ACT-8
CERN-TH.5777/90
CTP-TAMU-52/90
UGVA-DPT 1990/06-674
June, 1990

CONFINEMENT OF FRACTIONAL CHARGES YIELDS INTEGER-CHARGED RELICS IN STRING MODELS

JOHN ELLIS

CERN Theory Division, 1211 Geneva 23, Switzerland

JORGE L. LOPEZ[(a)] and D. V. NANOPOULOS[(b)(c)]

Center for Theoretical Physics, Department of Physics

Texas A & M University, College Station, TX 77843-4242, USA

and

Astroparticle Physics Group, Houston Advanced Research Center (HARC)

The Woodlands, TX 77381, USA

ABSTRACT

We demonstrate that all fractionally charged particles in the revamped flipped $SU(5)$ model are confined by a hidden analogue of QCD. This result applies to states with mass $\sim M_{Pl}$ as well as to the lighter states discussed previously. We give sufficient conditions on other string models to confine light fractionally charged particles. The hidden sector of the revamped flipped $SU(5)$ model contains metastable integer-charged hadron-like states. Such "cryptons" are likely to be generic features of string models that confine unwanted fractional charges, and should be present as dark matter in the Universe.

(a) Supported by an ICSC–World Laboratory Scholarship.

(b) Theory Division, CERN, 1211 Geneva 23; and Départment de Physique Théorique, Université de Genève, 1211 Geneva 4, Switzerland.

(c) Supported in part by DOE grant DE-AS05-81ER40039, and by the Swiss National Foundation.

1. Introduction

There have been many experimental searches for fractionally charged particles, largely inspired by the fractional charge assignments of the quarks in the Standard Model. No reported observation of a fractionally charged particle has ever been confirmed, and there are upper bounds on the abundance of any such particles in the range of 10^{-19} to 10^{-26} [1] of the nucleon abundance for charges between 1/3 and 1. Comparison with the abundances calculated [2] in conventional Big Bang cosmology implies that there can be no fractionally charged particle state with mass below the reheating temperature when a large amount of entropy was last dumped into conventional matter particles. An upper bound on this reheating temperature is presumably set by the string Hagedorn temperature, though it might be somewhat lower in inflationary cosmology or if there were some late-decaying particle. Thus there are no light free fractionally charged particle states, although some might be allowed if their masses were $\gg m_W$.

The experimental lack of fractionally charged particles is a potential embarrassment to superstring models, since many of them are known to predict the existence of particles with charges $1/n$, where n is some model-dependent integer [3]. Indeed, Schellekens has recently argued [4] that their presence is a generic feature of models derived from the superstring with a level-one Kac-Moody current algebra. One could hope to avoid this phenomenological embarrassment by going to a higher-level Kac-Moody algebra. However, no viable higher-level model is known, and the construction of any such model is beset with constraints [5, 6] . Fortunately, there is another way to avoid detectable fractionally charged particles, namely to confine them analogously to quarks in QCD.

This confinement solution was in fact proposed [7] before the full gravity of the fractional charge problem became apparent, in the context of the revamped flipped $SU(5) \times U(1)$ model [7] . It was shown that all the light ($m \ll M_{Pl}$) fractionally charged hidden states had charges $|Q_{em}| = \frac{1}{2}$ and were in **4** or $\overline{\mathbf{4}}$ representations of a hidden $SU(4)$ gauge group which became strong at some large mass scale Λ_4. According to standard lore, one would expect all the **4** and $\overline{\mathbf{4}}$ states to be bound into "hidden hadrons"—called here "cryptons"—with charges $Q_{em} \in \mathbf{Z}$. However, in that paper [7] the existence of any massive ($m \sim M_{Pl}$) fractionally charged states was left open, and this possibility was indeed subsequently raised by Schellekens [4].

In this note we first prove that there are in fact no free fractionally charged states at any mass level of the revamped flipped $SU(5) \times U(1)$ model. We also give sufficient conditions under which other models that confine light fractionally charged states also confine their massive sisters. On the other hand, the hidden sector of the revamped flipped $SU(5) \times U(1)$ model contains integer-charged "cryptons", with global quantum numbers that are conserved by the renormalizable field-theoretical interactions. Hence, their lightest representatives are likely to be metastable with lifetimes much longer than the age of the Universe, and are therefore possible candidates for the dark matter which seems to abound in the Universe.

2. Absence of fractionally-charged states at all mass levels

Our solution [7] to the charge quantization problem can be encoded in the following experimentally motivated charge quantization dogma: All (massless) fractionally charged particles must have nontrivial quantum numbers under unbroken nonabelian gauge groups, such that when confinement sets in and thus only gauge singlets are observable, the resulting physical states are integrally charged. In a string-derived model this condition can be conveniently implemented in the language of simple currents [8] of the Kac-Moody algebra underlying the gauge group [4,6]. Compatibility with the string rules places severe restrictions on the gauge groups and their Kac-Moody levels, which could enforce such a quantization condition.

The revamped flipped $SU(5)$ model of [7] has gauge group $G = G_{obs} \times G_{hidden}$, with $G_{obs} = SU(5) \times U(1) \times U(1)^4$, $G_{hidden} = SO(6)_h \times SO(10)_h$, and all subgroups are realized at level one. It can be shown [6] that the following charge quantization rule can be consistently imposed on the string-derived spectrum of the model,

$$\alpha = \frac{1}{3}t_3 + Q + \frac{1}{2}t_4 + \frac{1}{2}c \in \mathbf{Z}, \tag{1}$$

where t_3, t_4, and c are the triality, quadrality, and conjugacy class of the respective $SU(3)_C$, $SO(6)_h \approx SU(4)_h$, and $SO(10)_h$ representations. From the spectrum in [7] one can readily verify that this condition is satisfied for all massless states in the model. In fact, this condition holds at *all* mass levels, as we now show.

The full spectrum of the model falls into representations of the Virasoro algebra. These are obtained by acting on the primary states with the Virasoro creation operators, thus generating the so-called Verma module [9] . The primary states allowed in a conformal field theory are restricted by unitarity. For the $SO(2n)$ Kac-Moody algebras only the singlet, vector, spinor, and conjugate spinor matter representations are unitary at level one. The spectrum of the revamped flipped $SU(5)$ model indeed contains the $\mathbf{1},\mathbf{4},\overline{\mathbf{4}}$, $\mathbf{6}$ representations of $SO(6)_h$, and $\mathbf{1}$ and $\mathbf{10}$ representations of $SO(10)_h$; the $SO(10)_h$ $\mathbf{16},\overline{\mathbf{16}}$ representations are expected to arise at the massive level. The Virasoro creation operators, represented by fermion oscillator modes, correspond to vector representations of these groups, since they have conformal dimension 1/2, and it is known that the conformal dimension of the vector representation of $SO(2n)$ at level one is 1/2 also.* We then get a product of vector representations times the original highest weight primary field representation. However, vector representations of these groups have $t_4 = c = 2$, and hence they do not change the quadrality or conjugacy class of the state $mod\, 2$. Eq. (1) then implies that if the highest weight primary fields satisfy this condition, so will all states in the corresponding Verma modules. Hence, there are no free fractionally charged particles at any mass level of the revamped flipped $SU(5)$ model.

* Indeed, the conformal dimension of the vector representation of an $SO(2n)$ Kac-Moody algebra is given by $h_v = C_v/(2k + C_A)$. With $C_v = 2n - 1, C_A = 4n - 4$, and $k = 1$ we get $h_v = 1/2$ [9].

In general, unified models which do not have the electric charge operator fully within their $SU(2)_L \times U(1)_Y$ subgroup, will have *a priori* fractional charged states in their spectrum. This is notably the case in more general flipped-type models [10] and in other models in which the weak hypercharge is a linear combination of various $U(1)$ subgroups [11]. In these models, if the massless spectrum (*i.e.*, a subset of all the primary fields) satisfies a charge quantization condition, we then expect the massive states to do so as well. This expectation seems reasonable since after all we expect the Virasoro generators to commute with the electric charge operator. However, in many of these models the first condition is not satisfied due to an inadequate hidden sector, and hence fractionally charged states are expected to occur at both the massless and massive levels.

In string-derived models the magnitude of the electric charge quantum can be related to the two-dimensional dynamics that originated the particular pattern of gauge symmetry breaking, *e.g.*, the Wilson line configuration [3]. Indeed, in orbifold models with discrete symmetry group $\mathbf{Z}_n$, the electric charges are multiples of $1/n$ [3]. The revamped flipped $SU(5)$ model can be interpreted as such a model with a $\mathbf{Z}_4$ discrete symmetry* in the fermionic degrees of freedom representing the $SU(5) \times U(1)$ gauge group. In this model an $SO(10)$ gauge symmetry is broken down to $SU(5) \times U(1)$ by a boundary condition vector of order 4, and the electric charges of the states are indeed seen to be multiples of $1/4$ [6]†.

3. Integer-charged "cryptons"

Models which do not have a hidden sector that is sticky enough to confine the existing fractionally charged states will be doomed as explained above. On the other hand, potentially realistic models that confine fractional charges will then generally contain integer-charge "hidden hadrons" as their "solution" to the

* For a connection between the free fermionic language and the orbifold language see [12].

† Here we are referring to states that exist in the $SU(5) \times U(1)$ model but *not* in the original $SO(10)$ $\mathbf{10}$ and $\mathbf{16},\overline{\mathbf{16}}$ representations. The observable sector states form $SO(10)$ multiplets and hence have their traditional charge assignments.

charge quantization problem. The phenomenology of these "cryptons" depends crucially on their lifetime. Should they be long-lived, then they could affect significantly the whole evolution of the Universe. Should they be charged also, then additional observational constraints could exist on their relic abundances. We now address these issues in the context of the revamped flipped $SU(5)$ model.

The spectrum of massless ($m \ll M_{Pl}$) fields in the revamped flipped $SU(5)$ model [7] is given in Table 1. From this spectrum one can see that there are three kinds of hidden $SO(6) \times SO(10)$ invariant bound states possible, as follows

$$\text{"cryptons"} \quad \begin{cases} \text{"hidden mesons"}: & T_i T_j,\ \Delta_i \Delta_j,\ \tilde{F}_i \tilde{\bar{F}}_j; & (0, \pm 1), \\ \text{"hidden baryons"}: & \tilde{F}_i \tilde{F}_j \Delta_k,\ \tilde{\bar{F}}_i \tilde{\bar{F}}_j \Delta_k; & (0, \pm 1), \\ \text{"hidden tetrons"}: & \tilde{F}_i \tilde{F}_j \tilde{F}_k \tilde{F}_l,\ \tilde{\bar{F}}_i \tilde{\bar{F}}_j \tilde{\bar{F}}_k \tilde{\bar{F}}_l; & (0, \pm 1, \pm 2), \end{cases} \tag{2}$$

where we have indicated in parentheses the possible electromagnetic charges for each class of crypton.

The trilinear superpotential involving hidden and hidden/observable couplings can be straightforwardly calculated following the rules given in [12]. We obtain*

$$\begin{aligned} W = g\sqrt{2}\{ & \Delta_1^2 \bar{\Phi}_{23} + \Delta_2^2 \Phi_{31} - \Delta_4^2 \bar{\Phi}_{23} + \Delta_5^2 \bar{\Phi}_{31} + \frac{1}{\sqrt{2}} \Delta_4 \Delta_5 \bar{\phi}_3 \\ & + T_1^2 \bar{\Phi}_{23} + T_2^2 \Phi_{31} + T_4^2 \Phi_{23} + T_5^2 \Phi_{31} + \frac{1}{\sqrt{2}} T_4 T_5 \phi_2 \\ & + \frac{1}{\sqrt{2}} (\tilde{F}_1 \tilde{\bar{F}}_2 \phi_4 + \tilde{F}_2 \tilde{\bar{F}}_1 \phi_1) + \tilde{F}_2 \tilde{\bar{F}}_2 \phi^- - \frac{1}{2} \tilde{F}_3 \tilde{\bar{F}}_3 \Phi_3 \\ & + \tilde{F}_5 \tilde{\bar{F}}_5 \bar{\Phi}_{12} + \tilde{F}_3 \tilde{F}_6 \Delta_1 + \tilde{F}_3 \tilde{\bar{F}}_4 l_2^c \}. \end{aligned} \tag{3}$$

The gauge singlet fields appearing in W generally acquire vacuum expectation values (v.e.v.'s) of $\mathcal{O}(\xi)$, $\xi \equiv (\mathrm{Tr} U_A(1) g^2 \sqrt{2\alpha'}/192\pi^2)^{1/2}$ [7,14], thus effecting the breaking $SU(5) \times U(1) \times U(1)^4 \to SU(5) \times U(1)$ near the string scale. This

---

* This result is in basic agreement with a similar calculation in [13], where some $\mathcal{O}(1)$ factors were not kept.

is in order to preserve D- and F-flatness in the presence of a Fayet-Iliopoulos D-term for the anomalous $U_A(1)$ in the model. Specifically, we find that, *e.g.*, $\langle \mathbf{\Phi}_{23}, \mathbf{\Phi}_{31}, \bar{\mathbf{\Phi}}_{23}, \bar{\mathbf{\Phi}}_{31}, \phi_{45}, \bar{\phi}_{45}, \phi^+, \bar{\phi}^+ \rangle \neq 0$, whilst $\langle \mathbf{\Phi}_3, \mathbf{\Phi}_{12}, \bar{\mathbf{\Phi}}_{12} \rangle = 0$. This pattern of v.e.v.'s implies that $\Delta_{1,2,4,5}$ and $T_{1,2,4,5}$ acquire large masses at lowest-order in superpotential interactions, as can readily be seen from W. Mass terms of the form $\Delta_i \Delta_j \langle \phi^n \rangle$, $T_i T_j \langle \phi^n \rangle$: i and/or $j = 3$, where ϕ^n is a product of gauge singlets in the model, could in principle arise at the nonrenormalizable level. However, conservation of global $U(1)$ symmetries [12] in the model, together with a peculiar global $U(1)$ charge assignment for Δ_3 and T_3, make these terms vanish identically for all n. One can consider more general forms for $\langle \phi^n \rangle$, *e.g.*, hidden sector condensates $\left\langle T_i T_j, \Delta_i \Delta_j, \tilde{F}_i \tilde{\tilde{F}}_j, \ldots \right\rangle$, and/or v.e.v.'s arising after $SU(5) \times U(1)$ symmetry breaking. However, all these alternatives can be shown to give no Δ_3 or T_3 masses $\gtrsim \Lambda_4$.

The $\tilde{F}_i \tilde{\tilde{F}}_j$ mass matrix can be analyzed similarly. In this case only linear combinations of $\tilde{F}_1, \tilde{\tilde{F}}_1, \tilde{F}_2, \tilde{\tilde{F}}_2$ acquire $\mathcal{O}(\xi)$ masses. At the level of quartic non-renormalizable interactions, one finds generic mass terms of the forms $\tilde{F}_4 \tilde{\tilde{F}}_4 \langle \phi^2 \rangle$, $\tilde{F}_6 \tilde{\tilde{F}}_6 \langle \phi^2 \rangle$ which give $\mathcal{O}(\xi^2)$ masses to these fields. One can show that mass terms of the form $\tilde{F}_i \tilde{\tilde{F}}_j \langle \phi^n \rangle$, i and/or $j \in \{3, 5\}$, vanish identically for all n, due to global $U(1)$ obstructions as above. Also, more general forms for $\langle \phi^n \rangle$ do not yield any mass terms $\gtrsim \Lambda_4$.

Since a renormalization group analysis [13] suggests that $\Lambda_4 \simeq 10^{12}\,\mathrm{GeV} \ll \Lambda_{10} \simeq 10^{15}\,\mathrm{GeV}$, we expect $SU(4)$ bound states to be much lighter than $SO(10)$ bound states in general. The lightest $SU(4)$ bound states are in turn expected to be those made out of the $\mathbf{4}/\bar{\mathbf{4}}$ fields $\tilde{F}_3, \tilde{F}_5$, $\tilde{\tilde{F}}_3$, and $\tilde{\tilde{F}}_5$, which have an approximate $SU(2)_F$ flavour symmetry as in QCD: among these states are two-constituent mesons $\tilde{F}_{3,5} \tilde{\tilde{F}}_{3,5}$ and four-constituent "tetrons" $\tilde{F}^4_{3,5}$ and $\widetilde{\tilde{F}^4_{3,5}}$. Because of its larger crypto-colour charge, one would expect bound states of the $\mathbf{6}$ field Δ_3, namely the two-constituent meson Δ_3^2 and the three-constituent baryons $\tilde{F}_{3,5} \tilde{F}_{3,5} \Delta_3$ and $\tilde{\tilde{F}}_{3,5} \tilde{\tilde{F}}_{3,5} \Delta_3$ to be somewhat heavier. We will not consider further $SU(4)$ bound states made out of massive constituents, nor $SO(10)$ bound states.

We expect the lightest hidden $SU(4)$ meson to be analogous to the π^0: $\pi_4^0 \simeq \frac{1}{\sqrt{2}}(\widetilde{F}_3\widetilde{\bar{F}}_3 - \widetilde{F}_5\widetilde{\bar{F}}_5)$, with charged $\pi_4^{\pm} \simeq (\widetilde{F}_3\widetilde{\bar{F}}_5, \widetilde{\bar{F}}_3\widetilde{F}_5)$ states slightly heavier because of electromagnetic mass splitting analogous to that in QCD:

$$m^2_{\pi_4^0} \simeq \Lambda_4 \times m_{\widetilde{F}_{3,5}/\widetilde{\bar{F}}_{3,5}}, \tag{4}$$

$$m^2_{\pi_4^{\pm}} - m^2_{\pi_4^0} \simeq (\frac{\alpha_{em}}{\pi})\,\Lambda_4^2\,\ln(\Lambda_4^2/m^2_{\pi_4^0}). \tag{5}$$

As in QCD, we expect the $\eta_4^0 \simeq \frac{1}{\sqrt{2}}(\widetilde{F}_3\widetilde{\bar{F}}_3 + \widetilde{F}_5\widetilde{\bar{F}}_5)$ state to be significantly heavier, because of a $U_A(1)$ anomaly. By analogy with the nucleon and Δ states in QCD, the lightest tetrons are expected to be the neutral $\widetilde{F}_3^2\widetilde{F}_5^2$ and $\widetilde{\bar{F}_3^2}\widetilde{\bar{F}_5^2}$ states, with singly-charged $\widetilde{F}_3^3\widetilde{F}_5$, $\widetilde{F}_3\widetilde{F}_5^3$, $\widetilde{\bar{F}_3^3}\widetilde{\bar{F}}_5$, and $\widetilde{\bar{F}}_3\widetilde{\bar{F}_5^3}$ states somewhat heavier, and doubly-charged $\widetilde{F}_3^4$, $\widetilde{F}_5^4$, $\widetilde{\bar{F}_3^4}$, and $\widetilde{\bar{F}_5^4}$ states even heavier.

4. Crypton decays

A complete discussion of the probable decay modes and likely lifetimes of all these hidden sector bound states would be very model-dependent and take us beyond the scope of this paper. However, we think it interesting to point out that although some of these bound states are expected to be very short-lived, some may have very long lifetimes and could be present as dark matter in the Universe [15] . As an example of a very short-lived state, consider the π_4^0, which can decay via the renormalizable trilinear $\widetilde{F}_3\widetilde{\bar{F}}_3\Phi_3$ and/or $\widetilde{F}_5\widetilde{\bar{F}}_5\bar{\Phi}_{12}$ interactions in the superpotential (3). To discuss longer-lived states one must look for non-renormalizable higher-order interactions using the rules of Ref. [12].

In the case of the mesonic cryptons, we found nonzero decay interactions via N^{th} order superpotential terms with the following matrices of values of N corresponding to the different combinations of mesonic constituents (i^+ means

$N \geq i$)

$$T_iT_j \quad (i,j=1,\ldots,5) \qquad \mathbf{N} = \begin{pmatrix} 3 & 6^+ & 6^+ & 5 & 6^+ \\ 6^+ & 3 & 6^+ & 5 & 6^+ \\ 6^+ & 6^+ & 6^+ & 6^+ & 6^+ \\ 5 & 5 & 6^+ & 3 & 3 \\ 6^+ & 6^+ & 6^+ & 3 & 3 \end{pmatrix}, \tag{6}$$

$$\Delta_i\Delta_j \quad (i,j=1,\ldots,5) \qquad \mathbf{N} = \begin{pmatrix} 3 & 6^+ & 6^+ & 6^+ & 4 \\ 6^+ & 3 & 6^+ & 6^+ & 5 \\ 6^+ & 6^+ & 6^+ & 6^+ & 6^+ \\ 6^+ & 6^+ & 6^+ & 3 & 3 \\ 4 & 5 & 6^+ & 3 & 3 \end{pmatrix}, \tag{7}$$

$$\widetilde{F}_i\widetilde{\overline{F}}_j \quad (i,j=1,\ldots,6) \qquad \mathbf{N} = \begin{pmatrix} 5 & 3 & 6^+ & 4 & 6^+ & 6^+ \\ 3 & 3 & 6^+ & 5 & 6^+ & 6^+ \\ 5 & 5 & 3 & 3 & 6^+ & 6^+ \\ 5 & 5 & 6^- & 4 & 6^+ & 6^+ \\ 6^+ & 6^+ & 6^- & 6^+ & 3 & 6^+ \\ 4 & 5 & 6^- & 6^+ & 6^+ & 4 \end{pmatrix}. \tag{8}$$

In the case of baryons, we find the $N=3$ term $\widetilde{F}_3\widetilde{F}_6\Delta_1$, no $N=4$ terms, and the following $N=5$, $\widetilde{F}_i\widetilde{F}_j\Delta_k$ and $\widetilde{\overline{F}}_i\widetilde{\overline{F}}_j\Delta_k$ terms, denoted by ijk and $\bar{i}\bar{j}k$ respectively: $141, 145, 162, 223, 231, 232, 235, 342, 344, 364, 443, 464, 663$; $\bar{1}\bar{4}5, \bar{2}\bar{2}3, \bar{4}\bar{4}3, \bar{6}\bar{6}3$. In the case of tetrons, we find no $N=4$, $\widetilde{F}_i\widetilde{F}_j\widetilde{F}_k\widetilde{F}_l$ or $\widetilde{\overline{F}}_i\widetilde{\overline{F}}_j\widetilde{\overline{F}}_k\widetilde{\overline{F}}_l$ interactions, and the only nonzero interactions for $N=5$ are $\widetilde{F}_2^2\widetilde{F}_3^2\Phi_{23}$ and $\widetilde{\overline{F}}_1^2\widetilde{\overline{F}}_3^2\Phi_{31}$. As for the lightest $\widetilde{F}_{3,5}$ and $\widetilde{\overline{F}}_{3,5}$ tetrons, we find no $\widetilde{F}_{3,5}^4\mathcal{O}_i^n$, $(n \leq 3)$ interactions, where the $\mathcal{O}_i$ are generic light observable fields. Therefore, the decays of the lightest tetrons could only arise via N^{th} order interactions (where $N \equiv n+4$) with $N \geq 8$, or possibly via non-perturbative interactions.

For general N^{th} order interactions, we expect strengths of order g^{N-2}/m_K^{N-3}, where g is the gauge coupling and $m_K \simeq 10^{17}$ GeV. This gives a decay rate $\Gamma_\chi \simeq$

$m_\chi(m_\chi/m_K)^{2(N-3)}$ for a generic state χ of mass $m_\chi \simeq \Lambda_4$. When $4 \leq N < 8$, we expect such a state to decay when $\Gamma_\chi \simeq H$, the Hubble expansion rate calculated assuming that non-relativistic χ states dominate the energy density. The decays therefore occur at a characteristic temperature

$$T_\chi \simeq m_\chi \left(\frac{m_\chi}{m_K}\right)^{4(N-3)/3} \left(\frac{M_{Pl}}{m_\chi}\right)^{2/3}, \tag{9}$$

and reheat the Universe to a temperature

$$T_{R_\chi} \simeq m_\chi \left(\frac{m_\chi}{m_K}\right)^{N-3} \left(\frac{M_{Pl}}{m_\chi}\right)^{1/2}, \tag{10}$$

generating an increase in the entropy of the Universe by a factor

$$X_\chi = \left(\frac{T_R}{T_\chi}\right)^3 \simeq \left(\frac{m_K}{m_\chi}\right)^{N-3} \left(\frac{m_\chi}{M_{Pl}}\right)^{1/2}. \tag{11}$$

When $N \geq 8$, we expect the χ lifetime to exceed $10^6 y$, in which case at least some χ states could be around today as cosmological relics from the Big Bang.

Combining the analyses of the two previous paragraphs, we see that the lightest neutral tetrons are possible candidates to be the dark matter of the Universe. Although part of our discussion has been specific to the flipped $SU(5)$ model, we believe that this conclusion is generic. String models that avoid fractionally charged relic particles must confine them, and this means that they should expect integrally charged metastable states. These are neutral in the flipped $SU(5)$ case, but could have non-zero charges in other string models.

5. Cosmological constraints

Finally we turn to the cosmological constraints on, and implications of, metastable cryptons like those found above in the flipped $SU(5)$ model. Although other constraints can be considered [16] , we concentrate on the following key ones: big bang baryogenesis (BBB), big bang nucleosynthesis (BBN), and the present density of cryptons in the Universe relative to the critical value (Ω_c). Although BBB has normally been discussed in the context of GUTs at a temperature $T_{BBB} \sim 10^{15}$ GeV there exist at least two possible scenarios for BBB at a temperature of order 100 to 1000 GeV, namely the decays of coherent squark and/or slepton fields [17] , and non-perturbative electroweak effects [18]. Moreover, crypton decays themselves could be the mechanism for BBB, in which case BBB could occur at any time before BBN. Accordingly, we will not impose any constraints on cryptons or other particles that, when they decay, reheat the Universe to a reheating temperature $T_R > (T_{BBB})_{min} \simeq 100$ GeV, and even consider the possibility of BBB by crypton decay at some T_R as low as the characteristic temperature for BBN, namely $T_{BBN} \sim 0.1$ to 1 MeV. The agreement between BBN calculations and the observed light element abundances corresponds to a baryon-to-entropy ratio $\eta \sim 3 \times 10^{-10}$ at T_{BBN} [19] . Accordingly, unless the cryptons are responsible for BBB, the maximum possible increase in the entropy due to the decays of cryptons or other particles giving reheating temperatures $T_R \in ((T_{BBN})_{min} \simeq 0.1\,\mathrm{MeV}, (T_{BBB})_{min} \simeq 100\,\mathrm{GeV})$ is about 10^9. Subsequent to this, an upper limit on the maximal entropy increase is provided by the agreement between η as measured by astrophysicists and as estimated on the basis of BBN, namely a factor $\lesssim 10$. As for the present density of the Universe, we assume for reasons of naturalness and inflation that the total $\Omega = \Omega_c + \Omega_B + \ldots = 1$, and hence that the present crypton density $\Omega_c \lesssim 1$.

This last constraint can be compared with the naive estimate of Ω_c which would be obtained by taking over technicolour-inspired calculations [20] of the relic density Ω_T of particles with mass m that annihilate strongly via a cross-section $\sigma \simeq 1/m^2$:

$$\Omega_T \simeq \left(\frac{m}{100\,\mathrm{TeV}}\right)^2. \tag{12}$$

Since our metastable cryptons are likely to weigh $\sim 10^{12}$ GeV, as discussed above, we would find $\Omega_T \simeq 10^{14}$! However, the naive technicolour-inspired estimate (12) does not take into account possible entropy generation subsequent to annihilation.

We interpret (12) as telling us that a factor $\gtrsim 10^{14}$ of entropy should be generated by one or more mechanisms after crypton annihilation. As discussed above, a factor $\simeq 10$ could have been generated after BBN, a factor $\gtrsim 10^{9}$ between BBN and BBB, and there is no constraint on entropy release before BBB.

There are many possible mechanisms for entropy release in a string model with metastable hidden sector particles, such as inflation, thermalization of coherent oscillations of some other scalar fields [21] , gravitino decay [22] , and the decays of some other more unstable hidden sector particles. We now discuss each one of these in turn, with specific comments on their realizability in the flipped $SU(5)$ model.

No general mechanism for *inflation* in superstring-inspired models has been found [23] , but this does not mean that one does not exist, at least in some specific models. Inflation with a reheating temperature between Λ_4 and T_{BBB} would certainly solve the Ω problem (12).

Next we consider entropy generation by the *thermalization of coherent oscillations* of some non-inflaton scalar field, such as the one (ν^c_H) that breaks $SU(5) \times U(1) - SU(3) \times SU(2) \times U(1)$ in the flipped $SU(5)$ model [21]. It is well known that thermal effects could trap such a field in the symmetric phase until it eventually escapes and produces an enormous entropy increase [24] . In the specific flipped $SU(5)$ model, strong non-perturbative $SU(5)$ effects are likely to eject ν^c_H from the symmetric phase, generating an entropy increase factor [21]

$$X_{\nu^c_H} \simeq M^3_{GUT} \widetilde{m}^{1/2} / (M^{1/2}_{Pl} T^3_i), \tag{13}$$

where $\widetilde{m}$ is the scale of supersymmetry breaking in the observable sector and $T_i \lesssim \Lambda_4$ is the temperature of the Universe when ν^c_H oscillations begin. The corresponding reheating temperature is $T_R \simeq \widetilde{m}^{3/2} M^{1/2}_{Pl} / M_{GUT}$. Taking $M_{GUT} \simeq 10^{16}\,\mathrm{GeV}$, $\widetilde{m} \simeq m_W \simeq 10^2\,\mathrm{GeV}$, and $T_i \sim 10^{12}\,\mathrm{GeV}$, we find $X_{\nu^c_H} \simeq 10^{3.5}$ and $T_R \simeq 0.3\,\mathrm{MeV}$, which would need to be combined with some other other mechanism for entropy generation.

In models where supersymmetry breaking originates in a hidden sector, so that $m_{3/2} \gg m_W$, *gravitino decay* is a well-known possible source of entropy [22] . Assuming that non-relativistic gravitinos dominate the Universe's energy density until they decay, while the density of relativistic particles decreases $\propto T^4$, we find that gravitinos decay at a characteristic temperature $T_G \simeq m_{3/2}^{5/3}/M_{Pl}^{2/3}$. They then reheat the Universe to a temperature $T_{R_G} \simeq m_{3/2}^{3/2}/M_{Pl}^{1/2}$, and thereby generate an entropy increase factor

$$X_{3/2} \simeq \left(\frac{M_{Pl}}{m_{3/2}}\right)^{1/2} \simeq \left(\frac{M_{Pl}}{T_{R_G}}\right)^{1/3}. \tag{14}$$

For the range $(T_{BBN})_{min} \simeq 0.1\,\mathrm{MeV} \lesssim T_{R_G} \lesssim 100\,\mathrm{GeV} \simeq (T_{BBB})_{min}$, or equivalently $5 \times 10^3\,\mathrm{GeV} \lesssim m_{3/2} \lesssim 5 \times 10^7\,\mathrm{GeV}$, we get $5 \times 10^7 \gtrsim X_{3/2} \gtrsim 5 \times 10^5$. This entropy release is within our above limits, although not by itself enough to reduce Ω_c below 1. However, gravitino decay could easily be combined with one of the other entropy generation mechanisms discussed here. It should be noted that $m_{3/2} \gtrsim 10^3\,\mathrm{GeV}$ is compatible with the uncertainties in $m_{3/2}$ in the flipped $SU(5)$ model.

As a final possible source of entropy, we discuss the *decays of some more unstable hidden sector particle* χ. As discussed above, if χ decays through an N^{th} ($4 \leq N < 8$) order nonrenormalizable interaction, then it will generate an increase in the entropy of the Universe by a factor of X_χ. In the flipped $SU(5)$ model, with $m_\chi \simeq 10^{12}\,\mathrm{GeV}$, equations (9)–(11) become

$$T_\chi \simeq 10^{10(11-2N)/3}\,\mathrm{GeV}, \quad T_{R_\chi} \simeq 10^{5(6.1-N)}\,\mathrm{GeV}, \quad X_\chi \simeq 10^{5(N-3.62)}. \tag{15}$$

For $N = 4,5$ we obtain $T_{R_\chi} > 100\,\mathrm{GeV}$ and $X_\chi \lesssim 10^7$, which could be useful when combined with one of the other mechanisms described above. For $N = 6$, we get $T_{R_\chi} \sim 3\,\mathrm{GeV}$ and $X_\chi \sim 10^{12}$, which could be acceptable, within the uncertainties, if BBB occurred in the decays of these cryptons. The case $N = 7$ seems to be unacceptable.

The above brief discussion indicates that suppressing the relic crypton density so that $\Omega_c < 1$ is not impossible, and we defer a detailed discussion to a future publication. However, there is one point we would like to mention: in view of the likelihood and necessity of generating a large entropy increase after crypton-anticrypton annihilation, there is no natural reason why $\Omega_c \sim 0.1$ to 1 should be realized.

6. Conclusions

The existence of fractionally charged particles has been emphasized as a distinctive, even unavoidable feature of many string models. We have shown previously how this conclusion can be evaded by confining the unwanted fractionally charged particles. In this paper we have shown how this confinement mechanism may also apply to massive string states. We have also emphasized that this confinement leads to the existence of metastable integer-charged states which could be present as dark matter in the Universe. This possibility is severely constrained by cosmology and astrophysics, but could serve as the long-sought "smoking gun" of string.

REFERENCES

1. Particle Data Group, Phys. Lett. B **204** (1988) .
2. S. Wolfram, Phys. Lett. B **82** (1979) 65.
3. X. G. Wen and E. Witten, Nucl. Phys. B **261** (1985) 651; G. Athanasiu, J. Atick, M. Dine, and W. Fischler, Phys. Lett. B **214** (1988) 55; P. Huet, City College of CUNY preprint CCNY-HEP-90/9 (1990).
4. A. Schellekens, Phys. Lett. B **237** (1990) 363.
5. A. Font, L. Ibanez, and F. Quevedo, CERN preprint TH.5666/90 (1990).
6. J. Ellis, J. L. Lopez, and D. V. Nanopoulos, Texas A & M University preprint CTP-TAMU-35/90 (1990).
7. I. Antoniadis, J. Ellis, J. Hagelin, and D. V. Nanopoulos, Phys. Lett. B **231** (1989) 65.
8. A.N. Schellekens and S. Yankielowicz, Nucl. Phys. B **327** (1989) 673; CERN preprint TH.5662/90 (1990).
9. For a review see *e.g.*, P. Goddard and D. Olive, Int. J. Mod. Phys. A **1** (1986) 303.
10. See, *e.g.*, S. Barr, Phys. Rev. D **40** (1989) 2457, and references therein.
11. See, *e.g.*, J. Kim, Phys. Lett. B **207** (1988) 434; A. Faraggi, D. V. Nanopoulos, and K. Yuan, Nucl. Phys. B **335** (1990) 347.
12. S. Kalara, J. L. Lopez, and D. V. Nanopoulos, Texas A & M University preprint CTP-TAMU-34/90 (1990).
13. G. Leontaris, J. Rizos, and K. Tamvakis, University of Ioannina preprint IOA-237/90 (1990).
14. J. Lopez and D. V. Nanopoulos, Texas A & M University preprint CTP-TAMU-60/89 (1989), Nucl. Phys. B (in press).
15. E.W. Kolb, D. Seckel, and M. S. Turner, Nature **314** (1985) 415.

16. P.H. Frampton and S.L. Glashow, Phys. Rev. Lett. **44** (1980) 1481; J. Ellis, T.K. Gaisser and G. Steigman, Nucl. Phys. B **117** (1981) 427.

17. I. Affleck and M. Dine, Nucl. Phys. B **249** (1985) 361; J. Ellis, D. V. Nanopoulos, and K.A. Olive, Phys. Lett. B **184** (1987) 37.

18. M.E. Shaposhnikov, Nucl. Phys. B **287** (1987) 757, Nucl. Phys. B **299** (1988) 797.

19. K.A. Olive, D.N. Schramm, G. Steigman, and T. Walker, Phys. Lett. B **236** (1990) 454.

20. R.S. Chivukula and T.P. Walker, Nucl. Phys. B **329** (1990) 445.

21. B.A. Campbell, J. Ellis, J.S. Hagelin, D.V. Nanopoulos, and K.A. Olive, Phys. Lett. B **197** (1987) 355.

22. J. Ellis, J.E. Kim, and D.V. Nanopoulos, Phys. Lett. B **145** (1984) 181; S. Weinberg, Phys. Rev. Lett. **48** (1982) 1303; D.V. Nanopoulos, K.A. Olive, and M. Srednicki, Phys. Lett. B **127** (1983) 30.

23. J. Ellis, K. Enqvist, D.V. Nanopoulos, and M. Quiros, Nucl. Phys. B **277** (1986) 231.

24. K. Yamamoto, Phys. Lett. B **168** (1986) 341; K. Enqvist, D.V. Nanopoulos, and M. Quiros, Phys. Lett. B **169** (1986) 343.

Table 1: The spectrum of massless hidden fields in the revamped flipped $SU(5)$ model. The superscripts indicate the $U(1)_{\widetilde{Y}}$ charge; the electric charge is given by $Q = \frac{2}{5}\widetilde{Y}$.

State	$SO(6)_h \times SO(10)_h$	$U_i(1)$				State	$SO(6)_h \times SO(10)_h$	$U_i(1)$			
Δ_1	$(\mathbf{6},\mathbf{1})^0$	0	$-\frac{1}{2}$	$\frac{1}{2}$	0	T_1	$(\mathbf{1},\mathbf{10})^0$	0	$-\frac{1}{2}$	$\frac{1}{2}$	0
Δ_2	$(\mathbf{6},\mathbf{1})^0$	$-\frac{1}{2}$	0	$\frac{1}{2}$	0	T_2	$(\mathbf{1},\mathbf{10})^0$	$-\frac{1}{2}$	0	$\frac{1}{2}$	0
Δ_3	$(\mathbf{6},\mathbf{1})^0$	$-\frac{1}{2}$	$-\frac{1}{2}$	0	$\frac{1}{2}$	T_3	$(\mathbf{1},\mathbf{10})^0$	$-\frac{1}{2}$	$-\frac{1}{2}$	0	$-\frac{1}{2}$
Δ_4	$(\mathbf{6},\mathbf{1})^0$	0	$-\frac{1}{2}$	$\frac{1}{2}$	0	T_4	$(\mathbf{1},\mathbf{10})^0$	0	$\frac{1}{2}$	$-\frac{1}{2}$	0
Δ_5	$(\mathbf{6},\mathbf{1})^0$	$\frac{1}{2}$	0	$-\frac{1}{2}$	0	T_5	$(\mathbf{1},\mathbf{10})^0$	$-\frac{1}{2}$	0	$\frac{1}{2}$	0
$\widetilde{F}_1$	$(\mathbf{4},\mathbf{1})^{+5/4}$	$-\frac{1}{4}$	$\frac{1}{4}$	$-\frac{1}{4}$	$\frac{1}{2}$	$\bar{\widetilde{F}}_1$	$(\bar{\mathbf{4}},\mathbf{1})^{-5/4}$	$-\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{2}$
$\widetilde{F}_2$	$(\mathbf{4},\mathbf{1})^{+5/4}$	$-\frac{1}{4}$	$\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{2}$	$\bar{\widetilde{F}}_2$	$(\bar{\mathbf{4}},\mathbf{1})^{-5/4}$	$-\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{4}$	$-\frac{1}{2}$
$\widetilde{F}_3$	$(\mathbf{4},\mathbf{1})^{-5/4}$	$\frac{1}{4}$	$\frac{1}{4}$	$-\frac{1}{4}$	$\frac{1}{2}$	$\bar{\widetilde{F}}_3$	$(\bar{\mathbf{4}},\mathbf{1})^{+5/4}$	$-\frac{1}{4}$	$-\frac{1}{4}$	$\frac{1}{4}$	$-\frac{1}{2}$
$\widetilde{F}_4$	$(\mathbf{4},\mathbf{1})^{+5/4}$	$\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{2}$	$\bar{\widetilde{F}}_4$	$(\bar{\mathbf{4}},\mathbf{1})^{-5/4}$	$-\frac{1}{4}$	$\frac{1}{4}$	$\frac{1}{4}$	$-\frac{1}{2}$
$\widetilde{F}_5$	$(\mathbf{4},\mathbf{1})^{+5/4}$	$-\frac{1}{4}$	$\frac{3}{4}$	$\frac{1}{4}$	0	$\bar{\widetilde{F}}_5$	$(\bar{\mathbf{4}},\mathbf{1})^{-5/4}$	$-\frac{3}{4}$	$\frac{1}{4}$	$-\frac{1}{4}$	0
$\widetilde{F}_6$	$(\mathbf{4},\mathbf{1})^{+5/4}$	$-\frac{1}{4}$	$\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{2}$	$\bar{\widetilde{F}}_6$	$(\bar{\mathbf{4}},\mathbf{1})^{-5/4}$	$\frac{1}{4}$	$-\frac{1}{4}$	$\frac{1}{4}$	$-\frac{1}{2}$