A Practical Model of Quark Confinement

Apparently hadrons are made of quarks. By now most physicists know that the quantum states of hundreds of baryons and mesons can be classified by the arrangements of (at present) four types of spin-1/2 quarks (and antiquarks) which differ in mass and charge and in four other quantum numbers (upness, downness, strangeness and charm,—"flavors" in the current jargon) which are conserved in strong interactions. Nevertheless so far no quarks have been observed. The forces between them must be quite different from those which bind atoms, molecules or nuclei, since there a constituent can be separated by providing an energy which is a fraction of its rest mass. In contrast, if free quarks existed, they would have to be at least an order of magnitude heavier than the hadrons they compose.[1] Finding a local, relativistic and quantum-mechanical description of quark confinement is the outstanding problem of hadronic physics.

In recent years progress in this direction has been substantial. It is now believed that quarks carry an additional set of quantum numbers, "colors", and that they interact by the exchange of colored, massless, vector gluons. It is hoped that this interaction will explain confinement and all the other properties of composite hadrons. This hope has been difficult to realize in the framework of conventional field theory because non-perturbative methods must be used *ab initio*. (A perturbative formulation begins with free quarks and gluons, which if the program is successful, must miraculously disappear in the finale!) Further, it is not just the effects of a known interaction which must be handled non-perturbatively, but it is the form of the interaction itself which must be modified in a non-perturbative way (for example the potential between heavy colored objects must be changed from $1/r$ to $\sim r$). Symptomatic of the difficulty, present attempts to use the quark-gluon field theory to describe physical systems are burdened with *ad hoc* assumptions and lack any underlying unity.

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We wish to describe here our own work on confinement which has been motivated by the belief that the starting point of conventional Lagrangian field theory is too distant from the phenomena to be useful. We have developed a formulation in which quark confinement is manifest at the outset and which because of its conceptual simplicity offers the hope of providing a unified description of hadronic phenomena. Before we describe our approach it is necessary first briefly to review the historical development of ideas about quark confinement.

Originally—to the extent quarks were taken seriously at all—their absence from the physical spectrum was explained by assuming them to be very heavy and to be bound by very strong forces to make light (~ 1 GeV) hadrons. This picture was difficult to reconcile with the calculations which originally suggested the model. These treated quarks as free and light (~ 300 MeV). Properties of hadrons were obtained simply by adding up the properties of the appropriate quarks.[2] Further, even in the early days of the quark model there was dynamical evidence that they are confined by an unconventional mechanism. This came from the observation of "linear Regge trajectories". Hadrons with the same flavor quantum numbers belong to families whose spins J and masses M follow the formula

$$J = \alpha' M^2 + \alpha_0, \tag{1}$$

where $\alpha'(\sim 0.9 \text{ GeV}^{-2})$ is a universal constant.

The forces between a small number of particles which arise from conventional local, relativistic dynamics fall off exponentially with distance. The addition of a small amount of relative angular momentum keeps the particles sufficiently far away from one another so that they do not bind. In that sort of system a spectrum like Eq. (1) is impossible. On the other hand, if the forces prevent the constituents from getting out, this spectrum is quite natural. When angular momentum is added to a system one expects it to become more extended in space as it spins. The rotational spectrum of Eq. (1) emerges if the mass is proportional to the spatial separation of the constituents ($M \sim L$) and the system has the largest angular momentum compatible with its mass. For if the system spins as fast as consistent with relativity, then the relative momentum p will be of order M so $J \sim pL \sim M^2$. As J increases, instead of coming to pieces the hadron spins more slowly ($\omega \sim 1/L \sim 1/M$) and becomes more massive.[3]

A second striking feature of Eq. (1) is the universality of the slope α'. It is difficult to see how this could emerge in a natural way in a system where the energy and angular momentum are carried predominantly by the quarks, since quarks with different flavors generally have different masses. Furthermore, mesons and baryons have different numbers of constituents but also share the same slope α'.

Thus a model based on heavy quarks bound by very strong but conventional forces was already in trouble with the observed rotational spectra. It was finally laid to rest and replaced with our current picture of permanently confined, light and relatively weakly interacting quarks as a result of a series of beautiful experiments performed by an MIT–SLAC collaboration at the SLAC electron accelerator.[4]

The quark and lepton worlds are united by the weak and electromagnetic interactions. The weak and electromagnetic currents probe the local structure within hadrons. Inelastic electron–proton scattering at high momentum transfer measures the correlation function for current fluctuations at short distances and relative times within the proton. The experiments showed that at high momentum transfer, electrons appear to scatter *elastically* from *light*, spin-1/2 constituents within the proton. This phenomenon, known as Bjorken scaling,[5] has an analog in inelastic electron *nucleus* scattering. There, at momentum transfers large compared to nuclear binding energies (but still small compared to hadronic excitations) the electrons appear to scatter elastically from the constituent *nucleons*. For hadrons, this is the basis of the so called "parton model".[6] This experiment and all subsequent inelastic electron and neutrino scattering studies suggest the proton's "parts" are quarks. In fact there is no serious alternative interpretation of deep inelastic lepton scattering.

It is difficult to exaggerate the theoretical implications of these experiments. First, they identify quarks as the local dynamic variables in terms of which a microscopic theory of hadrons must be constructed. It is too much to expect that a non-field-theoretic description should conspire to impersonate a local quark current! Secondly, they force us to abandon the picture of heavy quarks bound by extremely strong forces—apparently the masses of quarks in a proton target are negligible compared to momentum transfers of roughly 1 GeV. Likewise at distances characteristic of these momenta quark–quark interactions are also negligible. We are *forced* to entertain the seemingly paradoxical notion that hadrons are loosely but permanently bound aggregates of light quarks.

In most relativistic quantum field theories interactions become stronger at short distances, in contrast to what was observed in the MIT–SLAC experiments. A familiar example is QED: as one penetrates the polarization cloud surrounding the electron one "sees" the bare charge, which is infinite. In 1973 it was found[7] that interactions mediated by a non-Abelian gauge theory grow weaker at short distances, and vanish asymptotically at zero separation. (Weisskopf has suggested an intuitive explanation of this phenomenon which is too lengthy to present here.[8]) Non-Abelian gauge theories are generalizations of electrodynamics based on sets of non-commuting charges.[9] In this "asymptotically free" theory quarks have three additional quantum numbers, colors, which we may take as Redness, Greenness and Blueness. (The reason for three will become apparent in a moment.) A quark interacts by the exchange of massless vector gluons coupled in a gauge invariant manner to color changing currents. As long as there are not too many quark flavors such a theory is asymptotically free. It is therefore tempting to adopt the non-Abelian group SU(3)-color as the basis for a gauge theory of quark interactions.

The need for three additional quark quantum numbers had been suspected for quite some time. One of the earliest quark mysteries was how they managed to obey the exclusion principle. The lightest baryon multiplets (the octet, N, Λ, Σ, Ξ and decuplet, Δ, Σ^*, Ξ^*, Ω^-) were classified as three quarks *symmetrized* in spin and flavor. In any simple model the lightest states are expected to be S-waves (spatially symmetric). If quarks were to be fermions they had to carry an additional quantum number in which the baryon states are antisymmetrized.

Color was originally introduced to provide such a quantum number.[10] In order not to increase the number of kinds of hadrons it is postulated that all are color singlets. For baryons, the only singlet combination of three quarks is one totally *anti*symmetrized in the quark's color indices. Hence the resolution of the statistics problem.

Quark confinement can now be rephrased as a color rule: quark interactions must be such that all hadrons are color singlets. One or two quarks cannot form a singlet but three quarks or a quark and an antiquark can. In fact all color singlets have only conventional quantum numbers (integral charge and baryon number) in contrast to the quarks themselves. The problem of quark confinement is now: Why cannot colors be separated? The asymptotically free color gauge theory of quarks and gluons hints at an answer: Since the coupling becomes weak at short distances, it becomes strong at long distances. This is the basis for the hope expressed by the slogan "infrared slavery": Perhaps the force becomes so strong that colors cannot be separated. It remains little more than a hope since at present there exist few methods for handling the strong coupling problem in conventional relativistic quantum field theory. It also may be that when sufficiently powerful techniques are developed to handle the long distance behavior of this theory then it will be found that the coupling gets strong but not strong enough to confine color.

Our approach is a more radical one.[11] We abandon conventional local field theory. Rather than starting with quarks and gluons everywhere, and then trying to confine them, we *start* by assigning colored fields only to the *inside* of hadrons (since that is where they seem to be). This is a radical assumption since we now must distinguish the space inside of a particle from the space outside of it—the hadron must be an "extended" object. If the hadron is extended, the equations which define its spatial boundary must be local in order not to violate causality. It is possible to distinguish an otherwise "empty" region of space in a relativistically invariant way by associating with it locally a constant stress-energy tensor, $T^{\mu\nu} = -Bg^{\mu\nu}$, where $g^{\mu\nu}$ is the ordinary Minkowski metric. $T^{00} = B$ is an energy density, and $T^{kl} = -B\delta^{kl}$ is the stress: We assign such a term to the space *inside* a hadron. These are just the stress and energy one would associate with the interior of a bubble in a uniform isotropic perfect fluid with pressure B. Thus, we confine quark and gluon fields by enclosing them within an extended hadron which, of necessity, also carries an energy B per unit volume. This term keeps the fields together. B acts as a uniform pressure on the surface of the hadron, squeezing the constituent fields. We call this extended space region a "bag".

Quarks and gluons are the quanta of colored fields. A bag with one quark in it would be a physical quark and would also carry a net color. It is easy to see that states such as this never appear in this theory—by construction, color electric fields, like quarks, exist only inside the bag; but by the analog of Gauss' law, color electric fields must emanate from a bag with net color. The only consistent solution is for all bags to be color singlets. It is easy to see that an infinite energy barrier prevents a color singlet bag from fissioning into two colored bags. For suppose in a color singlet bag, the colors C and $-C$ are separated by a finite distance. The color electric field lines which connect the charges C and $-C$ exist only in a finite volume V since (in our model) it costs an energy $BV > 0$ to create the space which carries them. Hence at a point between the charges, the fields span a finite cross-section A and by Gauss' law, $C = EA$, where E is the mean color field strength in the cross-section. Therefore the energy per unit length is $\frac{1}{2}E^2A = \frac{1}{2}C^2/A$. This diverges for fission ($A \to 0$) into two colored bags. Thus color confinement, and hence quark-confinement, is automatic and comprehensible in our model. The mechanism does not rely on "infrared slavery". Note (for later use) that if the tube of color flux is in equilibrium, we would expect that the energy per unit length of the flux lines and the bag volume energy per unit length would be the same, so according to the previous formula, $\frac{1}{2}C^2/A \sim BA$ or $A^2 \sim \frac{1}{2}C^2/B$. Thus, the flux lines are squeezed and fill a cross-section scaled by $1/\sqrt{B}$. So the energy per unit length of the flux tube is given by $(\frac{1}{2}E^2 + B)A \sim 2BA \sim (2BC^2)^{1/2}$.

As the simplicity of this argument illustrates, one virtue of our approach is the promise of a relatively simple and systematic calculation of hadron properties. We have found in studying the theory that there are two broad classes of hadrons, which we shall call spherical and deformed. For the lightest states of up, down and strange quarks (with Compton wavelength large compared to the bag size) one would expect the bag to be spherical with the quarks in S-states to minimize the kinetic energy. To lowest order such a hadron is a bag of free quarks (as seen in deep inelastic electron scattering). To the next order, gluon mediated interactions split degenerate states, renormalize static moments and induce small deviations from parton model results in electron scattering.

Hadrons possessing large amounts of angular momentum or containing heavy (charmed) quarks will deform. Consider, for example, a meson consisting of a quark and an antiquark in a color singlet state. We expect the quark and antiquark to be in a circular orbit well separated from one another. From the discussion in the previous section, the color flux lines which connect the quark and antiquark are expected to form a tube with cross-section $(\frac{1}{2}C^2/B)^{1/2}$ with energy per unit length $k = (2C^2B)^{1/2}$. Therefore the bag will be deformed and take the shape indicated in Figure 1. In a system consisting of a heavy (charmed) quark and heavy antiquark we expect deformation even without high angular momentum. When the quarks are heavy enough that their Compton wavelength is small ($m \simeq 1.5$ GeV for the charmed quark) compared to the size set by the color field pressure ($\sim B^{-1/4}$ which is about 1 fm), then the quarks will act like point localized charges and again the bag will deform to the shape of the quasi-dipole field lines established by the slow-moving, heavy quarks. Again the meson looks like Figure 1.

If the quarks are light, the colors must be kept separated by a finite angular momentum J, and when the resulting state has the maximum spin consistent with its mass (the leading Regge trajectory), the energy and angular momentum will be carried principally in the colored fields. The quark and antiquark in Figure 1 move opposite to each other and at the velocity of light. Such a

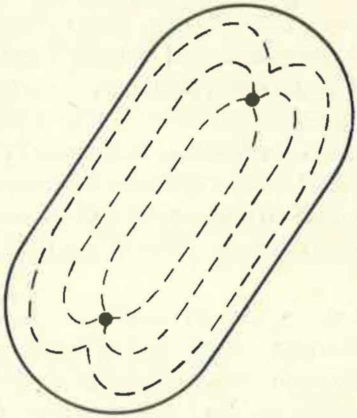

FIGURE 1 A deformed $Q\bar{Q}$ bag, showing the color electric flux lines connected to quark and antiquark.

relativistic system has the spectrum of Eq. (1) with $\alpha' = 1/2\pi k$, where k is the energy per unit length in the instantaneous rest frame of a point between the ends. Since the slope parameter α' depends only on *color* (and not flavor) we find the universality mentioned above.

This description may be pursued more quantitatively.[12] We find

$$\frac{1}{\alpha'} = 8\pi^{3/2}(2\alpha_c BC^2)^{1/2} \tag{2}$$

where α_c is the color fine-structure constant ($\alpha_c = g^2/4\pi$), B is the bag constant and C^2 is the squared-color charge (quadratic Casimir operator) of the quark on one end (or more generally of whatever color is on either end). α_c and B are determined from light hadron spectroscopy (see below) to be 0.55 and 145 MeV, respectively. For a single quark $C^2 = \frac{4}{3}$ and we find $\alpha' = 0.88$ GeV^{-2} in excellent agreement with the observed value of 0.9 GeV^{-2}.

Baryons have similar deformed high angular momentum states. In this case there is a single quark on one end and two quarks with antiquark color on the other, the same color arrangement as in mesons. Hence we obtain the same slope for baryon Regge trajectories as for mesons, in agreement with observation.

Since heavy quarks carry the same colors as light quarks, the energy per unit length in this system is the same as in the light mesons. Here, however, in low mass states the field energy acts as a potential energy which governs the motion of the slowly moving quarks (cf. Born–Oppenheimer). The potential is then predicted to have the form kr with $k = 1/2\pi\alpha'$. With $\alpha' = 0.9$ GeV^{-2}, this gives $k \simeq 0.18\ \text{GeV}^2$, which is very close to the value determined by a study of the J/ψ spectrum.[13]

The model predicts the existence of additional deformed states with more quarks and antiquarks. $Q^2\bar{Q}^2$ is the simplest example. These states should be important in baryon–antibaryon annihilation. Before discussing them in depth we must develop the bag description of round, S-wave hadrons.

In states of light quarks without net orbital angular momentum the confining pressure B is balanced by the quantum pressure of the quarks. To lowest order the mass of a hadron is the sum of the kinetic energies of the quarks, their masses (the up and down quarks are massless and the strange quark weighs 280 MeV) and the bag energy BV: $M(R) = \sum_i (m_i^2 + x_i^2/R^2)^{1/2} + BV$, where x_i/R is the momentum of the ith quark. x_i is fixed by boundary conditions at the bag surface ($x_i = 2.04$ for $m_i = 0$; $x_i = \pi$ for $m_i \to \infty$). The volume is determined dynamically by balancing the field pressure against the confining pressure B, $\partial M/\partial R = 0$. For a proton of radius one fermi we require $B^{1/4} = 145$ MeV. At this stage we find that baryons are roughly one and a half times as massive as mesons and hadron masses grow roughly linearly with the number of strange quarks. There are many spin degeneracies—states with the same quarks in the same spatial state are degenerate. Thus the π and ρ are degenerate as are the nucleon and Δ.

To first order in the color coupling constant quarks exchange gluons. This interaction is primarily magnetic in character. (The local color electric charge density inside a hadron is zero when all the quarks are in the same spatial state.) Like the hyperfine interaction of atomic physics gluon exchange generates a spin–spin force between quarks. The color "hyperfine" interaction splits the π from the ρ and the N from the Δ in the right direction and by nearly the right amount. Color is essential to this result.[14] For example in an *Abelian* (vector) theory the interaction would be proportional to $\sum_{i \neq j} a_i a_j \boldsymbol{\sigma}_i \cdot \boldsymbol{\sigma}_j$ with Abelian charges $a_i = \pm 1$ for quarks and antiquarks. For mesons ($a_i = -a_j$) higher spin states are heavier, but for baryons ($a_i = a_j$) the state of highest spin lies lowest in disagreement with observation. With non-Abelian charges the product $a_i a_j$ is replaced by $\boldsymbol{\lambda}_i \cdot \boldsymbol{\lambda}_j$—$\{\boldsymbol{\lambda}_i\}$ being the color matrices of the *i*th quark analogous (in SU(3)) to its spin matrices $\{\boldsymbol{\sigma}_i\}$. Since both mesons and baryons are color singlets the quarks in both have "opposite" colors ($\boldsymbol{\lambda}_i \cdot \boldsymbol{\lambda}_j < 0$ for both) and the proper splittings are obtained. The spectrum of S-wave baryons and mesons calculated in this approximation to the bag model is shown in Figure 2.[15] The fit involved

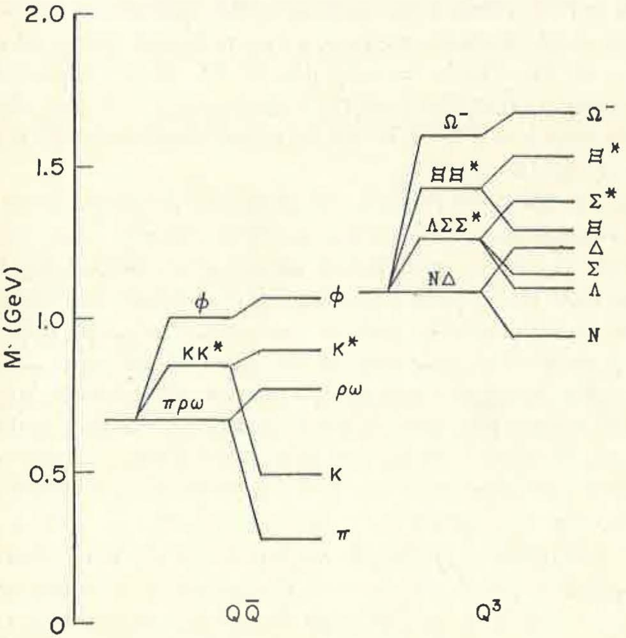

FIGURE 2 S-wave baryons and mesons in the bag model.[14]

four parameters: B, α_c, m_s—the strange quark mass—and z_0—a measure of the zero point fluctuations of the confined fields.

This picture is readily extended to hadrons containing more than three quarks. Theorists have long been puzzled by the absence of hadrons with exotic quantum numbers (e.g. mesons with charge or strangeness two), which are by definition impossible to build from just $Q\bar{Q}$ or Q^3. We find[16] that some multiquark hadrons are light enough to be prominent but that the color and spin dependent forces dictate that these states will have ordinary quantum numbers (many combinations of $Q^2\bar{Q}^2$, for example, have quantum numbers also accessible to $Q\bar{Q}$). Truly exotic multiquark states are heavier and much harder to see. This raises the possibility that some $Q^2\bar{Q}^2$ hadrons have been misclassified as $Q\bar{Q}$ configurations. Indeed, we believe we have found an example of this.

The bag model predicts that the lightest $Q^2\bar{Q}^2$ multiplet is a flavor-SU(3) nonet (octet plus singlet) with spin-parity 0^+ and masses ranging from 650 to 1100 MeV. Just such a multiplet is known from $\pi\pi$ and $K\bar{K}$ scattering experiments. It has always been classified as a $Q\bar{Q}$ nonet, but the systematics of couplings and decays have never agreed well with that classification. In contrast they agree rather well with the $Q^2\bar{Q}^2$ assignment. Work has begun on higher configurations[17] ($Q^4\bar{Q}$ baryons, Q^6 dibaryons) and is limited only by the complexity of the group theory for such systems.

As a final example of bag phenomenology we wish to consider deformed states of two quarks and two antiquarks.[12,18] These might be expected to be prominent in baryon–antibaryon ($B\bar{B}$) annihilation. If a quark in the (right-moving) baryon annihilates as an antiquark in the (left-moving) antibaryon a deformed $Q^2\bar{Q}^2$ state remains. These states will have a $J = \alpha' M^2 + \alpha_0$ spectrum. The slope α' will be given by Eq. (2). It depends critically on the color charge of the QQ and $\bar{Q}\bar{Q}$ pairs left on ends of the spinning meson. In general the QQ system can be symmetric or antisymmetric in color, but in a baryon every quark pair is antisymmetric in color. Corresponding to this symmetry we find $C^2 = \frac{4}{3}$, identical to a single quark. We conclude that the Regge trajectories of deformed $Q^2\bar{Q}^2$ resonances seen in $B\bar{B}$ annihilation will have the same slope as ordinary baryons and mesons.

We can go further. The deformed $Q^2\bar{Q}^2$ states are orbital excitations of the spherical $Q^2\bar{Q}^2$ mesons we have just discussed. We know the masses of these states in the spherical approximation which allows us to calculate the intercepts α_0 for $Q^2\bar{Q}^2$ trajectories of various flavors. There is a complication: the spherical states are not eigenstates of quark and antiquark color. They are mixtures of the symmetric and antisymmetric QQ and $\bar{Q}\bar{Q}$ configurations. To find the needed intercepts we form states antisymmetric in QQ and $\bar{Q}\bar{Q}$ color and weight the masses accordingly. Now with both the slopes and intercepts in hand we can predict the masses, spins and flavors of all these $B\bar{B}$

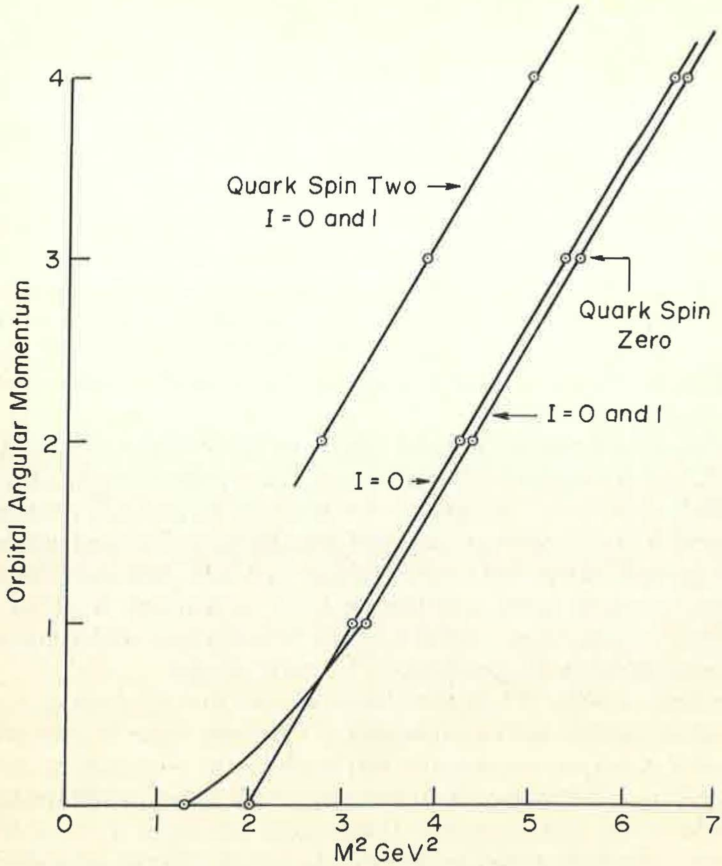

FIGURE 3 $Q^2\bar{Q}^2$ states which couple strongly to $P\bar{P}$ annihilation channels: (*a*) states with quark spin 0 and 2.

resonances. Their masses and quantum numbers are shown in Figure 3 *a* and *b*.

This analysis is by no means complete. The couplings of specific $B\bar{B}$ channels to specific resonances can be calculated with the aid of the Clebsch-Gordon coefficients for the symmetry groups involved. Notice must be taken of the partial wave in which the $B\bar{B}$ resonance appears: a state coupled only to the $B\bar{B}$ S-wave (for example) will not be strongly excited if it lies far above threshold where high partial waves dominate.

There are many striking predictions lurking in Figure 3*a* and *b*. First there should be resonances in $p\bar{p}$ annihilation at or just below threshold with quantum numbers $J^PI^G = 1^-0^-, 1^-1^+$. Our deformed $Q^2\bar{Q}^2$ states decay

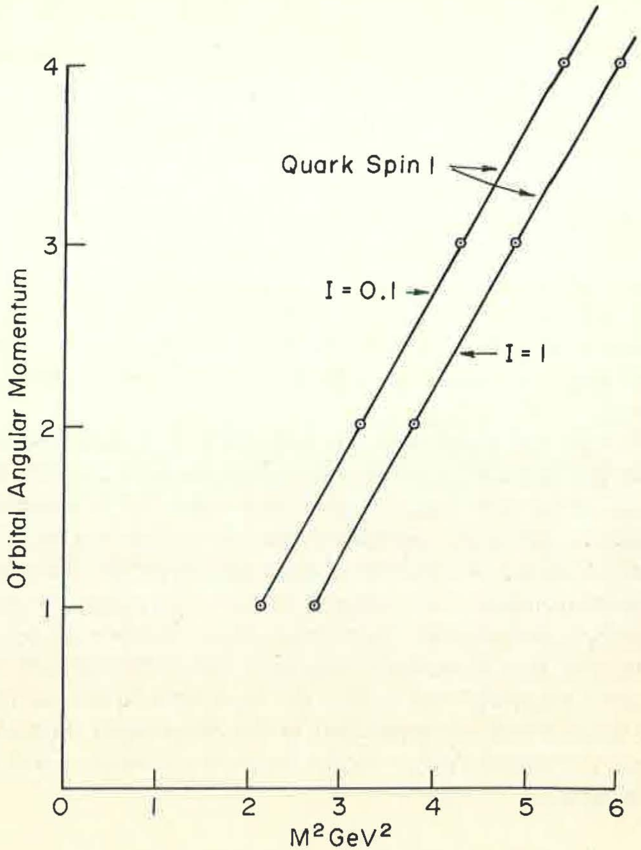

FIGURE 3 $Q^2\bar{Q}^2$ states which couple strongly to $P\bar{P}$ annihilation channels: (*b*) states with quark spin 1. Physical states, denoted by circles, occur when the Regge trajectories cross integer values of *L*. Warning: these are not Chew–Frautschi plots; the vertical axis is the orbital angular momentum. To obtain the total *J*, this must be coupled to the quark spin.

predominantly back into $B\bar{B}$ (there is an angular momentum barrier preventing a regrouping into mesons), so a state below $B\bar{B}$ threshold should be quite narrow. There are persistent reports of such states.[19] Secondly, there are low-lying resonances of high spin. Such states have recently been reported.[20] Finally, there are exotic states with rather low mass which unlike their S-wave siblings may be narrow (and therefore prominent) on account of their angular momentum. These are harder to see (they do not couple to $p\bar{p}$ or $n\bar{n}$) but might be important in virtual processes such as $p\bar{n} \rightarrow \pi^+$ (fast) + anything where observation of a fast π^+ selects $\bar{\Delta}^{--}$ exchange.[21]

Much remains to be done. Several important regimes of hadronic phenomena are not easily incorporated into a confinement based phenomenology such as ours. We lack a quantum description of bag fission and fusion making it difficult to discuss the dynamics of resonance formation, elastic form factors, and multiparticle processes at all energies. Our present understanding of the quantum mechanics of the bag's surface is meagre and too formal.[22] The dynamics of non-spherical excited hadrons of low angular momentum involves the surface in an essential way.[22,23] Without this we cannot provide much insight into the multitude of baryon and meson resonances with masses between 1 and 2 GeV. (One advantage of a fairly precise dynamical picture is that it allows us to see which phenomena are likely to be simple and which are, like chemistry, both more complex and less fundamental.)

Not everyone will agree with our prejudices. It is indicative of the odd state of the field that while there is wide agreement as to the basic ingredients for a theory of hadrons (quarks, unbroken color gauge interactions, confinement and so forth), few can agree to rules for translating these theoretical notions into concrete calculations of hadronic properties. The bag model is both phenomenological and rather precisely formulated. It provides an explicit example that the basic features on which there is wide agreement can be implemented in a framework consistent with relativity and locality. In this position it may help teach us how the ingredients in our emerging theory of strong interactions are manifested in the phenomena themselves. What the ultimate theoretical framework for the theory of hadrons will be remains to be determined.

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