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New Approach to Constructing the Operator on the Lattice for the Calculation of the Glueball Masses*

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Abstract We develop a new approach to constructing the lattice operators for the calculation of the glueball mass, which is based on the connection between the continuum limit of the chosen operator and the quantum number J^{PC} of the state. The spin of the state is then determined uniquely and directly in numerical simulation. Furthermore, the approach can be applied to the calculation of the mass of glueball states with any spin, J. Under the quenched approximation, we present our preliminary results in $SU(3)$ pure gauge theory for the mass of 0^{++} state and 2^{++} state, which are 1754(85)(86) MeV and 2417(56)(117) MeV, respectively.

Key words lattice QCD, glueball mass, Wilson operator

1 Introduction

During the past two decades, there have been extensive lattice calculation of the glueball spectroscopy in Lagrangian scheme[1-5] or in Hamiltian scheme[6]. Most of previous works use two key steps: one is the choice of glueball operators with certain quantum number J^{PC} based on the method introduced in Ref. [1] and the other is the application of variational principle. Meanwhile, with a great amount of the improvement, such as fuzzying and smearing, etc., these approaches work well and the errors are under well controlled. The choice of lattice operators[1-5] is based on the correspondence of the irreducible representation R of the cubic point group and the representation J of the rotation group. However, since the correspondence between R and J is not one-to-one, there exist some ambiguities in the choice.

Meanwhile, basing on the representation theory of $O(4)$ group and the hypercubic group, Mandula et al.[7] develop an elegant scheme for the choice of glueball operators. Decomposing the lattice color electric and magnetic fields into certain representations of the hypercubic group, they construct operators with definite J^{PC}. However, the correspondence between irreducible representation of the hypercubic group and spin J is also not one-to-one. The 'leading spin' is then assumed when $a \to 0$. Even so, this assumption also cannot determine the spin uniquely. For example, one does not know how to separate 'leading spin' $J = 1^-$ from 'leading spin' $J = 2^+$ in $6^{(+)}$ representation and one can not get the content of the non-leading spin[7].

In this paper, we would like to show a possible solution to these troubles. By expanding the chosen operator according to power of lattice spacing a, we require that the leading term of the

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chosen operator belongs to the irreducible representation J^{PC} of $SO(3)^{PC}$ group. We assume that the leading term will give the main effect to the state when $a \to 0$, and the contribution should be only given by the leading term in the continuum. Therefore, the spin of the corresponding state is uniquely determined by the leading term of the expansion when the lattice tends to continuum.

Some observations are shown in the forthcoming section. We introduce our method in Sect. 3 and give an example of some preliminary results to verify our method in Sect. 4. Sect. 5 is a short summary.

2 Some Observations

In the continuum, the glueball states with definite quantum number J^{PC} make up of the basis of certain irreducible representation J^{PC} of $SO(3)^{PC}$ group. But, on the lattice, there only exists its finite point subgroup, O^{PC}, and its corresponding irreducible representations R^{PC} ($R = A_1, A_2, E, T_1$ and T_2). Then to measure glueball mass in lattice QCD, there arises such a problem as how we get correct results by only utilizing O^{PC} group. To solve the problem, following Berg and Billoire[1], authors make the following continuum limit assumption ($\beta = \infty$)[1,2,4,5]:

$$\begin{aligned} &m(0^{PC}) = m(A_1^{PC}), \qquad m(1^{PC}) = m(T_1^{PC}), \\ &m(2^{PC}) = m(E^{PC}) = m(T_2^{PC}), \quad m(3^{PC}) = m(A_2^{PC}), \end{aligned} \tag{1}$$

where $m(R^{PC})$ is the mass of the state extracted from operators in the irreducible representation R^{PC} on the lattice and $m(J^{PC})$ is the mass of state with certain spin J^{PC} in the continuum.

But as Morningstar and Peardon show[3], this assumption is not always right. From their simulation results, for example, T_1^{++} channel is not interpreted as $J^{PC} = 1^{++}$ states but most likely as $J^{PC} = 3^{++}$ state (less likely $J = 6, 7, 9, \cdots$, interpretation cannot be ruled out), since it seems this channel and A_2^{++} are degenerated in the continuum.

Meanwhile, on a $D = 2 + 1$ lattice, Johnson and Teper[8] find that in A_1^{++} channel there exist two states with different masses. They interpret the higher one as the 4^{++} state and the lower one as 0^{++} state. Therefore, they also believe that one needs to develop a systematic and general procedure to construct operators of arbitrary spin as $a \to 0$[8].

Now, we present a possible procedure to solve these problems here. Let us begin the discussion with some observations.

1) An arbitrary state $|\psi\rangle$ can be generated by the current o acting on vacuum $|0\rangle$:

$$|\psi\rangle = o|0\rangle. \tag{2}$$

Since $|0\rangle$ is invariant under Poincare group and $SU_C(3)$ group, the character of $|\psi\rangle$ can be described by o. For simplicity, we only consider currents with mass dimension 4 here, saying $B_i^a(x)\ B_j^b(x)$, where B_i^a are the color magnetic fields and $i, j = 1, 2, 3$.

Obviously, both $|\psi\rangle$ and o are color singlet. One gets such 6 color singlet currents:

$$2\mathrm{Tr}(B_iB_j) = \sum_{a=1}^{8} B_i^a B_j^a,$$

where $B_i = \sum_{a=1}^{8} B_i^a \frac{\lambda^a}{2}$.

We also require $|\psi\rangle$ and o transform as certain representation J^{PC} under $SO(3)^{PC}$ group. Since $\boldsymbol{B}$ transforms as 1^{+-} under this group, the P, C of $\mathrm{Tr}(B_iB_j)$ are $++$. Using Clebsch-Gordan coefficients, we can decompose these basis into $J = 0$ and $J = 2$ pieces (due to the color sin-

glet, there is no basis with $J=1$)[1]. Then, in the subduced representation $J \downarrow O$ of the rotation group $SO(3)$ restricted to subgroup O, we find that the basis in $J=0$ is the basis of representation A_1, and we can further reduce another five basis in $J=2$ according to irreducible representations E and T_2 of the cubic point group. So, we can categorize the basis as

$$J=0: \quad a_{11} = \mathrm{Tr}(B_1B_1 + B_2B_2 + B_3B_3); \tag{3a}$$

$$J=2: \quad e_1 = \mathrm{Tr}(B_1B_1 - B_2B_2), e_2 = \mathrm{Tr}(B_1B_1 + B_2B_2 - 2B_3B_3); \tag{3b}$$

$$t_{21} = \mathrm{Tr}(B_2B_3), t_{22} = \mathrm{Tr}(B_1B_3), t_{23} = \mathrm{Tr}(B_1B_2). \tag{3c}$$

Here a_{11} is the basis of representation A_1, e_1 and e_2 construct basis of representation E, while t_{21}, t_{22} and t_{23} just make up of basis of representation T_2.

This is just what Table 2 in Ref. [1] tells us: the subduced representation $J=2$ of the rotation group can be decomposed into representation E and T_2 in the cubic group; the subduced representation $J=0$ is just representation A_1.

We can do similar analysis for higher mass-dimensional gauge invariant operators consisting of color magnetic fields and its covariant derivatives.

2) Now, we consider how to construct the glueball operator. By expanding the chosen operator according to the power of spacing a, we require that the leading term of the chosen operator belongs to and only belongs to certain irreducible representation J^{PC} of $SO(3)^{PC}$ group. A simple example is for the plaquette operator

$$O_{ij} = \sum_n O_{ij}(n) = \sum_n \mathrm{Tr}[1 - U(n,i)U(n+\hat{i},j)U^{-1}(n+\hat{j},i)U^{-1}(n,j)]. \tag{4}$$

The link variable U is a connector defined by

$$U(n,i) = P\exp\left(i\int_0^a \mathrm{d}tA_i(an + a\hat{i}t)\right), \tag{5}$$

where $\hat{i}$ is the i-th positive direction and P is the path-order operator.

There are two methods to expand the operator. One is the application of non-Abelian Stokes theorem[10,11] and another method is introduced by Luscher and Weisz in Ref. [12]. We find that both methods lead to the same results:

$$O_{ij} = \sum_n \left\{\frac{a^4}{2}\mathrm{Tr}(F_{ij}F_{ij})(n) + \frac{\mathrm{i}a^6}{6}\mathrm{Tr}(F_{ij}F_{ij}F_{ij})(n) - \frac{a^6}{24}\mathrm{Tr}[F_{ij}(D_i^2 + D_j^2)F_{ij}](n)\right\} + O(a^8), \tag{6}$$

where $F_{ij} = \partial_i A_j - \partial_j A_i - \mathrm{i}[A_i, A_j]$ is the field strength tensor and $D_i \cdot = \partial_i \cdot - \mathrm{i}[A_i, \cdot]$ is the covariant derivative.

We now consider the $PC = ++$ sector of operators, or real part in Eq. (6) with ignoring the second term of r.h.s.. Due to $O_{ji} = O_{ij}^*$, there are three non-zero independent operators $\mathrm{Re}O_{12}$, $\mathrm{Re}O_{23}$, $\mathrm{Re}O_{31}$. Restricting oneself into the cubic group, one can combine these operators into representation A_1^{++} and E^{++}:

$$A_1^{++}: \mathrm{Re}(O_{23} + O_{13} + O_{23}) = \frac{a^4}{2}\sum_n \mathrm{Tr}(B_1B_1 + B_2B_2 + B_3B_3)(n) + O(a^6);$$

$$E^{++}: \mathrm{Re}(O_{23} - O_{13}) = \frac{a^4}{2}\sum_n \mathrm{Tr}(B_1B_1 - B_2B_2)(n) + O(a^6),$$

1) The systemic decomposition is well discussed by Jaffe et al [9] in the study of the qualitative features of the glueball spectrum. They suggest to construct glueball operators for certain J^{PC} states with color magnetic and electric fields in the continuum case.

$$\mathrm{Re}(O_{23} + O_{13} - 2O_{12}) = \frac{a^4}{2}\sum_n \mathrm{Tr}(B_1B_1 + B_2B_2 - 2B_3B_3)(n) + O(a^6), \qquad (7)$$

where the color magnetic field is $B_i = -\frac{1}{2}\sum_{jk}\varepsilon_{ijk}F_{jk}$.

We suppose that the leading term gives the most contribution of the operator when a is small enough, or, only the leading term gives the contribution in the continuum. While comparing Eq.(7) with Eq.(3), it is assured that, in the continuum limit, the state extracted from such operator A_1^{++} is $J^{PC} = 0^{++}$, and the state extracted from the operator E^{++} corresponds to $J^{PC} = 2^{++}$ state.

We should emphasis again, in the general case, the continuum limit of the operator in representation E or in T_2 is not always corresponding to $J = 2$, i.e., the parallelism in Eq.(1) does not always hold. Only after expanding the chosen operator as we do above, we are then able to affirm or disaffirm the parallelism.

By the way, we should point out here that the non-leading terms in the expansion of the operator do not always belong to the same J^{PC} as that of leading term, which will bring up the mixing with different spin J. But, this artificial mixing will decrease with the decreasing of lattice spacing a so that the mixing should vanish when $a \to 0$ in despite that it will affect our error estimate. On the other hand, we can utilize the non-leading terms to explore high-spin states.

These two examples tell us that to calculate the mass of the definite J^{PC} state, we should require the continuum limits of our operators belong to and only belong to J^{PC} representation of $SO(3)^{PC}$ group. One can achieve this aim by using the combination of the different operators which belong to the same R^{PC}.

3 The Construction of the Operator

We examplify here how to construct operator 0^{++} and 2^{++} up to a^4. Gauge-invariant current o with 0^{++} corresponding to the scalar glueball can be written as

$$o = \sum_n \Big\{ a^4 \sum_{i=1}^{5} \mathrm{Tr}(B_iB_i)(n) + a^6 \times (\text{current with mass dimension 6}) + \cdots \Big\}, \qquad (8)$$

where the current with mass dimension 6 is the combination of $\sum_{i,j}^{3}\mathrm{Tr}(D_iF_{ij}D_iF_{ij})$, $\sum_{i,j,k}^{3}\mathrm{Tr}(D_iF_{jk}D_iF_{jk})$ *and* $\sum_{i,j,k}^{3}\mathrm{Tr}(D_iF_{ik}D_jF_{jk})$. For simplicity, we only consider the current o up to mass dimension 4 in this paper. Then, let us observe the sum of the planar special 2×1 rectangular over all lattice sites:

$$\begin{aligned} O'_{ij} = \sum_n O'_{ij}(n) = \frac{1}{2}\sum_n \mathrm{Tr}\{&[1 - U(n,i)U(n+\hat{i},i)U(n+2\hat{i},j)U^{-1}(n+\hat{i}+\hat{j},i)\cdot \\ &U^{-1}(n+\hat{j},i)U^{-1}(n,j)] + [1 - U(n,i)U(n+\hat{i},j)U(n+\hat{i}+\hat{j},j)\cdot \\ &U^{-1}(n+2\hat{j},i)U^{-1}(n+\hat{j},j)U^{-1}(n,j)]\}. \end{aligned} \qquad (9)$$

The real part of the expansion for the operator O'_{ij} is

$$\mathrm{Re}\,O'_{ij} = \sum_n \Big\{ \frac{a^4}{2} 4\mathrm{Tr}(F_{ij}F_{ij})(n) - \frac{a^6}{24} 10\mathrm{Tr}(F_{ij}(D_i^2 + D_j^2)F_{ij})(n) \Big\} + O(a^8).$$

Then, we define

$$\Theta_{ij}(n) \equiv \mathrm{Re}\Big(O_{ij}(n) - \frac{1}{10}O'_{ij}(n) \Big).$$

Continuum limit of operator Θ_{ij} is

$$\Theta_{ij} = \sum_n \frac{3a^4}{10}\mathrm{Tr}(F_{ij}F_{ij}) + O(a^8). \tag{12}$$

Decomposing Θ_{ij} into A_1^{++} according to traditional method, we get the basis of representation A_1^{++}:

$$F = \Theta_{12} + \Theta_{13} + \Theta_{23} = \sum_n \frac{3a^4}{10}\mathrm{Tr}(B_1B_1 + B_2B_2 + B_3B_3) + O(a^8). \tag{13}$$

Apparently, the quantum number in the continuum limit of F is 0^{++}. In other words, F transforms as 0^{++} under $SO(3)^{PC}$ group up to a^4. We expect that the symmetry of $SO(3)$ has been restored when $a \to 0$ and the extracted state should be mainly given by the leading term of F. So that, the extracted state is 0^{++} in the continuum limit. Operator F is our aimed operator for 0^{++} state.

We may also choose the basis G_1 and G_2 of the representation E^{++} to measure the tensor glueball mass as follows. The operators and their expansions are

$$G_1 = \mathrm{Re}(\Theta_{23} - \Theta_{13}) = \sum_n \frac{3a^4}{10}\mathrm{Tr}(B_1B_1 \quad B_2B_2) + O(a^8), \tag{14}$$

and

$$G_2 = \mathrm{Re}(\Theta_{23} + \Theta_{13} - 2\Theta_{12}) = \sum_n \frac{3a^4}{10}\mathrm{Tr}(B_1B_1 + B_2B_2 - 2B_3B_3) + O(a^8). \tag{15}$$

According to (3b), they belong to basis of the representation 2^{++} up to $O(a^4)$.

4 Simulation Results

In the $SU(3)$ pure gauge theory and under the quenched approximation, we perform our calculation on an anisotropic lattice with improved gluonic action as chosen in Ref.[13]

$$S_{\mathrm{II}} = \beta\left\{\frac{5\Omega_{\mathrm{sp}}}{3\xi u_s^4} + \frac{4\xi\Omega_{\mathrm{tp}}}{3u_s^2u_t^2} - \frac{\Omega_{\mathrm{sr}}}{12\xi u_s^6} - \frac{\xi\Omega_{\mathrm{str}}}{12u_s^4u_t^2}\right\}, \tag{16}$$

where $\beta = 6/g^2$, g is the QCD coupling constant; u_s and u_t are mean link renormalization parameters, u_s is given by the fourth root of the average plaquette, and we set $u_t = 1$; $\xi = a_s/a_t$ is the aspect ratio; Ω_{sp} includes the sum over all spatial plaquettes on the lattice; Ω_{tp} indicates the temporal plaquettes; Ω_{sr} denotes the planar 2×1 spatial rectangular loops and Ω_{str} refers to the short temporal rectangles (one temporal and two spatial links). More detail is given in Ref.[13]. We adopt parameters in Ref.[3] as our simulation ones which are given in Table 1.

Table 1. The glueball simulation parameters[3]. Here we assume $r_0 = 410(20)$ MeV.

β	ξ	u_s^4	Lattice	r_s/r_0	a_s/fm
1.7	5	0.295	$6^3\times30$	0.8169	0.39
1.9	5	0.328	$8^3\times40$	0.727	0.35
2.2	5	0.378	$8^3\times40$	0.5680	0.27
2.4	5	0.409	$8^3\times40$	0.459	0.22
2.5	5	0.424	$10^3\times50$	0.407	0.20

As argued above, we choose operator F to calculate the scalar glueball mass and operator G_1 and G_2 to calculate the tensor glueball mass. As usual, we calculate the vacuum expectation of the correlation function $C(t) = \langle 0|\, o^R(t)o^R(0)\,|0\rangle$ to determine masses of the corresponding glueball states by fitting its decaying exponential, where $o^R(t) = o(t) - \langle 0|\, o(t)\,|0\rangle$ is the vacuum-sub-

tracted form of the chosen operator. Configuration ensembles were generated using Cabibbo-Marinari pseudo-heatbath and $SU(2)$ subgroup over-relaxation methods. We set 2800 heatbath sweeps to make configurations reach to equilibrium. At the same time, following the mean field theory[14], we also replace link variant U by U/u_s in the chosen operators to suppress the tadpole contribution. Improvements such as fuzzying and smearing are also used to suppress the fluctuations of the fields. Four heatbath updating sweeps and one canonical sweep were performed between two measurements. We group 5600 measurements into 80 bins for making a correct statistic error estimate. The results for $a_t m_G$ are shown in Table 2.

Table 2. Glueball energy $m_G a_t$ for each β.

β	1.7	1.9	2.2	2.4	2.5
scalar	0.609(4)	0.515(8)	0.412(7)	0.315(6)	0.322(2)
tensor	1.019(3)	0.95(1)	0.71(2)	0.548(6)	0.519(4)

Now we comment on the error estimate. First, our action breaks the rotation symmetry to $O(a_s^4, a_t^2)$, i.e., the upper limit of the precision in the calculation is $O(a_s^4, a_t^2)$. Since as argued by many authors, the contribution of $O(a_t^2)$ can be ignored, the upper limit of the precision here is $O(a_s^4)$. Second, we ignore terms (currents) with mass dimension 6 in Eq.(8). Due to the dimensional analysis, the contribution of the currents to error should have a square mass suppression[9], which will make two effects on our mass measurement. One is that we should include it in systematic error in the continuum limit, which needs further calculation to get its accurate value. Here we simply expect that it is about $(\Lambda_{QCD}/m)^2$, where we set $\Lambda_{QCD} \approx 250$MeV and m is measured mass. The second one is that it will take $O(a_s^2)$ error when $a_s \neq 0$. Since it is not statistical error, its contribution to error can be fitted by $c_2 a_s^2 + c_4 a_s^4 + \cdots$.

As expected, we get the same simulation results from operators G_1 and G_2, since they correspond to the same state 2^{++}.

From the argument and calculated data, we use the formula $m(0^{++}, a_s) = 1.754 - 1.514(a_s/r_0)^2 + 1.773(a_s/r_0)^4$ and $m(2^{++}, a_s) = 2.417 + 0.783(a_s/r_0)^2 - 0.787(a_s/r_0)^4$ (unit: GeV) to fit our data. We present our data and fitting curves in Fig.1.

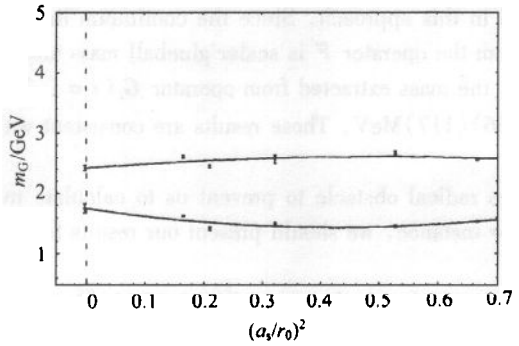

Fig.1. Masses of scalar and tensor glueball against the lattice spacing $(a_s/r_0)^2$
The fitting curves are $m(0^{++}, a_s) = 1.754 - 1.514(a_s/r_0)^2 + 1.773(a_s/r_0)^4$ for scalar glueball mass and $m(2^{++}, a_s) = 2.417 + 0.783(a_s/r_0)^2 - 0.787(a_s/r_0)^4$ (unit: GeV) for tensor glueball mass, respectively. The masses in the continuum limit are 1.754(76) GeV and 2.417(44) GeV if we only consider the statistical error.

The statistical error is 0.076 GeV for scalar glueball and 0.044 GeV for tensor glueball. According to Ref.[3], systematic error is 1 percent (from aspect ratio). But since our method also

gives about 2 and 1 percent systematic error for 0^{++} and 2^{++} states respectively, the total systematic error is about 2.2 percent (39MeV) and 1.4 percent (34MeV) respectively. Therefore, the mass of scalar glueball is 1.754(85)GeV and the mass of tensor glueball is 2.417(56) GeV. Including the uncertainty in $r_0^{-1} = 410(20)$MeV, our final results are $m_G(0^{++}) = 1754(85)(86)$ MeV and $m_G(2^{++}) = 2417(56)(117)$ MeV.

For comparsion, we list our results and that from UKQCD[2], IMB[5] and Morningstar and Peardon[3] in Table 3.

Table 3. 0^{++} and 2^{++} glueball masses (unit: MeV).

	UKQCD[2]	IBM[5]	Ref.[3]	our results
scalar	1550(50)	1740(71)	1730(50)(80)	1754(85)(86)
tensor	2270(100)	2359(128)	2400(25)(120)	2417(56)(11)

We set our simulation parameters as those in Ref.[3]. Therefore, one can find that our simulation results are more consistant with that from Ref.[3]. It conforms our approach. But, unfortunately, Morningstar and Peardon[3] apply the variational principle in their numerical simulation, we cannot compare their operators with ours by expanding it.

5 Conclusion

Basing on the connection between the asymptotic expansion of the operators and the quantum number J^{PC} of the extracted state, we have presented a new approach to constructing operator on lattice for the calculation of the glueball mass, which may solve the ambiguity in the simulation. In general, to calculate the mass of definite J^{PC} glueball states, first one should write out these currents which transform as the representation J^{PC} under the $SO(3)^{PC}$ group in continuum and decompose them into irreducible representations R^{PC} of the group O^{PC} in the subduced representation, then one should construct corresponding operators which belong to the representation R^{PC} of the O^{PC} group on the lattice just as the paper mentioned above.

To verify our approach, we have calculated the $SU(3)$ scalar and tensor glueball mass under the quenched approximation in this approach. Since the continuum limit of operator F is 0^{++}, we affirm the mass extracted from the operator F is scalar glueball mass and its value is 1754(85)(86) MeV. For the same reason, the mass extracted from operator $G_i\ (i = 1,2)$ is that of the tensor glueball and its value is 2417(56)(117)MeV. These results are consistent with those obtained in Refs. [3,5,13,16].

Apparently, there is no radical obstacle to prevent us to calculate the mass of states with any spin J in this approach. For instance, we should present our results for the mass of the ground 4^{++} glueball elsewhere.

Of course, the operator, which transforms as J^{PC} of $SO(3)^{PC}$ group in the continuum, is not unique. For example, one can also construct the operator including color electric field. With these operators, one can determine their relative weights of contribution by variational principle. But, we did not make such treatment here.

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在格点上构造计算胶球质量的新途径

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摘要　发展了一种为了计算胶球质量而构造格点算符的新途径. 基于所选用算符的连续极限与状态量子数 J^{PC} 两者之间的联系,状态的自旋就可以在数值模拟中唯一和直接地被确定下来. 进而,这一途径可以被应用于计算任意自旋 J 的胶球质量. 在淬火近似下,给出在 $SU(3)$ 纯规范场中 0^{++} 态和 2^{++} 态胶球质量的初步结果,它们分别是1754(85)(86) MeV 和 2417(56)(117)MeV.

关键词　格点量子色动力学　胶球质量　Wilson 算符

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