Astrophysical Production of Fractional Charge in Broken Quantum Chromodynamics

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The production of states of fractional charge in the early universe and by cosmic rays is calculated when quantum chromodynamics is broken at a scale less than that of the quantum chromodynamics scale parameter Λ. It is shown that it is possible to produce a detectable abundance of fractional charge states by both mechanisms.

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Since Gell-Mann[1] and Zweig[2] first postulated the existence of quarks with fractional electric charge ($\pm\frac{1}{3}e$ + integer) as the fundamental constituents of hadrons there have been many searches for fractionally charged states (FCS). Although all but one of the searches have reported null results,[3] potentially severe chemical and electrostatic effects make it impossible to interpret the null results as proof that FCS do not exist.[4] La Rue and co-workers[5] continue to find evidence for the existence of FCS; however, translation of their result into a universal abundance of FCS is not possible until the chemical effects are understood. In the absence of any chemical enhancement, their results suggest a FCS-to-nucleon ratio of $\gtrsim 10^{-20}$.

We consider production of FCS in the early universe and by high-energy cosmic rays in the solar system. We find that for reasonable values of the parameters, the abundance of FCS from the early universe can be as large as 10^{10} per nucleon, or can be much smaller than 10^{-50} per nucleon. We also find that cosmic-ray production can be an important terrestrial source of FCS, and should lead to an even larger FCS abundance on the moon.

We will assume that unbroken quantum chromodynamics (QCD) is perfectly confining and that the existence of FCS signals a breaking of QCD. (Other possibilities exist, e.g., fractionally charged leptons.) It is attractive to assume that at short distances QCD is exact, has the usual triality–electric-charge assignments, and that FCS exist because $SU(3)_C$ is spontaneously broken at some energy scale μ into a theory that allows the existence of free FCS.

If QCD is broken at a scale μ which is less than the QCD scale parameter Λ,[6] then in the early universe at a temperature $T \simeq \Lambda$ ($T > \mu$) when quarks condensed into hadrons, QCD would have been exact. Since we are assuming that QCD confines, only integrally charged, color singlets would be produced in the quark-hadron transition. Therefore, any FCS must be produced at temperatures $T \lesssim \mu$. In previous work, it has been assumed that FCS is produced during the quark-hadron transition.[7]

The first model of broken QCD was proposed by De Rújula, Giles, and Jaffe[8] (DGJ), in which the local $SU(3)_C$ is broken to a global SU(3). Recently Slansky, Goldman, and Shaw[9] (SGS) have proposed breaking $SU(3)_C$ to $SO(3)_g$, where $SO(3)_g$ is the maximal subgroup of $SU(3)_C$. The results we obtain apply to both models, although for convenience we specialize to the SGS model for our calculations.

In the $SU(3)_C \rightarrow SO(3)_g$ breaking scheme proposed by SGS, the symmetry is spontaneously broken by a $\underline{27}_C$ of Higgs bosons; three gluons are left massless, while the other five gluons obtain a mass μ. At short distances ($r \ll 1/\mu$) the effective theory is QCD. Quarks and gluons are confined, since they all carry $SO(3)_g$ charge, and it is assumed that $SO(3)_g$ confines. The $SO(3)_g$ triquark and quark-antiquark singlets are also $SU(3)_C$ singlets and hence have integer charge. Thus the breaking $SU(3)_C \rightarrow SO(3)_g$ does not upset the usual hadron spectroscopy.[9]

The simplest $SO(3)_g$ singlets with fractional charge are quark-gluon bound states and the diquark bound states, d. SGS argue that the diquark states should have the lowest mass, and that their mass may not be too different from typical meson masses. Although we shall discuss production in this context, neither assumption affects our results since we find that production is independent of the mass of the FCS and whether it is a diquark or quark-gluon bound state. Accelerator production of the low-mass diquarks is energetically possible, but the production process is suppressed. Since the same suppression mechanisms also occur in the context of astrophysical production, we will now discuss diquark production in detail.

Consider the production of a diquark-antidiquark ($d\bar{d}$) pair in e^+e^- annihilation. Production proceeds via virtual photon production of a quark-antiquark ($q\bar{q}$) pair, followed by separation of the $q\bar{q}$ to a distance $r \gtrsim \mu^{-1}$ (beyond the range of the massive gluons) where they may be dressed to become $SO(3)_g$ singlet [but $SU(3)_C$ nonsinglet] states. At short distances ($\Lambda^{-1} < r < \mu^{-1}$) QCD confining forces oppose the separation of the $q\bar{q}$ pair. When $r \simeq \mu^{-1}$ these forces decrease and fall to zero [for $SO(3)_g$ singlet states] because of the gluon's finite mass. We shall assume that for $\Lambda^{-1} \leq r \leq \mu^{-1}$ the potential between the $q\bar{q}$ pair is linear: $V_{q\bar{q}} \simeq kr$, where r is the separation between q and $\bar{q}$ and k is the string constant, $k \simeq 1$ GeV fm^{-1}.[10] At a separation $r = \mu^{-1}$, the potential energy in the string is large; $V^{\max} \simeq (200\text{ GeV})\,\mu_{\text{MeV}}^{-1}$ where μ_{MeV} is μ in units of megaelectronvolts. Therefore the quarks must either tunnel through the barrier to $r \geq \mu^{-1}$ or have energy greater than $V^{\max}$. Once the $q\bar{q}$ pair have separated (either through tunneling or by being energetic enough to go over the barrier) they can be dressed as a diquark-antidiquark ($d\bar{d}$) state of mass m_d in the SGS model, or propagate undressed as a free quark and antiquark of mass $m_q \simeq V^{\max}/2$ in the DGJ model.

The probability that the $q\bar{q}$ pair tunnels through the linear QCD potential barrier to a distance $\simeq \mu^{-1}$ (where they can become diquarks in the SGS model) may be estimated for initial c.m. energies much less than the barrier height as

$$p_T = \exp(-2S_{\text{cl}}) = \exp(-2\times 10^5 \mu_{\text{MeV}}^{-2}), \qquad (1)$$

where S_{cl} is the classical Euclidean action for the bounce solution of the relativistic quarks.

In either tunneling or going over the top, there is an additional effect that hinders $d\bar{d}$ production. As the string is stretched it is likely that it will break and produce mesons before reaching a length μ^{-1}. The probability that the string stretchs to a distance μ^{-1} without breaking has been estimated based on experimental information on the width of high-spin mesons,[11] and on the width of charmonium states.[12] The probability p_V of stretching the string to a length μ^{-1} based on the high-spin-meson widths is $p_V \simeq \exp(-2\times 10^4 \mu_{\text{MeV}}^{-2})$, while the charmonium potential suggests $p_V \simeq \exp(-10^3 \mu_{\text{MeV}}^{-2})$. If these estimates are reliable, then even if it is energetically possible to stretch the string far enough to produce diquarks, for $\mu \leq 50$ MeV the branching ratio may be small enough to account for the lack of FCS detection at accelerators. The smallness of the vacuum persistence p_V and of the tunneling probability p_T also hinder production in the early universe.

The scenario for diquark production in the early universe proceeds as follows: At temperatures T much greater than Λ, quarks and gluons behave as an ideal gas. Above same critical temperature T_c, the thermal effects of the ambient background gas should restore the symmetry so that $SU(3)_C$ is unbroken. A reasonable estimate for the critical temperature is $T_c = \mu g^{-1}$, where g is of the order of a gauge coupling constant.[13] Since it is necessary that $\mu \lesssim 50$ MeV $< \Lambda$, and g should be of $O(1)$, T_c should be less than Λ. Therefore, when the confinement transition takes place ($T \simeq \Lambda$) $SU(3)_C$ will be unbroken, and only color singlets (hence integer charge) hadrons will form in the transition.[14] Hence, when it is possible to form FCS ($T \leq T_c$), production must proceed from color-singlet initial states.

The evolution of the diquark abundance is given by the Boltzmann equation

$$\dot{n}_d = \dot{n}_{\bar{d}} = (n_{\text{eq}}^2 - n_d^2)\,|v|\,\sigma - 3(\dot{R}/R)n_d\,, \qquad (2)$$

where n_d is the number density of diquarks, $|\sigma v|$ is the annihilation cross section, n_{eq} is the equilibrium diquark abundance, and the $\dot{R}/R$ term represents the dilution of their density due to the expansion of the universe. If annihilations are ignored then the diquark density is given by

$$n_d/n_\gamma = \int n_\gamma^{-1} n_{\text{eq}}^2 |v|\,\sigma\,dt, \qquad (3)$$

where n_γ is the photon number density [$n_\gamma = 2\zeta(3)T^3/\pi^2$], and the limits of integration are from the time of $SU(3)_C \to SO(3)_g$ symmetry breaking, t_c, to the age of the universe t_U. The cross section in Eq. (3) is equal to the tunneling cross section if the total c.m. energy, $\sqrt{s}$, is less than the height of the barrier $V^{\max}$, and is equal to a constant if $\sqrt{s} \geq V^{\max}$. Therefore Eq. (3) becomes

$$n_d/n_\gamma = \int_{t_c}^{t_U} n_\gamma^{-1}[n_{\text{eq}}^2(\sqrt{s} \leq V^{\max})\,|v|\,\sigma_T + n_{\text{eq}}^2(\sqrt{s} \geq V^{\max})\,|v|\,\sigma_B]\,dt, \qquad (4)$$

where the tunneling cross section is $|v|\,\sigma_T = p_T p_V \sigma_0$ with $\sigma_0 \simeq \Lambda^{-2} \simeq m_\pi^{-2}$, and if $s \geq V^{\max}$ the cross section is $|v|\,\sigma_B = p_V \sigma_0$. With p_T given by Eq. (1), the contribution of the first term is negligible if $g \leq 1$, and Eq. (4) may be easily integrated to yield a diquark-to-photon ratio of

$$n_d/n_\gamma = 10^{26} \mu_{\text{MeV}}^{-3} g p_V \exp(-2\times 10^5 g \mu_{\text{MeV}}^{-2})\,. \qquad (5)$$

Here g relates the temperature of symmetry breaking, T_c, to μ: $T_c = g^{-1}\mu$.

Traditionally the FCS abundance has been expressed relative to nucleons. Since $n_N \simeq 10^{-10} n_\gamma$, the diquark-to-nucleon ratio, n_d/n_N, may easily be obtained from Eq. (5). The relative diquark abundance is shown as a function of μ in Fig. 1 for several values of g.[15] We have integrated the Boltzmann equation including annihilations and have found that annihilations are unimportant if $n_d/n_N \lesssim 1$, so that Eq. (5) is a good approximation.[16] Since the c.m. energy required to go over the barrier is large ($\simeq V^{\max}$), as long as the FCS mass is $\lesssim V^{\max}/2$, FCS production is independent of the mass [cf. Eq. (5)]. Thus Eq. (5) is also valid for the DGJ model where $m_q \simeq V^{\max}/2$. [If $m_q = V^{\max}/2$, there will be kinematic threshold factors which also suppress production, $(s - 4m_g^{\,2})^n$; however, these factors are unimportant compared to the uncertainties in the exponential suppression factors, and so we shall neglect them.]

Now consider diquark production by cosmic rays. Diquarks (or quarks in the DGJ model) can be produced in the debris of high-energy collisions if the c.m. energy is sufficient to produce a quark jet of energy greater than $V^{\max}$.[17] It seems likely that for $\sqrt{s} \geq 10V^{\max}$, there will be a quark jet produced with energy greater than $V^{\max}$. Therefore for incident proton energies greater than $E_0 = (10V^{\max})^2/2m_{\text{proton}}$, cosmic-ray collisions will produce $d\bar{d}$ pairs with probability p_V. The flux of cosmic rays with energy greater than E_0 is[18] $J(E > E_0) = 14.4[E_0/(1\text{ GeV})]^{-1.75}\text{ cm}^{-2}\text{ s}^{-1}$. The total number of diquarks produced is

$$N_d = J(E > E_0)A\tau = 2.5\times10^{25}\mu_{\text{MeV}}^{3.5}p_V, \qquad (6)$$

where we have used the cross-sectional area of the earth ($A = 1.3\times10^{18}\text{ cm}^2$), and the age of the earth ($\tau = 1.4\times10^{17}$ s), and again have used $V^{\max} = (200\text{ GeV})\mu_{\text{MeV}}^{-1}$. Translation of the total number of diquarks produced into a diquark-to-nucleon ratio depends on the depth to which the surface material of the earth has been mixed during the lifetime of the earth. The number of nucleons on the "surface" of the earth is about $1.7\times10^{50}[l_M/(100\text{ km})]$, where l_M is the mixing depth. The FCS abundance expected for $l_M = 100$ km is shown in Fig. 1 for $p_V = \exp(-2\times10^4\mu_{\text{MeV}}^{-2})$. Although the cross-sectional area of the moon is only about 7% of that of the earth, the fact that the mixing depth on the moon is probably only $O(1\text{ m})$ means that the diquarks produced have been mixed with very few nucleons, and so the diquark-to-nucleon abundance expected is greater than the terrestrial value:

$$\left(\frac{n_d}{n_N}\right)^{\text{moon}} = \left(\frac{n_d}{n_N}\right)^{\text{earth}} \frac{(l_M\rho)^{\text{earth}}}{(l_M\rho)^{\text{moon}}} \simeq 1.6\times10^5\left(\frac{n_d}{n_N}\right)^{\text{earth}}. \qquad (7)$$

The cosmic-ray production of FCS as a function of μ is indicated by the dashed curves in Fig. 1. Once again, our results apply both in the SGS model and the DGJ model.

To summarize, if QCD is spontaneously broken at a scale $\mu \ll \Lambda$, then production of FCS in the early universe is qualitatively different than if $\mu \geq \Lambda$.[7] For reasonable values of μ (10 MeV $\leq \mu \leq$ 50 MeV) and g (order unity), early-universe production of FCS can be enormous ($n_d \simeq n_\gamma \simeq 10^{10}n_N$) or insignificant ($n_d/n_N \lesssim 10^{-50}$). However, it is clear from Fig. 1 that if the lack of FCS production at accelerators should require $\mu < 20$ MeV, astrophysical production of FCS is insignificant. Cosmic-ray production of FCS may be the dominant process if $g \simeq 1$, and since the surface of the moon has been geologically dead for the past 4×10^9 years, one expects the diquark-to-nucleon abundance in lunar samples to be enhanced by about 10^5 relative to terrestrial samples. Finally, we note that in both astrophysical production scenarios discussed here, d and $\bar{d}$ abundances should be equal—a fact that may be

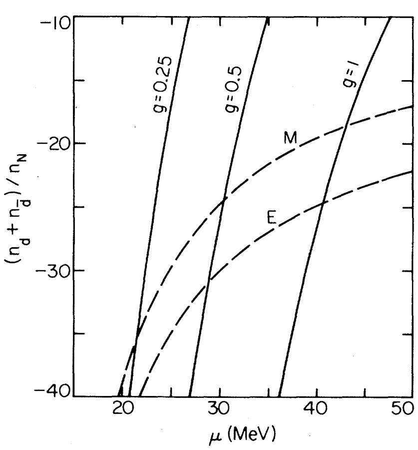

FIG. 1. The FCS-to-nucleon ratio expected from early-universe production (solid curves), and from cosmic-ray production (dashed curves) on earth (E) and on the moon (M) as a function of the QCD-breaking scale parameter μ. The temperature at which QCD breaks is related to μ by $T_c = g^{-1}\mu$. For both sets of curves we have assumed that $p_V = \exp(-2\times10^4\mu_{\text{MeV}}^{-2})$.

relevant if diquarks are to serve as a fusion catalyst.[19]

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[13]In calculating the transition temperature, nonperturbative effects may make any estimate based on perturbation theory meaningless.

[14]This will be true, of course, if $T_c \lesssim \Lambda$.

[15]If $g \gtrsim 1$, tunneling becomes important, and the result is approximately the same as for $g = 1$.

[16]We have assumed that $m_d \lesssim V^{\max}/2$.

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Upper Limit for *T*-Invariance Violation in Elastic *pp* Scattering

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A model-independent upper limit on the time-reversal-invariance (TRI) violating amplitude in elastic *pp* scattering based on a direct experimental reconstruction of the scattering matrix is presented. This analysis is made possible by a set of sixteen polarization parameters measured at the Swiss Institute for Nuclear Research for 579-MeV elastic *pp* scattering at six center-of-mass angles between 66° and 86°. Less than 1% is found for the fraction of TRI-violating cross section.

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Since the discovery of time-reversal-invariance (TRI) violation in the weak decay of the long-lived neutral kaon,[1] no further evidence of TRI violation in other processes has been found. In the electromagnetic interactions, measurements of the neutron electric dipole moment[2] have shown no violation to a level of about 0.03%, and until recently, tests of detailed balance and polariza-

