Light-meson spectroscopy

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The authors survey the current status of light-meson spectroscopy, beginning with a general introduction to meson spectroscopy and its importance in understanding the physical states of quantum chromodynamics. Phenomenological models of hadron spectroscopy are described with particular emphasis on the constituent-quark model and the qualitative features it predicts for the meson spectrum. The authors next discuss expectations for hadrons lying outside the quark model, such as hadron states with excited gluonic degrees of freedom. These states include so-called hybrids and glueballs, as well as multiquark states. The established meson states are compared to the quark-model predictions and most meson states are found to be well described by the quark model. However, a number of states in the light-quark sector do not fit in well, suggesting the existence of hadronic states with additional degrees of freedom. The review ends with a brief description of future directions in meson spectroscopy. [S0034-6861(99)00805-3]

CONTENTS

I. INTRODUCTION

Meson physics and the strong interactions have been intimately connected since pions were first introduced by Yukawa to explain the internucleon force (Yukawa, 1935). Since that time, our knowledge of mesons and, in parallel, our understanding of the strong interactions have undergone several major revisions. Our present understanding of the strong interactions is that they are described by the non-Abelian gauge-field theory quantum chromodynamics (QCD) (Fritzch *et al.*, 1971; Gross and Wilczek, 1973; Weinberg, 1973), which describes the interactions of quarks and gluons. It appears that mesons are the ideal subjects for the study of strong interactions in the strongly coupled nonperturbative regime. Even though in QCD we have a theory of the strong interactions, we know very little about the physical states of the theory. Until we can both predict the properties of the physical states of the theory and confirm these predictions by experiment we can hardly claim to understand QCD, which has implications beyond hadron physics. For example, it is possible that at high energies the weak interactions become strong, so that strongly interacting field theories may be relevant to the mechanism of electroweak symmetry breaking. In QCD we have an example of such a theory where we can test our understanding against experiment. The study of meson spectroscopy will hopefully elucidate this theory.

To a large extent our knowledge of hadron physics is based on phenomenological models, in particular, the constituent-quark model (Gell-Mann, 1964; Zweig, 1964).[1] Meson and baryon spectroscopy is described surprisingly well as composite objects made of constituent objects—valence quarks. We shall refer to these hadrons, described by only valence-quark configurations, as "conventional." Most QCD-motivated models, however, predict other types of hadrons with explicit glue degrees of freedom. These are the *glueballs*, which have no constituent quarks in them at all and are entirely described in terms of gluonic fields, and *hybrids*, which have both constituent quarks and excited gluon degrees of freedom.[2] It is the prospect of these new forms of hadronic matter that has led to continued excitement among hadron spectroscopists.

To be able to understand the nature of new resonances it is important that we have a template against which to compare observed states with theoretical predictions. The constituent-quark model offers the most complete description of hadron properties and is probably the most successful phenomenological model of hadron structure. But to use it as a template to find new physics, it is very important that we test the quark model against known states to understand its strengths and weaknesses. At one extreme, if we find too large a discrepancy with experiment, we may decide that it is not such a good model after all, and we should start over again. On the other hand, if it gives general agreement with experiment, discrepancies may indicate the need for new physics, either because approximations to the model are not appropriate, or because we are dealing with new types of hadrons that cannot be explained by the quark model. To understand our reliance on this very simple, and perhaps naive, model it is useful to look at the historical evolution of our understanding of hadron physics.

Mesons were first introduced by Yukawa (1935) with pions acting as the exchange bosons responsible for the strong interactions between nucleons. With the advent of higher-energy accelerators, a whole zoo of mesons and baryons appeared, leading to great confusion. Eventually, when the various mesons and baryons were arranged into multiplets according to their quantum numbers, patterns started to emerge. It was recognized that hadrons of a given J^{PC} arranged themselves into representations of the group $SU(3)$ (Gell-Mann, 1961; Ne'eman, 1961), although none of the observed states seemed to correspond to the fundamental triplet representation.[3] In an important conceptual leap, Zweig (1964) and Gell-Mann (1964) postulated that mesons and baryons were in fact composite objects, with mesons made of a quark-antiquark pair and baryons made of three quarks. Zweig referred to these constituent spin-$\frac{1}{2}$ fermions as aces and Gell-Mann referred to them as quarks. This simple picture explained the qualitative properties of hadrons quite well.[4] Serious problems remained, however. In the "naive" quark model, the spin-$\frac{3}{2}$ baryons, the constituent quarks' spin wave functions were symmetric, as were their flavor wave functions. Being based upon fermions, the baryon wave function should be antisymmetric in the quark quantum numbers. This would imply that either quarks obey some sort of bizzare parastatistics or that the ground-state spatial wave function is antisymmetric. Yet no reasonable models could be constructed to give this result. To avoid this result, Greenberg postulated that quarks had another quantum number, which was later named color, with respect to which the quark wave functions could be antisymmetrized (Greenberg, 1964). The serious shortcoming of this model was that no quarks were observed. Most physicists took the view that if they could not be observed they were nothing more than a convenient bookkeeping device.

In addition to the spectroscopic evidence for hadronic constituents, in the late 1960s deep-inelastic scattering experiments (Bloom *et al.*, 1969; Breidenbach *et al.*, 1969) analogous to Rutherford's scattering experiment suggested that the proton was composed of constituents (Bjorken, 1967; Bjorken and Paschos, 1969; Feynman, 1969).

By the beginning of the 1970s it was becoming clear that the weak interactions could be explained by gauge theories (Glashow, 1961; Weinberg, 1967; Salam, 1968).

[1]Introductions to the quark model are given by Isgur (1980) and Rosner (1981).

[2]Some recent reviews on this subject are given by Close (1988), Godfrey (1989), and Isgur (1989).

[3]In fact, Sakata (1956) postulated that all hadrons could be constructed from the triplet p,n,Λ and their antiparticles but this model led to spurious baryon states.

[4]A proper history of the quark model requires a review in itself. Lacking the room to do justice to this subject we direct the interested reader to the monograph by Kokkedee (1969), which includes reprints of early quark-model papers. A historical perspective from a personal point of view is given by Zweig (1980).

If this was the case, it seemed reasonable that the strong interactions should also be described using the same formalism. "Gauging" the color degree of freedom leads to quantum chromodynamics, a non-Abelian gauge theory based on the group $SU(3)$, as the theory of the strong interactions (Fritzch *et al.*, 1971; Gross and Wilczek, 1973; Weinberg, 1973).

Nevertheless, there was still considerable skepticism about the existence of quarks since they had never been directly seen. This situation changed when, in November 1974, the discovery of very narrow hadron resonances was announced simultaneously at Brookhaven National Laboratory (Aubert *et al.*, 1974) and the Stanford Linear Accelerator Center (Augustin *et al.*, 1974). These states, named J/ψ, and later the ψ', χ_c, and η_c, were quickly interpreted as being bound states of a new heavy quark—the charm quark. Quark models that incorporated the qualitative features of QCD (asymptotic freedom and confinement) were able to reproduce the charmonium ($c\overline{c}$) spectrum rather well (Appelquist *et al.*, 1975; Appelquist and Politzer, 1975a, 1975b; Eichten *et al.*, 1975). These developments, both experimental and theoretical, convinced all but a few that quarks were real objects and were the building blocks of hadronic matter. In a seminal paper on the subject, De Rújula, Georgi, and Glashow (1975) showed that these ideas could successfully be used to describe the phenomenology of light-quark spectroscopy.

With the acceptance of QCD as the theory of the strong interactions comes the need to understand its physical states. Interpretation of the spectrum of hadrons reveals information on the nonperturbative aspects of QCD. Unfortunately, calculating the properties of hadrons from the QCD Lagrangian has proven to be a very difficult task in this strongly coupled nonlinear theory. In the long term, the most promising technique is formulating the theory on a discrete space-time lattice (Kogut 1979, 1983; Creutz, 1983; Creutz *et al.*, 1983; Montvay and Munster, 1994). By constructing interpolating fields with the quantum numbers of physical hadrons and evaluating their correlations on the lattice, one is able to calculate hadron properties from first principles. Although a great deal of progress has been made, it has been slow since these calculations take enormous amounts of computer time. Additionally, a disadvantage of this approach is that one may obtain numerical results without any corresponding physical insight.

A less rigorous approach, which has proven to be quite useful and reasonably successful, has been to use phenomenological models of hadron structure to describe hadron properties. These models predict, in addition to the conventional $q\overline{q}$ mesons and qqq baryons of the quark model, multiquark states, glueballs, and hybrids. Probably the most pressing question in hadron spectroscopy is whether these states do in fact exist and what their properties are. However, the predictions of the various models can differ appreciably so that experiment is needed to point the model builders in the right direction.

We shall often refer to glueballs and hybrids as exotics because they lie outside the constituent-quark model. However, this is somewhat misleading as they are not exotics in the sense that if they do exist, they are simply additional hadronic states expected from QCD. Nevertheless, from the historical development of the field we see that the quark model provides a good framework on which to base further study. If we find discrepancies everywhere it obviously fails and we should abandon it as a tool. On the other hand, since it does work reasonably well, it gives us a basis on which to decide if we have discovered the new forms of hadronic matter we are interested in, namely, glueballs and hybrids.

The present situation in light-meson spectroscopy is that the constituent-quark model works surprisingly well in describing most observed states. At the same time there are still many problems and puzzles that need to be understood and that might signal physics beyond the quark model. Although most QCD-based models expect glueballs and hybrids and there is mounting evidence that some have been found, thus far no observed state has unambiguously been identified as one. The best candidates are states with "exotic quantum numbers," that is, states that cannot be formed in the quark model. Part of the problem and confusion is that the conventional mesons are not understood well enough to rule out new states as conventional states, and part is that these exotics may have properties that have made them difficult to find up to now. Despite these qualifications, there has been considerable recent progress in understanding the properties of exotic mesons that could help distinguish them from conventional mesons. With sufficient evidence, a strong case can be made to label an experimentally observed state as an exotic hadron. Thus, to have any hope of distinguishing between conventional and exotic mesons, it is crucial that we understand conventional-meson spectroscopy very well.

The purpose of this review is to summarize the present status of meson spectroscopy and identify puzzles, perhaps pointing out measurements that could help resolve them. To this end we shall begin with a discussion of the theoretical ingredients relevant to this article. In the course of this review we shall refer to numerous experiments. In Sec. III we briefly survey relevant experiments along with the attributes that contribute useful information to the study of mesons. Since the eventual goal is to identify discrepancies between the observed meson spectrum and conventional quark-model predictions, in Sec. IV we shall compare the predictions of one specific quark model with experiment. This will allow us to identify discrepancies between the quark model and experiment that may signal physics beyond conventional hadron spectroscopy. In Sec. V we shall go over these puzzles in detail to help decide whether the discrepancy is most likely a problem with the model or the experiment, or whether it most likely signals some interesting new physics. In Sec. VI we shall briefly outline some future facilities for the study of meson spectroscopy that are under construction or that are being considered. Finally, in Sec. VII we shall attempt to summarize our most interesting findings. Our hope is that the reader will see that meson spectroscopy is a vibrant field.

Because of the breadth of this review we can only touch upon many interesting topics. There are a number of recent reviews of meson spectroscopy and related topics with emphases somewhat different from that of the present one. We strongly encourage the interested reader to refer to these reviews for further details.[5]

II. THEORETICAL OVERVIEW

A. Quantum chromodynamics

Quantum chromodynamics (QCD), the theory of the strong interactions (Fritzch *et al.*, 1971; Gross and Wilczek, 1973; Weinberg, 1973), may be thought of as a generalization of quantum electrodynamics (QED), our most successful physical theory. QCD is described by the Lagrangian

$$\mathcal{L}_{QCD}=\overline{q}_i\left(i\partial_\mu \gamma^\mu \delta_{ij} - g\frac{\lambda^a_{ij}}{2}A^a_\mu \gamma^\mu - m\,\delta_{ij}\right) q_j - \frac{1}{4}F^a_{\mu\nu}F^{a\mu\nu}, \tag{1}$$

where

$$F^{\mu\nu}_a = \partial^\mu A^\nu_a - \partial^\nu A^\mu_a - g f_{abc} A^\mu_b A^\nu_c, \tag{2}$$

A^μ_a are the gluon fields, which transform according to the adjoint representation of $SU(3)$ with $a=1,\ldots,8$, q_i are the quark fields with color indices $i=1,2,3$, g is the bare coupling, m is the quark mass, and $\lambda^i/2$ are the generators of $SU(3)$. One immediately observes that quarks couple to gluons in much the same way as electrons couple to photons with the $e\gamma^\mu$ of QED replaced by $g\gamma^\mu(\lambda/2)$ of QCD. The significant difference between QED and QCD is that in QCD the quarks come in colored triplets and the gluons in a color octet, where color is labeled by the Latin subscripts. The non-Abelian group structure of $SU(3)$ leads to nonlinear terms in the field strength $F^{\mu\nu}$, which give rise to trilinear and quadratic vertices in the theory, so that gluons couple to themselves in addition to interacting with quarks. This makes the theory nonlinear and very difficult to solve and leads to the confinement of color. A consequence of this behavior appears to be the existence of new forms of hadronic matter with excited gluonic degrees of freedom known as glueballs and hybrids (Close, 1988).

Because of the difficulties in solving QCD exactly to obtain the properties of the physical states of the theory, we have resorted to various approximation methods. The most promising of these is to redefine the problem on a discrete space-time lattice, in analogy to the approach one might take in the numerical solution of a difficult differential equation (Wilson, 1974; Kogut, 1979, 1983; Creutz, 1983, Creutz *et al.*, 1983; Montvay and Muenster, 1994; Michael, 1995, 1997). For QCD, one formulates the problem in terms of the path integral in Euclidean space-time and evaluates expectation values of the appropriate operators using a Monte Carlo integration over the field configurations. Although advances are being made on the problem, it requires enormous computer capacity so that progress is slow in making precise, detailed predictions of the properties of the physical states of the theory. As a consequence, our understanding of hadrons continues to rely on insights obtained from experiment and QCD-motivated models in addition to lattice QCD results.

In later sections we shall use the predictions of QCD-inspired models as the basis for interpreting the nature of the observed mesons.[6] It is therefore useful to sketch the QCD motivation for these models. We start with the quark-antiquark ($Q\overline{Q}$) potential in the limit of infinitely massive quarks, which can be used in the Schrödinger equation to obtain the spectroscopy of heavy quarkonium. This is analogous to the adiabatic potentials for diatomic molecules in molecular physics, with the heavy quarks corresponding to slow-moving nuclei and the gluonic fields corresponding to fast-moving electrons. The $Q\overline{Q}$ potential is found by calculating the energy of a fixed, infinitely heavy, quark-antiquark pair given by the expectation value of what is known as the Wilson loop operator (Wilson, 1974). The resulting potential is referred to as the static potential since the massive quarks do not move. Limiting cases of the static potential are given by

$$V(r)=br \quad \text{for} \quad r \gg \frac{1}{\Lambda}, \tag{3}$$

where Λ is a parameter that sets the QCD scale with $\Lambda \sim 200$ MeV, and the constant b, the "string tension," is numerically found to be $b \simeq 0.18\ \mathrm{GeV}^2 \simeq 0.9$ GeV/fm and

$$V(r) \sim \frac{-4}{3}\frac{\alpha_s}{r} \quad \text{for} \quad r \ll \frac{1}{\Lambda}, \tag{4}$$

where α_s is the strong coupling. The result of one such lattice calculation is shown in Fig. 1, taken from Bali *et al.* (1997). The lattice potential $V(r)$ can then be used to determine the spectrum of $b\overline{b}$ mesons by solving the Schrödinger equation since the b-quark motion is approximately nonrelativistic.

There are also spin-dependent forces between the quarks analogous to the spin-dependent forces in the hydrogen atom that give rise to the fine and hyperfine

[5]These reviews include Close (1988), Cooper (1988), Diekmann (1988), Burnett and Sharpe (1990), Törnqvist (1990), Amsler and Myhrer (1991), Königsmann (1991), Landau (1996), Amsler (1998), and Barnes (1998). In addition, the Review of Particle Physics (Caso, 1998) contains a wealth of information on the properties of mesons in its tables of properties and minireviews on topics of special interest and should be consulted for further information.

[6]For a recent review see Barnes (1996).

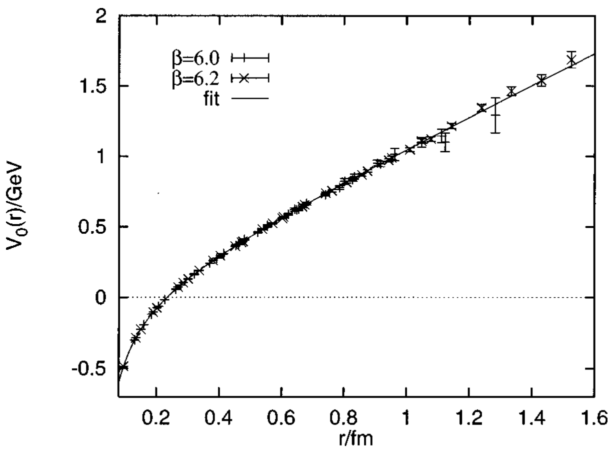

FIG. 1. The static $Q\overline{Q}$ potential from Bali, Schilling, and Wachter (1997). β is related to the inverse of the QCD coupling.

structure in atomic spectroscopy. To obtain the spin-dependent potentials in QCD the Wilson loop is expanded in the inverse quark mass, which gives the next-order terms in an expansion in v^2/c^2 (Eichten and Feinberg, 1979, 1981; Gromes, 1984a, 1984b):

$$V_{spin}(r)=\left(\frac{1}{2m_1^2}\vec{L}\cdot\vec{S}_1+\frac{1}{2m_2^2}\vec{L}\cdot\vec{S}_2\right)\frac{1}{r}\frac{d}{dr}[V(r)+2V_1(r)]+\frac{1}{m_1m_2}\vec{L}\cdot(\vec{S}_1+\vec{S}_2)\frac{1}{r}\frac{dV_2(r)}{dr}+\frac{1}{m_1m_2}\left(\hat{r}\cdot\vec{S}_1\hat{r}\cdot\vec{S}_2-\frac{1}{3}\vec{S}_1\cdot\vec{S}_2\right)V_3(r)+\frac{1}{3m_1m_2}\vec{S}_1\cdot\vec{S}_2V_4(r), \tag{5}$$

where $\vec{S}_1$, $\vec{S}_2$, and $\vec{L}$ are the quark and antiquark spins and relative orbital angular momentum and $V(r)$ is the interquark potential defined by the Wilson loop operator. The spin-dependent potentials, V_1, V_2, V_3, and V_4, can be related to correlation functions of the color electric and color magnetic fields. For example, V_3 and V_4 are given by (Eichten and Feinberg, 1979, 1981; Gromes, 1984a, 1984b)

$$\left(\frac{r_ir_j}{r^2}-\frac{1}{3}\delta_{ij}\right)V_3(r)+\frac{1}{3}\delta_{ij}V_4(r)=\lim_{T\to\infty}\frac{g^2}{T}\int_0^T\frac{dtdt'\langle B_i(0,t)B_j(r,t')\rangle}{\langle 1\rangle}, \tag{6}$$

which can be evaluated using nonperturbative techniques, in particular, lattice QCD. The lattice results can be compared to the phenomenological expectations that the magnetic correlations are short range, as expected from one-gluon exchange:

$$V_3(r)=\frac{4\alpha_s}{r^3} \quad \text{and} \quad V_4(r)=\frac{8\pi}{3}\alpha_s\delta^3(r). \tag{7}$$

The lattice results do indeed agree with these spin-dependent potentials (Campostrini *et al.*, 1986, 1987; Bali *et al.*, 1997). Bali, Schilling, and Wachter (1997) have studied the $b\overline{b}$ and $c\overline{c}$ spectra using the potentials calculated using lattice QCD. They solved the Schrödinger equation using the spin-independent potential and treated the spin-dependent and other relativistic corrections (not discussed here) as perturbations. The resulting beautyonium spectrum shown in Fig. 2, is found to be in reasonable agreement with the experimental Υ spectrum. The main deviations between experiment and prediction are due to the quenched approximation, which neglects internal quark loops in the lattice calculation, and the neglect of higher-order relativistic corrections. Direct lattice calculations of spin-dependent splittings also agree with the measured splittings (Davies, 1998). For heavy quarkonium an exact expression has been obtained for the interaction up to order $(v/c)^2$ using a technique known as nonrelativistic QCD (for a recent review see Brambilla, 1998).

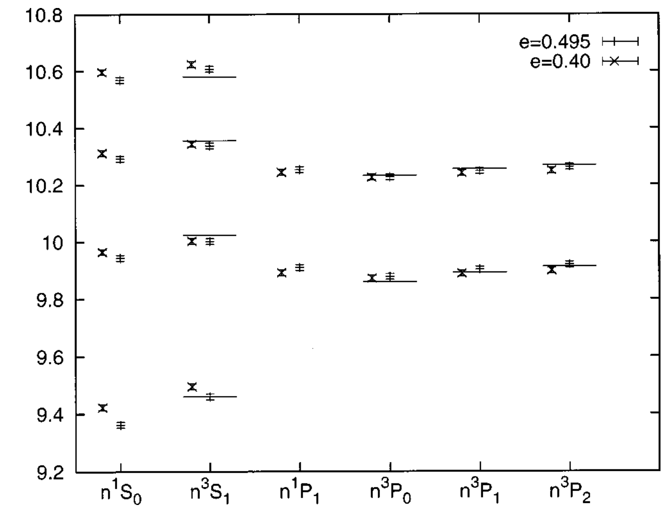

FIG. 2. The bottomonium spectrum (in GeV) calculated using potentials from lattice QCD. The horizontal lines are experimental results. From Bali, Schilling, and Wachter (1996). e is the coefficient of the $1/r$ piece of the $Q\overline{Q}$ potential obtained from a fit to the lattice results. The different values correspond to different approximations in the lattice calculation.

Although historically the spin-dependent potentials were obtained phenomenologically by comparing the observed quarkonium spectra ($c\overline{c}$ and $b\overline{b}$) to the predictions of potential models, it turns out that the resulting potentials are in reasonable agreement with those obtained from lattice QCD and nonrelativistic QCD. The phenomenological spin-dependent potential typically assumes a Lorentz vector one-gluon exchange for the short-distance piece, which results in terms analogous to the Breit-Fermi Hamiltonian in atomic physics, and a Lorentz scalar linear confining piece. The resulting spin-dependent Hamiltonian is then of the form

$$H_{spin}=H_{ij}^{hyp}+H_{ij}^{s.o.(cm)}+H_{ij}^{s.o.(tp)}, \tag{8}$$

where

$$H_{ij}^{hyp}=\frac{4\alpha_s(r)}{3m_im_j}\left\{\frac{8\pi}{3}\vec{S}_i\cdot\vec{S}_j\delta^3(\vec{r}_{ij})+\frac{1}{r_{ij}^3}\left[\frac{3\vec{S}_i\cdot\vec{r}_{ij}\vec{S}_j\cdot\vec{r}_{ij}}{r_{ij}^2}-\vec{S}_i\cdot\vec{S}_j\right]\right\} \tag{9}$$

is the color hyperfine interaction,

$$H_{ij}^{s.o.(cm)}=\frac{4\alpha_s(r)}{3r_{ij}^3}\left(\frac{1}{m_i}+\frac{1}{m_j}\right)\left(\frac{\vec{S}_i}{m_i}+\frac{\vec{S}_j}{m_j}\right)\cdot\vec{L} \tag{10}$$

is the spin-orbit color magnetic piece arising from the one-gluon exchange, and

$$H_{ij}^{s.o.(tp)}=\frac{-1}{2r_{ij}}\frac{\partial V(r)}{\partial r_{ij}}\left(\frac{\vec{S}_i}{m_i^2}+\frac{\vec{S}_j}{m_j^2}\right)\cdot\vec{L} \tag{11}$$

is the spin-orbit Thomas precession term where $V(r)$ is the interquark potential given by the Wilson loop. Note that the contribution arising from one-gluon exchange is of opposite sign to the contribution from the confining potential. $\alpha_s(r)$ is the running coupling constant of QCD.

B. Color singlets in QCD

Because of confinement, only color singlet objects can exist as physical hadrons. Colored quarks form the fundamental triplet (3) representation of the $SU(3)$ color gauge group and antiquarks the conjugate $\overline{3}$ representation. Therefore a quark-antiquark pair can combine to form a color singlet as can three quarks, while a quark-quark pair cannot. Other states are also possible, for example $q\overline{q}q\overline{q}$, and it is a dynamical question whether such multiquark systems are realized in nature as single multiquark states, as two distinct $q\overline{q}$ states, or as a loosely associated system of color singlet mesons analogous to a diatomic molecule. Color singlets can also be constructed with gluons (g). Glueballs are hadrons with no valence-quark content and hybrids are made up of valence quarks and antiquarks and an explicit gluon degree of freedom. Of course, life is not so simple, and one expects the physical mesons to be linear combinations of the various Fock-space components: $q\overline{q}$, $q\overline{q}q\overline{q}$, gg, $q\overline{q}g$, etc.

C. The constituent-quark model

In the constituent-quark model, conventional mesons are bound states of a spin-$\frac{1}{2}$ quark and a spin-$\frac{1}{2}$ antiquark bound by a QCD-motivated phenomenological potential such as the one described above. The quark and antiquark spins combine into a spin singlet or triplet with total spin $S=0$ or 1, respectively. S is coupled to the orbital angular momentum L resulting in total angular momentum $J=L$ for the singlet state and $J=L-1$, L, $L+1$ for the triplet states. In spectroscopic notation the resulting state is denoted by $n^{2S+1}L_J$ with S for $L=0$, P for $L=1$, D for $L=2$, and F, G, H, for $L=3,4,5$, etc. Parity is given by $P(q\overline{q},L)=(-1)^{L+1}$ and C parity is also defined for neutral self-conjugate mesons and is given by $C(q\overline{q},L,S)=(-1)^{L+S}$. Thus the ground-state vector meson with $J^{PC}=1^{--}$ is the 1^3S_1 quark-model state.

The light-quark quarkonia are composed of u, d, or s quarks. Since the u and d quarks are quite similar in mass (~5–10 MeV, which is much smaller than the intrinsic mass scale of QCD), it is convenient to treat them as members of an "isosopin" doublet with the resulting $SU(2)$ isospin an approximate symmetry of the strong interactions. Combining u,d and $\overline{u},\overline{d}$ into mesons forms isospin singlet and triplet multiplets. We shall use the symbol n (for nonstrangeness) generically to stand for u or d. Thus one should read

$$n\overline{n}=(u\overline{u}\pm d\overline{d})/\sqrt{2} \tag{12}$$

with + for the isoscalar mesons and − for the neutral member of the isovector multiplet.

When dealing with the charged members of an isospin triplet, it is customary to refer to their C parity as the C parity of the neutral member of the multiplet C_n. It is convenient, however, to introduce a new quantum number, $G\equiv C_n(-1)^I=\pm 1$. The so-called G parity is defined, and has the same value, for all members of the multiplet. It is important to note, however, that unlike C parity, G parity is *not* an exact symmetry of the strong interaction because of the inherently approximate nature of isospin.

Hadrons containing s quarks have similar properties to the (u,d) systems so that mesons are arranged into $SU(3)$ flavor nonets: three isovector states $(-u\overline{d},u\overline{u}-d\overline{d},d\overline{u})$, two isoscalar states $(u\overline{u}+d\overline{d},s\overline{s})$ and four strange $I=1/2$ states $(u\overline{s},d\overline{s},-s\overline{d},s\overline{u})$. With the heavier strange-quark mass, the $s\overline{s}$ isoscalar states are sufficiently heavier than the (u,d) $q\overline{q}$ states that there is little mixing between $s\overline{s}$ and the light $n\overline{n}$ states, with the exception of the $\eta-\eta'$ system, where

$$|\eta\rangle=\cos(\phi)|n\overline{n}\rangle-\sin(\phi)|s\overline{s}\rangle, \tag{13}$$

$$|\eta'\rangle=\sin(\phi)|n\overline{n}\rangle+\cos(\phi)|s\overline{s}\rangle, \tag{14}$$

where the flavor mixing angle $\phi\simeq 45°$. These states are often also expressed as linear combinations of flavor $SU(3)$ octet and singlet states,

$$|\eta\rangle=\cos(\theta)|8\rangle-\sin(\theta)|1\rangle, \tag{15}$$

$$|\eta'\rangle=\sin(\theta)|8\rangle+\cos(\theta)|1\rangle. \tag{16}$$

The two angles are trivially related by $\phi=\theta+\tan^{-1}(\sqrt{2})$.

Combining the spin and orbital angular momentum wave functions with the quark flavor wave functions results in the meson states of Table I, where we have used the Particle Data Group's naming conventions (Caso *et al.*, 1998). States not fitting into this picture are considered to be "exotics." Thus a meson with $J^{PC}=1^{-+}$ would be forbidden in the constituent-quark model, as would a doubly charged meson m^{++}.

To obtain the meson spectrum one solves for the eigenvalues of the Schrödinger equation with a $q\overline{q}$ poten-

TABLE I. The quantum numbers and names of conventional $q\bar{q}$ mesons. (Note that η and η' are linear combinations of $n\bar{n}$ and $s\bar{s}$.)

		J^{PC}	$I=1$	$I=0$ $(n\bar{n})$	$I=0$ $s\bar{s}$	Strange
$L=0$	$S=0$	0^{-+}	π	η	η'	K
	$S=1$	1^{--}	ρ	ω	ϕ	K^*
$L=1$	$S=0$	1^{+-}	b_1	h	h'	K_1
	$S=1$	0^{++}	a_0	f_0	f_0'	K_0^*
		1^{++}	a_1	f_1	f_1'	K_1
		2^{++}	a_2	f_2	f_2'	K_2^*
$L=2$	$S=0$	2^{-+}	π_2	η_2	η_2'	K_2
	$S=1$	1^{--}	ρ	ω	ϕ	K_1^*
		2^{--}	ρ_2	ω_2	ϕ_2	K_2
		3^{--}	ρ_3	ω_3	ϕ_3	K_3^*
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tial, including the spin-dependent potentials, and tunes the constituent-quark masses to give agreement with experiment. There is nothing fundamental about the values assigned to the constituent-quark masses; specific values are chosen simply to improve the predictions of the model. The relative positioning of the multiplets, i.e., the $1S$,$1P$,$2S$,$1D$,$2P$, . . . , levels, is sensitive to the details of the potential. Figure 3 shows the $b\bar{b}$ spectrum predicted by one representative model which gives reasonably good agreement with experiment. The phenomenological, QCD-motivated, linear-plus-Coulomb potential gives the observed multiplet positioning and is consistent with the lattice potential described above. The spin-dependent potentials split the multiplets by giving the spectrum "fine" and "hyperfine" structure analogous to their counterparts in QED and reproduce the observed $b\bar{b}$ spectrum quite well.

The preceding discussion may seem like a lengthy digression, but it demonstrates an important result: phenomenological models starting with experimental measurements, and lattice calculations starting from the underlying theory, find a common ground in the language of potential models. That phenomenological models of heavy quarkonium work well and agree, at least qualitatively, with the potentials predicted by quantum chromodynamics using lattice QCD, is strong support for this approach, at least for heavy-quark systems. The success for heavy quarkonia begs the question about whether we can extend potential models to light-quark systems where the use of the static potential is questionable.

In Fig. 4, we show the evolution of the $1^3P_2-1^3S_1$ and the $1^3S_1-1^1S_0$ splittings as a function of quark masses (Godfrey and Isgur, 1985). In heavy-quark systems the former splitting is a nonrelativistic orbital excitation analogous to the Lyman α line in hydrogen, while the latter is a Breit-Fermi (order p/m) relativistic correction analogous to the 21-cm line of hydrogen. One can see that there is a smooth evolution going from the heavy $b\bar{b}$ system, where we believe potential models to be approximately valid, to the relativistic light-quark systems. We take this as evidence that, qualitatively, the basic structure in heavy and light systems are identical, the main difference being that in the light-quark systems

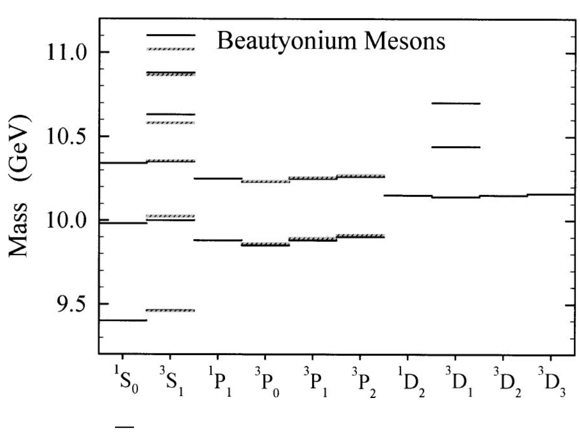

FIG. 3. The $b\bar{b}$ spectrum from a quark potential model (Godfrey and Isgur, 1985). The solid lines are the quark-model predictions and the shaded regions are the experimental states.

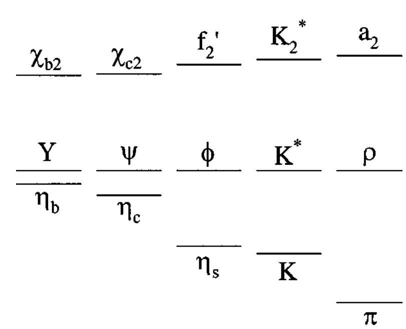

FIG. 4. The evolution of the $1^3P_2-1^3S_1$ and the $1^3S_1-1^1S_0$ splittings as a function of quark masses (Godfrey and Isgur, 1985). The splittings are drawn to scale. Note that the η_b and the η_s are in fact calculations.

the relativistic splittings are comparable to the orbital splittings. Thus to describe the light-quark hadrons we should include relativistic effects and the characteristics expected from QCD (Capstick *et al.*, 1986). Ideally this would be done by deriving the correct relativistic equations from QCD and solving them. A less ambitious approach is to model the exact equations by including various parameters. In the following sections, to interpret the spectrum of mesons with light-quark content, we shall use the predictions of one such attempt at a "relativized" quark model (Godfrey and Isgur, 1985, 1986; Godfrey, 1985b), which we take to be representative of the many similar models in the literature (Stanley and Robson, 1980; Carlson *et al.*, 1983a, 1983b; Gupta *et al.*, 1986; Brayshaw, 1987; Crater and Van Alstine, 1988; Olson *et al.*, 1992; Fulcher *et al.*, 1993; Jean *et al.*, 1994).

In quark models, mesons are approximated by the valence-quark sector of the Fock space, in effect integrating out the degrees of freedom below some distance scale μ. This results in an effective potential $V_{q\bar{q}}(\vec{p},\vec{r})$ whose dynamics are governed by a Lorentz vector one-gluon-exchange interaction at short distance and a linear Lorentz scalar confining interaction. In the relativized quark model, mesons are described by the relativistic rest-frame Schrödinger-type equation

$$H|\Psi\rangle=[H_0+V_{q\bar{q}}(\vec{p},\vec{r})+H_A]|\Psi\rangle=E|\Psi\rangle, \tag{17}$$

where H_0 is the kinetic-energy operator, H_A is the annihilation amplitude, which must be included for self-conjugate mesons where $q\bar{q}$ annihilation via gluons can contribute to the masses, and $V_{q\bar{q}}(\vec{p},\vec{r})$ is the effective quark-antiquark potential, which is found by equating the scattering amplitude of free quarks with the potential $V_{\rm eff}$ between bound quarks inside a meson (Gromes, 1984b). To first order in $(v/c)^2$, $V_{q\bar{q}}(\vec{p},\vec{r})$ reduces to the standard nonrelativistic result:

$$V_{q\bar{q}}(\vec{p},\vec{r})\rightarrow V(\vec{r}_{ij})=H_{ij}^{conf}+H_{ij}^{hyp}+H_{ij}^{s.o.}, \tag{18}$$

where

$$H_{ij}^{conf}=-\frac{4}{3}\frac{\alpha_s(r)}{r}+br+C \tag{19}$$

includes the spin-independent linear confinement and Coulomb-like interaction, and H_{ij}^{hyp} and $H_{ij}^{s.o.}$ are given by Eqs. (9)–(11).

In the confinement potential, Eq. (19), the Coulomb piece dominates at short distance while the linear piece dominates at large distance. Because heavy quarkonia have smaller radii they are more sensitive to the short-range color-Coulomb interaction while the light quarks, especially the orbitally excited mesons, have larger radii and are more sensitive to the long-range confining interaction. Thus measurements of both heavy quarkonia and mesons with light-quark content probe different regions of the confinement potential and complement each other. The linear character of the Regge trajectories[7] of the orbital excitations is a direct consequence of the linear confining potential, so that the experimental measurement of these masses is a measure of the slope of the potential and will give information about the nature of confinement (Godfrey, 1985b).

The spin-dependent parts of the potential consist of the color hyperfine interaction [Eq. (9)] and the spin-orbit interaction [Eqs. (10) and (11)]. The color hyperfine interaction is responsible for $^3S_1-{}^1S_0$ splitting in $\rho-\pi$, K^*-K, and $J/\Psi-\eta_c$. In addition to multiplet splitting, the tensor term can cause mixings between states with the same quantum numbers related by $\Delta L=2$, such as $^3S_1$ and $^3D_1$. The spin-orbit terms contribute to the splitting of the $L\neq 0$ multiplets. For states with unequal-mass quark and antiquark, where C parity and G parity are no longer good quantum numbers, the spin-orbit terms can also contribute to $^1L_J-{}^3L_J$ mixing. The spin-orbit interaction has two contributions, $H_{ij}^{s.o.(cm)}$ and $H_{ij}^{s.o.(tp)}$, given by Eqs. (10) and (11). Since the hyperfine term is relatively short distance, it becomes less important for larger radii, so that multiplet splittings become a measure of the spin-orbit splittings, with contributions of opposite sign coming from the short-range Lorentz vector one-gluon exchange and the long-range Lorentz scalar linear confinement potential. The ordering of states within a multiplet of given orbital angular momentum gives information on the relative importance of the two pieces (Schnitzer, 1984a, 1984b; Godfrey, 1985a, 1985b; Isgur, 1998). Thus the multiplet splittings act as a probe of the confinement potential, providing information on nonperturbative QCD.

D. Meson decays

While the quark potential model makes mass predictions for $q\bar{q}$ mesons, the couplings of these states are sensitive to the details of the meson wave functions and consequently provide an important test of our understanding of the internal structure of these states. Knowledge of expected decay modes is also useful for meson searches; comparing the observed decay properties of mesons to the expectations of different interpretations, $q\bar{q}$ vs hybrid, for example, is an important means of determining what they are. Thus the strong, electromagnetic, and weak couplings of mesons can give important clues to the nature of an observed state.

As a consequence, a successful model of strong decays would be a very useful tool in determining the nature of observed resonances. A large number of models exist in the literature. In an important subset of models, a quark-antiquark pair materializes and combines with the quark and antiquark of the original meson to form two new mesons. This process is described in Fig. 5(a). The models differ in the details of how the quark pair creation process occurs. In one variant the $q\bar{q}$ pair originates in an intermediate gluon and is therefore formed in a $^3S_1$ state with $J^{PC}=1^{--}$, while in the other major variant, the $^3P_0$ model of LeYaouanc *et al.* (1973, 1974, 1975), the quark pair creation process is viewed as an inherently nonperturbative process in which the $q\bar{q}$ pair is formed with the quantum numbers of the vacuum,

[7]Plots of meson spin J vs their mass squared.

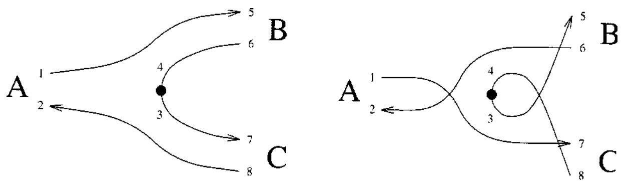

FIG. 5. Diagrams contributing to the meson decay $A \rightarrow BC$: (a) OZI allowed, (b) OZI suppressed.

$J^{PC}=0^{++}$, and therefore is in a $^3P_0$ state. Geiger and Swanson (1994) performed a detailed study of these models and concluded that the decay-width predictions of the $^3P_0$ model give better agreement with experiment. A variation of this model is the flux-tube-breaking model of Kokoski and Isgur (1987), which assumes that the $q\bar{q}$ pair is most likely to be created in the region between the original quark and antiquark. In practice the predictions of this variation do not differ significantly from those of the $^3P_0$ model, as the overlap of the original and final-state mesons is greatest in this central region. Detailed decay predictions, which can be compared with experiment are given by Kokoski and Isgur (1987), Blundell and Godfrey (1996), and Barnes, Close, Page, and Swanson (1997). Comparing the partial decay widths of non-$q\bar{q}$ candidates to quark-model predictions provides an important tool for understanding the nature of observed resonances when we discuss candidates for non-$q\bar{q}$ "exotic" mesons.

In addition to the aforementioned strong decays described by the connected quark lines of Fig. 5(a) there are strong decays that can be represented by the disconnected quark diagrams of Fig. 5(b). The decays proceeding via the disconnected diagrams are significantly suppressed compared to the connected diagrams. This is known as the OZI rule (Okubo, 1963; Zweig, 1964; Iizuka, 1966), and the two processes are referred to as OZI allowed and OZI suppressed, respectively. The breaking of OZI suppression is considered to be a strong signal for physics outside the quark model, such as glueball production.

Electromagnetic couplings are another source of useful information about resonances. An example is two-photon couplings, which are measured in the reaction $e^+e^- \rightarrow e^+e^- +$hadrons. In the cross section to the final state $\pi^0\pi^0$, for example, the $f_2(1270)$ can be seen as a clear bump. From the cross section, the two-photon width times the branching fraction, $\Gamma_{\gamma\gamma}(f_2)\cdot B(f_2 \rightarrow \pi^0\pi^0)$, can be determined. In principle the absolute two-photon widths can be calculated from quark-model wave functions (Godfrey and Isgur, 1985), but for light-$q\bar{q}$ mesons the results are sensitive to relativistic effects (Ackleh and Barnes, 1992). To some extent this sensitivity can be evaded when testing for $q\bar{q}$ candidates by comparing the relative rates of the possible members of a u,d,s multiplet with the same J^{PC}. The decay amplitude for $\gamma\gamma$ couplings involves the charge matrix element of two electromagnetic vertices,

$$A(q\bar{q} \rightarrow \gamma\gamma) \propto \langle q\bar{q}|e_q^2|0\rangle. \tag{20}$$

This gives, for example, the relative amplitudes

$$\langle f{:}a{:}f'|e_q^2|0\rangle = \frac{(2/3)^2+(-1/3)^2}{\sqrt{2}} : \frac{(2/3)^2-(-1/3)^2}{\sqrt{2}} : (-1/3)^2, \tag{21}$$

which result in the relative $\gamma\gamma$ decay rates for $I=0{:}I=1{:}s\bar{s}$ mesons of the same state, neglecting phase-space differences, of

$$\Gamma_{\gamma\gamma}(f{:}a{:}f')=25{:}9{:}2. \tag{22}$$

This agrees reasonably well with measurements (Caso *et al.*, 1998) of $\Gamma_{\gamma\gamma}(f_2)/\Gamma_{\gamma\gamma}(a_2)$ and agrees qualitatively for $\Gamma_{\gamma\gamma}(f_2')$, which is the lowest-lying $J^{PC}=2^{++}$ nonet, i.e., the $f_2(1270)$, $a_2(1320)$, and the $f_2'(1525)$. The discrepancy for $\Gamma_{\gamma\gamma}(f_2')$ can be attributed to its sensitivity to its $n\bar{n}$ and $s\bar{s}$ content. The $L=2$ η_2 state has probably been observed in two-photon production (Karch *et al.*, 1992). Because glueballs do not have valence-quark content, the observation or nonobservation of a state in two-photon production provides another clue about the nature of an observed state.

Measurement of single-photon transitions $(q\bar{q})_i \rightarrow \gamma(q\bar{q})_f$ is also useful for identifying $q\bar{q}$ states. These have the characteristic pattern of rates based on flavor,

$$\Gamma[(q\bar{q})_i \rightarrow \gamma(q\bar{q}_f)]=9{:}4{:}1 \tag{23}$$

for $\Delta I=1{:}s\bar{s}{:}\Delta I=0$. Although measurements of radiative transitions would be useful for the classification of higher-$q\bar{q}$ and non-$q\bar{q}$ states, only two such transitions have been measured, $a_2 \rightarrow \pi\gamma$ and $a_1 \rightarrow \pi\gamma$. Since measurements of radiative transitions could determine the nature of controversial states such as the $f_1(1420)$, which will be discussed in Sec. V, they should be carried out if possible. These measurements could be made in electroproduction, which is the inverse reaction.

E. Mesons with gluonic excitations

In addition to conventional hadrons it is expected that other forms of hadronic matter exist with excited gluonic degrees of freedom—glueballs, which are made primarily of gluons, and hybrids, which have both valence quarks and gluonic degrees of freedom (Jaffe and Johnson, 1976; Chanowitz and Sharpe, 1983a, 1983b; Barnes, 1984, 1985a; Close, 1988; Godfrey, 1989; Isgur, 1989).

1. Glueballs

Many different QCD-based models and calculations make predictions for such states: bag models (Barnes, 1977; Hasenfratz *et al.*, 1980; Barnes *et al.*, 1983; Chanowitz and Sharpe, 1983a; DeTar and Donoghue, 1983), constituent-glue models (Horn and Mandula, 1978), flux-tube models (Isgur and Paton, 1983, 1985; Isgur *et al.*, 1985), QCD sum rules (Latorre *et al.*, 1984), and lattice gauge theory. Recent lattice QCD calculations are converging towards agreement (Michael and Teper, 1989; Schierholz, 1989; Bali *et al.*, 1993; Sexton *et al.*, 1995b; Teper, 1997; Morningstar and Peardon, 1999), although there is still some variation between the calcula-

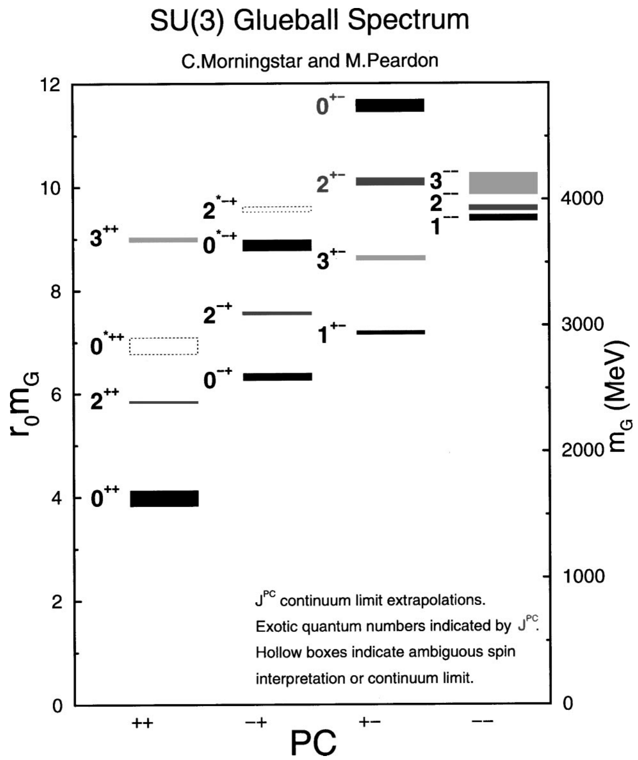

FIG. 6. The mass of the glueball states. The scale is set by r_0 with $1/r_0=410(20)$ MeV. From Morningstar and Peardon, 1999.

tions. We expect that ultimately the lattice results will be the most relevant, since they originate from QCD.

Lattice QCD predictions for glueball masses from one representative calculation are shown in Fig. 6. One should be cautioned that these results are in the so-called quenched approximation, which neglects internal quark loops. The lightest glueball is found to be a 0^{++} state with the following masses from the different collaborations: 1550 ± 50 MeV (Bali *et al.*, 1993), 1600 ± 160 MeV (Michael, 1998), 1648 ± 58 (Lee and Weingarten, 1998b), and 1730 ± 100 MeV (Morningstar and Peardon, 1999). The difference between these results lies mainly in how the mass scale is set. The next-lightest states are the 2^{++} with mass estimates of 2232 ± 220 MeV (Michael, 1998), 2270 ± 100 MeV (Bali *et al.*, 1993), 2359 ±128 MeV (Chen *et al.*, 1994), and 2400 ± 120 MeV (Morningstar and Peardon 1999), and the 0^{-+} state with a similar mass. Mixings with $q\bar{q}$ and $q\bar{q}q\bar{q}$ states could modify these predictions.

We concentrate on glueballs with conventional quantum numbers, since the first glueballs with exotic quantum numbers, "oddballs," $(0^{+-},2^{+-},1^{-+})$ do not appear until ~3 GeV. Because the lowest glueballs have conventional quantum numbers with masses situated in a dense background of conventional $q\bar{q}$ states it is difficult to distinguish them from conventional mesons. It is therefore a painstaking process to identify a glueball by comparing a candidate's properties to the expected properties of glueballs and conventional mesons.

Sexton *et al.* (1996) have estimated the width of the 0^{++} glueball to decay to all possible pseudoscalar pairs to be 108 ± 29 MeV. This number combined with reasonable guesses for the effects of finite lattice spacing, finite lattice volume, and the remaining width of multibody states yields a total width small enough for the lightest scalar glueball to be easily observed if the partial width to two pseudoscalar mesons is not too small. A significant property of glueball decays is their expected flavor-symmetric coupling to final-state hadrons. This gives the characteristic flavor singlet branching fraction for pseudoscalar pairs (factoring out phase space)

$$\Gamma(G\rightarrow\pi\pi:K\bar{K}:\eta\eta:\eta\eta':\eta'\eta')/(\text{phase space}) =3:4:1:0:1. \tag{24}$$

Of course, one should also expect some modifications from phase space, the glueball wave function, and the decay mechanism (Sexton *et al.*, 1996).

Measurements of electromagnetic couplings to glueball candidates would be extremely useful for the clarification of the nature of these states. The radiative transition rates of a relatively pure glueball would be anomalous relative to the expectations for a conventional $q\bar{q}$ state. Similarly, a glueball should have suppressed couplings to $\gamma\gamma$. The former could be measured in electroproduction experiments at, say, an energy-upgraded CEBAF, while the latter would be possible at B factories or a tau-charm factory.

There are three production mechanisms that are considered optimal for finding glueballs. The first is the radiative decay $J/\Psi\rightarrow\gamma G$, where the glueball is formed from intermediate gluons (Cakir and Farrar, 1994; Close, Farrar, and Li, 1997). The second is in central production $pp\rightarrow p_f(G)p_s$ away from the quark beams and target, where glueballs are produced from pomerons, which are believed to be multigluon objects. The third is in proton-antiproton annihilation, where the destruction of quarks can lead to the creation of glueballs. Because gluons do not carry electric charge, glueball production should be suppressed in $\gamma\gamma$ collisions. By comparing two-photon widths to J/ψ production of a state Chanowitz created a measure of glue content he calls "stickiness" (Chanowitz, 1984):

$$S=\frac{\Gamma(J/\psi\rightarrow\gamma X)}{PS(J/\psi\rightarrow\gamma X)}\times\frac{PS(\gamma\gamma\rightarrow X)}{\Gamma(\gamma\gamma\rightarrow X)}, \tag{25}$$

where PS denotes phase space. A large value of S reflects enhanced glue content. The idea of stickiness has been further developed by Close, Farrar, and Li (1997).

The simple picture presented above is likely to be muddied by complicated mixing effects between the pure glueball and $q\bar{q}$ states with the same J^{PC} quantum numbers (Amsler and Close, 1995, 1996). Lee and Weingarten (1998a, 1998b) have calculated the mixing energy between the lightest $q\bar{q}$ scalar state and the lightest scalar glueball in the continuum limit of the quenched approximation on the lattice. With this motivation, they performed a phenomenological fit that found the $f_0(1710)$ to be $\sim74\%$ glueball and the $f_0(1500)$ to be $\sim98\%$ quarkonium, mainly $s\bar{s}$. Although these results are not rigorous, they do remind us that physical states are most likely mixtures of underlying components with the same quantum numbers. As we shall see in subsequent sections, mixings can significantly alter the properties of the underlying states, which makes the interpretation of observed states difficult and often controversial.

2. Hybrids

Given the discussion of the previous subsection, the conventional wisdom is that it would be more fruitful to search for low-mass hybrid mesons with exotic quantum numbers than to search for glueballs. Hybrids have the additional attraction that, unlike glueballs, they span complete flavor nonets and hence provide many possibilities for experimental detection. In addition, the lightest hybrid multiplet includes at least one J^{PC} exotic. The phenomenological properties of hybrids have been reviewed elsewhere (Barnes, 1984, 1985a, 1996; Chanowitz, 1987; Close, 1988; Godfrey, 1989; Barnes *et al.*, 1995; Close and Page, 1995a; Page, 1997c). In this section we briefly summarize hybrid properties, such as quantum numbers, masses, and decays, which may help in their discovery.

In searching for hybrids there are two ways of distinguishing them from conventional states. One approach is to look for an excess of observed states over the number predicted by the quark model. The drawback to this method is that it depends on a good understanding of hadron spectroscopy in a mass region that is still rather murky; the experimental situation is sufficiently unsettled that the phenomenological models have yet to be tested to the extent that a given state can be reliably ruled out as a conventional meson. The situation is further muddied by expected mixing between conventional $q\bar{q}$ states and hybrids with the same J^{PC} quantum numbers. The other approach is to search for quantum numbers that cannot be accommodated in the quark model. The discovery of exotic quantum numbers would be irrefutable evidence of something new.

To enumerate the hybrid J^{PC} quantum numbers in a model-independent manner obeying gauge invariance, one forms gauge-invariant operators (Barnes, 1985b; Jaffe *et al.*, 1986) from a color octet $q\bar{q}$ operator and a gluon field strength. The resulting lowest-lying $q\bar{q}g$ states with exotic quantum numbers not present in the constituent-quark model have $J^{PC}=2^{+-}$, 1^{-+}, 0^{+-}, and 0^{--}. We label these states with the same symbol as the conventional meson with all the same quantum numbers except for the C parity and add a hat to the symbol. For example, an isospin 1 $J^{PC}=0^{--}$ meson would be a $\hat{\pi}$, an isospin 0 $J^{PC}=0^{--}$ meson would be an $\hat{\eta}$ or $\hat{\eta}'$, an isospin 1 $J^{PC}=1^{-+}$ meson would be a $\hat{\rho}$, an isospin 1 $J^{PC}=2^{+-}$ an $\hat{a}_2$, etc. The discovery of mesons with these exotic quantum numbers would unambiguously signal hadron spectroscopy beyond the quark model.

To gain some physical insight into hybrids, it is useful to turn to lattice results in the heavy-quark limit before turning to predictions of specific models and calculations. A useful approach is to use the leading Born-Oppenheimer approximation to map out the adiabatic surfaces corresponding to the non-ground-state gluon configurations (Griffiths *et al.*, 1983; Perantonis and Michael, 1990; Morningstar and Peardon, 1997; Manke *et al.*, 1998; Juge *et al.*, 1999). This is analogous to the calculation of the nucleus-nucleus potential in diatomic molecules, with the slow heavy quarks and fast gluon fields in hybrids corresponding to the nuclei and electrons, respectively, of the diatomic molecules. One treats the quark and antiquark as spatially fixed color sources and determines the energy levels of the glue as a function of the $Q\bar{Q}$ separation. Each of these energy levels defines an adiabatic potential $V_{Q\bar{Q}}(r)$. The ground-state potential has cylindrical symmetry about the interquark axis while less symmetric configurations correspond to excitations of the gluonic flux joining the quark-antiquark pair. For example, the lowest-lying gluonic excitation corresponds to a component of angular momentum of one unit along the quark-antiquark axis. The adiabatic potentials are determined using lattice QCD. One such set of adiabatic surfaces is shown in Fig. 7. The quark motion is then restored by solving the Schrödinger equation in each of these potentials.

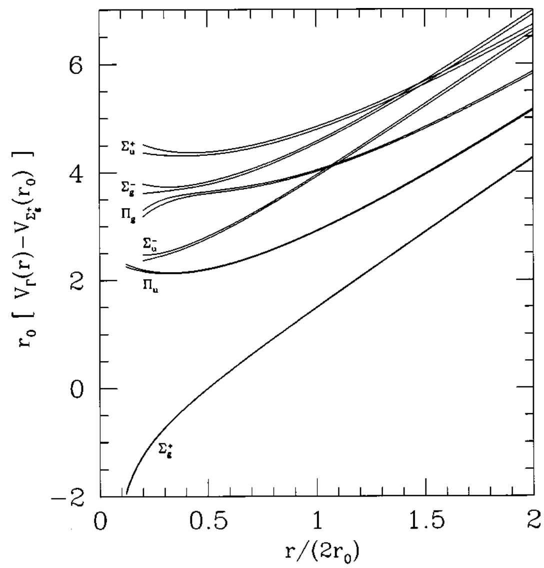

FIG. 7. A set of hybrid adiabatic surfaces for static central potentials. $\Lambda=\Sigma,\Pi,\Delta,\ldots$ corresponds to the magnitude of $J_{glue}=0,1,2,\ldots$ projected onto the molecular axis. The superscript $=\pm$ corresponds to the evenness or oddness under reflections in a plane containing the molecular axis, and the subscript u/g corresponds to odd/even charge conjugation plus spatial inversion about the midpoint. The familiar $q\bar{q}$ potential is labeled as Σ_g^+ and the first-excited potential is the Π_u, so the lowest-lying hybrid mesons should be based on this potential. The double lines on the excited surfaces indicate the calculational uncertainty in determining the potential. From Juge *et al.*, 1998.

Conventional mesons are based on the lowest-lying potential, and hybrid states emerge from the excited potentials. Combining the resulting flux-tube spatial wave functions, which have $L^{PC}=1^{+-}$ and 1^{--}, with the quark and antiquark spins yields a set of eight degenerate hybrid states with $J^{PC}=1^{--}$, 0^{-+}, 1^{-+}, 2^{-+}, and 1^{++}, 0^{+-}, 1^{+-}, 2^{+-}, respectively. These contain the J^{PC} exotics with $J^{PC}=1^{-+}$, 0^{+-}, and 2^{+-}. The degeneracy of the eight J^{PC} states is expected to be broken by the different excitation energies in the $L^{PC}=1^{+-}$ (magnetic) and 1^{-+} (pseudoelectric) gluonic excitations, spin-orbit terms, and mixing between hybrid states and $q\bar{q}$ mesons with nonexotic spins.

While this picture is appropriate for heavy quarkonium, it is not at all clear that it can be applied to light-quark hybrids. Nevertheless, given that the constituent-quark model works so well for light quarks, it is not unreasonable to extend this flux-tube description to light quarks. The flux-tube model, developed by Isgur and Paton, is based on the strong-coupling expansion of lattice QCD (Isgur and Paton, 1983, 1985; Isgur *et al.*, 1985). It predicts eight nearly degenerate nonets around 2 GeV—$J^{PC}=2^{\pm\mp}$, $1^{\pm\mp}$, $0^{\pm\mp}$, and $1^{\pm\pm}$. In the flux-tube model, the glue degree of freedom manifests itself as excited phonon modes of the flux tube connecting the $q\bar{q}$ pair, so the first excited state is doubly degenerate. This picture of gluonic excitations appears to be supported by lattice calculations (Perantonis and Michael, 1990; Michael and Stephenson, 1994).

Other models exist, for example the bag model (Barnes and Close, 1983a, 1983b; Barnes *et al.*, 1983; Chanowitz and Sharpe, 1983a, 1983b), which, in contrast to the flux-tube model, expects only the four nonets, 2^{-+}, 1^{-+}, 1^{--}, and 0^{-+} to be similar in mass, while the other four are expected to have considerably higher masses. This difference is symptomatic of the differences between the two models. In the bag model, the gluon degrees of freedom are either transverse electric or transverse magnetic modes of the bag, with the transverse magnetic mode considerably higher in mass than the transverse electric mode.

A recent Hamiltonian Monte Carlo study of the flux-tube model (Barnes *et al.*, 1995) finds the lightest $n\bar{n}$ hybrid masses to be 1.8–1.9 GeV. This result is consistent with the lattice QCD (quenched approximation) results of the UKQCD Collaboration (Perantonis and Michael, 1990; Lacock *et al.*, 1996, 1997, 1998), who find $M_{\hat{\rho}}\simeq 1.88$ GeV and $M_{\hat{\phi}}\simeq 2.09$ GeV. It is also consistent with the results of the MILC Collaboration (Bernard *et al.*, 1996, 1997), who find $M_{\hat{\rho}}\simeq 1.97$ GeV and $M_{\hat{\phi}}\simeq 2.17$ GeV. Lacock *et al.* (1996) find $M_{\hat{a}_0}\sim 2.09$ GeV and $M_{\hat{a}_2}\sim 2.09$ GeV for the next-lightest hybrids.

Hybrid decays appear to follow the almost universal selection rule that gluonic excitations cannot transfer angular momentum to the final states as relative angular momentum. Rather, the momentum must appear as internal angular momentum of the $q\bar{q}$ pairs (Kalashnikova, 1994; Page, 1997a, 1997b). The selection rule suppresses decay channels likely to be large and may make hybrids stable enough to appear as conventional resonances. Unfortunately, this selection rule is not absolute; in the flux-tube and constituent-glue models it can be broken by wave-function and relativistic effects, while the bag model adds a qualifier that it is also possible that the excited quark loses its angular momentum to orbital angular momentum. In this case the 1^{-+} could decay to two S-wave mesons such as $\pi\eta$ or $\pi\rho$ in a relative P wave, which would dominate due to the large available phase space. In any case, if we take this selection rule seriously, it explains why hybrids with exotic J^{PC} have yet to be seen; they do not couple strongly to simple final states. Thus in the list of possible $\hat{\rho}$ decays,

$$\begin{aligned}&\hat{\rho}\to[\underline{\pi\eta},\underline{\pi\eta'},\pi\rho,K^*K,\eta\rho,\ldots]_P,\\&[\pi b_1,\pi f_1,\eta a_1,KK_1\cdots]_S,\end{aligned}\tag{26}$$

most models expect the $b_1\pi$ and $f_1\pi$ modes to dominate. (The underlined modes to two distinct pseudoscalars provide a unique signature of the 1^{-+} state.) For states with conventional quantum numbers we expect mixing between hybrids and conventional $q\bar{q}$ states, even in the quenched approximation, which could significantly modify the properties of these states.

The decay predictions of the flux-tube model (Isgur *et al.*, 1985; Close and Page, 1995a) are given in Table II. These predictions suggest that many hybrids are too broad to be distinguished as a resonance, while a few hybrids should be narrow enough to be easily observable. In particular, of the hybrids with exotic quantum numbers, the flux-tube model predicts that the $\hat{a}_0$, $\hat{f}_0$, and $\hat{f}_0'$ are probably too broad to appear as resonances. The $\hat{\omega}_1$ decays mainly to $[a_1\pi]_S$ with partial width $\Gamma\approx 100$ MeV, which would make it difficult to reconstruct the original hybrid given the broad width of the final-state mesons. Similar problems could also make the $\hat{\phi}_1$ difficult to find. According to the flux-tube model, the best bets for finding exotic hybrids are

$$\begin{aligned}&\hat{\rho}_1\to[b_1\pi]_S\quad(\Gamma\approx 150\ \text{MeV}),\\&\quad\to[f_1\pi]_S\quad(\Gamma\approx 50\ \text{MeV}),\\&\hat{a}_2\to[a_2\pi]_P\quad(\Gamma\approx 200\ \text{MeV}),\\&\hat{f}_2\to[b_1\pi]_P\quad(\Gamma\approx 250\ \text{MeV}),\\&\hat{f}_2'\to[K^*(1430)_2\bar{K}]_P\quad(\Gamma\approx 90\ \text{MeV}),\\&\quad\to[\bar{K}K_1]_P\quad(\Gamma\approx 100\ \text{MeV}),\end{aligned}\tag{27}$$

where the partial widths for the specified decays are given in parentheses. Finally, some "forbidden" decays such as $\hat{\rho}(1900)\to\rho\pi$ have small but finite partial widths due to differences in the final-state spatial wave functions. Thus it may be possible to observe hybrids in these simpler decay modes in addition to the favored but more difficult to reconstruct final states such as $b_1\pi$ and K_1K.

So far we have concentrated on the J^{PC} exotic members of the lowest flux-tube hybrid multiplet; one might wonder whether the nonexotic hybrids might be narrow enough to be observable. According to the results of Close and Page (1995a) reproduced in Table II, many of the nonexotic hybrids are also so broad as to be effectively unobservable. There are several notable exceptions. The first is a 1^{--} ω hybrid with a total width of

TABLE II. The dominant hybrid decay widths for $A \to [BC]_L$ for partial-wave L calculated using the flux-tube model. From Close and Page (1995a). Hybrid masses before spin splitting for $n\bar{n}$ are 2.0 GeV except for 0^{+-} (2.3 GeV), 1^{+-} (2.15 GeV), and 2^{+-} (1.85 GeV) and for $s\bar{s}$ are 2.15 GeV except for 0^{+-} (2.25 GeV), following Merlin and Paton (1987). The partial widths are given in MeV.

	$I=1$		$I=0$ $n\bar{n}$		$I=0$ $s\bar{s}$	
A	$[BC]_L$	Γ	$[BC]_L$	Γ	$[BC]_L$	Γ
2^{-+}	$[f_2(1270)\pi]_S$	40	$[a_2(1320)\pi]_S$	125	$[K_2^*(1430)K]_S$	100
	$[f_2(1270)\pi]_D$	20	$[a_2(1320)\pi]_D$	60	$[K_1(1270)K]_D$	20
	$[b_1(1235)\pi]_D$	40	$[f_2(1270)\eta]_S$	$\sim$50		
	$[a_2(1320)\eta]_S$	$\sim$40	$[K_2^*(1430)K]_S$	$\sim$30		
	$[K_2^*(1430)K]_S$	$\sim$30				
	$[\rho\pi]_P$	8	$[K^*K]_P$	2	$[K^*K]_P$	6
	$[K^*K]_P$	2				
2^{+-}	$[a_2(1320)\pi]_P$	200	$[b_1(1235)\pi]_P$	250	$[K_2^*(1430)K]_P$	90
	$[a_1(1260)\pi]_P$	70	$[h_1(1170)\eta]_P$	30	$[K_1(1270)K]_P$	30
	$[h_1(1170)\pi]_P$	90			$[K_1(1400)K]_P$	70
	$[b_1(1235)\eta]_P$	$\sim$15				
			$[\rho\pi]_D$	1	$[K^*K]_D$	1
0^{+-}	$[a_1(1260)\pi]_P$	700	$[b_1(1235)\pi]_P$	300	$[K_1(1270)K]_P$	400
	$[h_1(1170)\pi]_P$	125	$[h_1(1170)\eta]_P$	90	$[K_1(1400)K]_P$	175
	$[b_1(1235)\eta]_P$	80	$[K_1(1270)K]_P$	600		
	$[K_1(1270)K]_P$	600	$[K_1(1400)K]_P$	150		
	$[K_1(1400)K]_P$	150				
1^{+-}	$[a_2(1320)\pi]_P$	175	$[b_1(1235)\pi]_P$	500	$[K_2^*(1430)K]_P$	70
	$[a_1(1260)\pi]_P$	90	$[h_1(1170)\eta]_P$	175	$[K_1(1270)K]_P$	250
	$[h_1(1170)\pi]_P$	175	$[K_2^*(1430)K]_P$	60	$[K_0^*(1430)K]_P$	125
	$[b_1(1235)\eta]_P$	150	$[K_1(1270)K]_P$	250		
	$[K_2^*(1430)K]_P$	60	$[K_0^*(1430)K]_P$	70		
	$[K_1(1270)K]_P$	250				
	$[K_0^*(1430)K]_P$	70				
	$[\omega\pi]_S$	15	$[\rho\pi]_S$	40	$[K^*K]_S$	20
	$[\rho\eta]_S$	20	$[\omega\eta]_S$	20	$[\phi\eta]_S$	40
	$[\rho\eta']_S$	30	$[\omega\eta']_S$	30	$[\phi\eta']_S$	40
	$[K^*K]_S$	30	$[K^*K]_S$	30		
1^{++}	$[f_2(1270)\pi]_P$	175	$[a_2(1320)\pi]_P$	500	$[K_2^*(1430)K]_P$	125
	$[f_1(1285)\pi]_P$	150	$[a_1(1260)\pi]_P$	450	$[K_1(1270)K]_P$	70
	$[f_0(1300)\pi]_P$	$\sim$20	$[f_2(1270)\eta]_P$	70	$[K_1(1400)K]_P$	100
	$[a_2(1320)\eta]_P$	50	$[f_1(1285)\eta]_P$	60		
	$[a_1(1260)\eta]_P$	90	$[K_2^*(1430)K]_P$	$\sim$20		
	$[K_2^*(1430)K]_P$	$\sim$20	$[K_1(1270)K]_P$	40		
	$[K_1(1270)K]_P$	40	$[K_1(1400)K]_P$	$\sim$20		
	$[K_1(1400)K]_P$	$\sim$20				
	$[\rho\pi]_S$	20	$[K^*K]_S$	15	$[K^*K]_S$	10
	$[K^*K]_S$	15				
1^{-+}	$[f_1(1285)\pi]_S$	40	$[a_1(1260)\pi]_S$	100	$[K_1(1270)K]_S$	40
	$[f_1(1285)\pi]_D$	20	$[a_1(1260)\pi]_D$	70	$[K_1(1270)K]_D$	60
	$[b_1(1235)\pi]_S$	150	$[f_1(1285)\eta]_S$	50	$[K_1(1400)K]_S$	25
	$[b_1(1235)\pi]_D$	20	$[K_1(1270)K]_S$	20		
	$[a_1(1260)\eta]_S$	50	$[K_1(1400)K]_S$	$\sim$125		
	$[K_1(1270)K]_S$	20				
	$[K_1(1400)K]_S$	$\sim$125				

TABLE II. (*Continued*).

	$I=1$		$I=0\ n\overline{n}$		$I=0\ s\overline{s}$	
A	$[BC]_L$	Γ	$[BC]_L$	Γ	$[BC]_L$	Γ
	$[\rho\pi]_P$	8	$[K^*K]_P$	2	$[K^*K]_P$	6
	$[K^*K]_S$	2				
0^{-+}	$[f_2(1270)\pi]_D$	20	$[a_2(1320)\pi]_D$	60	$[K_2^*(1430)K]_D$	20
	$[f_0(1300)\pi]_S$	~ 150	$[f_0(1300)\eta]_S$	~ 200	$[K_0^*(1430)K]_S$	400
	$[K_0^*(1430)K]_S$	200	$[K_0^*(1430)K]_S$	200		
	$[\rho\pi]_P$	30	$[K^*K]_P$	8	$[K^*K]_P$	30
	$[K^*K]_P$	8				
1^{--}	$[a_2(1320)\pi]_D$	50	$[K_1(1270)K]_S$	40	$[K_2^*(1430)K]_D$	20
	$[a_1(1260)\pi]_S$	150	$[K_1(1400)K]_S$	60	$[K_1(1270)K]_S$	60
	$[a_1(1260)\pi]_D$	20			$[K_1(1400)K]_S$	125
	$[K_1(1270)K]_S$	40				
	$[K_1(1400)K]_S$	~ 60				
	$[\omega\pi]_P$	8	$[\rho\pi]_P$	20	$[K^*K]_P$	15
	$[\rho\eta]_P$	7	$[\omega\eta]_P$	7	$[\phi\eta]_P$	8
	$[\rho\eta']_P$	3	$[\omega\eta']_P$	3	$[\phi\eta']_P$	2
	$[K^*K]_P$	4	$[K^*K]_P$	4		

only $\sim$100 MeV which decays to $K_1(1270)K$ and $K_1(1400)K$. The ϕ is also relatively narrow, with Γ_{tot} $\sim$225 MeV. Two more interesting hybrids are the π_2 with $\Gamma_{tot}\sim$170 MeV and its $s\overline{s}$ partner, the η_2', with $\Gamma_{tot}\sim$120 MeV decaying dominantly to K_2^*K. In addition, there are several other hybrids that have total widths around 300 MeV and so should also be observable.

To determine whether an observed state with nonexotic quantum numbers is a conventional $q\overline{q}$ state or a hybrid, one would make use of the detailed predictions we have described above for the two possibilities (Barnes *et al.*, 1997). For example, a selection rule of the $^3P_0$ decay model forbids the decay of a spin-singlet $q\overline{q}$ state to two spin-singlet mesons in the final state (Page, 1997a). This selection rule forbids the decay $\pi_2(^1D_2)\rightarrow b_1\pi$ while, in contrast, the decay is allowed for the hybrid π_2 and is in fact rather large. A second illustration is a 0^{-+} state with $M\simeq 1800$ MeV. The largest decay modes for the $\pi(3^1S_0)$ and a hybrid with the same mass and quantum numbers, π_H, are shown in Table III. Both states decay to most of the same final states, albeit with much different partial widths. A discriminator between the two possibilities is the $\rho\omega$ channel, which is dominant for $\pi(3^1S_0)$, whereas it is predicted to be absent for the π_H. There are many such examples.

TABLE III. Decay of quark model and hybrid $\pi(1800)$.

	Partial widths to final states					
State	$\pi\rho$	$\omega\rho$	$\rho(1465)\pi$	$f_0(1300)\pi$	$f_2\pi$	K^*K
$\pi_{3S}(1800)$	30	74	56	6	29	36
$\pi_H(1800)$	30	—	30	170	6	5

The essential point is that although the two states may have the same J^{PC} quantum numbers they have different internal structure, which will manifest itself in their decays. Unfortunately, nothing is simple, and we once again point out that strong mixing is expected between hybrids with conventional quantum numbers and $q\overline{q}$ states with the same J^{PC}, so that the decay patterns of physical states may not closely resemble those of either pure hybrids or pure $q\overline{q}$ states.

The final ingredient in hybrid searches is the production mechanism. Just as in glueball searches, a good place to look is in the gluon-rich J/ψ decays. A second reaction that has attracted interest is $p\overline{p}$ annihilation. Finally, photoproduction is potentially an important mechanism for producing hybrids. Hybrids could be produced copiously at an upgraded CEBAF at Jefferson Lab via an off-shell ρ, ω, or ϕ via vector-meson dominance interacting with an off-shell exchanged π (Close and Page, 1995b). The moral is that what is really needed is careful high-statistics experiments in all possible reactions.

F. Multiquark hadrons

The notion of color naturally explained Nature's preference for $q\overline{q}$ and qqq colorless systems. However, it also appears to predict multiquark states such as $q^2\overline{q}^2$ and $q^3q\overline{q}$ which could have exotic quantum numbers, thus indicating non-$q\overline{q}$ and qqq states (Jaffe, 1977a, 1977b, 1978; Lipkin, 1978). Upon considering $qq\overline{q}\overline{q}$ systems we find that the color couplings are not unique as they are in mesons and baryons. For example, we can combine two color triplet q's into a color 6 or $\overline{3}$. Likewise we can combine two antitriplet $\overline{q}$'s into a 3 or a $\overline{6}$.

Therefore there are two possible ways of combining $qq\overline{qq}$ into a color singlet: $3\overline{3}$ or $6\overline{6}$. In addition, since we could have combined a q and $\overline{q}$ into a color singlet 1 or color octet 8, we could also have combined the $qq\overline{qq}$ into color 11 and 88. Since two free mesons (in a color 11) are clearly a possible combination of $qq\overline{qq}$, the $3\overline{3}$ $6\overline{6}$ couplings can mix to give the 11-88 color configurations, which further complicates the details of the calculation. Thus, whether or not multiquark states exist is a dynamical question. It is possible that multiquark states exist as bound states, but it is also possible that $qq\overline{qq}$ configurations lead to hadron-hadron potentials (Barnes *et al.*, 1987; Weinstein and Isgur, 1990; Swanson, 1992). Both possibilities must be taken into account when attempting to unravel the hadron spectrum.

As in the case of hybrid mesons, states not fitting into the $q\overline{q}$ framework are the most unambiguous signature for multiquark states. In particular, flavor exotics, which have quantum numbers not allowed by $q\overline{q}$ or qqq such as fractionally charged or doubly charged mesons, are our best bet for finding genuine multiquark states. There is a large literature on the physics of multiquark states that attempts to predict their masses, explain their properties, and interpret observed hadron structures as multiquark states. One should take much of what exists in the literature with a grain of salt, as few of these predictions are based on full dynamical calculations. An exception is a quark-model study of the $J^{PC}=0^{++}$ sector of the $qq\overline{qq}$ system (Weinstein and Isgur, 1982, 1983). It found that weakly bound $K\overline{K}$ "molecules" exist in the isospin-zero and -one sectors, in analogy to the deuteron. It was suggested that these two bound states be identified with the $f_0(980)$ and $a_0(980)$. The meson-meson potentials that come from this picture, when used with a coupled-channel Schrödinger equation, reproduce the observed phase shifts for the a_0 and f_0 in $\pi\pi$ scattering. The $K\overline{K}$ molecules are the exception, however, as the model predicts that in general the $qq\overline{qq}$ ground states are two unbound mesons.

So far only pseudoscalar mesons in the final state have been considered in detail, so the next logical step is to extend the analysis to vector-vector (Dooley *et al.*, 1992) and pseudoscalar-vector channels (Caldwell, 1987; Longacre, 1990). One such possibility is a threshold enhancement in K^*K.

There are a number of distinctive signatures for the multiquark interpretation of a resonance. For molecules one expects strong couplings to constituent channels. For example, the anomalously large coupling of the $f_0(980)$ to $K\overline{K}$, despite having almost no phase space, is a hint that it is not a conventional $q\overline{q}$ state. Electromagnetic couplings are another clue to nonstandard origins of a state. Barnes found that the two-photon widths for a $q\overline{q}$ state are expected to be much larger than those of a $K\overline{K}$ molecule (Barnes, 1985b). Radiative transitions can also be used to distinguish between the two possibilities. For the case of a $f_1(K^*K)$ object one would expect the dominant radiative mode to arise from the radiative transition of the K^* constituent in $K^*\to K\gamma$, while an $f_1(s\overline{s})$ state would be dominated by the transition $f_1(s\overline{s})\to\gamma\phi$. The two cases would be distinguished by different ratios of the $\gamma K^0\overline{K}^0$ and γK^+K^- final states. Likewise, Close, Isgur, and Kumano (1993) suggest a similar test for the $f_0(980)$ and $a_0(980)$ involving the radiative decays $\phi\to\gamma f_0\to\gamma\pi\pi$ and $\gamma a_0\to\gamma\pi\eta$, which can distinguish between the quarkonium and $K\overline{K}$-molecule interpretation of these states.

Whether or not multiquark states exist, it is still extremely important to understand hadron-hadron potentials arising from multiquark configurations (Swanson, 1992) so that the observed experimental structure can be understood. There is, in fact, evidence for meson-meson potentials. In the reaction $\gamma\gamma\to\pi^0\pi^0$, the meson-meson potentials are needed to reproduce the $\gamma\gamma\to\pi^0\pi^0$ cross-section data (Blundell *et al.*, 1998). Enhancements in the production of low-invariant-mass $\pi\pi$ pairs have been observed in other processes as well; $\eta'\to\eta\pi\pi$, $\psi'\to J/\psi\pi\pi$, $\Upsilon(nS)\to\Upsilon(mS)\pi\pi$, and $\psi\to\omega\pi\pi$. Similar enhancements have also been seen in some $K\pi$ channels in $\overline{p}p\to K\overline{K}\pi$. These enhancements contribute to the confusion in attempting to understand the resonance structure underlying the experimental cross sections. The lesson is that, even if multiquark states do not exist, final-state interactions arising from hadron-hadron potentials will play an important role in our understanding of meson spectroscopy in the 1-to-2-GeV mass region.

III. EXPERIMENTS

The meson spectrum consists of an increasingly large number of broad, overlapping states, crowded into the mass region above around 1 GeV/c^2. Experiments disentangle these states using a combination of three approaches. First, states have different sensitivity to the various available production mechanisms. Comparing the results from two or more sources can be particularly useful when trying to enhance one state relative to another. Second, experiments using the same production mechanism (or at least the same beam and target) may be sensitive to a number of different final states. This naturally leads to a sensitivity to differing isospin and G-parity channels and can provide consistency checks for states within one particular experiment.

The third approach aims at unraveling the different J^{PC} combinations within a specific experiment and final state. Such a "partial-wave analysis" is crucial when overlapping states are produced in the same reaction (Chung and Trueman, 1975; Aston *et al.*, 1985; Sinervo, 1993; Cummings and Weygand, 1997). In all cases, it is important to understand fully the "acceptance" of the detector over the whole of phase space, so that the basis orthogonality conditions are correctly exploited. One clear example is demonstrated (Adams *et al.*, 1998) in Fig. 8, where the $\pi^-\pi^+\pi^-$ mass spectrum from the reaction $\pi^-p\to\pi^-\pi^+\pi^-p$ clearly shows a rich structure of strongly overlapping peaks. In this case, the partial-wave analysis is able to cleanly separate the contribu-

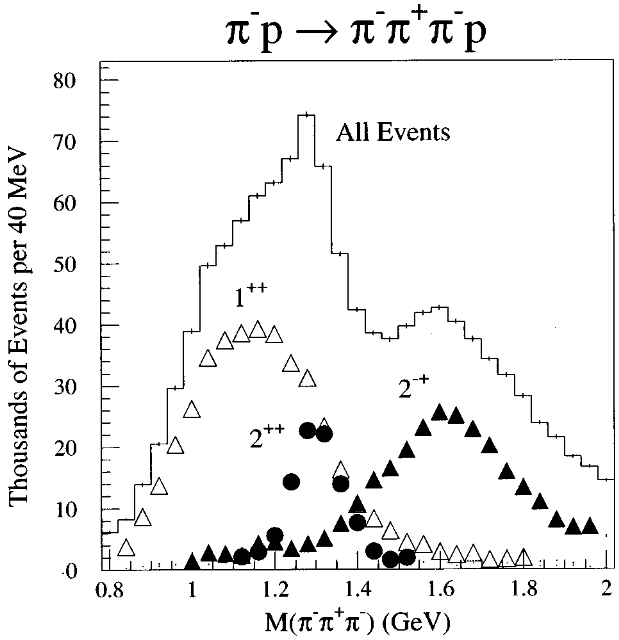

FIG. 8. Three-pion mass distribution for the reaction $\pi^- p \to \pi^- \pi^+ \pi^- p$ at 18 GeV/c, from experiment E852 at BNL. A partial wave analysis is used to decompose the spectrum into its dominant J^{PC} components, clearly showing the $a_1(1260)$, the $a_2(1320)$, and the $\pi_2(1680)$.

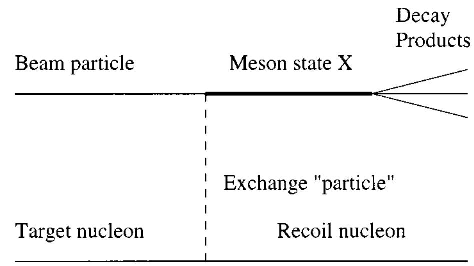

FIG. 9. Schematic diagram of a hadronic peripheral production process. Momentum is exchanged through an off-mass-shell particle, which may or may not be charged.

tions from the various J^{PC}. Furthermore, one can examine the phase motion of one "wave" relative to the other and demonstrate that the intensity peak is in fact resonant.

Different experiments measuring different properties are taken together to unravel the meson spectrum. In the following sections, we describe the various types of experiments, along with their advantages and disadvantages. We also mention some specific experiments whose data will be described in later sections.

A. Hadronic peripheral production

Most of the data on light-meson spectroscopy have come from multi-GeV pion and kaon beams on nucleon or nuclear targets, where the beam particle is excited and continues to move forward, exchanging momentum and quantum numbers with a recoiling nucleon.

Meson-nucleon scattering reactions at high energy are strongly forward peaked, in the direction of the incoming meson. Typically, the forward-going products are mesons, with a ground- or excited-state baryon recoiling at a large angle. This mechanism is shown schematically in Fig. 9. The excited meson state X has quantum numbers determined by the exchange, and subsequently decays to two or more stable particles. Two typical examples in modern experiments include $K^- p \to K^- K^+ \Lambda$ by the LASS collaboration (Aston *et al.*, 1988d) and $\pi^- p \to X^- p \to \rho^0 \pi^- p \to \pi^- \pi^+ \pi^- p$ by the E852 collaboration (Adams *et al.*, 1998). Other recent experiments with π^- beams include VES at IHEP/Serpukhov (Beladidze *et al.*, 1993; Gouz *et al.*, 1993) and BENKEI at KEK (Fukui *et al.*, 1991; Aoyagi *et al.*, 1993). Of particular note is the GAMS Collaboration (Alde *et al.*, 1986, 1988a, 1988b, 1992, 1997), which detected all neutral final states. Peripheral reactions with positively charged beams on proton targets are also possible. There is no reason to expect *a priori* that any particular type of hadronic state ($q\bar{q}$, multiquark, glueball, or hybrid) should be preferred over any other in this mechanism.

Peripheral reactions are characterized by the square of the four-momentum exchanged, called $t \equiv (p_{\rm beam} - p_X)^2 < 0$ (see Fig. 9). The forward-peaking nature can be seen in an approximately exponentially falling cross section with t, i.e., e^{bt} with $b \sim 3-8\ \mathrm{GeV}^{-2}$. For example, in $\pi^- p$ charge-exchange reactions at small values of $-t$, one-pion exchange dominates and is fairly well understood. It provides access only to states with J^{PC} = even^{++} and odd^{--}, the so-called "natural-parity" states. Other states such as $J^{PC}=0^{-+}$ can be produced by neutral $J^{PC}=0^{++}$ "pomeron" exchange or ρ^+ exchange, but these are not as well understood. Often the analysis is performed independently for several ranges of t, to try to understand the nature of the production (exchange) mechanism.

In the spirit of both Regge phenomenology and field theory, the diagram in Fig. 9 is taken literally when interpreting a partial-wave analysis of peripheral reactions. That is, the result of the partial-wave analysis is used to infer that the exchange particle exists and that it couples to the beam particle and excited meson state, conserving angular momentum, parity, and charge conjugation. As shown by Chung and Trueman (1975), the analysis is naturally divided into two sets of noninterfering waves on the basis of positive or negative "reflectivity," which in turn corresponds to "natural" $[P=(-1)^J]$ or "unnatural" $[P=(-1)^{J+1}]$ parity of the exchange particle with spin J.

The generality of this production mechanism and the high statistics available result in several advantages. One possibility is to use triggers that are as unrestrictive as possible to give large, uniform acceptance, and to choose particular final states in the analysis stage (Aston *et al.*, 1990). On the other hand, many experiments design the trigger to choose only a particular final state since the events of interest may occur very infrequently. A difficulty with this approach is that the detection efficiency for the final state is usually not uniform in the kinematic variables and one must be careful in modeling the experiment when performing the partial-wave analysis.

Detection of final-state photons, to identify π^0 and η and the objects that decay to them, has come of age in recent experiments. The GAMS Collaboration made use of all-photon final states in particular, but fine-grained electromagnetic calorimeters have been combined with charged-particle tracking and particle identification by the E852 and VES Collaborations. This opens up a large number of final states that can be studied simultaneously in a single experiment. Comparing decay branches of various states and searching for decay modes that were not previously accessible can be very powerful.

B. Peripheral photoproduction

Peripheral hadronic reactions have been the workhorse of meson spectroscopy, mainly because of the wide range of kinematics available, along with the accessibility and high cross section of hadron beams. Unfortunately, however, there is little selectivity for specific meson states. Except for the ability to do some selection on t and to bring in strange quarks by using K instead of π beams, one is limited to exciting the spin singlet ground-state $q\overline{q}$ combination (i.e., pions and kaons) because only these are stable against strong decay and therefore live long enough to produce beams for experiments.

Peripheral *photo*production reactions provide a qualitative alternative. The hadronic properties of the photon are essentially given by vector dominance (Bauer *et al.*, 1978). That is, the photon couples to hadrons as if it were a superposition of vector-meson states. In this case, Fig. 9 still applies, but the incoming "beam" particle is a spin *triplet* ground-state $q\overline{q}$. Consequently the series of preferred excitations is likely to be quite different. This mechanism has, in fact, been argued to be the most likely way of producing hybrid mesons with exotic quantum numbers by means of flux-tube excitation (Isgur, Kokoski, and Paton, 1985; Afanasev and Page, 1998).

Peripheral photoproduction has further advantages. The vector-dominance model allows non-OZI-suppressed excitation of heavy-quark states, such as $s\overline{s}$ and $c\overline{c}$, through production of the associated vector meson(s) (the ϕ and ψ states, respectively).

Unfortunately, there is a dearth of data from peripheral photoproduction. This is mainly due to the lack of high-quality, high-intensity photon beams and associated experimental apparatus, although this situation will change in the near future (see Sec. VI.F). A thorough review of the experimental situation through the mid-1970s is available in Bauer *et al.* (1978). The most significant contributions to meson spectroscopy since that time have been made by the LAMP2 experiment at Daresbury (Barber *et al.*, 1978, 1980), with photon-beam energies up to 5 GeV, and the Ω-Photon Collaboration at CERN (Aston *et al.*, 1982), using energies between 20 and 70 GeV. Spectroscopy in exclusive photoproduction has also been carried out by the E401 group at Fermilab (Busenitz *et al.*, 1989), who also used an electronic detector with a relatively open trigger, and the E687 group at Fermilab (Frabetti *et al.*, 1992), a successor to E401 that concentrated on heavy-quark physics.

One set of measurements performed with the SLAC hybrid bubble chamber and a 19-GeV laser-backscattered photon beam (Condo *et al.*, 1990, 1993; Blackett *et al.*, 1997) has yielded tantalizing results despite having very weak statistics. For example, Condo *et al.* (1990, 1993) see evidence for a narrow state at 1775 MeV decaying to $(\rho\pi)^{\pm}$ and produced in charge-exchange photoproduction. Only with new, precision beams and detectors, now in the planning stage, can these data be thoroughly investigated.

C. $\bar{p}p$ and $\bar{N}N$ reactions

Annihilation of antiquarks on quarks can be accomplished straightforwardly by using antiproton beams on hydrogen or deuterium targets. States that decay directly to $\overline{p}p$ and $\overline{p}n$ can be studied by measuring inclusive and exclusive annihilation cross sections as a function of the beam energy. This approach is obviously limited to states with masses greater than 2 GeV/c^2. It has been used quite effectively to study the charmonium system by the Fermilab E769 Collaboration (Armstrong *et al.*, 1997, and references therein) and to study some relatively massive light-quark states in the JETSET experiment at CERN (Bertolotto *et al.*, 1995; Buzzo *et al.*, 1997; Evangelista *et al.*, 1997, 1998). However, most of the contributions to light-meson spectroscopy have come from $\overline{p}p$ annihilations at rest with the Crystal Barrel experiment (Aker *et al.*, 1992) at CERN. This experiment has been reviewed quite recently (Amsler, 1998). Significant contributions have also come from the OBELIX experiment (Bertin *et al.*, 1997a, 1997b, 1998a, 1998b), in particular, using the $\overline{n}p$ annihilation reaction. These reactions were also studied by the ASTERIX Collaboration (May *et al.*, 1989, 1990a, 1990b; Weidenauer *et al.*, 1993).

The annihilation process is clearly complicated at a microscopic level. However, in annihilation at rest, in which a state X recoils against a light, stable meson, one may expect many different components to the wave function of X, including non-$q\overline{q}$ degrees of freedom. In fact, this process has been suggested as a fine way to excite gluonic degrees of freedom, in which case X might have large glueball or hybrid content, so long as the mass is not much larger than ~1700 MeV/c^2.

Antiproton annihilation in liquid hydrogen proceeds almost entirely through a $\overline{p}p$ relative S state (Amsler, 1998). This lends enormous power to the partial-wave analysis because the initial state is tightly constrained. Annihilation into three stable, pseudoscalar mesons (for example, $\overline{p}p \to \pi^0\pi^0\pi^0$ and $\overline{p}p \to \eta\eta\pi^0$) has been particularly fruitful. In these reactions, one studies the two-body decay of a meson resonance that recoils off of the third meson. These data have had their greatest impact on the scalar-meson sector.

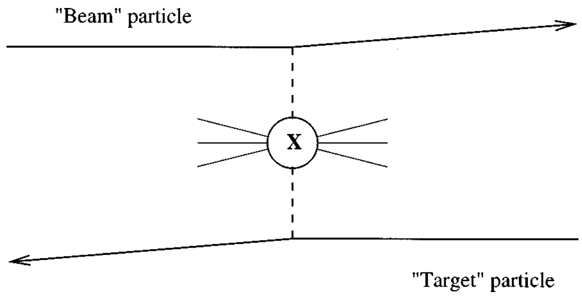

FIG. 10. Schematic diagram of a central production process. Two off-mass-shell particles collide, while the ''beam'' and ''target'' move off essentially unscathed. The exchange particles are understood to be very rich in ''glue'' when the transverse-momentum kick to the beam and target is small.

D. Central production

Peripheral processes are viewed as exciting the ''beam'' particle by means of an exchange with the ''target'' particle, leaving the target more or less unchanged. Central production refers to the case in which there is a collision between exchange particles, shown schematically in Fig. 10. Experimentally, using a proton beam and target, one observes the reaction $pp \rightarrow p_f(X^0)p_s$ where p_s and p_f represent the slowest and fastest particles in the laboratory frame. At high energies and low transverse momentum to p_s and p_f, this process is believed to be dominated by double Pomeron exchange. As the Pomeron is believed to have large gluonic content, one might expect X^0 to be a state dominated by gluonic degrees of freedom (Close and Kirk, 1997).

This technique has been exploited extensively at CERN, in the WA76 (Armstrong *et al.*, 1989a, 1989b, 1991a, 1991b, 1992), WA91 (Abatzis *et al.*, 1994), and WA102 experiments (Barberis *et al.*, 1997a, 1997b, 1997c, 1998).

E. Results from e^+e^- storage rings

High-luminosity ($\mathcal{L} \geqslant 10^{32}/\mathrm{cm}^2 \cdot \mathrm{sec}$) e^+e^- storage rings have been in operation for close to three decades. Known primarily for their contributions to heavy-quark spectroscopy, they have shed valuable light on the light-quark mesons in a variety of ways. These include direct production and spectroscopy of isovector and isoscalar vector mesons (i.e., ρ, ω, and ϕ states), states produced in the radiative decay of the J/ψ, and indirect production of various mesons in ''two-photon'' collisions.

1. Vector-meson spectroscopy

The e^+e^- annihilation process is mediated by a single virtual photon with the quantum numbers $J^{PC}=1^{--}$. These reactions therefore produce vector-meson resonances, and the isospin and other dynamic features are studied through the appropriate final states. By varying the $e^{\pm}$ beam energy, experiments scan through the center-of-mass energy and trace out the resonance shape, modified by interferences with overlapping states. One interesting recent application of this technique is the observation of the isospin-violating decay $\phi \rightarrow \omega\pi^0$ by the SND group at VEPP-2M (Achasov *et al.*, 1999).

The reaction $e^+e^- \rightarrow \pi^+\pi^-$ has been carefully studied in the region up to around 2-GeV center-of-mass energy (Barkov *et al.*, 1985; Bisello, 1989). These data have clearly established three ρ resonances (Bisello, 1989) and have also been used to study ρ/ω mixing precisely. Other reactions have also been studied by scanning over center-of-mass energies in this mass region, mainly by the DM2 Collaboration at ORSAY (Castro *et al.*, 1994).

2. Two-photon collisions

Typically, e^+e^- colliders are used to acquire large amounts of data at high-energy resonances, such as $c\bar{c}$ or $b\bar{b}$ states, or (in the case of LEP) at the Z^0 pole. One very fruitful source of data on meson spectroscopy (as well as other physics) in high-energy e^+e^- collisions is the reaction $e^+e^- \rightarrow e^+e^-X$, where the state X is produced by the collision of two photons radiated from the beam electron and positron. This is analogous to the central production process (Sec. III.D, Fig. 10) in which the photons replace the less-well-understood Pomeron. The field of ''two-photon'' physics has been rather extensively reviewed (Cooper, 1988; Morgan, Pennington, and Whalley, 1994). Data continue to be acquired and analyzed at operating e^+e^- storage ring facilities.

Some particular features of meson spectroscopy in two-photon collisions are immediately apparent. First, it is clear that only self-conjugate, $C=+1$ meson states X will be formed in the collision. Second, to the extent that the photons couple directly to the q and $\bar{q}$ that are formed, the production rate will be proportional to the fourth power of the quark charge. Thus u (and c) quarks will be preferred, relative to d and s quarks. This fact has been used to determine the singlet/octet mixing in the η and η' (Cooper, 1988). Also, if a state is dominated by gluonic degrees of freedom (a ''glueball''), then there is no valence charge to couple to photons, so we expect glueballs *not* to be produced in these reactions. This was discussed in Sec. II.E.1 and quantified in Eq. (25). Thirdly, two-photon reactions are a powerful tool for spectroscopy because of the way the scattered $e^{\pm}$ are detected.

Spectroscopic data from two-photon collisions are generally separated into ''untagged'' and ''tagged'' samples. The virtual-photon spectrum is sharply peaked in the forward direction, since the photon propagator is essentially proportional to $1/q^2$ where q is the photon four-momentum. Consequently, if the incident $e^{\pm}$ is scattered through a large enough angle to be ''tagged'' by the detector, the exchanged photon will have large enough q^2 to be strongly ''virtual.'' That is, it will have a significant component of longitudinal polarization. On the other hand, if neither the electron nor the positron is tagged, one can safely assume that the exchanged photons are essentially ''real.''

This leads to powerful selection rules (Yang, 1950) on the quantum numbers of the meson formed in the collision. In particular, for real ($q^2=0$) photons, all spin-1 states and odd-spin states with negative parity are forbidden. Therefore only states with $J^{PC} = 0^{\pm +}, 2^{\pm +}, 3^{++}, \ldots$ are produced in untagged events, and states with other quantum numbers should show a very strong dependence on q^2 for tagged events.

3. Radiative J/ψ decays

All decays of the form $J/\psi \to \gamma X$ (except $J/\psi \to \gamma \eta_c$) involve the annihilation of the $c\bar{c}$ pair into a photon and a hadronic state of arbitrary mass. To first order in perturbative QCD, this proceeds through $J/\psi \to \gamma gg$, so one might expect the hadronic state to couple strongly to two gluons. Consequently radiative J/ψ decay has long been regarded as a fertile hunting ground for glueballs. In this manner, at least, it is quite complementary to two-photon production. Here again, the state X must be self-conjugate with $C = +1$.

The branching ratio for radiative J/ψ decay is typically between $\sim 10^{-4}$ and $\sim 10^{-3}$. The best experiments produce on the order of several million J/ψ, so only a few thousand accepted events can be expected for each of these states. Consequently the statistical power is meager, and complete partial-wave analyses are difficult.

This subject has not been reviewed for some time (Königsmann, 1986; Hitlin and Toki, 1988). Since then, however, new results have been presented from the DM2 Collaboration at Orsay (Augustin *et al.*, 1988, 1990) and the Mark III experiment at SPEAR (Bai *et al.*, 1990; Bolton, 1992a, 1992b; Dunwoodie, 1997), and new data are being collected and analyzed by the BES Collaboration in Beijing (Bai, 1996a, 1996b, 1998a, 1998b).

IV. THE QUARK MODEL: COMPARISON WITH EXPERIMENT

A. Heavy quarkonia

It is useful to start with heavy quarkonium where there is theoretical justification for using potential models to calculate spectra and where there is some validity in identifying the static quark potential one obtains from lattice QCD calculations with the phenomenological potential obtained empirically from heavy-quarkonia spectra. What is surprising is that the general spectroscopic features evident in the heavy-quarkonia spectra persist to light-meson spectroscopy, where the quark model is on shakier ground. In Fig. 3 we compared quark-model predictions for the $b\bar{b}$ system to experiment and found the agreement to be good. In Fig. 11 we show a similar comparison for the $c\bar{c}$ spectrum, also with good agreement between prediction and experiment. The basic experimentally observed level structure consists of a tower of 1^- radial excitations with only one P-wave orbitally excited multiplet for the $c\bar{c}$ mesons and two for the $b\bar{b}$ mesons. This is because the 1^- states are directly produced in e^+e^- collisions, while the higher orbital states require decays from the produced 1^- states and are therefore much more difficult to produce and observe.

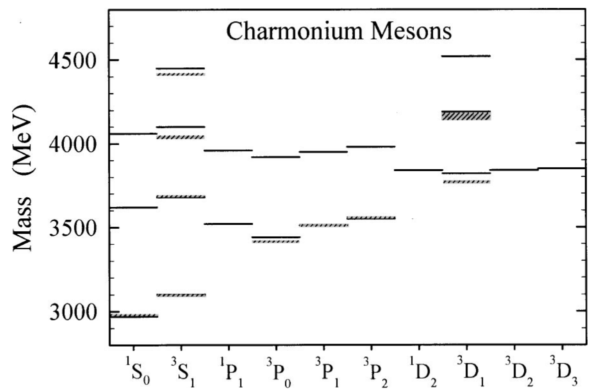

FIG. 11. The $c\bar{c}$ spectrum. The solid lines are the quark model predictions (Godfrey and Isgur, 1985) and the shaded regions are the experimental states. The size of the shaded regions approximates the experimental uncertainty. Note that the interpretation of the $\psi(4040)$ and $\psi(4160)$ as a single resonance is uncertain because of the substantial threshold effects in this energy region. Although they are included as $^3S_1$ states, their classification is ambiguous.

There have been suggestions that D-wave quarkonium can be produced via gluon fragmentation at hadron colliders (Qiao *et al.*, 1997a), in Z^0 decays (Qiao *et al.*, 1997b), in B decays (Yuan *et al.*, 1997; Ko *et al.*, 1997), and in fixed-target experiments (Yuan *et al.*, 1999). There is some evidence for the D-wave 2^{--} charmonium state with mass 3.836 ± 0.013 GeV in Fermilab experiment E705 (Antoniazzi *et al.*, 1994), although this result is questioned by other experiments (Gribushin *et al.*, 1996). It is also possible that, with the high statistics available at a B factory, the $b\bar{b}$ D waves might be observed via cascade radiative transitions from $\Upsilon(3S) \to 2P\gamma \to 1D\gamma\gamma$ (Kwong and Rosner, 1988). For both the $b\bar{b}$ and $c\bar{c}$ spectra, the states that differ measurably from the predicted masses are the 1^{--} states near open bottom and charm threshold, respectively, where the neglect of coupling to decay channels may not have been justified (Eichten *et al.*, 1975, 1978, 1980). Note that the quark-model classifications of some high-mass 1^{--} resonances are ambiguous. It has also been suggested that the $\psi(4040)$ and $\psi(4160)$ are mixtures of the $3S(c\bar{c})$ and charmonium hybrid (Close and Page, 1996).

B. Mesons with light quarks

Given the successful description of heavy quarkonia by the quark potential model we proceed to the light-quark mesons. Following the argument of Sec. II that the basic structure in heavy and light systems is qualitatively identical, we use the quark model to interpret the spectra of mesons with light-quark content. As already noted, studies of mesons with light quarks complement those of heavy quarkonium in that they probe a different piece of the $q\bar{q}$ potential, which allows the study of the strength and Lorentz structure of the long-range confining part of the potential. In addition, the hadroproduction mechanism is sufficiently different from production in e^+e^- colliders that the experimentally accessable excitations are nearly orthogonal.

In our survey of mesons with light quarks we begin by establishing the global validity of the quark-model predictions. With this backdrop, in Sec. V we shall focus on, and study in detail, the states that are either poorly understood or that pose a problem for the quark model.

1. Mesons with one light quark and one heavy quark

We start with mesons that contain one heavy quark, such as the charmed and beauty mesons (Rosner, 1986; Godfrey and Kokoski, 1991). These systems are an interesting starting point because, as pointed out long ago by De Rújula, Georgi, and Glashow (1976), as the heavy quark's mass increases, its motion decreases, so the meson's properties will increasingly be governed by the dynamics of the light quark and will approach a universal limit, referred to as the heavy-quark limit. Accordingly, these states become the hydrogen atoms of hadron physics. Mesons with one heavy quark provide a spectroscopy as rich as charmonium but, because the relevant scales are governed by the light quark, they probe different regimes. They bridge the gap between heavy quarkonium and light hadrons, providing an intermediate step on the way to the more complicated light-quark sector, where one searches for exotica like glueballs and hybrids. A growing number of excited charmed and beauty mesons have been observed by the ARGUS, CLEO, Fermilab E691 and E687, and more recently the LEP and CDF Collaborations. The experimental situation and quark-model predictions for the $c\overline{u}$, $c\overline{s}$, $b\overline{u}$, $b\overline{s}$, and $b\overline{c}$ states are summarized in Fig. 12 [see also Kwong and Rosner (1991) and Eichten and Quigg (1994)].

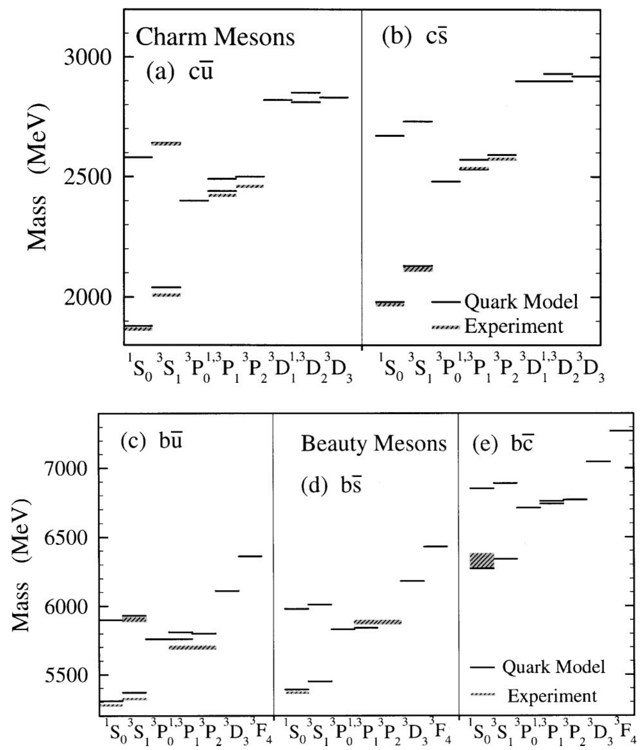

FIG. 12. Spectra for (a) $c\overline{u}$, (b) $c\overline{s}$, (c) $b\overline{u}$, (d) $b\overline{s}$, and (e) $b\overline{c}$. The solid lines are the quark-model predictions (Godfrey and Isgur, 1985) and the shaded regions are the experimental measurements, with the size representing the approximate experimental uncertainty. The experimental results shown for the $^{1,3}P_1$ and $^3P_2$ $b\overline{u}$ and $b\overline{s}$ represent broad bumps interpreted as superpositions of more than one state, assumed to be the B_1 (B_{s1}) and B_2 (B_{s2}) (see the text and Fig. 13).

For mesons composed of an unequal-mass quark and antiquark, charge-conjugation parity is no longer a good quantum number, so the triplet and singlet states of the same total angular momentum can mix via the spin-orbit interaction or some other mechanism (Lipkin, 1977). For example, the physical $J=1$ states are linear combinations of $^3P_1$ and $^1P_1$ with mixing angle θ, and the partial widths of the $J=1$ states are very sensitive to this angle. In the heavy-quark limit one of the $J=1$ states is degenerate with the $^3P_0$ state and the other is degenerate with the $^3P_2$ state. The decays of the P-wave mesons can be described by two independent amplitudes, S-wave and D-wave. One finds that the $J=1$ state degenerate with the $^3P_0$ state decays into final states in a relative S wave, the same as the $^3P_0$-state decay, while the $J=1$ state degenerate with the $^3P_2$ decays into final states in a relative D wave, the same as the $^3P_2$-state decay. Thus in the heavy-quark limit the four P-wave states split into two pairs of degenerate states.

These patterns in spectroscopy and decays can be extended to general principles. Recognition that the heavy-quark limit results in a new symmetry of QCD has led to considerable progress in our understanding of QCD through the study of mesons containing a single heavy quark (Isgur and Wise, 1989, 1990, 1991a, 1991b; Voloshin and Shifman, 1987; see also Neubert, 1994 for a recent review). 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 and the heavy quark acts as a static source of chromoelectric field as far as the light degrees of freedom are concerned. Thus, heavy-hadron spectroscopy differs from that of hadrons containing only light quarks because we may separately specify the spin quantum number of the light degrees of freedom and that of the heavy quark. That is, $\vec{S}_Q$ and $\vec{j}_q=\vec{S}_q+\vec{L}$ are separately conserved so that each energy level in the excitation spectrum is composed of degenerate pairs of states with total angular momentum given by $j=j_q\pm 1/2$. This new symmetry in the QCD spectrum in the heavy-quark limit leads to relations between hadrons containing a single heavy quark. The significance of these results cannot be overstated as they follow rigorously from QCD in the heavy-quark limit.

The heavy-quark flavor symmetry also says that strong-decay amplitudes arising from the emission of light quanta like π, η, ρ, $\pi\pi$, etc., are independent of heavy-quark spin states $\pm 1/2$, leading to a number of predictions. For one, the two D_1's both have $j^P=1^+$ and are only distinguished by j_q, which is a good quantum number in the limit $m_c\to\infty$. Since strong decays are entirely transitions of the light-quark degrees of freedom,

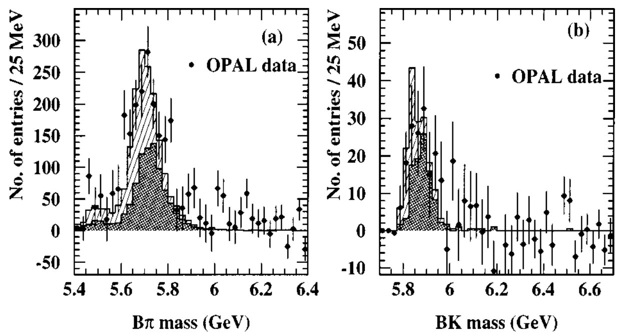

FIG. 13. The $B\pi$ and BK invariant-mass distributions. The solid histograms show Monte Carlo results for the B_2^* and B_{s2}^* states, respectively, and the hatched histograms show Monte Carlo results for B_1 and B_{s1}. From Akers *et al.*, 1995.

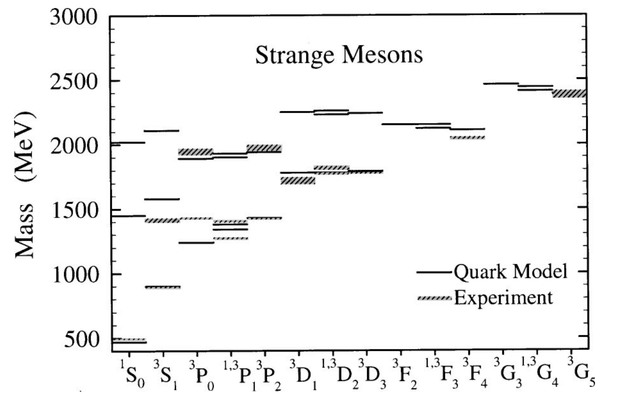

FIG. 14. Level diagram for the strange mesons. The quark-model predictions (Godfrey and Isgur, 1985) are given by the solid lines and the experimental measurements are given by the shaded regions with the size representing the approximate experimental uncertainty.

the decays from both members of a doublet with given j_q to the members of another doublet with j_q' are all essentially a single process and hence the two excited states should have exactly the same widths. The decays $D_0^*\rightarrow D\pi$ and $D_1\rightarrow D^*\pi$ are via the S wave and can be quite broad and therefore difficult to identify experimentally. In contrast $D_1\rightarrow D^*\pi$ and $D_2^*\rightarrow D^*\pi$, $D\pi$ proceed via the D wave and are much narrower. These states are identified as the $D_1(2420)$ and $D_2^*(2460)$. These decay predictions are the same as those obtained by the quark model. There are numerous other predictions resulting from heavy-quark effective theory (HQET) that are relevant to weak decays, and we encourage the interested reader to consult one of the more specialized reviews on this important subject (Neubert, 1994).

While the P-wave charmed mesons have been known for some time, the OPAL (Akers *et al.*, 1995), ALEPH (Buskulic *et al.*, 1996), and DELPHI (Abreu *et al.*, 1995) Collaborations at LEP have recently reported the discovery of P-wave beauty mesons. The OPAL results are shown in Fig. 13. Broad bumps are seen in $B\pi$ ($M=5.68\pm0.011$ GeV, $\Gamma=116\pm24$ MeV) and BK ($M=5.853\pm0.015$ GeV, $\Gamma=47\pm22$ MeV). The $B\pi$ results are consistent with similar results by ALEPH and DELPHI. In both cases the widths are larger than the detector resolution of 40 MeV and the bumps are interpreted as superpositions of several states and/or decay modes. In $B\pi$ the bump is assumed to be the B_1 and B_2 superimposed, and in BK the bump is assumed to be the B_{2s} and B_{1s} superimposed.

The LEP Collaborations have reported several candidates for the B_c state. The mean value for its mass averaged over the $\psi\pi$ decay mode is $m_{B_c}=6.33\pm0.05$ GeV (Abreu *et al.*, 1997; Ackerstaff *et al.*, 1998). More recently CDF has also reported an observation of B_c (Singh, 1998).

The DELPHI Collaboration (Abreu *et al.*, 1998a; Ehret, 1998) has reported evidence for radial excitations of the D^* and B^* (denoted D_r^* and B_r^*). From the invariant-mass distribution $M(D^*\pi\pi)$ DELPHI obtains the mass measurement of $M(D_r^*)=2637\pm2\pm6$ MeV, which is in good agreement with the quark-model prediction (Godfrey and Isgur, 1985). However, the observation of this state has not been confirmed by the OPAL Collaboration (Krieger, 1999). DELPHI also reports evidence for the B_r^* from $Q(B^{(*)}\pi^+\pi^-)$ with $M_{B_r^*}=5904\pm4\pm10$, which is also in good agreement with the quark model. We stress that these results are preliminary, but, together with the $L=1$ beauty-meson observations, demonstrate the potential of high-energy colliders for contributing to our understanding of hadron spectroscopy.

2. The strange mesons

An important ingredient in studying light-meson spectroscopy is the study of the level structure of the anticipated $q\bar{q}$ states. The strange mesons are a good place to start (Aston *et al.*, 1986a, 1986b, 1987) for a number of reasons; from an experimental perspective there is a reasonable amount of data on strange mesons to test the model. From a theoretical perspective strange mesons exhibit explicit flavor and therefore do not have the additional complication of annihilation mixing, which makes the isoscalar mesons much more difficult to unravel. It also eliminates the possibility of misidentifying a new meson as a glueball or $K\bar{K}$ molecule.

The strange-meson spectrum is shown in Fig. 14 and the strong-decay widths in Fig. 15. Many of the data come from the Large-Aperture Superconducting Solenoid Collaboration (LASS) at SLAC, a high-statistics study of K^-p interactions at 11 GeV/c using the LASS detector. First note the qualitative similarity to the $b\bar{b}$ (Fig. 3) and $c\bar{c}$ (Fig. 11) spectra. An important difference is that, in the strange-meson spectrum, there are few candidates for radial excitations while the complete leading orbitally excited K^* series are observed up to $J^P=5^-$. A substantial number of the expected underlying states are also observed in $K^-\pi^+$, $K_s\pi^+\pi^-$, and $K\eta$ final states. This reflects the importance of the production mechanism in determining features of the spectrum. Another important difference between the heavy-

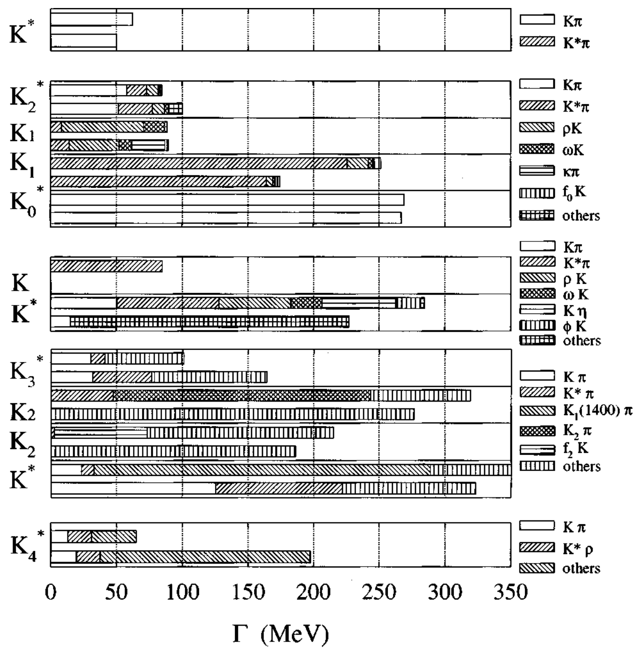

FIG. 15. Strange-meson strong decays. For each state the quark-model predictions are given by the upper bar [from Kokoski and Isgur (1987)] and the experimental results are given by the lower bar [from the Review of Particle Physics (Caso *et al.*, 1998) unless otherwise noted]. Some of the more important decay modes are indicated by the hatching. When only the total experimental width or at best only some of the partial widths have been measured we denote the unknown partial widths as "other."

quarkonia spectra and the light-quark meson spectra is the relative importance of the electromagnetic and hadronic transitions in the two systems. In light-meson systems, OZI-allowed decays are kinematically allowed and dominate, while in heavy-quarkonia the lower-mass states are below the threshold for OZI-allowed decays so that electromagnetic transitions dominate between these states. The important strong decays are shown in Fig. 15 for both the quark-model $^3P_0$-model predictions (Kokoski and Isgur, 1987) and the experimentally measured widths. Qualitatively, the predictions are in reasonable agreement with experiment. If a state is expected to have a broad width it generally does, and the pattern of partial widths is approximately reproduced by experiment. Thus, although quantitatively the predictions could be better, the predicted decay patterns should provide useful information when trying to decide on the nature of a newly found meson.

In $K^-p\to K^-\pi^+n$, a partial-wave analysis reveals the natural spin-parity states $J^P=1^-$ $K^*(892)$, 2^+ $K_2^*(1430)$, 3^- $K_3^*(1780)$, 4^+ $K_4^*(2060)$, and 5^- $K_5^*(2380)$ (Aston, 1986a). The agreement between the quark model and experiment is good for the masses of the leading orbital excitations, supporting the picture of linear confinement.

Partial waves of the reaction $K^-p\to\overline{K}^0\pi^+\pi^-n$ indicate structure in the 1^- wave around 1.4 and 1.8 GeV and in the 2^+ wave around 2.0 GeV. The individual K^* and ρ contributions reveal two 1^- Breit-Wigner resonances; a higher state with $M=1717\pm27$ MeV and $\Gamma=322\pm110$ couples to both channels, while the lower-mass state with $M=1414\pm15$ and $\Gamma=232\pm21$ is nearly decoupled from the $K\rho$ channel (Aston *et al.*, 1987). It is simplest to associate the higher state with the 1^3D_1 state based on the small triplet splitting; the lower state would be mostly the first radial excitation of the $K^*(892)$. However, the 2^3S_1 $K^*(1410)$ lies much lower in mass than quark-model predictions, and its small coupling to $K\pi$ indicates a breakdown in the simple $SU(3)$ model of decay rates. This is likely due to mixing between the two states via decay channels that would push the lower state down in mass and the upper state higher. This mixing would also cause one of the states to couple strongly to one decay channel and the other state to another decay channel.

There is also evidence for two structures in the S wave (Aston *et al.*, 1988e). The first, with mass around[8] $M[K_0^*(1430)]=1412\pm6$ MeV and $\Gamma[K_0^*(1430)]=294\pm23$ MeV, is classified as the $^3P_0$ partner of the $K_2^*(1430)$. Although this is $\sim$170 MeV higher than the quark-model prediction, the predicted 1^3P_0 mass is particularly sensitive to the details of the model, so not too much should be read into the discrepancy. There is a second S-wave structure at around 1.9 GeV with parameters $M=1945\pm22$ MeV and $\Gamma\sim201\pm86$ MeV. This structure can only be classified as a radial excitation of the 0^+ member of the $L=1$ multiplet, most probably the 2^3P_0 state. The 2^+ also demonstrates resonance behavior in this same mass region with $M=1973\pm26$ MeV and $\Gamma=373\pm68$ MeV, most probably the 2^3P_2.

One can probe the internal dynamics of these states by studying their decays. For example, $SU(3)$ predicts that the $K\eta$ branching ratio will be very small from even-spin K^* states and large from odd-spin states (Lipkin, 1981). The branching ratios are related to the ratios as

$$R_2=\frac{\Gamma(K_2^*\to K\eta)}{\Gamma(K_2^*\to K\pi)}=\frac{1}{9}(\cos\theta_P+2\sqrt{2}\sin\theta_P)^2\left(\frac{q_{K\eta}}{q_{K\pi}}\right)^5, \tag{28}$$

$$R_3=\frac{\Gamma(K_3^*\to K\eta)}{\Gamma(K_3^*\to K\pi)}=(\cos\theta_P)^2\left(\frac{q_{K\eta}}{q_{K\pi}}\right)^7, \tag{29}$$

where θ_P is the $SU(3)$ singlet-octet pseudoscalar mixing angle and has a value of $\theta_P\sim-20°$. R_2 suffers a significant suppression due to the cancellation between the two terms. This was studied in the reaction $K^-p\to K^-\eta p$ by the LASS Collaboration, who found (Aston *et al.*, 1988b)

$$BR(K_3^*\to K\eta)=9.4\pm3.4\%,$$

[8]The mass and width quoted in Table 2 of Aston *et al.* (1988e) are in fact incorrect due to a simple recording error. We quote the correct values here, as communicated to us by W. Dunwoodie.

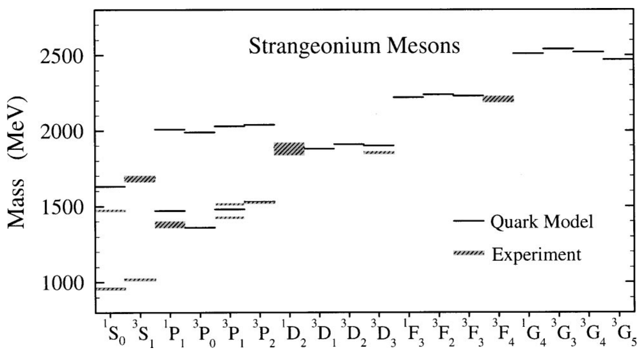

FIG. 16. Level diagram for the strangeonium mesons. Note that there are two candidate 1^3P_1 states, the $\eta(1470)$ is not unambiguously identified as the 2^1S_0 state, and the $f_4(2220)$ is an unconfirmed report by the LASS Collaboration. From Aston *et al.*, 1988d. See Fig. 14 for further details.

$$BR(K_2^* \rightarrow K\eta) < 0.45\% \quad (95\% \quad \text{C.L.})$$

in agreement with $SU(3)$. For the $K\eta'$ channel the situation is reversed, with the even-spin K^*'s expected to couple preferentially to $K\eta'$ and the odd-spin K^*'s couplings to $K\eta'$ expected to be suppressed. The relative phases of the different decay amplitudes also agree with the quark-model predictions.

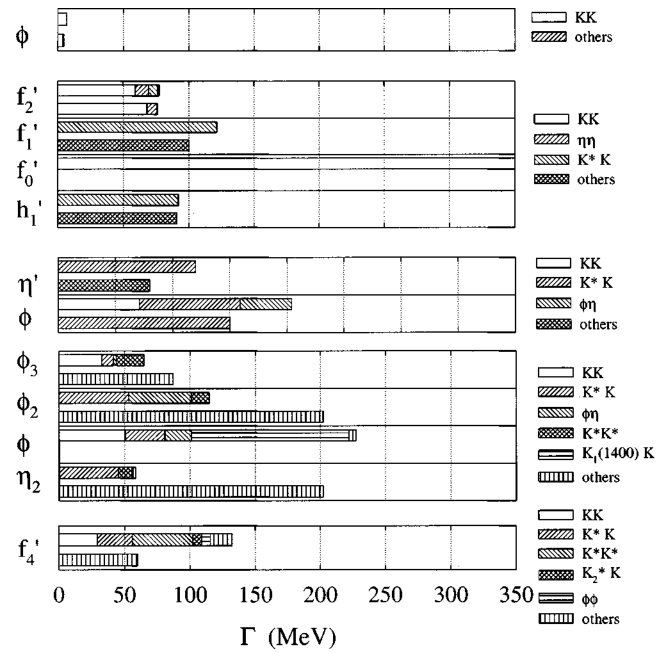

FIG. 17. Strangeonium-meson strong decays. For each state the quark-model predictions are given by the upper bar (Kokoski and Isgur, 1987; Blundell and Godfrey, 1996; Barnes *et al.*, 1997) and the experimental results are given by the lower bar. From the Review of Particle Physics, Caso *et al.*, 1998. See Fig. 15 for further details.

3. The strangeonium mesons

The strangeonium ($s\bar{s}$) states provide an intermediate mass between the heavier systems where the quark model is approximately valid and the lighter mesons where it is on less firm foundation. For this reason strangeonium states provide important input for hadron spectroscopy. It is also important to understand these states since a number of exotic candidates have been observed in final states where strangeonium might be expected.

As in the case of the strange mesons, many of the data on the $s\bar{s}$ states have come from the high-statistics study ($\sim$113 million triggers) of strangeonium mesons produced in the LASS detector by an 11-GeV/c K^- beam. The channels of interest are dominated by hypercharge-exchange reactions such as $K^-p \rightarrow K^-K^+\Lambda$, $K^-p \rightarrow K_s^0 K^{\pm}\pi^{\mp}\Lambda$, and $K^-p \rightarrow K_s^0 K_s^0 \Lambda$, which strongly favor the production of $s\bar{s}$ mesons over glueballs. The study of strangeonium in hypercharge-exchange reactions can provide revealing comparisons with the same final states produced in gluon-rich channels such as J/ψ radiative decays. The level diagram for the strangeonium spectrum is given in Fig. 16 and their decay widths are shown graphically in Fig. 17. The $s\bar{s}$ spectrum is similar to that of the strange-meson spectrum. Except for the ground-state pseudoscalar mesons, the observed states fit into $SU(3)$ multiplets consistent with magic mixing (pure $s\bar{s}$) and agree reasonably well with quark-model predictions. The observed leading states lie on an essentially linear orbital ladder that extends up through the 4^+ f_4', and there are good candidates for the $^3P_0$ and $^3P_1$ partners of the $f_2'(1525)$. The couplings of natural-parity states agree well with the $SU(3)$ predictions and the phases of the decay amplitudes are consistent with the quark-model predictions. Overall, the parameters and decay transitions agree well with the predictions of the quark model. Some of the open issues in $s\bar{s}$ spectroscopy are the $\eta(1440)$, $f_J(1710)$, and the $f_J(2220)$, which we discuss below and in Sec. V.

An amplitude analysis of the reactions $K^-p \rightarrow K_S^0 K_S^0 \Lambda_{seen}$ and $K^-p \rightarrow K^-K^+\Lambda_{seen}$ indicates S-wave structure around the $f_2'(1525)$ mass, Fig. 18 (Aston *et al.*, 1988a), suggesting the existence of a 0^+ resonance, which is most naturally interpreted as the $^3P_0$ partner of the $f_2'(1525)$. The approximate mass degeneracy of these states would imply that the spin-orbit interaction in the strangeonium sector is weaker than predicted, which is consistent with the strange-meson sector. This implies that the $f_0(980)$, sometimes interpreted as the $^3P_0(s\bar{s})$ state, is not a conventional $q\bar{q}$ state. There have been numerous sightings of another scalar meson with mass $\sim$1500 MeV but with different properties than the state possibly observed by LASS in $K^-p \rightarrow K\bar{K}\Lambda_{seen}$. We conclude that they are different states and discuss the latter in detail in the following section.

There are two candidates for the 1^{++} $s\bar{s}$ state, with masses $\sim$1510 MeV and $\sim$1420 MeV. Disentangling them requires consideration of the E/ι region, to be discussed in Sec. V.D. In the $K\bar{K}^*$+c.c. modes, a partial-wave analysis reveals structure around 1.5 GeV which is dominated by the 1^+ K^* wave and around 1.85 GeV in the 2^- and 3^- waves. The 1^+ waves can be combined to form eigenstates of G parity as shown in Figs. 19(a) and

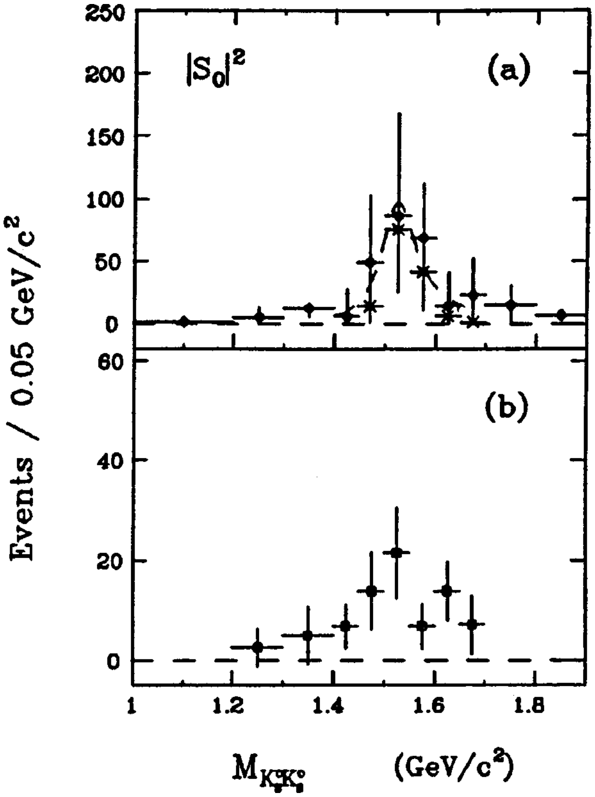

FIG. 18. S-wave intensity distribution for (a) the reaction $K^- p \to K^0_S K^0_S \Lambda_{seen}$ and (b) $K^- p \to K^- K^+ \Lambda_{seen}$. From Aston *et al.*, 1988a.

(b) (Aston *et al.*, 1988c). These distributions are well described by Breit-Wigner curves as shown, and, assuming $I=0$, they represent good evidence of two $s\bar{s}$ axial-vector-meson states—$J^{PC}=1^{++}$ with $M\simeq 1530$ MeV and $J^{PC}=1^{+-}$ with $M\simeq 1380$ MeV. These states are good candidates for the mainly $s\bar{s}$ members of the respective nonets since they are strongly produced in the Kp hypercharge-exchange reaction. The 1^{+-} state has also been reported by King (1991) in the partial-wave anaysis

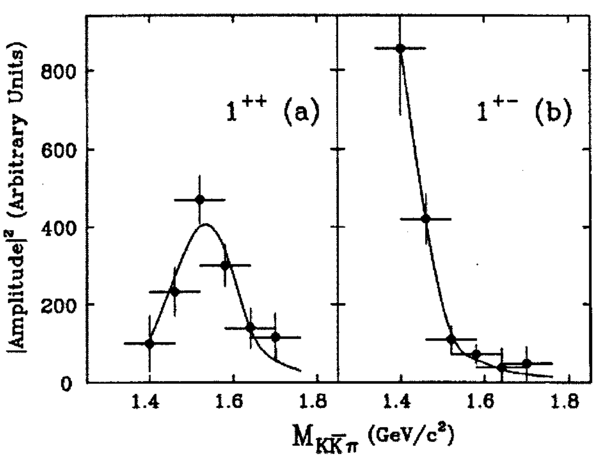

FIG. 19. The mass dependence of the total production amplitude intensities for the axial-vector G-parity eigenstates. The curves correspond to the Breit-Wigner fits described in the text. From Aston *et al.*, 1988c.

of the $K^+ K_S \pi^-$ system from the $K^- p$ interactions at 8 GeV/c (BNL-771 experiment) and the Crystal Barrel Collaboration (Abele *et al.*, 1997b). The 1^+ state, observed at around 1.5 GeV by LASS, would be identified with the $f_1'(1530)$ claimed by Gavillet *et al.* (1982) although no evidence for this state was found by Fermilab E690, who see the $f_1(1420)$ (Berisso *et al.*, 1998). These results, taken together with the $f_0(1530)$, indicate that the tensor and spin-orbit mass splittings are small. LASS shows no evidence for significant production of the $f_1(1420)$. If this is the case, the $q\bar{q}$ interpretation of the $f_1(1420)$, which has generally been taken as the 1^{++} strangeonium state, comes into question, indicating that the $f_1(1420)$ must be something else, a $K\bar{K}^*$ molecule perhaps. An alternative view advocated by the Particle Data Group (Caso *et al.*, 1998) assigns the $f_1(1420)$ as the $s\bar{s}\,1^{++}$ state and concludes that the $f_1(1510)$ is not well established.

There is another $s\bar{s}$ candidate, the $f_J(1710)$, with width $\Gamma=150$ MeV, seen in the invariant-mass distributions of $K^0_s \overline{K^0_s}$, $K\bar{K}$, and $\eta\eta$ pairs in J/ψ radiative decay (Baltrusaitis *et al.*, 1987). There is no evidence for such a state by the LASS Collaboration indicating that the $f_J(1710)$ is not a conventional strangeonium state. This state will be discussed in detail in Sec. V.

The $f_J(2220)$ was also first observed in a similar decay by the MARK III Collaboration ($M=2231\pm 8$ GeV, $\Gamma=21\pm 17$ MeV; Baltrusaitis *et al.*, 1986a; Einsweiler, 1984). The LASS group sees a similar object in the reaction $K^- p \to K^+ K^- \Lambda_{seen}$ G-wave amplitude, which is evidence for a 4^{++} state (Aston *et al.*, 1988d). This would be a member of the 4^{++} nonet [$f_4(2030)$, $a_4(2040)$, Cleland *et al.* (1982), and $K_4^*(2060)$, Aston *et al.* (1988e)] predicted by the quark model. The LASS analysis yields mass and width values of 2209^{+17}_{-15} and 60^{+107}_{-57} MeV/c^2 for this $J^{PC}=4^{++}$ state. There is also evidence for this state by the GAMS Collaboration in $\pi^- p \to \eta\eta' n$ (Alde *et al.*, 1986). This implies that the $f_J(2220)$, which has been conjectured to be an exotic hadron of some sort (Chanowitz and Sharpe, 1983a), is instead the $s\bar{s}$ member of the quark-model $^3F_4$ ground-state nonet (Godfrey *et al.*, 1984, Blundell and Godfrey, 1996). On the other hand the BES Collaboration (Bai *et al.*, 1996a) finds a decay width too narrow to be easily accommodated as the 1^3F_4 $s\bar{s}$ state. More importantly they find that the decays are approximately flavor symmetric, which supports a glueball interpretation. One intrepretation that could accommodate the contradictory experimental evidence is that there are in fact two resonances; the first is a broader conventional $s\bar{s}$ state and the second is a narrow glueball that is not observed in hadronic reactions.

To summarize, the $s\bar{s}$ spectrum agrees reasonably well with the predictions of the constituent-quark model. With this clarity, a number of important puzzles have been revealed. For example, it now seems clear that there are too many low-mass 0^{++} states and the $f_1(1420)/\eta(1440)$, $f_J(1710)$, and $f_J(2220)$ regions con-

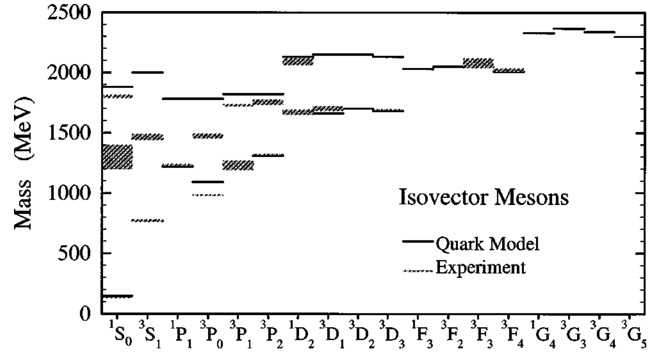

FIG. 20. Level diagram for the isovector mesons. See Fig. 14 for further details.

tain intriguing hints of non-quark-model physics. These puzzles will be explored in Sec. V.

4. The isovector mesons

Isovector mesons are made of the light up and down quarks, so they have the complications that arise from relativistic effects but do not have the additional complication of annihilation mixing that contributes to their isoscalar partners. The isovector-meson spectrum is shown in Fig. 20 and the decay widths are given in Fig. 21. For the most part there is good agreement between experiment and the quark-model predictions. In the isovector sector, the orbitally excited states extend up as far as the $L=5$ $J^{PC}=6^{++}$ and are consistent with the quark-model predictions, although the higher-mass states need confirmation. The multiplet splittings for the P- and D-wave mesons for the most part behave as expected. Given the general good agreement, it is the discrepancies that draw attention to themselves and require further discussion.

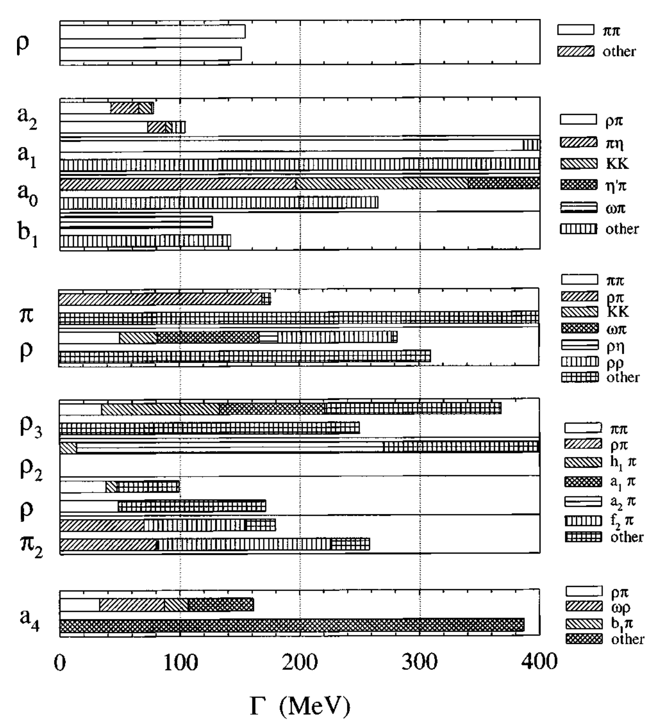

FIG. 21. Isovector-meson strong decays. See Fig. 15 for further details.

Although the $a_0(980)$ scalar meson $J^{PC}=0^{++}$ lies about 100 MeV below the quark-model prediction for the $^3P_0$ state expected in this mass region, it is the measured width of ~ 50 MeV that is much more difficult to reconcile with the quark-model prediction of ~ 500 MeV. This large discrepancy has led to numerous conjectures that the $a_0(980)$ is not a $q\overline{q}$ state but rather a more complicated four-quark object (Jaffe, 1977a, 1977b, 1978), most probably a $K\overline{K}$ molecule (Weinstein and Isgur, 1982, 1983). The observation of the $a_0(1450)$ in $p\overline{p}$ annihilation by the Crystal Barrel Collaboration (Amsler *et al.*, 1994a, 1995a) and the OBELIX Collaboration (Bertin *et al.*, 1998b) is naturally assigned to be the isovector member of 1^3P_0 nonet, reinforcing the interpretation that the $a_0(980)$ is a non-$q\overline{q}$ state. We shall return to the scalar-meson sector in the next section.

There is growing evidence for members of the radially excited $L=1$ multiplet. Both the CLEO (Kass, 1998) and DELPHI (Abreu *et al.*, 1998b) Collaborations observed a signal in $\tau \to a_1' \nu_\tau \to \pi^- \pi^+ \pi^- \nu_\tau$ with $M_{a_1} \simeq 1750$ MeV and $\Gamma_{a_1} \simeq 300$ MeV. Other observations of this state have been reported by BNL E818 (Lee *et al.*, 1994), BNL E852 (Ostrovidov, 1998), and VES (Amelin *et al.*, 1995). There is also evidence for the $J=2$ partner of this state, the a_2', observed by the L3 Collaboration in two-photon production with $\Gamma_{\gamma\gamma}(a_2) \times BR(\pi^+\pi^-\pi^0) = 0.29 \pm 0.04$ keV (Hou, 1998), making it the first radial excitation reported in $\gamma\gamma$ collisions. The measured resonance parameters are $M(a_2') = 1752 \pm 21$ MeV and $\Gamma(a_2') = 150 \pm 115$ MeV. In addition, the Crystal Barrel Collaboration reports an observation of an a_2 in $p\overline{p}$ annihilation in the final states $a_2 \to \eta\eta\pi^0$, $\eta 3\pi^0$ with a much lower mass, although the evidence for this state is not particularly strong (Degener, 1997).

The $\pi(1800)$ has been observed by the VES Collaboration (Amelin *et al.*, 1995). Despite the fact that its mass is consistent with the 3^1S_0 state there is speculation that it may be a hybrid because of its weak $\rho\pi$ decay mode. However, the 3^1S_0 interpretation predicts that the main decay mode would be $\rho\omega$ (Barnes *et al.*, 1997). VES has studied the $\rho\omega$ final state and does indeed find evidence for a large $\pi(1800)$ signal (Ryabchikov, 1998), which supports the 3^1S_0 assigment.

Finally, we mention the excited vector mesons. At first only one excited vector-meson state was observed, the $\rho(1600)$, whose properties did not agree well with the quark-model predictions. Godfrey and Isgur (1985) surmised that the observed state was a mixture of two broad overlapping resonances, the 2^3S_1 and the 1^3D_1, which gave rise to the observed properties. Subsequently, Donnachie and Clegg (1987) performed a full analysis of the data for the annihilation reactions $e^+e^- \to \pi^+\pi^-$, $2\pi^+2\pi^-$, $\pi^+\pi^-\pi^0\pi^0$ and the photoproduction reactions $\gamma p \to \pi^+\pi^- p$, $2\pi^+2\pi^- p$, $\pi^+\pi^-\pi^0\pi^0 p$ and came to the conclusion that a consistent picture required two states, the 2^3S_1 at 1465 ± 25

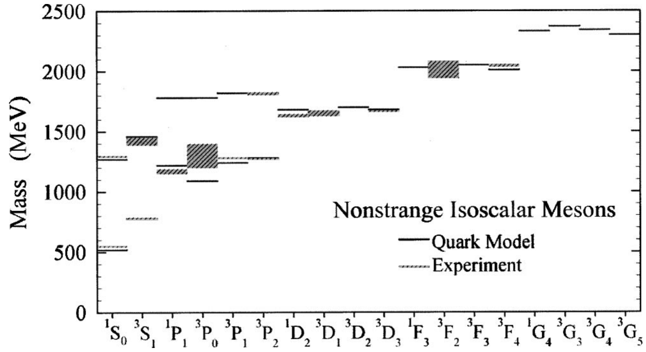

FIG. 22. The nonstrange isoscalar meson spectrum. See Fig. 14 for further details.

MeV with width 220±25 MeV and the 1^3D_1 at 1700±25 MeV with width 220±25 MeV. These results have been confirmed by recent results in $\overline{p}p$ annihilation by the Crystal Barrel Collaboration (Abele *et al.*, 1997c) and by the OBELIX Collaboration (Bertin *et al.*, 1997b, 1998b). The properties of these two states are in reasonable agreement with quark-model predictions, and more recent results examining the 4π decays of these states support the quark-model assignments (Thoma, 1998).

5. The nonstrange isoscalar mesons

The nonstrange isoscalar mesons are the last of the light-quark meson families and are also the most problematic. Phenomenologically, there is the problem of $n\overline{n}$ and $s\overline{s}$ mixing and, perhaps, the possibility of glueballs mixing with ordinary isoscalars. The spectrum and decay widths are shown in Figs. 22 and 23, respectively. Although there is generally good agreement between the experimental values and the quark-model predictions, there are also numerous puzzles that may point to the existence of "exotic" hadrons lying outside the constituent-quark model. We shall defer their discussion until Sec. V.

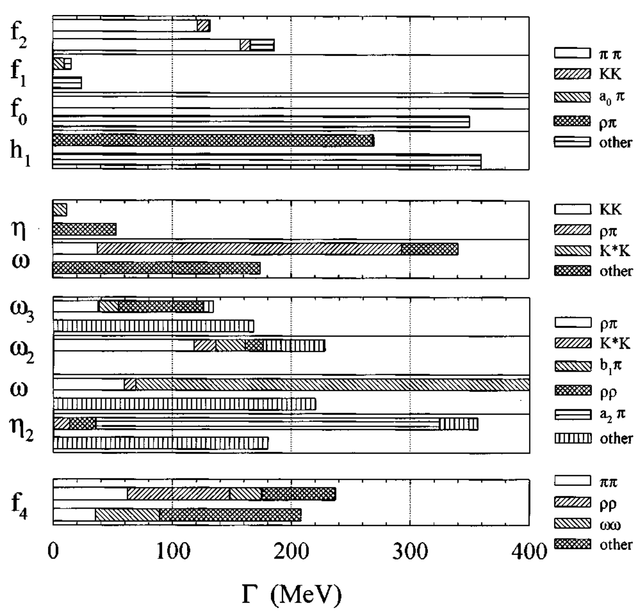

FIG. 23. The nonstrange isoscalar-meson decays. See Fig. 15 for further details.

As in the case of the $s\overline{s}$ mesons, the $^1P_1$ state $h_1(1170)$ and the $^3P_1$ $f_1(1285)$ differ slightly from their respective quark-model predictions, most likely indicating the need for further study of the annihilation mixing mechanism in the $^1P_1$ and $^3P_1$ sector of the model.

The Crystal Barrel Collaboration has reported two η_2 states in $p\overline{p}\to\eta 3\pi^0$ (Amsler *et al.*, 1996). The lighter state has a total width of about $\Gamma\sim 180$ MeV and was seen in $a_2\pi$. This is consistent with the quark-model expectations for the $I=0$ $^1D_2$ $n\overline{n}$ state, which should appear near the 1D multiplet mass of ~ 1670 MeV with $\Gamma\sim 260$ MeV and $a_2\pi$ the dominant decay mode. The second state, $\eta_2(1875)$, is too high in mass to be a 1D $n\overline{n}$ state, and the strong $f_2\eta$ mode argues against a mainly $s\overline{s}$ state. This state has been reported previously by the Crystal Ball Collaboration, who measured the $\eta\pi^0\pi^0$ mass spectrum in $\gamma\gamma\to 6\gamma$ (Karch *et al.*, 1990) and found $M(\eta_2)=1876\pm 35\pm 50$ MeV, $\Gamma_{total}(\eta_2)=228 \pm 90\pm 34$ MeV, and $\Gamma_{\gamma\gamma}(\eta_2)\cdot BR(\eta_2\to\eta\pi\pi)=0.9\pm 0.2 \pm 0.3$ keV. The angular distribution of the $\eta\pi^0$ subsystem gives $J^{PC}=2^{-+}$. CELLO also observed an enhancement in the cross section for $\gamma\gamma\to\eta\pi^+\pi^-$ yielding a mass of 1850±50 MeV, $\Gamma_{total}=380\pm 50$ MeV, and $(2J_X+1)\cdot BR(X\to\eta\pi\pi)=15\pm 5$ keV (Feindt, 1990). The best single resonance assignment is 0^{-+} $f_0\eta$ followed by 2^{-+} $a_2\pi$. This was the first new resonance to be discovered in $\gamma\gamma$ collisions. One possible explanation is that its higher-than-expected mass, like the π_2, might be understood in terms of final-state interactions. A second possibility is that it is a hybrid candidate (Barnes, 1998). This could be tested by searching for the $a_2\pi$ decay mode, which Close and Page (1995a) predict to be dominant.

The excited pseudoscalar mesons, in particular in the E/ι mass region, remain a puzzle (see Sec. V.D). Although the quark model predicts two excited pseudoscalar mesons, it is difficult to reconcile the properties of the $\eta(1295)$ and $\eta(1440)$ with the quark-model predictions. In addition, the $\eta(1440)$ is now considered to be composed of two separate resonances (Caso *et al.*, 1998). They are referred to as $\eta(1410)$ and $\eta(1490)$. One could identify the $\eta(1295)$ as the radial excitation of the η and the $\eta(1490)$ as the mainly $s\overline{s}$ radial excitation of the η'. This leaves the $\eta(1410)$ as an extra state. There is some speculation that it is somehow related to the $f_1(1420)$.

The 0^{++} state, the $f_0(980)$, has problems with its quark-model assignment similar to those of the $a_0(980)$ and is also interpretated as a four-quark state. If the $f_0(1370)$ and the $f_0(1525)$ reported by LASS are identified with the quark-model ground-state isoscalars, the discrepancies between the observed and predicted masses would be due to an overestimate of the quark-model spin-orbit splittings as in the isovector and strange-meson sectors.

There is also some confusion in the 2^{++} sector, which to be understood will have to be studied in conjunction with similar puzzles in the $s\overline{s}$ 2^{++} mesons. In particular, the $f_J(1710)$ has been mentioned previously in Sec. IV.B.3. In addition, there are three tensor mesons, $f_2(2010)/g_T$, $f_2(2300)/g'_T$, and $f_2(2340)/g''_T$, produced in the reaction $\pi^- p \to \phi\phi n$, which do not fit in with the quark-model predictions. Because they are produced in an OZI-forbidden process, it has been argued that they are strong candidates for gluonium states. These are discussed further in Sec. V.E.1.

6. Summary of light mesons

For the most part the observed meson properties are consistent with predictions of the quark model. The experimental regularities fit the quark model well, with orbital excitations on linear trajectories and triplet splittings at least qualitatively described by the quark model. To make further progress in the isovector and nonstrange isoscalar sectors, high-statistics experiments would be useful in finding higher-mass states.

Given the generally good agreement between the quark model and experiment, it is the discrepancies that suggest interesting physics. With the rather complete picture of the low-mass $q\bar{q}$ states described above it is becoming increasingly clear that several states have no obvious home in the $q\bar{q}$ sector. For example, the low-mass 0^{++} systems have been confusing for many years and it now appears that there are too many such states. Figures 16 and 22 show that two radial excitations of the isoscalar pseudoscalars should occur in the mass region 1300–1600 MeV. The $\eta(1295)$, $\eta(1410)$ and the $\eta(1470)$ cannot all fit into this picture and therefore at least one must appear as a spurious state. Two ground-state isoscalar 1^{++} states occur at 1240 and 1480 MeV in the quark model and are filled in by $f_1(1285)$ and $f_1(1510)$ so that the $f_1(1420)$ clearly appears as an extra state. Similar observations apply to the $f_0(1500)$, the $f_J(1710)$, the $f_2(2010)$, the $f_2(2300)$, and the $f_2(2350)$. These states point to a need for a better understanding of hadronic structure, perhaps to be gained by studying the relation between the $q\bar{q}$ meson properties and experimental observations or perhaps by enlarging the quarkonium picture to include gluonic degrees of freedom and multiquark states.

In addition, there are other puzzles that do not fit in the picture of $q\bar{q}$ spectroscopy and have not yet been discussed, for example, the state with exotic J^{PC} quantum numbers 1^{-+} $\hat{\rho}(1405)$, seen in $\pi^- p \to \pi^0 \eta n$ and structures in $\gamma\gamma \to VV$.

In the next section we shall examine the properties and possible interpretations of these anomalous meson resonances.

V. PUZZLES AND POSSIBILITIES

The quark model compares very well with the light-quark meson spectrum and meson decays. Certainly, there are disagreements, but many of these can be ascribed to the natural limitations of the model and the inherent complexity of QCD.

However, there is another class of disagreements that cannot be so ascribed. These cases are more profound and presumably point to the features of QCD *not* contained in the quark model. They may indeed point to the fundamental degrees of freedom needed to fully describe hadron structure. These may help us identify the necessary conditions that will one day have to be met by any purported solution of the full theory.

This class of disagreement is identified by one (or more) of three features. The most unequivocal of these is the establishment of states with "exotic" quantum numbers, that is, J^{PC} which cannot be accommodated with only $q\bar{q}$ degrees of freedom. The second is the overpopulation of states in well-defined mass regions. This is slightly ambiguous because the "expected" mass of a state will depend on the dynamics of the model, and we can never be entirely certain that the overpopulation does not arise from an "intruder" state based on higher radial or orbital excitations. Finally, we can identify puzzling states by their specific dynamical characteristics, such as their mass and partial decay widths. This is most difficult, since we need to rely on specific quark-model calculations to claim a fundamental disagreement, yet there are some cases in which this disagreement is indeed quite profound.

A. Exotic quantum numbers

As discussed in Sec. II (see Table I), the quantum numbers of a $q\bar{q}$ system must have either $P=C$ for all J, except $J^{PC}=0^{--}$, or $J^{PC}=0^{-+},1^{+-},2^{-+},\ldots$. Specifically, a state with any of the quantum numbers $J^{PC}=0^{--},0^{+-},1^{-+},\ldots$ would be manifest evidence for non-$q\bar{q}$ degrees of freedom. Because of their unique role, these states have been sought after for quite some time. However, only recently has some clear evidence been obtained experimentally. There are a number of reasons for this, but the main one is that their branching ratios to conventional final states are small. This means that one needs either to search very carefully amidst the forest of well-established states or to build and operate a dedicated experiment to search in more complicated final states, or both.

The flux-tube model (Isgur, Kokoski, and Paton, 1985; Isgur and Paton, 1985; Close and Page, 1995a) predicts that exotic hybrid mesons preferentially decay to pairs of S- and P-wave mesons, while decays to two S-wave mesons are suppressed on rather general grounds (Page, 1997b). Examples of preferred final states include $\pi b_1(1235)$, $\pi f_1(1285)$, $\pi a_2(1320)$, and $KK_1(1270)$. Essentially all of the $S+P$ decay modes are very difficult to observe experimentally, since they involve a large number of final-state particles that can arise from different quasi-two-body states. For example, $\pi^- f_1(1285)$ leads to $\pi^-\pi^+\pi^-\eta$, which might also be due to such things as $\pi^-\eta(1295), \eta a_2^-, \eta a_1^-, \rho^0 a_2^-$, and so forth. Unless the particular reaction enhances *production* of hybrid exotics (Isgur, Kokoski, and Paton, 1985) it will likely be very hard to identify clearly the exotic state amidst the

TABLE IV. Peripheral hadronic production experiments with $\eta\pi$ final states.

Experiment	Laboratory	Reaction	p_{beam} (GeV/c)	Reference
NICE	IHEP	$\pi^- p \to \eta\pi^0 n$	40	Apel *et al.*, 1981
GAMS	CERN	$\pi^- p \to \eta\pi^0 n$	100	Alde *et al.*, 1988a
BENKEI	KEK	$\pi^- p \to \eta\pi^- p$	6.3	Aoyagi *et al.*, 1993
VES	IHEP	$\pi^- N \to \eta\pi^- X$	37	Beladidze *et al.*, 1993
E852	BNL/AGS	$\pi^- p \to \eta\pi^- p$	18	Thompson *et al.*, 1997

debris of highly excited, conventional $q\bar{q}$ mesons (Barnes, Close, Page, and Swanson, 1997).

Experimentally, however, it is reasonable to look first at less complex final states, and this is how the field has progressed so far. This is particularly true in the case of peripheral production of mesons with exotic quantum numbers. For example, $\eta\pi$ and $\eta'\pi$ final states are attractive since any resonant odd-L partial wave would be manifestly exotic. Also, states with relatively few particles, all charged, are also attractive since photon detection (which is expensive and leads to poorer resolution and complex acceptance) is not needed and the partial-wave analysis is less complicated.

We take a historical approach to reviewing the experimental evidence for mesons with exotic quantum numbers. We are just now starting to see clear evidence of such phenomena, and it will take some time to sort out the misleading evidence of the past.

1. $\eta\pi$ final states and the $\hat{\rho}(1400)$

Any state decaying to $\eta\pi$ must have $J=L$ and $P=(-1)^L$ (since both the η and π are spinless) and $C=+$. Therefore any resonant, odd-L partial wave is good evidence for a meson state with exotic quantum numbers.[9]

This is obviously attractive experimentally, and a number of experiments have acquired data, mainly in peripheral production with π^- beams. As it happens, the $\eta\pi$ mass spectrum in $\pi^- N \to (\eta\pi)X$ reactions is dominated by the $a_2(1320)$ and this is both an advantage and a disadvantage. Obviously, one is relegated to picking out an exotic state from a spectrum dominated by a conventional $q\bar{q}$ meson. However, the D-wave a_2 provides a convenient wave against which an odd-L wave would interfere. In the $\eta\pi$ rest frame, this must lead to a forward-backward asymmetry in the decay, relative to the direction of the incoming π^-. Although such an asymmetry unambiguously implies the presence of an odd-L wave, it can only be shown to be resonant using a complete partial-wave analysis.

[9]We tacitly assume that C is a good quantum number, even for charged final states in which case C refers to the neutral member of a multiplet. In principle, a nonexotic 1^- meson could decay to $\eta\pi^\pm$ if isospin symmetry, and so G parity, were violated, but this should be a negligible effect for the cases discussed here.

Experiments searching for exotic resonance structure in $\eta\pi$ systems are listed in Table IV. Every experiment listed sees a clear forward-backward asymmetry, and therefore strongly implies the need for an odd-L partial wave. The asymmetry changes quickly in the region of the $a_2(1320)$ signal that dominates the mass spectrum, although details of the asymmetry differ from experiment to experiment. Each experiment except the first one listed (Apel *et al.*, 1981) suggests or claims outright evidence for an exotic $J^{PC}=1^{-+}$ resonance, although the only two consistent results are from VES (Beladidze *et al.*, 1993) and E852 (Thompson *et al.*, 1997). Finally, recent results from the Crystal Barrel at CERN (Abele *et al.*, 1998a, 1999) are consistent with these two peripheral experiments, yielding a mass of around 1400 MeV/c^2 and a width between 300 and 400 MeV/c^2. This is the state we call $\hat{\rho}(1400)$.

Both the NICE and GAMS experiments (see Table IV) work with all neutral final states, detecting only the four photons from $\eta\to\gamma\gamma$ and $\pi^0\to\gamma\gamma$. The recoil neutron is undetected, and in principle the target itself could be excited, decaying as $N^\star\to n\pi^0$ where the recoil π^0 decays to low-energy photons at large angles and is therefore undetected. The experimenters perform a constrained fit on the four observed photons, and only accept events with a reasonably high probability for the exclusive final state.

This is not a straightforward procedure, and therein lies one of the most crucial problems in these $\eta\pi$ experiments. If the detector acceptance is not well understood, then after applying this correction, one may end up with a forward-backward asymmetry that is due to instrumental effects and not physics. This is particularly difficult in the case of photon detection, and is less of a problem for BENKEI, VES, and E852, since they examine the $\eta\pi^-$ final state. Furthermore, VES detects the η via $\eta\to\pi^+\pi^-\pi^0$ decay, and E852 obtains results using both η decay modes.

The first positive result was claimed by GAMS (Alde *et al.*, 1988a), generating great interest in its theoretical interpretation (Close and Lipkin, 1987; Iddir *et al.*, 1988a, 1988b; Tuan, Ferbel, and Dalitz, 1988). From the start, however, it was clear that there were some internal inconsistencies in the experiment (Tuan, Ferbel, and Dalitz, 1988). To be brief, Alde *et al.*, (1988a) saw that both the $a_2(1320)$ and $\hat{\rho}(1400)$ were produced mainly in the unnatural-parity exchange waves, namely, the D_0 and P_0 waves, respectively. (Thompson *et al.*, 1997, include a detailed explanation of the notational convention.) This would imply $a_1(1260)$ exchange, as the a_1 is the lightest particle with the appropriate quantum numbers. However, as the $a_2(1320)$ is known to decay predominantly to $\rho\pi$, one would have expected production through natural-parity ρ^+ exchange with a π^- beam. Consequently their result was met with a good deal of skepticism. In fact, a reanalysis of their data by two of their collaborators (Prokoshkin and Sadovsky, 1995a, 1995b) obtains different conclusions.

The KEK Collaboration (Aoyagi *et al.*, 1993), which studied the $\eta\pi^-$ system at a much lower beam energy, analyzed their asymmetry to find the $a_2(1320)$ produced predominantly in the natural-parity D_+ wave, consistent with ρ^+ exchange. A P-wave resonance was suggested based on the shape of the intensity distributions, but it was seen with about equal strength in both the P_0 and P_+ waves. Furthermore, the mass and width of the "exotic" resonance were completely consistent with those of the $a_2(1320)$, and the relative phases of the P_+ and D_+ waves did not vary across the mass region of interest. These facts strongly suggest that the P_0 and P_+ waves were the result of "leakage" from the much stronger D_+ wave, caused by an imperfect knowledge of the experimental acceptance.

VES obtained a large-statistics data sample of both $\eta\pi^-$ and $\eta'\pi^-$ final states, using topologies with three charged pions and two photons (Beladidze *et al.*, 1993). The beam was incident on a beryllium target, however, and the recoils were not identified. Nevertheless, a clean signal was observed for $a_2(1320)$ production, strictly in the D_+ wave. A weak exotic P_+ wave was also observed, close to the mass of the $a_2(1320)$ but much broader, and with a significant phase motion relative to the P_+. No claims were made as to the existence of an exotic meson, but the data were certainly suggestive of this.

In E852, exclusivity of the final state was well established by tracking the recoil proton, rejecting events with a π^0 associated with the recoil, and requiring high probability with a constrained kinematic fit to the full final state (Thompson *et al.*, 1997). Statistics were achieved that were comparable to those of VES, and consistent D_+ and P_+ intensities and phase motion were observed. Figure 24 compares the results from the two experiments. The E852 group fit the waves simultaneously to two relativistic Breit-Wigner curves, and determined the mass and width of the $\hat{\rho}$ to be 1370 $\pm 50\,\mathrm{MeV}/c^2$ and $385\pm 100\,\mathrm{MeV}/c^2$, respectively. These fits are shown in Fig. 24.

A recent Crystal Barrel result (Abele *et al.*, 1998a) comes from $\bar{p}$ capture in liquid deuterium, the reaction being $\bar{p}d\rightarrow\eta\pi^-\pi^0 p$ where the final-state proton is a "spectator." Resonances in either $\eta\pi$ system recoil against the other pion, and $\pi^-\pi^0$ recoil against the η. The Dalitz plot is dominated by $\rho^-\rightarrow\pi^-\pi^0$, and there is a clear $\eta\pi$ P wave which interferes with it. The Dalitz plot is fit with a combination of resonances, and the χ^2 is 1.29 per degree of freedom when the $\hat{\rho}$ is included and 2.69 when it is removed. Abele *et al.* determine the mass

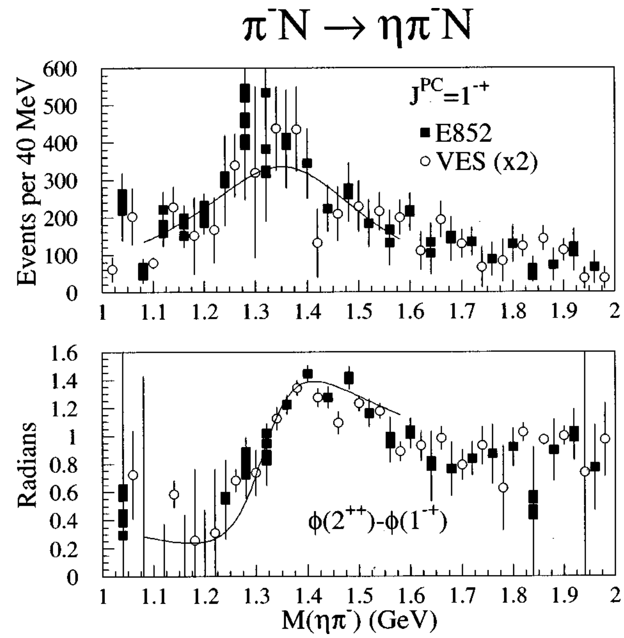

FIG. 24. A comparison of the exotic $J^{PC}=1^{-+}$ signal observed by the E852 (Thompson *et al.*, 1997) and VES (Beladidze *et al.*, 1993) experiments. The E852 data explicitly show the ambiguous solutions of the partial-wave analysis. The VES intensity distribution is multiplied by a factor of 2, and the phase difference is offset by $9\pi/10$. The solid line is a simultaneous fit to the E852 data of Breit-Wigner forms for the $a_2(1320)$ and $\hat{\rho}(1400)$, assuming appropriate interference between the D_+ and P_+ partial waves. The fit implies a mass and width of 1400 MeV/c^2 and 350 MeV/c^2, respectively, for the $\hat{\rho}(1400)$.

and width of the $\hat{\rho}$ to be 1400 ± 30 MeV/c^2 and 310 ± 70 MeV/c^2, respectively, quite consistent with the E852 result. Even more recently Abele *et al.* (1999) studied $\bar{p}p\rightarrow\eta\pi^0\pi^0$ in both liquid and gaseous hydrogen targets and found a consistent fit including a resonant P wave with mass 1360 ± 25 MeV and width 220 ± 90 MeV.

Curiously, the decay of an exotic qqg hybrid meson to $\eta\pi$ is highly suppressed if one assumes $SU(3)$ flavor symmetry and nonrelativistic degrees of freedom (Close and Lipkin, 1987; Iddir *et al.*, 1988a, 1988b; Tuan, Ferbel, and Dalitz, 1988; Page, 1997a). Furthermore, the contribution expected from breaking these assumptions should be small (Page, 1997a). If the $\hat{\rho}(1400)$ is confirmed as an exotic meson, there is therefore the theoretical prejudice that it may be a $q\bar{q}q\bar{q}$ state. A different suggestion (Donnachie and Page, 1998) is that the $\hat{\rho}(1400)$ is in fact an artifact of the $\hat{\rho}(1600)$ (see the next section), which manifests itself through the $\rho\pi$ and $b_1(1235)\pi$ thresholds.

2. The $\hat{\rho}(1600)$ in $\rho\pi$ and $\eta'\pi$

The E852 Collaboration has recently put forth evidence for another $J^{PC}=1^{-+}$ exotic meson (Adams *et al.*, 1998), decaying in this case to $\rho\pi$, in the reaction $\pi^-p\rightarrow\pi^-\pi^+\pi^-p$ at 18 GeV/c. This state, the $\hat{\rho}(1600)$, has a mass and width of 1593 ± 8 MeV/c^2 and 168 ± 20 MeV/c^2, respectively. Tentative evidence of the $\hat{\rho}(1600)$

was put forward by the VES Collaboration (Gouz *et al.*, 1993), who also see evidence for a broad but resonant P_+ wave near this mass in the $\eta'\pi^-$ system (Beladidze *et al.*, 1993). Although the $\hat{\rho}(1600)$ data do not agree particularly well with recent lattice QCD results and phenomenological expectations for a hybrid meson, it appears that it can still be accommodated as a hybrid meson (Page, 1997c).

The $\pi^-\pi^+\pi^-$ mass spectrum is dominated by the $a_1(1260)$, the $a_2(1320)$, and the $\pi_2(1670)$ (see Fig. 8). Each of these states is dominated by natural-parity exchange and decay to $\rho^0\pi^-$ or $f_2(1270)\pi^-$. The $\hat{\rho}(1600)$ is seen clearly in natural-parity exchange as well, decaying to $\rho^0\pi^-$, and its interference with the dominant waves shows clear phase motion, consistent with resonant behavior, in all cases. The peak intensity of the $\hat{\rho}(1600)$ is about 5% of the nearby $\pi_2(1670)$. There is also indication of a peak in intensity in the unnatural-parity waves, but there is not enough intensity in other unnatural-parity waves to check the phase motion.

One important feature of the E852 result is that "leakage" was checked explicitly. A data set was simulated, using the waves employed in the fit procedure but excluding the small $\hat{\rho}(1600)$. Then, a full partial-wave analysis was performed, including the 1^{-+} waves, to see if finite resolution and acceptance effects would lead to a spurious signal. Indeed, there was a significant leakage of the large $a_1(1260)$ signal into the 1^{-+} channel in the 1200–1300 MeV/c^2 region (Adams *et al.*, 1998), but the region near 1600 MeV/c^2 was clean.

While measuring the $\eta\pi^-$ final state, the VES Collaboration (Beladidze *et al.*, 1993) also studied the $\eta'\pi^-$ system in the reaction $\pi^-N\rightarrow\eta'\pi^-N$ with $\eta'\rightarrow\eta\pi^+\pi^-$. The $\eta'\pi^-$ mass spectrum shows a clear peak at the $a_2(1320)$ (despite the limited phase space) superimposed on a broad distribution peaked at ~1.6 GeV/c^2. In a situation rather similar to the $\eta\pi^-$ system, the partial-wave analysis finds the reaction almost completely dominated by the natural-parity exchange P_+ and D_+ waves. The peak in the mass distribution near the $a_2(1320)$ is completely absorbed into the D_+ wave, and a Breit-Wigner parametrization yields the correct values for the mass and width. (This analysis also yields the relative branching ratio for $a_2\rightarrow\eta'\pi$ to $a_2\rightarrow\eta\pi$ to be ~5%.) The relative phase motion of the P_+ and D_+ waves also supports the existence of the $\hat{\rho}(1400)$ seen in $\eta\pi^-$, and the matrix elements they extract are quite comparable in this mass region. The matrix elements, however, are considerably stronger for $\eta'\pi^-$ than for $\eta\pi^-$ over most of the accepted $\eta'\pi$ masses, particularly in the ~1.6-GeV/c^2 region. However, the structure appears to be quite broad and the phase motion is not distinctive so it is difficult to associate this with the $\hat{\rho}(1600)$ seen in $\rho\pi$ by E852 (Adams *et al.*, 1998).

The E852 data indicate the $BR[\hat{\rho}(1600)\rightarrow\rho\pi]$ is about 20% (Page, 1997c), while there is no indication of this state in the $\eta\pi^-$ data (Beladidze *et al.*, 1993; Thompson *et al.*, 1997), and clearly not all of the $\eta'\pi^-$ signal (Beladidze *et al.*, 1993) is consistent with the $\hat{\rho}(1600)$. It is therefore likely that its dominant decay mode has yet to be observed, and Page (1997c) suggests that this might be $b_1(1235)\pi$ or $f_1(1285)\pi$ since b_1 or f_1 exchange would account for the unnatural-parity exchange signal observed by E852.

3. Searches for $S+P$ decays

As noted previously, rather general symmetrization selection rules (Page, 1997b) argue that hybrid mesons, exotic or otherwise, should not decay predominantly to a pair of S-wave mesons. Consequently the cases discussed in the previous two sections, namely, $\eta\pi$, $\eta'\pi$, and $\rho\pi$, should not represent the dominant decay modes of the $\hat{\rho}(1400)$ or $\hat{\rho}(1600)$, if those states are indeed hybrid, exotic mesons.

Decays to $S+P$ pairs of mesons are rather complex, however, and until now only a very limited number of experiments have attempted to search for exotic mesons in these signatures. In response to the suggestion of Isgur, Kokoski, and Paton (1985) that exotic hybrids decaying to $b_1(1235)\pi$ should be produced with good signal-to-noise in peripheral photoproduction, the Ω spectrometer group at CERN reexamined their previous data (Atkinson *et al.*, 1983) on the reaction $\gamma p\rightarrow\omega\pi^+\pi^-p$. This group showed that production of $b_1(1235)\rightarrow\omega\pi$ was enhanced when $1.6\leqslant\text{mass}(\omega\pi\pi)\leqslant 2.0\,\text{GeV}/c^2$. Although the number of events was limited, their reanalysis (Atkinson *et al.*, 1987) showed what appeared to be production of $\omega_3(1670)$ and a new state at $\sim 1.9\,\text{GeV}/c^2$, decaying to $b_1(1235)\pi$. This was based strictly on the invariant-mass distribution; no partial-wave or angular distribution analyses were performed.

One would prefer to search for $S+P$ wave decays for which the P-wave meson is not too broad, and which decay mainly to experimentally accessible final states. Good candidates include $b_1(1235)\rightarrow\omega\pi$ (with total width Γ=142 MeV/c^2), $a_2(1320)\rightarrow\rho\pi$ (Γ=107 MeV/c^2), and $f_1(1285)\rightarrow a_0(980)\pi$ (Γ=25 MeV/c^2).

At this time, results have only been presented for the $f_1(1285)\pi$ final state. These are from the VES collaboration (Gouz *et al.*, 1993), who made use of the $f_1\rightarrow\eta\pi^+\pi^-$ final state, and BNL experiment E818 (Lee *et al.*, 1994), for which $f_1\rightarrow K^+\overline{K}_0\pi^-$. In both cases, there are not very many events, and the VES result remains preliminary. However, the results of E818 show that the intensity distribution and phase motion of the $J^{PC}=1^{-+}$ signal are suggestive of the presence of an exotic meson with mass around 1.9 GeV/c^2 and a $J^{PC}=1^{++}$ resonance near 1.7 GeV/c^2. As stated by Lee *et al.* (1994), additional data are required for a more complete understanding of the $J^{PC}=1^{-+}$ wave.

There have been two attempts to search for $S+P$ final states in high-energy photoproduction since the original CERN data (Atkinson *et al.*, 1987). One of these, experiment E687 at Fermilab, performed an inclusive search for $f_1(1285)\pi^{\pm}$ states (Danyo, 1995) in the reaction $\gamma\text{Be}\rightarrow f_1(1285)\pi^{\pm}X$ with $f_1(1285)\rightarrow a_0(980)^{\pm}\pi^{\mp}$ and $a_0(980)^{\pm}\rightarrow K_SK^{\pm}$. A tagged bremsstrahlung beamline provided photons with energies near 200 GeV. A peak in the $f_1\pi$ mass distribution is clear above the background, which was measured using combinations of f_1 and π taken from different events. The fit yields a mass and width of 1748±12 MeV/c^2 and 136±30 MeV/c^2, respectively, not inconsistent with that suggested by E818. The angular distribution in the helicity frame is not inconsistent with $J^{PC}=1^{-+}$, although the statistics are quite poor.

The only other such photoproduction experiment is a reanalysis by the Tennessee group (Blackett *et al.*, 1997) of data taken with the SLAC Hybrid Photoproduction experiment (Abe, 1984). Data were collected for the reaction $\gamma p \rightarrow p\omega\pi^+\pi^-$ in a search for states decaying to $b_1^{\pm}(1235)\pi^{\mp}$ with $b_1^{\pm}(1235)\rightarrow\omega\pi^{\pm}$. However, most of the sample was in fact more consistent with the charge-exchange reaction $\gamma p\rightarrow\Delta^{++}b_1^-(1235)$ and no significant structure was observed in the $b_1^{\pm}(1235)\pi^{\mp}$ mass spectrum when these events were removed.

B. The scalar mesons

The scalar ($J^P=0^+$) mesons have long been a source of controversy (Morgan, 1974). As discussed in Secs. IV.B.3, IV.B.4, and IV.B.5, it is difficult to compare observed states to the theoretical predictions in this sector. For masses below 2 GeV/c^2, the two-body S-wave structure is very complicated, and overlapping states interfere with each other differently in different production and decay channels. Interpreted in terms of resonances, the quark model is clearly oversubscribed, and only very recently has a semblance of order emerged. This subject has already been touched on in Sec. IV; a brief technical review is given by the Particle Data Group (Caso *et al.*, 1998).

We shall take a different approach here and focus directly on the observed states where the $J^{PC}=0^{++}$ assignment is well established. We note that this discussion is linked to the $f_J(1710)$ and $f_J(2220)$, which are discussed separately in Sec. V.C.1.

The present bias is that the quark-model nonet is tentatively filled, with only the predominant $s\bar{s}$ state in need of confirmation, and that there are two manifestations of degrees of freedom beyond the quark model. These are the $f_0(1500)$, suggested as being dominated by gluonic degrees of freedom (Amsler and Close, 1996), and the $a_0(980)$ and $f_0(980)$, which are candidates for multiquark states (Jaffe, 1977a; Rosner, 1983; Weinstein and Isgur, 1983). Still, there is significant controversy especially regarding the possibility of a broad, low-mass isoscalar state (Janssen *et al.*, 1995; Törnqvist and Roos, 1996).

The strange members of the scalar nonet are the $K_0^\star(1430)$ states, clearly established by the LASS collaboration (Aston *et al.*, 1988e). The $I=1$ member is most likely the $a_0(1450)$, observed (Amsler *et al.*, 1994a) by the Crystal Barrel Collaboration, who also measured its decay branches to $\eta'\pi$ (Abele *et al.*, 1997a) and to $K\bar{K}$ (Abele *et al.*, 1998b). The broad, isoscalar $\pi\pi$ S wave has within it three clear resonances, namely, the $f_0(980)$, $f_0(1370)$, and $f_0(1500)$. We focus our discussion on the isoscalar states near 1500 MeV/c^2, mainly the $f_0(1500)$, and on the enigmatic $a_0(980)$ and $f_0(980)$ mesons. See also Amsler (1998).

1. The $f_0(1500)$

The $f_0(1500)$ has been touted as a likely candidate for a "glueball" (Amsler and Close, 1995, 1996). The evidence is circumstantial, based on its peculiar width and decay properties. This suggestion is supported by various theoretical calculations (see Sec. II.E.1 as well as Szczepaniak *et al.*, 1996) which indicate that the lightest glueball should have $J^{PC}=0^{++}$ and a mass near 1500 MeV/c^2. The scalar nonet is *not* overpopulated [discounting the $a_0(980)$ and $f_0(980)$] until the $s\bar{s}$ state is unambiguously identified.

Most of the data on the $f_0(1500)$ are from the Crystal Barrel Collaboration, who resolved two scalar states in this mass region (Anisovich *et al.*, 1994) and who determined its decay branches to a number of final states, including $\pi^0\pi^0$ and $\eta\eta$ (Amsler and Close, 1995; Abele *et al.*, 1996c), $\eta\eta'$ (Amsler *et al.*, 1994b), K_LK_L (Abele *et al.*, 1996b), and $4\pi^0$ (Abele *et al.*, 1996a), all using $\bar{p}p$ annihilation at rest in liquid hydrogen. The $f_0(1500)$ has also been observed by the OBELIX experiment (Bertin *et al.*, 1998a) in the $\bar{n}p\rightarrow\pi^+\pi^+\pi^-$ reaction for neutrons in flight, by the WA91 (Antinori *et al.*, 1995) and WA102 (Barberis *et al.*, 1997b) Collaborations, in "glue-rich" central production reactions $pp\rightarrow p_f(\pi^+\pi^-\pi^+\pi^-)p_s$ and $pp\rightarrow p_f(\pi^+\pi^-)p_s$, and most recently in Fermilab experiment E690 (Reyes *et al.*, 1998). There is some evidence of $J/\psi\rightarrow\gamma f_0(1500)$ with $f_0(1500)\rightarrow\pi^+\pi^-\pi^+\pi^-$ (Bugg *et al.*, 1995) and perhaps also with $f_0(1500)\rightarrow\pi^+\pi^-$ (Dunwoodie, 1997), but S-wave structure in the K^+K^- and K_SK_S final states seems absent in this mass region for J/ψ radiative decay (Bai *et al.*, 1996a; Dunwoodie, 1997). For a collective discussion of these data and the evidence that the $f_0(1500)$ is a glueball, see Amsler and Close (1995, 1996), Close (1997); Close, Farrar, and Li (1997); and Amsler (1998). No evidence has been reported in $\gamma\gamma$ collisions [as would be expected to be the case if the $f_0(1500)$ were predominantly glue], but the $\pi\pi$ and $K\bar{K}$ mass distributions are dominated by the $f_2(1270)$ and $f_2'(1525)$ in this mass region (Morgan, Pennington, and Whalley, 1994), making a search for scalar structure quite difficult.

The mass and width of the $f_0(1500)$ are given (Caso *et al.*, 1998) as 1500±10 MeV/c^2 and 112±10 MeV/c^2, respectively. The mass is rather close to the nominal $I=1$ and $I=1/2$ members of the scalar nonet, i.e., the $a_0(1450)$ and the $K_0^\star(1430)$. However, it is decidedly narrower, with $\Gamma[a_0(1450)]=265\pm13$ MeV/c^2 and $\Gamma[K_0^\star(1450)]=287\pm20$ MeV/c^2, which are more in line with the very broad $f_0(1370)$. This comparison and the penchant of the $f_0(1500)$ to be produced in "glue-rich" environments hint at the special status of $f_0(1500)$ as something different from a standard $q\bar{q}$ meson.

Of course, if we take the $f_0(980)$ and $a_0(980)$ to be multiquark states, then there should be *two* isoscalar 0^{++} mesons with a mass typical of the 0^+ nonet, and it is tempting to take these to be the $f_0(1370)$ and $f_0(1500)$. This is clearly not workable, however, even disregarding the narrow width of the $f_0(1500)$, since neither can convincingly be associated with the predominantly $s\bar{s}$ member, based on the decay patterns of the $f_0(1370)$ and $f_0(1500)$.

It is difficult to assign specific branching ratios to the $f_0(1370)$ because not only is it broad, but it also interferes strongly with the even broader underlying S-wave $\pi\pi$ structure as well as with the $f_0(980)$ (Amsler, 1998; Caso *et al.*, 1998). It is nevertheless clear that this state couples mainly to pions, and $K\bar{K}$ is suppressed (Amsler, 1998). Given that its width is in line with the $a_0(1450)$ and $K_0^\star(1430)$, we associate it with the $n\bar{n}\equiv[u\bar{u}+d\bar{d}]/\sqrt{2}$ member of the nonet.

It is possible to derive branching ratios for the $f_0(1500)$ (Amsler, 1998). From a coupled-channel analysis (Amsler *et al.*, 1995a) of $\bar{p}$ annihilation at rest in liquid hydrogen to $\pi^0\pi^0\pi^0$, $\eta\pi^0\pi^0$, and $\eta\eta\pi^0$ final states, the Crystal Barrel Collaboration determine the following branching fraction for production and decay of $f_0(1500)$:

$$B[\bar{p}p\to f_0(1500), f_0(1500)\to\pi^0\pi^0]=(12.7\pm3.3)\times10^{-4},$$

$$B[\bar{p}p\to f_0(1500), f_0(1500)\to\eta\eta]=(6.0\pm1.7)\times10^{-4}.$$

An analysis (Amsler *et al.*, 1994b) by the Crystal Barrel Collaboration of the reaction $\bar{p}p\to\pi^0\eta\eta'$ gives

$$B[\bar{p}p\to f_0(1500), f_0(1500)\to\eta\eta']=(1.6\pm0.4)\times10^{-4}.$$

They also determine (Table II in Amsler *et al.*, 1998)

$$B[\bar{p}p\to f_0(1500), f_0(1500)\to K_LK_L]$$
$$=(1.13\pm0.09)\times10^{-4}$$

from data (Abele *et al.*, 1996b) on the reaction $\bar{p}p\to\pi^0K_LK_L$, using data from the reaction $\bar{p}p\to K_LK^\pm\pi^\mp$ (Abele *et al.*, 1996a) to fix the contribution from $a_0(1450)\to K_LK_L$.

A simple phenomenological model makes it possible to derive couplings of the $f_0(1500)$ to various pairs of pseudoscalar mesons (Amsler and Close, 1996). The model incorporates $SU(3)$ flavor symmetry breaking and meson form factors, and is tested on decays of the well-understood 2^{++} nonet. Following this model and incorporating two-body phase space, we arrive at the following relative decay rates for $f_0(1500)$ into two-body, pseudoscalar meson pairs:

$$\pi\pi:KK:\eta\eta:\eta\eta'$$
$$=(5.1\pm2.0):(0.71\pm0.21):(\equiv1.0):(1.3\pm0.5), \qquad (30)$$

where the value for $\pi\pi$ (KK) multiplies the branching ratio for $\pi^0\pi^0$ (K_LK_L) by 3 (4) to account for charge combinations. These are obviously *inconsistent* with the $f_0(1500)$'s being the $s\bar{s}$ member of the nonet. In fact, detailed considerations of isoscalar mixings (Amsler and Close, 1996; Amsler, 1998) show that they are consistent with the $f_0(1500)$'s being the $n\bar{n}$ member, although we have already argued that the $f_0(1370)$ is a better candidate.

Therefore the circumstantial evidence for $f_0(1500)$'s being the scalar glueball is clear. It is produced primarily in glue-rich environments; its mass and width are consistent with theoretical predictions; and it overpopulates the $q\bar{q}$ states in this mass region if we assume it is not the $s\bar{s}$ state. Direct evidence, however, is lacking. If the $f_0(1500)$ were a pure glueball, one would naively expect "flavor-blind" decays to all available $SU(3)_f$ singlets, and as discussed in Sec. II.E.1 the couplings (30) would be 3:4:1:0 (ignoring single/octet mixing in the η and η'). Most notable here is the strongly suppressed $K\bar{K}$ coupling relative to $\pi\pi$, whereas the naive prediction is that it should be comparable if not larger.

Amsler and Close (1996), argue that this problem is linked to the other outstanding problem in the isoscalar 0^+ nonet, namely, the missing $s\bar{s}$ state, i.e., the $f_0'(\sim1600)$. Using first-order perturbation theory, one finds that $f_0(1500)\to K\bar{K}$ can be strongly suppressed by mixing between $f_0(1370)$, $f_0(1500)$, and the hypothetical $f_0'(\sim1600)$. In fact, they find that if pure glue is indeed flavor blind with respect to $s\bar{s}$ and $n\bar{n}$, then $f_0(1500)\to K\bar{K}$ goes to zero if the $f_0(1500)$ lies exactly between the other two states in mass. Turning this analysis around, the $f_0(1500)\to K\bar{K}$ branch above infers two possible values for the mass of the f_0', namely, 1600 or 1900 MeV/c^2 (Amsler, 1998).

It is essential to confirm the existence of an $s\bar{s}$ f_0' and to measure its decay properties in order to clearly establish the $f_0(1500)$ as the scalar glueball (see Sec. IV.B.3). There are two candidates at present. One possibility is the tentative identification (Aston *et al.*, 1988a) of an S-wave resonance, produced and decaying through $K\bar{K}$ in the reaction $K^-p\to K_S^0K_S^0\Lambda$, directly underneath the dominant $f_2'(1525)$. Not only is this a weak observation, however, it may in fact be an observation of the $f_0(1500)$ itself. Another candidate is the $f_J(1710)$ (Sec. V.C.1). Although its spin assignment is somewhat uncertain and controversial, it shows some characteristics of being a glue-dominated state.

2. The $a_0(980)$ and $f_0(980)$

These states have been known for a very long time (Morgan, 1974) but their nature continues to generate controversy (Janssen *et al.*, 1995). With the establishment of the $a_0(1450)$, $f_0(1370)$, and $f_0(1500)$ (Amsler, 1998), it is no longer feasible to argue that the $a_0(980)$ and $f_0(980)$ are members of the $q\bar{q}$ scalar nonet. Their near degeneracy in mass, as well as their proximity to the $K\bar{K}$ threshold and their propensity to decay to $K\bar{K}$, strongly suggest that they are $I=1$ and $I=0$ bound states of $K\bar{K}$ (Weinstein and Isgur, 1983).

Because of their very peculiar decay properties, it is difficult to quantify even the mass and width of these states (Caso *et al.*, 1998). The states are somewhere between 50 and 100 MeV/c^2 wide, so their nominal mass allows the width to straddle the $K\bar{K}$ threshold at 990 MeV/c^2. The non-$K\bar{K}$ decays are fully dominated by $f_0(980)\to\pi\pi$ and $a_0(980)\to\eta\pi$, which are not significantly suppressed by their own kinematic thresholds. A recent measurement by the E852 Collaboration (Teige *et al.*, 1999) gives the charged $a_0(980)$ mass and width as 995.8 ± 1.6 MeV/c^2 and 62 ± 6 MeV/c^2, respectively, when the $\eta\pi^{\pm}$ final state is fit to a relativistic Breit-Wigner curve, and 1001.3 ± 1.9 MeV/c^2 and 70 ± 5 MeV/c^2, based on the coupled-channel description developed by Flatté (1976).

The $\gamma\gamma$ decay widths (Caso *et al.*, 1998) have been measured in photon-photon collisions (Morgan, Pennington, and Whalley, 1994), but theoretical estimates vary widely and it is difficult to make a definitive statement (Barnes, 1985b; Antreasyan *et al.*, 1986).

A novel measurement to elucidate the nature of these states was suggested by Close, Isgur, and Kumano (1993). By determining the radiative decay rate $\phi\to a_0(980)\gamma$ or $\phi\to f_0(980)\gamma$, one could infer the $s\bar{s}$ content of the a_0 or f_0 wave function, since the rate is proportional to the overlap with the ϕ, a well-known $s\bar{s}$ state. They calculate that $BR(\phi\to a_0\gamma)\approx BR(\phi\to f_0\gamma)\approx4\times10^{-5}$ if the a_0 and f_0 are indeed $K\bar{K}$ molecules, whereas the branching ratio should be 10^{-6} or less for $q\bar{q}$ or other multiquark configurations. Similar results are obtained by Achasov, Gubin, and Shevchenko (1997) and Achasov and Gubin (1997), who incorporate the peculiar line shape of the a_0 and f_0. Recent results from the SND Collaboration running at the VEPP-2M storage ring in Novosibirsk (Achasov, 1998b) yield a value $BR(\phi\to f_0\gamma)=(3.42\pm0.30\pm0.36)\times10^{-4}$, much larger than expected for a $K\bar{K}$ molecule. This collaboration also reports (Achasov *et al.*, 1998a) that $BR(\phi\to\eta\pi^0\gamma)\approx(0.83\pm0.23)\times10^{-4}$ but with no suggestion of a peak at the $a_0(980)$. The CMD-2 Collaboration (Akhmetshin *et al.*, 1997), also at VEPP-2M, searched for $\phi\to f_0\gamma\to\pi^+\pi^-\gamma$ and reported consistent upper limits, including interference from radiative corrections to $e^+e^-\to\pi^+\pi^-$ (Achasov, Gubin, and Solodov, 1997). A new experiment using ϕ photoproduction on hydrogen (Dzierba, 1994) will take data on $\phi\to\gamma X$ in 1999. It will take some time to sort out all the new information, but it is reasonable to expect that important new results will be available soon.

C. Other possible glueballs: The $f_J(1710)$ and the $f_J(2220)$

Focus on the $f_0(1500)$ as the lightest glueball comes partly because of the wealth of information provided on this state by the Crystal Barrel experiment. A number of branching ratios have been measured, many including complicated, all-neutral final states, and this has made detailed analyses possible. These analyses have led to clear inconsistencies with what is expected from the scalar $q\bar{q}$ nonet, and that has fueled the conjecture.

It is worth noting, however, that the search for glueballs began much earlier, using J/ψ radiative decays. As discussed in Sec. III.E.3, the annihilation of a $c\bar{c}$ pair

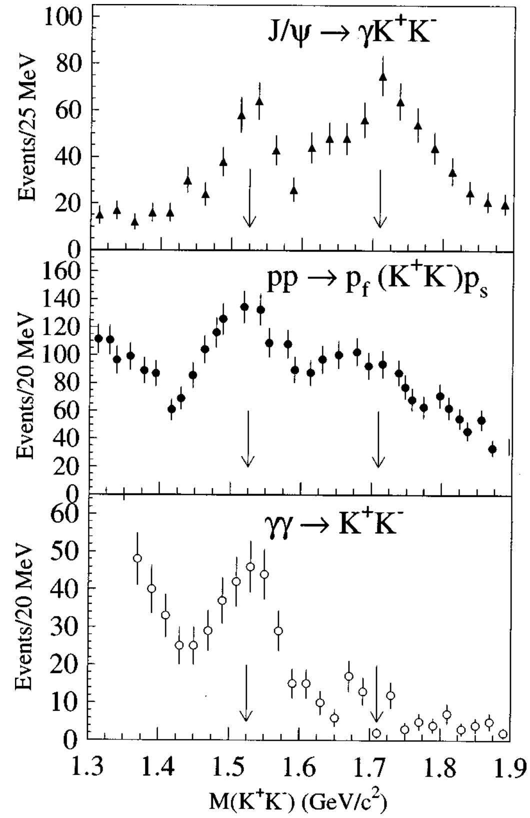

FIG. 25. Possible glueball sensitivity in different reactions, for states decaying to K^+K^-. Shown are the K^+K^- invariant-mass distributions for $J/\psi\to\gamma K^+K^-$ (Bai *et al.*, 1996b), $pp\to p_f(K^+K^-)p_s$ (Armstrong *et al.*, 1991a), and $\gamma\gamma\to K^+K^-$ (Albrecht *et al.*, 1990). The arrows mark the positions of the $f_2'(1525)$ and the $f_J(1710)$.

into a photon and two gluons is certainly expected in lowest-order QCD, so one is naturally led to search for gluonic states in these decays. This was discussed in some detail recently by Close, Farrar, and Li (1997) and Page and Li (1998). Further information is provided by the coupling, or its upper limit, of the candidate state to $\gamma\gamma$ (Sec. III.E.2). Pure glue can only couple to photons through the creation of an intermediate $q\bar{q}$ pair and is therefore suppressed relative to quark-model states. On the other hand, central production in pp collisions (Sec. III.D) may be a rich source of gluonic states, based on speculation that it proceeds through double-Pomeron exchange (Close, 1997; Close and Kirk, 1997).

These processes are shown in Fig. 25 for excited mesons decaying to K^+K^-. The figure shows the invariant mass of K^+K^- pairs produced in radiative J/ψ decay, $J/\psi\to\gamma K^+K^-$ (Bai *et al.*, 1996b); central pp collisions, $pp\to p_f(K^+K^-)p_s$ (Armstrong *et al.*, 1991a); and two-photon collisions, $\gamma\gamma\to K^+K^-$ (Albrecht *et al.*, 1990). For both radiative J/ψ decay and for central pp collisions, two enhancements are clear, one near 1500 MeV/c^2 and the second near 1700 MeV/c^2. The 1500-MeV/c^2 structure is consistently found to be dominated by $J=2$ and is identified as the $f_2'(1525)$, the mainly $s\bar{s}$ member of the tensor nonet. The 1700-MeV/c^2 structure contains the $f_J(1710)$. Note that this second structure is *not* seen in two-photon collisions, and this has fueled speculation that the $f_J(1710)$ is a glueball. [There is a preliminary result for $\gamma\gamma \to f_J(1710)$, suggesting that $J=0$, from the L3 experiment at LEP; see Braccini (1998).] The relative ratios of the abundances of $f_2'(1525)$ and $f_J(1710)$ for radiative J/ψ decay and two-photon collisions can be used to evaluate the "stickiness" (Chanowitz, 1984) of these two enhancements. Recall from Sec. II.E.1 [Eq. (25)] that this quantity is proportional to the ratio of squared matrix elements for coupling of the state to gg and $\gamma\gamma$. For a state of pure glue, there would be no coupling to photons, and the state would have infinite stickiness.

Both $J/\psi \to \gamma X$ and $\gamma\gamma \to X$ can only produce states X with $C=+1$. Furthermore, in this section, our discussion is mainly limited to decays to pairs of identical psuedoscalar particles, so $P=+1$ and the total spin J must be even. There is in fact considerable controversy regarding the total spin J of the $f_J(1710)$ and $f_J(2220)$, hence the indeterminate nomenclature.

1. The $f_J(1710)$

The $f_J(1710)$ is the main competitor of the $f_0(1500)$ for status as the lightest glueball, assuming that $J=0$. Our best estimates for glueball properties are from lattice gauge theory calculations, and although they all agree that the lightest glueball should have $J^{PC}=0^{++}$, there is some disagreement on the mass. For example, two comprehensive studies find $M(0^{++})=1550\pm 50$ MeV/c^2 and $M(2^{++})=2270\pm 100$ MeV/c^2 (Bali *et al.*, 1993) and $M(0^{++})=1740\pm 71$ MeV/c^2 and $M(2^{++}) =2359\pm 128$ MeV/c^2 (Chen *et al.*, 1994) for the lowest-mass scalar and tensor glueballs. In fact, one of these groups compares the measured properties of the $f_J(1710)$ to their calculations and directly argue that it must be the lightest scalar glueball (Sexton, Vaccarino, and Weingarten, 1995a; see Sec. II.E.1 for more details.) In all cases, however, the tensor mass remains in the region near 2.2 GeV/c^2.

It is also important to recall that in order to accommodate a scalar glueball anywhere in the 1.5–1.7 GeV/c^2 region, one needs to identify the $s\bar{s}$ partner to the $n\bar{n}$ $f_0(1370)$. If $J=0$, then the $f_J(1710)$ and $f_0(1500)$ might well represent the glueball and the $s\bar{s}$ state, or more likely each is a mixture of both. Recenty, the IBM group has computed mixing with quarkonia (Weingarten, 1997; Lee and Weingarten, 1998a, 1998b) and again claim good agreement with the $f_J(1710)$ as mainly the 0^{++} glueball, while establishing that the $f_0(1500)$ is a good candidate for the mainly $s\bar{s}$ member of the nonet. (It remains to find a mixing scheme to explain the peculiar branching ratios; see, however, Burakovsky and Page, 1999.) If $J=2$, however, it will be difficult to assign a glueball status to the $f_J(1710)$ since that would be at odds with all current lattice gauge calculations.

The Particle Data Group (Caso *et al.*, 1998) estimates the mass and width of the $f_J(1710)$ to be 1712 ± 5 MeV/c^2 and 133 ± 14 MeV/c^2, respectively. The differing experimental results for the spin of this state are clearly intertwined with determination of its other properties. We shall therefore go through the experimental evidence for the $f_J(1710)$ and point out the various important agreements and disagreements. Controversy still remains, and there is some suggestion that the $f_J(1710)$ is actually more than one state.

a. Radiative J/ψ decay

This state was first observed by the Crystal Ball Collaboration in radiative J/ψ decay (Edwards *et al.*, 1982b). It was immediately recognized as a glueball candidate. Called $\theta(1640)$, it was seen as a peak in the $\eta\eta$ mass distribution of 39 ± 11 events over background. The width was large (~ 220 MeV/c^2) and the $\theta \to \eta\eta$ angular distribution favored $J=2$, but the analysis did not include the presence of the $f_2'(1525)$, which decays $\sim 10\%$ of the time to $\eta\eta$. Soon afterwards, however, a consistent peak was observed in K^+K^- mass by the Mark II Collaboration (Franklin, 1982) in the reaction $J/\psi \to \gamma K^+K^-$ and this analysis did include the $f_2'(1525)$, again slightly favoring $J=2$. The simultaneous observation of a state decaying both to $\eta\eta$ and to $K\bar{K}$ fueled speculation that this was a glueball. A reanalysis of the Crystal Ball data that included the $f_2'(1525)$ (Bloom and Peck, 1983; Königsmann, 1986) led to a larger mass, near 1700 MeV/c^2, and a smaller width, but $J=2$ was still preferred.

The Mark III Collaboration followed up with higher-statistics measurements of the charged particle decays $\pi^+\pi^-$ and K^+K^- in the reactions $J/\psi \to \gamma K^+K^-$ and $J/\psi \to \gamma\pi^+\pi^-$ (Baltrusaitis *et al.*, 1987). The $f_J(1710)$ was observed in both modes. The K^+K^- mode was particularly clean, apparently obstructed only slightly by the nearby $f_2'(1525)$, and once again the angular distribution preferred $J=2$. Peaks were confirmed by the DM2 Collaboration in the $\pi^+\pi^-$ (Augustin *et al.*, 1987) and K^+K^- (Augustin *et al.*, 1988) channels, as well as in K_SK_S (Augustin *et al.*, 1988).

The published Mark III result suggesting that $J=2$ (Baltrusaitis *et al.*, 1987) has actually been called into question in unpublished reports by the same collaboration (Chen, 1990, 1991; Dunwoodie, 1997). This is the result of a separate analysis, using a somewhat larger event sample (5.8×10^6 J/ψ decays as compared to 2.7×10^6) and also including $J/\psi \to \gamma K_SK_S$ decays. The key difference in the analyses, however, is that Baltrusaitis *et al.*, (1987) assumed that the peaks at 1525 MeV/c^2 and 1700 MeV/c^2 (Fig. 25) consisted of pure resonances and compared observed and predicted angular distributions for $J=0$ and $J=2$. That is, interference effects in the amplitudes were unaccounted for. However, an analysis of moments (see Dunwoodie, 1997) clearly demands the presence of S wave in the 1710 MeV/c^2 region for both $\pi\pi$ and $K\bar{K}$ final states. The full-amplitude analysis, shown in Fig. 26, finds that $K\bar{K}$ is nearly entirely S wave in the $f_J(1710)$ region [with a clear D-wave signal for the $f_2'(1525)$]. In $\pi^+\pi^-$ this analysis clearly sees the

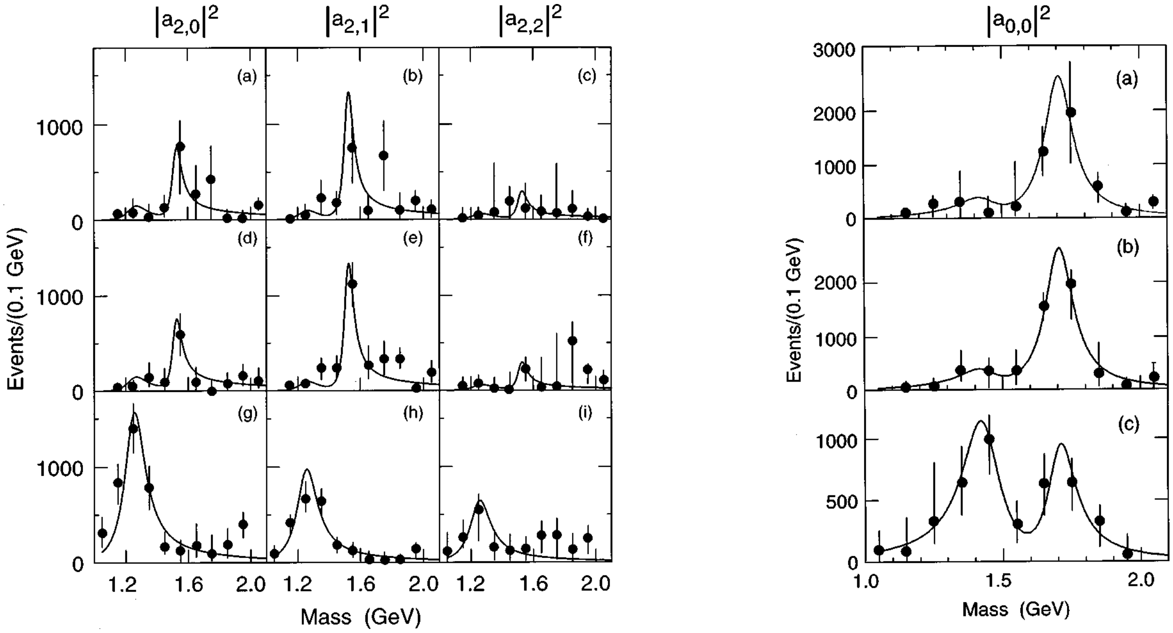

FIG. 26. Partial-wave analysis of J/ψ radiative decay to two pseudoscalar mesons, from the Mark III Collaboration. The analysis allows all partial wave components to vary independently for both D wave (left) and S wave (right). The top row is the analysis of $J/\psi \to \gamma K_S K_S$, middle row for $J/\psi \to \gamma K^+K^-$, and bottom row for $J/\psi \to \gamma\pi\pi$. The structure near 1700 MeV/c^2 is clearly dominated by an S-wave state.

$f_2(1270)$ in D wave, and confirms the $K\bar{K}$ S-wave result near 1700 MeV/c^2 while also identifying what appears to be a scalar with mass near 1400 MeV/c^2. This result in particular would argue for the $f_J(1710)$'s consisting mainly of the scalar glueball, as its decays to $\pi\pi$ preclude identifying it as the $s\bar{s}$ scalar nonet member.

Other analyses of J/ψ radiative decay, leading to 4π decay modes of the $f_J(1710)$, likewise did not support the original $J=2$ assignment. Baltrusaitis *et al.* (1986a) observed a strong enhancement of $\rho^0\rho^0$ and $\rho^+\rho^-$ masses in the 1.4–2.0 GeV/c^2 region, well above phase space. The angular analysis strongly argued that a single $J^{PC}=0^{-+}$ component dominated this region. A recent reanalysis of these data (Bugg *et al.*, 1995) includes $\sigma \equiv (\pi\pi)_S$ components, as well, and finds this region populated with both $J^{PC}=2^{++}$ states, decaying mainly to $\rho\rho$, and $J^{PC}=0^{++}$ states, decaying to $\sigma\sigma$.

Two-body decays of the $f_J(1710)$ have been recently reexamined with new data on K^+K^- from the BES Collaboration (Bai, 1996b). BES finds the $f_J(1710) \to K^+K^-$ region dominated by a 2^{++} state near 1700 MeV/c^2, but also resolves a 0^{++} state at 1780 MeV/c^2 with about half the strength of the 2^{++}. However, this analysis is more tightly constrained than the Mark III analysis by Chen (1990, 1991) and Dunwoodie (1997). In particular, the D-wave amplitudes are forced to be relatively real, reducing the number of unknowns that need to be determined from the fit. Bai *et al.*'s analysis shows that there are indeed more than one overlapping D-wave state in this region, and the assumption that the D waves are relatively real is therefore questionable.

We conclude that the $f_J(1710)$ region, as observed in two-body modes in J/ψ radiative decay, is most likely dominated by a single $J^{PC}=0^{++}$ state. Branching ratios for different modes are given by Chen (1990, 1991), and Dunwoodie (1997). Their analysis (Fig. 26) gives

$$\frac{\Gamma[f_{J=0}(1710)\to\pi\pi]}{\Gamma[f_{J=0}(1710)\to K\bar{K}]}=0.27^{+0.17}_{-0.12}.$$

b. Central production in pp collisions

Supposing that central pp collisions through double-Pomeron exchange, one might suspect that glue-rich states are produced in these reactions (see Sec. III.D). The $f_J(1710)$ has been studied in this way at CERN by WA76 (Armstrong, 1989a, 1991a) and WA102 (Barberis *et al.*, 1997b), and also at Fermilab (Reyes *et al.*, 1998). Final states include $\pi^+\pi^-$ (Armstrong *et al.*, 1991a), K^+K^- and K_SK_S (Armstrong *et al.*, 1989a; Reyes *et al.*, 1998), and $\pi^+\pi^-\pi^+\pi^-$ (Barberis *et al.*, 1997b). As can be seen in Fig. 25, there is a clear enhancement of K^+K^- in the region of the $f_J(1710)$. The shape of the mass spectrum is quite sensitive to momentum transfer, with the $f_J(1710)$ region enhanced for more peripheral reactions, i.e., where Pomeron exchange is expected to dominate. The 4π spectrum shows a clear peak associated with $f_1(1285)$ and other peaks at 1440 and 1920 MeV/c^2. Again, the shape depends very much on the momentum transfer.

The angular distribution of the two-body decays (Abatzis *et al.*, 1994) seems to favor $J^P=2^+$, although a partial-wave analysis leads to ambiguous results (Reyes *et al.*, 1998). The 4π system is analyzed assuming the contribution of a number of isobars, including ρ, $f_2(1270)$, $a_1(1260)$, $a_2(1320)$, and $(\pi\pi)_S$. Both 1^{++} and 2^{++} structures, as well as others, are found throughout this region.

c. Peripheral hadronic reactions

The $f_J(1710)$ has generally been unobserved in peripheral hadronic reactions. A measurement of the reaction $\pi^- p \rightarrow K_S K_S n$ with 22 GeV/c pions, which included a systematic study of the 2^{++} meson spectrum for states decaying to $\pi\pi$ and $K\bar{K}$ (Longacre *et al.*, 1986), found no evidence for the $f_J(1710)$ except in the J/ψ radiative decay data. A measurement of $K^- p \rightarrow K_S K_S \Lambda$ by the LASS Collaboration (Aston *et al.*, 1988a) sees a clear signal for the $f_2'(1525)$ with no evidence for any structure near the $f_J(1710)$.

Interestingly, however, a rather old measurement of $\pi^- p \rightarrow K_S K_S n$ at BNL (Etkin *et al.*, 1982b) and a detailed analysis of the $K_S K_S$ S-wave (Etkin *et al.*, 1982c) reveals two states that are more or less consistent with both the $f_0(1500)$ and $f_0(1710)$. It is reasonable to assume that these states produced in charge exchange with a π beam, also couple to $\pi\pi$. It might also be plausible to argue that they are in fact significant mixtures of the scalar glueball and the $s\bar{s}$ member of the 0^{++} nonet. These states were confirmed in a subsequent experiment at Serpukhov (Bolonkin *et al.*, 1988). In fact, a reanalysis of these data, in combination with the data from both LASS and J/ψ radiative decay (Lindenbaum and Longacre, 1992), shows clearly that $J=0$ and derives branching ratios to $\pi\pi$, $K\bar{K}$, and $\eta\eta$.

The GAMS Collaboration (Alde *et al.*, 1992) observes a state that may or may not be the $f_J(1710)$. Called $X(1740)$, it is observed decaying to $\eta\eta$ in the reaction $\pi^- p \rightarrow \eta\eta N^\star$, for 38 GeV/$c$. The $\eta\eta$ mass distribution shows a significant peak at 1744 ± 15 MeV/c^2 when the recoiling nucleon is accompanied by photons at large angle. That is, the signal is present for $\pi^- p \rightarrow \eta\eta N^\star$ with $N^\star \rightarrow n+\gamma$'s, but not for $\pi^- p \rightarrow \eta\eta n$. The peak is narrower than has been observed for the $f_J(1710)$, with $\Gamma < 80$ MeV/c^2. No structure is observed in the $\pi^0\pi^0$ or $\eta\eta'$ mass spectra in the same experiment.

d. Two-photon collisions

One expects glueballs to be absent in two-photon production. In studies of $\gamma\gamma \rightarrow K\bar{K}$ (Althoff *et al.*, 1985; Behrend *et al.*, 1989c; Albrecht *et al.*, 1990) a clear signal for $f_2'(1525)$ is evident, but only upper limits are put on $\Gamma_{\gamma\gamma}$ for $f_J(1710)$ (see, however, Braccini, 1998). The analysis in this case is difficult, because overlap from the various amplitudes producing $f_2'(1525)$ must be taken into account. This is particularly true for K^+K^- where the broad $a_2(1320)$ also contributes to the sample. A high-statistics measurement of $\gamma\gamma \rightarrow K_S K_S$ would be particularly useful along with a complete partial-wave analysis in the 1400-to-1800-MeV/c^2 region.

2. The $f_J(2220)$

The $f_J(2220)$, also known as $\xi(2220)$ or $\xi(2230)$, is a candidate for the lightest tensor glueball. However, this association is tenuous for a number of reasons. As listed by the Particle Data Group (Caso *et al.*, 1998), its mass and width are 2231 ± 4 MeV/c^2 and 23 ± 8 MeV/c^2, respectively. The mass is close to that expected for the 2^{++} glueball from lattice gauge calculations (Bali *et al.*, 1993; Chen *et al.*, 1994) but the width is very small. The state has been seen mainly in J/ψ radiative decay, with a number of decay channels, but never with a strong statistical significance.

This state was first observed in $J/\psi \rightarrow \gamma K^+K^-$ and $J/\psi \rightarrow \gamma K_S K_S$ by the Mark III Collaboration (Baltrusaitis *et al.*, 1986b), based on a sample of 5.8×10^6 J/ψ decays. In both decays, the $K\bar{K}$ mass distribution rises near 2 GeV/c^2, producing a broad enhancement at high masses. Superimposed on this enhancement is a narrow signal of $\sim 3-4$ standard deviations in *each* channel, consistent with the mass and width of one state, the $f_J(2220)$. The state was unobserved in a number of other two-body channels, and upper limits are quoted. No attempt was made to identify the spin.

The DM2 Collaboration searched through a sample of 8.6×10^6 J/ψ for radiative decays to $\pi^+\pi^-$ (Augustin *et al.*, 1987), K^+K^-, and $K_S K_S$ (Augustin *et al.*, 1988) and did not see the $f_J(2220)$ in any of these three channels. They quote a limit on the product branching ratio, that is, $B[J/\psi \rightarrow \gamma f_J(2220); f_J(2220) \rightarrow K^+K^-]$, which is incompatible with the value determined by Mark III. However, they observe the same broad high-mass enhancement in $K\bar{K}$ and suggest it may represent a state at 2197 ± 17 MeV/c^2 with width ~ 200 MeV/c^2.

Recent measurements of J/ψ radiative decay by BES, also with $\sim 8\times 10^6$ J/ψ events, claim observation of $f_J(2220)$ at the level of several standard deviations in the $\pi^+\pi^-$, K^+K^-, $K_S K_S$, $p\bar{p}$ (Bai *et al.*, 1996a), and $\pi^0\pi^0$ (Bai *et al.*, 1998a) channels. Mass distributions in the region of the $f_J(2220)$ are reproduced from Bai *et al.* (1996a) in Fig. 27. The product branching ratios for these channels are marginally consistent with those determined by Mark III.

Stringent limits have been placed on the two-photon coupling of the $f_J(2220)$ by the CLEO Collaboration in the reactions $\gamma\gamma \rightarrow K_S K_S$ (Godang *et al.*, 1997) and $\gamma\gamma \rightarrow \pi^+\pi^-$ (Alam *et al.*, 1998). Of course, one expects the two-photon width of glueballs to be small. These results determine the "stickiness" (Chanowitz, 1984) of the $f_J(2220)$ to be 100 times as great as that of the $f_2(1270)$. Despite the clear presence of the $f_2'(1525)$, Godang *et al.* (1997) see very few events for $K_S K_S$ masses above 2 GeV/c^2.

Narrow structures have been reported at 2220 MeV/c^2 in peripheral hadron production. GAMS reported (Alde *et al.*, 1986) a small but significant signal decaying to $\eta\eta'$ in $\pi^- p \rightarrow \eta\eta' n$ interactions at 38 GeV/c and at 100 GeV/c. The angular distribution argues strongly that J is greater than or equal to 2. The LASS group (Aston *et al.*, 1988d) report a narrow $J^{PC}=4^{++}$ state decaying to $K\bar{K}$ in both the reactions $K^- p \rightarrow K^+K^-\Lambda$ and $K^- p \rightarrow K_S K_S \Lambda$, at 11 GeV/$c$ beam momentum. Both the mass and width of the GAMS and LASS states are consistent with the $f_J(2220)$ as seen in J/ψ radiative decay. A moments analysis of the LASS results makes it clear that spins greater than $J=2$ are required to describe the data.

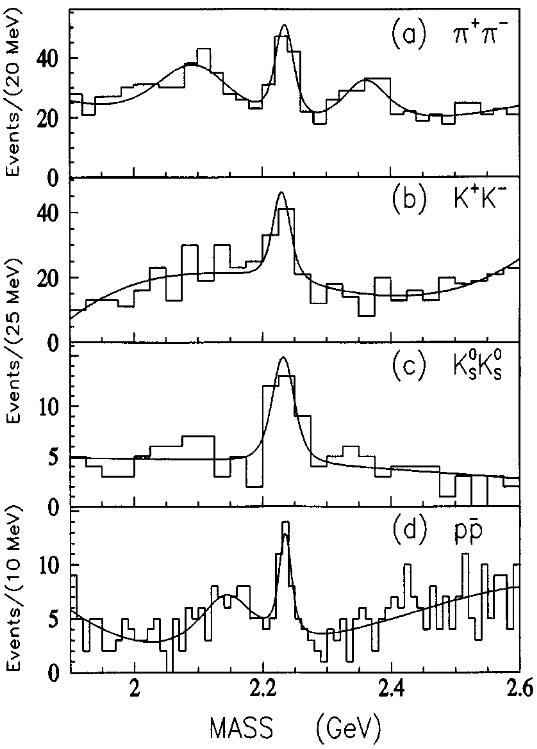

FIG. 27. Figure 2 from Bai *et al.* (1996a), showing various two-particle mass distributions observed in J/ψ radiative decay. Each shows a signal for the $f_J(2220)$. The lines represent fits to a smooth background and one or more Breit-Wigner resonance shapes, convoluted with the appropriate Gaussian resolution function.

The production in high-mass $K\bar{K}$ states in peripheral hadronic reactions prompted a recent study of $s\bar{s}$ quark-model states with $J \geqslant 2$ and $C=P=+$ (Blundell and Godfrey, 1996). These are the $L=3$, (i.e., ${}^3F_2$ or ${}^3F_4$) states, and the quark model does indeed predict rather large widths. Although these states may explain at least some of the structure observed in the 2.2-GeV/c^2 region, it would be difficult to identify them with one state having a width as small as the $f_J(2220)$.

Since the $f_J(2220)$ lies above the $p\bar{p}$ threshold, it is possible to search for it in $p\bar{p}$ annihilation in flight. This is particularly interesting in light of the positive result observed by BES (Bai *et al.*, 1996a) for $J/\psi \rightarrow \gamma f_J(2220)$ followed by $f_J(2220) \rightarrow p\bar{p}$. Several annihilation-in-flight searches have in fact been carried out, including $p\bar{p} \rightarrow \pi^+\pi^-$ (Hasan and Bugg, 1996), $p\bar{p} \rightarrow K^+K^-$ (Bardin *et al.*, 1987; Sculli *et al.*, 1987), $p\bar{p} \rightarrow K_S K_S$ (Evangelista *et al.*, 1997), $p\bar{p} \rightarrow \phi\phi$ (Evangelista *et al.*, 1998), and $p\bar{p} \rightarrow p\bar{p}\pi^+\pi^-$ (Buzzo *et al.*, 1997). No evidence for a narrow state at 2220 MeV/c^2 is seen in any of these experiments. Since the branching ratio $B[J/\psi \rightarrow \gamma f_J(2220)]$ is not known and in principle is unconstrained, it is not possible to make a model-independent consistency check of these data. However, if we combine the JETSET result (Evangelista *et al.*, 1997)

$$B[f_J(2220) \rightarrow p\bar{p}] \times B[f_J(2220) \rightarrow K_S K_S] \leqslant 7.5 \times 10^{-5}$$

(95% C.L.),

with the values from BES (Bai *et al.*, 1996a),

$$B[J/\psi \rightarrow \gamma f_J(2220)] \times B[f_J(2220) \rightarrow K_S K_S] = (2.7 \pm 1.1) \times 10^{-5},$$

$$B[J/\psi \rightarrow \gamma f_J(2220)] \times B[f_J(2220) \rightarrow p\bar{p}] = (1.5 \pm 0.6) \times 10^{-5},$$

then we can infer the lower bound

$$B[J/\psi \rightarrow \gamma f_J(2220)] \geqslant (2.3 \pm 0.6) \times 10^{-3}.$$

This is not only an inordinately large branch for a radiative decay, it also implies that all the branches reported by BES, $B[J/\psi \rightarrow \gamma f_J(2220); f_J(2220) \rightarrow X] \approx 1.5 \times 10^{-4}$, represent only about 10% of the total decay modes of the $f_J(2220)$. We conclude that the evidence for a narrow state $f_J(2220)$ is rather suspect and the branching ratio to $p\bar{p}$ must be checked. Preliminary results of a high-sensitivity search in $p\bar{p} \rightarrow \eta\eta$ and $p\bar{p} \rightarrow \pi^0\pi^0$ have been reported by the Crystal Barrel Collaboration (Seth, 1997) with no evidence of the $f_J(2220)$.

D. $J^{PC}=0^{-+}$ and 1^{++} states in the E region

In principle, studying the $K\bar{K}\pi$ final state is a good way to study mesons that couple to $s\bar{s}$ but that are forbidden to decay to $K\bar{K}$. This would include mesons with $J^P = odd^+$ or $J^P = even^-$. Indeed, experiments that produce $K\bar{K}\pi$ show significant structure in the mass spectra. As an example, Fig. 28(a) shows a histogram of the $K\bar{K}\pi$ mass for the reaction $\pi^- p \rightarrow K^+ K_S \pi^- n$ at 18 GeV/c (Cummings, 1995). Relatively narrow structures are clear at $\sim$1300 MeV/c^2 and $\sim$1400 MeV/c^2. These structures are traditionally referred to as the "D" and "E" regions, respectively. The names D and E have come to mean $J^{PC}=1^{++}$ states within these peaks. In fact, we know that these regions are more complex. The D region is actually well understood, and contains the $f_1(1285)$ and $\eta(1295)$ mesons. However, the structure within the E region is considerably more complicated and remains controversial. For example, although peripheral experiments see the presence of $J^{PC}=1^{++}$, $p\bar{p} \rightarrow E\pi\pi$ experiments (Baillon *et al.*, 1967; Foster *et al.*, 1968; Amsler *et al.*, 1995b) observe mainly $J^{PC}=0^{-+}$. The name given to the $J^{PC}=0^{-+}$ strength within this peak was originally $\iota(1450)$, as observed in J/ψ radiative decay (Scharre *et al.*, 1980; Edwards *et al.*, 1982a). For these reasons, many still refer to this controversy as the E/ι puzzle.

The discovery of the $\iota(1450)$ led to enormous excitement over the possibility that the first glueball may have been found (Scharre *et al.*, 1980; Ishikawa, 1981; Edwards *et al.*, 1982a, Aihara *et al.*, 1986a), but this is no longer the generally accepted viewpoint. For one thing, the controversy over 1^{++} or 0^{-+} quantum numbers makes this interpretation difficult to support. Further-

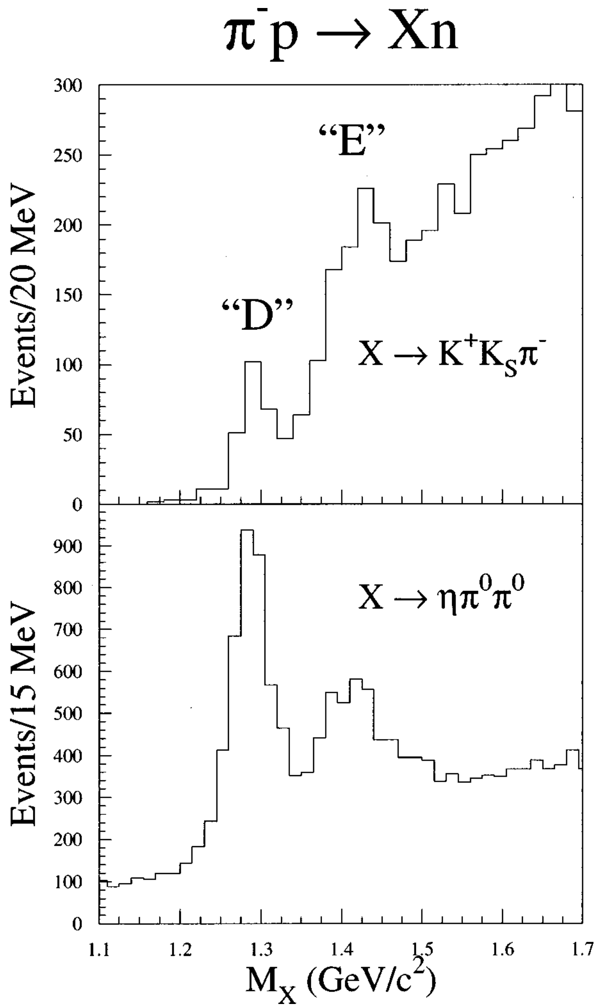

FIG. 28. Invariant-mass distributions for states produced in peripheral production with high-energy π^- beams; (a) the $K^+K_S\pi^-$ distribution at 18 GeV/c (Cummings, 1995); (b) the $\eta\pi^0\pi^0$ result at 100 GeV/c (Alde *et al.*, 1997). In each case, peaks in the the D and E regions are clear.

more, lattice gauge theory would have great difficulty accommodating such a low-mass glueball with these quantum numbers.

A large part of the difficulty is that these structures lie very close to the $K^\star\overline{K}$ threshold,[10] nominally just below 1400 MeV/c^2. In fact, it is clear that the background beneath the D and E peaks in Fig. 28(a) rises very sharply right at this point. In principle, this would favor $J^{PC}=1^{+\pm}$ states, since they would decay to an S-wave $K^\star\overline{K}$ pair. However, the limited phase space leads to very small relative momenta between the K and $\overline{K}$. As a consequence, coupling through the $a_0(980)$, which has a great affinity for $K\overline{K}$, is very likely, leading to S- and P-wave $a_0(980)\pi$ decays, that is to, $J^{PC}=0^{-+}$ and $J^{PC}=1^{++}$ states. It is therefore not surprising that the $K\overline{K}\pi$ system can get very complicated.

A byproduct of nearness to the $K^\star\overline{K}$ threshold, and a possible tool for unraveling this structure, is suggested by the decay $a_0(980)\rightarrow\eta\pi$. That is, one might expect similar structure in $\eta\pi\pi$ final states. Figure 28(b) plots the $\eta\pi\pi$ invariant-mass distribution for the reaction $\pi^-p\rightarrow\eta\pi^0\pi^0n$ at 100 GeV/c (Alde *et al.*, 1997). The same D and E structures are apparent here as well. We note that the E peak in the mass spectrum does not appear very clearly in peripheral hadroproduction of states decaying to $\eta\pi^+\pi^-$ (Fukui *et al.*, 1991; Manak, 1997), because of additional quasi-two-body states such as $\eta\rho$.

$K\overline{K}\pi$ and $\eta\pi\pi$ final states in the E region have been studied not only in peripheral hadroproduction, but also in J/ψ radiative decay, central production, two-photon collisions, and $\overline{p}p$ annihilation. Although there is still significant controversy, one can identify a number of experimental consistencies:

The $J^{PC}=0^{-+}$ strength is resolved into two components. One, decaying to $\eta\pi\pi$ and to $K\overline{K}\pi$ through $a_0(980)\pi$, has mass between 1400 and 1420 MeV/c^2. The other decays to $K^\star\overline{K}$ and has a mass between 1470 and 1480 MeV/c^2. [Note that the Particle Data Group (Caso *et al.*, 1998) tabulates all the $J^{PC}=0^{-+}$ decays under the single heading $\eta(1440)$.]

There is at least one $J^{PC}=1^{++}$ component to the E structure called $f_1(1420)$. This is seen most clearly in single-tagged two-photon collisions, which identify the quantum numbers unambiguously. It is also clearly seen in central production and in $p\overline{p}$ annihilation at rest in gaseous hydrogen.

A second $J^{PC}=1^{++}$ state, the $f_1(1510)$, has been identified. It decays to $K^\star\overline{K}$ and is produced in peripheral hadroproduction with both K^- and π^- beams (see Sec. IV.B.3).

There is no evidence for the $\eta(1440)$ in untagged two-photon collisions. It is considerably more "sticky" than the η or η'. This enhanced stickiness is attributed to a very large J/ψ radiative decay width and may be the strongest evidence for significant gluonic degrees of freedom in these states.

There is also fair evidence for a $J^{PC}=1^{+-}$ state, the $h_1(1380)$, decaying to $K^\star\overline{K}$ and produced in K^-p interactions and in $p\overline{p}$ annihilations at rest (see Sec. IV.B.3).

Even putting aside various experimental inconsistencies (which we detail below), there is already serious difficulty accommodating these results. In the quark model, one might expect two states in the E region: the $s\overline{s}$ partners to the $f_1(1285)$ and $\eta(1295)$. The evidence is, however, that there are two states of each J^{PC}. Accommodating the overpopulation with states having gluonic degrees of freedom is problematic, since this is at odds with nearly all models. A variety of models have been suggested (for example, see Longacre, 1990), but it is difficult to find clear, testable predictions of these models outside of the E region. It may be that the proximity of this region to the $K^\star\overline{K}$ threshold is at the heart of the difficulties, both in the experiments and in their interpretation.

[10]We use the notation $K^\star\overline{K}$ to denote the combination of $K^\star(892)\overline{K}+$complex conjugate.

We discuss separately in greater detail the individual evidence and a possible interpretation for the $J^{PC}=1^{++}$ and 0^{-+} states.

1. $J^{PC}=1^{++}$

The $f_1(1420)$ and the $f_1(1510)$ are well separated in mass and well resolved in the different experiments although there are no experiments (or production reactions) in which *both* states are observed. The proximity to the $K^\star\bar{K}$ threshold makes it plausible that a $K^\star\bar{K}$ $L=0$ molecular bound state is mixing with the $s\bar{s}$ state, but so far there is no good explanation of why the different production mechanisms so strongly favor one component over the other.

In experiments performed since around 1980, the lower mass is favored by all measurements *other than* peripheral hadron production (Caso *et al.*, 1998), including central production, two-photon collisions, $p\bar{p}$ annihilation, and J/ψ radiative decay. The higher mass comes only from production in $\pi^- p\to Xn$ (Birman *et al.*, 1988; Cummings, 1995) and $K^- p\to Xn$ (Aston *et al.*, 1988c; King, 1991) reactions. [A measurement of $\gamma\gamma^\star\to\pi^+\pi^-\pi^0\pi^0$ (Bauer, 1993) suggests a mass near 1510 MeV/c^2, but the final state is not fully reconstructed.]

A natural way to study the mainly $s\bar{s}$ partner of the $f_1(1285)$ would be hadronic peripheral production in $K^- p\to K\bar{K}\pi\Lambda$. That is, one would consider hypercharge exchange leading to $K^\star\bar{K}$ final states, so the intermediate state couples to $s\bar{s}$ on both the input and output channels. This experiment was in fact carried out thoroughly by the LASS Collaboration (Aston *et al.*, 1988c) using the $K^- p\to K_S K^\pm\pi^\mp\Lambda$ reaction at 11 GeV/c. A clean sample of 3900 events was obtained with the K_S and Λ both clearly identified with decay vertices separated from the primary interaction vertex. The $K^\star\bar{K}$ mass spectrum clearly shows the D region heavily suppressed relative to the E, suggesting that the E is dominated by $s\bar{s}$. A partial-wave analysis of the E, shown in Fig. 19, revealed a predominance of $J^P=1^+$, and also showed that this cannot be the result of a single resonance. When the final state was symmetrized on the basis of charge conjugation, the $J^P=1^+$ strength was clearly resolved into a $J^{PC}=1^{++}$ state at 1530±10 MeV/c^2, and a $J^{PC}=1^{+-}$ state at 1380±20 MeV/c^2. The experimenters explicitly point out that these are good candidates for the mainly $s\bar{s}$ members of their respective nonets. Furthermore, these results are confirmed by a BNL/MPS experiment (King, 1991), which studied the reaction $K^- p\to K_S K^+\pi^-(\Lambda,\Sigma^0)$ at 8 GeV/c.

Single-tagged two-photon production (Sec. III.E.2) of $K\bar{K}\pi$ final states (Aihara *et al.*, 1986b; Gidal *et al.*, 1987b; Behrend *et al.*, 1989b; Hill *et al.*, 1989) all show clear evidence of the $J^{PC}=1^{++}$ state, although the mass is generally more consistent with 1420 than with 1510 MeV/c^2, with a $\gamma\gamma^\star$ width (times $K\bar{K}\pi$ branching ratio) of 1.7±0.4 keV (Caso *et al.*, 1998). On the other hand, single-tagged production of $\eta\pi\pi$ (Gidal *et al.*, 1987a; Aihara, 1988b) shows a clear signal for $f_1(1285)$ and a $\gamma\gamma^\star$ width (times $\eta\pi\pi$ branching ratio) of 1.4 ±0.4 keV (Caso *et al.*, 1998). This appears to be consistent with the $f_1(1285)$ and $f_1(1420)$'s being $n\bar{n}$ and $s\bar{s}$ partners, but the comparison is probably misleading. It is quite difficult to make good comparisons of the $\gamma\gamma^\star$ width alone, because of the proximity to threshold and to the inherent difficulties in predicting the width at all. In fact, two-photon couplings to molecular states are not likely to be profoundly different (Barnes, 1985b).

Radiative J/ψ decay into $K\bar{K}\pi$ shows considerable strength in the E region, although it is dominated by $J^{PC}=0^{-+}$ (see the following section). However, a partial-wave analysis (Bai *et al.*, 1990; Augustin *et al.*, 1992) shows a signal for $J^{PC}=1^{++}$ consistent with a mass closer to 1420 MeV/c^2, and $B[J/\psi\to f_1(1420)\gamma]\times B[f_1(1420)\to K\bar{K}\pi]=8.3\pm1.5\times10^{-4}$ (Caso *et al.*, 1998). A small $J^{PC}=1^{++}$ contribution decaying to $K\bar{K}\pi$ has been seen in $p\bar{p}$ annihilation in gaseous hydrogen by the OBELIX Collaboration (Bertin *et al.*, 1997a) for the reaction $p\bar{p}\to K^\pm K^0\pi^\mp\pi^+\pi^-$ (where the K^0 is not seen), with mass 1425±8 MeV/c^2. There is also recent evidence from the Crystal Barrel Collaboration (Abele *et al.*, 1997b) for the $J^{PC}=1^{+-}$ state observed by LASS (Aston *et al.*, 1988c). Central production of $K\bar{K}\pi$ (Armstrong *et al.*, 1989b, 1992; Barberis *et al.*, 1997c) and of $\eta\pi\pi$ (Armstrong *et al.*, 1991b) has been performed many times, and a clear peak appears at 1420 MeV/c^2, consistent with $J^{PC}=1^{++}$. The behavior of this signal with transverse momentum suggests that it, like the $f_1(1285)$, is a conventional $q\bar{q}$ meson.

These results have led some to question the existence of two separate states (Close and Kirk, 1997). It is clearly of high importance to observe these two states simultaneously. There is a chance that the next generation of $\gamma\gamma^\star$ measurements will be able to see both states and resolve them separately.

2. $J^{PC}=0^{-+}$

Once again, we might normally expect one $J^{PC}=0^{-+}$ state in the E region, namely, the $s\bar{s}$ partner of the $\eta(1295)$. These two states would be taken to be linear combinations of the radial excitations of the ground states η and η'. There again seems to be clear evidence of two $J^{PC}=0^{-+}$ states in the E, but these are seen simultaneously in the same experiments, unlike the $f_1(1420)$ and $f_1(1510)$. There are other oddities about the $J^{PC}=0^{-+}$ as well.

Table V shows results from several experiments that see pseudoscalar resonances in the E region, decaying to $K\bar{K}\pi$. Nearly all measurements of this final state, including peripheral production with pion beams, $p\bar{p}$ annihilation, and J/ψ radiative decay, distinguish two states. [One notable exception is a recent BES measurement (Bai *et al.*, 1998b) of $J/\psi\to\gamma K^+K^-\pi^0$.] The masses differ by about 50 MeV/c^2, with the heavier decaying to $K^\star\bar{K}$, except for the experiment of Augustin

TABLE V. Masses, in MeV/c^2, of $J^{PC}=0^{-+}$ $K\overline{K}\pi$ resonances in the E region.

Reaction	Reference	Quasi-two-body mode $a_0\pi$	$K^\star\overline{K}$
$p\overline{p}\rightarrow K^{\pm}K^0\pi^{\mp}\pi^+\pi^-$	Bertin *et al.*, 1997a	1407±5	1464±10
$J/\psi\rightarrow\gamma K_S K^{\pm}\pi^{\mp}$	Bai *et al.*, 1990	1416±10	1490±20
$J/\psi\rightarrow\gamma K\overline{K}\pi$	Augustin *et al.*, 1992	1459±5	1421±14
$\pi^-p\rightarrow K_S K_S\pi^0$ at 21 GeV/c	Rath *et al.*, 1989	1413±5	1475±4
$\pi^-p\rightarrow K^+K_S\pi^-$ at 18 GeV/c	Cummings, 1995	1412±2	1475±6
Particle Data Group	Caso *et al.*, 1998	1418.7±1.2	1473±4

et al. (1992). Furthermore, the mass of the lower state agrees very well with measurements of states decaying to $\eta\pi\pi$, again in a number of different types of experiments. Most recently, measurements of the reactions $\pi^-p\rightarrow\eta\pi^0\pi^0$ at 100 GeV/c (Alde *et al.*, 1997) and $\pi^-p\rightarrow\eta\pi^+\pi^-$ at 18 GeV/c (Manak, 1997) show that most of the $J^{PC}=0^{-+}$ $\eta\pi\pi$ signal is concentrated in $\eta(\pi\pi)_S$ instead of $a_0(980)\pi$. This is not necessarily consistent, however, with $J^{PC}=0^{-+}$ $\eta\pi\pi$ signals in $\overline{p}p$ annihilation (Amsler *et al.*, 1995b).

Clearly, this leads us to guess that one of these states is the $s\overline{s}$ partner of the $\eta(1295)$, and the other is some manifestation of non-$q\overline{q}$ degrees of freedom. That would imply, however, that one of the states (presumably the one that decays to $K^\star\overline{K}$) should appear in untagged $\gamma\gamma$ collisions, while the other might show some anomalous dependence on transverse momentum in central pp and πp collisions. In fact, there is *no* evidence that *either* state is produced in either of these reactions.

The most stringent limits on two-photon production of an $s\overline{s}$ pseudoscalar meson in the E region were obtained by Behrend *et al.* (1989b) in the reaction $\gamma\gamma\rightarrow K_SK^{\pm}\pi^{\mp}$. They find $\Gamma_{\gamma\gamma}[\eta(1440)]\times B[\eta(1440)\rightarrow K\overline{K}\pi]<1.2$ keV at 95% C.L. They further determine that this implies that the $\eta(1440)$ is at least 20 times as "sticky" as the η'. Unless there is a fortuitous cancellation due to quark mixing angles, this would argue that the $\eta(1440)$ (or both pseudoscalars in this region) have large glue content, and it would leave the $s\overline{s}$ partner of the $\eta(1295)$ unidentified. There may be more to this than meets the eye, however. The $\eta(1295)$ has also been unidentified in two-photon collisions (Caso *et al.*, 1998).

There is some evidence of $J^{PC}=0^{-+}$ production in $K_SK^{\pm}\pi^{\mp}$ final states in pp and π^+p central collisions (Armstrong *et al.*, 1992), but the signal in the E region is dominated by 1^{++}. Also, the E recoiling against $\pi\pi$ in $\overline{p}p$ annihilation appears to be dominated by 0^{-+} (Amsler *et al.*, 1995b).

It is obviously very difficult to draw a clear, consistent picture from the $J^{PC}=0^{-+}$ results in the E region. It may simply be that its proximity to the $K^\star\overline{K}$ threshold brings in more complicated mechanisms than can be treated with the formalisms presently at our disposal.

E. Other puzzles

Thus far we have discussed the states that have received the most attention in recent years. In addition to these, numerous other extraneous states have been reported, which have received much less attention of late, primarily because there has been no new information on them. In this subsection, for completeness, we briefly discuss some additional examples.

1. Extra $J^{PC}=2^{++}$ states

The ground-state 2^{++} nonet, consisting of the $a_2(1320)$, $f_2(1270)$, $f_2'(1525)$, and $K^*(1430)$, has been complete for some time. However, several additional isoscalar 2^{++} states that are inconsistent with quark-model predictions have been observed. These include the three g_T states in $\phi\phi$ at 2011, 2297, and 2339 MeV, the $f_2(1565)$ seen in $\overline{p}N$ annihilation, and the $f_2(1430)$ and $f_2(1480)$ observed in $\pi\pi$ and $K\overline{K}$ spectra between the $f_2(1270)$ and the $f_2'(1525)$.

a. OZI suppression and states in $\pi^-p\rightarrow\phi\phi n$

In the high-mass region, three states, sometimes known as g_T, have been observed in the OZI-suppressed reaction $\pi^-p\rightarrow\phi\phi n$ at 22 GeV/c by a BNL group (Etkin *et al.*, 1978a, 1978b, 1982a, 1985, 1988) using the Multi Particle Spectrometer (MPS) facility at the AGS. The three are distinguishable by their decay couplings to different $\phi\phi$ partial waves. The two higher-mass states have also been seen as peaks in the two 2^{++} components of inclusive production from π^- Be interaction at 85 GeV/c in the WA67 experiment at CERN (Booth *et al.*, 1986).

The $\phi\phi$ mass spectrum observed by the BNL group shows a broad enhancement from threshold to 2.4 GeV, while the experimental acceptance remains flat up to 2.6 GeV. A partial-wave analysis of the bump reveals that it consists of three distinct 2^{++} states, 2^{++} $f_2(2010)/g_T$, $f_2(2300)/g_T'$, and $f_2(2340)/g_T''$ (Etkin *et al.*, 1985; Longacre *et al.*, 1986). Using the notation L_S where L is the orbital angular momentum and S is the total intrinsic spin for $\phi\phi$, the states are g_T with $M=2011\pm70$ MeV/c^2 and $\Gamma=202\pm65$ MeV/c^2 (about 98% S_2); g_T' with $M=2297\pm28$ MeV/c^2 and $\Gamma=149\pm41$ MeV/c^2 (about 25% D_2 and 69% D_0); and g_T'' with $M=2339$

$\pm 55\,\mathrm{MeV}/c^2$ and $\Gamma = 319\pm\ \mathrm{MeV}/c^2$ (about 37% S_2 and 59% D_0). From a study of the production characteristics, it was concluded that the g_T states are produced by one-pion-exchange processes. If so, the g_T states should also couple to $\pi\pi$ channels but they are difficult to observe due to large backround in these channels. As they are produced in the OZI-forbidden channel (Landberg *et al.*, 1996), they are thought to be candidates for tensor glueballs, although this interpretation is controversial (Lindenbaum and Lipkin, 1984).

The WA67 group at the CERN Ω-Spectrometer (Booth, 1986) studied inclusive $\phi\phi$ production from π^-Be interactions at $85\,\mathrm{GeV}/c$. They saw general enhancement at the $\phi\phi$ threshold followed by a second peak at 2.4 GeV. Assuming that they saw the second and third g_T states, they fitted their mass spectrum with two Breit-Wigner forms, one with 50-50% S and D waves and the other 100% D wave over a smooth background. The resulting masses and widths are 2231 ± 10 MeV/c^2 and 133 ± 50 MeV/c^2 for the second $f_2(2300)/g_T'$ and 2392 ± 10 MeV/c^2 and 198 ± 50 MeV/c^2 for the third $f_2(2340)/g_T''$, respectively. They also carried out a joint moment analysis and found that the $\phi\phi$ system up to 2.5 GeV/c^2 was mainly 2^{++} (Booth *et al.*, 1986; Armstrong *et al.*, 1989b) although the statistics are limited.

One might expect the g_T's to couple to the $\rho\rho$ and $\omega\omega$ channels as well, if they were indeed glueballs. The GAMS group (Alde *et al.*, 1988b) observed two 2^{++} resonances in $\pi^- p \to \omega\omega n$ at 38 GeV/c. The state at 1956 ± 20 MeV/c^2 with width 220 ± 60 MeV/c^2 is not inconsistent with the lightest g_T.

The DM2 Collaboration (Bisello, 1986) carried out an analysis of the $\phi\phi$ system produced in J/ψ radiative decays. Although 2^{++} was found to be the main wave, no threshold enhancement in the $\phi\phi$ system was observed, in contrast to the hadronic production. However, they found a narrow peak at around 2.2 GeV with a preferred spin parity of 0^-. The Mark III Collaboration also studied their $\phi\phi$ spectrum in J/ψ radiative decays, with possible structures in the 2.1–2.4-GeV mass region (Mallik, 1986; Blaylock, 1987; Toki, 1987). However, no spin-parity has been given for the $\phi\phi$ structures and it is not clear if they are to be associated with their 0^{-+} structures near threshold in $\rho\rho$ and $\omega\omega$, which are also seen in J/ψ radiative decays, or with the BNL g_T states.

As these states are above $p\bar{p}$ threshold, they can in principle be observed in annihilation-in-flight reactions. The JETSET Collaboration (Evangelista *et al.*, 1998) studied the reaction $\bar{p}p \to \phi\phi$ with antiproton beams having between 1.1 and $2.0\,\mathrm{GeV}/c$ momentum. This experiment was sensitive to intermediate states formed in the annihilation channel with masses between 2.1 and $2.4\,\mathrm{GeV}/c^2$. No evidence was found for the $f_2(2300)/g_T'$ or $f_2(2340)/g_T''$.

Finally, the $\phi\phi$ system has been studied in pp central production by the WA102 Collaboration (Barberis *et al.*, 1998). Some weak structure was observed in the $\phi\phi$ mass distribution, but the statistics were poor. The angular distribution favored $J^{PC} = 2^{++}$.

b. The $f_2(1565)$

This state has a long history going back to the 1960s (Bettini , 1966; Conforto *et al.*, 1967). In 1990, ASTERIX at LEAR presented evidence for the production of a resonance $f_2(1565)$ in P-wave $p\bar{p} \to \pi^0\pi^+\pi^-$ annihilation in hydrogen gas (May *et al.*, 1989, 1990a). A state with $M = 1565\pm 10\,\mathrm{MeV}/c^2$ and $\Gamma = 170 \pm 20\,\mathrm{MeV}/c^2$ was observed decaying to $\pi^+\pi^-$, recoiling against the π^0. No enhancement was visible in the $\pi^\pm\pi^0$ invariant mass, indicating that it was an $I = 0$ resonance. A Dalitz-plot analysis showed clear evidence for $J^{PC} = 2^{++}$. This resonance could not be identified with the $f_2'(1525)$ meson, which decays mostly to $K\bar{K}$. Otherwise it would be produced strongly in the final-state $K\bar{K}\pi$ where it has not been observed (Conforto *et al.*, 1967). In a separate analysis (May *et al.*, 1990b), selecting initial $p\bar{p}$ S states, no indication of a resonance at 1.5 GeV was observed.

The Crystal Barrel experiment at LEAR subsequently studied all neutral events from $\bar{p}p$ annihilation (Aker *et al.*, 1991), which gave information on the $3\pi^0$, $\eta\eta\pi^0$, and $\eta\eta'\pi^0$ channels. The $f_2(1270)$ and $f_2(1565)$ resonances were clearly visible in the $\pi^0\pi^0$ invariant-mass projections. As more data was acquired and the analysis matured, however, it became clear that the prominent feature in $\pi^0\pi^0$ near 1500 MeV/c^2 was in fact the $f_0(1500)$. Still, the analysis required the presence of some 2^{++} strength decaying to $\pi^0\pi^0$ in the same region (Amsler, 1998). Recently, the OBELIX Collaboration provided new evidence for the $f_2(1565)$ in $\bar{n}p \to \pi^+\pi^-\pi^+$ (Bertin *et al.*, 1998a).

c. The $f_2(1430)$ and $f_2(1480)$

Evidence for these two states was found in the data on the double-Pomeron-exchange reaction in experiment R807 at CERN ISR (Cecil, 1984; Akesson *et al.*, 1986). The reaction involved exclusive $\pi^+\pi^-$ production in $pp \to p_f(\pi^+\pi^-)p_s$ at $\sqrt{s} = 63$ GeV. The recoil protons were detected with $-t$ less than $0.03\,(\mathrm{GeV}/c)^2$, thus ensuring nearly pure Pomeron exchanges at both vertices. The resulting $\pi\pi$ spectrum exhibited a set of remarkable bump-dip structures near 1.0, 1.5, and 2.4 GeV, respectively (Akesson *et al.*, 1986). Another striking feature was that the $\rho(770)$ and the $f_2(1270)$ were not seen in the data. One might expect that the $I = 1$ $\rho(770)$ should not have been seen in a Pomeron-Pomeron interaction; however, the apparent absence of the $f_2(1270)$ is noteworthy. The second dropoff near 1.4 GeV was partly due to the $f_0(1370)$ and a D-wave structure which was attributed to the $f_2(1480)$. The data in fact favor a 2^{++} structure above the $f_2(1270)$ with mass and width of 1480 ± 50 and 150 ± 40 MeV, respectively. A full understanding of the nature of the $f_2(1480)$ will probably also require an explanation for the absence of the $f_2(1270)$ in the R807 data. An explanation may follow from Close and Kirk's (1997) glueball filter.

2. Structure in $\gamma\gamma \rightarrow VV$

Structures in $\gamma\gamma \rightarrow VV'$ have generated considerable interest. A recent review of results from the ARGUS experiment (Albrecht *et al.*, 1996) includes a thorough discussion of this reaction.

Interest in this subject began with the original observation by the TASSO Collaboration (Brandelik *et al.*, 1980) of structures in the $\gamma\gamma \rightarrow \rho^0\rho^0$ cross section. This was subsequently confirmed by other experiments (Burke *et al.*, 1981; Althoff *et al.*, 1982; Behrend *et al.*, 1984; Aihara, 1988a), in which one would normally expect the rate to be suppressed in the threshold region due to the reduced phase space. Results are now available for numerous final states, including $\rho^0\rho^0$, $\rho^+\rho^-$, $\omega\omega$, $\omega\rho^0$, $K^{0*}\bar{K}^{0*}$, $K^{+*}\bar{K}^{-*}$ $\rho^0\phi$, $\omega\phi$, and $\phi\phi$. The cross sections for the different final states vary in their relative size and in the energy at which they peak.

The large difference in cross section between the $\rho^0\rho^0$ and $\rho^+\rho^-$ channels, $\sigma(\gamma\gamma \rightarrow \rho^0\rho^0) \simeq 4\sigma(\gamma\gamma \rightarrow \rho^+\rho^-)$ (Albrecht *et al.*, 1996; Behrend *et al.*, 1989a) rules out a simple s-channel resonance explanation for which one expects the decay of a conventional resonance into $\rho^+\rho^-$ to occur twice as often as decay into $\rho^0\rho^0$. The two ρ^0 mesons can only be in a state with $I=0$, $I=2$, or a mixture of the two. The large ratio of the $\rho^0\rho^0$ to $\rho^+\rho^-$ cross section cannot be accounted for by a pure $I=0$ or a pure $I=2$ state with the same spin-parity quantum numbers, but must come from interference between the two. The interference is observed to be constructive in $\gamma\gamma \rightarrow \rho^0\rho^0$ and destructive in $\gamma\gamma \rightarrow \rho^+\rho^-$. A demonstration that the $\rho\rho$ cross section is predominantly resonant, together with this isospin argument, would imply the existence of exotic $I=2$ states, possibly $q\bar{q}q\bar{q}$ resonances (Achasov *et al.*, 1982; Li and Liu, 1983).

A number of models have been invoked to explain the structure in the original $\gamma\gamma \rightarrow \rho^0\rho^0$ and have made predictions for other channels: $q\bar{q}q\bar{q}$ exotica were first suggested by Jaffe (1977a, 1977b, 1978) and were explored as an explanation of the structure in $\gamma\gamma \rightarrow VV'$ by Achasov *et al.* (1982) and by Li and Liu (1983). The vector-dominance model with factorization in the t channel attempts to identify specific t-channel exchanges and extract them from photoproduction data (Alexander *et al.*, 1982). Perturbative QCD with Coulombic rescattering corrections was examined by Brodsky *et al.* (1987). In the one-meson-exchange model (Achasov and Shestakov, 1988; Törnqvist, 1991) the structure in $\gamma\gamma \rightarrow \rho\rho$ is explained by bound states formed by one-pion-exchange potentials analogous to those used in nuclear physics. These models do reasonably well for the process for which they were constructed but for the most part they fail to explain subsequent VV' data.

3. The $C(1480)$

The Lepton-F Collaboration at Serpukhov (Bityukov *et al.*, 1987) examined the reaction $\pi^- p \rightarrow \phi\pi^0 n$ with a 32-GeV/c beam. A very strong peak in the cross section was observed near 1500 MeV/c^2 in $\phi\pi$ mass. The production cross section was substantial. Interpreted as a resonance, called $C(1480)$, the peak had mass 1480 ± 40 MeV/c^2, width 130 ± 60 MeV/c^2, and $\sigma(\pi^- p \rightarrow Cn) \times B(C \rightarrow \phi\pi^0) = 40 \pm 15$ nb. The existence of such a narrow, isovector state is clearly peculiar given this decay mode, and a number of interpretations have been offered (Kubarovski *et al.*, 1988; Kopeliovich and Predazzi, 1995).

This state has not been observed in other reactions, including pp central production (Armstrong *et al.*, 1992) and $p\bar{p}$ annihilation at rest (Reifenröther *et al.*, 1991). Recent data taken by the E852 collaboration including K identification will check this reaction with pion beams, however, and photoproduction experiments are planned (Dzierba, 1994).

F. Missing states

Clearly one can only discuss "overpopulation" if all the expected $q\bar{q}$ states have in fact been identified. This is not the case for some multiplets. Particularly glaring examples are the missing isoscalar and isovector $J^{PC} = 2^{--}$ (i.e., ω_2 and ρ_2) states. These would be the $^3D_2$ partners of the relatively well established $^3D_3$ [$\rho_3(1690)$ and $\omega_3(1670)$] and $^3D_1$ [$\rho(1700)$ and $\omega(1600)$] $q\bar{q}$ combinations. The strange $J^P=2^-$ members appear to be the two K_2 states near 1800 MeV/c^2 (Aston *et al.*, 1993), i.e., $^3D_2$ and $^1D_2$, although they need to be confirmed (Caso *et al.*, 1998).

It is important to fill in both the orbitally and radially excited multiplets, although as we go higher in mass the states become broader as well as more numerous. As an example of how a search for the missing states would proceed, we examine the missing states of the $L=2$ meson multiplet. Quark-model predictions for these states are listed in Table VI.

Consider the $\eta_2(1^1D_2)$. There is some evidence that this state has in fact been observed (Sec. IV.B.5). We expect it to be almost degenerate in mass with its nonstrange isovector partner, the $\pi_2(1670)$. From Table VI we see that it is expected to be rather broad and it decays predominantly through the $a_2(1320)$ isobar, which in turn decays to $\rho\pi$. The 4π final state is complicated to reconstruct. Since $a_2 \rightarrow \eta\pi$, other final states should be checked. The $\rho_2(1^3D_2)$ will also decay dominantly to a 4π final state. The ω_2 decays to the simpler $\rho\pi$ final state with a moderate width, but since it has a mass similar to that of $\pi_2(1680)$, which also decays to $\rho\pi$, it is possible that it is masked by the π_2.

Similar situations exist for the other $J^P=2^-$ mesons. Clearly, it is important to search thoroughly through data that are already in hand in order to try and identify the missing states. However, it is at least equally important to understand phenomenologically why these states are produced less copiously than their $J^P=1^-$ and $J^P=3^-$ counterparts, if that is indeed the case.

One can perform a similar analysis of other multiplets. However, as we go higher in mass there are more channels available for decay, so that the meson widths become wider and wider (Barnes *et al.*, 1997). In general, given how complicated the meson spectrum is, it appears

TABLE VI. Quark-model predictions for the properties of the missing $L=2$ mesons. The masses and widths are given in MeV.

Meson state	Property	Prediction
$\eta_2(1^1D_2)$	mass	1680
	width	$\sim$400
	$BR(\eta_2\rightarrow a_2\pi)$	$\sim$70%
	$BR(\eta_2\rightarrow\rho\rho)$	$\sim$10%
	$BR(\eta_2\rightarrow K^*\bar{K}+\text{c.c.})$	$\sim$10%
$\eta_2'(1^1D_2)$	mass	1890
	width	$\sim$150
	$BR(\eta_2'\rightarrow K^*\bar{K}+\text{c.c.})$	$\sim$100%
$\omega_1(1^3D_1)$	mass	1660
	width	$\sim$600
	$BR(\omega_1\rightarrow B\pi)$	$\sim$70%
	$BR(\omega_1\rightarrow\rho\pi)$	$\sim$15%
$\rho_2(1^3D_2)$	mass	1700
	width	$\sim$500
	$BR(\rho_2\rightarrow[a_2\pi]_S)$	$\sim$55%
	$BR(\rho_2\rightarrow\omega\pi)$	$\sim$12%
	$BR(\rho_2\rightarrow\rho\rho)$	$\sim$12%
$\omega_2(1^3D_2)$	mass	1700
	width	$\sim$250
	$BR(\omega_2\rightarrow\rho\pi)$	$\sim$60%
	$BR(\omega_2\rightarrow K^*\bar{K})$	$\sim$20%
$\phi_2(1^3D_2)$	mass	1910
	width	$\sim$250
	$BR(\phi_2\rightarrow K^*\bar{K}+\text{c.c.})$	$\sim$55%
	$BR(\phi_2\rightarrow\phi\eta)$	$\sim$25%

that a good starting place would be the strange mesons. The reason for this is that in the strange-meson sector we do not have the additional problem of deciding whether new states are glueballs or conventional mesons, and we do not have the complication of mixing between isoscalar states due to gluon annihilation. Following this, a detailed survey of ϕ states would be useful, since they form a bridge between the heavy quarkonia ($c\bar{c}$ and $b\bar{b}$) and the light-quark mesons and would help us understand the nature of the confinement potential.

These of course are guidelines for the next generation of experiments, which we now discuss.

VI. FUTURE DIRECTIONS

To solve the remaining puzzles will require new data with significantly higher statistics. It is also important that the data come from different processes and channels to produce hadronic states with different quantum numbers and production mechanisms.

Among the highest-priority goals of hadron spectroscopy is to establish the existence, and to study the properties, of gluonic degrees of freedom in the hadron spectrum. New evidence discussed in earlier sections just begins to scratch the surface of this field. There is surely much new physics to be learned from a new generation of experiments and theoretical calculations and modeling. Another important step is to find some of the missing $q\bar{q}$ states, including both the orbitally and the radially excited multiplets.

It is unlikely that all this can be done simply by bump hunting. Rather, we shall need experiments with unprecedented statistics and uniform acceptance so that high-quality partial-wave analyses can filter by J^{PC}. A useful guide to the expected properties of excited quarkonia has been produced by Barnes, Close, Page, and Swanson (1997). As an example, in Sec. V.F, we discussed the quark-model predictions for the properties of the missing $L=2$ mesons.

One impediment to progress in this field is the lack of a consistent treatment of resonance phenomenology and its parametrization. A variety of different approaches to partial-wave analysis have been used, and each of these incorporates different ways of expressing the mass and width of a relativistic Breit-Wigner resonance. We emphasize the need for a more global approach to analyzing these data so that results from different experiments can be more carefully compared.

We point out that a number of results have not yet appeared in the journals. A good source for these are the proceedings of the Hadron '97 conference (Chung and Willutzki, 1998). In particular, we refer the reader to results on light-quark spectroscopy from Fermilab in central pp collisions from E690 (Berisso *et al.*, 1998) and in high-energy photoproduction from E687 (LeBrun, 1998).

A number of complementary new facilities and experiments are on the horizon (Seth, 1998). A partial list of new and planned facilities is given in Table VII. These are described in some detail in the following sections.

A. DAΦNE at Frascati

The DAΦNE ϕ factory (Zallo, 1992), a high-luminosity e^+e^- collider operating at $\sqrt{s}\sim 1$ GeV, is nearing completion at INFN-Frascati. The main goal is to study CP violation, but it will also examine the nature of the $a_0(980)$ and $f_0(980)$ scalar mesons through the radiative transitions $\phi\rightarrow\{a_0,f_0\}\gamma$. (See Achasov Gubin, and Shevchenko, 1997).

ϕ factories also offer the possibility of studying low-mass $\pi\pi$ production via two-photon production. The combination of relatively high luminosity with detectors optimized for detecting low momenta and energies should allow very detailed measurements of both charged and neutral modes to be made from threshold up to nearly 1 GeV. This process can shed additional light on low-mass scalar resonances.

B. B factories at CESR, SLAC, and KEK

Although the primary motivation for a B factory (Goldberg and Stone, 1989; Bauer, 1990, 1992) is to

TABLE VII. Some future facilities for studying quark gluon spectroscopy.

Facility	Beams	Energy	Principle reactions	Approximate date
DAΦNE	e^+e^-	~1 GeV	ϕ decays	1998
CLEO-III	e^+e^-	10 GeV	B decays; $\gamma\gamma\to X$	1998
BaBAR	e^+e^-	10 GeV	B decays; $\gamma\gamma\to X$	1999
BELLE	e^+e^-	10 GeV	B decays; $\gamma\gamma\to X$	1999
COMPASS	p	400 GeV	central production	1999
RHIC	p, nuclei	~200 GeV/nucleon	central production	1999
JHF	p	50 GeV	$\{\pi,K,p\}\to X$	~2004
CEBAF	e^-,γ	12 GeV	$\gamma p\to Xp$	~2004
BEPC	e^+e^-	3–4 GeV	$J/\psi\to\gamma X$	~2004

study CP violation in the B-meson system, light-meson spectroscopy will also be addressed in several ways. Final states produced in the strong *and* weak decays of b and c quarks will elucidate the structure of the light-meson daughters. Further, these high-luminosity facilities will study two-photon processes with very high statistics. Lastly, exclusive radiative processes $\Upsilon\to\gamma X$ will provide information complementary to J/ψ radiative decay.

Three such facilities in the final stages of construction are the CLEO-III upgrade at CESR, the new BaBAR detector at the SLAC asymmetric collider, and the new BELLE detector at KEK. CLEO-III is an upgraded detector to go along with a luminosity boost in the CESR storage ring, a classic and well-understood symmetric e^+e^- collider facility. The SLAC B factory is an asymmetric system (designed to boost B mesons in the lab frame) that has some implications for detection of lower-mass systems, such as those produced in two-photon annihilation. The KEK facility (Kurokawa, 1997) is similar to the SLAC experiment.

We also note that the LEP Collaborations have made contributions to the subject of light-meson spectroscopy and we should expect this to continue for at least the next several years.

C. COMPASS at CERN

CERN is about to commission a new experiment on fixed targets to be operated in the years prior to turn-on of LHC. The Common Muon and Proton Apparatus for Structure and Spectroscopy (COMPASS) is a large magnetic system designed to detect multiparticle final states using either muon or proton beams (von Harrach, 1998). The muon-beam program is mainly directed at measurements of spin structure functions, but the proton beams will be used for a number of hadronic production experiments, in particular, central production.

Operating such a multipurpose apparatus has specific challenges, mainly trying to optimize sensitivities to the different situations. However, using a common apparatus with a large directed collaboration has led to a state-of-the-art high-rate detector with excellent momentum and particle identification capabilities. First runs with hadron beams are expected in 1999. Central production measurements, including a RICH detector for particle identification, are presently planned for 2002.

D. RHIC at BNL

The Z^4 dependence of two-photon production from charged particles implies an enormous cross section in heavy-ion interactions. However, the luminosities at heavy-ion colliders are substantially less than at the e^+e^- factories so that in the end these two factors tend to balance out, leaving the rates for two-photon physics generally lower than at the present e^+e^- facilities. Still, reasonable event totals may be expected when integrating over the long running periods, with a suitably triggered detector. The real strength of heavy-ion colliders may be their ability to compare the gluon-rich Pomeron-Pomeron interactions with gluon-poor photon-photon interactions in the same experiment.

The Relativistic Heavy Ion Collider (RHIC) at BNL is nearing completion, and at least one detector (STAR) will be able to trigger on low-mass, low-multiplicity events from two-photon interactions and central production (Nystrand and Klein, 1988).

E. The Japanese Hadron Facility

Peripheral hadronic production experiments still have plenty to contribute, particularly using state-of-the-art detection systems. This is particularly true if they can look at production mechanisms and final states that have been largely ignored in the past, as has been aptly demonstrated by E852 at BNL, for example. It is also important to recognize the lessons taught by the LASS experiment at SLAC which, among other things, demonstrated the effectiveness of a programmatic approach and the value of analyzing many different channels in the same experiment under conditions of uniform and well-understood acceptance. Future experiments should follow these models.

The most promising scenario for such developments is the Japanese Hadron Facility (JHF) that has been proposed for KEK. This would be a high-intensity (10 μA) 50-GeV proton synchrotron, including polarized beams.

For comparison, the AGS at Brookhaven, one of the primary workhorses of spectroscopy through peripheral hadroproduction, produces a few-μA proton beam, but with energy around 25 GeV. The higher energy is worth large factors in the secondary $\pi^{\pm}$ and $K^{\pm}$ beams at 20 GeV or so.

F. CEBAF at Jefferson Laboratory

Peripheral photoproduction is conspicuously absent from among the studied production mechanisms in hadron spectroscopy. This is despite the fact that one expects profound new results from such experiments, as discussed in Sec. III.B. The reason, however, is obvious. Until recently, virtually no suitable facilities have existed for detailed work in this area.

The Continuous Electron Beam Accelerator Facility (CEBAF) at the Thomas Jefferson National Accelerator Facility (Jefferson Lab) is a very-high-intensity electron-beam facility with 4–6-GeV electron beams. The beam energy was chosen to pursue the laboratory's original primary goal, i.e., the study of nuclear structure. However, because of landmark improvements in the RF acceleration cavities, beams with energies up to ~ 8 GeV will be possible with minimal cost. In fact, a plan is in place to push the electron-beam energy initially to 12 GeV and hopefully later to 24 GeV. High-energy photons can be produced by either thin-radiator bremsstrahlung or by a variety of other means.

An aggressive experimental plan is underway that will keep pace with the accelerator improvements. This involves the construction of a new experimental hall at the site, which would include a dedicated experimental facility. Initial designs of this facility are taken directly from the LASS and E852 experiences, with necessary modifications for photon beams. State-of-the-art high-rate data acquisition electronics and computing will be an integral part of this new experiment.

G. A τ-charm factory at BEPC

τ-charm factories, e^+e^- colliders operating in the energy region of 3–4 GeV, have been proposed although none have yet been approved. These factories would produce large numbers of J/ψ and D mesons. The BEPC ring in China, presently home to the BES experiment, has worked well and is considered a prototype for a higher-luminosity facility at that laboratory. Present designs call for a factor of 10–100 improvement in the luminosity over the present storage ring. The project is not yet approved, however.

Hadron spectroscopy can be pursued in several ways at a τ-charm factory. These include J/ψ and ψ' decays, decays of $\tau \to \nu_\tau +$hadrons where the hadrons are produced by a virtual W boson, and semileptonic and leptonic decays of $D^{\pm}$, D^0. Probably the most important question in light-meson spectroscopy is the existence of gluonic excitations, and, as demonstrated earlier in this review, $J/\psi \to \gamma X$ is a critical reaction for the complete understanding of these states.

With the high statistics available it may even be possible to perform a partial-wave analysis of the χ_{cJ} decay products produced in $\psi' \to \gamma \chi_{cJ}$ radiative decays. For example, $\chi_{c1} \to \pi H$ is sensitive to the hybrid exotic sector H ($J^{PC}=1^{-+}$), while $\chi_{c0} \to f_0(980)X$ would be a source of 0^{++} mesons.

In addition to high-statistics searches for gluonic excitations, a τ-charm factory will also study the properties of the $a_1(1270)$, a_1', ρ, ρ', K^*, $K_1(1270)$, and $K_1(1400)$ mesons via hadronic decays of the τ lepton. These observations will help disentangle the radial excitations of the vector mesons.

VII. FINAL COMMENTS

In this review we hope to have conveyed the sense that meson spectroscopy is a lively, exciting subject with even the most basic questions still unresolved. We find it remarkable that after over 20 years of QCD we still do not know what the physical states of the theory are. Understanding QCD, and non-Abelian gauge theories in general, is one of the most important problems facing high-energy and nuclear physics. We have outlined the important issues, the puzzles, and the open questions, in light-meson spectroscopy. We have attempted to show where the field is heading with the next generation of experiments and how they can advance our knowledge of meson spectroscopy. In many cases, it is not only our lack of experimental data but a lack of answers to theoretical questions that has hindered progress in the field.

Our survey of established meson states shows the basic validity of the constituent-quark model. However, the survey also highlights some possible discrepancies with the predictions of the quark model which may point to the need to go beyond the $q\bar{q}$ states and include gluonic excitations and multiquark states in the light-quark sector. In this sense we may be on the verge of opening up a new frontier of hadron spectroscopy in the 1.5–2.5-GeV mass region of the meson sector where the vast majority of the complications seem to occur. Many $q\bar{q}$ states remain to be discovered. It is important to find them and then to pin down the details of the $q\bar{q}$ spectrum to pave the way for the search for exotic states.

Most of the significant deviations from the quark model occur in the light-quark isoscalar sector, and much of the current excitement is with the scalar mesons. Even if we relegate the $f_0(980)$ [along with the isovector $a_0(980)$] to multiquark status, we still have to contend with the $f_0(1500)$ and the $f_J(1710)$ (if indeed $J=0$). It is quite possible that these two states are mixtures of the ground-state glueball and the $s\bar{s}$ scalar meson, but we will have to wait for better experiments and clearer phenomenology to establish a consistent pattern.

The long-standing E/ι problem remains with us. The evidence clearly points to two separate states each, for $J^{PC}=1^{++}$ and 0^{-+}, whereas exactly one of each is expected based on $q\bar{q}$ degrees of freedom. Ascribing hybrid, glueball, or multiquark status to the extra states is problematic, because no clear picture emerges from the current set of experimental data. New high-statistics measurements of $\gamma\gamma$ and $\gamma\gamma^\star$ production may be very helpful. Certainly a better phenomenological understanding of reactions near the $K^\star \overline{K}$ threshold is crucial for progress here.

Despite its long history, the $f_J(2220)$ is still enigmatic. It may or may not be a peculiar, narrow meson representing the 2^{++} glueball, and it may or may not have an underlying broad structure that is a manifestation of 2^{++} or 4^{++} strangeonium. We also continue to deal with a possible overpopulation of 2^{++} states, in particular the various f_2 states observed in $\phi\phi$ decay.

Finally, there are exotic mesons that necessarily imply a meson state beyond $q\overline{q}$ models. Evidence is finally beginning to sort itself out in the $\eta\pi$ channel, and now does seem to point towards an exotic $J^{PC}=1^{-+}$ state near 1380 MeV/c^2. It is difficult to accommodate this, however, in present models of excited gluonic degrees of freedom. Some evidence is now emerging for higher-mass exotic states, but more data will be necessary, particularly in reactions that are expected to enhance the production rate. A good candidate laboratory is peripheral photoproduction, and new facilities are on the horizon.

Hadron spectroscopy is undergoing a renaissance, taking place at many facilities worldwide. We anxiously look forward to new results.

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