Production of $0^{++}$ glueball from double diffractive process in high energy $p + p(\bar{p})$ collision

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Abstract

Motivated by the recent experimental data about candidates for glueball from different processes, we discuss in this paper the production of $0^{++}$ glueball from double diffractive scattering at momentum transfer $|t| \lesssim 1 GeV^2$ in high energy $p + p(\bar{p})$ collision. We apply the phenomenology of Pomeron ($IP$) of Donnachie-Landshoff, the field theory model of $IP$ of Landshoff-Nachtmann and the relevant calculating approaches. We assume while $IP$ coupling with glueball, the $0^{++}$ glueball can be considered as a bound state of two non-perturbative massive gluons. We evaluate the dependence of cross section for $0^{++}$ glueball production on system energy $\sqrt{s}$ and show that it could be tested experimentally.

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1 Introduction

Since the color $SU(3)$ gauge theory possesses self-coupling characteristic between gluon fields, it had predicted long before the quantum chromodynamics (QCD) became the basic dynamics of strong interaction, that there would exist states in hadronic spectroscopy which are formed from purely gluon field—glueball. Though now there are still lack of strict theoretical proving and conclusive experimental evidence to confirm its existence, in recent couple of years there are a lot of reports from relevant data analysis that several candidates of glueball and/or states mixing with quark pair have been observed experimentally. It is well believed that glueballs should easily be produced in processes with rich source of gluon constituent, thus experimentalists concentrate their attention to seek glueballs on $p\bar{p}$ annihilation, $J/\psi$ radiative decay, and high energy central $h^{-}h^{+}$ collisions. Glueball production from $J/\psi$ radiative decay has been analyzed in a model independent manner by Close et al., but yet there has not been any analysis in central $h^{-}h^{+}$ collision and in $p\bar{p}$ annihilation.

At high energy central $h^{-}h^{+}$ collisions the increase in cross section with increased center of mass energy is consistent with the double pomeron exchange (DPE) mechanism, which was predicted to be a source of gluonic states. A large glueball production cross section in the central region is predicted by Gershtein and Logunov, they related the rise of total cross section with increasing energy to the exchange of glueballs in the $t$-channel or to the colliding of the soft gluon seas of the interacting particles. These observations tell us that glueball production in high energy $h^{-}h^{+}$ central collisions would be intimately connected with the non-perturbative mechanism.

In this paper, we discuss the glueball production from double diffractive scattering in high energy $p+p(\bar{p})$ collisions, as a first attempt, we consider $J^{PC} = 0^{++}$ glueball only.

Our approach is on the basis of the phenomenology of $IP$ from Donnachie-Landshoff, the field theory model of $IP$ from Landshoff-Nachtmann. Based on the above approach, a more refined and sophisticated model of diffractive scattering had proposed by Cudell and Hernandez, concentrating on Higgs diffractive production in pp collision $p + p \rightarrow p + p + H$, they obtained a result close to that of Bialas and Landshoff using the L-N model of DPE.

Since the glueballs are produced through DPE, it could believed that they are laid predominately in the central region of final state rapidity distribution and with large symmetry rapidity gaps from each of the final diffractive protons (or antiprotons) which are the characteristic properties of DPE processed as Eq.(1). In principle, it should be easily detected and according to it we could distinguish it from all other strong backgrounds. We will come back to this point in last section.

Of how to link L-N field theory of $IP$ with the phenomenology of Regge behaviour, we follow the approach of Bialas and Landshoff. Since the information of coupling of $0^{++}$ glueball with non-perturbative gluon coupling is very scarely, we adopt the color singlet approximation, which always used in heavy quark physics. Under this approximation the coupling structure is $Bg^{\alpha\beta}\delta_{\delta\gamma}$, where $\alpha e etc.$ are the Lorentz and color indexes of the two non-perturbative gluons respectively, the parameter $B$ can be estimated from the branching ratio of $J/\psi$ radiative decay to $0^{++}$ glueball.

In the following we describe in detail our calculating scheme in Sec.II, we give the formalism and derive the coupling constant $B$ of the two nonperturbative gluons with $0^{++}$ glueball in Sec.III, our results and some related discussions are given in last Sec.IV.

2 Calculating Scheme

In phenomenology of $IP$ given by Donnachie-Landshoff, the nucleon is to be treated as a bound state composed of three clusters, they are formed from non-perturbative QCD effects and their size are much smaller than nucleon radius. The core of each cluster is a valence quark, with sea quarks and gluons surrounded. Furthermore D-L have demonstrated that in coupling with nucleon, to a good approximation, the $IP$ behaves like an isoscalar ($C = +1$) photon. So it means that the coupling of $IP$ with a nucleon is actually
that of IP with the three clusters (constitute quark) in an incoherent manner. In this picture IP model one does not need to consider the interaction of IP with any other parton (sea quark or gluons) in the nucleon. With this model Donnachie and Landshoff have successfully explained a lot of experimental effects in high energy soft strong interaction processes.

In discussing glueball production from DPE, we adopt the same point of view: we should not start such a problem from parton model and perturbative QCD, but instead from non-perturbative QCD at all and assume the glueball to be a bound state of two small flavorless clusters (or “constituent” gluons) which are formed from gluons by non-perturbative QCD effects, and are massive. Like the nucleon case, we assume the coupling of IP with glueball is actually that of IP with these “constituent” gluons. Also, we do not need to consider the interaction of IP with any other massless gluons in the glueball.

From the status both of the available data for elastic diffractive scattering and the effectuality of D-L IP model in high energy h-h collisions, we restricted in this paper the momentum transfer of scattered proton (antiproton) to be small, e.g. $|t_i| \lesssim 1.0\text{GeV}^2$. The whole picture is sketched in Fig.1.

From Regge pole theory\(^1\) and IP model of D-L, when $s \ll s_1$, $s_2 \ll m_N^2$, and $|t_i| \lesssim 1\text{GeV}^2$, the asymptotic form of the amplitude for Fig.1 is

$$
\left(\frac{s}{s_2}\right)^{\alpha_{IP}(t_1)-1}\left(\frac{s}{s_1}\right)^{\alpha_{IP}(t_2)-1} F^{(G)}(\alpha_{IP}(t_1), \alpha_{IP}(t_2), z) F_1(t_1) F_1(t_2) \gamma_\lambda \otimes \gamma^\lambda
$$

where $t_i = (p_i - p_i')^2$, $z = \frac{s_i}{m^2}$, and $m$ is the 0++ glueball mass, $\alpha_{IP}(t)$ is Regge pole trajectory of IP. From data analysis\(^2\) when $|t|$ is small, $\alpha_{IP}(t) \approx \alpha_{IP}(0) + \alpha'_{IP} t = 1.086 + 0.25\text{GeV}^2 t$. $F_1(t)$ is the Dirac form factor of proton. The direct product of gamma matrix shows that the external proton lines will be traced when calculate the cross section.

Since the diffractive scattering condition requires

$$
\frac{s_1}{s} \frac{s_2}{s} < \delta
$$

and usually assume $\delta = 0.1$, hence in the asymptotic form Eq.(2) one could neglect lower powers of $s/s_1$ and $s/s_2$. $F^{(G)}(\alpha_{IP}(t_1), \alpha_{IP}(t_2), z)$ is the IP $- IP - G$ vertex function, its general structure is known\(^3\), but its concrete form must be further considered.

Since IP-nucleon coupling should be considered as pomeron incoherently coupled with three constituent quarks, and according to field theory model, the IP exchange is just the exchange of two non-perturbative gluons\(^4\)\(^5\), thus Fig.1 should be concreted as Fig.2.

In the field theory model of pomeron, the IP $- IP - G$ vertex function $F^{(G)}(\alpha_{IP}(t_1), \alpha_{IP}(t_2), z)$ for $|t_i| \lesssim 1.0\text{GeV}^2$ can be calculated at by the sum of Feynman amplitudes of Fig.2 which as shown in \(^4\) is equal to its $s$-channel discontinuity of the first diagram, thus as Fig.3. This simplify our calculation greatly. The black blob in Fig.3 is the vertex function of 0++ glueball with gluons of IP which will be discussed in next section.

In the following we first get the forward diffractive scattering amplitude for Eq.(1) in L-N field theory model of IP, then identify it with the corresponding Regge behaviour formula Eq.(2) at $t_i = 0$ to obtain the normalized value of IP $- IP -$ glueball vertex function $F^{(G)}(\alpha_{IP}(t_1), \alpha_{IP}(t_2), z)|_{t_i=t_2=0}$. After this has done we continue Eq.(2) to $|t_i| \neq 0$ and discuss DPE production of glueball.

### 3 Formalism and Coupling Constant $B$

Using the Sudakoff variables to calculate the loop integral of Fig.3 is very convenient, here

$$
k = \frac{\vec{p}_1}{s} + \frac{\vec{p}_2}{s} + v
$$

\(^1\)Since in D-L model, $\alpha_{IP}(t)|_{t=0} \equiv 1 + \varepsilon \approx 1.086$, so it seems having some inconsistent with this identification, which has explained in \(^5\). Furthermore, this problem has studied by Ross\(^6\) who using a hybrid model in which one uses PQCD to treat interactions between gluons, but takes a D-L type non-perturbative propagator for the gluon. In the leading logarithm approximation, by summing a dominant subset of diagrams, the phenomenonolog required Regge behaviour is obtained.
\[ p'_1 = x_1 p_1 + \frac{y_1 p_2}{s} + v_1 \]
\[ p'_2 = \frac{x_2 p_1}{s} + y_2 p_2 + v_2 \]  
(4)

We have put the light quark mass equal zero, the \( v, v_1, \) and \( v_2 \) are transverse to \( p_1 \) and \( p_2 \) and so effectively are two dimensional, \( v^2 = -v^2, v'_2 = -v'_2 \). Then

\[ t_1 = (p_1 - p'_1)^2 = -v^2_1/x_1 \]
\[ t_2 = (p_2 - p'_2)^2 = -v^2_2/y_2 \]
\[ s_1 \sim s(1 - y_2) \]
\[ s_2 \sim s(1 - x_1) \]  
(5)

As mentioned above, we first need the \( s- \) channel discontinuity of Fig.3 at \( t_1 \approx t_2 \approx 0 \), so we set \( v_1 \approx v_2 \approx 0 \), then

\[ \int d^4k \delta(q^2_1) \delta(q^2_2) \sim \frac{1}{2s} \int d\bar{x}dyd^2v \delta(\bar{y} - v^2) \delta(-\bar{x} - v^2), \]
\[ \int d^4p'_1 \delta(p'^2_1) \sim \frac{1}{2} \int d\bar{x}_1d\bar{y}_1 \delta(\bar{x}_1\bar{y}_1)d^2v_1, \]
\[ \int d^4p'_2 \delta(p'^2_2) \sim \frac{1}{2} \int d\bar{x}_2dy_2 \delta(\bar{x}_2y_2)d^2v_2. \]  
(6)

Here we explain how we can get from Fig.3 the direct product from \( \gamma_\lambda \otimes \gamma^\lambda \) as in Eq.(2). The upper quark line in Fig.3 has gamma matrices \( \gamma^\mu_1 \gamma \cdot q_1 \gamma_\lambda \), for large \( s \) its asymptotic form equivalent to

\[ 2q^\mu_1 \gamma^\lambda, \]  
(7)

because in calculating the differential cross section, we also need to multiply \( \gamma \cdot p_1 \) on the left and \( \gamma \cdot p'_1 \) on the right of this expression. When these are included, the difference between (7) and the original expressions gives contribution to cross section are of order \( \delta \). Similarly, from the lower quark line we obtain \( 2q^\mu_2 \gamma_\lambda \).

Now from Fig.3, the amplitude of \( q - q(\bar{q}) \) diffractive scattering through DPE at \( t_1 = t_2 = 0 \) is

\[ M^{(G)} = \gamma^\lambda \otimes \gamma^\lambda, \]
\[ M^{(G)} = \frac{ig^6}{2\pi^2s} \int d^2v W^{(G)} D(-v^2)D(-x_1v^2)D(-y_2v^2), \]  
(8)

where \( D(q^2) \) is non-perturbative gluon propagator of \( IP, \) \( g \) is the coupling constant between non-perturbative gluon with quark,

\[ g^2 D(q^2) = -A \exp(\frac{q^2}{\mu^2}), \]  
(9)

where \( \mu = 1.1 GeV, A^2\mu^2 = 72\pi\beta^2, \beta^2 = 3.93 GeV^{-2} \) is \( IP-nucleon \) coupling constant in the pomeron D-L model[20].

Using the approximation matrix form Eq.(7), we found from Fig.3, the \( W^{(G)} \) in Eq.(8),

\[ W^{(G)} = q^\mu_1 A^{(G)}_{\mu \nu} q^\nu_2, \]  
(10)

where vertex function \( A^{(G)}_{\mu \nu} \) connecting the gluons of \( IP \) with \( 0^+ \) glueball. The simplest diagrams for this vertex function shown in Fig.4 and Fig.5.
In deriving these vertex functions, for calculating simplicity but not necessary we have used color singlet model approximation for the vertex, which require gluon lines cutting by vertical dashed line in these figures are limited on mass shell. Thus from Fig.4 we get

\[ A_{\mu
u,bc}^{G_3} = 2BF_{\mu\nu,\lambda}(k_1, q, \frac{1}{2} P)D(q^2)g^{\alpha\beta}F_{\beta\nu\rho}(-\frac{1}{2} P, k_2, -q)g^{\rho\lambda} :f_{beg} f_{hec} : \delta_{eh}, \]

where

\[ F_{\mu_1\mu_2\mu_3}(r_1, r_2, r_3) = (r_1 - r_2)_{\mu_1} g_{\mu_1\mu_2} + (r_2 - r_3)_{\mu_2} g_{\mu_2\mu_3} + (r_3 - r_1)_{\mu_3} g_{\mu_3\mu_1} \]

\[ r_1 + r_2 + r_3 = 0. \]

From Fig.5 we get

\[ A_{\mu
u,bc}^{G_4} = -12B g_{\mu\nu} f_{beg} f_{hec} \delta_{eh}. \]

From Eqs.(10) to (13) we get

\[ W^{G_3} = \frac{-B D(q^2)}{4s} \cdot (5m^2s^3 + 10m^2s^2v^2 + 16m^2sv^4 + 15s^3v^2x_1 + 15s^3v^2y_2 - 10s^3v^2 + 6s^2v^4x_1 + 6s^2v^4y_2 + 4s^2v^4) \]

\[ q^2 = \frac{1 - 1}{4s}(m^2s + 2sv^2x_1 + 2sv^2y_2 + 4v^4) \]

\[ W^{G_4} = -12B \left( \frac{1}{2}s + \frac{v^4}{s} \right) \]

\[ W^{(G)} = W^{G_3} + W^{G_4} \]

Form Eqs.(8)(9)(14), we get \( M^{(G)} \), for Fig.3, which is \( F^{(G)}(\alpha_{IP}(t_1), \alpha_{IP}(t_2), z)|_{t_1=t_2=0} \) as we explained in section II.

For \( t_1, t_2 \neq 0 \), the general structure of \( F(\alpha_{IP}(t_1), \alpha_{IP}(t_2), z) \) can be analyzed from Regge theory and this has been done by Drummon et. al.[17]. They have shown that when \( |t_1|, |t_2| \) are small, \( F(\alpha_{IP}(t_1), \alpha_{IP}(t_2), z) \) are finite for any \( z \). This point is very important for us since the D-L L\( IP \) model, in strictly speaking, is valid only for \( |t_1|, |t_2| \leq 1 \text{GeV}^2 \). But we can also see from Eqs.(2) and (3), all contributions from large \( |t_1|, |t_2| \) are higher powers of \( \delta \), and they can be neglected safely. So, for \( t_1, t_2 \neq 0 \), we have

\[ F(\alpha_{IP}(t_1), \alpha_{IP}(t_2), z)|_{t_1,t_2 \leq 1 \text{GeV}^2} \simeq F(\alpha_{IP}(t_1), \alpha_{IP}(t_2), z)|_{t_1=t_2=0} = M^{(G)} \]

Thus the amplitude \( T^{(G)} \) for glueball production in \( p - p(\bar{p}) \) high energy double diffractive scattering from D-L and L-N L\( IP \) model and formalism, in a good approximation, can be got from Eqs.(2) and (15) and is

\[ T^{(G)} = 9M^{(G)} \left( \frac{s}{s_2} \right)^{\alpha_{IP}(t_1)-1} \left( \frac{s}{s_1} \right)^{\alpha_{IP}(t_2)-1} F_1(t_1) F_1(t_2) \gamma_{\lambda \otimes \gamma^{\lambda}}. \]

where factor 9 comes from 3 quarks in nucleon, \( F_1(t) \) is the Dirac form factor of proton, using the exponential approximation for \( F_1(t) \) at small \( |t| : F_1(t) \simeq e^{\lambda t}, \lambda \simeq 2 \text{GeV}^2 \), the cross section formula is

\[ \sigma^{(G)} = \frac{F^G_c}{2(2\pi)^3} \left( \frac{s}{m^2} \right)^{2\epsilon} \int \frac{dx_1 dy_1}{x_1 y_1} \frac{M^{(G)}}{2} \delta( (1 - x_1)(1 - y_2) - \frac{m^2}{s} ) \]

\[ \cdot \int d^2v_1 d^2v_2 (1 - x_1)^{2\alpha' v_1^2/x_1} (1 - y_2)^{2\alpha' v_2^2/y_2} e^{-2\lambda(v_1^2/x_1 + v_2^2/y_2)} \]

where \( F^G_c = \frac{s}{s} \) is the color factor.
Let us now consider the coupling constant $B$ of constituent gluons with $0^{++}$ glueball, which could be fixed by radiative decay mode $J/\psi \rightarrow \gamma + f_0(1500)$ as shown in Fig.6, where

$$\begin{align*}
p_i &= \frac{1}{2}(p_j - p_G) \\
p_j &= k - \frac{1}{2}p_j \\
p_k &= \frac{1}{2}(p_G - p_f) \\
g_i &= g_2 = \frac{1}{2}p_G
\end{align*}$$

(18)

It is well known that the $J/\psi$ is dominantly an $S$-wave state and the non-relativistic quark potential mode is very successful in describing the static properties of the $J/\psi$. Inclusive quarkonium decays into light hadrons are accessible to PQCD. For heavy-quarkonium, the annihilation time $B$ gets coupling constant $g_2$. Then decay width of $J/\psi$ gets

$$\Gamma_{J/\psi} \sim \frac{1}{m_Q}$$

from $\Gamma_{J/\psi}$, we get

$$\Gamma_{J/\psi} = \frac{g_2^2}{16\pi m_j^2} \sum_i |M^j|^2$$

(20)

and

$$\begin{align*}
M_{\mu\nu}^{(a)} &= \frac{1}{\sqrt{2}} \epsilon_\mu \cdot \epsilon_\nu \gamma_\mu \gamma_5 \gamma_\nu \\
M_{\mu\nu}^{(b)} &= \frac{1}{\sqrt{2}} \epsilon_\mu \cdot \epsilon_\nu \gamma_\mu \gamma_5 \\
M_{\mu\nu}^{(c)} &= \frac{1}{\sqrt{2}} \epsilon_\mu \cdot \epsilon_\nu \gamma_\mu \gamma_5
\end{align*}$$

(21)

Put Eqs.(18) and (19) into (17), after sum and average with initial states and final states respectively, we get

$$\sum_i |M^j|^2 = \frac{2}{3}A^2G^2c_\psi^4 \frac{1024(m_6^2 - m_4^2 m_4 + 11m_2^2 m_4^2 + m_6^2)}{m_j^4(m_j^2 - m_2^2)^2},$$

(22)

then decay width of $J/\psi \rightarrow \gamma + f_0$ is

$$\Gamma_{J/\psi} = m_j^2 - m_2^2 \frac{\sum_i |M^j|^2}{16\pi m_j^2},$$

(23)

where color factor $F_2^J = \frac{2}{3}$.

Put Eqs.(18), (20), and (21) together, we get

$$B^2 = \frac{\sum_i |M^j|^2}{16\pi m_j^2} \frac{9\pi a m_j^4 (m_j^2 - m_2^2)}{128G^4 (m_6^2 + 11m_2^2 m_4^2 - m_2^2 m_4^2 + m_6^2)}$$

(24)

Decay width $\Gamma_{J/\psi} \rightarrow \gamma + f_0$ can be found in \[3\] and in \[21\] that the scalar $f_0(1500)$ may be a glueball-$q\bar{q}$ mixture, here for simplicity we assume it is a pure glueball.
4 Results and Discussions

We first give a brief comments on the parameters used in our paper.

1. 0++ glueball mass $m$.

From several models (e.g. lattice QCD, bag model, potential model, and QCD sum rule) analysis, the lowest state mass of $0^{++}$ all are fixed at $1.5 \sim 1.7 GeV$\[2\]. Experimentally, relevant data also manifest a clear signal of $0^{++}$ resonance state at about $1.5 GeV$ in the central region of high energy $p-p$ collision, thus we let $m = 1.5 GeV$.

2. Nonperturbative coupling constant $g$.

In the Abelian gluon theory of Landshoff and Nachtmann, the nonperturbative gluon only couples to quark, thus $g^2 D(q^2)$ always appear together and can be normalised by the constant $A$ in Eq.(9), the nonperturbative constant does not enter the calculation of LN model. But in non-Abelian case, especially including gluon self-interactions, as showed in Fig.4 and Fig.5, $g$ cannot be all absorbed by normalized condition Eq.(9), so $g$ enters indeed into our the calculation. Unfortunately our knowledge for the non-perturbative coupling constant $g$ is very poor now, in order to get a sensible answer for the gluon structure function at small $x$ \[23\] $\alpha_n = \frac{g^2}{4\pi} \sim 0.7$.

Putting these parameters into Eq.\[14\] the double diffractive production cross section of Eq.(1) for $\sqrt{s}$ from $20 GeV$ to $2 \times 10^4$ GeV, where $|t_1|, |t_2| < 1 GeV^2, \delta = 0.1$ are evaluated, as shown in Fig.7. We see for $\sqrt{s} = 23.7 GeV, 29 Gev, 630 GeV(SppS)$ and $1.8$ TeV (Tevatron) energies, the double diffractive production cross sections are $1.6 \times 10^2 nb, 2.5 \times 10^2 nb, 2.8 \mu b$ and $4.6 \mu b$ respectively. In the Joint CERN-IHEP experiment in $300 GeV$ central $\pi^- N$ collisions ($\sqrt{s} = 23.7 GeV$), the production cross section in the range $0 < x_F < 0.3$ (the experiment setup lead to the acceptance is zero for $x_F < 0$) is $0.2 \pm 0.1 \mu b$. Taking account of the additive quark rule, this corresponds to $0.4 - 1.4 \mu b$ in $P - P(\bar{P})$ central collisions at the same center-of-mass energy in the range $-0.3 < x_F < 0.3$. Since $f(1500)$ is produced dominantly at small $|t|$, our results are reasonable.

We have restricted ourselves to calculate elastic diffractive production process. It is easily extend to diffractive dissociation processes by removing the form factor $F_1(t_i)/(i = 1, 2)$ in Eq.(2) and thereafter, the resulting cross sections are also show in Fig.7. We can see that as center of mass energies increased the ratio of the cross section of diffractive dissociation to that of elastic one is reduced from 6 at $\sqrt{s} = 23 GeV$ to 3 at $\sqrt{s} = 20 TeV$.

Throughout the calculation we work in Feynman gauge. If one were to work in another gauge the function $D(q^2)$ must be substantially modified to ensure that physical observables are gauge independent.

In connection with the problem of experimentally detecting the process Eq.(1), we make following discussions:

1. Since glueballs in process Eq.(1) are produced through DPE in high energy $p - p(\bar{p})$ collision, they should be laid predominately in the central region of rapidity distribution of final particles and with large symmetry rapidity gaps from each of the final proton (antiproton) direction. However several effects would be weaken the gap rate. First, althropg the glueball and its decay product are colorless, as final state interacting there are get having survival soft color interactions between pre-glueball with outgoing pro-nucleons. The estimated gap survival rate to be $10^{-2}$ at Fermilab Tevatron \[23\]. Secondly, the produced glueball is a light object ($m_G \sim 1.5 GeV$), they will be produced in a broad rapidity range, for example, at Tevatron those glueballs are in the rapidity region $-5.9 < \eta < 5.9$.

2. We take the D-L IP model and DPE process to discuss glueball production process at small $|t|$ , like all other high energy diffractive processed discussed by using the same model, the common remarkable

\[3\] Even in process which are absent for glueon selfinteractions, may meet the same troubles. For example, in discussing the diffractive Higgs production process $p + p \rightarrow p + p + H_1$, only after supposing implicitly that coupling constant between top quark with IP equals to that of u, d quark with IP, then constant g can be all absorbed in Eq.(9).
character and parameter independent property are the energy dependence of total cross sections which are proportional to \( s^{2(\alpha_{I P}(0)^{-1})} \). So the cross section in this model always increases slowly with the increase of the center of mass energy, we could use this point to distinguish this production mechanism of glueball from others, especially from gluon-gluon fusion mechanism in parton model, since for later the production cross section would be decreased rapidly as energy increases, perhaps this is the most effective way to test our results if one could measure the total cross section of process Eq.(1) over a large energy range.

In conclusion, using the field theory model of pomeron exchange and the color singlet approximation of glueball, we obtain a parameter-free prediction of the cross section of \( 0^{++} \) glueball in diffractive production, which combineing with the experiment measuring of \( \sigma (hh \rightarrow hhG) \cdot BR (G \rightarrow h1 + h2 + \cdots) \) can output the important quantities \( BR (G \rightarrow h1 + h2 + \cdots) \), saying \( BR (G \rightarrow K\bar{K}) \), or if we knew \( BR (G \rightarrow h1 + h2 + \cdots) \) then glueball production would be a test of the approach we used.

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Figure Captions

Fig.1. Sketch diagram for Eq.(1) via DPE.
Fig.2. $G$ production in $q - q(\bar{q})$ scattering via DPE, dotted lines are non-perturbative gluon of $IP$.
Fig.3. $s-$channel discontinuity of the first diagram in Fig.2. $q_1$, $q_2$ lines are limited on their mass shell.
Fig.4. Vertex of glueball and gluons of $IP$. Vertical dashed line limit cutting lines on mass shell.
Fig.5. The same as Fig.4, but for four gluons vertex.
Fig.6. Glueball production in radiative decay of $J/\psi$.
Fig.7. Production cross section of $0^{++}$ glueball in double diffractive $P - P(\bar{P})$ collision. a) Elastic case, b) Including the diffractive dissociation contribution.
Fig. 1 Sketch diagram for Eq.(1) via DPE.
Fig. 2 $G$ production in $q - q(\bar{q})$ scattering via DPE, dotted lines are non-perturbative gluon of $IP$. 
Fig. 3 $s$—channel discontinuity of the first diagram in Fig. 2. $q_1$, $q_2$ lines are limited on their mass shell.
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Fig. 5 The same as Fig. 4, but for four gluons vertex.
Fig. 6 Glueball production in radiative decay of $J/\psi$. 

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Fig. 7. Ecm (GeV) Including diffractive dissociation contribution

Elastic case

Including diffractive dissociation contribution