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Observable Fractional Electric Charge in Broken Quantum Chromodynamics

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Quantum chromodynamics may be broken to $SO(3)^g$ with $SU(3)^c$ triplets becoming $SO(3)^g$ triplets. If so, there could 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 quantum chromodynamics can be done with a $\underline{27}^c$, 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). Possible relevance of diquarks for fusion catalysis is described.

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There is growing evidence that a tiny portion of the charged particles in nature have fractional electric charge, $\pm\frac{1}{3}e$ + integer.[1] It is possible that such particles are (closely related to) quarks, the subnuclear fermions hypothesized by Gell-Mann and Zweig to be the basic constituents of hadrons[2]; it is also conceivable that such particles behave like leptons, or that they are color-singlet hadrons containing constituents with electric charge outside the usual assignments, although these schemes appear to leave large abundances of relic fractional charge from the "big bang."[3] In any case it appears interesting to formulate gauge-theoretic phenomenologies that allow for observable states with fractional charge.[4] We argue here that such observations do not require a major revision of quantum chromodynamics (QCD) or unified models, but only a slight extension of the symmetry-breaking scheme.

The widely held belief that unbroken non-Abelian 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. 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, some of the implications of the model are quite different. The most important differences are as follows: Phenomenologically, observable fractionally charged states are not single quarks. The lowest-mass observable states are diquarks Q (two quarks bound together). It is possible that they are neither very massive nor strongly bound to nuclei. We argue qualitatively that their production is suppressed. Theoretically, the breaking scheme is easily embedded into unified models and requires neither extra symmetries nor abnormal electric-charge assignments. We then go on to speculate about the masses, production cross sections, the formation of bound states with nuclei, and the enormously important application of using the diquarks as a catalyst in fusion reactors.[5]

SU(3) has two maximal subgroups, SO(3) and $SU(2)\otimes U(1)$. We assume that $SU(3)^c$ is broken to its maximal subgroup $SO(3)^g$ (g for "glow"); this embedding is defined by the branching rule $\underline{3}^c=\underline{3}^g$. (Mathematically this is the same embedding that is used in nuclear physics.[6]) In breaking $SU(3)^c$ to $SO(3)^g$, 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 $SO(3)^g$ charge.

We now describe the spectrum of multiquark states, assuming that glow is confined.[7] One of the original motivations for selecting SU(3) as the color group was the statistics problem encountered in forming the observed $\underline{56}$, $l=0$ of low-mass baryons from three quarks.[8] The $\underline{56}$ is formed from the symmetrized product $(\underline{6}\otimes\underline{6}\otimes\underline{6})_s$ of SU(6), where SU(6) contains the $SU(3)^e\otimes SU(2)^s$ of the "eightfold way" (e) and rotational symmetry as $\underline{6}=(\underline{3}^e,\underline{2}^s)$, so that $(\underline{6}^3)_s=(\underline{8}^e,\underline{2}^s)\oplus(\underline{10}^e,\underline{4}^s)$. [All irreps, including those of SU(2), are labeled by their dimensions.] With the addition of the color quantum numbers the group becomes $SU(6)\otimes SU(3)^c$, and the triquark states satisfying the exclusion principle are

$$[(\underline{6},\underline{3}^c)^3]_a=(\underline{56},\underline{1}^c)\oplus(\underline{70},\underline{8}^c)\oplus(\underline{20},\underline{10}^c). \qquad (1)$$

The $SO(3)^g$ singlets exactly coincide with the $SU(3)^c$ ones, since the branching rules are $\underline{8}^c=\underline{3}^g\oplus\underline{5}^g$ and $\underline{10}^c=\underline{3}^g\oplus\underline{7}^g$. 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 $SO(3)^g$ breaking scheme is the only one that does not disturb the spectrum of triquark singlets.

One objection in the past to an $SO(3)^g$ "color" group is the existence of diquark singlets: The difermions are

$$[(\underline{6},\underline{3}^c)^2]_a=(\underline{21},\underline{3}^{*c})\oplus(\underline{15},\underline{6}^c), \qquad (2)$$

and the $(\underline{15},\underline{6}^c)$ has a glow singlet, since $\underline{6}^c=\underline{1}^g\oplus\underline{5}^g$. The $SU(3)^e\otimes SU(2)^s$ content of the $\underline{15}$ is $(\underline{6}^e,\underline{1}^s)\oplus(\underline{3}^{*e},\underline{3}^s)$. 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 $\bar{s}$.[9] 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 $-\frac{1}{4}$ that of the $\underline{1}^c$ $q\bar{q}$ system. However, if only the glow-magnetic spin splitting dominates, then the scalars are lower in mass. Thus, although it depends on the relative contribution of the $\underline{3}^g$ and $\underline{5}^g$ gluons, it is possible that the charge $\frac{4}{3}$ scalar is the least massive and is the only stable diquark state.

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

$$(\underline{6}^*,\underline{3}^{*c})\otimes[(\underline{6},\underline{3}^c)^2]_a=(\underline{6},\underline{6}^{*c})\oplus(\underline{120},\underline{6}^{*c})\oplus(\underline{6},\underline{3}^c)\oplus(\underline{6},\underline{3}^c)\oplus(\underline{84},\underline{3}^c)\oplus(\underline{120},\underline{3}^c)\oplus(\underline{6},\underline{15}^c)\oplus(\underline{84},\underline{15}^c). \qquad (3)$$

The $\underline{15}^c$ and $\underline{3}^c$ contain no glow singlets, but $\underline{6}^c$ does; thus the observable states, classified by $SU(3)^e\otimes SU(2)^s$, are $\underline{6}=(\underline{3}^e,\underline{2}^s)$ and $\underline{120}=(\underline{3}^e,\underline{2}^s)\oplus(\underline{6}^{*e},\underline{2}^s)\oplus(\underline{15}^e,\underline{2}^s)\oplus(\underline{3}^e,\underline{4}^s)\oplus(\underline{15}^e,\underline{4}^s)$. These are expected to be more massive than diquarks.

In a sense, the unbroken $SO(3)^g$ 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 $\underline{6}^{*c}\oplus\underline{3}^c\oplus\underline{15}^c$, where the $\underline{6}^{*c}$ has a glow-singlet piece that may be indistinguishable from a quark. However, in analogy with QCD, where glueball 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.[10]

The breaking of $SU(3)^c$ can be described in terms of an irreducible representation (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: $\underline{8}^c$, $\underline{10}^c\oplus\underline{10}^{*c}$, $\underline{27}^c$, $\underline{28}^c\oplus\underline{28}^{*c}$, $\underline{35}^c\oplus\underline{35}^{*c}$, $\underline{55}^c\oplus\underline{55}^{*c}$, $\underline{64}^c$, etc. However, before an irrep can break a group to a subgroup, it must have a subgroup singlet, and of these, only the $\underline{27}^c=\underline{1}^g\oplus\underline{5}^g\oplus\underline{5}^g\oplus\underline{7}^g\oplus\underline{9}^g$ and $\underline{28}^c\oplus\underline{28}^{*c}$ have SO(3) singlets. 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 $\underline{1}^g$ Higgs particle is suppressed as are all $\underline{1}^g$ that are not $\underline{1}^c$ states.

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 $(|\underline{27}|^2)^2$].[11] There are ranges of parameters for which the $\underline{27}$ can break $SU(3)^c$ either to SO(3) or to $SU(2)\otimes U(1)$, 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.[12] We may select the parameters in a range where the breaking is to $SO(3)^g$.

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.

In particular, for SU(5) the irrep with highest weight in $\underline{24}\times\underline{24}$ is the $\underline{200}$, which contains the $SU(2)^w\otimes SU(3)^c\,[\otimes U(1)]$ irreps

$$(\underline{1},\underline{1}^c)\oplus(\underline{3},\underline{1}^c)\oplus(\underline{5},\underline{1}^c)\oplus(\underline{2},\underline{3}^c)\oplus(\underline{2},\underline{3}^{*c})\oplus(\underline{4},\underline{3}^c)\oplus(\underline{4},\underline{3}^{*c})\oplus(\underline{3},\underline{6}^c)\oplus(\underline{3},\underline{6}^{*c})\oplus(\underline{1},\underline{8}^c)\oplus(\underline{3},\underline{8}^c)\oplus(\underline{2},\underline{15}^c)\oplus(\underline{2},\underline{15}^{*c})\oplus(\underline{1},\underline{27}^c),$$

where the U(1) charges of all irreps with zero $SU(3)^c$ triality are zero.[13] 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 $\underline{200}$ 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 $\underline{27}^c$ breaking is in the $\underline{770}$ only, which is the highest weight in $\underline{45}\times\underline{45}$; the only $\underline{27}^c$ 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 $\underline{27}^c$ is in the SO(10) $\underline{770}$. After spontaneous symmetry breaking the exact local symmetry of these theories could easily be $U(2)\sim SO(3)^g\otimes U(1)^{em}$, not $U(3)\sim SU(3)^c \otimes U(1)^{em}$, as is supposed in the usual approach.

We now turn to the much more speculative issues of the glow-singlet masses and production. 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, and so the color nonsinglet states are much heavier than the color singlets. In their picture it is μ and not the QCD scale parameter Λ that determines the mass of color nonsinglets; for small μ *all* fractionally charged states should be very heavy (of order $1/\mu$).[4,14] However, this limit is uncertain: Georgi has argued that QCD is broken in a first-order phase transition, in which case μ cannot be arbitrarily small,[15] the $\underline{1}^g$ mass spectrum is discontinuous at a finite μ, and there is no small-μ limit.

We examine here two possibilities, supposing that μ is less than, but not much less than, Λ. The first is similar to the DGJ scheme, where the shielding of the $\underline{5}^g$ gluons from a $\underline{1}^g$ state with color may require a gluon cloud with volume larger than the usual hadronic volume. In particular this could be a significant contribution to the mass of the $\underline{1}^g$ Higgs particle in the $\underline{27}^c$, but it is difficult at present to get a reliable estimate of this contribution. (Note that the Higgs particle could be mistaken for a glueball.)

A second possibility, because *all* the gluons carry glow, is for the volume of the gluon cloud of a $\underline{1}^g$ state to be determined by the confinement mechanism of glow. Then the $\underline{5}^g$ gluon contribution to the $\underline{1}^g$ masses could be small, since it would be controlled by the scale parameter of $SO(3)^g$ [which is Λ because for $Q^2\gg\mu^2$, the running coupling of $SU(3)^c$ is the same as its unbroken subgroups]. Then the 1^g diquarks in Eq. (2) have masses like the $q\bar{q}$ mesons, the $qq\bar{q}$ $\underline{1}^g$ states in Eq. (3) have masses like baryons, and the quark-gluon masses should lie between baryon and glueball masses.

We now speculate why production processes could be sufficiently suppressed, even if μ/Λ is not very small. Imagine separating a diquark and an antidiquark starting from a separation $\ll 1/\Lambda$. Note first that since the diquarks are composite particles, there is a form factor suppression to get into this channel, as exists for specific hadrons such as $p\bar{p}$. As the separation increases beyond $1/\Lambda$, but is still less than $1/\mu$, the confinement mechanism of QCD acts very strongly to rearrange the system into color singlets. Thus, when the separation of the system reaches $1/\mu$, few glow singlets have survived. Beyond $1/\mu$, the system freely decomposes into the asymptotic states of mesons, baryons, and glow singlets.

Copious production of diquarks might be achieved at lower energies in heavy-ion collisions. Let us suppose that 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 large distances should be greatly enhanced.

The diquark Q is a strongly interacting particle with interaction range similar to the color-singlet hadrons, unless μ is very small. 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 $\underline{5}^g$ 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 $\overline{Q}N$ (or $\overline{Q}$-nucleus) scattering is of special interest. If the $\overline{Q}N$ threshold is above the lowest-mass triality-one state, then the $\overline{Q}N$ 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 $\overline{Q}N$ threshold. Then the lowest-mass $\overline{Q}$ state is not only stable against weak decay, but it is not annihilated in interactions with nucleons. If this state is the $\overline{u}\overline{u}$ antiquark as speculated above, then it should form long-lived $\overline{Q}$-nucleus "atoms" of size around 50 fm.

Not all implications of the symmetry breakdown to glow are associated with fractional charge. For example, a glow singlet such as the electrically neutral Higgs particle and a nucleus could form a composite with shortened interaction length.

Finally, for such charge $-\frac{4}{3}$ $\overline{u}\overline{u}$ 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 $\overline{Q}$'s can be used as a catalyst in fusion reactors. As detailed by Zweig,[5] the perfect catalyst is a relatively light, stable, charged $-\frac{4}{3}$ "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 $\overline{Q}$ masses of order 600 MeV, the molecular size is small and the fusion rapid.[5] Typical temperatures projected for fusion reactors are ~ 10 keV, which should be sufficient to free a large fraction of the $\overline{Q}$'s from atomic states in the resultant fusion products.

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[9]It is a trivial matter to add other quarks to this classification scheme.

[10]The derivation of more quantitative estimates is hindered by the lack of experimental guidelines.

[11]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 $(\underline{27}^3)_s = \underline{1} \oplus \underline{650} \oplus \underline{3003}$ and $(\underline{27}^4)_s = \underline{27} \oplus \underline{351}' \oplus \underline{7722} + \underline{19\,305}'$, and count the $SU(3)$ singlets from the $E(6) \supset SU(3)$ branching rules. W. G. McKay and J. Patera, *Tables of Dimensions, Indices, and Branching Rules for Representations of Simple Lie Algebras* (Dekker, New York, 1981).

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the gluon force becomes very strong before its range $1/\mu$. Thus the bag model may misestimate the hadronic size for small μ. Nevertheless, it is attractive to impose confinement by supposing the bag size is $O(1/\mu)$ as $\mu \to 0$.

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Deciphering the Structure of Hadronic Final States

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We apply a theory of hadronic final states based on the notion of elongated quark-antiquark structure to the multiparticle production in e^+e^- annihilation. Our results compare very accurately with the available experimental information.

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The study of the structure of hadronic final states produced in high-energy collisions is recently attracting a considerable amount of experimental interest. In particular, the analyses carried out at PETRA[1] have demonstrated that the hadronic final states produced in high-energy e^+e^- annihilation show a characteristic jet structure, that is being tentatively interpreted in terms of quarks and gluon jets. Similar results appear to emerge in the analysis of inelastic states produced in proton-proton collisions at CERN intersecting storage rings at $\sqrt{s} \sim 50$ GeV,[2] revealing striking resemblance between these final states and the ones produced in e^+e^- annihilations.

In this Letter we wish to draw attention to the fact that one such theoretical framework exists in the massive-quark model (MQM)[3] and quark geometrodynamics (QGD),[4] and to show that it leads to an extremely accurate description of hadronic final states in e^+e^- physics from 3 to 30 GeV c.m. energy. We will also argue that the success of this approach casts doubts upon the current theoretical interpretation of high-sphericity (high-p_T) events in e^+e^- annihilation as being due to gluon jets. Furthermore, it should be stressed that, having at our disposal such a precise tool, to describe hadronic final states, we can use it in the future to analyze the mechanisms of production of final states in a variety of hadronic processes, from proton-proton to deep-inelastic lepton-hadron scattering.

Before presenting our results, let us recall briefly the main theoretical points of our approach.

(i) Multiparticle states originate from the decay of well-defined physical objects, elongated quark-antiquark structures.

(ii) These structures are coherent superpositions of physical $q\bar{q}$ (mesonic) states, with different angular momenta and approximately degenerate masses. In three-space they describe the notion of quark-antiquark waves moving freely inside a cylinder of length R_1 increasing with mass, and a constant (or logarithmically increasing) radius R_T.

(iii) The decay of a given structure is described by a sequential process, in which the initial structure decays into two mesonic states, and so on until all the energy of the initial $q\bar{q}$ system is spent in the production of hadronic stable particles (π's, K's, and η's).

The theoretical basis of (i) and (ii) has been described at length in several previous papers.[5,6] Here we shall only recall that the elongated structure is the *confined* $q\bar{q}$ state which most closely reproduces the structure of quark-partonic (unconfined) states.

In the QGD framework,[4] the basic decay [see Fig. 1(a)] of the elongated $q\bar{q}$ structure (S) into two such structures,

$$S(M) \to S_1(m_1) + S_2(m_2), \tag{1}$$

can be calculated by the space-time overlap of the wave functions (WF's), available from QGD.[4] In particular, in Ref. 6 we have given the expression for the kernel, K chain in the case of the kinematically favored decay

$$S \to \mu + S'(m'), \tag{2}$$

where μ denotes a low-mass mesonic state. It