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COSMIC STRINGS FROM HIDDEN-SECTOR FLUX TUBES

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ABSTRACT

We consider the possibilities for obtaining cosmic strings from an underlying $E_8 \times E_8'$ superstring theory. We show that the observable sector cannot provide cosmic strings with a tension large enough to give interesting density perturbations: $\delta\rho/\rho \sim 10^{-4}$. Neither can the hidden sector yield cosmic strings if E_8' is unbroken, while flux tubes give $\delta\rho/\rho$ too large if E_8' is broken to a subgroup bigger than $SU(4)$, and too small if it is broken to a subgroup smaller than $SU(4)$. However, flux tubes in the hidden sector can provide cosmic strings with $\delta\rho/\rho \sim 10^{-4}$ if E_8' is broken to an $SU(4)$ subgroup. If there was inflation, this mechanism would only be operative if the Universe subsequently reheated to a temperature above 10^{16} GeV.

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Cosmic strings have great appeal to astrophysicists for at least three reasons. One reason is that they provide [1] a welcome alternative to the adiabatic perturbations of inflationary cosmology for generating $\delta\rho/\rho \sim 10^{-4}$ in an almost scale-invariant way as desired for galaxy formation. Indeed, it has been claimed [2] that the two-point spatial correlation function for closed string loops matches closely that for clusters of galaxies. A second reason is that galaxy formation models without strings all possess some problem or another [3], an indication that galaxy formation is not very well understood as yet. A neutrino or a hot dark matter dominated Universe has trouble on small scales. Neutrino free streaming [4] wipes out density perturbations on small scales so that galaxies must fragment from larger objects. The trouble is that they tend to do so too late [5]. Cold dark matter, on the other hand, requires biasing [6]. In addition, unless the biasing is due to some scale-free process [7], clustering will not occur on large scales. Again cosmic strings are a nice alternative as they allow [8] structure on small scales, since strings are not free-streaming, and on large scales, if hot particles are also present [8] or if strings are formed in conjunction with inflation [9]. A third reason for liking cosmic strings is that the naked eye, and analyses of correlation functions, suggest [10] that the Universe has a one-dimensional (in the fractal sense) stringy structure. We are not necessarily convinced that cosmic strings are essential to understand the astrophysical data, but they are a sufficiently attractive hypothesis that we would like to know how plausible they are in the context of theoretical particle physics.

Nowadays all of particle physics is supposed to be a low-energy manifestation of the $E_8 \times E_8$ heterotic superstring [11], which is the only serious candidate [12] for a Theory of Everything (TOE). Therefore, it is natural to look for cosmic strings [13] in the four-dimensional theory abstracted from the superstring. Witten has written the definitive introduction [14] to this question. As he points out, there are two mechanisms which could

generate cosmic strings: vortex lines in the observable E_8 gauge sector and flux lines in the hidden E'_8 sector. Vortex lines can arise if E_6 is broken by the Hosotani mechanism [15] to a four-dimensional gauge group H containing an extra $U(1)$ factor orthogonal to the $U(1)$ of electromagnetism. Flux tubes can arise if the hidden E'_8 is broken by the Hosotani mechanism to a four-dimensional gauge group with a non-trivial centre. A third mechanism [13], which can give strings if a general group G is broken to a subgroup H and $\pi_1(G/H) \neq 0$, is not relevant to the superstring for reasons reviewed below.

In this paper we re-examine these possible origins for cosmic strings in the context of recently developed phenomenological intuition about the likely structure of the four-dimensional theory to be obtained from the superstring. For reasons discussed elsewhere [16] and recalled briefly here, we believe that the observable sector subgroup H must have rank 5 [$H = SU(3)_C \times SU(2)_L \times U(1)^2$] [17], in which case it gives uninteresting vortex lines with very small $\delta\rho/\rho$. Even if H were to have rank 6, e.g. $H = SU(3)_C \times SU(2)_L \times U(1)^3$, the vortex lines would be associated with a scalar field vacuum expectation value which the charge 1/3 quark mass matrix tells us [16] must be far too small to provide an interesting magnitude of $\delta\rho/\rho$. The hidden E'_8 could not give flux tubes if it were unbroken [14], but it has been argued elsewhere [18] and we re-emphasize here the desirability of breaking E'_8 by the Hosotani mechanism in order to ensure the gaugino condensation needed for dynamical supersymmetry breaking. If E'_8 is broken to subgroups H' smaller than $SU(4)$, the resulting flux tubes give $\delta\rho/\rho$ too small. If there was no inflation, or if there was inflation and the Universe subsequently reheated to a temperature above the scale of string formation [9], in this case the strong coupling scale for H' , the subgroup $H' = SU(4)$ has flux tubes which would give $\delta\rho/\rho \sim 10^{-4}$ as desired [1]. Subgroups H' larger than $SU(4)$ would have flux tubes giving $\delta\rho/\rho > 10^{-4}$, but it may be possible in some cases to bleach them out and replace them by flux tubes of a smaller subgroup H'' , which could either give $\delta\rho/\rho \sim 10^{-4}$ [if $H'' = SU(4)$] or smaller [if H'' were

smaller [if H' were smaller than $SU(4)$]. Although the discussion of each individual case in the above list is quite simple, there are many cases to consider, and the reader might become disoriented unless she consults the map of our arguments which is presented in Fig. 1.

We start our analysis with the observable sector. We assume that E_6 is broken by Wilson lines to some subgroup H of rank 5 or 6. We focus initially on the possibilities $H = SU(3)_C \times SU(2)_L \times U(1)^2$ or 3 , which have been favoured in most phenomenological work [12], and return later to other rank 6 possibilities. Vortex lines arise in both of these favoured groups when the additional $U(1)$ or $U(1)^2$ is broken by the expectation value of one or two scalar fields respectively. In the rank 5 case $SU(3)_C \times SU(2)_L \times U(1)^2$, this field is the $SU(5)$ - and $SO(10)$ -singlet field N in the $\underline{27}$ of E_6 , which gives (through $DD^C N$ couplings) masses to the new colour-triplet charge $-1/3$ 'quarks' D that come from the $\underline{10}$ of $SO(10)$ in the $\underline{27}$ of E_6 , and also mixes (through $H\bar{H}N$ couplings) the H and $\bar{H}$ Higgs fields of the supersymmetric Standard Model [12]. This mixing must be of order m_W if the correct electroweak vacuum with $v \equiv \langle 0|H|0 \rangle$ and $\bar{v} \equiv \langle 0|\bar{H}|0 \rangle = 0(m_W)$ is to be obtained, so $x \equiv \langle 0|N|0 \rangle = 0(m_W)$ also. The tension σ in the corresponding vortex lines would be

$$\sigma_x \sim x^2 \sim O(10^{-32}) m_P^2 \quad (1)$$

which would yield

$$\frac{\delta \rho}{\rho} \sim \sigma / m_P^2 \ll 10^{-4} \quad (2)$$

In the rank 6 case $SU(3)_C \times SU(2)_L \times U(1)^3$, one also needs $y \equiv \langle 0|\tilde{v}^C|0 \rangle \neq 0$, where $\tilde{v}^C$ is the $SU(5)$ singlet field contained in the $\underline{16}$ of $SO(10)$, which is commonly identified with the right-handed neutrino. As such, we have argued [16] that it is not possible to generate $y \neq 0$ without giving unacceptable

neutrino masses and/or lepton-number violation and/or flavour-changing neutral interactions, as well as having the unpleasant feature of violating weak universality in principle if not in practice. However, these arguments rely to some extent on our dynamical scenario for gauge symmetry breaking by radiative corrections [12], and one might wish to retain the possibility that $y \neq 0$. Even so, there is a simple argument that y cannot be large enough to give interestingly large perturbations $\delta\rho/\rho$. The reason [16] is that y induces (through a $D\bar{d}^c v^c$ coupling) mixing in the d/D mass matrix:

$$(\bar{d}^c, \bar{D}^c) \begin{pmatrix} O(\bar{v}) & O(y) \\ 0 & O(x) \end{pmatrix} \begin{pmatrix} d \\ D \end{pmatrix} \quad (3)$$

When we diagonalize (3) by a biunitary transformation we find that

$$m_d = O\left(\frac{\bar{v}x}{y}\right) \quad (4)$$

Since $m_d \geq m_e = O(\bar{v})$, we cannot allow y to be much larger than x . Therefore, $\sigma_y \sim y^2$ also implies $\delta\rho/\rho \ll 10^{-4}$. Although derived [16] in the context of the rank 6 group $SU(3)_C \times SU(2)_L \times U(1)^3$, similar arguments about the vanishing or smallness of $y = \langle 0 | \bar{v}^c | 0 \rangle$ would also apply in the context of other rank 6 observable groups. So these also cannot yield interestingly large perturbations $\delta\rho/\rho$.

One can also obtain cosmic strings in a conventional four-dimensional grand unified theory [13] if the original gauge symmetry G is broken down to a subgroup H : $\pi_1(G/H) \neq 0$. In our case the relevant group G is E_6 , which excludes this possibility. This follows from the facts that the following homotopy sequence is exact [13]

$$\rightarrow \pi_n(H) \rightarrow \pi_n(G) \rightarrow \pi_n(G/H) \rightarrow \pi_{n-1}(H) \rightarrow \pi_{n-1}(G) \rightarrow \dots \rightarrow \pi_0(G/H) \quad (5)$$

and $\pi_0(E_6) = \pi_1(E_6) = 0$, while the connectedness of the unbroken subgroup H means that $\pi_0(H) = 0$. Substituting these special cases into the general sequence (5), we see that

$$0 = \pi_1(E_6) \rightarrow \pi_1(E_6/H) \rightarrow \pi_0(H) = \pi_0(E_6) = 0 \quad (6)$$

which implies that $\pi_1(E_6/H) \stackrel{\sim}{=} \pi_0(H) = 0$. Thus the Wilson line breaking of E_6 down to H yields no vortices. The same argument also tells us that in the hidden sector $\pi_1(E_8/H') = 0$, so that the only mechanism for creating cosmic strings in the hidden sector is through flux tubes in the unbroken hidden-sector subgroup H' .

Such flux tubes cannot [14] arise unless E_8' is indeed broken to some smaller subgroup H' , because E_8 has no centre. Indeed, it has previously been argued that E_8' should be broken by Wilson lines in order to guarantee the hidden-sector gaugino condensation required for supersymmetry breaking [18]. We now repeat that argument and add some considerations which increase its weight. To break supersymmetry dynamically we need the condensation of gaugino pairs: $\langle 0 | \bar{\chi} \chi | 0 \rangle \neq 0$. Among the non-perturbative configurations which lead to condensates are instantons which have topological number 1, and torons which can appear [19] if the gauge group has a non-trivial centre, and have topological numbers k/N ($k = 1, 2, \dots, N - 1$) in the case of an $SU(N)$ gauge group with centre Z_N . The chiral selection rules tell us that instantons in unbroken E_8' would lead to the condensation of a product of $C(248) = 60$ gauginos [where here and subsequently quadratic Casimir invariants C are defined by $C(R)\delta_{ij} \equiv \sum_{a \in R} (T^a T^a)_{ij}$], which is not exactly what we want [18]. Instantons in $SU(N)$ would give us a condensate of $C(N^2 - 1) = 2N$ gauginos, again not what we want, but torons with $k = 1$ would give us a condensate of $1/N C(N^2 - 1) = 2$ gauginos as desired [18]. It is easy to check that torons in E_6 , E_7 , and $SO(M)$ groups do not yield $\langle 0 | \bar{\chi} \chi | 0 \rangle \neq 0$. Although

the presence of torons is sufficient to ensure $\langle 0 | \bar{\chi} \chi | 0 \rangle \neq 0$, one might question its necessity. In the case of colour $SU(3)$, we know that $\langle 0 | \bar{q} q | 0 \rangle \neq 0$, and might guess that quark condensation would be automatic in any confining theory. However, this is not necessarily the case, since the anomaly-matching conditions of 't Hooft [20] could, in principle, have been solved equally well with $\langle 0 | \bar{q} q | 0 \rangle = 0$ and $\langle 0 | (\bar{q} q)^n | 0 \rangle \neq 0$ for some $n > 1$. Perhaps this is exactly what would have happened if the gauge group of quantum chromodynamics (QCD) had not been one of the $SU(N)$ series? Our point of view is that gaugino condensation: $\langle 0 | \bar{\chi} \chi | 0 \rangle \neq 0$ and hence dynamical supersymmetry breaking have no reason to occur if E'_0 is not broken to some $SU(N)$ subgroup.

If, in view of the preceding arguments, we assume that E'_0 is Wilson line broken to a subgroup H' with non-Abelian factors, then we expect these non-Abelian factor groups to confine at lower scales determined by their renormalization group equations. If these factor groups have non-trivial centres, then their confinement will give rise to flux tubes of gauge field flux [21]. If Λ is the renormalization group scale parameter defining the scale of strong coupling, then we expect that the string tension σ of the resulting flux tubes will be $\sim (4\Lambda)^2$ as in QCD^{*}). If there was inflation, these flux tubes would not be present in significant numbers if they were formed beforehand, or if the Universe after inflation reheated to a temperature below Λ [9]. If these pitfalls are avoided, in order for these flux tubes not to disturb [23] the isotropy of the 3K background radiation we need $G_N \sigma \lesssim 4 \times 10^{-6}$, and to be useful [1, 23] for galaxy formation, we require $G_N \sigma \gtrsim (1 \text{ to } 4) \times 10^{-6}$ which suggests a scale parameter $\Lambda \gtrsim (2 \text{ to } 5) \times 10^{15}$ GeV. We will regard as interesting models for which Λ is within a factor 3 of this central value. To estimate the scale parameters for the various non-Abelian factors in the hidden sector H' we use [24] information coming from the observable sector. One uses renormalization group equations in the observable sector to

^{*}) This agrees with an estimate based on calculations [22] of the electric gluon mass in pure supersymmetric Yang-Mills theory at high temperature.

compute the unification scale at which the coupling constants merge, and the value of the coupling at unification. In the phenomenologically favourable class of models with three 27 families surviving to $\lesssim 1$ TeV, and an observable sector gauge group $SU(3) \otimes SU(2) \otimes U(1)^2$ or 3 , this gives [17] a unified coupling constant $\alpha_G \simeq 0.11$, and a unification scale $M_X \simeq 4 \times 10^{17}$ GeV, which is essentially also the scale of Wilson line symmetry breaking in the hidden sector. Hence the coupling there is the fundamental observable sector E_6 coupling, which is in turn equal to the hidden sector E'_8 coupling at the compactification scale M_C . This must be between $M_X \simeq 4 \times 10^{17}$ GeV and $M_P/\sqrt{8\pi} = 2 \times 10^{18}$ GeV. The Hosotani breaking of E'_8 to H' also occurs around $M_X \simeq 4 \times 10^{17}$ GeV. We ignore the possible variations in the E_6 and E'_8 couplings between M_X and M_C , which would in any case tend to compensate. Thus we take as the boundary condition for the couplings in the unbroken hidden sector subgroup H' that they should have the value $\alpha_G \simeq 0.11$ at 4×10^{17} GeV^{*)}.

There is a variety of possible subgroups H' , to which the hidden sector E'_8 could be Hosotani broken at compactification [18, 24]. Maximal subgroups of E_8 include: $SO(16)$, $SU(9)$, $E_7 \otimes SU(2)$, $E_6 \otimes SU(3)$, and $SU(5) \otimes SU(5)$. In view of our discussion above of using hidden-sector $SU(N)$ groups to provide supersymmetry breaking through toron-induced condensates, we shall focus on $SU(N)$ subgroups ($N \leq 9$) that may arise in the symmetry-breaking chain^{**)} . Now, having performed the Hosotani breaking in the hidden sector, the only massless particles remaining in that sector are the gauge bosons of H' and their gaugino partners (recall that on compactification, the only massless hidden-sector particles were E'_8 gauge bosons and fermions). The gauge bosons in E_8/H' and their fermion partners will have all gained a mass of order $\simeq M_X$

*) Analysis of the electric gluon mass [22] indicates that this is in the weak coupling regime for H' .

**) Since the only attribute of the group which will enter our subsequent analysis is the Casimir of the adjoint representation, it is trivial to redo our analysis for any other choice of factor group in H' .

from the symmetry-breaking. So in examining the renormalization group equation for any simple group in H' , the only matter representation which appears is the adjoint representation of Majorana gauginos.

Then for asymptotically free non-Abelian factors appearing in H' , we have the solution to the renormalization group equation:

$$\alpha_g(Q^2) \simeq \frac{24\pi}{9C \ln(Q^2/\Lambda^2)} \quad (7)$$

with C the Casimir invariant for the adjoint representation of the simple group in question. [In our normalization $C = 2N$ for gauge group $SU(N)$.] The boundary condition that $\alpha_G(M_C^2) \simeq 0.11$, when inserted in eq. (7), allows us to solve for M_C/Λ , for the non-Abelian factor group(s) in H' . Combined with the estimate [17] for $M_C \simeq 4 \times 10^{17}$ GeV, this gives us estimates for the scale parameters Λ associated with the H' factor groups. We show in the table the values of M_C/Λ , and Λ (assuming $M_C \simeq 4 \times 10^{17}$ GeV), for the different $SU(N)$ groups which could appear as factors in H' , as well as the magnitude of the cosmological density perturbations which would be produced by flux tubes of that gauge group. We see that for gauge groups smaller (in the sense of adjoint Casimirs) than $SU(4)$, the density perturbations produced by the flux tubes of the group are too small to be cosmologically relevant. A hidden-sector $SU(4)$ group would give flux tubes which could usefully seed galaxy formation. But for hidden-sector gauge groups larger (in the sense of adjoint Casimirs) than $SU(4)$, the density perturbations which would be produced by flux tubes of the gauge group would be too large, and would induce anisotropies in the cosmic microwave background radiation in conflict with experiment [23]. However, we should note at this point that while our estimates of M_C/Λ are reasonably insensitive to changes in $\alpha_G(M_C^2)$, the result we get for Λ does vary as M_C , and shares the numerical uncertainty in its determination. Hence models with a different value of

M_C might prefer a different subgroup for gaugino condensation. Varying the estimated ratio of σ/Λ^2 would have a similar effect.

One might wonder whether a process of 'bleaching' could occur, through which vacuum polarization effects could convert a flux tube of a gauge group that has a large adjoint Casimir, and hence high confinement scale and string tension, into a flux tube of a gauge group with lower Casimir and string tension. For example, in the breaking of E_8' to $SU(3) \otimes E_6$ the adjoint of E_8' decomposes as $\underline{248} = (\underline{8}, \underline{1}) + (\underline{1}, \underline{78}) + (\underline{3}, \underline{27}) + (\bar{\underline{3}}, \bar{\underline{27}})$, which contains states in the fundamental (antifundamental) representation under both $SU(3)$ and E_6 . If we now had an E_6 flux tube, with the fundamental unit of flux, it could be broken by vacuum polarization of a $(\underline{3}, \underline{27}), (\bar{\underline{3}}, \bar{\underline{27}})$ pair, as illustrated in fig. 2. The pair would act as the two ends of the broken $\underline{27}$ tube and be dragged off in opposite directions by the high $\underline{27}$ tension. Since these particles carry opposite $SU(3)$ fundamental charges, they will leave behind an $SU(3)$ flux tube as they separate. In such a manner high-tension flux tubes of one gauge group might be 'bleached' to low-tension flux tubes of a gauge group with a lower confinement scale.

That such a process is not significant in some cases is easy to see. Representations other than the adjoint gauginos, which are required to produce this effect [such as the $(\underline{3}, \underline{27})$ and $(\bar{\underline{3}}, \bar{\underline{27}})$ of our example], are generated by the Hosotani breaking to the low-energy group. As such, they pick up a mass of the order of the unification scale M_X . Then by exactly the same argument [25] used to discuss the breaking of magnetic vortices by monopole-antimonopole polarization effects, we find that the probability per unit length per unit time to break a flux tube of tension σ by polarization of pairs of particles each of mass M_X is:

$$P \approx M_C^2 f(M_C^2/\sigma) \exp(-\pi M_C^2/\sigma) \quad (8)$$

where f is some dimensionless function of order one.

Because of the exponential suppression in (8), this process is significant for subgroups in H' which confine close to the unification scale. High-tension flux tubes which would have led to $\delta\rho/\rho > 10^{-4}$ are not acceptable if they are bleached after recombination, i.e. after $t \sim 10^6$ s $\sim 3 \times 10^{47} M_X^{-1}$. This happens if the exponential factor in eq. (8) is larger than 10^{-95} , i.e. if

$$\frac{M_c^2}{\Lambda^2} \simeq \frac{16M_c^2}{\sigma} < \frac{16.220}{\pi} \sim 1100 \quad (9)$$

Thus the $SU(N)$ groups with $N \geq 5$ (and other groups with equally large confinement scales) can be bleached, and can lead to viable cosmological scenarios if the other simple subgroups of H' are no larger than $SU(4)$. Interesting density perturbations are again obtained if there is an $SU(4)$ factor.

So we conclude that it is possible to have flux tubes with cosmologically relevant tensions arising from confining non-Abelian gauge groups in the hidden sector of low-energy models based on the $E_8 \otimes E_8'$ heterotic superstring. On the other hand, our arguments have indicated that flux tubes of cosmological interest are unlikely to arise either from the Hosotani breaking of the hidden or observable sector E_8 's, or from spontaneous symmetry breaking in the observable-sector low-energy gauge group.

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Table

Indicated are the possible special unitary groups which could appear as hidden-sector gauge groups, the ratio of the compactification/unification scale to their confinement scale, their confinement scale assuming that the compactification/unification scale is $\simeq 4 \times 10^{17}$ GeV, the string tensions of their resulting flux tubes and the magnitudes of primordial density perturbations that they produce.

$N[SU(N)]$	M_C/Λ	Λ (GeV)	$G\sigma$	$\delta\rho/\rho \simeq 30 G\sigma$
9	7.0	6×10^{16}	3.5×10^{-4}	1.0×10^{-2}
8	8.9	4×10^{16}	2.2×10^{-4}	6.6×10^{-3}
7	12.	3×10^{16}	1.2×10^{-4}	3.5×10^{-3}
6	18.	2×10^{16}	5.0×10^{-5}	1.5×10^{-3}
5	33.	1×10^{16}	1.6×10^{-5}	5.0×10^{-4}
4	79.	5×10^{15}	2.8×10^{-6}	8.4×10^{-5}
3	340.	1×10^{15}	1.5×10^{-7}	4.6×10^{-6}
2	6,100	7×10^{13}	4.5×10^{-10}	1.4×10^{-8}

Figure captions

Fig. 1 : Guide-map to our studies of possible cosmic strings. Phenomenological scenarios that we disfavour are surrounded by parentheses (), while the reasons for disfavoursing them are shown in square brackets []. The preferred scenario is framed by a double box.

Fig. 2 : A high-tension flux tube (a) in the $\underline{27}$ of E_6 is sundered by the vacuum polarization (b) of a pair in the fundamental representations of both E_6 and $SU(3)$, which are pulled apart (c) by the higher string tension, leaving behind a low-tension flux tube in the $\underline{3}$ of $SU(3)$. This 'bleaching' process can also occur in other groups.

Fig. 1

Fig. 2