EVIDENCE FOR THE OBSERVATION OF A GLUEBALL

DON WEINGARTEN
IBM Research, Yorktown Heights
NY 10598, USA

I briefly review lattice QCD, the valence (quenched) approximation, and the application of both to the determination of the mass and two-body decay couplings of the lightest scalar glueball. Results in agreement with the observed properties of $f_J(1710)$ strongly suggest this resonance is largely a scalar glueball.

1 Introduction

The existence of chromoelectric field is the key hypothesis QCD adds to the quark model. The observation in experiment of a glueball would be a direct confirmation of the existence of chromoelectric field. So it would be nice to find one. Identifying a glueball in experiment is tough, however. The kinematics of QCD provides no clear footprint for glueballs comparable, say, to the footprint the Weinberg-Salam model gives for the W’s and the Z. The lightest states, which would be easiest to detect, are not expected to differ drastically in mass or decay properties from mesons containing quarks and antiquarks. On the other hand, $J^{PC}$ combinations which are impossible for quark-antiquark states but which can occur for heavier glueball excitations can also be realized with multiple quark-antiquark pairs. A reliable calculation of the consequences of QCD’s dynamics for the masses and decay properties of glueballs appears to be a necessary element in the identification of glueballs in experiment. At present such calculations can be done only numerically, using the lattice formulation of QCD in combination with the valence approximation.

As of three years ago, two independent calculations had been completed of the infinite-volume, continuum limit of the mass of the lightest glueball, which turns out to be a scalar. Hong Chen, Jim Sexton, Alessandro Vaccarino and I obtained a value of 1740(71) MeV using ensembles of 25000 to 30000 gauge configurations on each of several different lattices. An earlier valence approximation calculation by the UKQCD-Wuppertal collaboration yields 1625(94) MeV for the lightest scalar glueball mass. This calculation used several different lattices with ensembles of between 1000 and 3000 configurations each. If the two mass evaluations are combined, taking into account the correlations between their statistical uncertainties arising from a common procedure for converting lattice quantities into physical units, the result is 1707(64) MeV for the scalar glueball.
mass. Both the mass prediction with larger statistical weight and the combined mass prediction agree with the mass of $f_J(1710)$ and are inconsistent with all but $f_0(1500)$ among the established flavor singlet resonances which could be scalars. For $f_0(1500)$ the disagreement is still by more than three standard deviations. In addition, observed mass values for other scalar quark-antiquark states suggest a mass near 1500 MeV for the lightest $s\bar{s}$ scalar.

These calculations by themselves, however, do not make a strong case for taking $f_J(1710)$ as a scalar glueball. The key question which the mass results do not answer is whether the lightest scalar glueball has a decay width small enough for this particle actually to be identified in experiment. If the scalar glueball had a width of a GeV or more, the prospect of ever finding one seems remote. But a scalar glueball with a width of a few hundred MeV or less and mass in the neighborhood of 1700 MeV would be hard to miss and should already have been seen in experiment. A further question in the identification of $f_J(1710)$ as a glueball is raised by the argument that since glueballs are flavor singlets they should have the same decay couplings to $\pi \pi$, to $K\bar{K}$, and to $\eta\eta$. This equality is violated by $f_J(1710)$ decay couplings.

To address these questions, Jim Sexton, Alessandro Vaccarino and I have calculated the decay coupling constants of the lightest scalar glueball to pairs of pseudoscalar mesons. The calculation was done in the valence approximation on a lattice $16^3 \times 24$ with $\beta$, defined as $6/g^2$, set at 5.70, corresponding to lattice spacing, determined from the $\rho$ mass, of about 0.15 fm and lattice period of about 2.3 fm. For the total width of the scalar glueball to pairs of pseudoscalar quark-antiquark states we obtained $108 \pm 29$ MeV. The combined correction for the errors in our prediction arising from the valence approximation, from finite lattice spacing and from finite lattice volume we believe would be less than 50%. Based on our value for two-body decays, any reasonable guess for the partial width for multibody decay modes leads to a total width small enough for the scalar glueball to be hard to miss in experiment. In fact, the observed width of $f_J(1710)$ into pairs of pseudoscalars is $99 \pm 15$ MeV, consistent with our result. We obtain also a violation of the expected equality of glueball decay rates to $\pi_0\pi_0$, to $K\bar{K}$, and to $\eta\eta$ in agreement with the observed results for $f_J(1710)$.

Although so far I have simplified the story by supposing that physical resonances are either entirely glueballs or entirely quark-antiquark, another possibility is a state which is a linear combination of a glueball and a quark-antiquark state. In the valence approximation, however, glueballs contain no admixture of configurations with valence quarks or antiquarks. Thus we consider the agreement between the mass and decay couplings found in the valence approximation and the observed mass and decay couplings of $f_J(1710)$ to be
strong evidence that this state is largely a scalar glueball with at most some relatively smaller amplitude for configurations including valence quark-antiquark pairs.

The glueball calculations presented here were carried out on the GF11 parallel computer \[ \text{at IBM Research and took approximately 30 months to complete at a sustained computation rate of between 6 and 7 Gflops.} \]

In the remainder of this talk, I will briefly review lattice QCD, the valence approximation, the glueball mass calculation, the decay calculation, and close with a more detailed comparison of the lattice predictions with experiment.

## 2 Lattice QCD

Lattice QCD, as its name implies, approximates continuous space-time (at negative imaginary values of time) by a discrete lattice including only a finite number of points. Predictions for the real world of continuous, infinite-volume space-time are supposed to be recovered from lattice QCD by taking the limit of lattice predictions as lattice spacing goes to zero and lattice volume to infinity.

It is convenient, though not necessary, to take the lattice to be hypercubic with periodic boundary conditions. Living at each site are lattice versions of the chromoelectric potential, \( A_\mu^j(x) \), given by real numbers, and the quark and antiquark fields, \( \Psi_{sc}(x) \) and \( \Psi_{sd}(x) \), respectively, given by Grassmann variables. Here \( x \) is a lattice site, \( j, c \) and \( d \) are color indices for the adjoint, fundamental and conjugate representations, respectively, \( \mu \) is a lattice direction and \( s \) is a spin index.

Vacuum expectation values of time ordered products of fields are then defined by a path integral. A typical vacuum expectation value, which I offer in place of the slightly more complicated general rule, is

\[
\langle \bar{\Psi}(x) \Psi(x) \rangle = Z^{-1} \int \prod d\mu A \int \prod d\Psi d\bar{\Psi} \bar{\Psi}(x)\Psi(x) \exp(S_A + S_\Psi),
\]

\[
Z = \int \prod d\mu A \int \prod d\Psi d\bar{\Psi} \exp(S_A + S_\Psi), \tag{1}
\]

\[
S_A = -\frac{1}{4g^2} \int d^4xF_{\mu\nu}^jF_{\mu\nu}^j.
\]

\[
S_\Psi = \int d^4x \bar{\Psi}(\not{\bar{\Psi}} - \not{A} + im)\Psi,
\]

where \( \int \prod d\mu A \) can be thought of, but is not quite, a product over all \( x, \mu \) and \( j \) of the integral \( \int_{-\infty}^\infty dA_\mu^j(x) \) and \( \int \prod d\Psi d\bar{\Psi} \) is a product of Grassmann integration on each \( \Psi_{sd}(x) \) and \( \Psi_{sc}(x) \). For the continuum quantities and
integrals which appear in Eqs. (1), I actually intend lattice approximations but use the continuum expressions as more recognizable alias’s.

With respect to the quark and antiquark fields, Eqs. (1) are a Grassmann version of gaussian integrals. The integrals can be done analytically and give

\[ <\Psi(x)\Psi(x)> = Z^{-1} \int \prod d\mu_A \ tr((\not{p} - \not{A} + im)^{-1}(x,x)) \times \]

\[ det(\not{p} - \not{A} + im) \ exp(S_A), \]  

(2)

\[ Z = \int \prod d\mu_A det(\not{p} - \not{A} + im) \ exp(S_A), \]

where the trace is with respect to spin and color indices.

The lattice QCD industry does integrals like those in Eq. (2) by Monte Carlo. With a large enough ensemble of random \(A\) fields, \([A^j_\mu(x)]_1, \ldots [A^j_\mu(x)]_N\), generated by computer according to the differential probability

\[ dv = Z^{-1} \prod d\mu_A det(\not{p} - \not{A} + im) \ exp(S_A). \]

(3)

the vacuum expectation value of Eqs. (2) becomes

\[ <\Psi(x)\Psi(x)> = \frac{1}{N} \sum_k tr((\not{p} - \not{A}_k + im)^{-1}(x,x)) \]

(4)

3 The Valence Approximation

The problem in generating, with present computer hardware, ensembles of \(A\) according to \(dv\) of Eq. (3), and the origin of the valence approximation, is the factor \(det(\not{p} - \not{A} + im)\). For the moment forget about this factor. A simple algorithm to generate an ensemble of \(A\) begins with some arbitrarily chosen field, for example all \(A\) components set to 0, and then successively walks across all sites \(x\) and directions \(\mu\) on the lattice modifying the vector of \(A^j_\mu(x)\), for all \(j\), at each. The collection of \(A\) resulting from each complete sweep updating all \(x\) and \(\mu\) gives a possible Monte Carlo ensemble. The update at each \(x\) and \(\mu\) proceeds in two steps. First a random new trial vector of \(A^j_\mu(x)'\), for all \(j\), is generated according to a rule the details of which I will skip. Then the vector of \(A^j_\mu(x)'\) is either installed in place of the old \(A^j_\mu(x)\) or thrown away with a probability depending on the change \(\Delta S_A\) which this replacement would cause in \(S_A\). Now, the action \(S_A\) approximates the derivatives in \(F^j_\mu(x)\) with nearest-neighbor differences. So computing the change \(\Delta S_A\) arising from the trial replacement of \(A^j_\mu(x)\) by \(A^j_\mu(x)'\) leads to work involving only \(A^j_\mu(y)\) for \(y\) equal to \(x\) or nearest-neighbors of \(x\). Only a fixed number of arithmetic
operations are required independent of the size of the lattice. The total work to update $A^\mu_j(x)$ for all $x$ and $\mu$ to produce a new member of the Monte Carlo ensemble is therefore proportional to $V$, the number of sites in the lattices.

Now put $\det(\hat{p} - \hat{A} + i m)$ back in $d\nu$ in Eq. (5). The algorithm for generating a Monte Carlo ensemble is nearly unchanged, except that in place of the change $\Delta S_A$, we need the change $\Delta S_A + \Delta \log[\det(\hat{p} - \hat{A} + i m)]$. Since $\det(\hat{p} - \hat{A} + i m)$ couples together all $A^\mu_j(x)$ on the entire lattice, the work to find $\Delta S_A + \Delta \log[\det(\hat{p} - \hat{A} + i m)]$ in the update for a single $x$ and $\mu$ is already of order $V$. The work to update all sites becomes of order $V^2$. But recall, for a moment, that predictions for the real world are found from lattice QCD predictions by taking limits of zero lattice spacing and infinite lattice volume. Even a rough approximation to these limits, at least for the simplest lattice implementations of $S_A$ and $S_\Psi$ in Eqs. (1), requires lattice dimensions of at least $10 \times 10 \times 10 \times 10$ thus $V$ of $10^4$. The cost of including $\det(\hat{p} - \hat{A} + i m)$ therefore becomes a factor of $10^4$ or more. Actually, it turns out that somewhat fancier algorithms than what I just described handle $\det(\hat{p} - \hat{A} + i m)$ at lower cost, perhaps a factor of 100 to 1000 for the largest lattices. Unfortunately, present computer power is just barely sufficient to handle lattices large enough to give infinite-volume, continuum limit predictions in the absence of $\det(\hat{p} - \hat{A} + i m)$ in $d\nu$. The increase by a factor of 100 to 1000 in work required to include $\det(\hat{p} - \hat{A} + i m)$ in $d\nu$ is fatal.

What to do? A possible answer is suggested by viewing $\det(\hat{p} - \hat{A} + i m)$ through the eyes of weak coupling perturbation theory. In weak coupling perturbation expansions, $\det(\hat{p} - \hat{A} + i m)$ gives rise to closed quark loops interrupting gluon lines inside diagrams. So $\det(\hat{p} - \hat{A} + i m)$ is in charge of the QCD analogue of the particle-hole polarization process that occurs when an electromagnetic field propagates through a solid. In the case of electromagnetism, we know that for sufficiently weak, sufficiently low momentum electromagnetic fields the effect of particle-hole polarization is accurately approximated by simply replacing all charges and fields by screened values obtained by dividing each by the solid’s dielectric constant. A plausible hypothesis is that the same holds for QCD. Namely, for processes involving sufficiently weak, sufficiently low momentum chromoelectric fields, we can omit $\det(\hat{p} - \hat{A} + i m)$ from $d\nu$ and just replace the chromoelectric charge $g$ in $S_A$ of Eqs. (1) with a screened charge $g/\eta$, for some QCD analogue dielectric constant $\eta$. This replacement is the valence approximation, also called the quenched approximation.

In part as a test of the method, a calculation has been done of the infinite-volume, continuum limit of the valence approximation to light hadron masses. A total of 11 masses were calculated, three of which were used to determine input parameters. The $\pi$ mass was used to set the up and down quark masses,
taken to be equal. The K mass was used to set the strange quark mass. The ρ mass was used to determine the gauge coupling constant. QCD’s supply of free parameters is then exhausted. The remaining eight masses are predictions. The results in units of the ρ mass are shown in Figure 1. The triangles give zero-lattice-spacing results in a box with period of about 2.3 fm, which turns out to be nearly infinite volume for light hadron masses. The circles show infinite-volume predictions obtained by applying an additional correction to the 2.3 fm results. Out of eight numbers, the biggest disagreement between prediction and experiment is 6%. The statistical uncertainties in these numbers range up to 8%, however. For eight numbers with these uncertainties, one disagreement with experiment by 6% is expected. Thus 6% should be treated as the one sigma upper bound on the error in the valence approximation.

So taking the valence approximation as accurate to within 6% for light hadron masses, we applied it to predicting glueball properties.

4 Scalar Glueball Mass

To determine the scalar glueball mass, we evaluated the vacuum expectation value

\[ C(t) = \langle g(t)g(0) \rangle, \]
\[ g(t) = h(t) - \langle h(t) \rangle, \]
\[ h(t) = \int d^3x \ F^0_{ab}(x, t) F_{ab}(x, t), \]
where \(a\) and \(b\) are summed only over space directions, and as before, I actually intend lattice approximations for \(F^a_{\mu b}(\vec{x}, t)\) and \(\int d^3x\) but use the continuum expressions as more recognizable aliases. By inserting a complete set of energy eigenstates between the two glueball operators in \(C(t)\), it is easy to show that for large values of the euclidean time variable \(t\), \(C(t)\) has the asymptotic behavior

\[
C(t) \rightarrow \langle \text{vac}|g(0)|g\rangle < g|g(0)|\text{vac} > e^{exp(-m_g t)} + \ldots ,
\]

where \(|g\rangle\) is the zero-momentum state of the lightest scalar glueball, \(m_g\) is its mass, and the omitted terms come from scalar glueball excitations and fall off exponentially in \(t\) with a coefficient larger than \(m_g\). So for large \(t\), the effective glueball mass \(m_g(t)\) has asymptotic behavior

\[
m_g(t) = \log \frac{C(t)}{C(t+1)} \rightarrow m_g.
\]

Figure 2 shows \(m_g(t)\) in lattice units, for a lattice \(30 \times 32^2 \times 40\) at \(\beta\) of 6.4, obtained from an ensemble of 25440 configurations of chromoelectric field. The lattice spacing in this case, determined from the \(\rho\) mass, is about 0.05 fm and lattice period is about 1.6 fm. For \(t\) of 4 and greater, \(m_g(t)\) is consistent with a constant shown by the horizontal line, giving according to Eq. (6) a value for \(m_g\). From similar calculations for a range of different lattice periods, we found that 1.6 fm gives results within a fraction of a percent of the infinite-volume limit.

To obtain the zero-lattice-spacing limit of \(m_g\), we evaluated \(C(t)\) for five different values the lattice spacing, in all cases with lattice period of 1.6 fm or
greater. Now, the lattice version of the action $S_A$, defined in Eqs. (1), which occurs in a valence approximation calculation of $C(t)$, replaces derivatives by symmetric finite difference, obeying

$$\frac{f(x + a) - f(x - a)}{2a} = \frac{df(x)}{dx} + \frac{a^2}{6} \frac{d^3 f(x)}{dx^3} + \ldots,$$

where $a$ represents lattice spacing. The error in replacing an integral by a sum in $S_A$ goes to zero more rapidly than $a^2$. Thus for small $a$, $S_A$ will be the continuum action plus an error of order $a^2$. To leading order in $a$, $m_g$ measured in physical units will then also be its continuum value plus an error of order $a^2$. Figure 3 shows the five values of $m_g$ in physical units as a function of $a^2$ in physical units. The unit $\Lambda_{\text{MS}}$ in Figure 3 is the valence approximation $\Lambda_{\text{MS}}$, for which a value of $243.7 \pm 6.8$ MeV was obtained in the course of our calculation of quark-antiquark meson masses. To convert $m_g a$ and $a$ to physical units, values of $\Lambda_{\text{MS}} a$ in lattice units were found by the two-loop Callan-Symanzik equation.

As it turns out, the $\rho$ mass in lattice units $m_\rho a$ scales almost perfectly with $\Lambda_{\text{MS}} a$ for the range of parameters in Figure 3. So Figure 3 can also be thought of simply as a plot of $m_g$ and $a$ in units of $m_\rho$ but with the axes mislabeled by powers of $243$ MeV/$m_\rho$.

The three points at smallest lattice spacing in Figure 3 fit a straight line in $a^2$ quite well. Extrapolating to zero lattice spacing gives the limiting $m_g$ of $1740(71)$ MeV, in good agreement with the observed mass of $f_J(1710)$, as mentioned earlier in Section 1. This result is rather insensitive to how the extrapolation is done. Even if we had arbitrarily taken either of the last
two points as the continuum limit, a procedure which is certainly less reliable than extrapolation, the answer would change by less than half of the 71 MeV statistical uncertainty.

5 Decay Couplings

We calculated coupling constants both for the decay of the lightest scalar glueball to pairs of pseudoscalars at rest and to pairs of pseudoscalars with oppositely directed momenta of magnitude \(2\pi/L\) for lattice period \(L\). For the sake of simplicity, I am only going to give details for the decay to pseudoscalars at rest. To measure the decay at rest we evaluated the three-point function

\[
C(t_g, t_\pi) = < g(t_g) \sum_f \pi_f^\dagger(t_\pi) \pi_f(0) >, \tag{9}
\]

\[
\pi_f(t) = \int d^3x \overline{\Psi}_u(\vec{x}, t) \gamma_5 \lambda_f \Psi_d(\vec{x}, t),
\]

where \(g(t)\) is defined in Eqs. (5) and the \(\lambda_f\) are an orthonormal set of SU(3) flavor matrices. The \(u, d\) and \(s\) quark masses were set equal, and chosen so that the energy of the lightest flavor-singlet, two-pseudoscalar, zero-momentum state coincides with the glueball mass \(m_g\). How to apply this calculation to the real world with a very different set of quark masses, I will explain in the next section.

By inserting a complete set of energy eigenstates between \(g(t_g)\) and \(\pi_f^\dagger(t_\pi)\) in \(C(t_g, t_\pi)\), you can show that at large \(t_g - t_\pi\) \(C(t_g, t_\pi)\) has the asymptotic behavior

\[
C(t_g, t_\pi) \rightarrow < vac | g(0) | \pi\pi > < \pi\pi | \sum_f \pi_f^\dagger(t_\pi) \pi_f(0) | vac > \exp(-m_g t_g) + \ldots , \tag{10}
\]

where \(|\pi\pi\rangle\) is the lightest zero-momentum, flavor-singlet state of two pseudoscalars and the terms omitted fall off exponentially in \(t_g\) with coefficients larger than \(m_g\).

The coupling constant we are looking for can be found from \(< vac | g(0) | \pi\pi >\). To obtain \(< vac | g(0) | \pi\pi >\) from \(C(t_g, t_\pi)\) using Eq. (10), we need the value of \(< \pi\pi | \sum_f \pi_f^\dagger(t_\pi) \pi_f(0) | vac >\). This factor carries the corrections to the decay calculation arising from \(\pi - \pi\) final state interactions. We extracted

\[
< \pi\pi | \sum_f \pi_f^\dagger(t_\pi) \pi_f(0) | vac > \text{ from the four-point function}
\]

\[
C(t_{\pi3}, t_{\pi2}, t_{\pi1}) = < \sum_f \pi_f^\dagger(t_\pi) \pi_f(0) \sum_g \pi_g^\dagger(t_\pi) \pi_g(0) >. \tag{11}
\]
The job of evaluating $C(t_{\pi 3}, t_{\pi 2} t_{\pi 1})$ and extracting $< \pi \pi | \sum_f \pi_f^\dagger (t_\pi) \pi_f (0) | \text{vac} >$ I will not describe here. If we had simply taken for $< \pi \pi | \sum_f \pi_f^\dagger (t_\pi) \pi_f (0) | \text{vac} >$ its value assuming no $\pi - \pi$ interaction, however, the final coupling constant would be changed by less than 15%.

As a byproduct of finding the four-point function $C(t_{\pi 3}, t_{\pi 2} t_{\pi 1})$, we were also able to determine the first omitted term in Eq. (10). To accelerate the approach to large $t_g - t_\pi$ asymptopia, we subtracted this term from $C(t_g, t_\pi)$, giving $D(t_g, t_\pi)$. The large $t_g - t_\pi$ asymptotic behavior of $D(t_g, t_\pi)$ is then

$$D(t_g, t_\pi) \rightarrow \lambda K(t_g, t_\pi),$$

(12)

where $\lambda$ is the decay coupling constant for $g \rightarrow \pi \pi$ at rest and $K(t_g, t_\pi)$ is a kinematic factor determined in part from $C(t_{\pi 3}, t_{\pi 2} t_{\pi 1})$. Alternatively, defining from $D(t_g, t_\pi)$ an effective $\lambda(t_g, t_\pi)$,

$$\lambda(t_g, t_\pi) = \frac{D(t_g, t_\pi)}{K(t_g, t_\pi)},$$

(13)

it follows that for large $t_g - t_\pi$, $\lambda(t_g, t_\pi)$ approaches the decay coupling constant $\lambda$.

Values of $\lambda(t_g, t_\pi)$ were calculated on a lattice $16^3 \times 24$ at $\beta$ of $5.70$, corresponding to lattice spacing about $0.15$ fm and lattice period of about $2.3$ fm, from an ensemble of $10500$ configurations of $A$ field. For $t_g - t_\pi$ fixed, we found $\lambda(t_g, t_\pi)$ independent of $t_\pi$ for $t_\pi \geq 3$. Figure 4 shows $\lambda(t_g, t_\pi)$, for $t_\pi \geq 3$, as a function of $t_g - t_\pi$. Largely as a consequence of the subtraction in defining $D(t_g, t_\pi)$, the data is consistent with a constant from $t_g - t_\pi$ of zero on out. This constant is the final value of $\lambda$ for glueball decay to pseudoscalars at rest.

### 6 Comparison with Experiment

So far I have given u, d and s quarks degenerate, unphysical mass values. Here is how to fix that. An expansion to first order in the quark mass matrix taken around some relatively heavy SU(3) symmetric point gives glueball decay couplings to $\pi$'s, K's and the $\eta$'s which are a common linear function of each meson's average quark mass. Since meson masses squared are also nearly a linear function of average quark mass, the decay couplings are a linear function of meson masses squared. Therefore, a linear fit to our predictions for decay couplings as a function of pseudoscalar mass squared lets you extrapolate from unphysical degenerate values of quark masses to physical nondegenerate values of quark masses. Figure 5 shows predicted coupling constants as a function of predicted meson mass squared along with a linear extrapolation of
Figure 4: $\lambda$ as function of $t_g - t_\pi$.

the predicted values to the physical $\pi$, K and $\eta$ masses. Shown also are the observed couplings for decays of $f_J(1710)$ to pairs of $\pi$'s, K's and $\eta$'s. Everything is in units of the $\rho$ mass. The total predicted width for glueball decay to pseudoscalar pairs becomes 108(28) MeV, in comparison to 99(15) MeV for $f_J(1710)$.

How far are the valence approximation, finite lattice spacing decay couplings likely to be from the real world? From the comparison of finite lattice spacing valence approximation hadron masses with their values in the real world, I would expect an error of 15% or less in going to the continuum limit and another 6% or less arising from the valence approximation. The total predicted width for glueball decay to two pseudoscalars should then have an error of less than 50%. A 50% increase in our predicted two-body decay width, combined with any reasonable corresponding guess for multibody decays, gives a total glueball width small enough for the particle to be observed easily.

For the continuum limit glueball mass, a 6% valence approximation error would be 100 MeV, but according to an adaptation of an argument giving a negative sign for the valence approximation error in $f_\pi$, the sign of this error is also expected to be negative. Thus the only established resonance aside from $f_J(1710)$ with the correct quantum numbers and mass close enough to 1740 to be a candidate for the scalar glueball is $f_0(1500)$. The most likely interpretation of $f_0(1500)$, however, I think is as an $s\bar{s}$ quark-antiquark meson. The $s\bar{\pi}$ scalar and tensor are nearly degenerate at about 1430 MeV. So the $s\bar{s}$ scalar and tensor should lie close to each other somewhere above 1430 MeV. The $s\bar{s}$ tensor has been identified at 1525 MeV. An $s\bar{\pi}$ scalar around 1500 MeV...
Figure 5: Decay couplings.

seems to me hard to avoid.

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