PHYSICAL REVIEW D VOLUME 27, NUMBER 7 1 APRIL 1983

Colored models for anomalous nuclei

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(Received 23 August 1982; revised manuscript received 9 December 1982)

There seems to be good experimental evidence that anomalous nuclei are produced in heavy-ion collisions; they are anomalous in that they have an abnormally short mean free path, for example, in nuclear emulsions. Here we consider the possibility that anomalous nuclei are combinations of a colored anomalous particle fragment (based on theories with spontaneous breakdown of color symmetry) with ordinary nucleons. Phenomenological implications of various possible models in which the anomalous particle fragment is considered to be a colored particle with the color symmetry $SU(3)_c$ explicitly broken are given.

I. INTRODUCTION

It now seems possible to argue at a reasonable level of confidence that anomalous nuclei exist.[1–4] Anomalous nuclei (hereafter referred to as AN) are secondary nuclei produced in high-energy heavy-ion collisions with target nuclei, which so far have only been seen in nuclear-emulsion experiments. They demonstrate an anomalously short mean free path,[5] roughly 2.5 cm in emulsion, independent of charge, to be contrasted with a much longer mean free path

$$\lambda \sim 30Z^{-0.45}\ \text{cm} \tag{1}$$

for normal nuclei. Approximately 6% of the secondaries in heavy-ion collisions are anomalous, while the proportion in the third (tertiary) and later generations is higher.[2] As the anomalous nuclei are capable of traveling the distance 2.5 cm in nuclear emulsion, their lifetime must be greater than 10^{-10} sec; however, since interactions rather than decays are seen the lifetime can be assumed to be longer than 10^{-8} sec. The charge of the anomalous nuclei seems to be integral.[1] Within conventional nuclear-physics concepts, there seem to be no explanations for these anomalous nuclei.

Because of the lack of additional information, it is not surprising that these observations have spawned a number of incompatible theoretical models.[6–12] The models can be roughly separated into two categories: (a) nuclear models,[11,12] where the anomalous property is due to some peculiar collective excitation of the nuclear state (a crude analogy might be a giant dipole resonance) and (b) particle models,[6–10] in which the anomalous nuclei contain an anomalous particle whose interactions dominate the scattering process. The enhanced cross section might seem to suggest the presence of a new superstrong interaction. In this case, the natural candidate from particle physics would be some form of color excitation, either in the form of hidden color[6,7,9,10] or naked color.[8,10] In particle physics, a number of people have investigated the possibility that color may not be an exact symmetry, as assumed in QCD, but may be broken. Such investigations include those of Pati and Salam[13] (PS), De Rújula, Giles, and Jaffe[14] (DGJ), Slansky, Goldman, and Shaw[15] (SGS), and a model recently discussed by three of us[16] (SSW). In this paper we consider the phenomenological implications of the anomalous particle being some form of color excitation, hidden or naked.

The models involving color have a number of features in common. It is widely believed, although not proven, that if the local color gauge symmetry $SU(3)_c$ remains exact, then the massless colored gluons provide a confinement mechanism. The known hadrons are postulated to be color singlets. The interactions between color singlets would be produced by the color analog of the van der Waals force, and would thus be much weaker than the actual color force. In such a picture, however, failure to observe colored particles (color nonsinglets) does not necessarily imply nonexistence of such states; rather such failure may be a reflection of the difficulty in producing such states in ordinary hadronic collisions. The models involving broken color all suggest the idea of quasiconfinement, that is, confinement up to some moderately large distance, after which it breaks down. It is argued below that quasiconfinement would lead to suppression of the production of naked color in (e.g.) pp or e^+e^- collisions, while not affecting production in heavy-ion collisions.

The principal problem with the discussion of anomalous nuclei is the sheer lack of experimental information other than what has been outlined in the opening paragraphs. In Sec. II we discuss a number of features of colored models. Insofar as is possible, in Sec. III, we try to outline the predictions of each model for anomalous nuclear behavior and suggest new features to look for in further experiments. In this way it may be possible to establish or discard the broken-color model as a viable explanation for anomalous nuclear behavior.

II. SOME GENERAL FEATURES OF BROKEN-COLOR MODELS

As described in the Introduction, it does not seem possible to explain anomalous nuclear behavior in terms of conventional nuclear-physics or particle-physics concepts. The spectrum of states from the conventional quark model of elementary particles would have a baryon and a meson made up from three quarks and a quark-antiquark pair, respectively, bound into color-singlet states. Color-nonsinglet states are presumed to have infinite mass so that they are not observable. Nuclei built up from color-singlet hadrons have the normal behavior, namely have sizes given by $R=1.2A^{1/3}$ fm and geometric cross sections of the order of mb. Short mean free paths characterizing anomalous nuclear fragments, and hence larger than geometric cross sections, might signal either an anomalously large size for them or an extra-strong interaction or both. The extra-strong interaction might be attributable to color forces which are known to be very strong.

In broken-color models colored objects made up, for example, from new combinations of quarks, antiquarks, and gluons (which are not color singlets) but may have masses which are finite (though large), may be produced when the available energy is in excess of the production threshold for such states. We label such particles as anomalous particle fragments (APF's). We consider the possibility that anomalous nuclei are made up of such an APF bound with other ordinary nucleons. Such an object may exhibit anomalous features outlined above, like large size, large interaction, etc. Below, we outline in some detail the broken-color model discussed recently by three of us[16] and indicate the properties to be expected for the APF in such a model and hence for the anomalous nucleus of which the APF is a constituent. We also summarize the findings in the case of the other liberated or broken-color models.

In the SSW scheme,[16] the color SU(3) is assumed to be spontaneously broken down to SU(2)$\times$U(1) by an octet of Higgs bosons, and it is assumed[17] that the charge corresponding to the SU(2)$\times$U(1) subgroup remains confined so that only SU(2)$\times$U(1) singlets are observable. As all the representations of the SU(3) group which contain the SU(2)$\times$U(1) singlet have triality zero, the observable spectrum of particles is restricted to states which have integral electric charge. A natural candidate for the APF is a colored qqq or qq system from $\underline{8}_c$. Again the (assumed) large Higgs-boson mass (e.g., >10 GeV) forbids fast decay ($<10^{-10}$ sec) into color-singlet states and provides desirable partial conservation of color. For completeness, we outline the main features of this model here.

The (effective) Lagrangian which defines the theory has the form

$$\mathscr{L}=\mathscr{L}_{\rm QCD}+(D_\mu\phi)^\dagger(D^\mu\phi)-V(\phi)+g'\bar{\psi}\phi^iF^i\psi\ , \tag{2}$$

where $\mathscr{L}_{\rm QCD}$ is the standard Lagrangian of quantum chromodynamics, D_μ is the covariant derivative in the eight-dimensional representation of SU(3)$_c$, ϕ is the color octet of Higgs scalars, and the last term is quark—Higgs-boson interaction with a coupling constant g' [F_i, $i=1,\ldots,8$ are the standard generators of SU(3)]. The Higgs potential $V(\phi)$ has the form

$$V(\phi)=a\phi^i\phi^i+bd^{ijk}\phi^i\phi^j\phi^k+c\phi^i\phi^i\phi^j\phi^j\ , \tag{3}$$

where d^{ijk} is the symmetric SU(3) tensor.

We assume that the constants a,b,c satisfy the relation $3b^2>32ac$. Then $V(\phi)$ has absolute minimum at $\phi^8=v$, $\phi^i=0$, $i=1,\ldots,7$, where

$$v=\frac{\sqrt{3}b+(3b^2-32ac)^{1/2}}{8c}\ . \tag{4}$$

This minimum is not symmetric under the whole SU(3)$_c$ group but only under the SU(2)$\times$U(1) subgroup generated by F^1, F^2, F^3, and F^8. Now, following the standard prescription, we find that the gluons A^1,A^2,A^3,A^8 remain massless, while A^4,A^5,A^6,A^7 acquire a mass

$$\mu^2=3g^2v^2/2\ . \tag{5}$$

The mass of the Higgs field $\eta^8=\phi^8-v$ is

$$M_8{}^2=\sqrt{8}b\mu/g-6a\ , \tag{6}$$

the mass of the ϕ_1, ϕ_2, ϕ_3 fields is

$$M_1{}^2=M_2{}^2=M_3{}^2=\sqrt{18}b\mu/g\ . \tag{7}$$

As we do not have any theoretical arguments for the values of these parameters we have to assume (in order not to spoil temporary asymptotic freedom) that the gluon mass is small, e.g., less than 100 MeV. By choosing $a \ll b \ll c$, the mass of the Higgs boson can be made large (but to produce mass ratios of the desired magnitude one may have to retune the input parameters after the incorporation of radiative corrections in the masses). Notice that the SU(3) octet decomposes according to

$$8 = 2 \oplus 2 \oplus 3 \oplus 1 \tag{8}$$

under SU(2) and U(1), in other words, one gluon (A^8) is both massless and unconfined.

In order to estimate the cross section, we must go to a dynamical model. We may use the bag model to estimate the mass and radius of the state. Doing the calculation we find that the state, while larger than conventional hadrons ($R \sim 2.5$ fm) is still too small to explain the anomalously large cross section. The problem here may well lie with the bag model itself: from the point of view of two-body interactions, the bag is very hard.

We may obtain another estimate, avoiding the bag model, by considering a logarithmic confining potential superimposed on the one-gluon-exchange interaction. For the color-singlet meson, taking the wave function

$$\psi_s = (q_1\bar{q}_1 + q_2\bar{q}_2 + q_3\bar{q}_3)/\sqrt{3} ,$$

we get

$$V_s(r) = -\frac{2}{3}\frac{\alpha_s}{r} - \frac{2}{3}\frac{\alpha_s}{r}\exp(-\mu r) \tag{9}$$

[in the limit $\mu \to 0$ this reduces to the usual Coulomb potential $(-\frac{4}{3}\alpha_s/r)$]. Repeating the procedure for the colored meson with wave function

$$\psi_c = (q_1\bar{q}_1 + q_2\bar{q}_2 - 2q_3\bar{q}_3)/\sqrt{6} ,$$

we obtain

$$V = -\frac{1}{2}\frac{\alpha_s}{r} + \frac{2}{3}\frac{\alpha_s}{r}\exp(-\mu r) . \tag{10}$$

We see that the potential in the colored state differs from that in the color-singlet state by having a repulsive part and a weak attractive part. This implies that the colored state is more weakly bound and consequently it has larger size and mass. If the radius of the colored state is about 7 fm, then one can get an order-of-magnitude increase in the geometric cross section. Solving the Schrödinger equation numerically, we find that to obtain this radius the gluon mass has to be about 70 MeV. The mass difference between these two states is then about 0.6 GeV, which is in line with the observation that there is a threshold for the production of anomalous nuclei of approximately 1 GeV/nucleon.

Similarly, for the qqq system, taking the wave function for the $\underline{8}_c$ state

$$\psi = (q_1q_3q_2 + q_3q_1q_2 - q_2q_3q_1 - q_3q_2q_1)/2 ,$$

the potential is

$$V^c_{qqq} = -\frac{5}{12}\frac{\alpha_s}{r_{13}} - \frac{5}{12}\frac{\alpha_s}{r_{23}} - \frac{1}{6}\frac{\alpha_s}{r_{12}} + \frac{1}{2}\frac{\alpha_s}{r_{12}} e^{(-\mu r_{12})} .$$

Here we again see a repulsive part and a weak attractive part and such a potential naturally leads to the desired large size of the object.

The great difference between the bag- and potential-model descriptions should not be unexpected. Both of them have parameters optimized to describe a narrow range of phenomena: It is not surprising that they have quite different implications when used outside their region of validity.

Now we turn our attention to the possible decay of the colored states into color-singlet ones. Such transitions are possible because color does not have to be conserved, as the radial component of the Noether current does not vanish sufficiently rapidly at infinity. For simplicity we will consider only the transition $\underline{8}_c \to \underline{1}_c$ in the case of mesons. The relevant S-matrix element is

$$\left\langle \frac{1}{\sqrt{3}}(q_1\bar{q}_1 + q_2\bar{q}_2 + q_3\bar{q}_3) \,|\, S \,|\, \frac{1}{\sqrt{6}}(q_1\bar{q}_1 + q_2\bar{q}_2 - 2q_3\bar{q}_3) \right\rangle .$$

The lowest-order diagrams contributing to this process are shown in Fig. 1. The Feynman rules can be derived from the Lagrangian (2). The loop integrals are convergent and can easily be estimated. To get the rate of the process we multiply the nonrelativistic limit of the cross section for the process by $|f(0)|^2$, where $f(r)$ is the solution to the Schrödinger equation with the potential (9). Taking the gluon mass to be 70 MeV and the quark mass 330 MeV, the result is

$$\Gamma \lesssim \frac{g'^2}{{M_8}^4} 10^{-7}\ \text{GeV} .$$

It is obvious that one can choose plausible values for the parameters g' and M_8 to obtain lifetimes greater than 10^{-10} sec as apparently required for the

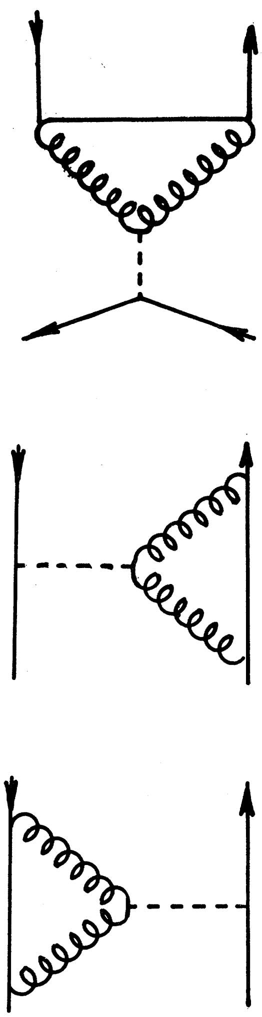

FIG. 1. Decay of colored particles in the SSW model. The dashed lines represent Higgs particles.

anomalous nuclei, e.g., $M_8 > 10$ GeV and $g' < 1$ will do that.[18]

As the large lifetime is caused mainly by the large radius of the colored state and especially by the large Higgs-boson mass, the lifetime of the colored baryons is also of the same order.

The suppression of the color-nonconserving processes is an important feature of this model; it explains why the anomalous nuclei show up in third-generation (tertiary) fragments even more strongly than in the secondaries. Once color is quasiconserved a colored secondary particle will almost certainly produce a colored tertiary.

This model is very rich in free colored states. In addition to the usual three-quark and quark-antiquark states used above, there is the massless colored gluon (A^8), and a Higgs particle which should have a mass $\gtrsim 1$ GeV if one is not to encounter fine tuning problems. The Higgs particle provides no problem: it would also form anomalous nuclei and these would obviously not be distinguishable from nuclei containing a colored quark system.

The massless colored gluon in this model may have consequences worth investigating. The strength of the color interaction might suggest that the gluon + nucleus would in fact be lighter than the nucleus itself. If such were the case, this model could clearly be rejected. We have found that within the bag model the colored *qqqg* (gluon + nucleon) state is indeed at a higher mass than the *qqq* state, the additional energy arising from the gluon zero-point energy. Also, since A^8 *is* colored, the quasiconfinement of color in this model would strongly suppress A^8 production in hadronic collisions. A^8 will also be difficult to detect: obviously it would not show up in emulsion experiments, except in combination with nucleons.

There is one more kind of colored particle in this model. It is well known that a spontaneous breakdown of symmetry to a group with a U(1) subgroup gives rise to "magnetic" monopoles with a mass

$$M \sim \frac{\mu}{\alpha_s}$$

(where μ is the mass of the gauge particles). This particle is then light ($\lesssim 1$ GeV), stable, and carries a chromomagnetic charge $g_s \sim 1/2\alpha_s$ but no electric charge. However, the quasiconfinement of color might suppress the production of these objects.

There is an interesting extension of this model which gives rise to another interpretation of APF. Suppose that the U(1) symmetry is broken with a triplet of Higgs bosons, so that the complete scheme is

$$\mathrm{SU(3)} \underset{\underline{8}}{\rightarrow} \mathrm{SU(2)} \times \mathrm{U(1)} \underset{\underline{3}}{\rightarrow} \mathrm{SU(2)}$$

with

$$\langle \underline{3} \rangle = \begin{bmatrix} 0 \\ 0 \\ V' \end{bmatrix} .$$

If the Higgs couplings are chosen such that

$$\Lambda_{\mathrm{QCD}} \gtrsim \langle \underline{8} \rangle \gg \langle \underline{3} \rangle ,$$

the color monopoles then become unstable, with exponentially long lifetimes because the decays occur by tunneling. They then become very attractive candidates for APF themselves: they are long-lived ($\gg 10^{-10}$ sec), strongly interacting, and relatively light. Furthermore, the Higgs boson H_8 can then be permitted to be relatively light, giving fast decays of colored states.

Another one of the most durable theories of liberated color in particle physics is the Pati-Salam (PS) model,[13,19,20] in which liberated quarks have integer charge. In this model also a number of colored states appear and some of their properties of relevance to APF will be summarized here. Color-nonsinglet hadrons in the model decay into color-singlet hadrons through mass-matrix mixings (arising from spontaneous symmetry breaking) between canonical color-nonsinglet gauge bosons (gluons or leptoquark bosons) and canonical color-singlet gauge bosons. An obvious example is the photon

$$A=A_0\cos\xi+A_8\sin\xi \tag{11}$$

(0 = color singlet, 8 = color octet). The photon necessarily has a color-octet component because red quarks of a given flavor have different electromagnetic charges than the corresponding yellow or blue quarks,

$$\begin{aligned}Q(u_r,u_y,u_b)&=(0,+1,+1)e\ ,\\ Q(d_r,d_y,d_b)&=(-1,0,0)e\ .\end{aligned} \tag{12}$$

The mixing angle ξ is given by

$$\tan\xi=4\alpha/(3\alpha_s-4\alpha)\ ,$$

which becomes small at values of Q^2 appropriate for hadronic binding.[21] The production of color-nonsinglet states in e^+e^- collisions is suppressed by a Glashow-Iliopoulos-Maiani-type[22] mechanism (if such collisions are above the energy threshold for the production of liberated color[23]) through interference between the photon and an orthonormal, predominantly colored gluon with nonzero mass[24]:

$$U=-A_0\sin\xi+A_8\cos\xi\ . \tag{13}$$

Similar mixing occurs between charged gluons $V^{\pm}$, connecting red quarks to yellow and blue quarks of the same flavor, and the $W^{\pm}$ bosons mediating weak charged-current interactions. The mixing angle characterizing the component of color-singlet flavor-changing W present in the (predominantly) color-octet-gluon eigenstates of the gauge-boson mass matrix is given by[19]

$$\tan\theta_{W-V}=gfc_1{}^2/m_W{}^2=(\alpha/\alpha_s)^{1/2}m_V{}^2/m_W{}^2\ , \tag{14}$$

where m_V is the gluon mass. The liberated gluon in the PS model is now estimated to have a mass in excess of 5 GeV, corresponding to $\theta_{W-V}=7\times10^{-4}$. It has been argued elsewhere, however, that the gluon mass appearing in the Lagrangian does not correspond to a liberated gluon mass, but rather, to the effective mass of a gluon inside a color-singlet hadron ($\sim$100 MeV), in which case θ_{W-V} may be as small as 3×10^{-7}. Consequently, some color-nonsinglet mesons may decay into color-singlet mesons with rates suppressed from strong-interaction decay rates by a factor of $\tan^2\theta_{W-V}$; such mesons could conceivably have lifetimes as long as 10^{-10} sec, allowing some possibility of being in line with observations on anomalous nuclei.

Such hadrons would be expected to be massive compared to their color-singlet decay products (provided such hadrons were in their own distinct color-nonsinglet bag) in which case the decay of AN into color-singlet matter would be quite noticeable in emulsion. However, this possibility (AN containing PS colored hadrons) depends on at least some colored hadrons having weak lifetimes; one would have to be very careful that higher-order diagrams (in which color could decay "electromagnetically") would not compete successfully with the lowest-order processes occurring through W-V mixing. In this sense, long-lived colored hadrons are an unnatural candidate for AN constituents.

There is, however, a second possibility in the PS model: quarks themselves are unstable and integer charged. An old problem for the model was to explain the empirical failure to observe such quarks. Quark lifetime decreases rapidly with increasing mass, so several early papers[19,25] argued that quarks with "light" liberated masses (2—3 GeV) would be sufficiently unstable to explain their nonobservation. However, as argued above, massive free quarks could have their masses reduced to near constituent mass (NCM) in the presence of nuclear matter.[23] Thus a short-lifetime massive free quark which would not yet have been observed could be consistent with a long-lived NCM quark in anomalous fragments.

In this model, the short mean free path could arise from some combination of (a) the decay of the NCM quark, (b) the increased cross section of the particle fragment. A heavy quark would be expected to decay into a neutrino and a number of hadrons,[25] giving a starburst for the target fragments. However, an NCM quark (mass $\sim$300 MeV) would decay predominantly into a single pion and a neutrino. It is even possible that much of the decay energy could be absorbed by the nucleus, providing an example of the invisible decay mentioned earlier. Such decays, however, would require quark masses of a few tens of MeV, a mass range difficult to reconcile with standard bound-state models.

If we suppose that NCM quarks provide an explanation for anomalous nuclei, we then have to reconcile light effective quark masses with proton stability. We have found such a reconciliation to be feasible within the parameter space of the PS model (we omit the details of the calculation here).

We describe briefly two other models[14,15] in which a further variety of APF occur. In the DGJ[14] scheme the local color SU(3) symmetry is broken down to global color SU(3) group by three triplets of Higgs particles. All eight gluons are massive. Both quarks (with fractional electric charge) and gluons are then unconfined. There is a relation between the gluon mass μ and the mass M and radius R of the colored state (derived from the bag model):

$$M \simeq \frac{0.076}{\mu}\sqrt{C^2}\ \mathrm{GeV}^2\ ,$$
$$R \simeq (6.4\ \mathrm{GeV}^{-4/3})M^{1/3}\ , \tag{15}$$

where C is the value of the quadratic Casimir operator for the given colored state.

Chapline[8] and Fowler *et al.*[10] have used this model to account for the anomalous nuclei assuming a free quark is bound to a color-singlet nucleus. The lowest-mass fractionally charged state is stable in this model and thus the model may not be compatible with the observations on cosmic-ray production rates of anomalous nuclei.

In the SGS[15] scheme the color SU(3) group is broken down to the SO(3) group by a 27-plet of Higgs particles, and it is assumed that only SO(3) singlets called glow singlets are observable. In addition to the normal color-singlet hadrons, the observable spectrum of hadrons contains qq, qg, qqq states from $\underline{6}$, qqg, $qqqg$, etc. from $\underline{27}$, etc. As a simplest candidate for the integrally charged APF, we can take the state containing three quarks and a gluon. We have shown that such a state is stable against decay into $qq+qg$ or $qq+qqg$ (by the bag-model argument) and, if the Higgs-boson mass is greater than 10 GeV, has a lifetime longer than 10^{-10} sec for decay into a color-singlet state. The glow-singlet state containing two quarks is also colored; such a state is fractionally charged, and hence may be an unlikely candidate for the APF.

In the SGS model, the APF's are colored glow singlets, and fractionally charged particles should also be produced in heavy-ion collisions. Although it is conceivable that the production of qq or qg is suppressed by a factor of 100 compared to the production of $qqqg$ (and therefore they would escape observation since the precision of charge measurements is low[2]), the fractionally charged particles should definitely be there. Further, if we suppose that the large-interaction cross section is due completely to the color of the APF (rather than its radius), it seems likely that the cross section will be proportional to the Casimir operator of the states. This implies that the fractionally charged states (in a $\underline{6}$) would have longer mean free paths than the integrally charged APF from the $\underline{27}$. A rough estimate for this mean free path is 20 cm.

In Table I, we summarize the pertinent facts concerning the different models which have been considered above and what the typical candidates are for the APF.

III. PHENOMENOLOGICAL IMPLICATIONS OF VARIOUS MODELS

In this section, we consider the phenomenological implications of the various models which have been proposed to explain anomalous nuclear behavior. Although we will be primarily considering AN as made up of colored APF's bound with ordinary nucleons, we sill also make a few comments pertaining to the hidden color models. First we consider the question of production of anomalous nuclear states. They seem to have been copiously produced in heavy-ion collisions, but apparently not in pp or e^+e^- collisions. There is also some evidence for an

TABLE I. Properties of models considered.

Model \ Property	Quark charges	Remaining local color symmetry	Higgs-boson contents	Observable charge contents	APF candidate
SSW	Fractional	$SU_{(2)}\times U_{(1)}$	Octet	Integer only	Colored mesons, baryons, and gluons
SGS	Fractional	$SO_{(3)}$	27-plet	Fractional and integer	Diquarks, $qqqg$, etc.
DGJ	Fractional		Octet	Fractional and integer	Any (including free quarks)
PS	Integer		Triplet	Integer only	Free quarks

energy threshold for production. We argue that if anomalous nuclei contain APF's, there may be enormous suppression factors for production in pp or e^+e^- collisions relative to production in heavy-ion collisions. Then we turn to the question of decays and interactions of anomalous nuclei. From the observations in nuclear emulsions in which secondary particles are observed to produce third- and higher-generation particles, it is not even clear whether the higher-generation particles arise as a result of decays of long-lived secondaries or whether they are due to interactions of secondaries. Below we describe briefly the phenomenology of decays and interactions and what one might expect on the basis of APF's being responsible for AN. The short mean free path for the AN may be due either to their large geometric size or due to the possibility that the APF's may have extra-strong long-range interactions. We comment briefly on the possibilities through further experiments of distinguishing between these cases.

A. Production mechanisms

It seems apparent that APF's are not produced in pp or e^+e^- collisions (although, of course, it is permissible to wonder if anyone has looked for them). We consider three possible mechanisms, which are hinted at by theoretical models, which would cause suppression in hadronic or leptonic collisions. These mechanisms are by no means exclusive.

(1) The APF is a multiple- (e.g., six-) quark state.[6,7,9] The production in

$$pp \to AX$$

would then be roughly akin to

$$pp \to d\pi ,$$

at least in the energy range of a few hundred MeV. This is described reasonably well by a single Regge-pole exchange

$$\frac{d\sigma}{dt} \sim \left[\frac{s}{s_0}\right]^{2\alpha(t)-2} . \tag{16}$$

The highest possible trajectory exchange is the N^*; if this is forbidden (by, e.g., isospin), then the nucleon is the leading trajectory with

$$\alpha(t) = 0.1 + t \ (t \text{ in GeV}^2) \tag{17}$$

and so the cross section falls off very rapidly. The Regge residue is related to the wave-function overlap: it requires very little juggling to make the cross section extremely small.

(2) Another possibility is that the APF may be "easily" pair-produced, but the potential between the associated colored pairs may be quasiconfining so that color nonsinglet states may not easily escape from the collision region before rearrangement into color singlets. This possibility is clearly suggested by broken-color models, but may well apply to others. In this context we may suppose that some (or all) gluons acquire a mass μ,[15] where $\mu < \Lambda$, the standard QCD scale. At short distances $\Lambda^{-1} < r < \mu^{-1}$, the mass is irrelevant and the standard QCD confinement mechanism acts to rearrange the system into color singlets. Thus when the system reaches a separation $R \sim 1/\mu$, almost no color-nonsinglet states survive.

We may contrast the above situation with the situation in heavy-ion collisions. It is plausible that here the interaction region is a quark-gluon plasma in approximate thermodynamic equilibrium. Then the effective confining potential will be smeared out to be roughly constant, and hence the probability of finding colored states at large distances would be greatly enhanced. This, of course, presupposes that the APF production is associated and color conserving.

We illustrate how dramatic the differences between pp and heavy-ion collisions could be in practice, by considering a toy model for APF (AN) production. This is based on the coupled-channel Schrödinger equation for the processes $NN \to NN$ [described by a potential $V_{11}(r)$], $NN \to AA$ [$V_{12}(r)$], and $AA \to AA$ [$V_{22}(r)$], where A stands for the APF or AN. To have an exactly soluble model, we take the AN mass to be equal to the nuclear mass, and consider only square-well potentials,

$$\begin{gathered}
V_{11}(r) = \begin{cases} \rho_{11}, & r < R_N \\ 0 , & r > R_N , \end{cases} \\
V_{12}(r) = V_{21}(r) = \begin{cases} \rho_{12}, & r < R_N , \\ 0 , & r > R_N , \end{cases} \\
V_{22}(r) = \begin{cases} 0, & r < R_0 , \\ \rho_A, & R_0 < r < R_A , \\ 0 , & R_A < r , \end{cases}
\end{gathered} \tag{18}$$

where R_N is the nuclear radius, R_0 the radius of the confining core, and R_A the AN radius. Obviously such a model is very arbitrary, but to get an idea of the effect we take the reasonably plausible values $R_A \sim 3$ fm, $\rho_{11} \sim \rho_{12} \sim 40$ MeV, and $\rho_A \sim 500$ MeV.

Typical results are shown in Fig. 2. Note that as the nuclear radius is increased by a factor of 3, the production cross-section is increased by 10^7. It is obviously easy to suppress AN production in collisions of light nuclei to an arbitrarily large amount.

Another feature of this toy model is worth men-

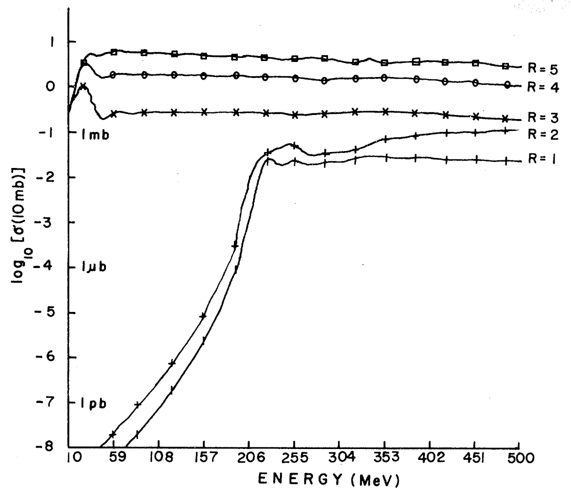

FIG. 2. The production cross section for anomalous nuclei as a function of energy using the potential given in Eq. (18).

tioning here. It is possible to reconcile a high apparent threshold ($E_T \sim 1$ GeV) with a low real threshold E_e. Production of AN will be suppressed until the energy is comparable to the barrier height *unless* the nuclear radius is comparable to the AN barrier radius. (Ironically, perhaps, this effect is known in potential theory as anomalous scattering.) In practice, of course, there will be many coupled channels, mostly normal ones, and the increasing inelasticity will serve to depress the AN production rate as energy increases.

(3) Another common aspect of broken-color models is that colored quarks embedded in nuclear matter are likely to have different properties from isolated free quarks. This has been dubbed the "Archimedes effect" in the PS model,[23] but it seems clear that it will occur in any liberated model.

The binding energy of a quark to the nucleons in the model is so large that the actual mass of the system is considerably below the sum of the masses. Thus the effective threshold (in energy/nucleon) for producing low-A (and hence low-Z) AN may be higher than that for high-Z states: again this might offer an explanation of why anomalous nuclei are only observed in heavy-ion collisions.[26]

It is, however, obvious that AN production cannot be completely eliminated in light nuclei in such models. In this context, it is worth mentioning the observation of a long-lived ($\tau > 10^{-10}$ sec), *negatively* charged ($Q=-1$) state with a mass of 4.3 ± 0.2 GeV produced with a very low cross section in p-Be collisions ($\sim 10^{-11}$ of the pion-production cross section).[27] It is obviously tempting to relate this to anomalous nuclei. It appears that the mass is probably too large to be reconciled with the threshold if this state is the AN itself. However, it would be quite consistent with a light (~ 1 GeV) $Q=-1$ APF attached to three neutrons.

We note, in passing, that the crucial question of whether AN production is associated or solo cannot be decided until the decays are understood. The conservation of color in strong interactions in all of the models imply both associated production and also an increased enhancement of AN in tertiaries.

B. Decays

We suppose (quite generally) that the anomalous nucleus A is the excited state of some normal nucleus N, with a lifetime τ_A. There is then an excitation energy

$$E_e = M_A - M_N ,$$

which can vary in the various models from a few MeV[12] to a few GeV[6]. The threshold suggested by Judek[1] and Jain[3] might imply $E_e \sim 1$ GeV, but, as in the toy model for production, this may be an apparent threshold much higher than the true threshold.

The lack of apparent decays indicates that τ_A is greater than 10^{-10} sec. We consider the three cases $\tau_A > 10^{16}$ sec, $10^{16}\ \mathrm{sec} \gtrsim \tau_A \gtrsim 10^5$ sec, and $\tau_A \lesssim 10^5$ sec. Assuming that the AN does, in fact, decay, the decay properties will depend crucially on whether E_e is large (~ 300 MeV), medium (~ 100 MeV), or small (~ 10 MeV). The decay may be weak[16] (although not necessarily mediated by the standard W or Z bosons), forbidden electromagnetic,[9] or tunneling-suppressed strong.[6]

The mean momentum transfer to AN in projectile fragments in emulsion is ~ 200 MeV/c[1], and hence it is reasonable to assume that an equal number of AN are produced in target fragments with similar momentum. If the AN have similar range-momentum relations to conventional particles, then approximately

$$R_A = A_A (p_A / M_A)^{\alpha} , \tag{19}$$

where

$$A_A = A_p M_p / Z_A{}^2 M_A , \tag{20}$$

M_p is the proton mass, and $A_p \sim 6\times 10^5$. Here p_A is measured in GeV/c, M_A in GeV, R_A in μm, and $\alpha \sim 3.6$. Then the stopping time is

$$\tau_{\mathrm{stop}} = \frac{\alpha}{\alpha-1}\frac{R_A}{p_A} M_A . \tag{21}$$

These give ranges of ~ 100 μm and stopping times of 10^{-11} sec, which implies that any visible decays will be at the end of fairly short, target-fragment tracks.

If the E_e is large, then, independent of the mechanisms of the decay, there will be several particles (pions) involved in the decay like the Λ_c which has only multibody decay modes. Hence this kind of particle would show a prominent star burst at the end of a target-fragment track, with the outgoing particles having momenta of ~ 100 MeV/c. The absence of such events would definitively rule out a number of models.[6,10]

Suppose that $E_e \sim 100$ MeV. Then the decay must be weak or electromagnetic. A weak decay would involve an energetic (~ 50 MeV) electron which would, again, be unmistakable. If the decay is electromagnetic, then it can only be seen via a recoil. The recoil range is

$$R = A_N \left[\frac{E_\gamma}{M_N} \right]^\alpha , \tag{22}$$

where A_N is the value of A from (20) for the recoiling object and M_N is its mass. It seems plausible that if the photon energy E_γ is greater than the nuclear binding energy, E_B, then the AN decay product will recoil in a quasifree fashion. If the decay product has protonic mass, then a 45-MeV photon would give a recoil range of 10 μm, which would be clearly visible. If the whole nucleus recoils, then for ^{4}He, a 10-μm recoil would correspond to 90 MeV.

A direct search for delayed γ decays of anomalous nuclear states has recently been carried out by Liss *et al.*[28] with negative results. They looked for a high-energy γ ray with its energy in the range 140 to 2000 MeV, and did not find any. This result would seem to rule out those models that predict the emission of one or more high-energy photons as the anomalous particle decays electromagnetically into normal matter. However, as the authors point out, this conclusion must be tempered by the following set of possibilities.

(1) The electromagnetic decay might proceed by means of a cascade in which none of the γ-ray energies exceeded 70 MeV.

(2) The branching ratio for electromagnetic decay could be small compared to the branching ratio to all other modes of decay.

(3) The lifetime of the anomalous state could exceed 1.8×10^{-10} sec in those models that predict that 6% of the primary collisions produce anomalous states.

(4) The threshold for production of the anomalous states of matter might be greater than 940 MeV/nucleon.

Each of these possibilities is consistent with the null result of Ref. 28. To this list of qualifiers we add yet another. If the apparent threshold E_T and the real threshold E_e are well separated in energy, so that $E_e < 70$ MeV, then the electromagnetic decay of the anomalous state to a normal state would only produce γ's with energy <70 MeV, which would not have been seen by Liss *et al.*[28] Finally, we would like to point out that the Romo-Watson model[6] is incorrectly included in Ref. 28 among models which predict the emission of high-energy γ's. Since the model predicts that electromagnetic transitions from anomalous states to normal states are highly suppressed, the only γ rays predicted by that model that one would expect to see are those arising from decays within the rotational band of the anomalous fragment, and those γ's would have energies <1 MeV.

If *no* evidence for decays of target fragments is found, then most of the models would be incorrect as associated production is involved in most of them. Decays of target fragments would not be detectable only if (a) the decay is weak or radiative and $E_e < 10$ MeV, or (b) the decay proceeds by a low-energy γ = cascade, or (c) $\tau_A > 1$ week (the development time for the emulsion). It would appear to be hard to reconcile the small E_e in (a) with the apparent 1 GeV/nucleon threshold. Furthermore, if E_e is small, it is extremely difficult to understand why the effect has not been seen in low-energy nuclear physics.

C. Interactions

It might appear obvious that a large interaction cross section would imply that the AN is a geometrically large object. For the sake of definiteness, we assume below that it is a uniform sphere 10 fm in radius. However, it is also possible that the AN is a particle of normal nuclear size (say 2–3 fm) with a long-range strong interaction. This is obviously suggested by color models. It would be useful to outline the possibilities of distinguishing these alternatives experimentally.

The simplest method in principle is to measure the differential cross section. We may try to estimate the differential cross section in the eikonal approximation, with two extreme possibilities: the black-sphere case

$$V(r) = \begin{cases} i\infty , & r < R_A' , \\ 0 , & r > R_A' , \end{cases} \tag{23}$$

and the long-range case

$$V(r) = \frac{\lambda R_1^{N-1}}{(r^2 + R_1{}^2)^{N/2}} . \tag{24}$$

[Note that R_A' in (23) need not necessarily be the R_A used previously.] The value of N would be expected to be 4 if the interaction is charge-induced dipole.

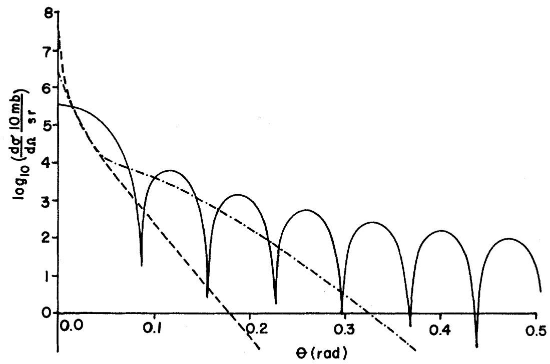

FIG. 3. Differential scattering cross sections for (a) a black sphere (solid curve); (b) long-range case, $N=3$ (dashed curve); (c) long-range case, $N=4$ (dash-dot curve), normalized to the same total cross section of 10 b.

In Fig. 3 we show the differential cross sections arising from the two cases $N=3$ and $N=4$, normalized to the same total cross section, and compared to the cross section for a 2-fm black sphere. Note that the 10-fm case is *much* more sharply forward peaked, as we would expect. At an angle of the order of 0.1 rad., the differential cross section for the 10-fm case is several orders of magnitude smaller than that for the 2-fm + long-range case, thus enabling one in principle to distinguish between these two possibilities.

Another test would be the dependence of the total cross section on the atomic number A of the target nucleus. If the cross sections are indeed geometrical, then

$$\sigma_{\text{tot}} \sim 2\pi(R_A' + R_N)^2 \;. \tag{25}$$

If $R_A' \gg R_N$, and $R_N \sim 1.2A^{1/3}$, then the cross section will depend only very weakly on A. A long-range force, however, would presumably depend on A (or A^2 if the process were coherent), which is a much stronger dependence than the usual $A^{2/3}$ dependence. Preliminary indications[29] favor the former, in that many secondary interactions in emulsion appear to be with hydrogen nuclei.

A further method of distinguishing between a large geometric size and a small object with a long-range strong interaction would, of course, be to measure the electromagnetic size of the AN via electron scattering. To measure a size R one needs a momentum transfer q in an electron scattering experiment, of order $q \sim R^{-1}$. For R in the range from 3 to 10 fm this translates to q in the range 70—20 MeV/c and one would expect to see form-factor effects for q in this range. Of course, such an electron scattering experiment on an AN is not possible, but the inverse process in which the AN scatters electrons may be clearly visible in the form of δ rays in emulsions. The energy distribution of δ rays can be related to the size of the AN. For an AN of 2 GeV/nucleon, the maximum momentum of the δ ray is expected to be about 4 MeV/c, which is considerably less than the 20 MeV/c required to probe a size of the order 10 fm for the AN. The maximum momentum of the δ ray reaches the range 20—70 MeV/c only when the AN has a total energy of 5 to 8 GeV per nucleon. Thus the absence of high-energy δ rays in the current experiments, where AN's have only 1 to 2 GeV per nucleon, is not an indicator of the size of the AN in the 10-fm range.

IV. CONCLUSIONS

It is clear from the foregoing that explanations for anomalous nuclei in which (colored) anomalous particle fragments are constituents of such nuclei is a hypothesis needing further investigation. Crucial questions to be answered are the following.

(a) Do anomalous nuclei exist with low-Z values, in particular, do negatively charged anomalous nuclei exist?

(b) Is there evidence for decays? If so, what are their characteristics?

(c) Is there evidence for associated production of anomalous nuclei?

If the answers to these questions are negative, then the models based on broken or liberated color cannot survive. The most serious theoretical problem for all the colored models is the question of whether quasiconfinement exists. The answer to this question will have to await further theoretical developments.

ACKNOWLEDGMENTS

These ideas have been stimulated by discussions with C. K. Hargrove, R. J. Hemingway, B. Judek, Y. Karant, L. Krauss, A. C. MacPherson, and J. C. Pati, among others. Financial support by the Natural Sciences and Engineering Research Council of Canada is gratefully acknowledged.

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