Independent Operators at Different Dimension

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Abstract
To apply lattice QCD in the calculation of glueball spectrum it is needed firstly to know associated operators acting on vacuum. We show how to find all the independent representations and operators, of group $SO(3)^{PC}$ at different dimension, since the work is not trivial. Then, we decompose these representation into irreducible representation of $O^{PC}$ group, which are listed in the note. At last we argue that $f_J(2220)$ and $g_T$ states can not be tensor glueball simultaneously.
This is a note when we calculate the glueball spectrum in quenched lattice QCD. We are focusing on the independent currents at each dimension here.

As we know, in path quantization, an arbitrary bounded particle, especially, glueball, with $J^{PC}$, where $J$ is spin of glueball, or total angular momentum of the composite particle, $P$ and $C$ are quantum number of parity and charge conjugate respectively, can be considered as a wave functional, which is an eigen-functional of Hamiltonian in QCD. With respect to the basis of eigenvalue of canonical fields, $A_i = A^{a}_i T_a$ and therefore $B_i = B^{a}_i T_a = \frac{1}{2} \epsilon_{ijk} F_{jk} = \frac{1}{2} \epsilon_{ijk} (\partial_j A_k - \partial_k A_j - i[A_j, A_k])$, a state wave functional $|\psi\rangle$, particularly, glueball state with $J^{PC}$, is generated by corresponding current acting on the vacuum functional $|0\rangle$,  
\begin{equation}
|\psi\rangle = O |0\rangle,
\end{equation}
where $O$ is gauge invariant operator with $J^{PC}$. In free QED, one can calculate operator $O$ for each state. However, in QCD, one can not, due to the non-perturbational and nonlinear behavior of QCD. Operator $O$ can be written as  
\begin{equation}
O = \sum_{n \geq 4} \sum_{i} c^{n}_{i} O^{(n)}_{i},
\end{equation}
where $O^{(n)}_{i}$ is the $i$'th operator with $J^{PC}$ in dimension $n$, and $c^{n}_{i}$ is coefficient to solve. Generally, if we set single or only finite $c^{n}_{i}$ nonvanishing, operator $O$ acting on vacuum functional will generate infinite eigenstates (physical states with definite energy) with the same $J^{PC}$, or vice versa, there are infinite nonvanishing $c^{n}_{i}$ for each eigenstate.

Therefore, to extract a definite glueball, one needs infinite operators $O^{(n)}_{i}$'s and calculates infinite coefficients $c^{n}_{i}$'s. This is an impossible mission, since in practice we can only use finite operators $O^{(n)}_{i}$'s. However, thanks to functional analysis, using these finite operators we can always obtain optimal approximation to the ultimate operator $O$ for definite state. Although eigen-currents under optimal approximation also generate infinite states, we expect the projections of these currents on heavier states are small and we can neglect contributions of those heavier states.

According to functional analysis, to extract a definite glueball, for example, quintet states $2^{++}$, we take such steps. Step one, one writes independent $2^{++}$ operators dimension by dimension. In practice, the operators are surely truncated at some dimension, for instance, dimension seven or dimension eight. Step two, since we are working on lattice, the rotation symmetry is broken into cubic symmetry. We thus decompose this operators into different irreducible representation of $O^{PC}$ group. For instance, we decompose $2^{++}$ operator into $E^{++}$ and $T^{++}_2$ representation. Step three, we represent these operators on lattice by Wilson loops, the continuum limits of which are just the corresponding operators. Step four, for instance, for those nonequivalent $E^{++}$ irreducible representations belonging to different $2^{++}$ representations in $SO(3)^{PC}$ group, we use variation principle to extract optimal approximations to different glueball $2^{++}$, at finite lattice spacing $a$. The last step, of course, we extrapolate the results into continuum limit.

This method has constraints. If we use three independent operators, which belong to three different representations of $2^{++}$, we at most extract three different $2^{++}$ glueballs. Secondly, these operators should be independent, otherwise we shall meet singularity in the utilizing of variation principle. Therefore, to find all the independent currents at each dimension is important on lattice, since the complete basis of glueball states is generated by all the gauge independent currents, which, as shown, are made up of chromomagnetic fields, $B_i$'s, and
covariant derivatives, $D_i$'s, where subscripts stand for directions and superscripts are color indices. This is just the topic of this note.

On one hand, for static glueball states one should sum operators over all spatial points on lattice to construct corresponding current. On the other hand, these operators should be local and gauge invariant products of $B_i$ and $D_i$ at the same time. For instance, $\sum_{\Lambda} B^a_i B^a_i$ generates 0++ (infinite) glueballs when it acts on vacuum. The summation over spatial points will be suppressed thereinafter.

For this aim we first need to know the behavior of $B^a_i$ and $D^{bc}_i$ under rotation. The rotation transformation is divided in two parts, one is the rotation in orbital space, and the other is the rotation in spin space. $B^a_i$ is vector operator under rotation transformation in spin space and scalar operator under rotation transformation in orbital space. Therefore, operator $B^2_i$ can be interpreted as spinor with unit spin. Whereas, $D^{bc}_i = \partial_i \delta^{bc} - i A^a_i t^a$, where $(t^b)_{ac} = i f^{abc}$ is generator in adjoint representation, is the vector operator only under total rotation transformation. This is because that the first term in covariant operator, $\partial_i \delta^{bc}$, is vector under orbital rotation and scalar under spin rotation, while the second term in covariant operator, $-i A^a_i t^a$, has a reverse character, for it is scalar under orbital rotation and vector under spin rotation. Thus, the covariant derivative operator can not be interpreted as orbital angular momentum vector which ascends or descends unit orbital angular momentum. But, if we ignore the term $-i A^a_i t^a$, it can. This is nearly true if the glueball is sufficient small and gauge fields are smooth enough, for one can always let $A^a_i(x) \equiv 0$ at one point, for instance the center of glueball mass, by performing a local gauge transformation. Attributed to this reason we sometimes call covariant derivative as orbital vector, although this will take many ambiguities.

It seems firstly that finding independent current at each dimension is a trivial mission. However, due to the so-called Bianchi identity, one always meets some subtleties. We shall show how to treat these subtleties by an example, i.e. finding all the independent operators at the case of dimension six. Such treatment can be generalized into other dimension directly.

All the operators at dimension six can be categorized into two types: type one is as $Tr((D_i D_j B_k) B_l)$ and type two is as $Tr(B_i B_j B_k)$. $Tr((D_i B_j)(D_k B_l))$ is not an independent type, for operators in such type can be rewritten as the form of type one up to a complete covariant derivative, which will be discarded at last.

We first consider operators in type one. Taking advantage of standard basis of chromo-magnetic fields $B_{\pm} = \mp(B_1 \pm i B_2)/\sqrt{2}$ and $B_0 = B_3$, we find the total spin of $B_k B_l$ can be 0, 1, 2. Similarly, the standard basis $D_{\pm} = \mp(D_1 \pm i D_2)/\sqrt{2}$ and $D_0 = D_3$ makes that the total ‘orbital angular momentum’ of operators $D_i D_j$ be 0, 1, 2. It seems that there are 81 operators. But, this is not true.

First, we rewritten such operator using a technique of symmetry-antisymmetry decomposition,

$$Tr((D_i D_j B_k) B_l) = \frac{1}{2} Tr\{[(D^s_{ij} B_k) B_l] + [(D^s_{ij} B_k) B_l] - (D^s_{ij} B_k) B_l]\} + [(D^s_{ij} B_k) B_l] + [(D^s_{ij} B_k) B_l] - (D^s_{ij} B_k) B_l]\}, \quad (3)$$

where $D^s_{ij} = \frac{1}{2}(D_i D_j + D_j D_i)$ and $D^a_{ij} = \frac{1}{2}(D_i D_j - D_j D_i)$. This technique is crucial. It reduces all the operators into the forms of symmetry, antisymmetry and admixture symmetry-antisymmetry. In each form all the operators at the same dimension make up of an invariant space of group $SO(3)^PC$. Eq. (3) tells us that operators in type one are divided into four
forms. But, terms in the last two square bracket are in fact operators in type two because of the identity \([D_i, D_j] = [F_{ij}, \cdot]\). (There is in fact only one independent nonvanishing operator \(\epsilon_{ijk} Tr(B_i B_j B_k)\) in the form \(Tr((D^2_b s B_k) B_l)\).) It is easy to check the operator in the second square bracket is vanishing up to a complete covariant derivative. Thus, for operator belonging to type one it is enough to consider the operators as the form shown in the first square bracket. Therefore, for operator of type one we only consider the case that the total spin and total orbital angular momentum are 0 or 2. There are total 36 operators to be considered now.

The \(J^{PC}\) of the total operators in type one are summarized in table 1, then,

| spin \ orbital | 0       | 2       |
|---------------|---------|---------|
| 0             | 0++(5S_0) | 2++(5D_2) |
| 2             | 2++(5S_2) | 0++(5D_0), 1++(5D_1), 2++(5D_2), 3++(5D_3), 4++(5D_4) |

Table 1: The \(J^{PC}\) of the total operators in type one. The state symbols in bracket are obtained by interpreting covariant derivative as orbital angular momentum vector. Surely such interpretation has ambiguities.

Now we turn to operators in type two. Type two can be divided into two forms under symmetry-antisymmetry decomposition, \(Tr(B_i (B_j B_k - B_k B_j))\) and \(Tr(B_i (B_j B_k + B_k B_j))\). In the first form there is only one independent operator, \(\epsilon_{ijk} Tr(B_i B_j B_k)\), which is 0++. This operator can also be regarded as the one in type one, as shown before. We pay attention on the second form here, the \(PC\) of which is +−. For the case \(i \neq j \neq k\), there is one independent operator, \(Tr(B_1 B_2 B_3 + B_2 B_3 B_1)\), which belongs to 3−+ in cubic group and therefore belongs to \(3^-\) in \(SO(3)^{PC}\) group concerning to the total number of operators, ten, in form two. There are three operators at the case \(i = j = k\), which construct to the representation \(T_i^{+−}\). For the case \(i \neq j = k\), there are six operators. Therefore, there are total eleven independent operators in type two. As argued, one of them belongs to 0++, and seven of them belong to 3++. The remnant three operators belong to 1+−, since three operators \(Tr(B_i (B_1 B_1 + B_2 B_2 + B_3 B_3))\)’s make up of the representation 1+− of the \(SO(3)^{PC}\) group.

We then have operators belong to three nonequivalent representations 0++ and 2++, one representation 1++, 3++, 4++, 1+− and 3+− respectively.

It seems the topic is finished, but the game is not so simple, due to the Bianchi identity, \(I = D_1 B_1 + D_2 B_2 + D_3 B_3 \equiv 0\).

Now we consider the constraint introduced by Bianchi identity. Bianchi identity is a scalar operator. Then, the operators containing Bianchi identity at dimension six are as the form \(Tr(B_i D_j I)\), i.e. they belong to type one and therefore, their \(PC\) are ++. Since \(B_i D_j\) can compose representation \(J = 0, 1\) and \(2\) under rotation transformation, we conclude that operators containing Bianchi identity can compose representations 0++, 1++ and 2++ under \(SO(3)^{PC}\) group. We emphasize here that for general case, we need also use symmetry-antisymmetry decomposition to study \(J^{PC}\)’s of the operators containing Bianchi identity, although at the case of dimension six or five we need not.

Therefore, operators belonging to 1++ are vanishing. The three representations 0++ are not independent and there are only two independent nonequivalent representations 0++. It is similar for 2++. For two independent representations 0++, we pick on states \(^1S_0\) and \(\epsilon_{ijk} Tr(B_i B_j B_k)\), while for 2++, states \(^1D_2\) and \(^5S_2\) respectively.
This statement can be checked by a directly calculation. For instance, suppose $\alpha_2$ is the operator of state $^5D_0$ up to a factor,

$$\alpha_2 = \text{Tr}\{ -B_1(-2D_1^2 + D_2^2 + D_3^2)B_1 - B_2(-2D_2^2 + D_1^2 + D_3^2)B_2 - B_3(-2D_3^2 + D_1^2 + D_2^2)B_3 \\
+ 3B_2(D_2D_1 + D_1D_2)B_1 + 3B_1(D_3D_1 + D_1D_3)B_3 + 3B_2(D_3D_2 + D_2D_3)B_3\},$$  \hspace{1cm} (4)

and $\alpha_0$ is the operator of state $^1S_0$ up to a factor,

$$\alpha_0 = \text{Tr}\{ B_1(D_1^2 + D_2^2 + D_3^2)B_1 + B_2(D_1^2 + D_2^2 + D_3^2)B_2 + B_3(D_1^2 + D_2^2 + D_3^2)B_3\},$$  \hspace{1cm} (5)

one has $\alpha_2 + \alpha_0 \equiv 0$ if he takes advantage of Bianchi identity.

The analysis for operators at dimension six can be generalized into other dimension. For instance, operators at dimension 5 belong to representations $0^{-+}$ and $2^{-+}$ respectively, and operators at dimension 4 belong to representations $0^{++}$ and $2^{++}$ respectively. Operators at dimension five has no nonvanishing representation $1^{-+}$ due to Bianchi identity.

In summary, all the operators at dimension four, five and six can be categorized into three independent representations $0^{++}$ and $2^{++}$, one representation $3^{++}$, $4^{++}$, $1^{-+}$, $3^{-+}$, $0^{-+}$ and one representation $2^{-+}$. One can make use of C-G coefficients to obtain these operators.

These representations are irreducible in $SO(3)^{PC}$ group, but they are reducible representations in group $O^{PC}$. Since on lattice we only have $O^{PC}$ symmetry, we shall furthermore reduce these representations into irreducible representations of group $O^{PC}$. The results are very lengthy, and we only list some of them, which will be used in our lattice simulations. In each irreducible representation only one operator is shown. This is enough even for representations $E^{PC}$, $T^{PC}_1$ and $T^{PC}_2$, because as long as one operator is found, other operator(s) belonging to the same representation can easily be obtained by suitable cubic transformations.

- Three $0^{++}$ representations:

1. The first representation $0^{++}$, which is of four dimension,

$$a_1 \propto \text{Tr}(B_1^2 + B_2^2 + B_3^2);$$  \hspace{1cm} (6)

2. The second representation $0^{++}$, which is of six dimension,

$$a_1 \propto \text{Tr}[B_1(D_1^2 + D_2^2 + D_3^2)B_1 + B_2(D_1^2 + D_2^2 + D_3^2)B_2 + B_3(D_1^2 + D_2^2 + D_3^2)B_3];$$  \hspace{1cm} (7)

3. The third representation $0^{++}$, which is of six dimension,

$$a_1 \propto \text{Tr}[B_1(B_2B_3 - B_3B_2)].$$  \hspace{1cm} (8)

- Three $2^{++}$ representations:

1. The first representation $2^{++}$, which is of four dimension, is decomposed into irreducible representation $E^{++}$ and $T_2^{++}$. For representation $E^{++}$,

$$e^1 \propto \text{Tr}(B_1^2 - B_2^2),$$  \hspace{1cm} (9)

while for $T_2^{++}$,

$$t_2^1 \propto \text{Tr}(B_2B_3);$$  \hspace{1cm} (10)
2. The second representation $2^{++}$ is of six dimension. For representation $E^{++}$,

$$e^1 \propto Tr[B_1(D_1^2 + D_2^2 + D_3^2)B_1 - B_2(D_1^2 + D_2^2 + D_3^2)B_2],$$

while for $T_2^{++}$,

$$t_2^1 \propto Tr[B_2(D_1^2 + D_2^2 + D_3^2)B_3];$$

3. The third representation $2^{++}$ is of six dimension. For representation $E^{++}$,

$$e^1 \propto Tr[B_1(D_1^2 - D_2^2)B_1 + B_2(D_1^2 - D_2^2)B_2 + B_3(D_1^2 - D_2^2)B_3],$$

while for $T_2^{++}$,

$$t_2^1 \propto Tr[B_1D_3D_2B_1 + B_2D_3D_2B_2 + B_3D_3D_2B_3].$$

- Representation $4^{++}$ only occurs in dimension six. This representation is reduced into $A_1^{++}$, $E^{++}$ and two $T_2^{++}$'s. For representation $A_1^{++}$,

$$a_1 \propto Tr[B_1(2D_1^2 - D_2^2 - D_3^2)B_1 + B_2(-D_1^2 + 2D_2^2 - D_3^2)B_2 + B_3(-D_1^2 - D_2^2 + 2D_3^2)B_3 - 2B_2(D_2D_1 + D_1D_2)B_1 - 2B_1(D_3D_1 + D_1D_3)B_3 - 2B_2(D_3D_2 + D_2D_3)B_3],$$

and for representation $E^{++}$

$$e^1 \propto Tr[B_1(D_1^2 - D_2^2)B_1 - B_2(D_2^2 - D_3^2)B_2 - B_3(D_1^2 - D_2^2)B_3 - 4B_1D_3D_1B_3 + 4B_2D_3D_2B_3].$$

- Representation $3^{++}$, which is reduced into $A_2^{++}$, $T_1^{++}$ and $T_2^{++}$, only occurs in dimension six. For representation $A_2^{++}$,

$$a_2 \propto Tr[B_1(D_2^2 - D_3^2)B_1 + B_2(D_3^2 - D_1^2)B_2 + B_3(D_1^2 - D_2^2)B_3],$$

and for representation $T_1^{++}$,

$$t_1^1 \propto Tr[-B_2D_2D_3B_2 + B_3D_2D_3B_3 + 2B_2D_1D_3B_1 - 2B_1(D_2D_1 + D_1D_2)B_3 + B_2(D_2^2 - D_3^2)B_3].$$

- We meet representation $0^{-+}$ in dimension 5,

$$a_1 \propto Tr[B_2D_1B_3 + B_3D_2B_1 + B_1D_3B_2].$$

- Representation $2^{-+}$ is met in dimension 5. For $E^{-+}$,

$$e^1 \propto Tr[B_2D_1B_3 + B_1D_2B_3],$$

and for $T_2^{-+}$

$$t_1^1 \propto Tr[B_1D_3B_3 + B_2D_2B_1].$$

- Representation $1^{+-}$ at dimension 6 is reduced into irreducible representation $T_1^{+-}$,

$$t_1^1 \propto Tr[B_1B_1^2 + B_1B_2^2 + B_1B_3^2].$$
• Representation $3^{+-}$ occurs in dimension six. We show the operator in $A_2^{++}$,

$$a_2 \propto Tr[B_2B_1B_3 + B_2B_3B_1].$$  \hspace{1cm} (23)

Now we have listed the results of decomposing twelve independent representations into irreducible representations in $O^{PC}$ group. If we use variation principle we obtain twelve glueballs. These glueballs, except excited scalar and tensor glueballs, have been measured by other authors. The listed results tell us the structures of different glueballs, which is useful, for instance, in glueball decay or in condensate of different current in glueballs.

Although we do not make any lattice simulation here, we can obtain some physical conclusions from the results, one of which will be shown in the next paragraph.

To apply variation principle we can obtain 12 glueballs, which are in fact the optimal approximation of the real glueballs. Attributed to dimension analysis, it is reasonable to regard $4^{++}$ glueball is the heaviest glueball with $PC = ++$ among those 12 glueballs. As pointed out by previous works\[^3\], the mass of which is about $m_{4^{++}} = 3 \sim 4 GeV$. We have shown elsewhere that it is enough to use chromomagnetic fields and their covariant derivative in the calculation of glueball spectrum merely. Therefore, there are three and only three $2^{++}$ states lower than $m_{4^{++}}$. One possibly debates that their masses will be corrected by correction from fermion-loop and from higher-dimension operators. On one hand, to consider the correction of fermion-loops we should go into full QCD. However, it is believed that the quenched approximation almost does not change the ratio between different states. We think the mass correction from fermion-loop is mainly on scale shift. On the other hand, we shall get decreased glueball masses on lattice if we consider the correction from operators at higher dimension. Such decrease is both on $4^{++}$ glueball and on $2^{++}$ glueballs. We conclude that mainly effect of higher dimension operator is also a scale shift. Therefore, the statement, there are three and only three $2^{++}$ states lower than $m_{4^{++}}$, is right or almost right even including the two corrections, possibly up to a scale shift. If $f_{J}(2220)$ (or $\xi (2230)$) is a tensor glueball, one must find other two tensor glueballs with mass lower than $m_{4^{++}}$. But, since there are only three tensor glueballs below $m_{4^{++}}$, it is difficult to place three $GT$ states, which are more like tensor glueballs. However, if $g_T$ are interpreted as three glueballs, $f_{J}(2220)$ should never be tensor glueball. The tensor glueball interpretations of $f_{J}(2220)$ and the tensor glueball interpretation of $g_T$ are conflictive interpretations, therefore. Since there are more supports of the tensor glueball interpretation of $g_T$, at least in nowadays, $f_{J}(2220)$ can not be tensor glueball, if it exists.

There are quenched lattice gauge calculation supporting the tensor glueball interpretation of $f_{J}(2220)$, for instance, reference\[^3\] claims that masses of ground scalar and ground tensor glueball are 1.73$\text{Gev}$ and 2.40$\text{Gev}$ respectively. The scale parameter is determined by phenominal Sommer potential\[^5\] in these works, that is, $r_0$ is determined by $r^2 \frac{dV}{dr}|_{r=0.49 fm} = 1.65$, since there is no exclusive experiment result on glueball mass. However, $r_0$ depends on different phenonmal potential models, for instance, as author pointed out in reference\[^3\], the logarithmic potential gives values that are lower by about 10%. Combining the argument in the last paragraph, we think that there may exist a global scale shift in these works. If this is the case and the ground tensor glueball is $g_T(2050)$, the mass of scalar glueball should be $1.73 \times 2.40 = 4.16 \text{Gev} = 1.48 \text{Gev}$, which looks very like $f_0(1500)$. If the ground scalar glueball is $f_0(1500)$, the dominant component should be $s\bar{s}$ in $f_0(1710)$. This interpretation is favored in many phenomenon analysis. For instance, we obtain a more natural conclusion that the $s\bar{s} - n\bar{n}$ mass difference is about 300$\text{MeV}$\[^6\]. In fact, a more recent work\[^7\] reveals that one
can not distinguish the scalar glueball interpretation of $f_0(1500)$ and $f_0(1710)$ nowadays in unquenched lattice QCD. From the above argument we also obtain the mass of $1^{--}$ glueball should be $3.85 \times \frac{2.25}{2.40} Gev = 3.29 Gev$. This means that $1^{--}$ glueball mass and $J/\psi$ mass are very close. Therefore, taking account into the lattice computing error, to solve the "$\rho - \pi$ puzzle" of $J/\psi$ and $\psi'$ decays, the suggestion\cite{8,9} that $J/\psi \rightarrow \rho \pi$ is enhanced by a mixing of the $J/\psi$ with an $1^{--}$ glueball that decays to $\rho \pi$ is also suitable.

Unfortunately, in simulations we found that the mass of first excited tensor glueball and the mass of second excited tensor tensor glueball are very close. Such unexpected occasion makes the traditional approach of variation principle, illustrated by C. Morningstar\cite{4}, has disadvantages in the calculation. We will show an alternative method in the future.

In summary, we have finished the first two steps in the calculation of glueball spectrum. The leavings will be shown in the future. Although we do not make any lattice simulation here, we obtain an interesting conclusion, that is, $f_J(2220)$ and $g_T$ states can not be tensor glueball simultaneously.

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