Flavored cosmic strings

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String solutions related to the spontaneously broken gauge symmetry $SU(3)_H$ of the quark and lepton generations are analyzed. Solutions of a new class are found: binary Z_3-string configurations. The string cosmic networks which they form might make up most of the dark matter in the universe.

1. It is generally believed[1] that a phase transition early in the evolution of the universe might have given rise to cosmic strings: linear topological defects with a tension $\mu \sim V^2$, where V is the vacuum expectation value of the scalar field. This expectation value causes a spontaneous breaking of some original symmetry G of the quarks and leptons. The symmetry G is usually understood as the grand unification symmetry, e.g., $SO(10)$ or $E(6)$ with $V_{GU} \sim 10^{15}$ GeV (Refs. 1 and 2). Heavy cosmic strings associated with this symmetry might have cosmologically significant density fluctuations which induce the formation of galaxies and lead to several other observable effects.[1,2]

In the present letter we examine the horizontal gauge symmetry of quark–lepton generations, G_H, whose spontaneous breaking might (as we will see below) lead to the appearance of flavored cosmic strings characteristic of the generation model. We assume the scale of this breaking to be $V_H \sim 10^4$ GeV, which is the value presently implied by experiments carried out to search for processes involving a nonconservation of quark and lepton flavors.[3]

2. Although our discussion applies equally well to any generational symmetry G_H which satisfies the topological conditions, we will be discussing gauge symmetries of the type $G_H = SU(3)_H$ [or $SU(3)_H \times U(1)_H$] with a chiral filling of quarks and leptons[4] on the basis of the standard $SU(2) \times U(1)$ electroweak model:

$$q_{L\alpha} = \begin{pmatrix} u \\ d \end{pmatrix}_{L\alpha}, \quad q_R^\alpha = (u, d)_R^\alpha, \quad \alpha = 1,\ 2,\ 3\ \ (SU(3)_H) \tag{1}$$

(with a corresponding expression for leptons). One reason for this choice is that these symmetries are the most attractive ones in terms of a possible extension to grand unification models. Another is that it allows one to describe the quark–lepton mass spectrum and quark mixing in a natural way (without any fine tuning).[5,6] In general, the theory should contain, along with the ordinary quarks and leptons, (1), their conjugate fermions Q_L^α and $Q_{R\alpha}$, so that there will be a complete cancellation of G_H^3 anomalies. The Q quarks (and leptons) acquire masses $M_Q \sim V_H$ as a result of the breaking of the symmetry G_H immediately after the scalars $\eta_n^{\alpha\beta}$ [the triplets $\eta_1^{[\alpha\beta]}$, $\eta_2^{[\alpha\beta]}$ and the sextet $\eta_2^{[\alpha\beta]}$ of $SU(3)_H$ symmetry] related to these masses develop their own vacuum expectation values. The masses of the ordinary quarks (the up and down quarks and also the charged leptons) induce effective constraints of the type

$$\frac{1}{M_F}\bar{q}^{\alpha}_{L}q^{\beta}_{R}\varphi\eta_{n\alpha\beta},\quad n=0,\ 1,\ 2, \tag{2}$$

which arise as a result of an exchange between the states $q_{L\alpha}\varphi$ (φ is an ordinary scalar Weinberg–Salam doublet) and the states $q^{\beta}_{R}\eta_{n\alpha\beta}$ of intermediate heavy fermions F_α ($F_{L\alpha}$,$F_{R\alpha}$) with a mass $M_F \sim V_H$ (Ref. 6).

The vacuum expectation values of the horizontal scalars $\eta^{\alpha\beta}_n$ determine the structure of the massive quark and lepton matrices, as well as the nature of the breaking of the symmetry G_H, through constraint (2). A solution of the form[5]

$$<\eta_0^{\alpha\beta}>=\eta_0\delta^{\alpha3}\delta^{\beta3},\quad <\eta_1^{\alpha\beta}>=\eta_1\delta^{\alpha2}\delta^{\beta3},\quad <\eta_2^{\alpha\beta}>=\eta_2\delta^{\alpha1}\delta^{\beta2} \tag{3}$$

for the vacuum expectation values of scalars (in their common Higgs potential), with the hierarchy $\eta_0 \gg \eta_1 \gg \eta_2$, leads to a specific Fritzsch ansatz[7] of massive matrices of quarks with an antisymmetric mixing of generations. We can say that this ansatz presently draws a basically correct picture of the hierarchy of quark masses and mixings observed experimentally, and we can deal with the vacuum expectation values in (3). The effective constraint in (2) and, analogously, all the Yukawa constraints of the model have an additional global chiral symmetry $U(1)_H$ with nonzero values of the horizontal hypercharge Y_H of the fermions q, Q, and F and of the scalars η_n (Ref. 6). We can think of this symmetry as the initial Peccei–Quinn symmetry[8] of all massless fermions (massless before the breaking of the generational symmetry). On the other hand, we might be able to localize it and obtain a complete generational gauge symmetry $SU(3)_H \times U(1)_H$. At any rate, as we will see below, the presence of an additional $U(1)_H$ is important for the topological properties of the model if we start from vacuum manifold (3).

3. We now consider string solutions which are associated with a spontaneously broken generational symmetry $SU(3)_H \times U(1)_H$. According to the general topological arguments, the presence of a string signifies the existence of a group orbit which cannot be contracted to a point in horizontal vacuum manifold (3). This orbit is

$$\Omega(\theta) = \exp[i(a\lambda_8 + b\lambda_3 + cI)\theta], \tag{4}$$

where λ_8, λ_3, and I are diagonal generators of $SU(3)_H \times U(1)_H$. Applying orbit (4) to vacuum expectation values (3), we obtain the conditions ($\eta_n(\theta) = \Omega(\theta)\eta_n(0)$, $n=0,1,2$)

$$-4a + c = l,\quad -a - b + c = m,\quad 2a + c = p, \tag{5}$$

where the indices l, m, and p by definition take on only the values ± 1 and 0, which correspond to the presence and absence, respectively, of a (main) string. Fractional values of the indices would imply the presence of domain walls on the string. It is easy to see from (5) that walls of this sort would necessarily arise in the absence of a hypercharge $U(1)_H (c=0)$:

$$SU(3)_H \xrightarrow{\eta_0} SU(2)_H \times Z_2 \xrightarrow{\eta_1} 1. \tag{6}$$

If there is instead a hypercharge $U(1)_H$, only one "hybrid" string arises, in the final stage of the symmetry breaking

$$SU(3)_H \times U(1)_H \xrightarrow{\eta_0} SU(2)_H \times \tilde{U}(1)_H \xrightarrow{\eta_1} \tilde{\tilde{U}}(1)_H \xrightarrow{\eta_2} 1\ , \tag{7}$$

after a successive mixing with the generators λ_8 and λ_3. In the case of a global $U(1)_H$, the symmetry $\tilde{\tilde{U}}(1)_H$ is the same as the actual Peccei–Quinn symmetry $U(1)_{PQ}$, but the latter is now related to only the fermions of the first generation. Spontaneous breaking of this symmetry gives rise to a corresponding (global) "hybrid" string, which subsequently becomes the boundary of a cosmologically safe axion domain wall as a result of a "hard" breaking of $U(1)_{PQ}$ by nonperturbative QCD effects. Such "hybrid" walls are well known (see the reviews[1,9]). They were recently discussed in connection with generational symmetry in Refs. 10.

Which string configurations are still possible in principle in a generational model with the set of vacuum expectation values in (3)? It is easy to see that in any case in which the scalars which spontaneously break $SU(3)_H$ also have a horizontal hypercharge $U(1)_H$ (whether local or global is irrelevant) we will have only the hybrid strings mentioned above. The explanation is that the discrete symmetry Z_N, which arises in an intermediate stage from the vacuum expectation value of some N-index symmetry scalar $\chi\{\alpha,\beta,...\rho\}$ [antisymmetric scalars—$SU(3)_H$ triplets—do not yield string solutions], is always "absorbed" by the continuous symmetry $\tilde{U}(1)_H$, as in the example considered above, with a two-index scalar $\eta_0^{\{\alpha\beta\}}$ [see Eq. (7)].

We now assume that the scalar χ does not have a hypercharge $U(1)_H$, and that its vacuum expectation value develops from one of the components $\chi_{\alpha\alpha...\alpha}$. We have the chain of symmetry breakings

$$SU(3)_H \times U(1)_H \xrightarrow{\chi} SU(2)_\alpha \times Z_N \times U(1)_H \xrightarrow{\eta_0} \tilde{U}(1)_H \xrightarrow{\eta_1} 1\ , \tag{8}$$

where $SU(2)_\alpha$ is one of the subgroups of $SU(3)_H$ ($\alpha = 1,2,3$). The string solution associated with the scalar χ thus has a Z_N configuration. Since all the previous scalars η_n have a horizontal hypercharge Y_H, domain walls do not form even in this case. Again, a hybrid string forms, but in this case it is associated not with the scalar η_2, as it was in (7), but with the more massive scalar η_1, which mixes the second and third generations. As we see, the condition that there are no domain walls completely determines the form of the new scalar χ and the type of string related to its vacuum expectation value. Applying group orbit (4) to it, we have, in addition to conditions (5),

$$-2Na = k \quad (\alpha = 3), \tag{9}$$

where for definiteness we have taken the vacuum expectation value of the field χ on the component $\chi_{33...3}$. Now using conditions (5), we finally find a relation among the topological indices of the string in this model:

$$l - p = 3\frac{k}{N}\,. \tag{10}$$

From this relation we find the unique nontrivial solution $N = 3$ for the main strings. The same solution arises for other subgroups of $SU(3)_H$ ($\alpha = 1,2$).

In essentially all cases we thus find characteristic topological defects in the generational model: binary Z_3-string configurations which subsequently (in the course of the evolution[1]) form string networks. A small "core" of these strings corresponds to the scalar χ, and a large core (a hybrid string) corresponds to the scalar η_1. If we assume that the symmetry $SU(3)_H$ first goes through a stage of the breaking of the octet of scalars, in a process accompanied by the creation of flavored monopoles, and if also we assume that this symmetry is later broken in accordance with (8), then there would be monopoles and antimonopoles at the nodes of our network.

The light ($V_H \ll V_{GU}$) flavored strings and string networks discussed here might constitute the physical basis of Vilenkin's[11] and Kibble's[12] string-dominated universe. We will discuss the astrophysical and cosmological implications of our model in detail elsewhere. Here we would like to briefly point out some possible manifestations of flavored strings in a string-dominated universe of this sort. These manifestations stem from the circumstance that because of the Yukawa constraints of the model, there will always be a localization of a certain anomaly-free set of transverse zero modes of the fermions q, Q, and F on a hybrid string,[10] in accordance with Witten's general criterion.[2] In the Z_3 case, the networks of this mode correspond to the second and third generations of ordinary and heavy ($M_Q \sim M_F \sim V_H$ off the string) quarks and leptons. After acquiring a sufficient momentum in the intergalactic magnetic fields, they would be injected from the strings of the network and would create very large cosmic-ray extensive air showers (with dimensions on the order of the dimensions of the earth), with an energy on the order of V_H. In a string-dominated universe, the number of such extensive air showers which could be seen at the earth would be on the order of one or two per year.

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[1]A. Vilenkin, Phys. Rep. **121**, 263 (1985).
[2]E. Witten, Nucl. Phys. B **249**, 557 (1985).
[3]Particle Data Group, Phys. Lett. B **239**, 1 (1990).
[4]Dzh. L. Chkareuli, Pis'ma Zh. Eksp. Teor. Fiz. **32**, 684 (1980) [JETP Lett. **32**, 671 (1980)].
[5]Z. G. Berezhiani and Dzh. L. Chkareuli, Yad. Fiz. **37**, 1043 (1983) [Sov. J. Nucl. Phys. **37**, 618 (1983)].
[6]Z. G. Berezhiani, Phys. Lett. B **129**, 99 (1983); **150**, 177 (1985).
[7]H. Fritzsch, Nucl. Phys. B **155**, 198 (1979).
[8]R. D. Peccei and H. R. Quinn, Phys. Rev. D **16**, 1791 (1977).
[9]J. Preskill, Ann. Rev. Nucl. Part. Sci. **34**, 461 (1984).
[10]T. M. Bibilashvili and G. R. Dvali, Phys. Lett. B **248**, 269 (1990); G. R. Dvali, Phys. Lett. B (in press).
[11]A. Vilenkin, Phys. Rev. Lett. **53**, 1016 (1984).
[12]T. W. B. Kibble, Phys. Rev. D **33**, 328 (1986).

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