CEBAF-TH-93-31

CEBAF
The Continuous Electron Beam Accelerator Facility
Theory Group Preprint Series

Additional copies are available from the authors.

The Southeastern Universities Research Association (SURA) operates the Continuous Electron Beam Accelerator Facility for the United States Department of Energy under contract DE-AC05-84ER40150

HADRON SPECTROSCOPY AND STRUCTURE†

Nathan Isgur

Continuous Electron Beam Accelerator Facility
12000 Jefferson Avenue
Newport News, VA 23606

In this talk I review and comment upon recent developments in hadron spectroscopy and structure. The talk is organized into three main sections dealing with heavy quarkonia ($Q\bar{Q}$), hadrons containing a single heavy quark ($Q\bar{q}$ and Qqq), and hadrons containing only light quarks and glue, although I will emphasize a surprising unity of the phenomena characterizing these systems.

---

† plenary talk given at the XXVI International Conference on High Energy Physics, Dallas, August 1992.

HADRON SPECTROSCOPY AND STRUCTURE*

Nathan Isgur
Continuous Electron Beam Accelerator Facility
12000 Jefferson Avenue
Newport News, VA 23606

Abstract

In this talk I review and comment upon recent developments in hadron spectroscopy and structure. The talk is organized into three main sections dealing with heavy quarkonia ($Q\bar{Q}$), hadrons containing a single heavy quark ($Q\bar{q}$ and Qqq), and hadrons containing only light quarks and glue, although I will emphasize a surprising unity of the phenomena characterizing these systems.

INTRODUCTION

To an audience like this it seems appropriate to begin by addressing the question of whether hadronic physics is still part of high energy physics and, whether it is or not, why you might still be interested in strong interaction physics.

To the first question I would give a qualified "yes" since I recognize in this audience—which must somehow define high energy physics—a substantial number of people whose work is dominantly in this area. My qualification is based on the fact that a growing fraction of the nuclear physics community is now devoted to the solution of the same basic problems.

Even if you are not a member of the strong interaction community of physicists, it seems to me that there are several reasons why you might still wish to follow this subfield:

(1) Our inability to understand from first principles strongly interacting matter, be it hadronic or nuclear, is a basic unsolved problem in describing the world around us. Of course, QCD solves this problem *in principle*, but that is of rather limited consolation to anyone who wants to know "how it works". I for one find this as unsatisfying as being told that QED *in principle* explains all of atomic, molecular, and biological science!

(2) Even if you are not bothered by leaving all of this science behind, you might worry that beyond the Standard Model is an underlying spontaneously broken strongly interacting gauge theory. Such a theory seems likely to be much more complex than QCD, so that it might be sensible to understand the basic physics of a theory we can study in detail experimentally before speculating about the next one.

(3) As a practical matter, it is often necessary to understand strong interaction effects before extracting information from experimental data relevant to other parts of the Standard Model. For example, a heavy quark weakly decays while embedded in a

*research supported in part by the U.S. Department of Energy under contract DE-AC05-84ER40150.

hadron, so that the extraction of the Cabibbo-Kobayaski-Maskawa matrix elements depends on being able to relate quark-level amplitudes to hadronic ones.

(4) Finally, and to my taste most importantly, there is every reason to believe that there is still a lot of good physics waiting to be done in the strong interactions. It is of course possible to take the point of view that between QCD and strong interaction phenomena one will encounter nothing but pure complexity, but I believe this is unlikely. Consider, for example, the success of the constituent quark model: surely this is a signal of underlying simplicity waiting to be understood. This is only one of a long list of critical issues waiting to be addressed, including the origin of confinement and chiral symmetry breaking (lattice gauge theory has already taught us a lot—but not yet enough—about these phenomena), the puzzling absence to date of the gluonic degrees of freedom in low energy spectroscopy, and the origin of the Okubo-Zweig-Iizuka rule.

$Q\bar{Q}$: HEAVY QUARKONIA

The discovery of the charmonium system in November 1974 was a watershed in the history of the Standard Model. It simultaneously contributed to electroweak theory by supplying the missing quark of the Glashow-Illiopoulos-Maiani mechanism and to the acceptance of both QCD and the quark model. Since then both the charmonium and Υ families have continued to be very fruitful testing grounds for the Standard Model; the last two years have been particularly productive in this regard.

$Q\bar{Q}$ Theory

One of the most important recent advances in the theoretical treatment of heavy quarkonia has been the development by Lepage and Thacker[1] and collaborators of *nonrelativistic QCD* (NRQCD). Nonrelativistic QCD is *not* a nonrelativistic potential model, but rather a systematic and in principle exact expansion of QCD field theory in inverse powers of the heavy quark mass m_Q. NRQCD is thus an effective field theory which is *equivalent* to QCD. Its great advantage is that it removes m_Q as a scale in the treatment of heavy quarkonia, leaving only $m_Q\alpha_s$ (the "Bohr radius") and $b^{\frac{1}{2}}$ (where b is the QCD string tension) as important scales, thereby enormously expediting the treatment of heavy quarkonia using lattice techniques. The advent of NRQCD marks the end of the era of the hegemony of potential models for $Q\bar{Q}$ systems, and the beginning of *rigorous*, precision tests of QCD based on these systems. An example of this is the addition of α_s determinations from the Υ spectrum to those available from other methods[2].

Another recent advance focuses on a long-standing problem in calculating the total widths of P-wave quarkonia. For example, from diagrams like that of Figure 1, one can calculate that[3]

$$\Gamma(Q\bar{Q}\,{}^1P_1) = \frac{|R'(0)|^2}{M_Q^4}\alpha_s^2\left\{0+\frac{\alpha_s}{\pi}\left[\frac{8}{3}\log\frac{M_Q}{\mu}+\text{constant}\right]+\ldots\right\} \quad (1)$$

where μ is an infrared cutoff. In this formula $R'(0)$ is the derivative of the radial $Q\bar{Q}$ wavefunction at the origin (which enters because $R(0)=0$ for a P-wave state, so that the $Q\bar{Q}$ annihilation is forced to live off the virtual intermediate heavy quark propagator) and the zero inside the bracket is to emphasize that this state has no α_s^2 term (in contradistinction to its ${}^3P_2$ and ${}^3P_0$ partners). It has usually been assumed that μ is an unknown nonperturbative parameter which spoils the predictability of the ${}^3P_1$ and ${}^1P_1$ hadronic widths. (The ${}^3P_2$ and ${}^3P_0$ widths also suffer from such infrared effects at order α_s^3, but these "corrections" have usually been ignored compared to

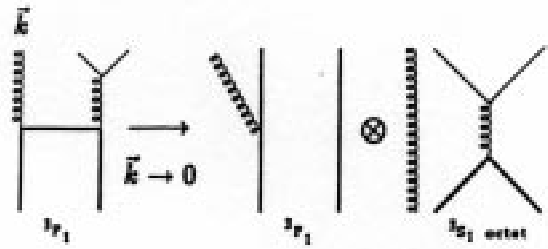

Figure 1. A graph contributing to the total hadronic width of a $^3P_1$ quarkonium state and the factorising part of this graph which is infrared singular.

the "leading" α_s^3 terms.) Bodwin, Braaten, and Lepage presented to this conference[4] a simple, predictive, and rigorous remedy for this situation. They showed that the infrared divergent part of Fig. 1 factorizes into the product of the amplitude for the process in which the Q or $\bar{Q}$ emits a soft gluon, leaving the $Q\bar{Q}$ system in a $^3S_1$ color octet state, and the amplitude for this state to annihilate through a single hard gluon into a light $q\bar{q}$ pair. The first amplitude is, they show, part of a Fock-space expansion of the state of the system, *i.e.*, the "$^3P_1$ quarkonia" have both a wavefunction for being $Q\bar{Q}$ in a $^3P_1$ state, and a wavefunction for being $gQ\bar{Q}$ with overall $J^{PC} = 1^{++}$ but with the $Q\bar{Q}$ subsystem in a color octet $^3S_1$ state. They go on to show that all of the nonperturbative information required (to this order in α_s) resides in $R'(0)$ and in a $gQ\bar{Q}$ "wavefunction at the origin" which absorbs the μ dependence of eq. (1).

In this very limited review of recent theoretical developments in heavy quarkonia, I will close with two advances in understanding the nonperturbative *gluonic* dynamics of QCD in the presence of static quark sources. The ground state of QCD in the presence of a static $Q\bar{Q}$ pair separated by a distance r is, in the approximation in which all dynamical quarks are ignored, the static $Q\bar{Q}$ potential which we associate with nonrelativistic potential models. Lattice QCD calculations of this gluonic ground state energy $V_0(r)$ are, indeed, in reasonable correspondence with the phenomenological potentials needed in such models. One can also consider the excited states of the glue in the presence of the static $Q\bar{Q}$ sources[5]. In the adiabatic approximation, these excited potentials $V_i(r)$ $(i > 0)$ will lead to entirely new spectroscopies in which the quarks and glue are simultaneously excited: the quarkonium hybrids[#1]. Perantonis and Michael[7] have recently been able to obtain clear signals on the lattice for such excited gluonic states and to measure their energies $V_i(r)$ and quantum numbers (see Fig. 2). They find that the first excited potential is doubly degenerate with the quantum numbers and energy of a *transverse phonon* as expected in a flux tube (or string) model[5]. It follows that exotic quarkonia with $J^{PC} = 0^{+-}$, 1^{-+}, and 2^{+-} should exist at modest excitation energies as predicted by Ref. 5.

In an independent development, Olson, Olsson, and Williams[8] presented to this conference analytic evidence for the relativistic flux tube model. They extend the earlier results of Eichten and Feinberg[9], Gromes[10], and Buchmuller[11] on the Wilson loop reduction of QCD. These earlier results had shown that a Lorentz scalar confinement potential matches the spin-dependence of QCD; they extend the comparison to the *spin-independent* v^2/c^2 corrections, finding that a scalar confinement potential fails to give these latter corrections, but that a relativistic flux tube model matches QCD to this order. The main effect is identi-

[#1] Hybrids were introduced in the context of the bag model: see Ref. 6 for some early papers.

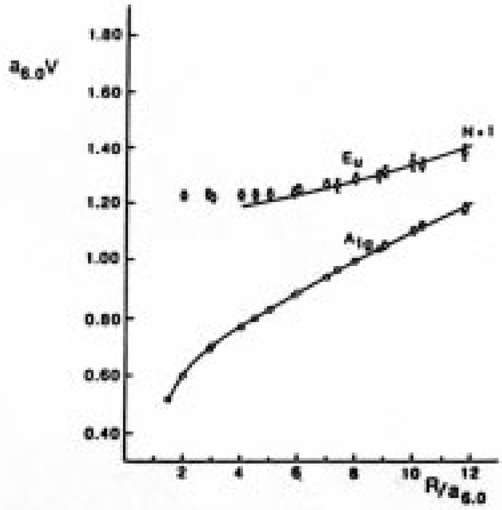

Figure 2. The ground and first excited gluonic energies in the presence of static triplet-antitriplet color sources at separation r (adapted from Ref. 7); the points are lattice data, while the curves are the adiabatic energies of the Nambu-Goto relativistic string. Note that $V_1(r) - V_0(r) = \pi/r$ for large r as expected for a transverse phonon.

Figure 3. The charmonium $^1P_1$ state completes the measurement of the lowest S- and P-wave levels of the positronium-like spectrum started in 1974.

fied as being the flux tube moment-of-intertia!

These two results, showing that at low energy the gluonic fields in a static $Q\bar{Q}$ system are string-like, indicate that heavy quarkonium hybrids may be as well described by potential models as ordinary heavy quarkonia. A search for the J^{PC} exotic hybrids of this type certainly now seems warranted. The more speculative extension of these ideas to light quark hybrids[8] may also now be considered to be on a better foundation. I will discuss these states below in the section on Light Quarks and Glue.

$Q\bar{Q}$ Experiment

The headline story of this section is certainly the report given to this conference of the discovery[13] of the long-missing $^1P_1$ charmonium by Fermilab's low energy $\bar{p}$-gas jet collider. They found the h_c at a mass of $3525.2 \pm 0.15 \pm 0.20$ MeV with a width $\Gamma <$ 1.1 MeV (at 90% c.l.). This mass is very close to the center-of-gravity (c.o.g.) of the ^3P_J states at 3525.3 ± 0.10 MeV, where it is naively expected. Halzen, Olson, Ollson, and Strong reported[13] to this conference that the one-loop perturbative correction to the $^1P_1 - {}^3P_{c.o.g.}$ is $+0.7 \pm 0.2$ MeV as observed. However, I am surprised by the smallness of the observed splitting, because I find it difficult to understand why the non-perturbative couplings of the P-waves to virtual decay channels couldn't produce relative shifts amongst these states of order 10 MeV. It would be interesting to understand whether the decoupling from these channels is really as complete as it would appear to be from this result (even though the continuum is only a few hundred MeV's away) or if this degeneracy is mainly accidental. (Altshuler and Silverman reported[14] finding other sources of splittings of this general size.) In any event, the discovery of the $^1P_1$ now completes in a significant way the picture started in the November Revolution of 1974 (see Fig. 3).

While this measurement stole the headlines, other significant results were presented at this conference, including another E760 result[13] giving $\Gamma(\chi_{c_2} \to \gamma\gamma) = 0.34 \pm 0.11$ KeV (compared to 0.8 ± 0.4 KeV expected) and an ARGUS measurement[15] of $\Gamma(\eta_c \to \gamma\gamma) =$ 12.2 ± 3.0 keV (compared to 7 keV expected).

Yet another $\gamma\gamma$ state, the η_c', remains obscure. Ref. 14 argues that the 90 MeV splitting of the current candidate from the ψ' is difficult to understand; clarifying the existence of this state should clearly be one of the important tasks on E760's agenda.

$Q\bar{q}$ and Qqq: HEAVY-LIGHT SYSTEMS

There are very few cases in which it is possible using analytic methods to make systematic predictions based on QCD in the low-energy, nonperturbative regime. Indeed, the theory has proved so intractable to analytic methods that all such predictions are based not on dynamical calculations, but rather on some symmetry of QCD. Isospin symmetry was the first such symmetry discovered, and we now understand that this approximate symmetry arises because the light quark mass difference $m_d - m_u$ is much smaller than the masses associated with confinement, which are set by the QCD scale Λ_{QCD}. Predictions based on isospin symmetry would, in a world with only strong interactions, be exact in the limit $m_d - m_u \to 0$; corrections to this limit can be studied systematically in an expansion in the small parameters $(m_d - m_u)/\Lambda_{QCD}$ and the electromagnetic fine structure constant α. $SU(3)$ flavor symmetry is similar, but the corrections are larger since $(m_s - m_d)/\Lambda_{QCD}$ is not small. Chiral symmetry $SU(2)_L \times SU(2)_R$ arises in QCD because *both* m_d and m_u are small compared to Λ_{QCD}; it is associated with the separate conservation of vector and axial vector currents. Although spontaneously broken in nature, the existence of this underlying symmetry allows the systematic expansion of chiral perturbation theory in which many low-energy properties of QCD are related to a few reduced matrix elements. If the strange quark mass is also treated as small compared with the QCD scale, then the chiral symmetry group becomes $SU(3)_L \times SU(3)_R$.

Over the last few years there has been progress in understanding systems containing a single heavy quark[16–26] (*i.e.*, a quark with mass m_Q much greater than the scale Λ_{QCD} of the strong interactions). It is now appreciated that there is a new symmetry of QCD, similar to isospin or chiral symmetry, in operation in such systems[21]. This symmetry arises because once a quark becomes sufficiently heavy, its mass becomes irrelevant to the nonperturbative dynamics of the light degrees of freedom of QCD. Consider, as an extreme example, two *very* heavy quarks of masses one and ten kilograms. Although these quarks will live in the usual hadronic "brown muck" of light quarks and glue, they will hardly notice it: their motion will fluctuate only slightly about that of a free heavy quark. Given that such quarks therefore define with great precision their own center-of-mass, we can study hadronic systems built on them in the frame where they act as static sources of color localized at the origin. The equations of QCD in the neighborhood of such an isolated heavy quark are therefore those of the light quark and gluonic degrees of freedom subject to the boundary condition that there is a static triplet source of color-electric field at the origin (*i.e.*, the heavy quark can be treated as a Wilson line). Since this boundary condition is the same for both of our hypothetical heavy quarks (in the static approximation which is essentially perfect given their masses), the solutions for the states of the light degrees of freedom in their presence will be the same (see Fig. 4). Thus *the light degrees of freedom will be symmetric under an isospin-like rotation of the heavy quark flavors into one another* even though the heavy quark masses are *not* almost equal. In particular, the heavy meson and baryon excitation spectra built on any heavy quark will be the same, as will be all amplitudes for the scattering of light hadrons off any state built on the heavy quark.

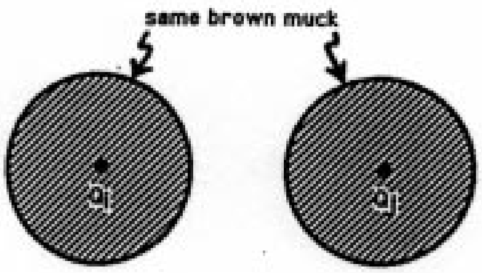

Figure 4. Q_i and Q_j are surrounded by identical brown muck

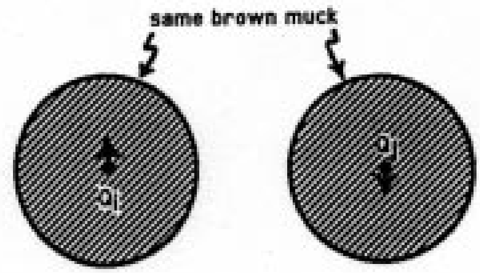

Figure 5. . . . even if the spin of Q_j is flipped

The preceeding comments ignored the spin of the heavy quark. This is appropriate in QCD since the spin of a heavy quark decouples from the gluonic field[18]: all heavy quarks look like scalar heavy quarks to the light degrees of freedom. Since the flavor *and* spin of the heavy quark are irrelevant, the static heavy quark symmetry is actually $SU(2N_h)$, where N_h is the number of heavy quarks (see Fig. 5). (The full symmetry group is actually much larger since heavy quarks moving with different velocities cannot be scattered into each other by the strong interactions.) At the spectroscopic level this additional symmetry means that each spectral level built on a heavy quark (unless it happens to have spin zero in its light degrees of freedom) will be a degenerate doublet in total spin.

Heavy quark flavor symmetry is thus analogous to the fact that different isotopes of a given element have the same chemistry: their electronic structure is almost identical because they have the same nuclear charge. The spin symmetry is in turn analogous to the near degeneracy of hyperfine levels in atoms: the electronic structure of the states of a hyperfine multiplet are almost the same because nuclear magnetic moments are small.

In the situation described above where the light degrees of freedom (*i.e.*, light quarks

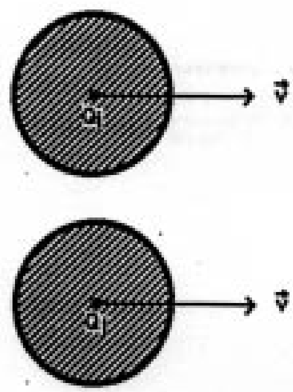

Figure 6. $Q_i(\vec{v})$ is related by the symmetry to $Q_j(\vec{v})$

and antiquarks and the gluons) typically have four-momenta small compared with the heavy quark mass, it is appropriate to go over to an effective theory where the heavy quark mass goes to infinity, with its four-velocity fixed[21,22] (see Fig. 6). The symmetry is based on heavy quarks with fixed *velocity* because the mass and therefore momentum of the heavy quark are irrelevant: the important variable is the motion of the imperturbable center to which the light quarks and gluons must respond dynamically.

The $SU(2N_h)$ spin-flavor symmetry of the heavy quark effective theory is not manifest in the full theory of QCD; it only becomes apparent in the effective theory where the heavy quark masses are taken to infinity. This situation is familiar from our experience with the light quark flavor symmetries of QCD mentioned above. The strong interactions of light quarks q (with masses m_q that are much less than the QCD scale) are greatly simplified by going over to an effective theory where the light quark masses are taken to zero. For N light quarks this effective theory has an $SU(N)_L \times SU(N)_R$ chiral symmetry that is spontaneously broken to the vector $SU(N)_V$ subgroup. Again, the symmetry is not immediately apparent in the full theory of QCD. However, as long as the quark masses are small compared with the QCD scale, they have only a small impact on strong interaction dynamics. Thus the effective theory, where the light quark masses are set to zero, is a good approximation to QCD. The heavy quark flavor-spin symmetry endows us with predictive power much in the same way that light quark chiral symmetry does. For light quark chiral symmetry, it is possible to treat the small quark masses as perturbations and consider the corrections of order m_q/Λ_{QCD} to predictions based on the effective theory where $m_q \to 0$. Similarly, for heavy quark spin-flavor symmetry, it is possible to treat as perturbations the Λ_{QCD}/m_Q corrections to the predictions based on the effective heavy quark theory where $m_Q \to \infty$.

The relationship between operators involving the heavy quarks (e.g., $\bar{Q}\gamma_\mu q$) in the full theory of QCD and operators in the effective theory where the heavy quark masses go to infinity involves some interesting applications of perturbative QCD. Contributions to matrix elements of these operators from loop graphs with virtual momenta comparable to or greater than the heavy quark mass are clearly not correctly reproduced by the effective theory. However, because of asymptotic freedom these differences can be handled by perturbative QCD. In what follows I will suppress the multiplicative matching factors required to connect the low energy effective theory which we are discussing with full QCD; see Ref. 26 for a discussion of these effects and for further references.

Most of the physics underlying heavy quark symmetry has been understood for a long time and has, to some extent, been incorporated into phenomenological models used to predict properties of hadrons containing a single heavy quark.[27,28] What is new is that we now understand that this physics arises from symmetries of an effective theory that is a systematic limit of QCD. Consequently, model-independent predictions are now possible, including important predictions for semileptonic B-meson decay form factors. These are expected to play a vital role in the accurate determination of the values of the Cabibbo–Kobayashi–Maskawa matrix elements V_{cb} and V_{ub} from experimental data. Before looking at these predictions, it is useful to set the stage by considering the implications of heavy quark symmetry for heavy-light spectroscopy[29].

In the limit $m_Q \to \infty$, the spin of the heavy quark $\vec{S}_Q$ and the spin of the light degrees of freedom (*i.e.*, the angular momentum of the light degrees of freedom in the heavy quark's rest frame)

$$\vec{S}_\ell = \vec{S} - \vec{S}_Q \ , \qquad (2)$$

are separately conserved by the strong interactions (here $\vec{S}$ is the angular momentum of both the heavy quark and the light degrees of freedom in the heavy quark's rest frame, *i.e.*, the total spin). Therefore, in this limit, s_Q, m_Q, s_ℓ, and m_ℓ are good quantum numbers. Since the dynamics are completely independent of the mass and spin of the heavy quark Q it is convenient to classify states containing a single heavy quark by s_ℓ. Then associated with each such state for the light degrees of freedom will

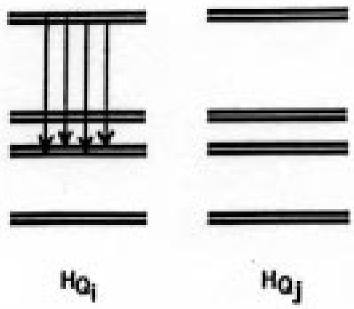

Figure 7. the spectra and transitions of the hadrons built on Q_i and Q_j; note that once the ground states are lined up, the full spectrum of other states built around each heavy quark will match up, *i.e.*, mesons, baryons, continua, and kitchen sinks are all included in this spectral diagram.

be a degenerate doublet of hadrons with total spins (formed from combining the spin of the heavy quark $s_Q = 1/2$ with the spin of the light degrees of freedom s_ℓ)

$$s_\pm = s_\ell \pm 1/2 \; , \tag{3}$$

(unless $s_\ell = 0$, in which case a single $s = 1/2$ state is obtained). The flavor symmetry ensures that the spectrum is identical for each flavor Q up to an overall constant mass shift associated with the mass of the heavy quark. Of course, states are also labeled by their parity π (which is the same as the parity of the light degrees of freedom, π_ℓ, since the heavy quark has positive parity) and by other "radial" quantum numbers (see Fig. 7).

To get a better picture of how this works, let's consider the mesons with $Q\bar{q}$ flavor quantum numbers. (Note that although we use the *language* of the constituent quark model, our conclusions will be completely general.) It is reasonable to assume that the ground state mesons with these flavor quantum numbers have $s_\ell = 1/2$ and negative parity, forming a doublet consisting of a spin zero state ($s_- = 0$) which we denote by P_Q and a spin one state ($s_+ = 1$) which we denote by P_Q^*. In the case $Q = c$, these are the D and D^* mesons, and in the case $Q = b$, these are the $\bar{B}$ and $\bar{B}^*$ mesons. In terms of the spin of the heavy quark and the spin of the light degrees of freedom, the states (at rest) are

$$|P_Q\rangle = \frac{1}{\sqrt{2}}[|\uparrow\downarrow\rangle - |\downarrow\uparrow\rangle] \; , \tag{4}$$

and

$$|P_Q^*\rangle = \frac{1}{\sqrt{2}}[|\uparrow\downarrow\rangle + |\downarrow\uparrow\rangle] \; , \tag{5}$$

where the state $|P_Q^*\rangle$ in eq. (5) has zero component of total spin along the quantization axis $\hat{z}$. In eqs. (4) and (5) the first arrow in a ket refers to the spin of the heavy quark along the z-axis, while the second arrow in a ket refers to that of the light degrees of freedom. Acting with the z-component of the heavy quark spin then gives

$$S_Q^z|P_Q\rangle = \frac{1}{2}|P_Q^*\rangle \; . \tag{6}$$

Since $\vec{S}_Q$ commutes with the Hamiltonian, the P_Q and P_Q^* states are degenerate in mass.

Since the heavy quark flavor symmetry applies to the complete set of n-point functions of the theory, not only mass splittings, but also all strong decay amplitudes arising from the emission of light quanta like $\pi, \eta, \rho, \pi\pi$, etc., are independent of heavy quark flavor. For a given heavy quark flavor the spin symmetry ensures that two states with spins $s_\pm$ must have the same total widths. This equality between total widths typically arises in a nontrivial way. The two states of a given multiplet can decay to both states of every available multiplet with distinct partial widths whose sum must be identical. The spin symmetry determines the ratios of these partial widths (see Fig. 7).

The heavy quark symmetry cannot, of course, tell us anything about the spectroscopy of the light degrees of freedom. It can only predict relationships between heavy quark systems involving given states of these degrees of freedom. As we have mentioned for mesons with $Q\bar{q}$ flavor quantum numbers, both the constituent quark model and experiment suggest that the ground states have $s_\ell^{\pi_\ell} = 1/2^-$ giving the $s_-^\pi = 0^-$ and $s_+^\pi = 1^-$ states P_Q and P_Q^*. The constituent quark model also suggests that the lowest lying excited states are likely to be those which correspond to giving the spin 1/2 constituent antiquark a unit of orbital angular momentum resulting in $s_\ell^{\pi_\ell} = 1/2^+$ and $3/2^+$ multiplets. It is easily shown that the $s_+^\pi = 2^+$ state of the $s_\ell^{\pi_\ell} = 3/2^+$ multiplet has decay amplitudes in the proportions $\sqrt{(2/5)} : \sqrt{(3/5)}$ to the states $[P_Q\pi]_{L=2}$ and $[P_Q^*\pi]_{L=2}$ respectively. Its multiplet partner, with $s_-^\pi = 1^+$, decays at the same total rate exclusively to $[P_Q^*\pi]_{L=2}$. Note that the $s_-^\pi = 1^+$ state does not decay to $[P_Q^*\pi]_{L=0}$ even though this is an allowed channel. Similarly, the $s_+^\pi = 1^+$ state of the $s_\ell^{\pi_\ell} = 1/2^+$ multiplet decays exclusively to $[P_Q^*\pi]_{L=0}$, and does not decay to $[P_Q^*\pi]_{L=2}$. Its $s_-^\pi = 0^+$ state decays to $[P_Q\pi]_{L=0}$ with the same total rate[*2]. These predictions are compatible with existing experimental information on mesons containing a charm quark (see below).

Let's now consider some matrix elements of operators in the effective heavy quark theory. The matrix elements we focus on are those that are likely to play an important role in determining the Cabibbo–Kobayashi–Maskawa matrix elements V_{cb} and V_{ub}. From the basic physics of heavy quark symmetry described above, and the cartoons given in Figures 8 and 9, it is not difficult to see that all of the form factors operative in $\bar{B} \to De\bar{\nu}_e$ and $\bar{B} \to D^* e\bar{\nu}_e$ are related. With an unconventional choice of normalizations and form factors which are suitable for a situation where four-velocities and not four-momenta are the relevant variables, namely[*3] $\langle P(v')|P(v)\rangle = 2\gamma\delta^3(v'-v)$, *etc.* and

$$\begin{aligned}&\langle P_{Q_j}(v')|\bar{Q}^{(j)}\gamma_\mu Q^{(i)}|P_{Q_i}(v)\rangle\\&= \tilde{f}_+(v+v')_\mu + \tilde{f}_-(v-v')_\mu \,. \qquad (7)\end{aligned}$$

for $\bar{B} \to D$ and

$$\begin{aligned}&\langle P_{Q_j}^*(v',\epsilon)|\bar{Q}^{(j)}\gamma_\mu\gamma_5 Q^{(i)}|P_{Q_i}(v)\rangle\\&= \tilde{f}\epsilon_\mu^* + (\epsilon^*\cdot v)[\tilde{a}_+(v+v')_\mu\\&\quad +\tilde{a}_-(v-v')_\mu] \,, \qquad (8)\end{aligned}$$

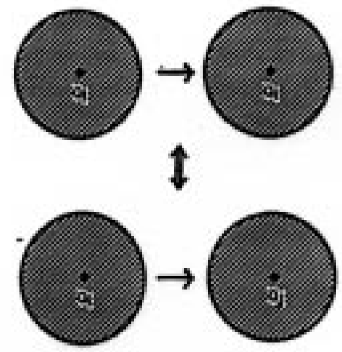

Figure 8. $Q_i \to Q_j$ is related to the elastic transition $Q_i \to Q_i$ by the symmetry

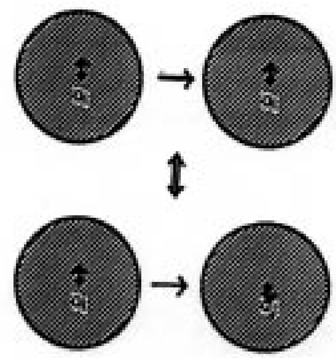

Figure 9. . . . even if a spin flips

[*2]These results were first noted by Rosner[50] who obtained them by taking the $m_Q \to \infty$ limit of a quark model calculation. Heavy quark symmetry allows us to see that they are model independent consequences of QCD in that limit.

[*3]Here $v^\mu = p^\mu/m$ is the four-velocity, and $\gamma = v^0$.

and

$$\langle P^*_{Q_j}(v',\epsilon)|\bar{Q}^{(j)}\gamma_\mu Q^{(i)}|P_{Q_i}(v)\rangle \\ = i\bar{g}\epsilon_{\mu\nu\lambda\sigma}\epsilon^{*\nu}v'^{\lambda}v^{\sigma} \ . \qquad (9)$$

for $\bar{B} \to D^*$, one can show that in fact the form factors characterizing the $P_{Q_i} \to P_{Q_j}$ and $P_{Q_i} \to P^*_{Q_j}$ matrix elements of the vector and axial vector currents are expressible in terms of a single universal function $\xi(w)$, where $w = v \cdot v'$, that is normalized to unity at zero recoil[N4]. Explicitly[31],

$$\bar{f}_+ = \xi \ , \ \ \bar{f}_- = 0 \ , \qquad (10)$$

$$\bar{f} = (1+w)\xi \ , \qquad (11)$$

$$(\bar{a}_+ - \bar{a}_-) = -\xi \ , \ \ (\bar{a}_+ + \bar{a}_-) = 0 \ , \qquad (12)$$

and

$$\bar{g} = \xi \ , \qquad (13)$$

with

$$\xi(1) = 1 \ . \qquad (14)$$

The function ξ is truly universal: it doesn't depend on the heavy quark's mass or spin, nor does it depend on the current which causes the $Q_i \to Q_j$ transition. The same function would even play a role in the physics of hadrons containing other heavy color triplet particles. Many extensions of the standard model (*e.g.*, supersymmetry and technicolor) contain such heavy spin zero color triplets.

The manipulations leading to eqs. (10) to (14) bear a striking resemblance to the method for deriving the predictions of light quark $SU(3)_V$ symmetry for form factors in the $K \to \pi$ matrix element of the current $\bar{s}\gamma_\mu d$. However, we see from eqs. (7) and (10) that heavy quark symmetry does not predict the conventionally defined f_- form factor to be zero and it provides a normalization, not at $q^2 \equiv (p'-p)^2 = 0$, but rather at the maximum value $q^2_{max} = (m_{P_i} - m_{P_j})^2$, corresponding to $w = 1$, where both the initial and final hadrons have the same four-velocity. We call this kinematic point zero recoil, since in the rest frame of the initial hadron the final hadron is also at rest.

Transition matrix elements involving the ground state (isospin-zero) baryons with Q_iud flavor quantum numbers are even easier to deduce than those involving the ground state mesons. These baryons are denoted by Λ_{Q_i} and we assume that they have $s_\ell^{\pi_\ell} = 0^+$ (this is suggested by the constituent quark model and in the case $Q = c$ is required by experiment). It is straightforward to prove that[31,32]

$$\langle \Lambda_{Q_j}(v',s')|\bar{Q}^{(j)}\Gamma Q^{(i)}|\Lambda_{Q_i}(v,s)\rangle \\ = \eta\bar{u}(v',s')\Gamma u(v,s) \ , \qquad (15)$$

where η is a universal function of w independent of the heavy quark masses. The heavy quark flavor symmetry implies once again that at zero recoil

$$\eta(1) = 1 \ . \qquad (16)$$

Apart from the overall factor of η, eq. (15) shows that the hadronic matrix element is like a heavy quark matrix element. This occurs because in a Λ_Q state the spin of the hadron is carried by the heavy quark.

Transition matrix elements between heavy and light states can also be considered. Here, the heavy quark flavor symmetry can be used to relate matrix elements involving different heavy quarks. For example[17,20], with our unconventional normalizations and as usual ignoring QCD matching factors,

---

[N4] This variable is called w after the French name for this letter.

$$\langle 0|\bar{q}\gamma_\mu\gamma_5 Q^{(i)}|P_{Q_i}(v)\rangle = \langle 0|\bar{q}\gamma_\mu\gamma_5 Q^{(j)}|P_{Q_j}(v)\rangle \; . \tag{17}$$

Similar relations hold to light final states like π and ρ provided, in the rest frame of the heavy quark, the final states have the same four-momenta[31]. As a result, heavy quark methods open an interesting avenue for determining the magnitude of the V_{ub} element of the Cabibbo–Kobayashi–Maskawa matrix. Ordinary isospin symmetry plus heavy quark symmetry implies that, for example,

$$\begin{aligned} &\langle\rho(k,\epsilon)|\bar{u}\gamma_\mu(1-\gamma_5)b|\bar{B}(v)\rangle \\ &= \left(\frac{m_B}{m_D}\right)^{1/2}\left[\frac{\alpha_s(m_b)}{\alpha_s(m_c)}\right]^{-6/25} \\ &\langle\rho(k,\epsilon)|\bar{d}\gamma_\mu(1-\gamma_5)c|D(v)\rangle \; , \end{aligned} \tag{18}$$

where, in a departure from the simplified discussion of this review, in this case the matching factor $\left[\frac{\alpha_s(m_b)}{\alpha_s(m_c)}\right]^{-6/25}$ has been included explicitly. Eq. (18) is valid in the rest frame of the $\bar{B}$ and D for momenta k not too large compared with the heavy quark masses.[*5] Since in the Cabibbo-suppressed semileptonic decay $D \to \rho\bar{e}\nu_e$ the weak mixing angles are known, the right side of eq. (18) can be determined experimentally. With this information, experimental data on $\bar{B} \to \rho e\bar{\nu}_e$ will allow a determination of $|V_{ub}|$ (see Fig. 10). If one uses light quark $SU(3)_V$ flavor symmetry instead of isospin, then the Cabibbo allowed semileptonic decay $D \to K^*\bar{e}\nu_e$ can be used. The form factors for this decay have already been determined experimentally[34]. Of course this strategy can be used for any convenient light hadronic final state.

It is possible to systematically improve order by order in $\alpha_s(m_c)$ and $\alpha_s(m_b)$ the matching between operators in the full theory of QCD and operators in the effective theory[35,36].

*5 Model calculations suggest that eq. (18) holds even for k comparable with the heavy charm quark mass[33].

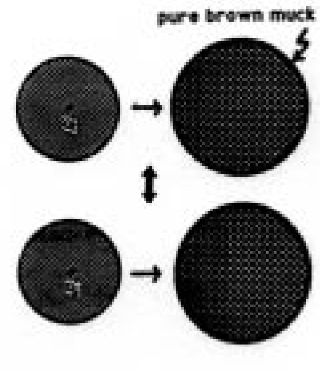

Figure 10. heavy-to-light transitions are also related by the symmetry

This gives calculable perturbative corrections and does not cause any loss of predictive power. Neubert[37] has recently examined the accuracy of the heavy quark symmetry predictions for the decays $\bar{B} \to De\bar{\nu}_e$ and $\bar{B} \to D^*e\bar{\nu}_e$, including both perturbative and $1/m_c$ corrections. As a result of Luke's theorem[38], which protects the zero recoil point from any $1/m_c$ corrections in leading order, he finds that the differential decay rates for these decays near zero recoil receive overall corrections of only $+7\%$ and -1%, respectively. There are therefore good reasons to believe that the determination of V_{cb} is now mainly an experimental problem.

Additional effects of $1/m_Q$ corrections have been studied in many models. Grinstein and Mende[39] and Burkhardt and Swanson[40] have shown how the symmetry is approached in $1+1$-dimensional QCD in the $N_c \to \infty$ limit; these issues have also been addressed in quark models too numerous to mention. The universal function $\xi(w)$ has also been calculated in many different schemes; at this conference we saw new results by Pirjol, Schilcher, and Wu[41] and de Rafael and Taron[42] to join earlier calculations by Radyushkin[43], Neubert[44], Bagan, Ball, Braun, and Dosch[45], and others[46].

Several groups have also realized that there is a natural way to marry heavy quark and chiral symmetry. This marriage was first arranged by M.B. Wise[47], T.-M. Yan *et al.*[48], and G. Burdman and J.F. Donoghue[49], and at this conference we heard of further developments in this direction from Cheng and Goity[50]. We also heard at this conference about a number of other interesting developments of the ideas and applications of heavy quark symmetry.

To me, one of the most exciting recent developments pertains to an ancient problem in particle phenomenology: can we understand the nonleptonic weak decays of hadrons? I had just finished some work on this subject at the time of the 1975 Palermo conference when Lipkin warned about the dangers of "Drinking Nonleptonic". He recalled how from time to time over the last 40 years there have been hallucinatory periods from "drinking nonleptonic", *i.e.*, spells of believing that the mystery of the $\Delta I = \frac{1}{2}$ rule and other puzzling features of these decays were on the verge of being solved, but that one always woke up from these giddy periods with a headache. We may still not fully understand the $\Delta I = \frac{1}{2}$ rule, but there is a good chance that we can now at least drink a little nonleptonic of the beauty flavor.

The new development is an extension of heavy quark effective theory by Dugan and Grinstein[51] to prove factorization *for a special class of weak hadronic decays* . Short distance QCD corrections create a low energy effective weak hadronic interaction

$$H_w^{eff} = \frac{G_F}{\sqrt{2}} V_{ud}^* V_{cb} [c_1 h_1 + c_8 h_8] \quad , \qquad (19)$$

where $h_1 = \bar{c}\gamma^\mu(1-\gamma_5)b\ \bar{d}\gamma_\mu(1-\gamma_5)u$ and $h_8 = \bar{c}\gamma^\mu(1-\gamma_5)\frac{\lambda^i}{\sqrt{2}}b\ \bar{d}\gamma_\mu(1-\gamma_5)\frac{\lambda^i}{\sqrt{2}}u$. If m_c/m_b is fixed, and $m_b \to \infty$, as in the heavy quark limit, then the c quark will recoil with fixed *velocity* against the residual hadronic system. Dugan and Grinstein show that in this limit the decay will be dominated by the c_1 term of H_w^{eff} so long as the $d\bar{u}$ system is at low mass. With this restriction to a collinear $d\bar{u}$ system, they show that the fast moving light quarks decouple from soft gluon exchanges; the octet light pairs therefore do not hadronize into a light $d\bar{u}$ meson, while the singlet pairs factorize from the $b \to c$ process. This leaves the latter process controlled by the same matrix elements as occur in semileptonic $b \to c$ decays, *i.e.* matrix elements that are governed by heavy quark symmetry. Dugan-Grinstein factorization therefore offers a real chance at a systematic expansion of this class of nonleptonic decays in powers of $1/m_q$ and α_s.

Heavy-light experiment

Given the venue of this meeting and the recent drama which unfolded in Congress, it is clear that this section should begin with the announcement that *ssc is confirmed.* Both ARGUS[52] and FNAL E687[53] reported at this conference the observation of the doubly strange Ω_c at masses of $2716 \pm 5 \pm 5$ MeV and $2707 \pm 2 \pm 5$ MeV respectively. These observations are consistent with both theoretical expectations and with an earlier claim[54] by the CERN hyperon beam experiment WA-62.

We also heard at this conference that ARGUS[55] confirms the $D_{s_1}(2536)^+$ in both the $D^{*0}K^+$ and $D^{*+}K^0_S$ channels. We now know that excited D_s spectroscopy matches excited D spectroscopy, and that both are consistent with heavy quark symmetry (see Fig. 11): they display the expected doublet pattern appropriate to an $s_\ell^{\pi_\ell} = \frac{3}{2}^+$ multiplet and the narrow D_{s_1} ($\Gamma_{D_{s_1}} < 3.9$ MeV) and D_1 required by the symmetry, which predicts that the kinematically allowed D^*K and $D^*\pi$ S-wave decays will be suppressed[29].

The LEP experiments are all beginning to try to exploit the fact that the Z is a good flavor factory, and we heard many re-

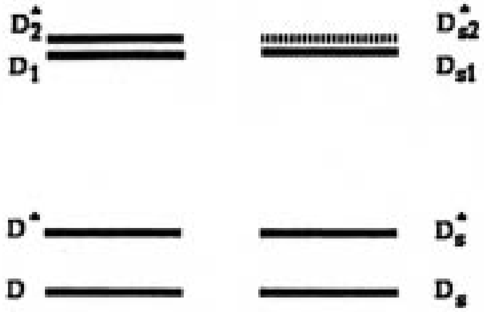

Figure 11. The observed $s_\ell^{\pi_\ell} = \frac{3}{2}^+$ spectra of excited D_s and D mesons, displaying the expected doublet structure. The reduced widths (*i.e.*, widths with phase spaced removed) of the multiplet partners are also nearly equal as expected.

ports of progress begin made in studying the B and Λ_b in a high energy environment[56]. Especially impressive was the strong evidence presented[57] by OPAL for the B_s. (They even report one candidate for the decay $B_s \to \psi\phi$ with a mass of 5.36 GeV.) We can expect a steady improvement in the masses, lifetimes, and branching ratios of all of these states from these experiments; they have an especially important role to play in Λ_b and B_s physics where we may be relying on high energy colliders for the bulk of our knowledge. There are many important issues to be resolved here. As a very first step, it is desirable to simply accurately measure the mass of the Λ_b (heavy quark symmetry predicts that $\Lambda_b = \left[\frac{B+3B^*}{4}\right] + \left[\Lambda_c - \left(\frac{D+3D^*}{4}\right)\right] + 0\left(\frac{\Lambda^2_{QCD}}{D}\right)$); eventually, the Λ_b semileptonic decays could prove to be the most accurate means of determining the CKM angles V_{cb} and V_{ub}.

CLEO[58,59] also reported a series of new measurements of isospin-breaking mass differences in the D, D^*, and Σ_c multiplets:

$$D^{*+} - D^{*0} = 3.32 \pm 0.08 \pm 0.05 \text{ MeV} \quad (20)$$
$$D^{+} - D^{0} = 4.80 \pm 0.10 \pm 0.06 \text{ MeV} \quad (21)$$
$$\Sigma_c^{+} - \Sigma_c^{0} = 1.1 \pm 0.6 \pm 0.5 \text{ MeV} \quad (22)$$
$$\Sigma_c^{++} - \Sigma_c^{+} = -0.2 \pm 0.6 \pm 0.5 \text{ MeV} \quad (23)$$

The first two splittings follow the general pattern of dominance of such splittings by the $d-u$ mass difference. However, there is a surprise in these splittings, in that heavy quark symmetry predicts that they would be equal in the limit $m_c \to \infty$, and (especially given that $D_s^* - D_s \simeq D^* - D$) the 1.5 MeV mass difference between them seems difficult to explain as a $1/m_c$ effect[60]. The second pair of splittings seems superficially within expectations, but I suspect there is a problem here: from them one can form the $\Delta I = 2$ combination $\Sigma_c^{++} + \Sigma_c^0 - 2\Sigma_c^+ = -1.3 \pm 0.7 \pm 0.6$ MeV which is pure electromagnetic and should turn out to be roughly $+1.3$ MeV.

Finally, many new tests of nonleptonic factorization were presented to the conference. There were, first of all, many reported *failures* of factorization for charm decays where Dugan-Grinstein (DG) factorization should not apply but where Bauer-Stech-Wirbel (BSW) factorization[61] has enjoyed considerable success. BSW factorization hypothesizes that a much broader applicability of factorization to *all* amplitudes, not just collinear ones as in DG factorization, can be obtained if the coefficients c_1 and c_2 of eq. (19) are treated as free parameters. BSW also assume that it is possible to reexpress H_w^{eff} in terms of two related hadronic operators with coefficients a_1 and a_2. In fact, the BSW model remains quite successful in those circumstances where only the c_1 term contributes. However, in other circumstances there now seem to be problems. For example, ARGUS[62] has measured

$$\frac{\Gamma(\Lambda_c^+ \to \Sigma^0\pi^+)}{\Gamma(\Lambda_c^+ \to \Lambda\pi^+)} = \frac{1.8 \pm 0.5 \pm 0.3}{2.2 \pm 0.3 \pm 0.4} \quad , \quad (24)$$

while factorization predicts *zero* since the $I = 0$ $c \to s$ current can't induce a $\Delta I = 1$ $\Lambda_c^+ \to \Sigma^0$ transition. ARGUS also reports[62]

$$\frac{\Gamma(D^0 \to \rho^0 \bar{K}^0)}{\Gamma(D^0 \to K^{*-}\pi^+)} = \frac{0.24 \pm 0.03 \pm 0.03}{0.70 \pm 0.04 \pm 0.04} \quad , \tag{25}$$

much larger than expected from BSW factorization since the process in the numerator proceeds through c_2 alone.

While BSW factorization is experiencing some difficulties, its less ambitious but more rigorous cousin DG factorization has experienced some successes. For example, CLEO has reported[59] an important set of new tests in B decay of DG factorization and heavy quark symmetry. Highlights of their results are given in Table 1. It is clearly too early to drink too heavily of nonleptonic, but there are reasons to remain optimistic.

LIGHT QUARKS AND GLUE

Finally we turn to systems of "pure brown muck". There is a recent revival of interest in these oldest hadronic systems, engendered by both new experimental and theoretical developments. It is an old refrain, but given its importance in structuring the world we live in, our ignorance of the nature of ordinary strongly interacting matter is an embarrassment, and it remains a challenge to close this glaring gap between our knowledge of the basic principles of nuclear and hadronic matter—QCD—and the phenomena which characterize it.

The theory of pure brown muck

One of the most productive theoretical tools in studying light quarks and glue remains the lattice. For example, within the last few years some basic characteristics of the quenched glueball spectrum have become established[63]:

(1) the lightest glueball has vacuum quantum numbers ($J^{PC} = 0^{++}$) and a mass of ~ 1.5 GeV.

(2) the next lightest state is probably a $J^{PC} = 2^{++}$ state with mass $\sim$ 2.2 GeV.

(3) the spectrum really begins above 2 GeV, but in the 2–3 GeV region there is so far no evidence for a J^{PC} exotic state.

These facts indicate that glueball hunting may be a sticky proposition. The 0^{++} meson sector between 1 and 2 GeV is the most poorly understood sector of hadron spectroscopy, and is *the* place that a malicious intelligence would hide a glueball. In addition, looking for a glueball with nonexotic quantum numbers above 2 GeV will be extremely difficult: the meson spectrum in this region will be densely populated with ordinary quarkonium states.

When combined with the recent evidence from the lattice discussed in the $Q\bar{Q}$ section for the flux tube picture of nonperturbative gluon dynamics, the bleak prospects for glueball hunting suggest that the most productive strategy for discovering a gluonic spectroscopy is to concentrate on the J^{PC} exotic hybrids predicted[5] (basically by extrapolation from the $Q\bar{Q}$ case) to be found in the vicinity of 1.8 GeV. Several searches are now underway for such states, which can be expected to be more readily produced than the $Q\bar{Q}$ hybrids.

In addition to this and other lattice activity, there have been continuing theoretical efforts of a more analytical nature. An important focus of activity has been to try to understand the nature of the constituent quark, and amongst the many papers on this subject, those of Manohar and Georgi[64], Kaplan[65], Weinberg[66], and Peris[66] received considerable attention. One of the central goals of this activity has been to attempt to "chiralize" the quark model. Karliner described at this

Table 1. Tests of DG factorization and heavy quark symmetry in B decays; the last two columns indicate which principles are being tested.

channel	expt[59]	theory	factorization?	h.q.s.?
$\bar{B}^0 \to D^{*+}\pi^-$	$1.28 \pm 0.19 \pm 0.30$	1.23 ± 0.17	$\surd$	
$\bar{B}^0 \to D^{*+}\rho^-$	$3.17 \pm 0.43 \pm 0.34$	3.26 ± 0.46	$\surd$	
$\frac{\Gamma_L}{\Gamma_T}(\bar{B}^0 \to D^{*+}\rho^-)$	$90 \pm 7 \pm 5\%$	88%	$\surd$	
$\frac{\bar{B}^0 \to D^+\pi^-}{\bar{B}^0 \to D^{*+}\pi^-}$	$0.96 \pm 0.19 \pm 0.25$	1	$\surd$	$\surd$
$\frac{\bar{B}^0 \to D^+\rho^-}{\bar{B}^0 \to D^{*+}\rho^-}$	$0.97 \pm 0.19 \pm 0.25$	1	$\surd$	$\surd$

meeting[67] a derivation in 1 + 1-dimensional QCD of Kaplan's *ansatz* in which the constituent quark is the "skyrmion" of the colored chiral Lagrangian after bosonization.

The experimental status of conventional brown muck

I am myself particularly fond of the much more mundane suggestion made by Bjorken: a constituent quark is what you get when you take a B meson and remove the heavy quark. To organize *conventional* light quark spectroscopy for you, let me adopt a version of this definition. First, in Figure 12, let's examine the onset of heavy quark symmetry; alternatively, this can be viewed as a demonstration that a constituent quark retains some of the characteristics of a heavy quark. We can see from this diagram that even the $\rho - \pi$ splitting roughly follows the $1/m_Q$ scaling law for the splitting of the $s_\ell^{\pi_\ell} = \frac{1}{2}^-$ multiplet. Note, moreover, that the center-of-gravity of the $s_\ell^{\pi_\ell} = \frac{3}{2}^+$ excited multiplet hardly moves with respect to that of the ground state multiplet, and that the $1/m_Q$ splitting within this excited multiplet remains small.

Figure 13 gives another look at the evolution of hadronic spectra within the (approximately) equal mass quarkonia from the Υ to the π. Once again we see the remarkable similarity of all of these spectra, suggesting that there are no dramatic changes in the underlying physics from the $b\bar{b}$ to the $u\bar{d}$ system, even though m_u, $m_d \ll \Lambda_{QCD}$. Somehow, a constituent light quark seems to behave much like a heavy quark[68].

There were several interesting additions to our knowledge of the spectra of Figure 13 made at this conference. I have already highlighted the discovery[12] of the $c\bar{c}$ $^1P_1$ state $h_c(3526)$. We also heard here that the Beijing e^+e^- machine has confirmed[69] the narrow $\xi(2200)$, which has an economical interpretation as one of the missing[68] F-wave $s\bar{s}$ states, namely $^3F_4$ or $^3F_2$. The LASS experiment, which has been enormously successful in clearing up many questions in strange particle spectroscopy, provided strong evidence[70] in the $s\bar{u}$ sector for both $^3D_2$ and $^1D_2$, the $K_2(1775)$ and $K_2(1820)$; this completes the discovery of expected[68] D-wave $S = -1$ quarkonia. In a related measurement, ARGUS[15] has provided a window on the nature of the $^1D_2$ state $\pi_2(1660)$ by measuring its $\gamma\gamma$ width to be 0.25 ± 0.10 keV. This width is much more understandable from quark potential models[71]

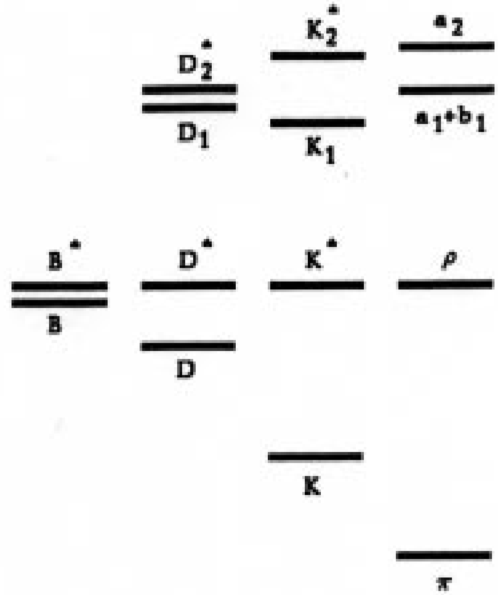

Figure 12. The onset of heavy quark symmetry

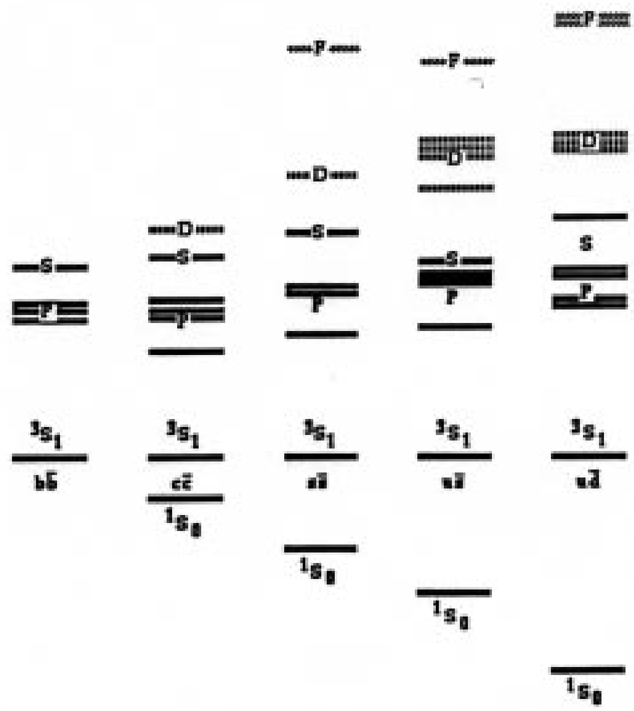

Figure 13. The experimental status of conventional quarkonia

than earlier reported results.

Many other new results in conventional spectroscopy were reported to this conference, but I will mention just one more which brings to a conclusion a long-standing uncertainty about the properties of the a_1 meson. The ARGUS collaboration has reported[72] an analysis of their new sample of $\tau \to \nu_\tau 3\pi$ data which pin down the elusive parameters of this state to be $m = 1211 \pm 7$ MeV and $\Gamma = 446 \pm 21$ MeV, with a D/S ratio for the dominant $\rho\pi$ decay of -0.11 ± 0.02. The mass once again leads to the conclusion that spin-orbit forces are suprisingly small; the width and D/S ratio are in good agreement with those expected in the flux tube model[73].

The experimental status of unconventional brown muck

The experts will have noticed that I have omitted a number of states from my picture in order to emphasize the simplicity of the patterns. There is of course a danger that in doing so I will have conveyed a misleading impression!

There are several classes of states which I have intentionally omitted. First, I have omitted states for which there is very weak evidence and no motivation; I assume that I will be excused for this. At the other extreme, I have omitted some states which are unambiguous, but which I believe are best interpreted within an entirely different framework. Let me begin with some *gedanken* states of this type: $\bar{p}p$ Coulombic bound states exist, but we should clearly omit them from the "elementary" meson spectrum. There are a number of states which are widely (but by no means universally) believed to be of this type—I'll call them "molecular states" since the hadrons themselves are already "quark atoms"—which I have excluded. Foremost amongst the candidates for such states are the old and well-established $J^{PC} = 0^{++}$ states $a_0(980)$ and $f_0(975)$ [formerly the S^* and δ]. Jaffe first suggested[74] that these states were bagged $qq\bar{q}\bar{q}$ states. Recent opinion has shifted to the closely related "$K\bar{K}$ molecule" picture[75] advocated by Weinstein *et al.*, in which these states have much more in common with the deuteron than the proton, and are as appropriately distinguished from "true" mesons as the deuteron is distinguished from a "true" $B = 2$ baryon. There are two other interesting candidates for such states which have been excluded from Figure 13. One is the narrow $f_1(1420)$ which was long interpreted as a member of the ${}^3P_1$ multiplet, but which now might be interpreted[76,77] as a (virtual) bound state of $K^*\bar{K}$. New experimental results on this state were presented[78] to this conference confirming its 1^{++} assignment and thus the need for a non-quarkonium interpretation of its nature. The other is the newly established $f_2(1520)$ state seen in $\bar{p}p$ collisions, for which convincing evidence was reported at this conference by Amsler[79] and Smith[80]. This state, first reported by the Asterix collaboration and called the AX (1560), has been interpreted by some[81] as a $\bar{p}p$ bound state, and by others[77,82] as a coupled channel bound state of the $\rho\rho - \omega\omega$ system. In either case, it would be appropriately excluded from the "true mesons."

These molecular states have their own interest, of course. Originally, $qq\bar{q}\bar{q}$ states were sought as a new type of hadron. (A "true" multiquark hadron would be a tightly bound system characterized by having all interquark distances comparable). Today I believe their main interest is as a testing ground in novel situations for our very rudimentary understanding of interhadronic forces. One of the goals of strong interaction physics is to understand such forces, since they are ultimately responsible for the existence and properties of nuclei. Of course, we all know that the nucleon-nucleon force can be very well parameterized in terms of meson-exchange-generated forces. However, this parameterization is unlikely to lead to a viable underlying explanation for this and other interhadronic forces. The forces between composite systems can arise not only from particle exchange, but also from constituent exchange. (In fact, in a heavy quark world meson exchange would be negligible compared to quark exchange forces[83].) Studying the forces in exotic hadron-hadron systems can help to delineate the role of these mechanisms (and others) in the mundane nucleon-nucleon case. These two competing mechanisms have been studied by a number of groups. For example, Barnes and Swanson[77,82] systematized the search for attractive meson-meson systems in a quark exchange picture (see also Refs. 74, 75, and 82). Tornquist reported to this conference[82] his interesting results on those channels which have an attractive pion-exchange-induced potential. It is interesting that all groups find attraction in the $\rho\rho - \omega\omega$ system that could be associated with the $f_2(1520)$.

Looking at Figure 13, some may wonder if there is life in spectroscopy beyond the quark model. In two senses, I would argue that there must be. In the first place, no matter how successful the quark model might be, its origin within QCD is not understood. So at the very least there lies ahead the vital task of explaining what the quark model is. This is the task being addressed by work like that I mentioned earlier on the nature of the constituent quark. However, there is more than just theoretical mopping up to be done: we *know* that glueballs and hybrids must exist as degrees of freedom of QCD so there *must* be life beyond the simple pattern of this spectrum. In fact, the absence of evidence for such states in the low energy spectrum seems likely to be a central clue to the existence of the quark model. Moreover, it may well be that as we discover these states experimentally, we will also see the quark model picture break down in such a way that it will be revealed to us what it was in the first place.

I should remark that I may be overstating the present situation. I have already mentioned how the search for glueballs in the 0^{++} sector is hampered by quarkonium chaos there. (For a good discussion, see the review by Close at this conference[84].) However, the GAMS collaboration has reported[85] an interesting candidate state at 1590 MeV which we heard more about at this conference[79]. Given the coexistence in this region of broad $^3P_0$ states, the probable existence of a "$K\bar{K}$ molecule" with these quantum numbers, and the possibility of maximal OZI-violation in this channel[86], much work will be needed to confirm the existence of a glueball in this mass region. Another contender for a glueball used to be the $\theta(1720)$ seen in ψ radiative decays. With the recent evidence that its quantum numbers are really 0^{++} and not 2^{++}, it is no longer such a serious candidate: it could now easily be the first radial excitation of the quarkonium scalar mesons[68]. For either set of quantum numbers, it could also be a $K^*\bar{K}^* - \omega\phi$ molecule[87]. A more interesting and persistent candidate has been proposed in the tensor sector in the 2.0–2.5 GeV mass range by Lindenbaum *et al.*[88]. This group has found a series of $\phi\phi$ resonances produced by $\pi\pi_{\mathrm{virtual}}$ scattering which would be difficult to accommodate as quarkonia. Given the OZI-violating nature of their production, there is also a *prima facie* case for associating these states with glue. Further study of these states is certainly warranted.

As I have already argued, it may be easier ultimately to establish the excitation of the gluonic degree of freedom by finding J^{PC} exotic hybrid mesons (or baryons). Some candidates have been reported to this conference[89], but so far there are no clear resonant signals. On the other hand, the lattice adiabatic potentials described in the $Q\bar{Q}$ section, plus the assumption of smooth evolution of quarkonium spectroscopy into the light quark sector, suggest where to look.

OUTLOOK

It is plausible that hadron spectroscopy, after several years of slow progress, is poised to make significant advances:

- heavy-light systems provide us with the "hydrogen atom" of brown muck, and in so doing severely constrain model building; at the same time our theoretical and experimental knowledge of such heavy-light states is increasing rapidly,
- There are reasons to believe that hybrids, and possibly glueballs, will be discovered in the next generation of hadron spectroscopy experiments; by watching the quark model fail, we should be able to understand better why it works so well, and, more generally,
- new accelerators and detectors, wielded by both high energy and nuclear physicists, will shed new light on this now old but still important problem.

REFERENCES

1. B.A. Thacker and G.P. Lepage, Phys. Rev. **D43**, 196 (1991).
2. See plenary talk PL-12 by R. Petronzio.
3. R. Barbieri, R. Gatto, and E. Remiddi, Phys. Lett. **61B**, 465 (1976); R. Barbieri, M. Caffo, and E. Remiddi, Nucl. Phys. **B162**, 220 (1980).
4. G. Bodwin, E. Braaten, and G.P. Lepage, "Rigorous QCD Predictions for Decays of P-Wave Charmonium", session PA-4B.
5. N. Isgur and J. Paton, Phys. Rev. **D31**, 2910 (1985).
6. T. Barnes, Z. Phys. **C10**, 275 (1981); T. Barnes, F.E. Close, and S. Monaghan, Nucl. Phys. **B198**, 380 (1982); M. Chanowitz and S. Sharpe, Nucl. Phys. **B222**, 211 (1983).
7. S. Perantonis and C. Michael, Nucl. Phys. **B347**, 854 (1990).
8. C. Olson, M.G. Olsson, and K. Williams, Phys. Rev. **D45**, 4307 (1992) and "QCD, Relativistic Flux Tubes, and Potential Models", paper 344.
9. E. Eichten and G. Feinberg, Phys. Rev. **D23**, 2724 (1981).
10. D. Gromes, Z. Phys. **C22**, 265 (1984); **C26**, 401 (1984).
11. W. Buchmüller, Phys. Lett. **112B**, 479 (1982).
12. J. Rosen for the Fermilab E760 collaboration, "Charmonium Results from E760", session PA-4B.
13. F. Halzen, C. Olson, M.G. Olsson, and M.L. Strong, "Mass of the Singlet P State", paper 419.
14. E. Altshuler and D. Silverman, "Predictability of Charmonium Levels for a Range of Fits", paper 428.
15. G. Kernel for the ARGUS collaboration, "Resonance Production by Two Photons", session PA-3B.
16. S. Nussinov and W. Wetzel, Phys. Rev. **D36**, 130 (1987); M.B. Voloshin and M.A. Shifman, Sov. J. Nucl. Phys. **47**, 511 (1988).
17. M.B. Voloshin and M.A. Shifman, Sov. J. Nucl. Phys. **45**, 292 (1987).
18. G.P. Lepage and B.A. Thacker, Nucl. Phys. **B4** (Proc. Suppl.) 199 (1988).
19. E. Eichten, Nucl. Phys. **B4** (Proc. Suppl.) 170 (1988).

20. H.D. Politzer and M.B. Wise, Phys. Lett. **B206**, 681 (1988); **B208**, 504 (1988).

21. N. Isgur and M.B. Wise, Phys. Lett. **B232**, 113 (1989); Phys. Lett. **B237**, 527 (1990).

22. E. Eichten and B. Hill, Phys. Lett. **B234**, 511 (1990).

23. H. Georgi, Phys. Lett. **B240**, 447 (1990).

24. B. Grinstein, Nucl. Phys. **B240**, 447 (1990); **B339**, 253 (1990).

25. J.D. Bjorken, in *Proceedings of the 4th Rencontres de Physique de la Vallee d'Aoste*, La Thuile, Italy, 1990, ed. M. Greco (Editions Frontieres, Gif-sur-Yvette, France, 1990).

26. N. Isgur and M.B. Wise, "Heavy Quark Symmetry", in *B Decays*, ed. S. Stone (World Scientific, Singapore, 1992), p. 158.

27. M. Suzuki, Nucl. Phys. **B258**, 553 (1985); B. Grinstein, M.B. Wise, and N. Isgur, Phys. Rev. Lett. **56**, 298 (1986); T. Altomari and L. Wolfenstein, Phys. Rev. Lett. **58**, 1583 (1987); N. Isgur, D. Scora, B. Grinstein, and M.B. Wise, Phys. Rev. **D39**, 799 (1989).

28. E.V. Shuryak, Phys. Lett. **93B**, 134 (1980); Nucl. Phys. **B198**, 83, (1982); Nucl. Phys. **B328**, 85 (1989).

29. N. Isgur and M.B. Wise, Phys. Rev. Lett. **66**, 1130 (1991).

30. J. Rosner, Comm. Nucl. Part. Phys. **16**, 109 (1986).

31. N. Isgur and M.B. Wise, Nucl. Phys. **B348**, 278 (1991).

32. H. Georgi, Nucl. Phys. **B348**, 293 (1991).

33. N. Isgur, Phys. Rev. **D43**, 810 (1991).

34. J.C. Anjos *et al.*, Phys. Rev. Lett. **62**, 722 (1989); **62**, 1587 (1989); Z. Bai *et al.*, SLAC-PUB-5341 (1990); D. Potter for the Fermilab E653 collaboration, "Charm Meson Production in 600 GeV/c π-Emulsion Interactions", session PA-10A; J. Spengler for the ARGUS collaboration, "New Results in Υ and Charm Physics from ARGUS", session PA-4B; G. Bellini for the Fermilab E678 collaboration "Semileptonic Charm Decays", session PA-1A.

35. A.F. Falk, H. Georgi, B. Grinstein, and M.B. Wise, Nucl. Phys. **B343**, 1 (1990); A. Falk and B. Grinstein, Phys. Lett. **B247**, 406 (1990); X. Ji and M. Musolf, Phys. Lett. **B257**, 409 (1991); G.P. Korchemsky and A.V. Radyushkin, Nucl. Phys. **B283**, 342 (1987); Phys. Lett. **B279**, 359 (1992); M. Neubert, Nucl. Phys. **B371**, 149 (1992).

36. See also B. Grinstein, "Heavy Quark Effective Theory: Application to Weak Decays", session PA-1B.

37. M. Neubert, SLAC preprint SLAC-PUB-5842, 1992; see also T. Mannel, "Leading and Subleading Logarithmic QCD Corrections to Bilinear Heavy Quark Currents", session PA-4A.

38. M. Luke, Phys. Lett. **B252**, 447 (1990).

39. B. Grinstein and P.F. Mende, "Heavy Mesons in Two Dimensions", paper 614.

40. M. Burkhardt and E.S. Swanson, to appear in Phys. Rev. **D**, (1992).

41. Y.L. Wu, K. Schilcher, and D. Pirjol, "QCD Sum Rule Calculation of the $B \to D\ell\nu$ Form Factor", paper 434.

42. J. Taron, "Constraints on Heavy Meson Form Factors", session PA-4A.

43. A.V. Radyushkin, Phys. Lett. **B271**, 218 (1991).

44. M. Neubert, Phys. Rev. **D45**, 2451 (1992).

45. E. Bagan, P. Ball, V.M. Braun, and H.G. Dosch, Phys. Lett. **B278**, 452 (1992).

46. N. Isgur, D. Scora, B. Grinstein, and M.B. Wise, Phys. Rev. **D39**, 799 (1989).

47. M.B. Wise, Phys. Rev. **D45**, R2188 (1992).

48. T.-M. Yan *et al.*, Phys. Rev. **D46**, 1148 (1992).

49. G. Burdman and J.F. Donoghue, Phys. Lett. **B280**, 287 (1992).

50. H. Cheng, "Heavy Quark Symmetry and Chiral Dynamics", session PA-4B; J. Goity, "Chiral Perturbation Theory for SU(3) Breaking in Heavy Mesons", submitted paper.

51. M.J. Dugan and B. Grinstein, Phys. Lett. **B255**, 583 (1991).

52. J. Stiewe for the ARGUS collaboration, "Production and Decays from ARGUS", session PA-10B.

53. H. Cheung for the Fermilab experiment E687 collaboration, "Charm Baryon Production", session PA-10B.

54. A. Biagi *et al.*, Z. Phys. **C28**, 175 (1985).

55. J. Spengler for the ARGUS collaboration, "New Results in Υ and Charm Physics from ARGUS", session PA-4B.

56. See plenary talks PL-1 by P. Drell and PL-3 by J. Butler.

57. The OPAL collaboration, "Evidence for the Existence of the Strange b-Flavoured Meson B_s in Z^0 Decays", submitted paper.

58. The CLEO collaboration, "Measurement of the $D^{*+} - D^+$ and $D^{*0} - D^0$ Mass Differences", paper 900; "Observation of the Charmed Baryon Σ_c^+", paper 901.

59. V. Jain for the CLEO collaboration, "Results on Charm and Bottom", session PA-4B.

60. J. Goity, private communication.

61. M. Bauer, B. Stech, and W. Wirbel, Z. Phys. **C34**, 103 (1987); M. Wirbel, Prog. in Nucl. Part. Phys. **22**, 33 (1988).

62. The ARGUS collaboration, "A Measurement of Asymmetry in the Decay $\Lambda_c^+ \to \Lambda\pi$", paper 278; "A Partial Wave Analysis of $D^0 \to K_S^0\pi^+\pi^-$", paper 105.

63. See, for example, C. Michael and M. Teper, Nucl. Phys. **B314**, 347 (1989).

64. A. Manohar and H. Georgi, Nucl. Phys. **B234**, 189 (1984); H. Georgi and A. Manohar, Nucl. Phys. **B310**, 527 (1988).

65. D.B. Kaplan, Phys. Lett. **B235**, 163 (1990); Nucl. Phys. **B351**, 137 (1991).

66. S. Weinberg, Phys. Rev. Lett. **65**, 1181 (1990); **67**, 3473 (1991); S. Peris, Phys. Rev. **D46**, 1202 (1992).

67. M. Karliner, "Quark Solitons as Constituents of Hadrons", session PA-6B.

68. For an explicit model of this type, see S. Godfrey and N. Isgur, Phys. Rev. **D32**, 189 (1985).

69. Zheng Zhipeng, "Results on $f_0(975)$ and $\xi(2.2)$ from BES", session PA-3B.

70. B. Ratcliff for the LASS collaboration, "The $K_2(1770)$ Region: Evidence for Two $J^P = 2^-$ Strange Mesons Decaying to $K\omega$, session PA-3B.

71. E.S. Ackleh and T. Barnes, Phys. Rev. **D45**, 232 (1992).

72. The ARGUS collaboration, "Analysis of the Hadronic Final State of the Decay $\tau \to \pi^-\pi^-\pi^+\nu_\tau$", paper 252.

73. R. Kokoski and N. Isgur, Phys. Rev. **D35**, 907 (1987); N. Isgur, C. Morningstar, and C. Reader, Phys. Rev. **D39**, 1357 (1989).

74. R.L. Jaffe, Phys. Rev. **D15**, 267, 281 (1977); **D17**, 1444 (1978); R.L. Jaffe and K. Johnson, Phys. Lett. **60B**, 2105 (1976).

75. J. Weinstein and N. Isgur, Phys. Rev. Lett. **48**, 659 (1982); Phys. Rev. **D27**, 588 (1983); **D41**, 2236 (1990).

76. D.O. Caldwell, Mod. Phys. Lett. **A2**, 771 (1987).

77. E.S. Swanson, "Intermeson Potentials from the Constituent Quark Model", to appear in Ann. Phys. (NY), (1992).

78. F. Votruba for the CERN WA76 collaboration, "A Study of the $E/f_1(1420)$ and of the $f_1(1285)$ in Central Production", session PA-3A.

79. C. Amsler for the Crystal Barrel collaboration, "New Light Mesons from the LEAR Crystal Barrel", session PA-3A.

80. G. Smith for the Fermilab E760 collaboration, "Neutral Light Quark Spectroscopy in $\bar{p}p$ Annihilation at $3000 < \sqrt{s} < 6000$ MeV", session PA-3A.

81. C.B. Dover, "Quasinuclear $N\bar{N}$ States", in proceedings of *Intersections between Particle and Nuclear Physics*, AIP Conf. Proc. 243, ed. W.T.H. van Oers (AIP, New York, 1992), p. 381.

82. Yu. S. Kalashnikova, "Four Quark Interpretation of AX", in proceedings of *Hadron '91*, ed. S. Oneda, and D.C. Peaslee (World Scientific, Singapore, 1992), p. 777.; N. Tornquist, Phys. Rev. Lett. **67**, 556 (1991) and paper 702; T. Barnes and E.S. Swanson, Phys. Rev. **D46**, 131 (1992).

83. N. Isgur, "Quark Exchange Forces from a Heavy Quark Perspective", in proceedings of *From Fundamental Fields to Nuclear Phenomena*, ed. J.A. McNeil and C.E. Price (World Scientific, Singapore, 1991), p. 46.

84. F.E. Close, "Scalars and Light Quark Mesons", session PA-3B.

85. D. Alde *et al.*, Phys. Lett. **B201**, 160 (1988).

86. P. Geiger and N. Isgur, "When Can Hadronic Loops Scuttle the OZI Rule?", CEBAF preprint CEBAF-TH-92-24.

87. K. Dooley, E.S. Swanson, and T. Barnes, Phys. Lett. **B275**, 478 (1992).

88. S.J. Lindenbaum, "The Status of Investigations of 0^{++} and 2^{++} Glueball States and New Searches for Exotic Hybrid and Glueball States", paper 249, sheduled for session PA-3A.

89. Y. Guoz for the VES collaboration, "Study of the Wave with $J^{PC} = 1^{-+}$ ... in π^--Nucleon Interactions at $p_\pi = 37$ GeV/c", session PA-3B.