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IS FRACTIONAL ELECTRIC CHARGE PROBLEMATIC FOR QCD?*

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ABSTRACT

A model of broken QCD is described here; SU_3^c is broken to SO_3^g ("g" for "glow") such that color triplets become₃glow triplets. With this breaking pattern, there should exist low-mass, fractionally-charged diquark states that are not strongly bound to nuclei, but are rarely produced at present accelerator facilities. The breaking of QCD can be done with a 27^c , in which case, this strong interaction theory is easily embedded in unified models such as those based on SU_5 , SO_{10} , or E_6 . This work was done in collaboration with Terry Goldman of Los Alamos and Gordon Shaw of U.C., Irvine.

The work described in this talk was done with Terry Goldman and Gordon Shaw.¹ In it we took seriously the growing experimental evidence that a tiny portion of the charged particles in Nature have fractional electric charge, $\pm \frac{1}{3} + \text{integer}$.² There has been more than the usual resistance to accepting this evidence, and not all of the grumbling is based on experimental objections, but rather on the understanding that the observation of fractional charge would likely require a major revision of QCD. Although it is not yet proven that unbroken QCD must confine, there is some good evidence for it. Thus if fractional charges do exist and are related to quarks, the sub-nuclear fermions hypothesized by Gell-Mann and Zweig to be the basic constituents of hadrons,³ then SU_3^c must be broken. There is an example of broken QCD,⁴ but the breaking mechanism described there requires other major modifications of the standard model: the breaking is done with color-triplet scalars. Usually $\underline{3}^c$ particles of any kind must have fractional charge, although these $\underline{3}^c$ are electrically neutral. Moreover, the origin of additional symmetries imposed on their model is not clear. So, some resistance to accepting Ref. 2 is based on the belief that a major revision of the theory is necessary to accommodate fractional charge. However, in the final analysis, we must await further experimental confirmation of those observations.

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Our object is to show that the observation of fractional charge does not require such drastic modifications of the standard model, but can be achieved by extending slightly the symmetry-breaking scheme of the theory. Before proceeding, several comments are in order. What we have tried to do is construct a gauge-theoretic phenomenology that allows for fractional charge, and at the same time disturbs the standard model very little. In a gauge-theory framework there are many restrictions, and so the model has many predictions for the properties and interactions of fractionally-charged particles. However, there are other ways that fractional charge could appear in a theory. For example, in many models we may simply modify the electric-charge operator. There could be color nonsinglets with charge assignments outside the usual triality-

electric charge constraints such as a $\underline{3}^c$ with charge $+1/3$, so that composite color singlets containing such a constituent would be fractionally charged. In the same vein, there could be color singlet "leptons" with fractional charge. It is possible in some unified models to construct such unusual electric charge operators, although it is not possible for theories based on the exceptional groups. It appears that these schemes would⁵ imply an unacceptably large abundance of relic fractional charge left from the "big bang."

The widely-held belief that unbroken nonabelian gauge theories confine the charges of the local-symmetry group is accepted here. We also assume that the fractionally-charged particles have constituents that are not color singlets, so that they are strongly interacting. We maintain the usual connection between charge and SU_3^c

triality. Consequently, in the context of QCD we must find a symmetry-breaking pattern that allows for unconfined fractional charge. Thus, our model is constructed in the spirit of that of De Rújula, Giles, and Jaffe (DGJ),⁴ but we propose a different symmetry-breaking scheme for SU_3^c , and as a result, the implications of the model are quite different. The most important differences are, phenomenologically, that some fractionally-charged states (called diquarks) need not be very massive and need not be strongly bound to nuclei, although their production at present machine energies is suppressed; theoretically, the breaking scheme is easily embedded into unified models and requires neither extra symmetries nor abnormal electric-charge assignments. After presenting the salient features of our version of broken QCD including the spectrum of fractionally charged particles, we discuss in the following order some questions regarding the symmetry breaking, masses, production cross sections, the formation of bound states with nuclei, and the highly speculative yet enormously important application of using the diquarks as a catalyst in fusion reactors.⁶

SU_3 has two maximal subgroups, SO_3 and $SU_2 \times U_1$, and several subgroups that are not maximal, including SU_2 , $U_1 \times U_1$, U_1 , and

nothing. We assume that SU_3^C is broken to its maximal subgroup SO_3^8 (g for "glow"); this embedding is defined by the branching rule $\underline{3}^C = \underline{3}^8$. (Mathematically this is the same embedding that is used in nuclear physics.)⁷ In breaking SU_3^C to SO_3^8 , three gluons are left massless and the other five gluons get equal masses μ . In the short distance limit $\ll 1/\mu$, this theory retains the features of unbroken QCD. At long distances all eight gluons are confined, since all have nonzero glow, the SO_3^8 charges.

We now describe the spectrum of multiquark states, assuming that glow is confined.⁸ One of the original motivations for selecting SU_3 as the color group was the statistics problem encountered in forming the observed $\underline{56}$, $k = 0$ of low-mass baryons from three quarks.⁹ The $\underline{56}$ is formed from the symmetrized product, $(\underline{6} \times \underline{6} \times \underline{6})_S$ of SU_6 , where SU_6 contains the $SU_3^e \times SU_2^s$ of the Eightfold Way (e) and rotational symmetry as $\underline{6} = (\underline{3}^e, \underline{2}^s)$, so $(\underline{6}^3)_S = (\underline{8}^e, \underline{2}^s) + (\underline{10}^e, \underline{4}^s)$. (All irreps, including those of SU_2 , are labeled by their dimensions.) With the addition of the color quantum numbers the group becomes $SU_6 \times SU_3^C$, and the triquark states satisfying the exclusion principle are

$$[(\underline{6}, \underline{3}^C)^3]_A = (\underline{56}, \underline{1}^C) + (\underline{70}, \underline{8}^C) + (\underline{20}, \underline{10}^C) \quad (1)$$

The SO_3^8 singlets exactly coincide with the SU_3^C ones, since the branching rules are $\underline{8}^C = \underline{3}^8 + \underline{5}^8$ and $\underline{10}^C = \underline{3}^8 + \underline{7}^8$. For any other subgroup of SU_3^C there are $\underline{70}$'s, $\underline{20}$'s, or both that are unconfined, since the $\underline{8}^C$, $\underline{10}^C$, or both will have subgroup singlets. The SO_3^8 breaking scheme is the only one that does not disturb the spectrum of triquark singlets.

One objection in the past to an SO_3^8 "color" group is the existence of diquark singlets: the difermions are

$$[(\underline{6}, \underline{3}^C)^2]_A = (\underline{21}, \underline{\bar{3}}^C) + (\underline{15}, \underline{6}^C) \quad (2)$$

and the $(\underline{15}, \underline{6}^C)$ has a glow singlet, since $\underline{6}^C = \underline{1}^8 + \underline{5}^8$. The $SU_3^e \times SU_2^s$ content of the $\underline{15} = (\underline{6}^e, \underline{1}^s) + (\underline{\bar{3}}^e, \underline{3}^s)$. The Q diquarks may have masses quite similar to the $q\bar{q}$ color-singlet masses.

The spectrum of "low-mass" observable scalar diquarks in Eq. (2) consists of a strong isotriplet with charges of uu , ud , and dd , and so forth. The three axial-vector diquarks have the SU_3^e quantum numbers of $\bar{u}$, $\bar{d}$, and $\bar{s}$.¹⁰ The relative mass of the scalar and axial vector states is model dependent. If the average mass difference is dominated by the color magnetic splitting, then its value is $-1/4$ that of the 1^C $q\bar{q}$ system. However, if only the glow-magnetic spin splitting dominates, then the scalars are lower in mass. Thus, it is possible that the charge $4/3$ scalar is the least massive and is the only stable diquark state. (For example, $M_{\Sigma^+} = 1189.3$ MeV, $M_{\Sigma^0} = 1192.5$ MeV, and $M_{\Sigma^-} = 1197.3$ MeV suggests this pattern may be correct, at least if the gluon field, which is flavor independent, is ignored.) Of course, the theory itself does not depend on the speculation that this is the correct ordering of masses.

The $q\bar{q}$ color-singlet states coincide with the glow singlets. The diquark-antiquark states do include glow singlets:

$$(\bar{6}, \bar{3}^C) \times [(6, 3^C)^2]_a = (6, \bar{6}^C) + (\underline{120}, \bar{6}^C) + (6, 3^C) + (6, \underline{3}^C) + (\underline{84}, \underline{3}^C) + (\underline{120}, \underline{3}^C) + (6, \underline{15}^C) + (\underline{84}, \underline{15}^C). \quad (3)$$

The $\underline{15}^C$ and $\underline{3}^C$ contain no glow singlets, but $\underline{6}^C$ does; thus the observable states, classified by $SU_3^e \times SU_2^s$ are $\underline{6} = (\underline{3}^e, \underline{2}^s)$ and $\underline{120} = (\underline{3}^e, \underline{2}^s) + (\bar{6}^e, \underline{2}^s) + (\underline{15}^e, \underline{2}^s) + (\underline{3}^e, \underline{4}^s) + (\underline{15}^e, \underline{4}^s)$. These are expected to be more massive than diquarks.

In a sense, the unbroken SO_3^8 theory cannot confine quarks because quark-gluon bound states always have a glow singlet. Thus a quark in $\underline{3}^C$ can bind with a gluon in $\underline{8}^C$ to form states in $\bar{6}^C + \underline{3}^C + \underline{15}^C$, where the $\bar{6}^C$ has a glow-singlet piece that may be indistinguishable from a quark. However, in analogy with QCD, where glue-ball masses are expected to be around 1.6 GeV (perhaps), the masses of the quark-gluon states should be at least as large as typical baryon masses and significantly larger than the diquark masses.¹¹

The breaking of SU_3^C can be described in terms of an irrep or a conjugate pair of complex irreps of (effective) Higgs spinless bosons. If the spinless bosons satisfy the usual connection between SU_3^C triality and electric charge, the breaking representation must have zero triality: $\underline{8}^C$, $\underline{10}^C + \bar{\underline{10}}^C$, $\underline{27}^C$, $\underline{28}^C + \bar{\underline{28}}^C$, $\underline{35}^C + \bar{\underline{35}}^C$, $\underline{55}^C + \bar{\underline{55}}^C$, $\underline{64}^C$, etc. However, before an irrep can break a group to a subgroup,

it must have a subgroup singlet, and of these, only the $\underline{27}^c$ and $\underline{28}^c$ + $\overline{\underline{28}}^c$ have SO_3 singlets. (The $\underline{28}$ is formed irreducibly as $[\underline{3}^6]_s$. The $\underline{27}^c$ has the SO_3 branching rule $\underline{27}^c = \underline{1}^8 + \underline{5}^8 + \underline{5}^8 + \underline{7}^8 + \underline{9}^8$.) The simplest candidate for breaking color to glow is a $\underline{27}^c$ of completely flavorless spinless bosons; the Lagrangian mass is expected to be $O(\mu)$. Production of the unconfined $\underline{1}^8$ "Higgs" is suppressed as are all $\underline{1}^8$ that are not $\underline{1}^c$ states. The radiative corrections to the Higgs mass could be large enough that its experimental signature would be similar to that of a glueball.

The (effective) Higgs potential depends on the second-order invariant in $(\underline{27}^2)_s$ (defined as $|\underline{27}|^2$), two third-order invariants in $(\underline{27}^3)_s$, and four independent fourth-order invariants in $(\underline{27}^4)_s$ [including $(|\underline{27}|^2)^2$].¹² There are ranges of parameters for which the $\underline{27}$ can break SU_3^c either to SO_3 or to $SU_2 \times U_1$, which are the two maximal little groups of the $\underline{27}$. If we assume that the fourth-order Higgs potential provides an exhaustive description of the possible breaking directions, then by Michel's conjecture these are the only two possible breaking directions for a single $\underline{27}^c$ with a nonzero vacuum expectation value.¹³ We may select the parameters in a range where the breaking is to SO_3 .

An important advantage of the $\underline{27}^c$ breaking scheme is the ease with which it can be embedded into unified models of electromagnetic, weak, and strong interactions. In most models the electric charge of a state is related to the triality of its color representation; only triality-zero irreps have neutral states. Also, the $\underline{27}^c$ is the irrep of highest weight in adjoint times adjoint; the (effective) scalars could conceivably be formed as a set of bound states. This possibility is easily generalized to unified models.

In particular, for SU_5 the irrep with highest weight in $\underline{24} \times \underline{24}$ is the $\underline{200}$, which contains the $SU_2^w \times SU_3^c (\times U_1)$ irreps $(\underline{1}, \underline{1}^c) + (\underline{3}, \underline{1}^c) + (\underline{5}, \underline{1}^c) + (\underline{2}, \underline{3}^c) + (\underline{2}, \underline{\bar{3}}^c) + (\underline{4}, \underline{3}^c) + (\underline{4}, \underline{\bar{3}}^c) + (\underline{3}, \underline{6}^c) + (\underline{3}, \underline{\bar{6}}^c) + (\underline{1}, \underline{8}^c) + (\underline{3}, \underline{8}^c) + (\underline{2}, \underline{15}^c) + (\underline{2}, \underline{\bar{15}}^c) + (\underline{1}, \underline{27}^c)$, where the U_1 charges of all irreps with zero SU_3^c triality are zero¹⁴. Thus the $(\underline{1}, \underline{27}^c)$ set of scalars carries no flavor at all. The one glow singlet in the $\underline{27}^c$, although strongly interacting, is rarely produced and would be difficult to detect; the rest of the $\underline{27}^c$ is confined. (We do not analyze the bound states with a scalar here: the quantum-number analysis is trivial and the mass estimates unreliable.) The $\underline{200}$ is by far the smallest irrep of SU_5 with a $\underline{27}^c$.

A similar discussion can be made for SO_{10} and E_6 . In SO_{10} the 27^c breaking is in the $\underline{770}$ only, which is the highest weight in $\underline{45} \times \underline{45}$; the only 27^c is in the SU_5 $\underline{200}$ piece of the $\underline{770}$. In E_6 the irrep of highest weight in adjoint times adjoint is the $\underline{2430}$ and the only 27^c is in the SO_{10} $\underline{770}$. After spontaneous symmetry breaking the exact local symmetry of these theories could easily be $U_2 \sim SO_3^8 \times U_1^{em}$, not $U_3 \sim SU_3^c \times U_1^{em}$, as is supposed in the usual approach.

Recall that in the breaking pattern suggested by DGJ, all eight gluons have equal nonzero masses and the vacuum has a global SU_3 symmetry. However, SU_3^c is completely broken so that all the multiquark states in $[(\underline{6}, \underline{3}^c)^k]_8$ ($k = 1, 2, \dots$) would be observable. DGJ expect from bag-model considerations that a gluon cloud of extension $1/\mu$ forms around any state that is not a color singlet. The self energy of the gluon field contributes to the mass, so the color nonsinglet states are much heavier than the color singlets. It is μ and not the QCD scale parameter that determines the mass of color nonsinglets; for small μ all fractionally charged states should be very heavy (of order $1/\mu$) in their model.^{4,15}

If SU_3^c is broken to SO_3^8 , the situation is very different. Of the eight colored gluons, the $\underline{3}^8$ is massless and the $\underline{5}^8$ has mass μ . It is important to keep in mind that all gluons carry glow, so the volume of the gluon cloud may be determined by the confinement mechanism of glow. However, the diquarks carry $\underline{6}^c$ of color, and the shielding of $\underline{5}^8$ gluons may require a slightly larger hadronic volume, making mass estimates difficult. For the $\underline{1}^8$ Higgs in the 27^c , this could be a significant contribution to the physical mass. For the moment we assume the $\underline{5}^8$ gluon contribution to the mass to be negligible, consistent with earlier comments on the bound-state mass splitting. Then, the glow-singlet masses should be controlled by the scale parameter of SO_3^8 , which must be equal to the QCD scale parameter for $Q^2 \gg \mu^2$, since the gauge coupling of SU_3^c is the coupling of any of its subgroups in this domain. Consequently we expect the glow-singlet diquarks in Eq. (2) to have masses of order of the $q \bar{q}$ mesons, and the $q q \bar{q}$ glow singlets in Eq. (3) to have typical baryon masses. Estimates of quark-gluon masses lie between baryon masses and glueball masses of QCD. Note that the proton mass is below the quark-gluon-diquark threshold. If the $\underline{5}^8$ field configuration is similar to that supposed by DGJ, then the diquark mass should go as $1/\mu$ as $\mu \rightarrow 0$. However, Georgi has argued that QCD

should be broken in a first-order phase transition, so μ cannot be arbitrarily small.¹⁶ Then the diquark mass spectrum is discontinuous, and there is no small μ limit. Nevertheless, we assume that μ is small enough that production processes are suppressed.

Although these glow-singlets appear in many ways as ordinary hadrons, their production cross sections are severely suppressed. As a typical example, a rough estimate of the production of light diquarks in $e^+ e^-$ annihilation can be obtained from the linear-potential model of confinement in QCD.¹⁷ At distances less than $1/\mu$ (≈ 200 fermi/ μ , where μ is now in MeV), all eight gluons participate in the confinement of color nonsinglets and the potential rises linearly to ~ 200 GeV/ μ . At distances greater than 200 f/ μ , the channel-dependent potential between states with glow continues to rise, but that between glow singlets drops rapidly to zero due to the finite range of the 5^8 gluons and the low mass of the free diquarks. In this picture a diquark must tunnel through the barrier or be energetic enough to go over the top of the barrier to reach the separation of 200 f/ μ needed to cut off the confinement contribution of the 5^8 gluons. In addition, the production is suppressed by its form factor, since the diquark is a composite particle. A gluon mass of order 50 MeV may be small enough to explain the experimental dearth of diquarks at PETRA. Similar arguments may be made for other scattering processes. Since μ can be large (but less than Λ), the gluon-cloud contribution to the diquark mass [$O(\Lambda^2/\mu)$] may be less than 1 GeV.

Copious production of diquarks might be achieved at lower energies in heavy-ion collisions. Let us suppose the interaction region can be viewed as a quark gas at an elevated temperature in approximate thermodynamic equilibrium. Then the effective confining potential of QCD may be constant over the interaction region, and the probability of finding glow singlets (that are not color singlets) separated by 200 f/ μ or more should be greatly enhanced.

The diquark Q is a strongly interacting particle with interaction range similar to the color-singlet hadrons. The QCD description of the collision region should be unaltered, since $\mu < 200$ MeV. Neither the color-singlet states formed during the final stages of the collision nor the final-state (glow-singlet) diquark couple to the 5^8 gluons, and so the character of the final-state interactions should be similar to those in meson-nucleon scattering. If the mass estimates above are also qualitatively correct, we should not expect any deeply bound states.

The stability of the triality ± 1 states, such as the quark-gluon or antiquark-diquark states, depends on mass relations. In this regard $\bar{Q} N$ (or $\bar{Q}$ nucleus) scattering is of special interest. If the $\bar{Q} N$ threshold is above the lowest-mass triality-one state, then the $\bar{Q} N$ may annihilate to it plus photons or mesons. The

lowest-mass triality-one state would then be nearly stable, since its decay to diquarks would have to proceed by interactions similar to those responsible for proton decay in unified models.

The more exciting possibility technologically is for the lowest-mass triality-one state to have a mass greater than the $\bar{Q} N$ threshold. Then the lowest-mass $\bar{Q}$ state is not only stable against weak decay, but it is not annihilated in interactions with nucleons. If this state is the $\bar{u} \bar{u}$ antidiquark as speculated above, then it should form long-lived $\bar{Q}$ -nucleus "atoms" of size around 50 fermi. An anomaly without fractional charge could involve a gluon singlet such as the Higgs particle and a nucleus. Such composites would have shortened mean free paths.

Finally, for such charge $-4/3 \bar{u} \bar{u}$ diquarks, we stress an enormously important application. In the event that they can be mined or produced in high-energy collisions and are not annihilated in collisions with nucleons, these $\bar{Q}$'s can be used as a catalyst in fusion reactors. As detailed by Zweig,⁶ the perfect catalyst is a relatively light, stable, charged $-4/3$ "quark" that does not bind very strongly to nuclei. A charge less than -1 greatly enhances the binding in, for example, D-D or D-T molecules. For $\bar{Q}$ masses of order 600 MeV, the molecular size is small and the fusion rapid.⁶ Typical temperatures projected for fusion reactors are ~ 10 KeV, which should be sufficient to free a large fraction of the $\bar{Q}$'s from atomic states in the resultant fusion products.

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10. It is a trivial matter to add other quarks to this classification scheme.
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12. A fairly simple way to derive these results is to use the maximal subgroup embedding, $E_6 \supset SU_3$, with $\underline{27} = \underline{27}$, compute the E_6 tensor products $(\underline{27}^3)_s = \underline{1} + \underline{650} + \underline{3003}$ and $(\underline{27}^4)_s = \underline{27} + \underline{351}' + \underline{7722} + \underline{19305}'$, and count the SU_3 singlets from the $E_6 \supset SU_3$ branching rules. Ref. 14 is an aid to this last step, but a few more tensor products are needed to count the singlets in the $\underline{19305}'$, which has just two singlets.
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