them and thence obtain an upper limit on the flux.

Once formed, strings preserve linear regions of space-time in the symmetry-unbroken phase since the boundary condition that the fields should be in their symmetry broken state at infinity places a topological restriction on their evolution. Thus, a network of cosmic strings stores a large amount of energy. How the network evolves determines the fate of this energy.

The string network may be described statistically by various physical length scales, for example the characteristic distance between strings. Such scales are time dependent, their time evolution describing how the network changes with time. The network can evolve in one of two possible ways. Causality prevents the physical length scales growing faster than $ct$ so either the length scales are all proportional to $ct$, which is referred to as scaling, or the length scales grow less fast than $ct$ and the strings come to dominate the energy density of the universe [13]. Which of these two alternatives describes a real string network depends on whether the network can lose energy fast enough to scale. If strings are to be a viable cosmological model then they must scale as string domination is ruled out by observation. The energy lost from a cosmic string network provides us with a way to look for strings and thus probe some of the GUT scale physics which would be responsible for them.

There are two ways in which cosmic strings can lose energy. Firstly they can radiate gravitationally, which provides one of the current bounds on $G\mu/c^4$. Alternatively, one can remove the topological stability from a length of string in some way. This is the basis of the mechanism whereby strings can produce cosmic rays. There are three possible ways of doing this, viz. inter-commutation, loop collapse and cusp evaporation. To calculate the expected flux of cosmic rays, it is necessary to know how many of each of these types of event occur at each epoch, each event producing...
quanta of either the scalar or gauge fields. Previously, the most general work had assumed a generic two parameter form for the event rate and thus the number of GUT quanta produced at each epoch. This was based on the idea of scaling \([15]\). However, the two parameters involved do not have obvious physical meanings. In the light of recent analytical work \([23]\), it is possible to improve on this generic fit and write expressions involving only physically meaningful parameters whose values and uncertainties are relatively well known from simulations.

GUT and gauge quanta are therefore produced and decay rapidly on moving through the symmetry broken phase outside the string. Whatever the actual GUT responsible for the physics, the final decay products must be stable particles with which we are familiar. Whatever the high energy physics, there is ultimately a very restricted range of decay routes once at the relatively low energies with which we are familiar. Three generic decay routes stand out. Firstly, the decay may end up producing gauge bosons; these do not propagate far as they either decay, like the Ws or Z, or interact with the inter-galactic medium. They tend to interact and produce other particles. A leptonic route tends to produce a large number of neutrinos and electrons, with a few hadrons and heavier charged leptons. These neutrinos could be important because of their comparatively long path lengths in the inter-galactic medium. Finally there is the QCD decay route in which a jet is formed from a quark-gluon source.

The theory of this scheme has already been extremely well studied \([17]\). It is the QCD jets which arouse the most interest as they produce a relatively large number of particles. They also have the advantage that there is a clear dichotomy between the observations which are characteristic of decay from GUT scale energies and those which are characteristic of the string network itself. The energy spectrum of the particles from a single jet depends almost exclusively on the QCD and is therefore characteristic of decay from a GUT particle, whereas the statistics of the jets themselves and the event rates, if they could be observed, would be characteristic of the particular defect responsible. On the other hand, if single high-energy particles were to be observed then it would be almost impossible to say conclusively whether these were from cosmic strings or from some other as yet unknown process involving decay from GUT energies.

Once one has an injection spectrum of ultra-high-energy particles, the propagation of these particles is a relatively well understood problem \([18]\) \([7]\), although it turns out that it is not really worth doing this part of the calculation since the event rate is so low.

### 3 Calculation Of The Cosmic Ray Flux

As stated previously, there are three types of events which remove topological stability from a length of string, viz. intercommutation, loop-collapse and cusp formation. All of these present difficulties for either numerical or analytic treatments of the field theory. In principle, a sufficiently good computer simulation would solve all of these problems, up to the freedom of choice of the two coupling constants and the symmetry breaking scale. In practice, however, such an approach is not tractable.
with current technology, mostly due to the difficulty of obtaining a sufficiently great dynamic range.

Consider an intercommutation event in which two lengths of string collide and swap partners [1]. Simulations, particularly of global strings, lead one to expect that energy is lost during the exchange. Physically, one expects this to be of the order of $\mu w$, where $w$ is the width of the strings and to be in the form of GUT scale gauge and scalar bosons. Henceforth these will be referred to collectively as $X$-particles; this terminology does not imply that they are GUT gauge bosons and is adopted merely to be consistent with other papers in the field. The number of $X$-particles produced per intercommutation event will be $(\mu w)/(m_X c^2)$, where $w \approx 10^{-31}$ m, is the width of the string, and $m_X$ is the mass of the GUT quanta. This is roughly of order one, within an order of magnitude or so, depending on exactly how one fixes $\mu$ and $m_X$.

Loops of cosmic string are formed when a length of string intersects with itself; intercommutation then chops off a loop of string which begins evolution independently of the original length. Loops of cosmic string normally radiate almost entirely gravitationally, shrinking as they do so. As a loop shrinks, it will probably chop itself up into smaller loops by inter-commutation events. Ultimately, loops reach a point when their diameter is of the order of the string radius and their topological stability is removed by field overlap. At this point, the energy remaining in the string is liberated as $X$-particles and the number of particles is again of the order $\mu w/m_X c^2$, neglecting factors of $\pi$.

It is not known how many inter-commutation events are involved in chopping up a loop prior to its demise or indeed when the loop ‘realises’ that it is unstable. This uncertainty may introduce a factor of a thousand or so but the important point is that the number of $X$-particles produced by the demise of a loop is fixed. It is not time dependent. Previously, it was sometimes assumed that a certain fixed fraction of the energy initially in a loop would end up as cosmic rays. Such an assumption does not accurately reflect the physics, however, as it would imply that the energy liberated in each loop demise scales since the average loop size is directly proportional to time. Instead of growing with time, the energy liberated per event is actually constant.

Cusps are points where the radius of curvature of the string becomes very small. Mathematically, the radius of curvature is undefined for a cusp but backreaction processes prevent this ever quite happening. The string may be approximately described by the Nambu-Goto action, from which it is possible to deduce that instantaneously a cusp moves along relative to the string at the speed of light.

Cusps have always been thought to radiate massive gauge and scalar bosons, which is a part of the backreaction mentioned above. According to a recent calculation by Mohazzab [19], the energy radiated by a cusp is of the order of $\frac{1}{7} g^{5/2} \eta$ where $\eta$ is the symmetry breaking scale and $g$ is the coupling constant between the string scalar field and whatever is radiated. This is much smaller than was previously held to be the case. Broadly speaking this is because the sizes of the incoming waves are much bigger than the thickness of the cosmic string. Choosing $g$ to be close to the unitarity limit, that is to be about one, would give a plausible upper bound on the cosmic-ray flux from cusps. However, this perturbative result is odd in that it is totally independent of the shape of the cusp and it might be speculated that
a higher order calculation would bring to light some small dependence of this sort, which would increase the energy released. To be conservative in the upper limit on the flux, therefore, old ideas about cusps will be employed which give much larger, time-dependent fluxes. According to the old ideas of cusps, the length of string which overlaps in a cusp is of the order of \((\zeta^2 w)^{1/3}\), where \(\zeta\) is the characteristic length of the small-scale structure on strings; essentially the interkink distance. Such a length scales, so one may write \(\zeta = c/\epsilon H = 3ct/2\epsilon\) in the matter era where the parameter \(\epsilon\) is dimensionless. This suggests that the number of X-particles radiated will be something like \((3ct/2\epsilon w)^{2/3}\). For the sake of a reliable upper bound, this is the expression we will use, rather than that of Mohazzab, although the latter may well be more accurate.

If we wish to calculate a cosmic ray flux then in addition to knowing the energy produced per topological stability removing event, one also needs some idea of how many such events there are likely to be and therefore some idea of how the network evolves. There exist three major simulations \[20\] \[21\] \[22\] and one analytic study \[23\], although these only give a rough indication of the evolution since this is a very difficult problem. A combination of these enables one to say something about the number of stability removing events and therefore the rate of production of X-particles; the analytic being used to extract the relevant physics from the simulations without the restrictions on the dynamic range. In what follows, much of the notation is the same as that used in the analytic study \[23\].

Consider first intercommutations. The length of string in a volume \(V\) is \(L = V/\xi^2\), where \(\xi\) is a length scale related to the overall density of strings. The probability that such a length intercommutes per unit time is \(p = \chi Lc/\xi^2\), where from the simulations \(\chi\) is probably about 0.03 and certainly no larger than 0.2. Thus the expected number of intercommutations per unit volume is:

\[
\langle n_{ic}(t) \rangle = \frac{\chi c}{\xi^4}.
\]

Following the same procedure as the aforementioned study, we define a dimensionless constant \(\gamma = c/H\xi\) where in the matter-dominated era \(\gamma\) is somewhere between one and four and \(H = 2/3t\), thus giving:

\[
\langle n_{ic}(t) \rangle = \frac{\chi (H\gamma)^4}{c^3} = \frac{16\chi}{81\epsilon^3} \left(\frac{\gamma}{t}\right)^4
\]

Loops are a better understood problem since they were once thought to be the crux of a successful structure formation model. Once formed, loops of cosmic string radiate energy in the form of gravitational waves until they reach the point when the overlap between the field on different sides of the loop causes the loop to realise that it is not topologically stable. This gives loops a life-time of \(l_b c^3/\Gamma G\mu\), where \(l_b\) is their initial length and \(\Gamma\) is a dimensionless constant of order 100, governing the rate of gravitational radiation.

Let \(l_b(t_b)\) be the mean length of loops born at \(t_b\). For simplicity, we shall assume that all loops are born with length \(l_b\). Their length at a later time \(t\) will then be \(l = l_b - \Gamma G\mu(t-t_b)\). It is convenient to define a dimensionless constant \(K\) by setting:

\[
l_b = (K - 1)\Gamma G\mu t_b/c^3.
\]
It then follows that the loops born at \( t_b \) shrink by the emission of gravitational radiation and finally expire at time \( t = K t_b \).

How large is \( K \)? In the analytic study mentioned above [23], it was concluded that the characteristic length of the small-scale structure on strings, \( \zeta \), is likely to be smaller than \( \xi \) by a factor of order \( \Gamma G \mu / c^4 \approx 10^{-4} \). Setting \( \epsilon = c / H \zeta \), as above, \( \epsilon \approx 10^4 \). Moreover, the simulations suggest that most loops are born with a length much smaller than \( \xi \), of the order of a few times \( \zeta \). It follows that we should expect \( K \) to be of the order of a few, probably between 2 and 10 say.

We also need to know how many loops are formed. The number of loops decaying to X-particles per unit time per unit volume at a time \( t \) is equal to the number born at a time \( t_b = t / K \) in a volume \( (a(t_b) / a(t))^3 \). Now the number of loops born at a time \( t \) in a volume \( V \) is given by:

\[
\dot{N} = \frac{\nu V c}{(K - 1) \Gamma G \mu t_b^4}.
\]

The parameter \( \nu \) is dimensionless and characterises the rate of loss of energy in long strings by loop formation. Thus the number of loops expiring per unit volume at a time \( t \) which were born at a prior time \( t_b \) is:

\[
\langle n(t) \rangle = \frac{\nu c}{(K - 1) \Gamma G \mu} \left( \frac{a(t_b)}{a(t)} \right)^3 t_b^{-4}.
\]

In this approximation where all loops are born with a length equal to the average length of a loop at birth, all of the loops born at the same time die at the same time and vice-versa so to find the total number of loops expiring at a particular time, there is no need to integrate over birthdays. Also, in the matter dominated era, \( a \propto t^{2/3} \), so one obtains the final result:

\[
\langle n(t) \rangle = \frac{K \nu c}{(K - 1) \Gamma G \mu t^4}.
\]

Cusps are the least well understood of the possible sources of cosmic rays, although there has been at least one convincing study [25, 26]. In a scaling regime, the number of cusps per unit time and length on a long string must be given by one over the square of some length scale, for example \( c / \xi^2 \), where \( \xi \) is the correlation length along the string. However, it is far from certain that \( \xi \) is the right scale to use. Arguably, cusps have more to do with the small-scale structure, governed by the smallest of the length scales, \( \zeta \), which is smaller than \( \xi \) by a factor of order \( \Gamma G \mu / c^4 \approx 10^{-4} \). In order to obtain an upper limit on the event rate the smallest length scale, \( \zeta \), will be used.

If \( CLS \) and \( CL \) are the numbers of cusps created per unit length per unit time on long string and loops respectively, then:

\[
\langle n_c(t) \rangle = L_{LS} C^{LS} + L_{L} C^{L},
\]

\( ^2 \)It is related to the parameters defined in [23] by \( \nu = 8 f c_0 \gamma \gamma^2 / 27 \), where \( f \approx 0.7 \) is the fraction of the loop energy that goes into gravitational radiation [24] and \( c_0 \) denotes the parameter designated \( c \) in the analytic study. Using these values, we have \( \gamma \gamma \approx 0.1 \), whence \( \nu \), also of order 0.1.
where $L_{LS}$ and $L_L$ are the lengths of string per unit volume in long string and in loops respectively. If one assumes that, on long string, the number of cusps per unit length per unit time is $C_{LS} = f_{LS} c / \xi^2$, then $f_{LS}$ will be a dimensionless constant at most of order unity. As before, the length of long string per unit volume is $L_{LS} = 1 / \xi^2$, where $\xi = c / H \gamma = 3ct / 2\gamma$ in the matter dominated era. Similarly one may assume $L_L C_L$ to be $f_{LC} / \xi^2 \xi^2$. Defining the constant $f = f_{LS} + f_L$, and exploiting the fact that during the matter era $H = 2 / 3t$, the number of cusp events per unit volume per unit time becomes:

$$\langle n_c(t) \rangle = \frac{16 f^2 \gamma^2}{81 c t^4}.$$

It will be seen that this also goes like $1/t^4$ as previously. There is in fact a general argument as to why such event rates should go like $1/t^4$. One expects $dn$ to be proportional to $V$ and $dt$. The only dimensionful quantities one can put in the denominator to make the dimensions correct are lengths or times. Since all of the relevant and physically meaningful lengths scale $^3$, the denominator must go like $t^4$.

It will now be shown that a calculation for the propagation of the rays is superfluous as the flux is far too small for spectra to be useful. Assume that the beam of cosmic ray particles from each X-particle event does not disperse too widely and counts as one event if observed. Also assume that all of the detected events are useful, which may not be the case with single particles. It turns out to be clearest to work with conformal time and space co-ordinates $\tau$ and $x$ as used in the metric for a spatially flat Friedmann-Robertson-Walker universe:

$$ds^2 = a^2(\tau) (c^2 d\tau^2 - dx^2).$$

In what follows, a zero subscript will always refer to the value of a quantity at the present day.

The number of events produced in a spherical shell between $x$ and $x + dx$ during a conformal time interval $\tau$ to $\tau + d\tau$ is of the form:

$$\langle n_t(t) \rangle 4\pi a^3 x^2 dx d\tau.$$

These will arrive during a time interval $dt = dt_0 = a_0 d\tau_0 = a_0 d\tau$ and the fraction which will be detected by a detector of area $A$ will be $A / 4\pi a_0^2 x^2$. Thus on integrating back to a distance equivalent to the time of equal energy and matter density, which is as far back as is physically feasible and summing over event types, the expected event rate will be:

$$\left\langle \frac{dN}{dt_0} \right\rangle = \sum_t A \int_0^{\tau_0 - \tau_q} N_t \langle n_t(x) \rangle \frac{a^4}{a_0^4} dx.$$

There is no point in integrating any further back because any cosmic rays from such distances will not only have been heavily red-shifted but also severely affected by the medium through which they have travelled and so would not be useful even in the

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$^3$A priori one might expect the width of the string, which does not scale, to be relevant to cusps but with the optimistic assumptions made this is not the case.
unlikely event that they should reach us intact. Thus, on substituting for $\langle n_t(x) \rangle$ and $N_t$ and changing variables to conformal time rather than $x$, one obtains:

$$\left\langle \frac{dN}{dt_0} \right\rangle = \frac{Ac}{t_0^{3/2}} \left[ t_0 - 1 \right] \left\{ N_w \frac{16 \chi \gamma^4}{81c^3} + N_t \frac{K \nu c}{(K - 1) \Gamma G\mu} \right\} + \frac{4Af\gamma^2}{27} \left( \frac{12 \epsilon}{c^2 w} \right)^{2/3} \frac{1}{t_{0}^{7/3}} \left\{ \left( \frac{t_0}{t_{eq}} \right)^{1/3} - 1 \right\}.$$  

To gain an idea of the order of magnitude of the event rate, it is necessary to substitute some of the values of the dimensionless constants obtained from the simulations. Using very optimistic values gives an upper limit on the eventual flux. At best, $\Gamma = 100$, $\gamma = 1$, $G\mu/c^4 = 10^{-6}$, $\chi = 0.1$, $w = 10^{-31} m$, $\nu \approx 0.1$ and $K = 10$. The event rate is then:

$$\left\langle \frac{dN}{dt_0} \right\rangle = \frac{A c^2}{t_0^{3/2}} \left[ t_0 - 1 \right] \left\{ 10^{-2} N_{ic} + 10^5 N_t \right\} + \frac{10^5 A}{(c^4 w^2 t_0^{1/3})^{1/3}} \left[ \left( \frac{t_0}{t_{eq}} \right)^{1/3} - 1 \right].$$

All of the numerical prefactors in this expression are dimensionless. Notice that cusps are much more significant that loop collapses which are in turn more significant than intercommutations. This is to be expected as near us there is far more string in loops than in long straight lengths and there is at least one cusp per loop in our scheme. Thus, only the cusp term needs to be considered.

Adjusting the numerical prefactors to obtain the event rate per year rather than per second and neglecting the intercommutation and loop terms, the event rate becomes:

$$\left\langle \frac{dN}{dt_0} \right\rangle = \frac{10^{12} A}{(c^4 w^2 t_0^{1/3})^{1/3}} \left[ \left( \frac{t_0}{t_{eq}} \right)^{1/3} - 1 \right] \approx 10^{-18} A \text{ year}^{-1}.$$  

This corresponds to an event rate of $10^{-10}$ per square kilometre per century from cosmic strings. The flux of primary cosmic rays at the highest energies is observed to be about one particle per square kilometer per year [29].

An alternative comparison with experiment may be obtained by the simple expedient of inserting a redshift factor into the flux calculation and calculating an energy flux, $F$. Omitting dimensionless factors of order unity, this turns out to be:

$$F = \left\langle \frac{dE}{dt_0} \right\rangle \approx Am \chi \left( \frac{c^2 \epsilon^4}{w t_0} \right)^{1/3}.$$  

If one then assumes that the energy spectrum is of the form $\alpha E^{-\beta}$ above about $10^{17}$ eV then the differential particle flux, which is the quantity normally plotted by experimentalists, is roughly $F/E^2$. The differential flux from cosmic strings is several orders of magnitude lower than that observed.

In interpreting both of these comparisons, it should be borne in mind that the predictions are an extreme upper limit on the flux due to strings and do not include any form of attenuation. Further, they include cosmic rays produced at the time of matter-energy equality whereas a few hundred megaparsecs is probably far more reasonable. The reason that the flux from cosmic strings is so small is that space is so big and there are so few relevant events in it.
4 Conclusions

The flux of ultra-high-energy cosmic rays from gauge cosmic strings is smaller than the observed fluxes by about ten orders of magnitude, too small to be detectable with current technology. Thus any cosmic rays observed above about a joule are not due to cosmic strings. They can only be explained by a modification of current ideas about acceleration mechanisms or some entirely new physics. The assumptions of previous authors which were incorrect and led to their conclusion that cosmic strings could be responsible for ultra-high-energy cosmic rays were:-

- The absolute amplitudes of the predicted spectra were found by matching to experiment, thus assuming that strings are responsible for ultra-high-energy cosmic rays, rather than from any knowledge of the evolution of a cosmic string network. If this is done then there is a tendency for the scenario to predict too large a proton flux at around a few hundred MeV, although it is possible to resolve this crisis [27] [28].

- The energy yielded by each event was sometimes expressed as a fraction of a physical length, such as the length of a loop. This makes scaling of the energy per event implicit in the calculation and vastly overestimates the flux.

Also, this paper has improved on previous work by considering all of the possible mechanisms for cosmic ray production at the same time and in the same formalism.

It is possible that other types of defect could be a source for cosmic rays. There are three which are cosmologically viable, viz. global strings, textures and superconducting strings. All of these suffer to some extent from the fact that they were invented more recently than gauge strings and therefore their properties and evolution have been less well studied than those of gauge strings. It is the authors’ conviction, however, that these defects will also fail to provide an observable flux for similar reasons to those for gauge strings. Also superconducting strings will tend to lose much energy to relatively low frequency synchrotron radiation which will not help.

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References

[1] Efstathiou G., in Physics of the Early Universe, ed. J.A. Peacock, A.F. Heavens, A.T.Davies. p.361 SUSSP(1990) Edinburgh U.K.

[2] P. Bhattacharjee in Astrophysical Aspects Of The Most Energetic Cosmic Rays, eds. M. Tagano and F. Takahara (World Scientific, Singapore, 1991), p382.
[3] Astrophysical Aspects Of The Most Energetic Cosmic Rays, eds. M. Tagano and F. Takahara (World Scientific, Singapore, 1991)

[4] Proceedings Of The 23rd. Cosmic Ray Conference, July 1993.

[5] A. Okada in : Astrophysical Aspects Of The Most Energetic Cosmic Rays, eds. M. Tagano and F. Takahara (World Scientific, Singapore, 1991)

[6] M. Srednicki and S. Theisen, Physics Letters B189 (1987) 397

[7] Astrophysical Aspects Of The Most Energetic Cosmic Rays, eds. M. Tagano and F. Takahara (World Scientific, Singapore, 1991)

[8] Chi et al. in Proceedings Of The 23rd. Cosmic Ray Conference, July 1993, Volume 9 p345

[9] The Formation And Evolution Of Cosmic Strings, eds. Gibbons, Hawking and Vachaspati, CUP 1990

[10] P. Langacker, Phys. Rep. 72 (1981) 187

[11] T.W.B.Kibble, J. Phys. A9 (1976) 1387

[12] T.W.B.Kibble in : Current Topics In Astrofundamental Physics, ed. N Sánchez and A. Zichichi (World Scientific, 1992), p.68.

[13] E. Copeland and T.W.B.Kibble, Physica Scripta, T36(1991) 153

[14] Bowick et al. SU-HEP-4241-512

[15] P. Bhattacharjee, C.T. Hill and D. Schramm, Phys. Rev. Lett. 69 (1992) 567

[16] Austin et al., Phys. Rev. D48 (1993) 5594

[17] F. A. Aharonian, P. Bhattacharjee and D. N. Schramm, Phys. Rev. D46 (1992) 4188

[18] C.T. Hill and D. N. Schramm, Phys. Rev. D31 (1985) 564

[19] M. Mohazzab, BROWN-HET-912

[20] B. Allen and E.P.S.Shellard, Phys. Rev. Lett. 64 (1990) 119

[21] A. Albrecht and N. Turok, Phys. Rev. D40 (1989) 973

[22] D.P.Bennett And F.R.Bouchet, Phy. Rev. Lett. 60 (1988) 257; 63 (1989) 2776

[23] Austin et al., Phys. Rev. D48 (1993) 5594

[24] B. Allen and R. R. Coldwell, Phys. Rev. D45 (1992) 3447

[25] J. H. MacGibbon and R.H. Brandenberger, Nucl. Phys. B331 (1990) 153

[26] J. H. MacGibbon and R.H. Brandenberger, Phys. Rev. D47 (1993) 2283

[27] Chi et al. Astroparticle Physics 1 (1992) 129

[28] Chi et al., Astroparticle Physics 1 (1993) 129

[29] P. Sokolsky in Introduction to Ultrahigh Energy Cosmic Ray Physics, Addison-Wesley 1989.