

COSMIC STRINGS FROM PSEUDO-ANOMALOUS $U(1)$'s

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We discuss the properties of cosmic strings based on a spontaneously broken anomalous $U(1)$ when the anomalies are cancelled via the Green-Schwarz mechanism. Such strings arise in a large class of superstring compactifications and act as a source for an axion-like field which has model dependent couplings to QCD and QED. If this field couples to QED, but not to QCD, these strings could exist today and would exhibit novel optical properties. Such strings could be distinguished from other types of cosmic strings by studying the relative polarization of the images they produce when acting as gravitational lenses. If the axion-like field produced by these strings does couple to QCD, these strings as well as macroscopic heterotic superstrings would become the boundaries of domain walls and would rapidly disappear. There would be, however, bound configurations of the two types of string which do not couple to QCD and which could exist in the present universe.

1. Introduction

Both superstrings and cosmic strings have been the subject of much attention in the last few years. If these objects could be shown to be one and the same it would be one of the most marvelous examples of the relationship between cosmology and particle physics. While macroscopic superstrings are possible, there seem to be difficult problems in associating them with the cosmic strings which might play a role in galaxy formation [1]. First of all, superstrings act as a source for a pseudo-scalar field with axion-like couplings. Because of this, in the QCD phase transition these strings become the boundaries of axion domain walls and rapidly disappear. Secondly, based on our current understanding, macroscopic superstrings would have a string tension which, while smaller than one might naively assume, is still large enough to be in apparent conflict with limits on the isotropy of the

microwave background [2]. On the other hand, macroscopic superstrings are always superconducting, a phenomenon which plays a role in some cosmic string scenarios for galaxy formation [3], and the possibility of such objects is certainly a generic prediction of superstring theory!

In this paper we will study a type of cosmic string with many connections to superstrings and axion strings. These strings are based on an "anomalous" $U(1)$ gauge symmetry which arises naturally in a large class of superstring compactifications. This $U(1)$ appears to be anomalous when only the fermion spectrum is considered, but in fact, the anomalies are cancelled by a variant of the anomaly cancellation mechanism of Green and Schwarz. This cancellation of anomalies is intimately related to modular invariance of the underlying superstring theory [4]. It is known that superstring compactifications can lead to only one such $U(1)$ factor, which we will refer to as the pseudo-anomalous $U(1)$.

In supersymmetric compactifications of the heterotic string this $U(1)$ factor leads to supersymmetry

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breaking at the one-loop level [5]. However, in all known cases it is possible to restabilize the vacuum and restore supersymmetry by giving appropriate expectation values to scalar fields which are charged under this $U(1)$ [5,6]. This spontaneous breaking of the $U(1)$ leads to topologically stable cosmic string solutions.

These strings turn out to have many of the same properties as axion strings. However, they are not the same as macroscopic superstrings; in particular, their string tension can be appreciably smaller than the superstring tension. In some models of string compactification, these strings could exist today as superconducting cosmic strings which would exhibit unique optical effects. In other models, we will argue that these strings would be confined by QCD effects as would macroscopic superstrings. In this case, however, there is a linear combination of these two types of string which does not couple to QCD and which would not be confined^{#1}. This hybrid string has a coaxial structure consisting of a superstring core with left-moving currents transforming under the $SU(3)$ color group and an outer sheath with the structure of a $U(1)$ vortex with right-moving currents that also carry $SU(3)$ quantum numbers and which cancel the color anomaly of the resulting 1+1 dimensional effective theory.

Because of the close similarities with axion strings, we will start with a brief review of some of their properties. We start by discussing a novel optical effect that is exhibited by axion strings.

2. Axion strings

As a simple model which serves to illustrate the main physical features, consider an unbroken $U(1)$ gauge symmetry Q , which we will think of as electromagnetism, and a chiral global symmetry A which is spontaneously broken by the expectation value of a scalar field Φ with $Q=0$ and $A=-2$. In addition we will have two left handed Weyl fermions, ψ_L with charges $Q=1$, $A=1$, and χ_L with charges $Q=-1$, $A=1$. (In four component notation, χ_L is the left-handed antiparticle of ψ_R .) These fermions obtain a

mass through their coupling to Φ . The lagrangian is given by

$$L = i\psi_L^\dagger \not{D}\psi_L + i\chi_L^\dagger \not{D}\chi_L - \lambda(\Phi^*\Phi - v^2)^2 + f(\Phi\epsilon^{\alpha\beta}\psi_{\alpha L}\chi_{\beta L} + \text{h.c.}) - (1/4e^2)F^{\mu\nu}F_{\mu\nu}, \quad (2.1)$$

where $F^{\mu\nu} = \partial^\mu A^\nu - \partial^\nu A^\mu$ is the field strength for the Q gauge symmetry.

As a result of the breaking of A , there exist topologically stable vortex solutions with $\Phi = f(\rho) \times \exp[i\theta(\phi)]$, where $f(\rho)$ approaches zero at the origin and v at infinity, v denoting the vacuum expectation value of Φ . For a vortex with winding number n , we can take $\theta(\phi) = n\phi$, where ϕ is the azimuthal angle around the vortex which we take to lie along the z axis. Because of the angular dependence of the phase of Φ , these vortices have a logarithmic divergence in their energy per unit length. In a cosmological context, this divergence would be cut off by the separation between strings of opposite winding number. The phase of Φ is a Nambu-Goldstone boson with axion-like couplings.

It is well known that there are striking electromagnetic effects in the presence of gradients in the axion field which couples to $\mathbf{E} \cdot \mathbf{B}$. This coupling arises via the four-dimensional chiral anomaly from a redefinition of the fermion fields which removes θ , the phase of Φ , from the fermion couplings. This results in a coupling

$$- (e^2 f / 32\pi^2) \theta \epsilon^{\mu\nu\lambda\rho} F_{\mu\nu} F_{\lambda\rho}, \quad (2.2)$$

where we have restored powers of the coupling constant and f is a model dependent factor ($f=1$ in the model presented here). In the presence of a static but spatially varying axion field, this coupling leads to additional contributions to the charge and current densities

$$\rho_\theta = (\alpha f / \pi) \nabla \theta \cdot \mathbf{B}, \quad \mathbf{j}_\theta = -(\alpha f / \pi) \nabla \theta \times \mathbf{E}, \quad (2.3)$$

with $\alpha = e^2 / 4\pi \sim 1/137$.

These terms lead to optical activity in the scattering of electromagnetic waves off axion domain walls [8] and axion strings [9]. In refs. [9,10] this effect was deduced by solving Maxwell's equations in the cylindrically symmetric axion field appropriate to an axion string. For our purposes the following approximate solution is sufficient.

If we assume that the gradient of the axion field

^{#1} The possibility of cancelling the axion field between different types of string has also been considered in ref. [7].

$I = \nabla\theta$ is small and approximately constant, then it is easy to solve the modified Maxwell equations

$$\begin{aligned}\nabla \cdot \mathbf{E} &= (f\alpha/\pi) \mathbf{I} \cdot \mathbf{B}, \\ \nabla \times \mathbf{B} - \partial \mathbf{E} / \partial t &= - (f\alpha/\pi) \mathbf{I} \times \mathbf{E}, \\ \nabla \cdot \mathbf{B} &= 0, \\ \nabla \times \mathbf{E} + \partial \mathbf{B} / \partial t &= 0.\end{aligned}\quad (2.4)$$

Without loss of generality we can take $\mathbf{I} = I\hat{z}$ and use translation invariance in the x - y plane and in time plus rotation invariance in the x - y plane to write the solution in the form

$$\begin{aligned}\mathbf{E} &= \mathcal{E}(z) \exp(ikx - i\omega t), \\ \mathbf{B} &= \mathcal{B}(z) \exp(ikx - i\omega t).\end{aligned}\quad (2.5)$$

Setting $\lambda_0 = \sqrt{\omega^2 - k^2}$ and $\lambda_{\pm}^2 = \lambda_0^2 = \alpha f l \lambda_0 / \pi$, the general solution can then be written in the form (assuming $\mathbf{I}$ constant)

$$\begin{aligned}\mathcal{E}_x &= A_+ \exp(i\lambda_+ z) + A_- \exp(i\lambda_- z), \\ \mathcal{E}_y &= -i(\lambda_0/\omega) [A_+ \exp(i\lambda_+ z) - A_- \exp(i\lambda_- z)], \\ \mathcal{E}_z &= -(k/\lambda_0^2) [\lambda_+ A_+ \exp(i\lambda_+ z) \\ &\quad + \lambda_- A_- \exp(i\lambda_- z)], \\ \mathcal{B}_x &= i(\lambda_0/\omega^2) [\lambda_+ A_+ \exp(i\lambda_+ z) \\ &\quad - \lambda_- A_- \exp(i\lambda_- z)], \\ \mathcal{B}_y &= (\omega/\lambda_0^2) [\lambda_+ A_+ \exp(i\lambda_+ z) \\ &\quad + \lambda_- A_- \exp(i\lambda_- z)], \\ \mathcal{B}_z &= -i(k\lambda_0/\omega^2) [A_+ \exp(i\lambda_+ z) \\ &\quad - A_- \exp(i\lambda_- z)].\end{aligned}\quad (2.6)$$

Here A_+ and A_- are arbitrary constants which represent the amplitudes for right- and left-handed polarized waves in the case $k=0$. Expanding this solution out to order α in this case shows that the direction of polarization of a linearly polarized wave rotates by an amount proportional to the gradient of θ along the direction of motion as the wave propagates. For the case of interest here, with $\mathbf{I}$ slowly varying, we can find the net rotation δ by integrating this rotation along the path of the plane wave, $\delta = (f\alpha/2\pi)\Delta\theta$.

This effect is analogous to Faraday rotation. If we diagonalize the wave equation following from (2.4) by working in terms of $\mathbf{F}_{\pm} = \mathbf{E} \pm i\mathbf{B}$ we see that right-

and left-hand circularly polarized waves obey different dispersion relations:

$$\omega_{\pm}^2(k) = k^2 \pm (f\alpha/\pi)(\mathbf{I} \cdot \mathbf{k}). \quad (2.7)$$

As in Faraday rotation, this leads to different indices of refraction for right- and left-handed circularly polarized waves and hence a net rotation of the plane of polarization for a linearly polarized wave. In contrast to Faraday rotation, the effect here is independent of frequency.

Local strings with string tension μ will give rise to a conical spacetime geometry with deficit angle $8\pi G\mu$. Because of this, such a string would act as a gravitational lens and one would see two images of a distant object with an angular separation of $8\pi G\mu$. The situation is somewhat different for global strings such as the axion strings considered here. Such strings have a non-vanishing energy density outside the core of the string. As a result, there is non-vanishing spacetime curvature outside the core of the string. In this case one can show that the gravitational lensing is quite similar to that of local strings, at least for relativistic particles, except that the string tension μ should be replaced by an effective string tension μ_{eff} which includes the energy inside a radius b , where b is the impact parameter [11]. Since μ_{eff} grows logarithmically with b , in a cosmological situation this leads to lensing effects with a deficit angle which is of order 100 larger than that expected based on the core string tension.

Since the light rays forming the two images pass on opposite sides of the string, the integrated gradient of the axion field differs by 2π along the two paths. Integrating the previous solution along the two different paths we thus find that the polarization of each image will be rotated by $f\alpha/2$ but in opposite directions so that there will be a net difference in the direction of polarization of the two images of $f\alpha$ radians^{#2}.

^{#2} This can also be seen directly in the cylindrical solution in ref. [9], in which δ_m ought to be defined as $(1 - \delta_{m,0}) \text{sgn } m$. The sign of the rotation depends on the sign of the angular momentum m of the partial wave. Light rays passing on opposite sides of the string have angular momenta of opposite signs. As for the $m=0$ component of the wave, which is not directly relevant here, an incoming wave polarized in the $\hat{\phi}$ direction is not scattered, but as pointed out in ref. [10], and contrary to ref. [9], a $\hat{z}$ polarized wave undergoes a phase shift $\exp(i f \alpha)$.

Although the amount of rotation is model dependent, it could easily be on the order of a few degrees. Luckily, the sources for the known gravitational lenses, namely quasars, are known to exhibit fairly strong polarization effects at radio wavelengths. If a cosmic string lens candidate were found, it appears feasible to measure the polarization to this accuracy after correcting for Faraday rotation [12].

It is worth emphasizing that effects such as these can only occur for global strings, not for local strings. The reason is that outside a local string one is in the vacuum so there can be no effect on the propagation of particles outside the string (except of course for the standard lensing effect). For a global string, on the other hand, the fields outside are not gauge equivalent to a vacuum configuration and these fields can affect particle propagation.

3. Pseudo-anomalous U(1) strings

In this section we consider gauging the global U(1) symmetry A of the previous section. To do this we must discuss cancellation of the anomalies involving A as well as any other gauge groups in the theory. We will then discuss the effect of the anomaly cancelling terms on the structure of the strings resulting from the spontaneous breaking of A.

We will consider a somewhat more general class of models which contain in addition to A a set of gauge symmetries which we denote collectively as G. We will assume that there are various anomalies involving A and fields in G, but no anomalies involving fields entirely in G. The left-handed fermions will be in some representation r of G (which may be reducible). If we decompose r into irreducible representations, $r = \oplus r_i$, then each irreducible component r_i will in general have a different charge a_i under A. Let λ_i be a representation matrix for the representation r_i of G. For simplicity we will assume that there is a single scalar field Φ which acquires a non-zero vacuum expectation value which leads to spontaneous breaking of A. G may also be broken to a subgroup by the expectation value of Φ but we will focus for now on the breaking of A.

In order to gauge the A symmetry we must first remove the anomalies involving A. One simple way to do this is to introduce additional fermion fields in

appropriate representations of G and A so as to cancel the anomalies. This is the model of cosmic string superconductivity with Fermi charge carriers considered in ref. [13]. There is, however, another possible mechanism for cancelling U(1) anomalies in four dimensions which is closely related to the Green-Schwarz mechanism. In this section we will give an effective field theory description of the anomaly cancellation and its effect on the structure of vortices created by the spontaneous breaking of A. In the next section we will see how models very similar to this occur naturally in a large class of compactifications of the heterotic superstring which preserve spacetime supersymmetry.

We can choose to define the anomalous variation of the measure coming from the fermions so that the measure is invariant under G and Lorentz transformations. There is then an anomalous variation of the effective action under A gauge transformations $\delta_A A_A = dA_A$ given by

$$\delta_A S = - \frac{1}{32\pi^2} \int (c_A F_A^2 + c_G \text{tr} F_G^2 - c_L \text{tr} R^2) A_A, \quad (3.1)$$

where the constants c_A , c_G , and c_L are given in terms of the fermion content by

$$c_A = \frac{1}{3} \sum_i a_i^3 \dim r_i, \quad c_G = \sum_i a_i \text{tr} \lambda_i^2, \\ c_L = \frac{1}{24} \sum_i a_i \dim r_i. \quad (3.2)$$

To cancel the anomaly (3.1) we follow ref. [14] and introduce a two-form potential B with a variation under G, A, and Lorentz transformations:

$$\delta B = -M^{-1} [\omega_{2A}^1 + (c_G/c_A) \omega_{2G}^1 \\ - (c_L/c_A) \omega_{2L}^1], \quad (3.3)$$

where M is an arbitrary parameter with dimensions of mass. The quantities ω_{2A}^1 , ω_{2G}^1 , ω_{2L}^1 are defined as usual by the descent equations, e.g. $F^2 = d\omega_3$, $\delta\omega_3 = d\omega_2^1$. The anomalies are then cancelled by the addition of the term

$$- \frac{1}{32\pi^2} \int M c_A B F_A + A_A (c_G \omega_{3G} - c_L \omega_{3L}). \quad (3.4)$$

The term involving $B F_A$ is independent of any ambiguity in defining the anomaly. The other terms de-

pend on an arbitrary choice of how the variation under Q and Lorentz transformations of the BF_A term is cancelled between the anomaly from the fermions and the $A_A\omega_{3Q}$, $A_A\omega_{3L}$ terms. As already mentioned, the choice made here is that the fermion anomaly has no contribution from G or Lorentz transformations.

To complete the lagrangian we must add a kinetic energy term for B of the form

$$\int H^*H, \quad (3.5)$$

where $H = dB + M^{-1}[\omega_3^A + (c_G/c_A)\omega_3^Q - (c_L/c_A)\omega_3^L]$ is the gauge invariant field strength for B . H satisfies a Bianchi identity

$$dH = M^{-1}[F_A^2 + (c_G/c_A) \text{tr} F_G^2 - (c_L/c_A) \text{tr} R^2], \quad (3.6)$$

and the equation of motion

$$d^*H + (1/64\pi^2)Mc_A F_A = 0. \quad (3.7)$$

This allows us to trade the two-form B for a zero-form b since (3.7) implies that on shell

$$^*H + (1/64\pi^2)c_A M A_A = db. \quad (3.8)$$

The Bianchi identity (3.6) then becomes the b equation of motion

$$d^*db = M^{-1}[F_A^2 + (c_G/c_A) \text{tr} F_G^2 - (c_L/c_A) \text{tr} R^2] + (1/64\pi^2)c_A M d^*A_A, \quad (3.9)$$

and the H equation of motion (3.7) is now the Bianchi identity for db .

Since H was gauge invariant, we see from (3.8) that b must vary under A gauge transformations as $\delta_A b = (1/64\pi^2)c_A M A_A$, where A_A is the parameter of the gauge transformation. We can now rewrite the anomaly cancelling terms and the H kinetic energy term in terms of b . The H kinetic energy term becomes

$$\int [^*db - (1/64\pi^2)c_A M^*A_A] \times [db - (1/64\pi^2)c_A M A_A], \quad (3.10)$$

which clearly indicates that the gauge field A_A becomes massive by eating the b degree of freedom [15]. The anomaly cancelling term can be rewritten using (3.8) as

$$\begin{aligned} & -2 \int [^*db - (1/64\pi^2)c_A M^*A_A] \\ & \times [db - (1/64\pi^2)c_A M A_A] \\ & + 2M \int b[F_A^2 + (c_G/c_A) \text{tr} F_G^2 - (c_L/c_A) \text{tr} R^2], \end{aligned} \quad (3.11)$$

and we see that the second term has a variation under A gauge transformations which cancels the anomalous variation (3.1).

In the models we are interested in, with A spontaneously broken, this is not the whole story since the vacuum expectation value of Φ also give a mass to A_A . This gives rise to a coupling $\int v^2(d\theta - gA_A)^* \times (d\theta - gA_A)$, where θ is the phase of the field Φ . The gauge field A_A therefore eats a linear combination of θ and b and acquires a mass of order $M + v$.

Gauging the A symmetry modifies the structure of the vortex solution since the A gauge field now plays a significant role. If A were an ordinary local $U(1)$ symmetry, the gauge field would assume the value $A_A = (1/g)d\theta$ outside the core of the vortex, corresponding to pure gauge away from the string and a magnetic field with total flux $2\pi/g$ running along it. This would ensure that the covariant derivative $(d - igA)\Phi$ vanishes outside the core, so the vortex solution describes a local string with finite energy per unit length.

In the case of a pseudo-anomalous symmetry A , however, the magnetic field at the core of the vortex acts as a source for the field strength H (3.7) which falls off as $1/\rho$ outside the string, leading to a logarithmic divergence in the energy per unit length of the string. Thus, breaking a local $U(1)$ symmetry leads to a vortex with a global structure because of the presence of the antisymmetric tensor field needed for anomaly cancellation. To see this another way, the core of the vortex acts as a source of θ but not of b , hence θ but not b has non-trivial winding number around the string. The gauge field A_A can be chosen so that either $[db - (1/64\pi^2)c_A M A_A]$ or $(d\theta - gA_A)$, but not both, vanishes outside the string. In fact, A_A assumes the intermediate value $(\lambda/g)d\theta$, $0 < \lambda < 1$, determined by the ratio of the mass scales M and v . Minimizing the coefficient multiplying the (logarithmically divergent) energy per unit length of the string yields $\lambda^{-1} = 1 + (Mc_A/64\pi^2gv)^2$. The induced

Q-gauge current (2.3), which is responsible for the optical activity described in the last section is independent of λ , however, and is exactly the same as that of an ungauged axion string. Of course, as mentioned before, it is because the vortices formed by the breaking of the pseudo-anomalous symmetry A by the field Φ are global rather than local that they can exhibit optical activity.

The introduction of the anomaly cancelling counterterms and the arbitrary mass scale M may appear rather artificial from the field theory point of view. They arise naturally, however, in superstring compactification as we will now discuss.

4. Pseudo-anomalous U(1) strings in superstring compactifications

In this section we will describe how the model discussed in the previous section can be imbedded in a large class of superstring compactifications. Although the simple field theory model was perfectly consistent, the introduction of the antisymmetric tensor field and its role in anomaly cancellation occurs most naturally in the context of superstring theory.

The simplest Calabi–Yau compactifications of the $E_8 \otimes E_8$ heterotic superstring and their orbifold counterparts do not lead to pseudo-anomalous U(1) factors in the low-energy effective field theory, but many other compactifications do. Examples include Calabi–Yau compactifications of the Spin(32) heterotic string and many orbifold compactifications of the $E_8 \otimes E_8$ heterotic string. In all these cases one can diagonalize the U(1) factors so that there is a single pseudo-anomalous U(1) factor A and a set of U(1) and non-abelian gauge groups G which have anomalies only with A. The anomaly cancelling terms described earlier can be shown to arise from string loop effects but with some stringent constraints on the spectrum. In particular, in string theory the transformation law of the two-form B must take the form $\delta B = \omega_{2A}^1 + \omega_{2G}^1 - \omega_{2L}^1$ and the fermion content must be such that the constants c_A , c_G , and c_L are all equal [4].

It was argued in ref. [5] that the presence of a pseudo-anomalous U(1) in the low-energy limit of a supersymmetric string compactification will lead to

the generation of a D term which will destabilize the vacuum. The terms in the effective lagrangian involving the auxiliary D field for the anomalous U(1) have the form

$$L_D = \frac{1}{2} D^2 + \exp(-3\phi) cD + \exp(-\phi) D \sum_i q_i \chi_i^* \chi_i, \quad (4.1)$$

where the χ_i are a set of scalar fields with charges q_i under the U(1), and ϕ is the dilaton field. The coefficient c can be determined by a one-loop string calculation [16,17], with the result $c = (g/192\pi^2) \times \sum_i n_i q_i h_i$, where g is the coupling constant and n_i is the number of fermions with U(1) charge q_i and helicity $h_i = \pm 1$. If c is non-zero then D has a one-loop expectation value which breaks supersymmetry and destabilizes the vacuum. In all known examples where this occurs it is possible to find a new stable, supersymmetric vacuum by giving expectation values to some subset of the scalar fields χ_i which carry charges under the pseudo-anomalous U(1) factor [5,6]. This spontaneous breaking of the U(1) symmetry will then lead to vortices of the type described in the previous section.

The structure of such a vortex is rather model dependent. The χ expectation values will in general give masses to some subset of the fermions which are in real representations of the remaining unbroken gauge group. The vortex will act as a source for an axion field with couplings to the unbroken gauge fields which are determined directly by the quantum numbers of the fermions which acquire a mass in the U(1) symmetry breaking and thus have zero-modes on the vortex. The coupling of the axion field outside the vortex to F^2 of the various gauge groups arises from the jacobian factor when one performs a chiral rotation on the massive fermions. Equivalently, we can argue that the fermion zero modes on the vortex will lead to anomalies in the (1+1)-dimensional theory on the vortex which must be cancelled by charge inflows resulting from the coupling of the vortex axion field to the appropriate gauge group [18]. So, for example, if none of the fermion zero modes carry SU(3) color quantum numbers then we know that the vortex does not couple to QCD. If the vortices couple to QED but not to any confining gauge forces then they could exist as cosmic strings today and would exhibit the optical effects discussed earlier. If these vortices

do couple to QCD, they would be confined, as would macroscopic heterotic strings. There will, however, be a linear combination of the two types of string which does not couple to QCD and which could thus exist today. This makes it at least conceivable that we could one day determine properties of superstring theory by astronomical observations. We will not address the question of whether such "coaxial" strings would be expected to form in a cosmological context.

Let us now work out the energy per unit length of such global strings in superstring theory. First consider macroscopic heterotic strings. As pointed out in ref. [19] the energy per unit length of such strings is not as large as one might imagine. This can be seen by considering a toroidal compactification of one spatial dimension on a circle of radius R , where we imagine R to be of astronomical size. The mass squared of a state with unit winding number is $M^2 = R^2/\alpha'^2$ and its length is $2\pi R$ so that the mass per unit length is $\mu = 1/2\pi\alpha'$, where in the heterotic string the parameter α' is given by $2\kappa^2/g^2$. Here g is the four-dimensional gauge coupling and $\kappa^2 = 8\pi G$. We thus have $G\mu = g^2/32\pi^2$ which is on the order of 10^{-3} for $g^2/4\pi \sim 1/40$.

For the global strings associated with breaking of the pseudo-anomalous $U(1)$ the energy per unit length is on the order of $|\langle\chi\rangle|^2$, where $\langle\chi\rangle$ is the vacuum expectation value of the scalar field which breaks the $U(1)$. Reinstating dimensions, we find from (4.1) that $\mu \sim (1/2\alpha')(g/192\pi^2)a$, where $a = \sum_i n_i q_i h_i$ is a model dependent factor of order 1. Using the previous results we thus find $G\mu \sim \text{few } 10^{-6}$. In many models the breaking of the pseudo-anomalous $U(1)$ is accompanied by breaking of additional non-anomalous $U(1)$ factors [20]. Thus there is the possibility for a rich spectrum of both global and local strings at the scale of interest for galaxy formation in these models.

Finally, it is worth mentioning that aside from macroscopic superstrings and the cosmic strings discussed here, the possibility for topologically stable global strings in superstring compactifications is quite limited. It was argued in ref. [21] that the only global symmetries in superstring compactifications are non-

linearly realized symmetries associated to axion-like fields. We thus expect that any global string arising in superstring compactifications will act as a source for fields with axion-like couplings.

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References

- [1] E. Witten, Phys. Lett. B 153 (1985) 243.
- [2] N. Kaiser and A. Stebbins, Nature 310 (1984) 391.
- [3] J.P. Ostriker, C. Thompson and E. Witten, Phys. Lett. B 180 (1986) 231.
- [4] A.N. Schellekens and N.P. Warner, Nucl. Phys. B 287 (1987) 317; Phys. Lett. B 177 (1986) 317; B 181 (1986) 339.
- [5] M. Dine, N. Seiberg and E. Witten, Nucl. Phys. B 289 (1987) 317.
- [6] L. Dixon and V. Kaplunovsky, unpublished.
- [7] J. Ostriker, C. Thompson and E. Witten, in preparation.
- [8] P. Sikivie, Phys. Lett. B 137 (1984) 353.
- [9] S.G. Naculich, Nucl. Phys. B 296 (1988) 837.
- [10] A. Manohar, Phys. Lett. B 206 (1988) 276; B 209 (1988) 543 (E).
- [11] M. Aryal and A.E. Everett, Phys. Rev. D 33 (1986) 333; D. Harari and P. Sikivie, Phys. Rev. D 37 (1988) 3438.
- [12] J. Gunn and E. Turner, private communication.
- [13] E. Witten, Nucl. Phys. B 249 (1985) 557.
- [14] M. Green and J.H. Schwarz, Phys. Lett. B 149 (1984) 117.
- [15] E. Witten, Phys. Lett. B 149 (1984) 351.
- [16] J.J. Atick, L.J. Dixon and A. Sen, Nucl. Phys. B 292 (1987) 109.
- [17] M. Dine, I. Ichinose and N. Seiberg, Nucl. Phys. B 293 (1987) 253.
- [18] C. Callan and J. Harvey, Nucl. Phys. B 250 (1985) 427.
- [19] N. Turok, Phys. Rev. Lett. 60 (1988) 549.
- [20] J.A. Casas, E.K. Katehou and C. Muñoz, University of Oxford preprint (1988).
- [21] L. Dixon and T. Banks, Princeton preprint PUPT-1086 (1988).