Light Meson Spectroscopy

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Submitted to Review of Modern Physics
November 12, 1998

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Abstract

We survey the current status of light meson spectroscopy. We begin with a general introduction to meson spectroscopy and and its importance in understanding the physical states of Quantum Chromo Dynamics (QCD). Phemenological models of hadron spectroscopy are described with particular emphasis on the constituent quark model and the qualitative features it predicts for the meson spectrum. We 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 we find that most meson states are 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. We end with a brief description of future directions in meson spectroscopy.

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1 INTRODUCTION

Meson physics and the strong interactions have been intimately connected since pions were first introduced by Yukawa to explain the inter-nucleon force (Yukawa, 1935). Since that time, our knowledge of mesons and in parallel, our understanding of the strong interactions, has undergone several major revisions. Our present understanding of the strong interactions is that it is described by the non-Abelian gauge field theory Quantum Chromodynamics (QCD) (Fritzch, 1971; Gross, 1973; Weinberg, 1973) which describes the interactions of quarks and gluons. Once again, it appears that mesons are the ideal laboratory for the study of strong interactions in the strongly coupled non-perturbative 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. The understanding of QCD 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 is the laboratory that will hopefully elucidate this theory.

To a large extent our knowledge of hadron physics is based on phenomenological models and in particular, the constituent quark model (G. Zweig, 1964; M. Gell-Mann, 1964)[1]. Meson and baryon spectroscopy is described surprisingly well as composite objects made of constituent objects — valence quarks. We will 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 interpret 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 that there is too much 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 new types of hadrons which 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.

[1]Introductions to the quark model are given in (Isgur, 1980) and (Rosner, 1984)

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

Mesons were first introduced by Yukawa (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 a confused state of understanding. Eventually, by arranging the various mesons and baryons 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) although none of the observed states seemed to correspond to the fundamental triplet representation. In an important conceptual leap Zweig (Zweig, 1964) and Gell-Man (Gell-Man, 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-Man referred to them as quarks. By taking this simple picture seriously, the qualitative properties of hadrons were explained quite well. Serious problems remained however. In the "naive" quark model the spin $\frac{3}{2}$ baryons, the constituent quarks' spin wavefunctions were symmetric as were their flavour wavefunctions. Being fermions, the baryon wavefunction should be antisymmetric in the quark quantum numbers. This would imply that either quarks obeyed some sort of bizzare para-statistics or that the ground state spatial wavefunction was 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 he named colour, with respect to which the quark wavefunctions 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.

By the beginning of the 1970's it was becoming clear that the weak interactions could be explained by gauge theories (Glashow 1961, Weinberg 1967, Salam 1968). If this was the case, it seemed reasonable that the strong interactions should also be described using the same formalism. "Gauging" the colour 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, 1971; Gross, 1973; Weinberg, 1973).

Nevertheless there was still considerable skepticism about the existence of quarks since they had never been seen. This situation changed when in November 1974 very narrow hadron resonances were discovered simultaneously at Brookhaven National Laboratory (Aubert 1974) and the Stanford Linear Accelerator Center (Augustin 1974). These states, named the J/ψ, were quickly interpreted as being bound states of a new heavy quark– the charm quark. Quark models which incorporated the qualitative features of QCD, asymptotic freedom and confinement, where able to reproduce the charmonium ($c\bar{c}$) spectrum rather well. (Appelquist 1975a, 1975b, 1975c, Eichten 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, deRujula, Georgi, and Glashow, (De Rújula 1975) showed that these ideas could successfully be used to describe the phenomenology of light quark spectroscopy.

With the acceptance that QCD is the theory of the strong interactions comes the need to understand its physical states. Understanding the spectrum of hadrons reveals information on the non-perturbative aspects of QCD. Unfortunately, calculating the properties of hadrons from the QCD Lagrangean has proven to be a very difficult task in this strongly coupled non-linear theory. In the long term, the most promising technique is formulating the theory on a discrete space-time lattice (Creutz 1983a, 1983b, Kogut 1979, 1983, Montvay 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, these calculations take enormous amounts of computer time and progress has been slow. 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 $q\bar{q}$ 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 will often refer to glueballs and hybrids as exotics because they lie outside the constituent quark model. However they are not exotics in the sense that if they 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 criteria 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 which 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 of the problem is that these exotics may have properties which 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 between conventional and exotic 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 which could help resolve them. To this end we will begin with a discussion of the theoretical ingredients relevant to this article. In the course of this review we will refer to numerous experiments so in section 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 section IV we will compare the predictions of one specific quark model with experiment. This will allow us to identify discrepancies between the quark model and experiment which may signal physics beyond conventional hadron spectroscopy. In section V we will go over these puzzles in detail to help decide whether the discrepancy is most likely a problem with the model, with a confused state in experiment, or whether it most likely signals some interesting new physics. In section VI we will briefly outline some future facilities for the study of meson spectroscopy that are under construction or that are being considered. Finally, in section VII we will 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 the surface of many interesting topics. There are a number of recent reviews of meson spectroscopy and related topics with emphasis somewhat different from that of the present one. We strongly encourage the interested reader to refer to these reviews for further details. The reader is referred to the reviews by F. Close (1988), S. Cooper (1988), B. Diekmann (1988), T. Burnett and S. Sharpe (Burnett 1990), N. Törnqvist (1990), C. Amsler and F. Myhrer (1991), K. Konigsman (1991), R. Landau (1996), C. Amsler (1998), and Barnes (1998). In addition, the Review of Particle Physics (Particle Data Group (PDG), Caso 1998) contains a wealth of information on the properties of mesons in its tables of properties and mini-reviews on topics of special interest and should be consulted for further information.

2 THEORETICAL OVERVIEW

2.1 Quantum Chromodynamics

Quantum Chromodynamics (QCD), the theory of the strong interactions, (Fritzch 1971, Gross 1973, Weinberg 1973) may be thought of as a generalization of quantum electrodynamics (QED), our most successful physical theory. QCD is described by the Lagrangean;

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

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,...,8, q_i are the quark fields with colour indices i=1,2,3, g is the bare coupling, m is the quark mass, and $\frac{\lambda^i}{2}$ are the generators of SU(3). One immediately observes that quarks couple to gluons in much the same way as the electron couples to photons with $e\gamma^\mu$ of QED replaced by $g\gamma^\mu \frac{\lambda}{2}$ of QCD. The significant difference between QED and QCD is that in QCD the quarks come in coloured triplets and the gluons in a colour octet where colour is labelled 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, very difficult to solve and leads to the confinement of colour. 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 spacetime lattice, in analogy to the approach one might take in the numerical solution of a difficult differential equation (Wilson 1974, Creutz 1983a, 1983b, Kogut 1979, 1983, C. Michael 1995, 1997, Montvay 1994). 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 progress is 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 will use the predictions of QCD inspired models as the basis for interpreting the nature of the observed mesons.[3] It is therefore useful to sketch the QCD motivation for these models. We start with the quark-antiquark ($Q\bar{Q}$) potential in the limit of infinitely massive quarks which can be used in the Schrödinger equation to obtain

[3]For a recent review see Barnes 1996.

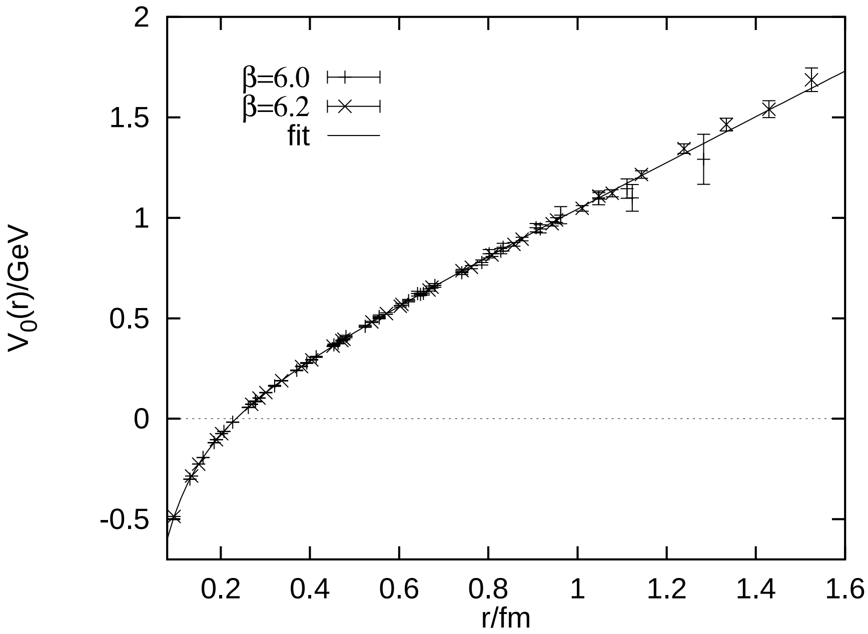

Figure 1: The static $Q\bar{Q}$ potential from Bali, Schilling, and Wachter (Bali 1997).

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 $Q\bar{Q}$ 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:

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

where the constant b, the "string tension", is numerically found to be $b \simeq 0.18$ GeV^2 $\simeq$ 0.9 GeV/fm and

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

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

There are also spin dependent forces between the quarks analogous to the spin dependent forces in the hydrogen atom which give rise to the fine and hyperfine 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 terms next order in an expansion in v^2/c^2 (Eichten

1979, 1981, Gromes 1984a, 1984b):

$$
\begin{aligned}
V_{spin}(r) \quad = \quad & \left(\frac{1}{2m_1^2}\vec{L}\cdot\vec{S}_1 + \frac{1}{2m_2^2}\vec{L}\cdot\vec{S}_2\right)\frac{1}{r}\frac{d}{dr}\left(V(r)+2V_1(r)\right) \\
& + \frac{1}{m_1 m_2}\vec{L}\cdot(\vec{S}_1+\vec{S}_2)\frac{1}{r}\frac{dV_2(r)}{dr} \\
& + \frac{1}{m_1 m_2}(\hat{r}\cdot\vec{S}_1\,\hat{r}\cdot\vec{S}_2 - \frac{1}{3}\vec{S}_1\cdot\vec{S}_2)V_3(r) + \frac{1}{3m_1 m_2}\vec{S}_1\cdot\vec{S}_2\,V_4(r)
\end{aligned}
\tag{5}
$$

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, $V_1 - V_4$, can be related to correlation functions of the colour electric and colour magnetic fields. For example V_3 and V_4 are given by (Eichten 1979, 1981, Gromes 1984a, 1984b):

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

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:

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

The lattice results do indeed agree with these spin dependent potentials (Bali 1997, Campostrini 1986, 1987). Bali Schilling and Wachter (Bali 1997) have studied the $b\bar{b}$ and $c\bar{c}$ 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 is shown in Fig 2 and 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 and the neglect of higher order relativistic corrections. Direct lattice calculations of spin-dependent splittings also agree with the measured splittings (Davies 1998).

Although historically, the spin dependent potentials were obtained phenomenologically by comparing the observed quarkonium spectrum ($c\bar{c}$ and $b\bar{b}$) to the predictions of potential models, it turns out that the resulting potentials are in reasonable agreement with those obtained from lattice 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

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

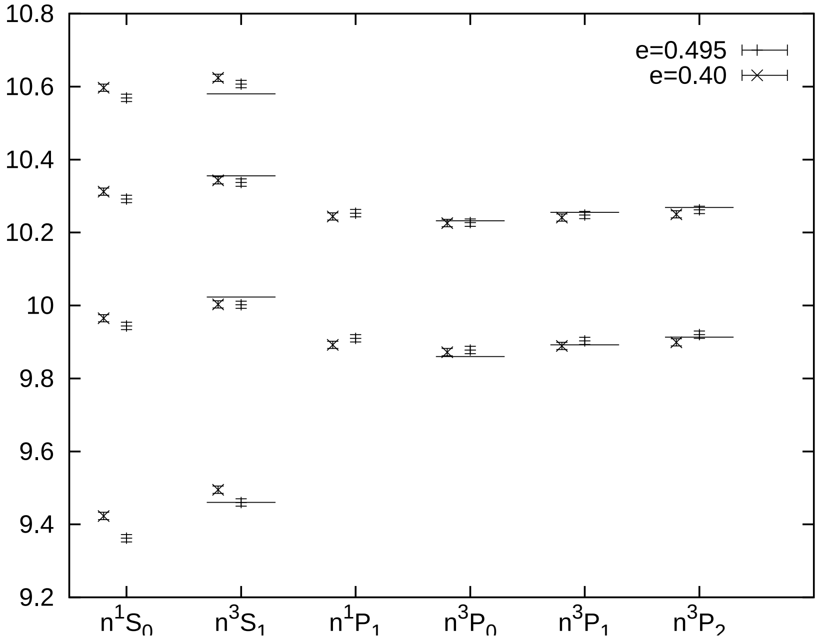

Figure 2: The bottomonium spectrum (in GeV) calculated using potentials from lattice QCD. The horizontal lines are experimental results. From Bali, Schilling, and Wachter (Bali 96).

is the colour 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 colour 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.

2.2 Colour Singlets in QCD

Because of confinement only colour singlet objects can exist as physical hadrons. Coloured quarks form the fundamental triplet (3) representation of the SU(3) colour gauge group and antiquarks the conjugate $\bar{3}$ representation. Therefore, a quark-antiquark pair can combine to form a colour singlet as can three quarks while a quark-quark pair cannot. Other states are also possible, for example $q\bar{q}q\bar{q}$, and it is dynamical question whether such multiquark systems are realized in nature as single multiquark states, as two distinct $q\bar{q}$ states, or as a loosely associated system of colour singlet mesons analogous to a diatomic molecule. Colour 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 there is apriori no reason that the physical mesons cannot be linear combinations of $q\bar{q}$, $q\bar{q}q\bar{q}$, gg, and $q\bar{q}g$.

2.3 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 $P(q\bar{q}, L) = (-1)^{L+1}$ and C-Parity is also defined for neutral self-conjugate mesons and is given by $C(q\bar{q}, L, S) = (-1)^{L+S}$. 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, $\sim 5 - 10$ 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 $\bar{u},\ \bar{d}$ into mesons forms isospin singlet and triplet multiplets. We will use the symbol n (for non-strangeness) to generically stand for u or d. Thus, one should read

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

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 inherent approximate nature of isospin.

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

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

Table 1: The quantum numbers and names of conventional $q\bar{q}$ mesons.

		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
.	.	.	.	.	.	.
.	.	.	.	.	.	.
.	.	.	.	.	.	.

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

where the flavour mixing angle $\phi \simeq 45^o$. These states are often also expressed as linear combinations of flavour 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 wavefunctions with the quark flavour wavefunctions results in the meson states of Table 1 where we have used the Particle Data Group naming conventions (PDG, Caso 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 $q\bar{q}$ potential, 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, 1D, 2S, 2P, ..., levels is sensitive to the details of the potential. Fig. 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..

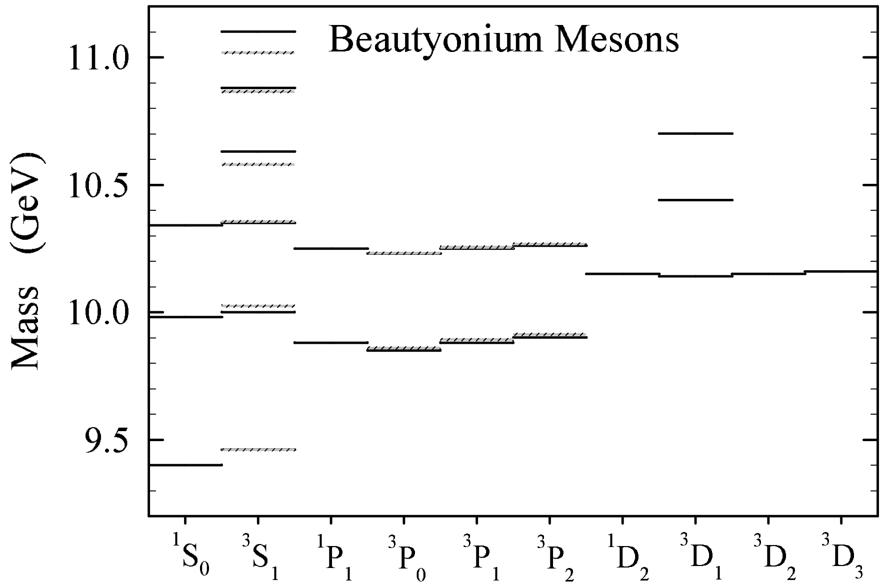

Figure 3: The $b\bar{b}$ spectrum from a quark potential model (Godfrey, 1985a). The solid lines are the quark model predictions and the shaded regions are the experimental states.

The preceeding discussion may seem like a lengthy digression but it demonstrates an important result; that 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 quarkonia 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 extending 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 1985a). In heavy quark systems the former splitting is a non-relativistic 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 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 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 will use the predictions of one such attempt at a "relativized" quark model (Godfrey, 1985a, 1985c, 1986), which we take to be representative of the many similar models in the literature (Stanley 1980, Carlson 1983a, 1983b, Gupta 1986, Brayshaw 1987,

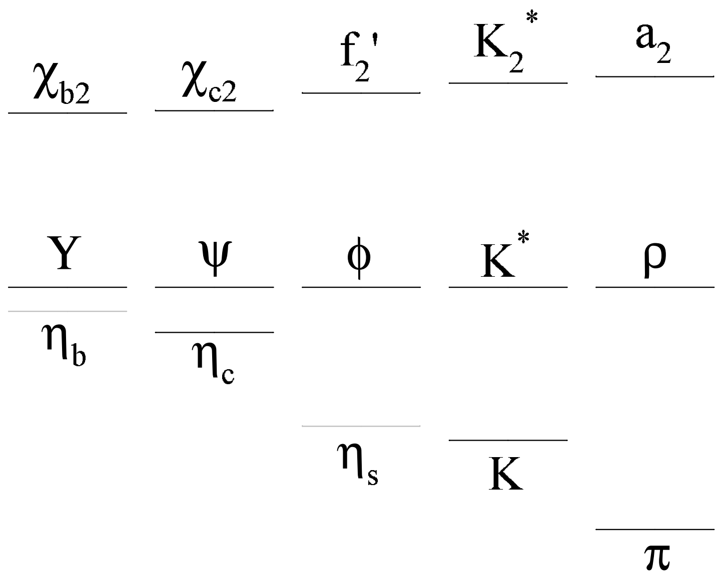

Figure 4: The evolution of the $1^3P_2 - 1^3S_1$ and the $1^3S_1 - 1^1S_0$ splittings as a function of quark masses (Godfrey 1985a). The splittings are drawn to scale. Note that the η_b and the η_s (dotted lines) are in fact calculations.

Crater 1988, Olson 1992, Fulcher 1993, Jean 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 we must consider in 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_{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 non-relativistic result:

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

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 equations 9-11.

The confinement potential, Eqn. 19, includes the spin independent linear confinement and Coulomb like interactions. The Coulomb piece dominates at short distance while the linear piece dominates at large distance. Because heavy quarkonium have smaller radii they are more sensitive to the short range colour-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, measurement of both heavy quarkonium and mesons with light quark content probe different regions of the confinement potential and complement each other. The linear character of the trajectories 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, 1985c).

The spin dependent parts of the potential consist of the colour hyperfine interaction (Eqn. 9) and the spin-orbit interaction (Eqn. 10,11). The colour hyperfine interaction is responsible for $^3S_1 -^1 S_0$ 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 $^1L_J -^3 L_J$ mixing. The spin-orbit interaction has two contributions, $H_{ij}^{s.o.(cm)}$ and $H_{ij}^{s.o.(tp)}$ given by Eqns. 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 1985b, 1985c, Isgur 1998). Thus, the multiplet splittings act as a probe of the confinement potential, providing information on non-perturbative QCD.

2.4 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 wavefunctions 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 and 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. 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

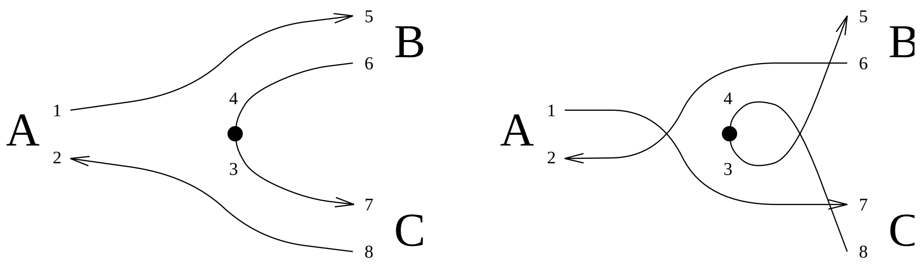

Figure 5: Diagrams contributing to the meson decay $A \to BC$. In many cases only one of these diagrams will contribute.

in a $^3S_1$ state with $J^{PC} = 1^{--}$ while in the other major variant, the $^3P_0$ model of LeYaouanc *et al.* (LeYaouanc 1973, 1974, 1975), the quark pair creation process is viewed as an inherantly nonperturbative process where the $q\bar{q}$ pair is formed with the quantum numbers of the vacuum, $J^{PC} = 0^{++}$, and therefore is in a $^3P_0$ state. Geiger and Swanson (Geiger 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 Isgur and Kokoski (Kokoski 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 are given by Kokoski and Isgur (Kokoski 1987), Blundell and Godfrey (Blundell 1996), and Barnes Close Page and Swanson (Barnes 1997) which can be used to compare with experiment. 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.

Electromagnetic couplings are another source of useful information about resonances. A first example is two photon couplings which are measured in the reaction $e^+e^- \to e^+e^- +$ 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, $\Gamma_{\gamma\gamma}(f_2) \cdot B(f_2 \to \pi^0\pi^0)$, can be determined. In principle the absolute two-photon widths can be calculated from quark model wavefunctions (Godfrey 1985a) but for light $q\bar{q}$ mesons the results are sensitive to relativistic effects (Ackleh 1992). To some extent this 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

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

This gives, for example, the relative amplitudes of

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

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

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

The $L = 2$ $\pi_2(1670)$ and probably the η_2 state have been observed in two-photon production (Antreasyan 1990, Behrend 1990, Karch 1992). Because glueballs do not have valence quark content, the observation or non-observation of a state in two photon production provides another clue about the nature of an observed state.

Single photon transitions $(q\bar{q})_i \to \gamma(q\bar{q})_f$ is another useful measurement for identifying $q\bar{q}$ states. These have the characteristic pattern of rates based on flavour;

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

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, $a_2 \to \pi\gamma$ and $a_1 \to \pi\gamma$. Since measurements of radiative transitions could determine the nature of controversial states such as the $f_1(1420)$, which will be discussed in Section V, they should be carried out if possible. These measurements could be made in electroproduction, the inverse reaction.

2.5 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 1976, Barnes 1984, 1985a, Chanowitz 1983a, 1983b, Close 1988, Godfrey 1989, Isgur 1989a).

2.5.1 Glueballs

Many different QCD based models and calculations make predictions for such states; Bag Models (Barnes 1977, 1983c, DeTar 1983, Chanowitz 1983a, Hasenfratz 1980). Constituent Glue Models (Horn 1978), Flux Tube Models (Isgur 1983, 1985a, 1985b), QCD Sum Rules (Latorre 1984), and Lattice Gauge Theory. Recent Lattice QCD calculations are converging towards agreement (Schierholz 1989, Michael 1989, Sexton 1995b, Bali 1993, Teper 1995, Morningstar 1997) although there is still some variation between the various calculations. 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 1993),

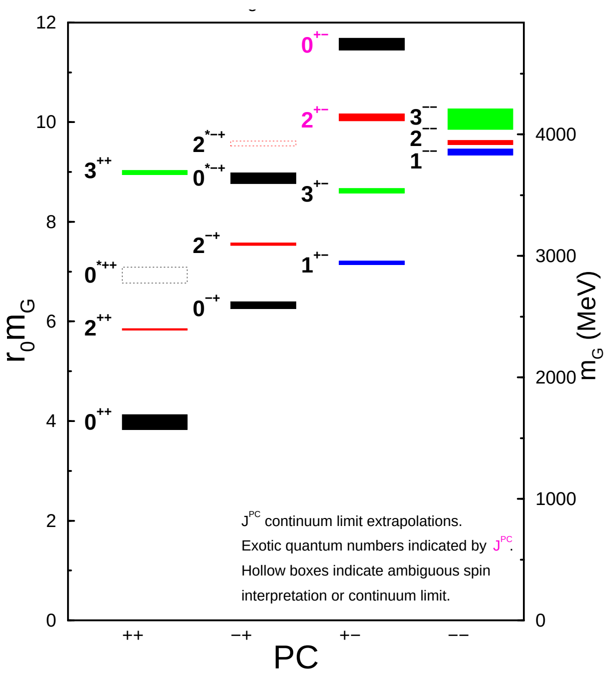

Figure 6: The mass of the glueball states. The scale is set by r_0 with $1/r_0 = 410(20)$ MeV (Morningstar 1997, Peardon 1998).

1600 ± 160 MeV (Michael 1998), 1648 ± 58 (Vaccarino 1998), and 1630 ± 100 MeV (Morningstar 1997). The difference between these results lies mainly in how the mass scale is set. The next lightest states are the 2^{++} with mass estimates 2232 ± 220 MeV (Michael 1998), 2270 ± 100 MeV (Bali 1993), and 2359 ± 128 MeV (Chen 1994) 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 $\sim$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 glueball candidate's properties to the expected properties of glueballs and conventional mesons.

Sexton *et al.* (Sexton 1996) have estimated the width of the 0^{++} glueball to all possible pseudoscalar pairs to be 108 ± 29 MeV. This number combined with reasonable guesses for the effect of finite lattice spacing, finite lattice volume, and the remaining width to multibody states yields a total width small enough for the lightest scalar glueball to be easily observed. The significant property of glueball decays is that one expects them to have flavour-symmetric couplings to final state hadrons. This gives the characteristic flavour-singlet branching fraction to pseudoscalar pairs (factoring out phase space)

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

Of course, one should also expect some modifications from phase space, the glueball wavefunction, and the decay mechanism (Sexton 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 and 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 $J/\Psi \to \gamma G$ where the glueball is formed from intermediate gluons (Cakir 1994, Close 1997c). The second is in central production $pp \to p_f(G)p_s$ 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):

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

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 (Close 1997c).

The simple picture presented above is likely to be muddied by complicated mixing effects between the pure LGT glueball and $q\bar{q}$ states with the same J^{PC} quantum numbers (Amsler and Close 1995, 1996). Lee and Weingarten (Lee 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 valence approximation on the lattice. With this motivation they perform a phenomenological fit which finds the $f_0(1710)$ to be $\sim 74\%$ glueball and the $f_0(1500)$ to be $\sim 98\%$ 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 will 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.

2.5.2 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 flavour 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, Chanowitz 1987, Close 1988, Godfrey 1989, Barnes 1995, Close 1995a, Barnes 1996, 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 which cannot be accomodated 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 1986) from a colour 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 $\hat{\pi}$ and an isospin 0 $J^{PC} = 0^{--}$ meson would be a $\hat{\eta}$ 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 insights 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 1983, Perantonis 1990, Morningstar 1997). This is analogous to the calculation of the nucleus-nucleus potential in diatomic molecules where the slow heavy quarks and fast gluon fields in hybrids correspond to the nuclei and electrons in diatomic molecules. One treats the quark and antiquark as spatially fixed colour 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. 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^{-+} (pseudo-electric) gluonic excitations, spin-orbit terms, as well as mixing between hybrid states and $q\bar{q}$ mesons with non-exotic spins.

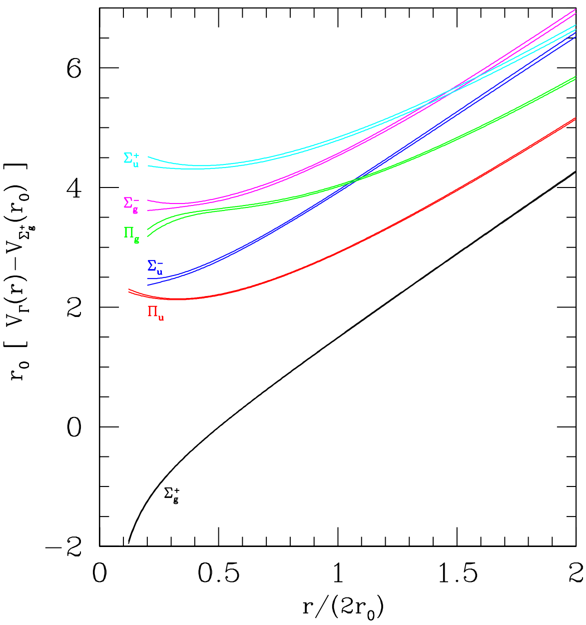

Figure 7: A set of hybrid adiabatic surfaces for static central potentials. $\Lambda = \Sigma,\ \Pi,\ \Delta, \ldots$ corresponds to the magnitude of $J_{glue} = 0,\ 1,\ 2, \ldots$ projected onto the molecular axis. The superscript $= \pm$ corresponds to the even or oddness under reflections in a plane 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 labelled 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. (Juge 1998)

While this picture is appropriate for heavy quarkonia 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 also 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 1983, 1985a, 1985b). It predicts 8 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 1990, Michael 1994).

Other models exist, for example the bag model (Chanowitz 1983a, 1983b, Barnes 1983a, 1983b, 1983c) which in contrast to the flux tube model, expects only four degenerate states, the 2^{-+}, 1^{-+}, 1^{--}, and 0^{-+}. This difference is symptomatic of the differences between the two models. In the bag model, the gluon degrees of freedom are either transverse electric (TE) or transverse magnetic (TM) modes of the bag, with the TM mode considerably higher in mass than the TE mode.

A recent Hamiltonian Monte Carlo study of the flux tube model (Barnes 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 1990, Lacock 1996, 1997) who find $M_{\hat{\rho}} \simeq 1.88$ GeV and $M_{\hat{\phi}} \simeq 2.09$ GeV and Bernard *et al.* (Bernard 1996, 1997) who find $M_{\hat{\rho}} \simeq 1.97$ GeV and $M_{\hat{\phi}} \simeq 2.17$ GeV. Lacock *et al.* (Lacock 1996) also 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, it must instead appear as internal angular momentum of the $q\bar{q}$ pairs (Page 1997a, 1997b, Kalashnikova 1994). 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,

$$\hat{\rho} \rightarrow [\underline{\pi\eta}, \underline{\pi\eta'}, \pi\rho, K^*K, \eta\rho, \ldots]_P, \quad [\pi b_1, \pi f_1, \eta a_1, KK_1 \ldots]_S, \tag{26}$$

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 1985b, Close 1995a) are given in Table 2. 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 $\Gamma \approx 100$ 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 hybrids are:

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

Finally, some "forbidden" decays such as $\hat{\rho}(1900) \rightarrow \rho\pi$ have small but finite partial widths due to differences in the final state spatial wavefunctions. Thus, it may be possible to observe hybrids in these simpler decay modes in addition to the favoured 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 and one might wonder whether the nonexotic hybrids might be narrow enough to be observables. According to the results of Close and Page (Close 1995a) reproduced in Table 2 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 only ~ 100 MeV which decays to $K_1(1270)K$ and $K_1(1400)K$. The ϕ is also relatively narrow

Table 2: The dominant hybrid decay widths for $A \to [BC]_L$ for partial wave L calculated using the flux tube model. From Close and Page (Close 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 (Merlin 1987).

	$I=1$		$I=0\ n\bar{n}$		$I=0\ s\bar{s}$	
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	$[f_2(1270)\eta]_S$	~ 50		
	$[a_2(1320)\eta]_S$	~ 40	$[K_2^*(1430)K]_S$	~ 30		
	$[K_2^*(1430)K]_S$	~ 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
	$[b_1(1235)\eta]_P$	~ 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		

Table 2, continued.

	$I=1$		$I=0\ n\bar{n}$		$I=0\ s\bar{s}$	
A	$[BC]_L$	Γ	$[BC]_L$	Γ	$[BC]_L$	Γ
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
	$[f_0(1300)\pi]_P$	~ 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	$[K_2^*(1430)K]_P$	~ 20		
	$[K_2^*(1430)K]_P$	~ 20	$[K_1(1270)K]_P$	40		
	$[K_1(1270)K]_P$	40	$[K_1(1400)K]_P$	~ 20		
	$[K_1(1400)K]_P$	~ 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	$[K_1(1400)K]_S$	~ 125		
	$[K_1(1270)K]_S$	20				
	$[K_1(1400)K]_S$	~ 125				
	$[\rho\pi]_P$	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		

Table 3: Decay of quark model and hybrid $\pi(1800)$.

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

with $\Gamma_{tot} \sim 225$ MeV. Two more interesting hybrids are the π_2 with $\Gamma_{tot} \sim 170$ MeV. and its $s\bar{s}$ 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 non-exotic quantum numbers is a conventional $q\bar{q}$ state or a hybrid one would make use of the detailed predictions we have described above for the two possibilities (Barnes 1997). For example, a selection rule of the $^3P_0$ decay model forbids the decay of a spin-singlet $q\bar{q}$ state to two spin-singlet mesons in the final state (Page 1997a). This selection rule forbids the decay $\pi_2(^1D_2) \to b_1\pi$ 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 3. 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. 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 $q\bar{q}$ 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 $q\bar{q}$ states.

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

2.6 Multiquark Hadrons

The notion of colour naturally explained nature's preference for $q\bar{q}$ and qqq colourless sytems. However, it also appears to predict multiquark states such as $q^2\bar{q}^2$ and $q^3q\bar{q}$ which could have exotic quantum numbers, thus indicating non $q\bar{q}$ and qqq states (Jaffe 1977a, 1977b, 1978, Lipkin 1978). Upon considering $qq\bar{q}\bar{q}$ systems we find that the colour couplings are not unique as they are in mesons and baryons. For example, we can combine two colour triplet q's into a colour 6 or $\bar{3}$. Likewise we can combine two antitriplet $\bar{q}$'s into a 3 or a $\bar{6}$. Therefore, there are two possible ways of combining $qq\bar{q}\bar{q}$ into a colour singlet: $3\bar{3}$ or $6\bar{6}$. In addition, since we could have combined a q and $\bar{q}$ into a colour 1 or 8 we could also have combined the $qq\bar{q}\bar{q}$ into colour 11 and 88. Since two free mesons (in a colour 11) are clearly a possible combination of $qq\bar{q}\bar{q}$ the $3\bar{3}$ $6\bar{6}$ couplings can mix to give the 11 - 88 colour 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 $qq\bar{q}\bar{q}$ configurations lead to hadron-hadron potentials (Barnes 1987, Weinstein 1990, Swanson 1992). Both must be taken into account when attempting to unravel the hadron spectrum.

As in the case of hybrid mesons, states not fitting into the $q\bar{q}$ framework are the most unambiguous signature for multiquark states. In particular, flavour exotics are our best bet for finding genuine multiquark states. There is a large literature on the physics of multiquark states which 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 $qq\bar{q}\bar{q}$ system (Weinstein 1982, 1983). It found that weakly bound $K\bar{K}$ "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 $f_0(975)$ and $a_0(980)$. The meson-meson potentials which 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 $K\bar{K}$ molecules are the exception however, as the model predicts that in general the $qq\bar{q}\bar{q}$ 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 1992) and pseudoscalar-vector channels (Caldwell 1987, Longacre 1990). One such possibility is a threshold enhancement in K^*K just above threhold.

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 $K\bar{K}$ despite having almost no phase space is a hint that it is not a conventional $q\bar{q}$ state. Electromagnetic couplings are another clue to nonstandard origins of a state. Barnes found that the two-photon widths for a $q\bar{q}$ state are expected to be much larger than that of a $K\bar{K}$ 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 $f_1(s\bar{s})$ state would be dominated by the transition $f_1(s\bar{s}) \to \gamma\phi$. The two cases would be distinguished by different ratios of the $\gamma K^0\bar{K}^0$ and γK^+K^- final states. Likewise, Close Isgur and Kumano (Close 1993) suggest a related test for the $f_0(980)$ and $a_0(980)$ involving the radiative decays $\phi \to \gamma f_0$ and γa_0.

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 unravelled and 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 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 $\bar{p}p \to K\bar{K}\pi$. The lesson is that final state interactions arising from hadron-hadron potentials will play an important role in understanding the 1 to 2 GeV mass region.

3 EXPERIMENTS

The meson spectrum consists of an increasingly large number of broad, overlapping states, crowded into the mass region above around one GeV/c^2. Experiments disentagle these states using a combination of three approaches. First, states have different sensitivity to the various available production mechanisms. This can be particularly useful when trying to enhance one state relative to another, comparing the results from two or more sources. 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 technique 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, 1985; Sinervo, 1993; Cummings and Weygand, 1997). In all cases, it is important to fully understand 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, 1998) in Fig. 8, where the $\pi^-\pi^+\pi^-$ mass spectrum from the reaction $\pi^-p \to \pi^-\pi^+\pi^-p$ shows a clearly rich structure of strongly overlapping peaks. In this case, the partial wave analysis is able to cleanly separate the contributions from the various J^{PC}. Furthermore, one can examine the relative phase motion of one "wave" relative to the other, and demonstrate that the intensity peak is in fact resonant.

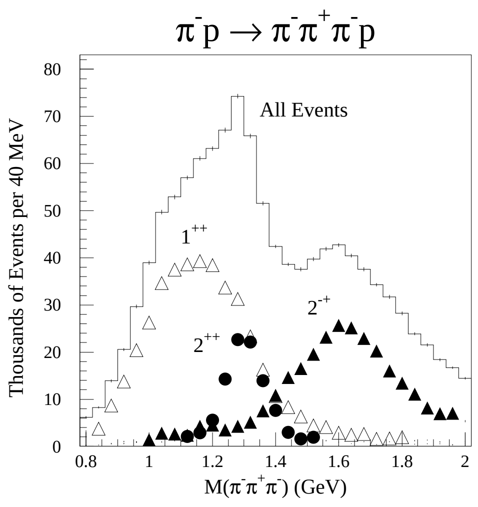

Figure 8: Three-pion mass distribution for the reaction $\pi^-p \to \pi^-\pi^+\pi^-p$ at 18 GeV/c, from experiment E852 at BNL. A partial wave analysis is used to decompose the spectrum into its dominant J^{PC} components, clearly showing the $a_1(1260)$, the $a_2(1320)$, and the $\pi_2(1680)$.

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 chapters.

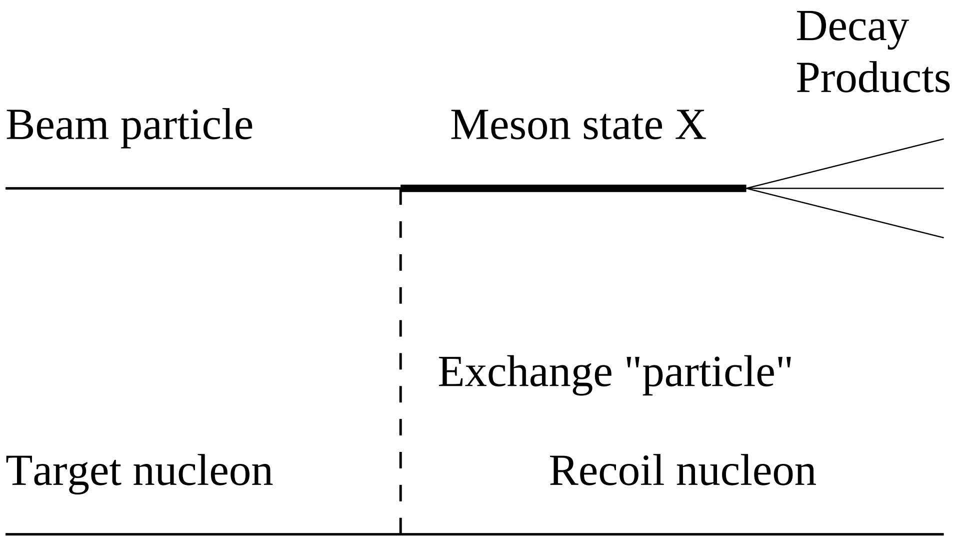

Figure 9: Schematic diagram of a hadronic peripheral production process. Momentum is exchanged through an off-mass-shell particle, which may or may not be charged.

3.1 Hadronic Peripheral Production

Most of the data on light meson spectroscopy has 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 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, 1988d) and $\pi^-p \to X^-p \to \rho^0\pi^-p \to \pi^-\pi^+\pi^-p$ by the E852 collaboration (Adams, 1998). Other recent experiments with π^- beams include VES at IHEP/Serpukhov (Beladidze, 1993; Gouz, 1993) and BENKEI at KEK (Fukui, 1991; Aoyagi, 1993). Of particular note is the GAMS collaboration (Alde, 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 apriori reason to expect 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 $t \equiv (p_{\text{Beam}} - p_X)^2 < 0$. (See Fig. 9.) The forward-peaking nature is seen in an approximately exponentially falling cross section with t, i.e. e^{bt} with $b \sim 3 - 8$ GeV^{-2}. For example, in charge exchange reactions at small values of $-t$, one pion exchange (OPE) dominates and is fairly well understood. It provides access only to states with $J^{PC} = \text{even}^{++}$ 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 undertood. 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 PWA is used to infer the exchange particle 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 (Chung 1975), the analysis is naturally divided into two sets of non-interfering 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 as unrestrictive triggers as possible to give large, uniform acceptance and choose particular final states in the analysis stage (Aston, 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 which decay to them, has come of age in recent experiments. GAMS made use of all-photon final states in particular, but fine-grained electromagnetic calorimeters have been combined with charged particle tracking and particle identification in E852 and in VES. This opens up a large number of final states that can be studied in a single experiment simultaneously during the same run. This can be very powerful by comparing decay branches of various states, as well as searching for decay modes that were not previously accessible.

3.2 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 $q\bar{q}$ 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, 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

$q\bar{q}$. 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 to produce hybrid mesons with exotic quantum numbers by means of flux tube excitation (Isgur, Kokoski, and Paton, 1985b; Afanasev and Page, 1998).

Peripheral photoproduction has further advantages. The vector dominance model allows non-OZI suppressed excitation of heavy quark states, such as $s\bar{s}$ and $c\bar{c}$, 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. 6.6.) A thorough review of the experimental situation through the mid-1970's is available in (Bauer 1978). The most significant contributions to meson spectroscopy since that time are by the LAMP2 experiment at Daresbury (Barber 1978, 1980) with photon beam energies up to 5 GeV, and the Ω-Photon collaboration at CERN (Aston, 1982) using energies between 20 and 70 GeV. Spectroscopy in exclusive photoproduction has also been carried out by E401 at Fermilab (Busenitz, 1989), also an electronic detector with a relatively open trigger; E687 at Fermilab (Frabetti, 1992), an evolution of E401 which concentrated on heavy quark physics; and by the SLAC hybrid bubble chamber and a laser-backscattered photon beam (Condo, 1993).

3.3 $\bar{p}p$ and $\bar{N}N$ Reactions

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

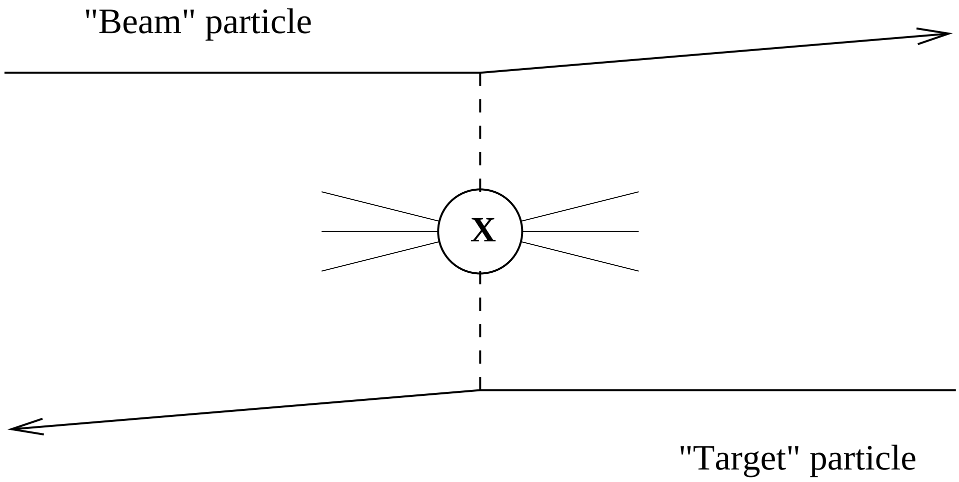

Figure 10: Schematic diagram of a central production process. Two off-mass-shell particles collide, while the “beam” and “target” move off essentially unscathed. The exchange particles are understood to be very rich in “glue” when the transverse momentum kick to the beam and target is small.

Antiproton annihilation in liquid hydrogen proceeds almost entirely through a $\bar{p}p$ 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 $\bar{p}p \to \pi^0\pi^0\pi^0$ and $\bar{p}p \to \eta\eta\pi^0$) has been particularly fruitful. In these reactions, one studies the two-body decay of a meson resonance which recoils off of the third meson. These data have had their greatest impact on the scalar meson sector.

3.4 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 where there is a collision between exchange particles. This is shown schematically in Fig. 10. Experimentally, using proton beam and target, one observes the reaction $pp \to p_f(X^0)p_s$ 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 that X^0 is a state dominated by gluonic degrees of freedom (Close, 1997b).

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

3.5 Results from e^+e^- Storage Rings

High luminosity ($\mathcal{L} \geq 10^{32}/\mathrm{cm}^2 \cdot \mathrm{sec}$) 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.

3.5.1 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.

The reaction $e^+e^- \to \pi^+\pi^-$ has been carefully studied in the region up to around 2 GeV center of mass energy (Barkov, 1985; Bisello, 1989). These data have clearly established three ρ resonances (Bisello, 1989) and have also been used to precisely study ρ/ω mixing. 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, 1994).

3.5.2 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 $e^+e^- \to e^+e^-X$ where the state X is produced by the collision of two photons radiated from the beam electron and positron. This is rather analogous to the central production process (Sec. 3.4, Fig. 10) where the photons replace the less well understood Pomeron. The field of "two-photon" physics has been rather extensively reviewed (Morgan, Pennington, and Whalley, 1994; Cooper, 1988). Data continues 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, and this 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 to *not* be produced in these reactions. This was discussed in Sec. 2.5.1 and quantified in Eq. 25. Thirdly, two-photon reactions are a powerful tool for spectroscopy in a way that is directly related to the way the scattered $e^{\pm}$ are detected.

Spectroscopic data from two-photon collisions is generally separated into "untagged" and "tagged" samples. The virtual photon spectrum is sharply peaked in the forward direction, since the photon propogator 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 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 $J^{PC} = 0^{\pm +}, 2^{\pm +}, 3^{++}, \cdots$ are produced in untagged events, and states with other quantum numbers should show a very strong dependence on q^2 for tagged events.

3.5.3 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 a few 10^{-4} and a few 10^{-3}. The best experiments acquire on the order of several million produced 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 (Hitlin and Toki, 1988; Königsman, 1986). Since then, however, new results have been presented from the DM2 collaboration at Orsay (Augustin, 1988, 1990) and the Mark III experiment at SPEAR (Dunwoodie, 1997; Burchell, 1991), and new data is being collected and analyzed by the BES collaboration at Beijing (Bai, 1996a, 1996b, 1998).

4 THE QUARK MODEL: COMPARISON WITH EXPERIMENT

4.1 Heavy Quarkonia

It is useful to start with heavy quarkonium where there is theoretical justification for using potential models to calculate their 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 grounds. In figure 3 we compared quark model predictions for the $b\bar{b}$ system to experiment and found the agreement to be good. In figure 11 we show a similar comparison for the $c\bar{c}$ spectrum, also with good agreement. The

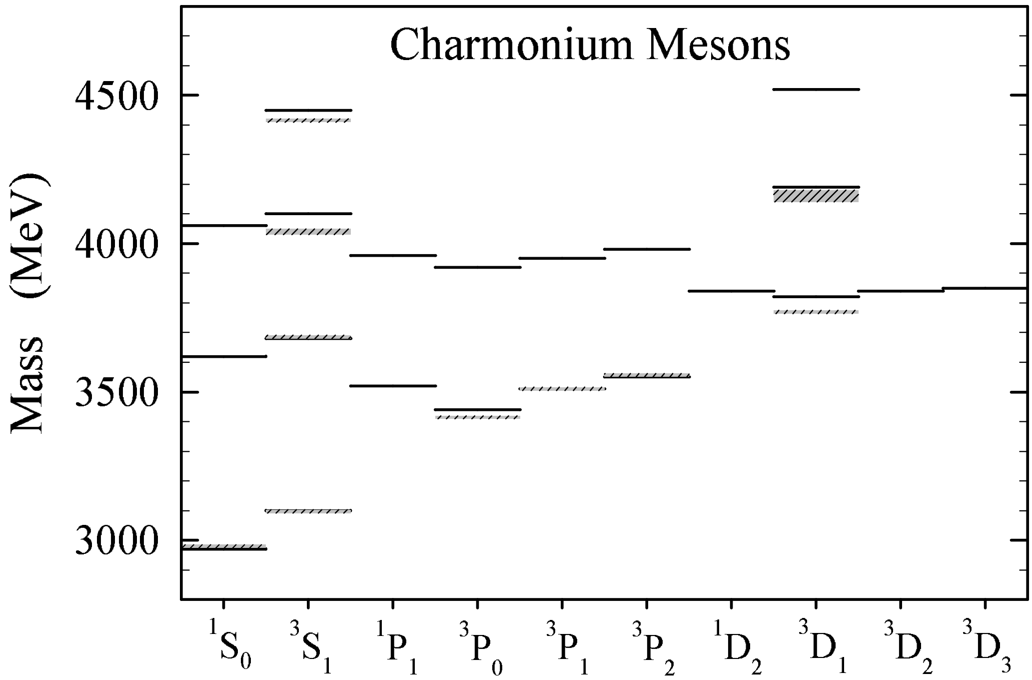

Figure 11: The $c\bar{c}$ spectrum. The solid lines are the quark model predictions and the shaded regions are the experimental states. The size of the shaded regions approximates the experimental uncertainty. Note that the interpretation of the $\psi(4040)$ and $\psi(4160)$ as a single resonance is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

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. There have been suggestions that D-wave quarkonium can be produced via gluon fragmentation at hadron colliders (Qiao 1997a), in Z^0 decays (Qiao 1997b), in B decays (Yuan 1997, Ko 1997) and in fixed target experiments (Yuan 1998). There is some evidence for the D-wave 2^{--} charmonium state with Mass 3.836 ± 0.013 GeV in Fermilab experiment E705 (Antoniazzi 1994) although this result is questioned by other experiments (Gribushin 1996). It is also possible that with the high statistics available at a B-factory that the $b\bar{b}$ D-waves might be observed via cascade radiative transitions from $\Upsilon(3S) \to 2P\gamma \to 1D\gamma\gamma$ (Kwong 1988). For both the $b\bar{b}$ and $c\bar{c}$ spectra, the states which 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 1975, 1978, 1980). Note that there are classification ambiguities in a few cases for the high mass 1^{--} resonances.

4.2 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 section 2 that the basic structure in heavy and light systems are 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 the 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 colliding e^+e^- machines that the experimentaly accessable excitations are nearly orthogonal.

In our survey of mesons with light quarks we begin with a general survey of these states to establish the global validity of the quark model predictions. With this backdrop, in section 5 we will focus on, and study in detail, the states which are either poorly understood or pose a problem for the quark model.

4.2.1 Mesons With One Light Quark and One Heavy Quark

We start with mesons which contain one heavy quark such as the charmed and beauty mesons (Rosner 1986, Godfrey 1991). These systems are an interesting starting point because, as pointed out long ago by De Rujula, Georgi, and Glashow (De Rujula 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. As such, 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 studying the more complicated light quark sector in search of 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 $c\bar{u}$, $c\bar{s}$, $b\bar{u}$, $b\bar{s}$, and $b\bar{c}$ states are summarized in fig. 12. See also Kwong and Rosner (Kwong 1991) and Eichten and Quigg (Eichten 1994).

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 which has an important effect on the meson decay properties. The OZI allowed decays for P-wave mesons can be described by two independent amplitudes, S-wave and D-wave. One state is degenerate with the ${}^3P_0$ state and the other is degenerate with the ${}^3P_2$ state. Furthermore, the 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 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 P-wave mesons form two degenerate doublets.

These patterns in spectroscopy and decays can be extended to general principles and this 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 1989b; 1990; 1991, Voloshin 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 $\vec{j}_l = \vec{S}_q + \vec{L}$ are separately conserved so that each energy level in the excitation spectrum is composed of degenerate pairs of states $\vec{J} = \vec{j}_q + \vec{S}_Q = \vec{j}_q \pm 1/2$. This represents a new symmetry in the QCD spectrum in the heavy quark limit and 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 flavour symmetry also applies to the complete set of n-point functions of the theory including all strong decay amplitudes arising from the emission of light quanta like π, η, ρ, $\pi\pi$, etc., are independent of heavy quark flavour so that two states with spins $s_\pm$ must have the same total widths. The heavy quark symmetry leads to a number of predictions. The two D_1's both have $J^P = 1^+$ and are only distinguished by J_l 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 the decays from both members of a doublet with given J_ℓ to the members of another doublet with J_ℓ^P are all essentially a single process. This leads to the simple prediction that the two excited states should have exactly the same widths. $D_0^* \to D\pi$ and $D_1 \to D^*\pi$ decay via S-wave and can be quite broad and therefore difficult to identify experimentally. In contrast $D_1 \to D^*\pi$ and $D_2^* \to D^*\pi,\ D\pi$ proceed via D-wave and

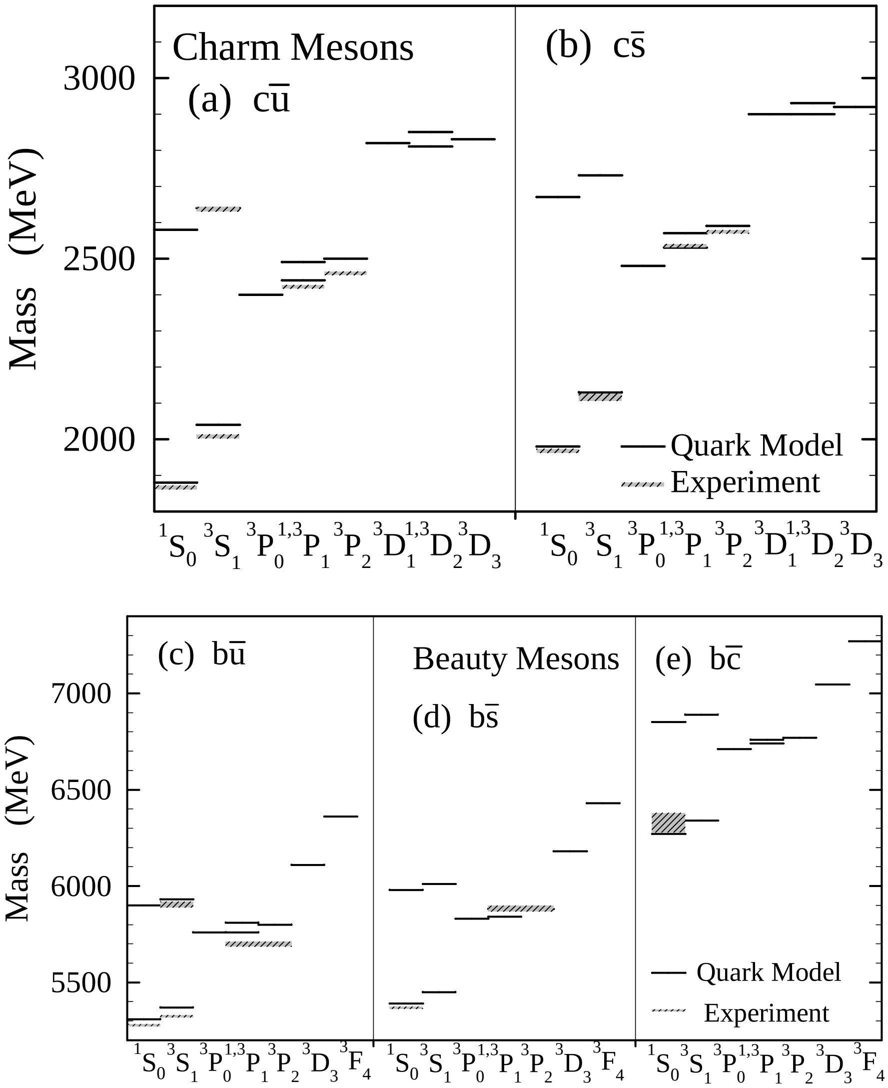

Figure 12: Spectra for (a) $c\bar{u}$, (b) $c\bar{s}$, (c) $b\bar{u}$, (d) $b\bar{s}$, and (e) $b\bar{c}$. The solid lines are the quark model predictions (Godfrey 1985a) and the shaded regions are the experimental measurements with the size representing the approximate experimental uncertainty. The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

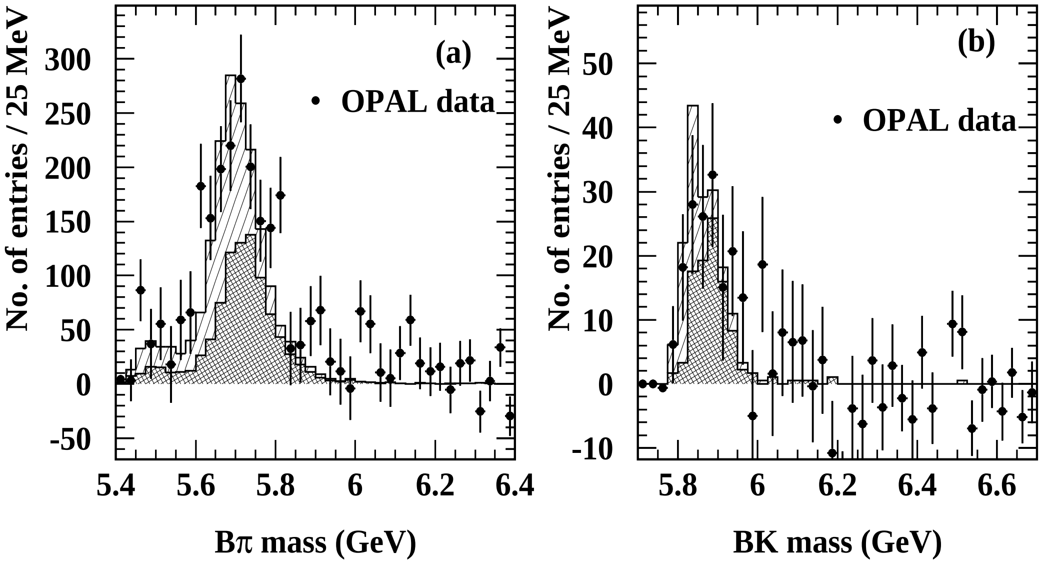

Figure 13: The $B\pi$ and BK invariant mass distributions. The solid histograms show Monte Carlo results for the B_2^* and B_{s2}^* states respectively and the hatched histograms show Monte Carlo results for B_1 and B_{s1} (Akers 1995).

are much narrower. These states are identified as the $D_1(2420)$ and $D_2(2460)$. These are the same conclusions obtained from the quark model. There are numerous other predictions resulting from HQET that are relevant to weak decays and we encourage the interested reader to read one of the more specialized reviews on the important subject of the Heavy Quark Effective Theory (Neubert 1994).

While the P-wave charmed mesons have been known for some time, the OPAL (Akers 1995), ALEPH (Buskulic 1996), and DELPHI (Abreu 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 $B\pi$ ($M = 5.68 \pm 0.011$ GeV, $\Gamma = 116 \pm 24$ MeV) and BK ($M = 5.853 \pm 0.15$ GeV, $\Gamma = 47 \pm 22$ 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 $m_{B_c} = 6.33 \pm 0.05$ GeV (Ackerstaff 1997, Abreu 1997). More recently CDF has also reported an observation of B_c (Singh 1998).

The DELPHI collaboration (Abreu 1998, Ehret 1998) has reported evidence for radial excitations of the D_r^* and B_r^*. From the invariant mass distribution $M(D^*\pi\pi)$ DELPHI obtains the mass measurement of $M(D_r^*) = 2637 \pm 2 \pm 6$ MeV which is in good agreement with the quark model prediction (Godfrey 1985a). However, this state has not been confirmed by the OPAL collaboration (Krieger 1998) or the CLEO collaboration (Shipsey 1998). DELPHI also reports evidence for the $B^{\star\prime}$ from $Q(B^{(*)}\pi^+\pi^-)$ with $M_{B_r^*} = 5904 \pm 4 \pm 10$ which is also in good agreement with the quark model. Although these results are preliminary they are promising and together with the B^{**} observations show the potential of high energy colliders for contributing to our understanding of hadron spectroscopy.

4.2.2 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 (Aston 1986a, 1986b, 1987) are a good place to start 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 flavour and therefore don't 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. Much of the data comes from the Large Aperture Superconducting Solenoid collaboration

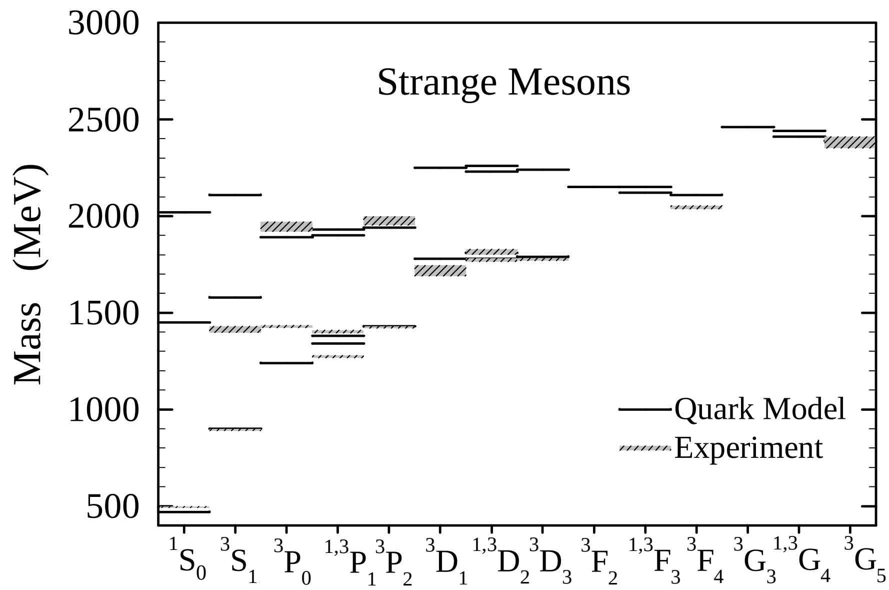

Figure 14: Level diagram for the strange mesons. The quark model predictions (Godfrey 1985a) are given by the solid lines and the experimental measurements are given by the shaded regions with the size representing the approximate experimental uncertainty.

(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 $J^P = 5^-$.

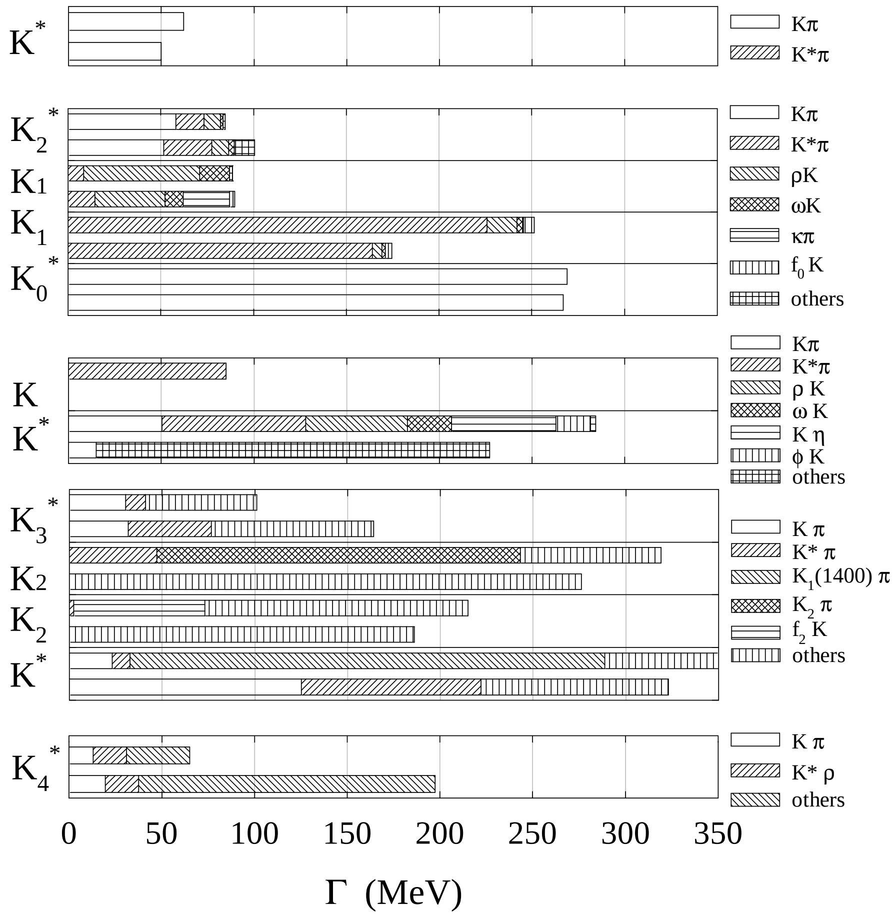

Figure 15: Strange meson strong decays. For each state the quark model predictions are given by the upper bar (from Kokoski and Isgur (1987)) and the experimental results are given by the lower bar (from the Review of Particle Physics (Caso 1998) unless otherwise noted). Some of the more important decay 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*.

A substantial number of the expected underlying states are also observed in $K^-\pi^+$, $\bar{K}^0_s\pi^+\pi^-$ 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 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 1987) and the 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 is good for the masses of the leading orbital excitations between the quark model and experiment supporting the picture of linear confinement.

Partial waves of the reaction $K^-p \to \bar{K}^0\pi^+\pi^-n$ 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 2 1^- Breit-Wigner resonances; a higher state with $M = 1717 \pm 27$ MeV and $\Gamma = 322 \pm 110$ couples to both channels while the lower mass state with $M = 1414 \pm 15$ and $\Gamma = 232 \pm 21$ is nearly decoupled from the $K\rho$ channel (Aston 1987). It is simplest to associate the higher state with the 1^3D_1 state based on the small triplet splitting and 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 which 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 2 structures in the S-wave (Aston 1988e). The first, with mass around[4] $M(K_0^*(1430)) = 1412 \pm 6$ MeV and $\Gamma(K_0^*(1430)) = 294 \pm 23$ MeV is classified as the $^3P_0$ partner of the $K_2^*(1430)$. Although this is ~ 170 MeV higher than the quark model prediction the 1^3P_0 state is particularly sensitive to the details of the model so not too much should be read into this. There is a second S-wave structure at around 1.9 GeV with parameters $M = 1945 \pm 22$ MeV and $\Gamma \sim 201 \pm 86$ 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 $M = 1973 \pm 26$ MeV and $\Gamma = 373 \pm 68$ MeV, most probably the 2^3P_2.

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

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:

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

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

where θ_P is the SU(3) singlet-octet 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 by the LASS collaboration in the reaction $K^-p \to K^-\eta p$ who found (Aston, 1988b)

$$BR(K_3^* \to K\eta) = 9.4 \pm 3.4\%$$

$$BR(K_2^* \to K\eta) < 0.45\% \ (95\% \ C.L.)$$

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 $K\eta'$ are expected to be suppressed. The relative phases of the different decay amplitudes also agree with the quark model predictions.

4.2.3 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 provides 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, much 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 $K^-p \to K^-K^+\Lambda$, $K^-p \to K_s^0K^\pm\pi^\mp\Lambda$, and $K^-p \to K_s^0K_s^0\Lambda$, which strongly favour 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 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

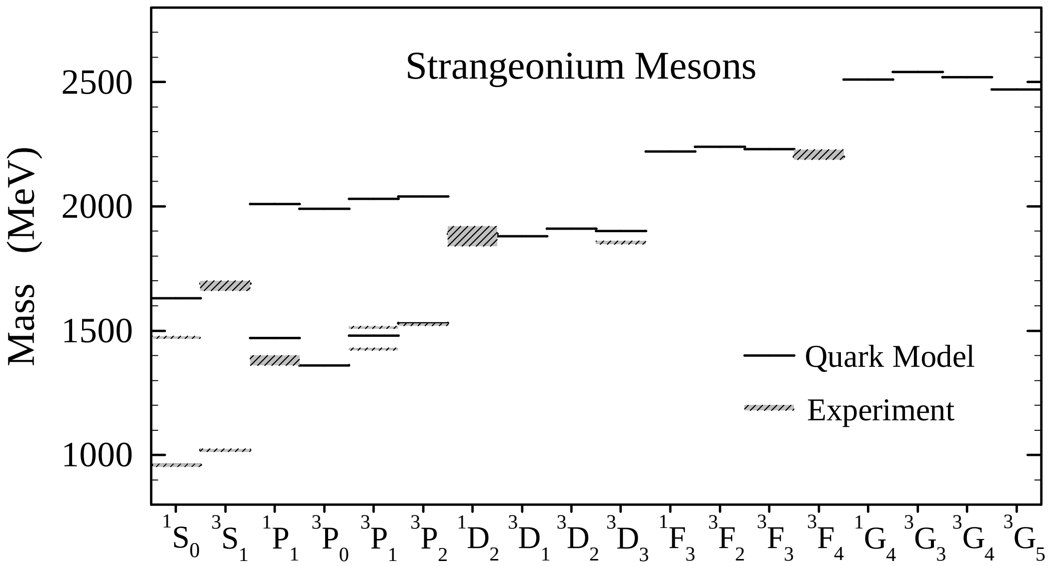

Figure 16: Level diagram for the strangeonium mesons. Note that there are 2 candidate 1^3P_1 states, the $\eta(1470)$ is not unambiguously identified as the 2^1S_0 state, and the $f_4(2220)$ is an unconfirmed report by the LASS collaboration (Aston 1988d). See Fig. 14 for further details.

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 Section 5.

An amplitude analysis of the reactions $K^-p \to K^0_S K^0_S \Lambda_{seen}$ and $K^-p \to K^-K^+\Lambda_{seen}$ indicates S-wave structure around the $f_2'(1525)$ mass, Fig. 18, (Aston, 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 that predicted which is consistent with the strange meson sector. This implies that the the $f_0(975)$, normally 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 ~ 1500 MeV but with different properties than the state possibly observed by LASS in $K^-p \to K\bar{K}\Lambda_{seen}$. We conclude that they are different states with the latter being discussed in detail in the following section.

There are two candidates for the 1^{++} $s\bar{s}$ state with masses ~ 1510 MeV and ~ 1420 MeV and disentangling them is tied up with the E/ι region to be discussed in Section 5.4. 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 Fig. 19(a) and (b). (Aston, 1988c) These distributions are well described by Breit-Wigner curves as shown, and, assuming I=0, represents 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

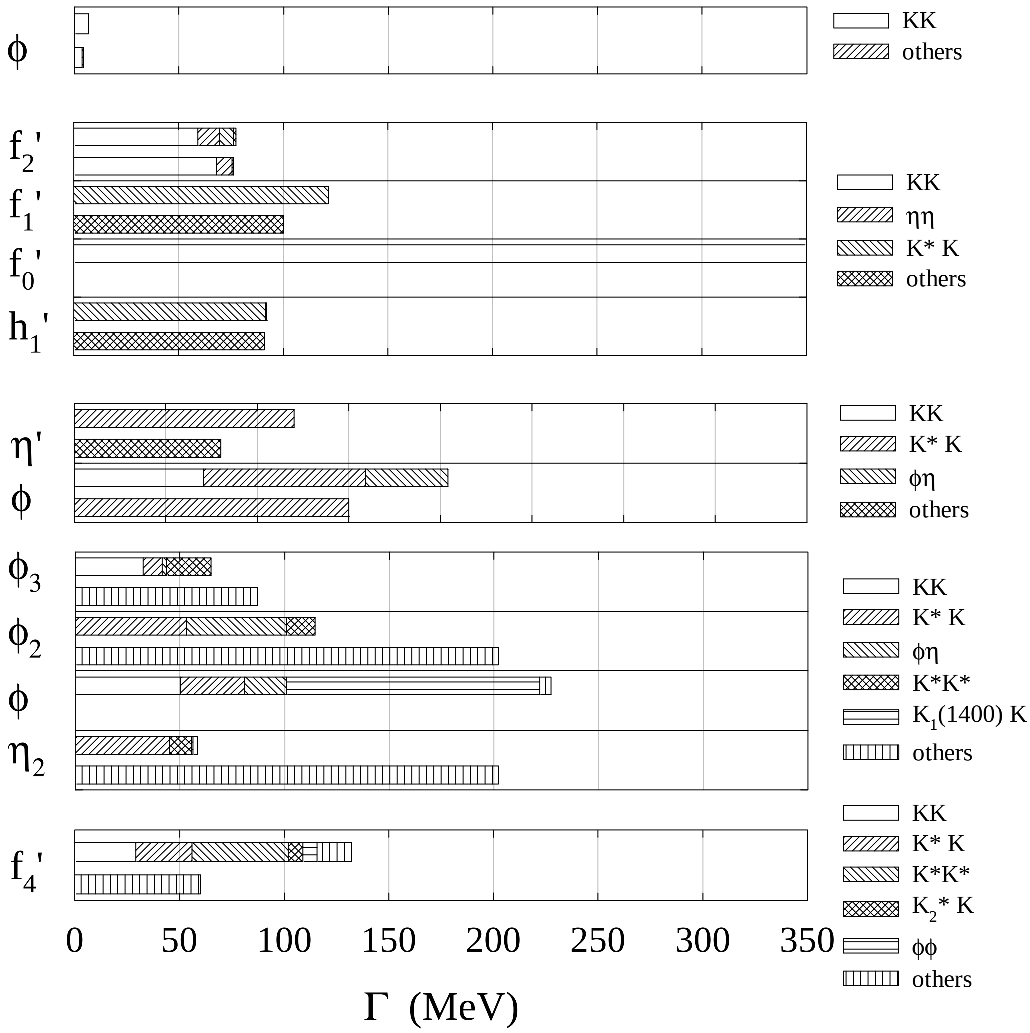

Figure 17: Strangeonium meson strong decays. For each state the the quark model predictions are given by the upper bar (Kokoski 1987, Blundell 1996, Barnes 1997) and the experimental results are given by the lower bar (from the Review of Particle Physics, Caso 1998). See Fig. 15 for further details.

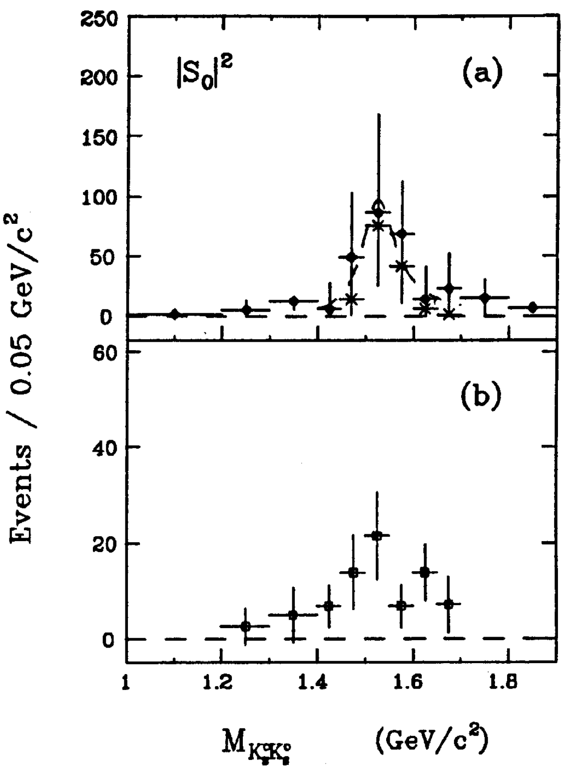

Figure 18: S-wave intensity distribution for (a) the reaction $K^-p \rightarrow K^0_S K^0_S \Lambda_{seen}$ and (b) $K^-p \rightarrow K^-K^+\Lambda_{seen}$. (From Aston 1988a.)

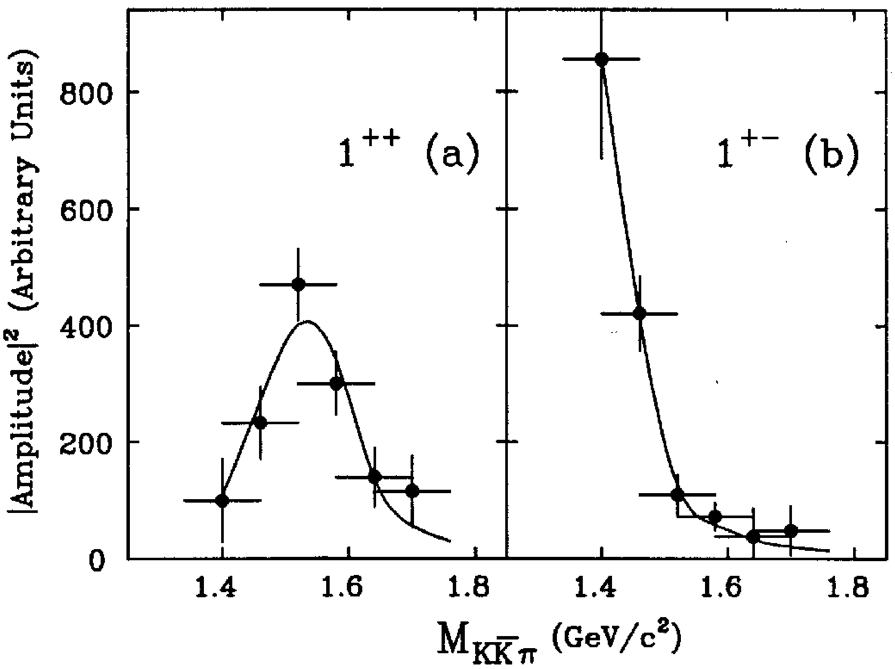

Figure 19: The mass dependence of the total production amplitude intensities for the axial vector G-parity eigenstates. The curves correspond to the Breit-Wigner fits described in the text. (From Aston 1988c.)

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 E. King *et al.* (King, 1991) in the partial-wave anaysis 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 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.*. (Gavillet, 1982) although there is no evidence for this state by Fermilab E690 who see the $f_1(1420)$ (Berisso 1998). These results, taken together with the $f_0(1530)$, indicate that the tensor and spin-orbit mass splittings appear to be 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 PDG (Caso, 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 $K^0_s\bar{K}^0_s$, $K\bar{K}$, and $\eta\eta$ pairs in J/ψ radiative decay. (Baltrusaitis, 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 section 5.

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 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, 1988d). This would be a member of the 4^{++} ($f_4(2030)$, $a_4(2040)$ (Cleland, 1982), and $K_4^*(2060)$) (Aston, 1988e) nonet 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 $\pi^-p \to \eta\eta'$ (Alde 1986). This implies that the $f_J(2220)$, which has been conjectured to be an exotic hadron of some sort (Chanowitz 1983a), is instead the $s\bar{s}$ member of the quark model $^3F_4$ gound state nonet (Godfrey 1984, Blundell 1996). On the other hand the BES collaboration (Bai, 1996a) finds a decay width too narrow to be easily accomodated as the 1^3F_4 $s\bar{s}$ state. More importantly they find that the decays are approximately flavour symmetric which supports a glueball interpretation. One intrepretation which could accomodate 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 which is not observed in hadronic reactions.

To summarize, the $s\bar{s}$ spectrum agrees remarkably 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 contain intriguing hints of non-quark model physics. These puzzles will be explored in section 5.

4.2.4 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 which contribute 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

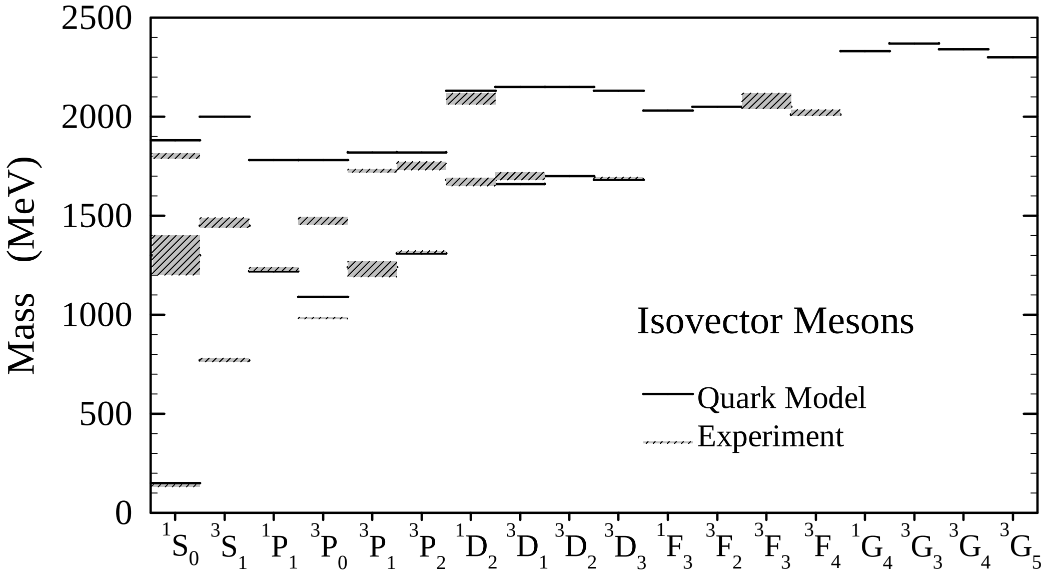

Figure 20: Level diagram for the isovector mesons. See Fig. 14 for further details.

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 which draw attention to themselves and require further discussion.

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 $\sim$ 50 MeV that is much more difficult to reconcile with the quark model prediction of $\sim$ 500 MeV. This large discrepancy has led to numerous conjectures that the $a_0(980)$ is not a $q\bar{q}$ state but rather, a more complicated 4-quark object (Jaffe, 1977a, 1977b, 1978), most probably a $K\bar{K}$ molecule (Weinstein and Isgur, 1982, 1983). The observation of the $a_0(1450)$ in $p\bar{p}$ annhilation by the Crystal Barrel Collaboration (Amsler 1994a, 1995) and the OBELIX Collaboration (Vecci 1998) is naturally assigned to be the isovector member of 1^3P_0 nonet reinforcing the interpretation that the $a_0(980)$ is a non-$q\bar{q}$ state. We will 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 (Andreazza, 1998) collaborations observed a signal in

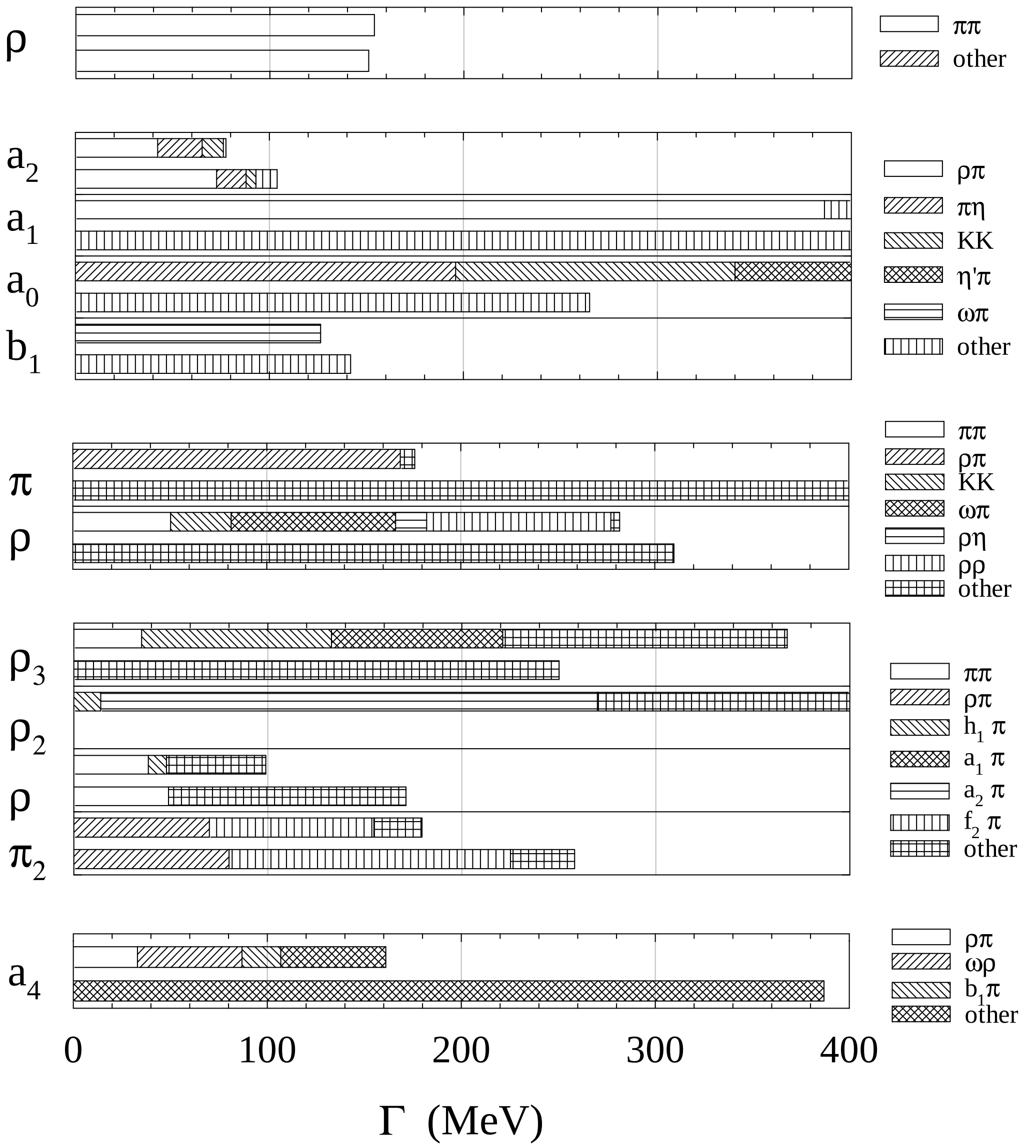

Figure 21: Isovector meson strong decays. See Fig. 15 for further details.

$\tau \to a_1' \nu_\tau \to \pi^- \pi^+ \pi^- \nu_\tau$ 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 1994), BNL E852 (Ostrovidov 1998) and VES (Amelin, 1995). There is also evidence for the $J = 2$ partner of this state, the a_2', by the L3 collaboration which observed it 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. BNL E852 also reports observing it in $\pi^- p \to \eta\pi^- p$ (Ostrovidov 1998) while the Crystal Barrel Collaboration reports an observation of an a_2 in $p\bar{p}$ 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 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 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 (Godfrey and Isgur, 1985a) surmised that the observed state was a mixture of two broad overlaping resonances, the 2^3S_1 and the 1^3D_1 which gave rise to the observed properties. Subsequently, Donnachie and Clegg (Donnachie 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 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 $\bar{p}p$ annihilation by the Crystal Barrel Collaboration (Abele 1997c) and by the OBELIX Collaboration (Bertin 1997b, 1998). The properties of these two states are in reasonable agreement and more recent results examining the 4π decays of these states support the quark model assignments (Thoma 1998). There have been additional claims of a $\rho(1250)$ seen in $\omega\pi$ but not in $\pi\pi$. Although this state has been claimed many times it has never been generally accepted and there are theoretical difficulties reconciling it with e^+e^- data so at present its status is unclear.

4.2.5 The Non-strange Isoscalar Mesons

The non-strange isoscalar mesons are the last of the light quark meson families and are also the most problematic. Experimentally, the isoscalars have been difficult to detect since neutral mesons tend to be more difficult to detect than their charged partners. Phenomenologically, there is the problem of $n\bar{n}$ and $s\bar{s}$ 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 which may point to the

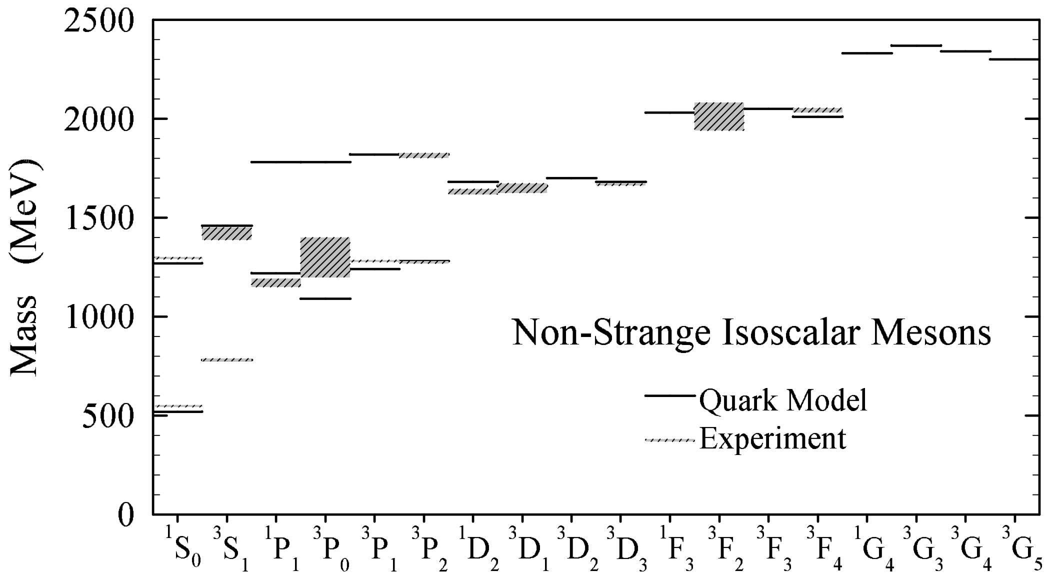

Figure 22: The non-strange isoscalar meson spectrum. See Fig. 14 for further details.

existence of "exotic" hadrons lying outside the constituent quark model. We will defer their discussion until the following section on puzzles and problems.

As in the case of the $s\bar{s}$ 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 $p\bar{p} \to \eta 3\pi^0$ (Amsler 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 $I = 0$ $^1D_2$ $n\bar{n}$ state which should appear near the 1D multiplet mass of ~ 1670 MeV with $\Gamma \sim 260$ MeV with $a_2\pi$ the dominant decay mode. The 2nd state, $\eta_2(1875)$ is too high in mass to be a 1D $n\bar{n}$ state and the strong $f_2\eta$ mode argues against a mainly $s\bar{s}$ 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, 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 (Close 1995a) predict to be dominant.

The excited pseudoscalar mesons, in particular the E/ι mass region, remains a puzzle.

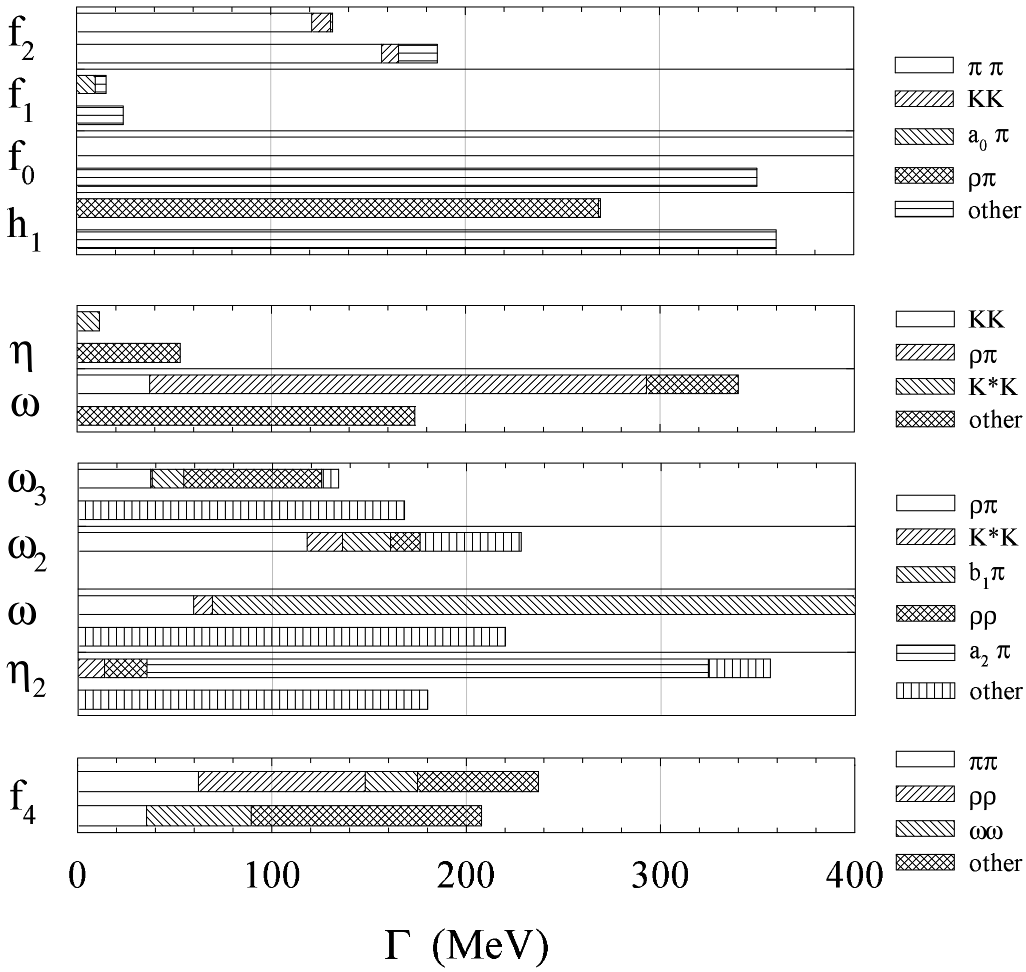

Figure 23: The non-strange isoscalar meson decays. See Fig. 15 for further details.

(See Sec. 5.4.) 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 (PDG, Caso 1998). They are referred to as $\eta(1410)$ and $\eta(1490)$. One could identify the $\eta(1275)$ as the radial excitation of the η and the $\eta(1490)$ as the mainly $s\bar{s}$ 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 that 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 isocalars 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 $s\bar{s}$ 2^{++} mesons. In particular, the $f_J(1710)$ has been mentioned previously in that subsection. 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 fit in with the quark model prodictions. 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. 5.5.1.

4.2.6 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 non-strange isoscalar sectors a high statistics would be useful in finding higher mass states.

Given the generally good agreement between the quark model and experiment it is the discrepancies which 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 shows 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 $f'_1(1530)$ so that the $f_1(1420)$ clearly appears as an extra state. Similar observations apply to the $f_0(1500)$, the $f_J(1710)$, and the $f_2(2010)$, $f_2(2300)$ and the $f_2(2350)$. These states point to a need for a better understanding of hadronic structure, perhaps 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 which do not fit in the $q\bar{q}$ spectroscopy at all and have not yet been discussed: for example, the state with exotic J^{PC} quantum numbers, 1^{-+} $\hat{\rho}(1405)$ seen in $\pi^- p \to \pi^0 \eta$ and structures in $\gamma\gamma \to VV$.

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

5 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 when compared to the inherent complexity of QCD.

However, there is another class of disagreements which cannot be so ascribed. These cases are more profound, and presumably point to the features of QCD *not* contained in the quark model, and may indeed point to the fundamental degrees of freedom needed to fully describe hadron structure. These may help us identify the necessary features which 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 aspects. The most unequivocal of these is the establishment of states with "exotic" quantum numbers, that is, J^{PC} which cannot be accomodated with only $q\bar{q}$ degrees of freedom. The second aspect 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 doesn't 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 where this disagreement is indeed quite profound.

5.1 Exotic Quantum Numbers

As discussed in Sec. 2 (see Table 1), the quantum numbers of a $q\bar{q}$ system must have either $P = C$ for all J except $J = 0$ 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, mainly because their branching ratio to conventional final states are small. This means that one needs 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, 1985b; Isgur, and Paton, 1985a; 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 would 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 experimentally, since they involve a large number of final state particles which can arise from different quasi-twobody states. For example, $\pi^- f_1(1285)$ leads to $\pi^-\pi^+\pi^-\eta$ which might also be due to such things as $\pi^-\eta(1295)$, ηa_2^-, ηa_1^-, $\rho^0 a_2^-$, and so forth. Unless the particular reaction enhances *production* of hybrid exotics (Isgur, Kokoski, and Paton, 1985b) it will likely be very hard to clearly identify the exotic state amidst the debris of highly excited, conventional $q\bar{q}$ mesons (Barnes, Close, Page, and Swanson, 1997).

Experimentally, however, it is reasonable to first look at less complex final states, and this is how the field has progressed so far. For example, $\eta\pi$ and $\eta'\pi$ final states are attractive since any resonant odd-L partial wave would be manifestly exotic. Alternatively, states with relatively few particles, all charged, are also attractive since photon detection (which costs money and leads to poorer resolution and complex acceptance) is not needed and the PWA 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.

5.1.1 $\eta\pi$ 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.[5]

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 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.

Experiments searching for exotic resonance structure in $\eta\pi$ systems are listed in table 4. 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 which dominates the mass spectrum, although details of the asymmetry differ from experiment to experiment. Each experiment except the first one listed (Apel, 1981) suggest or claim outright evidence for an exotic $J^{PC} = 1^{-+}$ resonance, although the only two consistent results are from VES (Beladidze, 1993) and E852 (Thompson, 1997). Finally, a very recent result from the Crystal Barrel at CERN (Abele, 1998a) is 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 $\hat{\rho}(1400)$.

[5] 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 non-exotic meson could decay to $\eta\pi^{\pm}$ if isospin symmetry is violated, but this should be a negligible effect for the cases discussed here.

Table 4: 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

Both the NICE and GAMS experiments (see table 4) 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 are 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 potential 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. This 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, 1988a) and this generated much interest in its theoretical interpretation (Close and Lipkin, 1987; Iddir, 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, 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. 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, 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

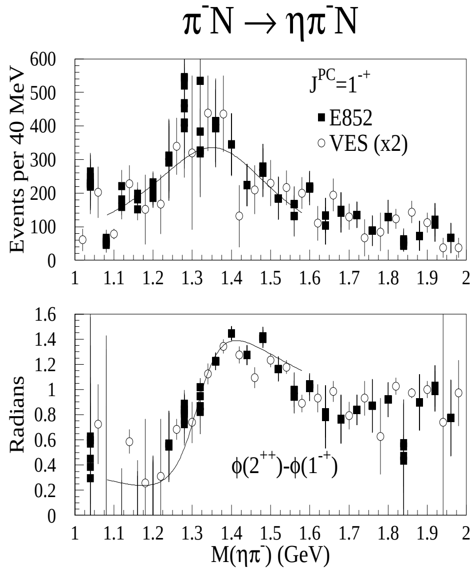

Figure 24: A comparison of the exotic $J^{PC} = 1^{-+}$ signal observed by the E852 (Thompson, et al., 1997) and VES (Beladidze, et al., 1993) experiments. The E852 data explicitly shows the ambiguous solutions of the partial wave analysis. The VES intensity distribution is multiplied by a factor of two, 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 $\hat{\rho}(1400)$.

D_+ waves did not vary across the mass region of interest. These 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, 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.

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, 1997). Statistics were achieved that were comparable to 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 determine the mass and width of the $\hat{\rho}$ to be approximately 1370 ± 50 MeV/c^2 and 385 ± 100 MeV/c^2 respectively. These fits are shown in Fig. 24.

The recent Crystal Barrel result (Abele, 1998a) comes from $\bar{p}$ capture in liquid deuterium, the reaction being $\bar{p}d \to \eta\pi^-\pi^0 p$ 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 $\rho^- \to \pi^-\pi^0$, 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. They determine the mass and width of the $\hat{\rho}$ to be approximately 1400 ± 30 MeV/c^2 and 310 ± 70 MeV/c^2 respectively, quite consistent with the E852 result.

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, 1988a, 1998b; 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.

5.1.2 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, 1998), decaying in this case to $\rho\pi$, in the reaction $\pi^- p \to \pi^-\pi^+\pi^- p$ 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, 1993), who also see evidence for a broad but resonant P_+ wave near this mass in the $\eta'\pi^-$ system (Beladidze, 1993). Although the $\hat{\rho}(1600)$ data does not agree particularly well with recent lattice QCD results and phenomenological expectations for a hybrid meson it appears that it can still be accomodated 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" has been checked explicitly. A data set was simulated, using the waves used 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 is 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 is clean.

Simultaneous with their measurement of the $\eta\pi^-$ final state, the VES collaboration (Beladidze, 1993) also studied the $\eta'\pi^-$ system in the reaction $\pi^- N \to \eta'\pi^- N$ with $\eta' \to \eta\pi^+\pi^-$. 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 PWA 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 parameterization yields the correct values for the mass and width. (This analysis also yields the relative branching ratio for $a_2 \to \eta'\pi$ to $a_2 \to \eta\pi$ to be $\sim 5\%$.) 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, 1998).

The E852 data indicates the $BR(\hat{\rho}(1600) \to \rho\pi) \approx 20\%$ (Page, 1997c), while there is no indication of this state in the $\eta\pi^-$ data (Beladidze, 1993; Thompson, 1997), and clearly not all of the $\eta'\pi^-$ signal (Beladidze, 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.

5.1.3 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, 1985b) 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, 1983) on the reaction $\gamma p \to \omega\pi^+\pi^- p$. This group showed that production of $b_1(1235) \to \omega\pi$ was enhanced when $1.6 \leq \text{Mass}(\omega\pi\pi) \leq 2.0$ GeV/c^2. Although the number of events was limited, their reanalysis (Atkinson, 1987) showed what appeared to be production of $\omega_3(1670)$ and a new state at ~ 1.9 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 demonstrated.

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

At this time, results have only been presented for $f_1(1285)\pi$ final state. These are from the VES collaboration (Gouz, 1993), who made use of the $f_1 \to \eta\pi^+\pi^-$ final state, and BNL Experiment E818 (Lee, 1994) for which $f_1 \to K^+\bar{K}_0\pi^-$. In both cases, there are not very many events, and the VES result remains preliminary. However, 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 in (Lee, 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, 1987). One of these, experiment E687 at Fermilab, performed an inclusive search for $f_1(1285)\pi^\pm$ states (Danyo, 1995) in the reaction $\gamma\text{Be} \to f_1(1285)\pi^\pm X$ with $f_1(1285) \to a_0(980)^\pm\pi^\mp$ and $a_0(980)^\pm \to K_S K^\pm$. A tagged bremsstrahlung beamline provided photons with energies near 200 GeV. A peak in the $f_1\pi$ mass distribution is clear above background. The background 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 statistics are quite poor.

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

5.2 The scalar mesons

The scalar ($J^P = 0^+$) mesons have long been a source of controversy (Morgan, 1974). As discussed in Sections 4.2.3, 4.2.4, and 4.2.5, 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 Section 4 and a brief techical review is given by the Particle Data Group (Caso, 1998).

We will 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 Section 5.3.1.

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 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 multi-quark states (Jaffe, 1977a; Weinstein and Isgur, 1983; Rosner, 1983). Still, there is significant controversy (Janssen, 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, 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, 1997a) and to $K\bar{K}$ (Abele, 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).

5.2.1 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 Section 2.5.1 as well as Szczepaniak, 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 $f_0(1500)$ is from the Crystal Barrel collaboration, who resolved two scalar states in this mass region (Anisovich, 1994), and who determine its decay branches to a number of final states, including $\pi^0\pi^0$ and $\eta\eta$ (Amsler, 1995; Abele, 1996c), $\eta\eta'$ (Amsler, 1994b), K_LK_L (Abele, 1996b) and $4\pi^0$ (Abele, 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, 1998) in the $\bar{n}p \to \pi^+\pi^+\pi^-$ reaction for neutrons in flight, and also in "glue rich" central production reactions $pp \to p_f(\pi^+\pi^-\pi^+\pi^-)p_s$ and $pp \to p_f(\pi^+\pi^-)p_s$, most recently by the WA91 (Antinori, 1995) and WA102 (Barberis, 1997b) collaborations. There is some evidence of $J/\psi \to \gamma f_0(1500)$ with $f_0(1500) \to \pi^+\pi^-\pi^+\pi^-$ (Bugg, 1995) and perhaps also with $f_0(1500) \to \pi^+\pi^-$ (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, 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; Amsler and Close, 1996; Close, Farrar, and Li, 1997c; Close, 1997a; Amsler, 1998). No evidence has been reported in $\gamma\gamma$ collisions (as would be expected to be the case if the $f_0(1500)$ is 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, 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 more narrow, with $\Gamma[a_0(1450)] = 265 \pm 13$ MeV/c^2 and $\Gamma[K_0^\star(1450)] = 287 \pm 20$ 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, hints to its special status as something different that 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 mass typical of 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. This information is 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 it is not only broad, but it also interferes strongly with the even broader underlying S-wave $\pi\pi$ structure as well as with the $f_0(980)$ (Caso, 1998; Amsler, 1998). It is nevertheless clear, however, 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, 1995) 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 $f_0(1500)$:

$$\begin{aligned} B[\bar{p}p \to f_0(1500), f_0(1500) \to \pi^0\pi^0] &= (12.7 \pm 3.3) \times 10^{-4} \\ B[\bar{p}p \to f_0(1500), f_0(1500) \to \eta\eta] &= (6.0 \pm 1.7) \times 10^{-4} \end{aligned}$$

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

$$B[\bar{p}p \to f_0(1500), f_0(1500) \to \eta\eta'] = (1.6 \pm 0.4) \times 10^{-4}$$

They also determine (Amsler, 1998; Table 11)

$$B[\bar{p}p \to f_0(1500), f_0(1500) \to K_L K_L] = (1.13 \pm 0.09) \times 10^{-4}$$

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

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:

$$\pi\pi : KK : \eta\eta : \eta\eta' = (5.1 \pm 2.0) : (0.71 \pm 0.21) : (\equiv 1.0) : (1.3 \pm 0.5) \tag{30}$$

where the value for $\pi\pi$ (KK) multiplies the branching ratio for $\pi^0\pi^0$ ($K_L K_L$) by 3 (4) to account for charge combinations. These are obviously *inconsistent* with the $f_0(1500)$ 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 $f_0(1500)$ being the $n\bar{n}$ member, although we've already argued that the $f_0(1370)$ is a better candidate.

Therefore, the circumstantial evidence for $f_0(1500)$ 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 Section 2.5.1 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 $f_0'(\sim 1600)$. 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 $f_0'(\sim 1600)$. 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 f_0' and to measure its decay properties in order to clearly establish the $f_0(1500)$ as the scalar glueball. See Section 4.2.3. There are two candidates at present. One possibility is the tentative identification (Aston, 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. 5.3.1). Although its spin assignment is somewhat uncertain and controversial, it shows some characteristics of being a glue-dominated state itself.

5.2.2 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, 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 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, 1998). For example, the states are somewhere between 50 and 100 MeV/c^2 wide, and 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, 1998) 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, and 1001.3 ± 1.9 MeV/c^2 and 70 ± 5 MeV/c^2 based on the coupled-channel description developed by Flatté (Flatté, 1976).

The $\gamma\gamma$ decay widths (Caso, 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, 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 $BR(\phi \to a_0\gamma) \approx BR(\phi \to f_0\gamma) \approx 4 \times 10^{-5}$ 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. Very recent results from the SND collaboration running at the VEPP-2M storage ring in Novosibirsk (Achasov, 1997a; Aulchenko, 1998) yield a value $BR(\phi \to f_0\gamma) = (3.42 \pm 0.30 \pm 0.36) \times 10^{-4}$, much larger than expected for a $K\bar{K}$ molecule. They also report that $BR(\phi \to \eta\pi^0\gamma) \approx (1.3 \pm 0.5) \times 10^{-4}$ but with no suggestion of a peak at the $a_0(980)$. 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.

5.3 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 lead 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. 3.5.3, the annihilation of a $c\bar{c}$ pair into a photon and two gluons is certainly expected in lowest order QCD, so one is naturally lead to search for gluonic states in these decays. This was discussed in some detail recently by (Close, Farrar, and Li, 1997c; Page and Li, 1998). Further information is provided by the coupling, or its upper limit, of the candidate state to $\gamma\gamma$ (Sec. 3.5.2). Pure glue can only couple to photons through the creation of an intermidiate $q\bar{q}$ pair and is therefore suppressed relative to quark model states. On the other hand, central production in pp collisions (Sec. 3.4) may be a rich source of gluonic states, based on speculation that it proceeds through double Pomeron exchange (Close, 1997a, Close and Kirk, 1997b).

These processes are shown in Fig. 25 for excited mesons decaying to K^+K^-. The figure

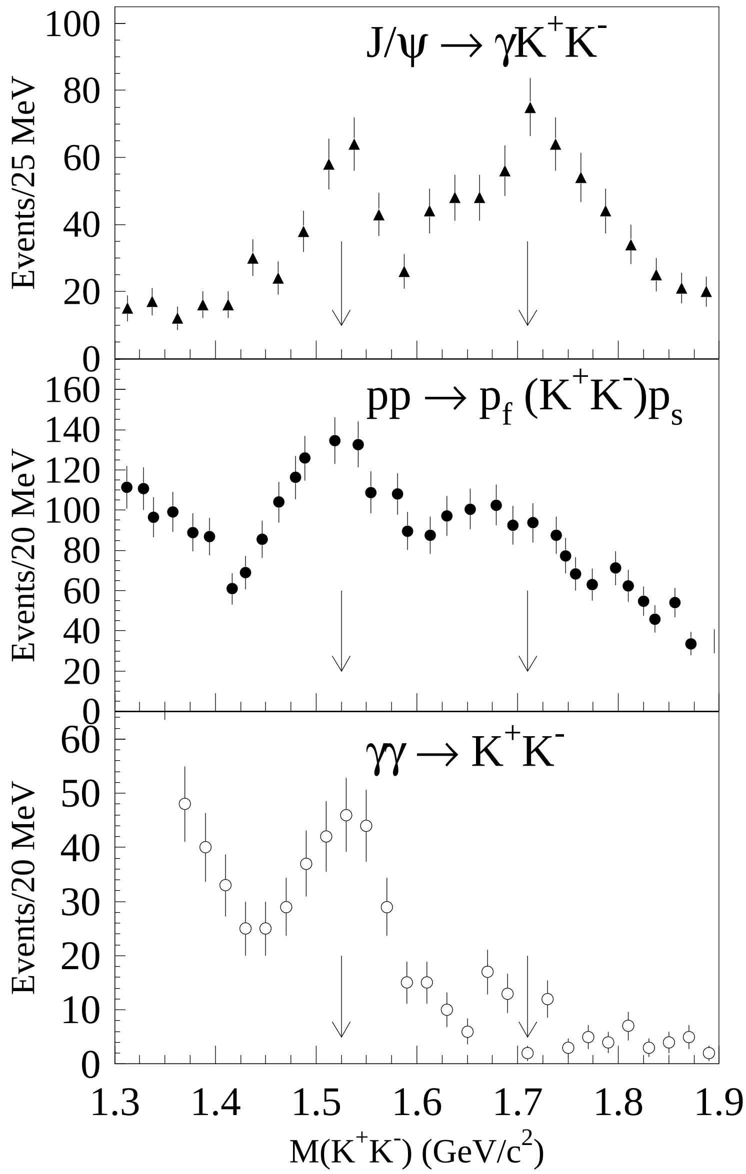

Figure 25: Possible glueball sensitivity in different reactions, for states decaying to K^+K^-. Shown are the K^+K^- invariant mass distributions for $J/\psi \to \gamma K^+K^-$ (Bai, et al., 1996b), $pp \to p_f(K^+K^-)p_s$ (Armstrong, et al., 1991a), and $\gamma\gamma \to K^+K^-$ (Albrecht, et al., 1990).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

shows the invariant mass of K^+K^- pairs produced in radiative J/ψ decay, $J/\psi \to \gamma K^+K^-$ (Bai, 1996b); central pp collisions, $pp \to p_f(K^+K^-)p_s$ (Armstrong, 1991a); and two-photon collisions, $\gamma\gamma \to K^+K^-$ (Albrecht, 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. 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 Section 2.5.1 (Eqn. 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.

5.3.1 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, 1993) and $M(0^{++}) = 1740 \pm 71$ MeV/c^2 and $M(2^{++}) = 2359 \pm 128$ MeV/c^2 (Chen, 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. 2.5.1 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 accomodate 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; Lee and Weingarten, 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. 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, 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 determining its other properties. We will 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.

Radiative J/ψ decay. This state was first observed by the Crystal Ball collaboration, in radiative J/ψ decay (Edwards, 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 ($\sim$ 220 MeV/c^2) and the $\theta \to \eta\eta$ angular distribution favored $J = 2$, however 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 $K\bar{K}$ fueled speculation that this was a glueball. A reanalysis of the Crystal Ball data which included the $f_2'(1525)$ (Bloom and Peck, 1983; Königsman, 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, 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, 1987) and K^+K^- (Augustin, 1988) channels, as well as in K_SK_S (Augustin, 1988).

The published Mark III result suggesting $J = 2$ (Baltrusaitis, 1987) has actually been called into question in unpublished reports by the same collaboration (Chen, 1990; Chen, 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, 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 $f_2(1270)$ 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 $f_J(1710)$ to consist mainly of the scalar glueball, as its decays to $\pi\pi$ preclude identifying it as the $s\bar{s}$ scalar nonet member.

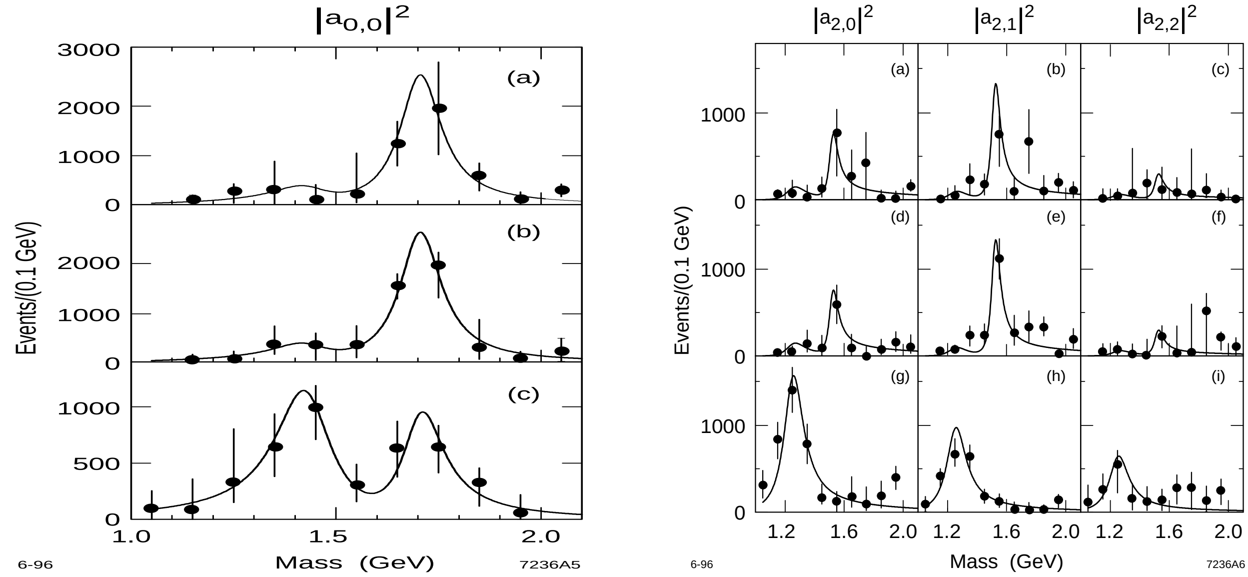

Figure 26: Partial wave analysis of J/ψ radiative decay to two pseudoscalar mesons, from the Mark III collaboration. Tha analysis allows all partial wave components to vary independently for both D-wave (left) and S-wave (right). The top row is the analysis of $J/\psi \to \gamma K_S K_S$, middle row 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.

Other analyses of J/ψ radiative decay, leading to 4π decay modes of the $f_J(1710)$, also 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 to 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 this data (Bugg, 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; Chen, 1991; Dunwoodie, 1997). In particular, the D-wave amplitudes are forced to be relatively real, reducing the number of unknowns which need to be determined from the fit. Their analysis, however, shows that there are indeed more than one overlapping D-wave state in this region, and the assumption 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; Chen, 1991; Dunwoodie, 1997). Their analysis (Fig. 26) gives

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

Central production in pp collisions. Supposing that it proceeds through double Pomeron exchange, one might suspect that glue-rich states are produced centrally in pp collisions. See Sec. 3.4. The $f_J(1710)$ has been studied in this way at CERN by the WA76 (Armstrong, 1989a, 1991a) and the WA102 (Barberis, 1997b). Final states include $\pi^+\pi^-$ (Armstrong, 1991a), K^+K^- and K_SK_S (Armstrong, 1989a), and $\pi^+\pi^-\pi^+\pi^-$ (Barberis, 1997b). As 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 the $f_1(1285)$, 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, 1994) seems to prefer $J^P = 2^+$. 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 are found throughout this region as well as other structures.

Peripheral hadronic reactions. The $f_J(1710)$ has generally been unobserved in peripheral hadronic reactions. A measurement of the reaction $\pi^-p \to K_SK_Sn$ 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, 1986), found no evidence for the $f_J(1710)$ except in the J/ψ radiative decay data. A measurement of $K^-p \to K_SK_S\Lambda$ by the LASS collaboration (Aston, 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 $\pi^-p \to K_SK_Sn$ at BNL (Etkin, 1982b) and a detailed analysis of the K_SK_S S-wave (Etkin, 1982c) reveals two states that are more or less consistent with both the $f_0(1500)$ and $f_0(1710)$. Produced in charge exchange with a π beam, it is reasonable to assume the states also couple to $\pi\pi$. It might be plausible to argue that these two states 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, 1988). In fact, a reanalysis of this 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, 1992) observes a state which may or may not be the $f_J(1710)$. Called $X(1740)$, it is observed decaying to $\eta\eta$ in the reaction $\pi^-p \to \eta\eta N^\star$, 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 $\pi^-p \to \eta\eta N^\star$ with $N^\star \to n + \gamma$'s, but not for $\pi^-p \to \eta\eta n$. 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.

Two-photon collisions. One expects glueballs to be absent in two-photon production. In studies of $\gamma\gamma \to K\bar{K}$ (Althoff, 1985; Behrend, 1989c; Albrecht, 1990) a clear signal for $f_2'(1525)$ is evident, but only upper limits are put on $\Gamma_{\gamma\gamma}$ for $f_J(1710)$. 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 $\gamma\gamma \to K_SK_S$ would be particularly useful along with a complete partial wave analysis in the 1400 to 1800 MeV/c^2 region.

5.3.2 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, 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, 1993; Chen, 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 $J/\psi \to \gamma K^+K^-$ and $J/\psi \to \gamma K_SK_S$ by the Mark III collaboration (Baltrusaitis, 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 is 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, 1987) and K^+K^- and K_SK_S (Augustin, 1988) and do not see the $f_J(2220)$ in any of these three channels. They quote a limit on the product branching ratio, that is $B[J/\psi \to \gamma f_J(2220); f_J(2220) \to K^+K^-]$, 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_SK_S, $p\bar{p}$ (Bai, 1996a) and $\pi^0\pi^0$ (Bai, 1998) channels. Mass distributions in the region of the $f_J(2220)$ are reproduced from (Bai, 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

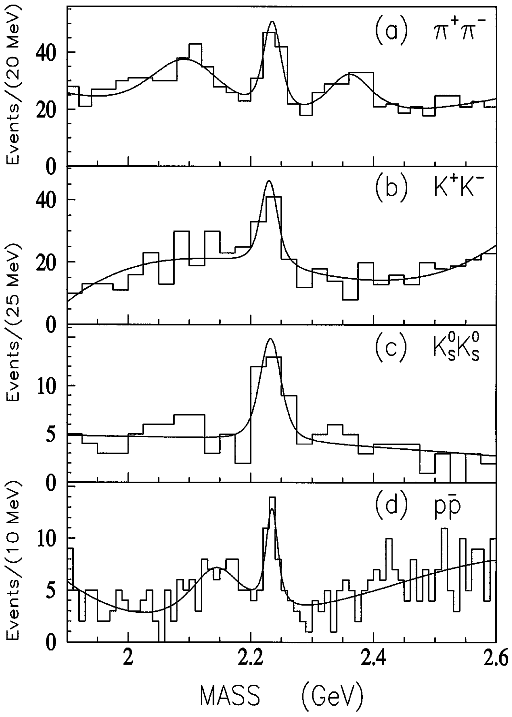

Figure 27: Figure 2 from (Bai, et al., 1996a) showing various two-particle mass distributions observed in J/ψ radiative decay. Each shows a signal for the $f_J(2220)$. The lines represent fits to a smooth background and one or more Breit-Wigner resonance shapes, convoluted with the appropriate Gaussian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

CLEO collaboration in the reactions $\gamma\gamma \rightarrow K_S K_S$ (Godang, 1997) and $\gamma\gamma \rightarrow \pi^+\pi^-$ (Alam, 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 large as the $f_2(1270)$. Despite the clear presence of the $f_2'(1525)$, Godang *et al.* (Godang, 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, 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 $J \geq 2$. 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 make it clear that spins greater than $J = 2$ are required to describe the data.

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 $J \geq 2$ 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 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, 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 $p\bar{p} \to \pi^+\pi^-$ (Hasan and Bugg, 1996), $p\bar{p} \to K^+K^-$ (Sculli, 1987; Bardin, 1987), and from the JETSET collaboration $p\bar{p} \to K_S K_S$ (Evangelista, et al, 1997), $p\bar{p} \to \phi\phi$ (Evangelista, et al, 1998), and $p\bar{p} \to p\bar{p}\pi^+\pi^-$ (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 $B[J/\psi \to \gamma f_J(1710)]$ 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, 1997)

$$B[f_J(2220) \to p\bar{p}] \times B[f_J(2220) \to K_S K_S] \leq 7.5 \times 10^{-5} \qquad (95\% C.L.)$$

with the values from BES (Bai, 1996a)

$$\begin{aligned} B[J/\psi \to \gamma f_J(2220)] \times B[f_J(2220) \to K_S K_S] &= (2.7 \pm 1.1) \times 10^{-5} \\ B[J/\psi \to \gamma f_J(2220)] \times B[f_J(2220) \to p\bar{p}] &= (1.5 \pm 0.6) \times 10^{-5} \end{aligned}$$

then we can infer the lower bound

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

This is not only an inordinately large branch for a radiative decay, it also implies that all the branches reported by BES, $B[J/\psi \to \gamma f_J(2220); f_J(2220) \to X] \approx 1.5 \times 10^{-4}$, 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.

5.4 $J^{PC} = 0^{-+}$ and 1^{++} states in the E region

In principle, the $K\bar{K}\pi$ final state is a good way to study mesons which couple to $s\bar{s}$ but which are forbidden to decay to $K\bar{K}$. This would include mesons with $J^P = \text{odd}^+$ or $J^P = \text{even}^-$. Indeed, experiments which produce $K\bar{K}\pi$ show significant structure in the mass spectra. As an example, Fig. 28(a) histograms the $K\bar{K}\pi$ mass for the reaction $\pi^- p \to K^+ K_S \pi^- n$ at 18 GeV/c (Cummings, 1995). Relatively narrow structures are clear at ~ 1300 MeV/c^2 and ~ 1400 MeV/c^2. These structures are traditionally referred to as the "D" and "E" regions, respectively. The names D and E were historically used 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. The name given to $J^{PC} = 0^{-+}$ strength within this peak was originally called $\iota(1450)$ and for that reason many still refer to this controversy as the E/ι puzzle. The discovery of the $\iota(1450)$ led to enormous excitement that the first glueball may have been found (Scharre, 1980; Ishikawa, 1981; Edwards, 1982a, Aihara, 1986a) but this is no longer the generally accepted viewpoint. There are a number of experimental inconsistencies which make this interpretation difficult. Furthermore, lattice gauge theory would have great difficulty accomodating such a low mass glueball with these quantum numbers.

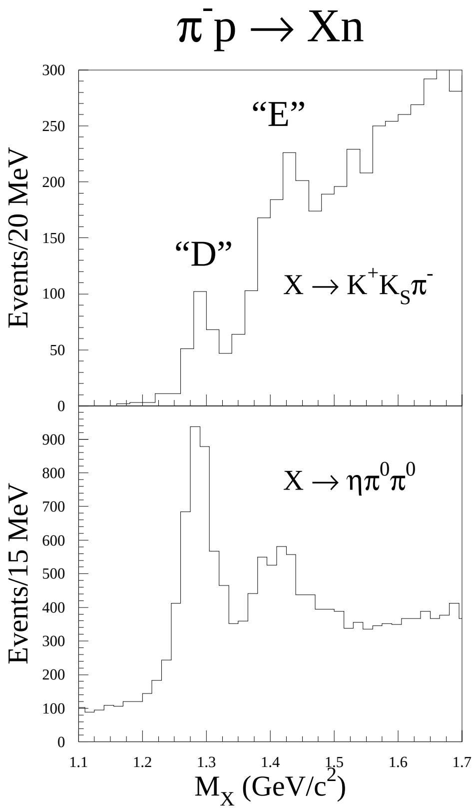

Figure 28: Invariant mass distributions for states produced in peripheral production with high energy π^- beams. In (a) is shown the $K^+K_S\pi^-$ distribution at 18 GeV/c (Cummings, 1995) and (b) shows the $\eta\pi^0\pi^0$ result at 100 GeV/c (Alde, et al., 1997). In each case, peaks in the the D and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

A large part of the difficulty is that these structures lie very close to the $K^\star\bar{K}$ threshold[6], 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 $K^\star\bar{K}$ pair. However, the limited phase space leads to very small relative momenta between the K and $\bar{K}$. As a consequence, coupling through the $a_0(980)$, which has a great affinity for $K\bar{K}$, is very likely, leading to S- and P-wave $a_0(980)\pi$ decays, that is $J^{PC} = 0^{-+}$ and $J^{PC} = 1^{++}$ states. It is therefore not surprising that the $K\bar{K}\pi$ system can get very complicated.

A byproduct of the nearness to $K^\star\bar{K}$ threshold, and a possible tool for unraveling this structure, is suggested by the decay $a_0(980) \to \eta\pi$. 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 $\pi^- p \to \eta\pi^0\pi^0 n$ at 100 GeV/c (Alde, 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, 1991; Manak, 1997), because of additional quasi-twobody states such as $\eta\rho$.

Besides peripheral hadroproduction, $K\bar{K}\pi$ and $\eta\pi\pi$ final states in the E region have been studied in J/ψ radiative decay, central production, and in two-photon collisions. 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 $K\bar{K}\pi$ through $a_0(980)\pi$, has mass between 1400 and 1420 MeV/c^2. The other decays to $K^\star\bar{K}$ and has mass between 1470 and 1480 MeV/c^2. (Note that the Particle Data Group (Caso, 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, called $f_1(1420)$. This is seen most clearly in single-tagged two-photon collisions, which identifies the quantum numbers unambiguously. It is also clearly seen in central production and in $p\bar{p}$ annihilation at rest in gaseous hydrogen.
- A second $J^{PC} = 1^{++}$ state, the $f_1(1510)$, has been identified. It decays to $K^\star\bar{K}$ and is produced in peripheral hadroproduction with both K^- and π^- beams. (See Sec. 4.2.3.)
- There is no evidence for the $\eta(1440)$ in untagged two photon collisions. It is considerably more "sticky" than the η or η'. This large stickiness is attributed to a very large J/ψ radiative decay width, and this 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 $K^\star\bar{K}$ and produced in $K^- p$ interactions and in $p\bar{p}$ annihilations at rest. (See Sec. 4.2.3.)

Even putting aside various experimental inconsistencies (which we detail below), there is already serious difficulty accomodating these results. In the quark model, one might expect

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

two states in the E region, the $s\bar{s}$ partners to the $f_1(1285)$ and $\eta(1295)$. The evidence, is, however, that there are two states of each J^{PC}. Accomodating the overpopulation with states having gluonic degrees of freedom, however, 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 $K^\star\bar{K}$ threshold is at the heart of the difficulties, both in the experiments and in their interpretation.

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

5.4.1 $J^{PC} = 1^{++}$

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 $K^\star\bar{K}$ 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, 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, 1988; Cummings, 1995) and $K^- p \to Xn$ (Aston, 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, 1988c) using $K^- p \to K_S K^\pm \pi^\mp \Lambda$ 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, shows a predominance of $J^P = 1^+$, and also shows that this cannot be the result of a single resonance. By symmetrizing the final state on the basis of charge conjugation, the $J^P = 1^+$ strength is 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. 3.5.2) of $K\bar{K}\pi$ final states (Aihara, 1986b; Gidal, 1987b; Hill, 1989; Behrend, 1989b) all show clear evidence of the $J^{PC} = 1^{++}$ state, although the mass is generally more consistent with 1420 than 1510 MeV/c^2, with a $\gamma\gamma^\star$ width (times $K\bar{K}\pi$ branching ratio) of 1.7$\pm$0.4 keV (Caso, 1998). On the other hand, single-tagged production of $\eta\pi\pi$ (Gidal, 1987a; Aihara, 1988) 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, 1998). This suggests a consistency with the $f_1(1285)$ and $f_1(1420)$ being $n\bar{n}$ and $s\bar{s}$ partners, but the comparison is likely 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, 1990; Augustin, 1992) shows a signal for $J^{PC} = 1^{++}$ consistent with a mass closer to 1420 MeV/c^2, and $B(J/\psi \to f_1(1420)\gamma) \times B(f_1(1420) \to K\bar{K}\pi) = 8.3 \pm 1.5 \times 10^{-4}$ (Caso, 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, 1997a) with 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 (Abele, 1997b) for the $J^{PC} = 1^{+-}$ state observed by LASS (Aston, 1988c). Central production of $K\bar{K}\pi$ (Armstrong, 1989b, 1992; Barberis, 1997c) and of $\eta\pi\pi$ (Armstrong, 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, 1997b). 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.

5.4.2 $J^{PC} = 0^{-+}$

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 as 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 5 shows results from several experiments which 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, determine two states. The masses differ by about 50 MeV/c^2, with the heavier decaying to $K^\star\bar{K}$, except for (Augustin, 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 $\pi^- p \to \eta\pi^0\pi^0$ at 100 GeV/c (Alde, 1997) and $\pi^- p \to \eta\pi^+\pi^-$ at 18 GeV/c (Manak, 1997) show that most of the $J^{PC} = 0^{-+}$ $\eta\pi\pi$ signal is

Table 5: Masses, in MeV/c^2, of $J^{PC} = 0^{-+}$ $K\bar{K}\pi$ resonances in the E region.

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

concentrated in $\eta(\pi\pi)_S$ instead of $a_0(980)\pi$.

Clearly, this leads us to guess that one of these states is the $s\bar{s}$ partner of the $\eta(1295)$, and the other is some manifestation of non $q\bar{q}$ degrees of freedom. That would imply, however, that one of the states (presumably the one that decays to $K^\star\bar{K}$) 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 $s\bar{s}$ pseudoscalar meson in the E region were obtained by (Behrend, 1989b) in the reaction $\gamma\gamma \to K_S K^\pm\pi^\mp$. They find $\Gamma_{\gamma\gamma}[\eta(1440)] \times B[\eta(1440) \to K\bar{K}\pi] < 1.2$ 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 $s\bar{s}$ 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, 1998).

There is some evidence of $J^{PC} = 0^{-+}$ production in $K_S K^\pm\pi^\mp$ final states in pp and π^+p central collisions (Armstrong, 1992) but the signal in the E region is dominated by 1^{++}.

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 $K^\star\bar{K}$ threshold brings in more complicated mechanisms that can be treated with the formalisms presently at our disposal.

5.5 Other Puzzles

Thus far we've 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 compeleteness, we briefly discuss some additional examples.

5.5.1 Extra $J^{PC} = 2^{++}$ States

The ground state 2^{++} nonet, consisting of the $a_2(1325)$, $f_2(1270)$, $f_2'(1525)$, and $K^*(1430)$, has been complete for some time. However, several additional isoscalar 2^{++} states have been observed which are inconsistent with quark model predictions. These include the three g_T states in $\phi\phi$ at 2011, 2297 and 2339 MeV; the $f_2(1565)$ seen in $\bar{p}N$ annihilation; and the $f_2(1430)$ and $f_2(1480)$ observed in $\pi\pi$ and $K\bar{K}$ spectra between the $f_2(1270)$ and the $f_2'(1525)$.

OZI suppression and states in $\pi^- p \to \phi\phi n$ In the high mass region three states, somtimes known as g_T, have been observed in the OZI-suppressed reaction $\pi^- p \to \phi\phi n$ at 22 GeV/c by a BNL group (Etkin, 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, 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 $f_2(2340)/g_T''$ (Longacre, 1986; Etkin, 1985). 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 $M = 2011 \pm 70$ MeV/c^2 and $\Gamma = 202 \pm 65$ MeV/c^2 (about 98% S_2); g_T' with $M = 2297 \pm 28$ MeV/c^2 and $\Gamma = 149 \pm 41$ MeV/c^2 (about 25% D_2 and 69% D_0); and g_T'' with $M = 2339 \pm 55$ MeV/c^2 and $\Gamma = 319\pm$ 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 are difficult to observe due to large backround in these channels. As they are produced in the OZI-forbidden channel (Landberg, 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 $\pi^- Be$ interactions at 85 GeV/c. They see general enhancement at the $\phi\phi$ threshold followed by a second peak at 2.4 GeV. Assuming that they see the second and third g_T states, they have 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 have also carried out a joint moment analysis and find that the $\phi\phi$ system up to 2.5 GeV/c^2 is mainly 2^{++} (Booth 1986, Armstrong 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 are indeed glueballs. The GAMS group (Alde, 1988b) observes 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^{++} is found to be the main wave, no threshold enhancement in the $\phi\phi$ system is observed, in contrast to the hadronic production. However, they find 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; Toki, 1987; Blaylock, 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 $\omega\omega$, 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, 1998) studied the reaction $\bar{p}p \to \phi\phi$ with antiproton beams between 1.1 and 2.0 GeV/c momentum. This experiment is sensitive to intermediate states formed in the annihilation channel with masses between 2.1 and 2.4 GeV/c^2. There is no evidence 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, 1998). Some weak structure is observed in the $\phi\phi$ mass distribution, but the statistics are poor. The angular distribution favors $J^{PC} = 2^{++}$.

The $f_2(1565)$ This state has a long history going back to the 1960's (Bettini 1966; Conforto 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, 1989, 1990a). A state with $M = 1565 \pm 10$ MeV/c^2 and $\Gamma = 170 \pm 20$ 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 is an $I = 0$ resonance. A Dalitz plot analysis showed clear evidence for $J^{PC} = 2^{++}$. This resonance cannot 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, 1967). In a separate analysis (May 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, 1991) which gives 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 are 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 requires 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, 1998).

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 an experiment R807 at CERN ISR (Akesson, 1986; Cecil, 1984). The reaction concerns the exclusive $\pi^+\pi^-$ production in $pp \to p_f(\pi^+\pi^-)p_s$ at $\sqrt{s} = 63$ GeV. The recoil protons have been detected with $-t$ less than 0.03 (GeV/c)^2, thus ensuring nearly pure Pomeron exchanges at both vertices. The resulting $\pi\pi$ spectrum exhibits a set of remarkable bump-dip structures near 1.0, 1.5 and 2.4 GeV, respectively (Akesson 1986). Another striking feature is that the $\rho(770)$ and the $f_2(1270)$ are not seen in the data. One may 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 drop-off near 1.4 GeV is partly due to the $f_0(1370)$ and a D-wave structure which is attributed to the $f_2(1480)$. The data in fact favours 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 of absence of the $f_2(1270)$ in the R807 data. An explanation may follow from Close and Kirk's (Close 1997b) glueball filter.

5.5.2 Structure in $\gamma\gamma \to VV$

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

This subject originated from the original observation by the TASSO collaboration (Brandelik, 1980) of structures in the $\gamma\gamma \to \rho^0\rho^0$ cross section. This was subsequently confirmed by other experiments (Burke, 1981; Althoff, 1982; Behrend, 1984; Aihara, 1988), where 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 the energy at which the cross sections peak also varies from one final state to another.

The large difference in cross section between the $\rho^0\rho^0$ and $\rho^+\rho^-$ channels, $\sigma(\gamma\gamma \to \rho^0\rho^0) \simeq 4\sigma(\gamma\gamma \to \rho^+\rho^-)$ (Albrecht, 1996; Behrend, 1989a) rules out a simple s-channel resonance explanation where one expects the decay of a conventional resonance into $\rho^+\rho^-$ to occur with a rate two times as often as that 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 from interference between the two. The interference is observed to be constructive in $\gamma\gamma \to \rho^0\rho^0$ and destructive in $\gamma\gamma \to \rho^+\rho^-$. 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, 1982; Li, 1983).

A number of models have been invoked to explain the structure in the original $\gamma\gamma \to \rho^0\rho^0$ 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 $\gamma\gamma \to$

VV' by (Achasov, 1982) and by Li and Liu (Li, 1983). The vector dominance model with factorization in the t-channel attempts to identify specific t-channel exchange and extract them from photoproduction data (Alexander 1982). Perturbative QCD with Coulombic rescattering corrections were examined by Brodsky (1987). In the one meson exchange model (Achasov, 1988; Törnqvist 1991) the structure in $\gamma\gamma \to \rho\rho$ 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.

5.5.3 The $C(1480)$

The Lepton-F collaboration at Serpukhov (Bityukov, 1987) examined the reaction $\pi^- p \to \phi\pi^0 n$ 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 is substantial. Interpreted as a resonance, called $C(1480)$, the peak has mass 1480 ± 40 MeV/c^2, width 130 ± 60 MeV/c^2, and $\sigma(\pi^- p \to Cn) \times B(C \to \phi\pi^0) = 40 \pm 15$ 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, 1988; Kopeliovich, 1995).

This state has not been observed in other reactions, including pp central production (Armstrong, 1992) and $p\bar{p}$ annihilation at rest (Reifenröther, 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).

5.6 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. A particularly glaring example 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, 1993), i.e. ${}^3D_2$ and ${}^1D_2$, although they need to be confirmed (PDG, Caso, 1998).

It will be 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 Tab. 6.

Consider the $\eta_2(1^1D_2)$. There is some evidence that this state has in fact been observed (Sec. 4.2.5). We expect it to be almost degenerate in mass with its non-strange isovector partner, the $\pi_2(1670)$. From Tab. 6 we see that it is expected to be rather broad and it

Table 6: 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	~ 400
	$BR(\eta_2 \to a_2\pi)$	$\sim 70\%$
	$BR(\eta_2 \to \rho\rho)$	$\sim 10\%$
	$BR(\eta_2 \to K^*\bar{K} + c.c.)$	$\sim 10\%$
η_2' (1^1D_2)	Mass	1890
	width	~ 150
	$BR(\eta_2' \to K^*\bar{K} + c.c.)$	$\sim 100\%$
ω_1 (1^3D_1)	Mass	1660
	width	~ 600
	$BR(\omega_1 \to B\pi)$	$\sim 70\%$
	$BR(\omega_1 \to \rho\pi)$	$\sim 15\%$
ρ_2 (1^3D_2)	Mass	1700
	width	~ 500
	$BR(\rho_2 \to [a_2\pi]_S)$	$\sim 55\%$
	$BR(\rho_2 \to \omega\pi)$	$\sim 12\%$
	$BR(\rho_2 \to \rho\rho)$	$\sim 12\%$
ω_2 (1^3D_2)	Mass	1700
	width	~ 250
	$BR(\omega_2 \to \rho\pi)$	$\sim 60\%$
	$BR(\omega_2 \to K^*\bar{K})$	$\sim 20\%$
ϕ_2 (1^3D_2)	Mass	1910
	width	~ 250
	$BR(\phi_2 \to K^*\bar{K} + c.c.)$	$\sim 55\%$
	$BR(\phi_2 \to \phi\eta)$	$\sim 25\%$

decays predominantly through the $a_2(1320)$ isobar which in turn decays to $\rho\pi$. The 4π final state is complicated to reconstruct. Since $a_2 \to \eta\pi$, 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 similar mass as the $\pi_2(1680)$ 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 will be important to search thoroughly through data that is 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 1997). In general, given how complicated the meson spectrum is, it appears that a good starting place would be the strange mesons. The reason for this is that in the strange meson sector we don't have the additional problem of deciding whether new states are glueballs or conventional mesons and we don't have the additional 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.

6 FUTURE DIRECTIONS

To understand 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 chapters just begins to scratch the surface of this field. There is surely much new physics to be gleaned 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 radially excited multiplets.

It is unlikely that all this could be done simply by bump hunting. Rather, we will 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 (Barnes 1997). As an example, in Sec. 5.6, we discussed the Quark Model predictions for the properties of the missing $L = 2$ mesons.

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; Reyes, 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 7. These are described in some detail in the following sections.

Table 7: 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
KEK	e^+e^-	10 GeV	B decays; $\gamma\gamma \to X$	1999
COMPASS	p	400 GeV	Central Production	1999
RHIC	p, Nuclei	~ 200 GeV$/A$	Central Production	1999
JHF	p	50 GeV	$\{\pi, K, p\} \to X$	~2004
CEBAF	e^-, γ	12 GeV	$\gamma p \to Xp$	~2004
BEBC	e^+e^-	3-4 GeV	$J/\psi \to \gamma X$	~2004

6.1 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 $\phi \to \{a_0, f_0\}\gamma$. See (Achasov, 1997a).

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.

6.2 B-factories at CESR, SLAC, and KEK

Although the primary motivation for a B factory (Goldberg and Stone, 1989; Bauer 1990, 1992) is to study CP violation in the B meson system, light meson spectroscopy will also be addressed 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 $\Upsilon \to \gamma X$ providing 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 a new facility 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) which has some implications for detection of lower mass sytems, 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.

6.3 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 **CO**mmon **M**uon and **P**roton **A**pparatus for **S**tructure and **S**pectroscopy (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 multi-purpose apparatus has specific challenges, mainly trying to optimize sensitivies to the different situations. However, using a common apparatus with a large directed collaboration has lead 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.

6.4 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 luminosiites at heavy ion colliders are substantially less than 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 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 the 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).

6.5 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. This has been aptly demonstrated by E852 at BNL, for example. It is also important to recognize the lessons taught by the LASS experiment as 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) which 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.

6.6 CEBAF at Jefferson Laboratory

Peripheral photoproduction is conspicuously absent 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. 3.2. 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 centered around the laboratory's original primary goal, i.e the study nuclear structure. However, because of landmark improvements in the the RF accleration cavities, beams with energies up to $\sim$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 bremstrahlung or through a variety of other means.

An aggressive experimental plan is underway which 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 take 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.

6.7 A τ-Charm factor at BEBC

τ-charm factories, e^+e^- colliders operating in the energy region of 3—4 GeV, have been proposed although none have yet been approved. τ-charm factories will 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 between a factor of 10 and 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 the 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 radially excitations of the vector mesons.

7 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 are the physical states of the theory. Understanding QCD, and non-Abelian gauge theories in general, is one of the most important problems facing high energy and nuclear physics. We outlined the important issues in light meson spectroscopy, the puzzles, and the open questions. We have attempted so show where the field is heading with the next generation of experiments and how they can advance our knowledge of mesons 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. We have attempted to point these out.

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 to 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 ground for the search for exotic states.

Most of the significant deviations from the quark model occur in the light-quark isoscalar sector. 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 this will have to wait for better experiments and clearer phenomenology to establish a consistent pattern.

The long outstanding E/ι problem remains with us. The evidence clearly points to two separate states each for $J^{PC} = 1^{++}$ and 0^{-+}, where exactly one of each are expected based on $q\bar{q}$ degrees of freedom. Ascribing hybrid, glueball, or multiquark status to the extra states is problematic, though, 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 $K^\star\bar{K}$ 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 not have an underlying broad structure which 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 which necessarily imply a meson state beyond $q\bar{q}$ models. Evidence is finally beginning to sort itself out in the $\eta\pi$ channel, and does now seem to point towards an exotic $J^{PC} = 1^{-+}$ state near 1380 MeV/c^2. It is difficult to accomodate 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 which 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.

Acknowledgments

This research was supported in part by the Natural Sciences and Engineering Research Council of Canada, and by the National Science Foundation in the US. The authors thank Gary Adams, Ted Barnes, Bob Carnegie, Suh-Urk Chung, John Cummings, Bill Dunwoodie, Alex Dzierba, Richard Hemingway, Nathan Isgur, Joseph Manak, Curtis Meyer, Colin Morningstar, Phillip Page, Vladimir Savinov, Eric Swanson, Don Weingarten, John Weinstein, and Dennis Weygand for helpful conversations and communications.

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