WHY $N^*$’s ARE IMPORTANT

NATHAN ISGUR
Jefferson Lab, 12000 Jefferson Avenue, Newport News, U. S. A.

The study of $N^*$’s can provide us with critical insights into the nature of QCD in the confinement domain. The keys to progress in this domain are the identification of its important degrees of freedom and the effective forces between them. I report on the growing evidence in support of the flux tube model, and comment on the connection between this model and spontaneous chiral symmetry breaking, on the spin-dependence of the long-range confining potential, on the evidence for short-range one gluon exchange, on instantons, and on the one pion exchange model.

1 Why $N^*$’s?

There are three main reasons (one for each quark) why I believe that $N^*$’s deserve the special attention they get from the series of Workshops to which $N^*$2000 belongs:

• The first is that nucleons are the stuff of which our world is made. As such they must be at the center of any discussion of why the world we actually experience has the character it does. I am convinced that completing this chapter in the history of science will be one of the most interesting and fruitful areas of physics for at least the next thirty years.

• My second reason is that they are the simplest system in which the quintessentially nonabelian character of QCD is manifest. There are, after all, $N_c$ quarks in a proton because there are $N_c$ colors.

• The third reason is that history has taught us that, while relatively simple, baryons are sufficiently complex to reveal physics hidden from us in the mesons. There are many examples of this, but one famous example should suffice: Gell-Mann and Zweig were forced to the quarks by $3 \times 3 \times 3$ giving the octet and decuplet, while mesons admitted of many possible solutions.

2 What are the Key Issues?

I am convinced that the keys to a qualitative understanding of “strong QCD” are the same as in most other areas of physics: identifying the appropriate degrees of freedom and the effective forces between them.

Let me remind you why the basic degrees of freedom of QCD and the elementary gluon-mediated interactions between them are not useful starting points for understanding QCD in the confinement regime. It is not difficult
to demonstrate that the string tension $b$ of QCD must depend on the strong coupling constant $g$ according to

$$b \propto \exp \left[ -\frac{16\pi^2}{11g^2} \right]. \tag{1}$$

Almost all other interesting properties (including the $N^*$ masses) have this same $e^{-1/g^2}$ dependence. This essential singularity in $g^2$ means that the “Feynman diagrammar” is useless: plane wave quarks and gluons are not a useful starting point for low-energy, confinement-dominated physics.

Of course the full story is contained in the elementary QCD Lagrangian, and even in this situation it can be solved numerically using lattice QCD. However, as always in physics, we need to understand the relevant degrees of freedom and the effective forces between them for the phenomena under study if we wish to understand QCD and be able to use our understanding to envision new phenomena and to extend the power of QCD beyond our very limited ability to perform exact calculations using the lattice.

### 3 All Roads Lead to Valence Constituent Quarks and Flux Tubes as the Dominant Low Energy Degrees of Freedom

From the preceding discussion on the inappropriateness of the perturbative quark and gluon degrees of freedom for describing the phenomena of strong QCD, it will come as no surprise that foremost among the puzzles we face is in fact a glaring “degree of freedom” problem: the established low energy spectrum of QCD behaves as though it is built from the degrees of freedom of spin-$\frac{1}{2}$ fermions confined to a $q\bar{q}$ or $qqq$ system. Thus, for mesons of various flavors we seem to observe “quarkonia” spectra, while for the baryons we seem to observe the spectra of the two relative coordinates of three spin-$\frac{1}{2}$ degrees of freedom.

In particular, the empirical $N^*$ spectrum seems to demand that we use these degrees of freedom. From the perspective of the QCD Lagrangian, this is a highly nontrivial observation. Where are the expected gluonic states? Where are the extra $q\bar{q}$ (meson cloud) degrees of freedom that would naively be expected to be as important as the valence degrees of freedom?

Compelling insights into these questions have come from four different directions and converged on a simple picture in which the appropriate degrees of freedom for strong QCD are valence constituent quarks and flux tubes. This picture is commonly referred to as the flux tube model.
3.1 Spontaneous Chiral Symmetry Breaking in QCD

Long before QCD was discovered, it was appreciated that the underlying theory of the strong interactions should have a chiral symmetry which is spontaneously broken. When chiral symmetry is spontaneously broken, QCD must, by the Goldstone mechanism, generate the octet of light Goldstone mesons like the pion. *Note that the forces which generate spontaneous chiral symmetry breaking and these mesons are not mysterious new forces: they are just the full $q \bar{q}$ gluonic forces of QCD.* These forces give the vacuum nonperturbative structure, in particular a nonzero value for the vacuum expectation value of $\langle 0 | \bar{q}(x)q(x) | 0 \rangle$ which leads directly to a shift of the low energy effective quark mass from near zero to around the standard constituent quark mass of 330 MeV. Given this, spontaneous symmetry breaking must also be associated with some effective spin-dependent force which splits the $\pi$ from the $\rho$, the $K$ from the $K^*$, and the $\eta_{ss}$ from the $\phi$. The venerable NJL model provides a prototypical example of how all of this happens. Very roughly speaking, the dynamics (which spontaneously breaks chiral symmetry by creating the $q \bar{q}$ condensate and thereby the constituent quark mass) “conspire” in the pseudoscalar channel, with attractive forces compensating exactly (in the limit that the light quark Lagrangian masses are zero) for the two constituent quark masses. In QCD, the role of the unspecified NJL dynamics must of course be played by gluonic forces, though one can entertain various options for the nature of these interactions. I shall return to a much more extensive discussion of these points and the related problem of the mass of the $\eta'$ (the “$U_A(1)$ problem”) below.

3.2 The Large $N_c$ Limit of QCD

It is now widely appreciated that many of the observed features of the strong interactions can be understood in QCD within the $1/N_c$ expansion. Moreover, there is growing evidence from lattice QCD that the main qualitative features of QCD are independent of $N_c$. This limit also provides the only known field theoretic basis for the success of not only the valence quark model, but also of Regge phenomenology, the narrow resonance approximation, and many of the systematics of hadronic spectra and matrix elements. It can be shown in the large $N_c$ limit that hadron two-point functions are dominated by graphs in which the valence quark lines propagate from their point of creation to their point of annihilation without additional quark loops. A form of the OZI rule also emerges naturally. Large $N_c$ QCD thus presents a picture of narrow resonances interacting weakly with hadronic continua. In this picture each resonance is made of the valence quarks of the quark model and glue *summed to all orders in g.*
3.3 Quenched QCD

Quenched lattice QCD provides other new insights into QCD. In quenched QCD the lattice sums amplitudes over all time histories in which no $q\bar{q}$ loops are present. It thus gives quantitative results from an approximation with many elements in common with the large $N_c$ limit. One of the most remarkable features of these calculations is that despite what would seem to be a drastic approximation, they provide a reasonably good description of low energy phenomenology. Indeed, for various intermediate quantities like the QCD string tension they provide very good approximations to full QCD results with the true lattice coupling constant replaced by an effective one. In quenched QCD, as in the large $N_c$ limit, two point functions thus seem to be well-approximated by their valence quark content.

3.4 Some Comments on the Relation between the Lessons of the Large $N_c$ Limit and of Quenched QCD

In comparing the large $N_c$ limit and quenched lattice QCD we note that:

- In both pictures all resonances have only valence quarks, but they have an unlimited number of gluons. Thus they support valence models for mesons and baryons, but not for glueballs or for the gluonic content of mesons and baryons.
- While both pictures tell us that hadrons are dominated by their valence quark structure, the valence quark propagators are not nonrelativistic propagators. I will elaborate on this point below.
- The large $N_c$ and quenched approximations are not identical. For example, the $NN$ interaction is a $1/N_c$ effect, but it is not apparently suppressed in the quenched approximation.

3.5 The Heavy Quark Limit

The fourth perspective from which there is support for the same picture is the heavy quark limit. While this limit has the weakest theoretical connections to the light quark world, it has powerful phenomenological connections: see Fig. [1]. We see from this picture that in mesons containing a single heavy quark, $\Delta E_{\text{orbital}}$ (the gap between, for example, the $J^{PC} = 1^{--}$ and $2^{++}$ states), is approximately independent of $m_Q$ while $\Delta E_{\text{hyperfine}}$ varies like $m_Q^{-1}$, both as expected in the heavy quark limit.

Recall that in the heavy quark limit a hadronic two-point function is dominated by a single valence $Q$ plus its associated “brown muck”, with neither $Q\bar{Q}$ loops nor $QZ$-graphs. The fact that heavy-quark-like behaviour persists all the
way down to light quark masses suggests that light quarks, like heavy quarks, behave like single valence quarks and thus by extension that the “brown muck” behaves like a single valence antiquark.

4 Comments on the Degree of Freedom Problem

4.1 Where is the Glue and Why Represent it by a Flux Tube?

Fig. 2 shows that heavy quarkonium (QQ) behaviour apparently persists into the light meson spectrum. This is surprising since the former spectrum depends on the decoupling of gluonic excitations (as opposed to glue) from the spectrum. The flux tube model provides an attractive explanation of the adiabatic approximation which leads to this decoupling and which is now receiving
strong support from lattice QCD studies. Most importantly, the strongly collimated chromoelectric flux lines of the flux tube model have been seen on the lattice, as have the predicted first excited adiabatic surfaces of the flux tube with an energy gap $\delta V(r) = \pi/r$ above the quarkonium potential and doubly-degenerate phonon quantum numbers. This basically requires that the $J^{PC}$ exotic hybrid mesons predicted fifteen years ago exist.

The flux tube model thus offers a possible explanation for one of the most puzzling apparent inconsistencies between the valence quark model and QCD. Moreover, as discussed above, in the large $N_c$ limit of QCD hadrons do indeed consist of just their valence quarks and the glue between them. Thus the flux tube model may be viewed as a realization of QCD in the large $N_c$ limit which is in addition consistent with insights into strong QCD which have emerged from spontaneous chiral symmetry breaking, the large $N_c$ limit, from quenched lattice QCD, and from heavy quark theory. It corresponds to a rigorous and in principle systematic expansion of QCD around the strong coupling limit rather than the weak coupling limit.

We should remember that many of these theoretical developments have their historical roots in data. In particular, the leading Regge trajectories ($1^-, 2^+, 3^-, 4^+, 5^-, 6^+$, etc. in mesons, $1/2^+, 3/2^-, 5/2^+, 7/2^-, 9/2^+, 11/2^-$ etc. in the $N^*$'s, and $3/2^+, 5/2^-, 7/2^+, 9/2^-, 11/2^+, 13/2^-$ etc. in the $\Delta^*$'s) led Nambu to postulate in the early 1970's that QCD was a relativistic string theory with the mass of states at high excitation dominated by the mass of the string. In QCD, Nambu's prophetic picture is realized via quarks confined to infinite towers of (narrow) bound states by a string-like flux tube.

Of course, more data is required to confirm that the flux tube model is indeed on the right path and to define the next steps down this path. This is the goal of Hall D of the 12 GeV Upgrade project, but there are very important implications as well for the current $N^*$ program: it has recently become clear that this picture also requires that the cross sections for hybrid $N^*$ electroproduction be large.

### 4.2 Constituent Quarks are Not Nonrelativistic Quarks

Discussions of the constituent quark model are sometimes confused by a failure to recognize that the nonrelativistic quark model provides a very convenient qualitative tool for understanding many of the features of strong QCD, but should not be interpreted literally. In particular, all of the QCD-based pictures described above have propagating valence quarks with contributions from not only a positive energy quark propagator, but also from “$Z$-graphs”. (A “$Z$-graph” is a time-ordered graph in which the interactions first produce a pair
and then annihilate the antiparticle of the produced pair against the original propagating particle). Cutting through a two-point function at a fixed time therefore would in general reveal not only the valence quarks but also a large \( q\bar{q} \) sea. This does not seem to correspond to the usual valence approximation. Consider, however, the Dirac equation for a single light quark interacting with a static color source. This equation represents the sum of a set of Feynman graphs which also include Z-graphs, but the effects of those graphs is captured in the lower components of the single particle Dirac spinor. I.e., such Z-graphs correspond to relativistic corrections to the quark model. That such corrections are important has been known for a long time. For us the relevant point is that while such effects have important quantitative effects on quark model predictions (e.g., they are commonly held to be responsible for much of the required reduction of the nonrelativistic quark model prediction that \( g_A = 5/3 \) in neutron beta decay), they do not qualitatively change the single-particle nature of the spectrum of the quark of our example, nor would they necessarily qualitatively change the spectrum of \( q\bar{q} \) or \( qq\bar{q} \) systems.

4.3 The Valence Approximation is an Approximation

The valence approximation can only be taken as a starting point for a systematic treatment of strong QCD, since we know that \( q\bar{q} \) pair creation plays an important role in many phenomena. In particular,

- While particle systematics, including the linearly rising Regge trajectories, require a narrow resonance approximation, we know that the resonances have finite but nonzero widths \( \Gamma << M \).
- In analyzing data, we are clearly faced with the fact that the \( N^* \)'s are immersed in a baryon-meson continuum, and that the interplay between the resonance spectra and these continua are important.
- Even amongst the more sophisticated analyses of the \( N^* \) spectra, there has historically been a great deal of confusion about the relationship between the predictions of the valence quark model for spectra and the narrow resonance approximation. The large \( N_c \) limit has begun to clarify matters: while the quark model spectrum is a narrow resonance approximation, it is a mixture of leading effects in the \( N_c \) expansion and \( 1/N_c \)-suppressed effects. Thus the connection of quark model predictions to the data is quite subtle, and cannot be treated by simply considering the interaction of narrow resonances with the continua via meson loop graphs. For example, many of the effects of meson loop graphs are already subsumed into the quark model’s string tension [10].
- The large \( N_c \) limit tells us that the valence approximation will be good to an accuracy of order \( 1/N_c \) except for \( SU(3) \) singlet channels where the valence
predictions will be multiplied by a correction factor of not \((1 + constant \ N_c^{-1})\), but rather \((1 + constant \ N_f N_c^{-1})\), where \(N_f\) is the number of light flavors. Thus, for example, the proton spin crisis may be attributed to the fact that the valence spin contribution is reduced coherently by relatively small effects from each of \(u\bar{u}\), \(d\bar{d}\), and \(s\bar{s}\) pairs from the quark-antiquark sea.

While the above points make it clear that the quark-antiquark sea must be added to the valence quarks to get an accurate picture of hadrons and their interactions, naively attempting to add \(q\bar{q}\) pair creation to the valence quark model leads to a number of very serious problems. These problems and potential solutions to them have been extensively discussed in a series of papers on “unquenching” the quark model. The most important lesson of these studies is that low energy truncations of the tower of meson loops corrections to the valence quark states (which generate the \(q\bar{q}\) sea; see Fig. 3) are usually misleading.

5 What Are the Effective Forces between the Low Energy Degrees of Freedom?

There are many contending pictures for the origin of the effective forces between the low energy degrees of freedom of strong QCD. I have observed, perhaps as a result of the variety of proposals, some confusion and unnecessary contentiousness over this key issue. In some cases there are real issues to be resolved, but not always.
5.1 The Confining Interaction

The confining interaction is automatically provided by the flux tube, and in the very low energy regime where the adiabatic approximation is valid, the flux tube degrees of freedom can be replaced by an adiabatic effective potential, at least for heavy enough quarks.

The central part of this potential is linear at large distances. This fact should not be confused with discussions of the basis space for solutions to spectroscopic problems, which is an issue of accuracy and not principle. For example, in the Isgur-Karl model\[11\], a harmonic basis is used to solve the linear potential problem. A harmonic potential is not assumed.

There have historically been many many proposals for the spin dependence of the confining potential; very recently, there has been much discussion of its spin-orbit components. The basic issues are discussed in the original Isgur-Karl papers, since in that model all spin dependence is assumed to arise from one gluon exchange (OGE) except for the unavoidable purely kinematic spin-orbit force from Thomas precession in the confining potential. The key point is that the central confining potential, whose eigenstates almost every modeller takes as the basis for spin-dependent perturbations of resonance spectra, must produce very strong spin-orbit forces through Thomas precession. To analyze the implications of this fact it is best to examine the meson sector. The mesons appear to have a “spin-orbit problem” as can be seen by examining the first band of positive parity excited mesons: the four P-wave mesons of every flavor are nearly degenerate. Most of the observed small non-degeneracies are due to hyperfine interactions, but the spin-orbit matrix elements can be extracted. For example, by taking the isovector meson combination \(\frac{1}{2}m_1 - \frac{1}{2}m_2 - \frac{1}{2}m_0\) one can isolate their spin-orbit matrix element of \(-3 \pm 20\) MeV. This matrix element is much smaller than would be obtained from either the standard short-distance OGE interaction or from the Thomas precession term alone. However, the large “normal” OGE spin-orbit matrix element tends to cancel the strong “inverted” spin-orbit matrix element from Thomas precession. Typical fits to the data, including the heavy quarkonia, give light quarkonium OGE and Thomas precession spin-orbit matrix elements of about \(\pm200\) MeV, respectively.

The physics behind this cancellation has received support recently from the theory of \(Q\bar{Q}\) systems, where both analytic techniques and numerical studies using lattice QCD have shown that the confining forces are spin-independent apart from the inevitable spin-orbit pseudoforce due to Thomas precession. Moreover, the data on charmonia require the inverted spin-orbit matrix element from Thomas precession in the confining potential to cancel part of the strength...
of their rigorously required positive OGE matrix element. If the charm quark were more massive, its low-lying spectrum would be totally dominated by one gluon exchange. Indeed, one observes that the $\Upsilon$ system is closer to this ideal, as expected. Conversely, as one moves from $c\bar{c}$ to lighter quarks, the $\ell = 1$ wave functions move farther out into the confining potential and the relative strength of the Thomas precession term grows. It is thus very natural to expect a strong cancellation in light quark systems, though from this perspective the observed nearly perfect cancellation must be viewed as accidental.

Given this rather satisfactory solution of the “meson spin-orbit problem”, it would be very surprising if the analogous “baryon spin-orbit problem” didn’t have an analogous solution, so this assumption was built into the Isgur-Karl model. I will return to a discussion of these matters below.

To summarize: there is very strong evidence that confinement may be approximated at low energy by replacing the flux tube with an effective potential with a linear central component. In the leading correction to this nonrelativistic approximation, the potential should be supplemented with only the minimally required Thomas precession term (which acting alone would produce strongly inverted spin-orbit multiplets) and with spin-independent but momentum-dependent terms which take into account the impact of the orbital angular momentum of the massive rotating flux tube on the spectrum of states.

5.2 The Short-Distance Interactions

One Gluon Exchange

The preceding discussion leaves little room for doubt about the confining interaction and its structure, at least to leading order. There is, however, room for doubt about the nature of the short-distance interactions, and in particular whether the standard OGE model for these interactions is correct. While healthy, this discussion has also sometimes been quite confused, and I want to take this opportunity to try to sort out the issues. Let me acknowledge from the beginning that I am a partisan of the OGE school, and so warn you that my organization of the issues (but not my facts!) may be prejudicial.

The discussion is once again best started in the simpler meson sector. Figure 2 shows what we know about the evolution of quarkonium spectroscopy as a function of the quark masses. In heavy quarkonia ($b\bar{b}$ and $c\bar{c}$) we know that hyperfine interactions are generated by OGE perturbations of wave functions which are solutions of the Coulomb plus linear potential problem. I find it difficult to look at this diagram and not see a smooth evolution of the wavefunction (characterized by the smooth evolution of the orbital excitation energy) convoluted with the predicted $1/m_Q^2$ strength of the OGE hyperfine interactions.
interaction. OGE-based quark models turn this qualitative impression into quite good quantitative predictions.

This same conclusion can be reached by approaching the light quarkonia from another angle. Figure 1 shows the evolution of heavy-light meson hyperfine interactions from the heavy quark limit to the same isovector quarkonia. In this case we know that in the heavy quark limit the hyperfine interaction is given by the matrix element of the operator $\vec{\sigma}_Q \cdot \vec{B} / m_Q$. The conclusion that heavy-light meson hyperfine interactions are controlled by OGE is thus also supported by the striking $1/m_Q$ behaviour of the ground state splittings in Fig. 1 as $m_Q$ is decreased from $m_b$ to $m_c$ (in which systems it may be rigorously applied) on down to $m_s$ and then to $m_d$: it certainly appears that for all quark masses the quark $Q$ interacts with $\vec{B}$ through its chromomagnetic moment $\vec{\sigma}_Q / 2m_Q$, as would be characteristic of the OGE mechanism.

Let us now turn to baryons. As shown in Fig. 5.2 the baryon analog of Fig. 1, experiment provides further strong evidence in support of the dominance of OGE in the baryons. It is clear from this Figure that in the heavy quark limit the splittings are once again behaving like $1/m_Q$ as demanded by heavy quark theory, and once again it is difficult to look at this diagram and not see a smooth evolution of this $1/m_Q$ behaviour from $m_b$ to $m_c$ to $m_s$ to $m_d$, where by SU(3) symmetry $\Sigma^*_{SU(3)} - \Lambda_{SU(3)} = \Delta - N$, the splitting relevant to $N^*$s. Indeed, using standard constituent quark masses ($m_d = m_u = m_s = 0.33$ GeV, $m_c = 1.82$ GeV, $m_b = 5.20$ GeV), the OGE mechanism with its natural $1/m_Q$ behaviour quantitatively describes these spectra. Thus one observes in both mesons and baryons the remnants of heavy quark spectroscopy in light quark systems. For example, in Fig. 4 the $D^* - D$ heavy quark spin multiplet is naturally identified with the $K^* - K$ multiplet, i.e., the basic

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Figure 4: Ground state baryon hyperfine splittings in heavy-light systems as a function of the mass $m_Q$ of the heavy quark. The spectra on the far right are the $m_Q \to \infty$ limits of heavy quark symmetry. The $\Sigma^*_Q - \Lambda_Q$ splitting and the positions of $\Sigma^*_b$ and $\Sigma_b$ are estimates from the quark model (for which there is now experimental support); all other masses are from experiment. The spectra are shown to scale and may conveniently be calibrated with the $\Sigma_c - \Lambda_c$ splitting of 169 MeV.
degrees of freedom seen in the spectrum are the same, and from the observed
$D^*-D$ splitting and the $1/m_Q$ heavy quark scaling law alone one expects a
$K^*-K$ splitting of 460 MeV, quite close to the actual splitting of 400 MeV
(including wave function effects improves the prediction). One might try to
escape this conclusion by arguing that between $m_c$ and $m_s$ the OGE-driven
$1/m_Q$ mechanism turns off and some other mechanism which imitates it turns
on. I do not know how to rule out this baroque possibility except on a case-
by-case basis.

The excited charmed baryon sector has recently provided further strong
evidence for the dominance of the OGE mechanism in baryons. From an
analysis of their decay patterns, it seems very likely that if the $\Lambda(1405)^{-\frac{1}{2}}$
and $\Lambda(1520)^{-\frac{3}{2}}$ are valence-quark-model-type $uds$ states, then they are the
expected spin-orbit partners of the quark model. Heavy quark symmetry demands
that in the heavy-light isospin zero $\Lambda_Q$ sector, the $\Lambda_Q^{-\frac{1}{2}}$ and $\Lambda_Q^{-\frac{3}{2}}$
be degenerate as $m_Q \to \infty$ and that their splitting open up like $1/m_Q$ as $m_Q$
decreases. The $\Lambda_c(2594)^{-\frac{3}{2}}$ and $\Lambda_c(2627)^{-\frac{1}{2}}$ appear to be such a nearly
degenerate pair of states in the charmed baryon sector. Their center-of-gravity
$\frac{1}{2}m_{\Lambda_Q^{-\frac{3}{2}}} + \frac{1}{2}m_{\Lambda_Q^{-\frac{1}{2}}}$ is 330 MeV above the $\Lambda_c(2285)$. This is to be compared
with the center-of-gravity of the $\Lambda(1520)$ and $\Lambda(1405)$ which lies 365 MeV
above the $\Lambda(1115)$, in accord with the expectation that the orbital excitation
energy of the negative parity excitations of $\Lambda_Q$ will be a slowly increasing
function of $1/m_Q$. (Recall Fig. 1: the same is true in mesons.) This alone
suggests that the strange quark analogs of the heavy quark spin multiplet
($\Lambda_c(2627)^{-\frac{3}{2}}, \Lambda_c(2594)^{-\frac{1}{2}}$) should exist just around the mass of the $\Lambda(1520)^{-\frac{3}{2}}$
and $\Lambda(1405)^{-\frac{1}{2}}$. Moreover, using the predicted $1/m_Q$ behaviour of the mul-
tiplet splitting, one would expect a splitting of $\simeq 110$ MeV in the $\Lambda$ sector
compared to the observed splitting of 115 MeV. It is thus difficult to avoid
identifying ($\Lambda(1520)^{-\frac{3}{2}}, \Lambda(1405)^{-\frac{1}{2}}$) as the strange quark analogs of a heavy
quark spin multiplet, and to avoid concluding that the $1/m_Q$ evolution of the
OGE mechanism is responsible for its splitting.

Of course in both Figs. 1 and 2 the extrapolation of the nonrelativistic
OGE mechanism to light constituent quark masses cannot be taken literally.
In $Q\bar{Q}$ systems, OGE is treated by using a nonrelativistic reduction to define
a low energy effective potential; such a treatment would not be accurate for
light quarks whose masses are of the order of the QCD scale. What quenched
lattice QCD and the large $N_c$ limit tell us is that mesons and baryons are
dominated by their valence quark structure, and what Figs. 1 and 2 strongly
suggest is that a smooth extension of the $Q\bar{Q}$ valence quark interaction is in
operation for all quark masses. For light quark masses this extended OGE
interaction will include many new effects, including not only straightforward relativistic corrections, but also those like Z-graphs which arise because the light valence quarks are embedded in a relativistic field theory and so have instantaneous projections into subspaces with additional $q\bar{q}$ pairs, and effects of vacuum structure that are not important at short distances.

This brings us to one of the most confused elements of the discussions of the quark model, that it is somehow inconsistent with chiral dynamics. This is certainly not true in principle. As discussed in Section 3.1, spontaneous chiral symmetry breaking generates the constituent quark mass, so this element of quark models is required by chiral dynamics. Moreover, the pion is not a magical massless Goldstone boson which appears out of nowhere and is thus outside of the domain of quark models, in stark contrast to the discussion (and the experimental data!) above. A massless pion automatically appears in both quenched lattice QCD and the large $N_c$ limit as the mass of the lightest state of a propagating quark-antiquark pair. I.e., even in the valence quark approximation, a massless pion is automatic!

**Instantons**

There is another possibility for the short-distance interactions of light quarks which has deservedly been receiving a great deal of attention recently: the instanton liquid model\textsuperscript{12}. This picture is inconsistent with any model based on a Feynman-diagrammatic treatment of QCD (even all-orders ones like the large $N_c$ limit).

Let me begin with a little history. For many years, QCD was plagued by the “$U_A(1)$ problem”: the equations of motion of QCD imply that spontaneous chiral symmetry breaking leads to nine and not just eight Goldstone bosons, a conclusion that is apparently inconsistent with the large mass of the $\eta'$ meson at about 1 GeV. However, the $U_A(1)$ current is anomalous, and by the late 1970’s it was understood through the study of instantons\textsuperscript{13,14} that the anomaly leads to a nonconservation of the $U_A(1)$ charge and thereby to the evasion of Goldstone’s theorem when chiral symmetry is spontaneously broken. The connection between the quark model and instantons was discussed by Witten\textsuperscript{15}, Veneziano\textsuperscript{16} and others, who explored more generally the conflict between instantons and the large $N_c$ expansion. It is certainly a mainstream, viable, and widely held belief that the $U_A(1)$ problem is solved through instanton contributions to the axial anomaly. However, as emphasized by Witten, instantons vanish like $e^{-N_c}$ and so do not appear in the large $N_c$ expansion. “Insofar as [instantons play] a significant role in the strong interactions, the large $N_c$ expansion must be bad. It is necessary to choose between the two.”
What is the essential origin of this conflict? Unlike QED, the nonabelian structure of QCD admits many possible classical vacua characterized by a “winding number” which classifies the topology of the nonabelian vector potential, but all of which have the zero color electric and magnetic fields and thus zero energy. (In QED there appear to be many choices for a vector potential which have zero electric and magnetic fields and zero energy, but they are all related by gauge transformations and are thus equivalent. A gauge transformation cannot change winding number (it is a topological property), and thus the infinite number of classical QCD vacua are distinct.) The instanton picture begins from the premise that the correct vacuum around which one should compute quantum corrections (i.e., on which a normal Feynman diagram expansion should be made) is the one resulting from tunneling between these energy-equivalent classical vacua (much like the tunneling between two potential wells separated by a classically forbidden region). When the vacuum tunnels between two of its possible states (characterized by different winding numbers), it does so quickly (by the uncertainty principle); such an event is called an instanton. The tunneling process thus dynamically generates the true vacuum of QCD as a superposition of classical vacua of various winding numbers. It turns out that this superposition can be characterized by single parameter called $\theta$.

The physics of $\theta$ is a very important story all unto itself. If $\theta$ is nonzero as naively expected from the previous discussion (it is after all a strong interaction parameter!), then the strong interactions would violate $CP$ invariance. Experimental constraints tell us that $\theta$ is roughly ten orders of magnitude smaller than we might have expected! This is known as the “strong CP problem”. Axions were invented to try to deal with this problem, but there is no established solution. At the moment, $\theta$ must be listed as one of the 17 basic parameters of the standard model (though its value is measured to be negligibly small).

Let us leave this interesting digression to return to our main story. Since the usual Feynman diagrammatic expansion of QCD is built on the perturbative vacuum with zero vector potential, it obviously has no knowledge of other vacua or of the instantons that connect them, so the effects of such instanton events have to be added “by hand” as a supplementary effective interaction. For us the relevant point is that this effective interaction includes an interaction between the quarks known as the ’t Hooft interaction. As noted by Witten, if instantons are important, then the large $N_c$ expansion would fail, since it assumes that all-orders properties of the QCD Feynman diagrammatic expansion are properties of QCD.

What would the impact be on the effective interaction we are trying to determine? The instantons (which are still a purely gluonic effect!) would
require the addition of the 't Hooft interaction to the quark model. This interesting possibility would indicate that the Coulomb plus linear picture, when extrapolated to light quarks, will fail quantitatively and that light quark spin-dependence, while still of gluonic origin in both mesons and baryons, will in general be a mixture in some proportions of OGE-type forces and the 't Hooft interaction. However, as Witten argued long ago, given the $U_A(1)$ anomaly and confinement, the large $N_c$ limit is also able to explain the $\eta'$ mass without instantons: large confinement-generated gluonic fluctuations in the vacuum also couple to the $U_A(1)$ anomaly and change the winding number. No matter which solution to the $\eta'$ mass is correct, there is still $\theta$-dependence in QCD and thus a strong $CP$ problem. But for strong interaction physics, understanding whether instantons are or are not important remains one of the most important outstanding issues.

The Pion Exchange Model for Interquark Forces

An alternative to the OGE model for the short-distance forces between quarks is the pion exchange model (OPE model), which posits that these forces are dominated by one pion exchange between the constituent quarks. More generally, the model assumes that the short-distance forces between the quarks in a baryon is mediated by the exchange of the octet of nearly massless Goldstone bosons generated by the spontaneous breaking of chiral symmetry. It is argued that the quark model requires such effects to be consistent the spontaneous breaking of chiral symmetry, and that the replacement of OGE by OPE solves many of the problems of the quark model for baryons. I have explained above why there is no need to invoke a special mechanism like OPE to make the OGE model consistent with the spontaneous breaking of with chiral symmetry. I have also presented my objections to the OPE model in conference talks. More recently, I have published these criticisms, taking into account recent attempts to overcome some of the problems of the original OPE model by extending it to include heavier meson exchanges and other effects. I briefly summarize the main points here:

- One of the original motivations for the OPE model was that it could solve the baryon spin-orbit problem since OPE, in contrast to OGE, produces no spin-orbit force. From the discussion of Section 5.1 on the inverted spin-orbit potential generated by confinement, one can see that this motivation was based on a misunderstanding of the nature of the problem, which is to arrange a sufficiently precise cancellation between the inevitable Thomas precession term generated by confinement and short-range dynamical spin-orbit forces. Recent elaborations of the original OPE model to include the exchange between quarks
of higher mass mesons can produce such dynamical spin-orbit forces, but in these circumstances this model clearly also has a “spin-orbit problem”, i.e., no advantage has been achieved. On the contrary, since higher mass mesons will have only short-range effects, it may be very difficult to arrange a cancellation of spin-orbit forces in both the $L = 1$ and $L = 2$ baryons as required by the data.

- Another of the original claims of the OPE model is that it produces a superior description of baryon masses. I will comment briefly on this somewhat complicated claim below, but in Ref. 19, I focus on a more straightforward matter: in a complex system like the baryon resonances, predicting the spectrum of states is not a very stringent test of a model. The prototypical example (and the first case in $N^*$ spectroscopy where this issue arises) is the two $N^{3/2}_-^-$ states found in the 1500-1700 MeV range. In any reasonable valence quark model, two $N^{3/2}_-^-$ states will be predicted in this mass range since totally antisymmetric states with overall angular momentum $\frac{3}{2}$ can be formed by coupling either quark spin $\frac{3}{2}$ or quark spin $\frac{1}{2}$ with $\ell = 1$. In the general case such a model will therefore give

\begin{align}
|N^{3/2}_-^-(upper)\rangle &= \cos\theta_{3/2} - |4P_N\rangle + \sin\theta_{3/2} - |2P_N\rangle \\
|N^{3/2}_-^-(lower)\rangle &= \cos\theta_{3/2} - |2P_N\rangle - \sin\theta_{3/2} - |4P_N\rangle
\end{align}

in an obvious notation. Since the masses of these resonances are only known (and currently interpretable) to roughly 50 MeV, it is not very difficult to arrange for a model to give a satisfactory description of the $N^{3/2}_-^-$ spectrum. However, among models which “perfectly” describe the spectrum there is still a continuous infinity of predictions for the internal composition of these two states since all values of $\theta_{3/2}$ from 0 to $\pi$ correspond to distinct states. The correct mixing angle has been known experimentally for 25 years. It is correctly predicted by the OGE model, but not by the OPE model.

This is the simplest example of examining the complete spectroscopy (and not just the masses) of baryons. In the positive parity states, an analysis which fails to examine the internal composition of states is nearly useless. For example, the valence quark model predicts five $N^{3/2}_+^+$ states, but only one is known. It would be an unlucky modeller who couldn’t identify one of their five predicted states with the observed state and claim success!

- There are two ways in which the OPE model is a disaster for mesons: since it doesn’t produce spin-dependent interactions in mesons (see below), it requires that different mechanisms are dominant in these two systems despite
the great similarity of their spin-independent spectra (and the implied similarity of their internal structure), and it predicts the existence of effects in mesons which are ruled out experimentally.

We know from quenched lattice QCD that the bulk of both meson and baryon hyperfine interactions occur in the quenched approximation (i.e., in the absence of closed $q\bar{q}$ loops). There are nevertheless Z-graph-induced meson exchanges between quarks that arise in this approximation that could in principle be the origin of those posited in OPE-type models. However, such meson exchange can only operate between two quarks and not between a quark and an antiquark, with the unsatisfactory consequence just mentioned.

The second problem in mesons is that the OPE model would produce unacceptably large violations of the OZI rule. While meson exchange cannot produce interactions between the quark and antiquark in a meson, it would lead to $q\bar{q}$ annihilation which will badly violate the OZI rule.

• The OPE model is not consistent with the heavy quark symmetry of QCD for $m_Q \to \infty$. Recall the discussion associated with Fig. 4. It is clear from these spectra that in the heavy quark limit the OPE mechanism is not dominant: exchange of the heavy pseudoscalar meson $P_Q$ would produce a hyperfine interaction that scales with heavy quark mass like $1/m_Q^2$, while for heavy-light baryons the splittings are behaving like $1/m_Q$ as in the heavy-light mesons.

Ref. 19 does not discuss the claim that the OPE models produce a better spectrum than the Isgur-Karl model. I have explained that without using mixing angle information, spectroscopy is a rather blunt tool. It is for this reason that I did not engage this argument in my paper, but I want to take this opportunity to clear up a misconception about the spectroscopy of the Isgur-Karl model that has been labelled the “level ordering problem”. Begin in the $S = 0$ states where we have the $N$, $N^*(1535)$, and $N^*(1440)$. 

In the harmonic limit, these states would be equally spaced in the order they are listed instead of the order in which they are seen. Despite the impression that the OPE papers might give, the observed pattern is automatically realized in the Isgur-Karl model. The inversion is due to two effects.

The first is that the true confining potential will rigorously break a harmonic oscillator spectrum (to leading order in the anharmonicity, which may be minimized by choice of the unphysical oscillator constant of the basis space) into the following pattern in terms of $SU(6)$ multiplets (note that I have arranged these equations so that the lightest state is at the bottom of the array):

\[
E(20, 1+) = E_0 + 2\Omega \quad (4) \\
E(70, 2+) = E_0 + 2\Omega - 1/5\Delta \quad (5)
\]
where the standard Coulomb plus linear potential gives $\Delta > 0$, and indeed gives $\Delta$ comparable to $\Omega$. Here $E_0$, $\Omega$, and $\Delta$ are three constants which are rigorously related by theory to the usual three parameters of the interquark potential: an overall constant, $\alpha_s$, and the string tension $b$. Since the $[56,0^+]$, $[70,1^-]$, and $[56,2^+]$ supermultiplet states are very well known, $E_0$, $\Omega$, and $\Delta$ are known, and one can confirm \textit{experimentally} that $\Delta$ is comparable to $\Omega$, i.e., that the $[56',0^+]$ supermultiplet should be close to the $[70,1^-]$.

The second effect is OGE. In the Isgur-Karl model, spin-spin interactions from OGE dominate the departures of the $S=0$ masses from the $SU(6)$ limit and have a very simple structure. Using the zeroth order harmonic oscillator wavefunctions (I have once again arranged these equations so that the lightest state is at the bottom of the array):

\[
\Delta' = E_0 + \Omega - \Delta + 5/8d \quad (11)
\]
\[
N' = E_0 + \Omega - \Delta - 5/8d \quad (12)
\]
\[
4^N* = 2^\Delta* = E_0 + \Omega + 1/4d \quad (13)
\]
\[
2^N* = E_0 + \Omega - 1/4d \quad (14)
\]
\[
\Delta = E_0 + 1/2d \quad (15)
\]
\[
N = E_0 - 1/2d . \quad (16)
\]

Thus since from $\Delta - N$ we know $d = 300$ MeV, with $\Omega$ approximately equal to $\Delta$, one expects the $N^*(1440)$ to be roughly $3/8d$ or about 110 MeV below the $N^*(1535)$ as observed.

While a very satisfactory starting point, this simple story certainly has some well-known problems. With $\Omega$ approximately equal to $\Delta$, the use of perturbation theory in the anharmonicity is suspect and indeed realistic potentials all have $E(56',0^+) > E(70,1^-)$. That $\Delta' - N'$ is larger than $\Delta - N$ is also an artifact of the harmonic limit. Thus more complete analyses of the OGE model tend to leave the $N^*(1440)$ too high in mass. I will comment on this discrepancy below, but for now note that these two states are naturally very close to each other, and are certainly \textit{not} expected to be split by a harmonic-like gap of $\Omega$.

Now consider the $S = -1$ sector, which is used as the key to defining the “level ordering problem” of the Isgur-Karl model. The OPE papers note
that in this case the lightest states of the $N = 0$, 1, and 2 bands, namely the $\Lambda$, $\Lambda^*(1405)$, and $\Lambda^*(1600)$, have the “normal” level ordering and claim that this reversal is inconsistent with the OGE picture. The argument given is incorrect. In addition to the SU(3) partners $^2\Lambda_8$ and $^4\Lambda_8$ of the octet $N^*$ states $^2N$ and $^4N$, the $N = 1$ band contains a new SU(3) singlet state $^2\Lambda_1$ that has no $N^*$ counterpart. The OGE spin-spin interaction in the SU(3) limit predicts that the $^4\Lambda_8$, $^2\Lambda_8$, and $^2\Lambda_1$ states will have hyperfine interactions of $+1/4d$, $-1/4d$, and $-3/4d$, respectively: the new singlet state $^2\Lambda_1$ is the most highly stabilized state in the entire negative parity spectrum. It automatically drops substantially relative to the octet states, and thus OGE naturally gives a low-lying negative parity state.

Thus the claim that the Isgur-Karl model (with its OGE) cannot explain the observed level orderings, and especially the reversal of the $\Lambda$’s relative to the $N^*$’s, is incorrect. A more productive discussion might focus on the accuracy of the predicted spectrum, since neither the $N^*(1440)$ nor the $\Lambda^*(1405)$ mass is very accurately predicted. As we have seen, a lot of physics is missing from the valence quark model. In particular, I believe that one should not expect flux tube model spectroscopy to work to better than $50\text{ MeV}$ until we have learned how to treat the couplings to continua correctly.

6 Conclusions

The study of $N^*$’s can provide us with critical insights into the nature of QCD in the confinement domain. It is becoming clear that, in this domain, constituent quarks and flux tubes are the appropriate low-energy degrees of freedom. However, there are still many very important questions to be resolved on the nature of the short-distance interactions of the quarks. I expect that $N^*$’s, as has always been the case historically, will in the future play a pivotal role in answering these questions.

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