Implications for $B \rightarrow \eta K$ and $B \rightarrow \text{glueball} + K$ modes from observed large $B \rightarrow \eta' K + X$

Xiao-Gang He ^a, Wei-Shu Hou ^b, C.-S. Huang ^c

^a School of Physics, University of Melbourne, Melbourne, Australia

^b Department of Physics, National Taiwan University, Taipei, 10764 Taiwan, ROC

^c Institute of Theoretical Physics, Academia Sinica, Beijing, China

Received 11 January 1998

Editor: H. Georgi

Abstract

The unexpectedly large branching ratios for $B \rightarrow \eta' K (\eta' X_s)$ decays could be of gluonic origin. We study the implications for $B \rightarrow \eta K (\eta X_s)$ and $PK (PX_s)$, where P is the pseudoscalar glueball. In the mechanism proposed by Fritzsche, large branching ratios are predicted for these modes. The $B \rightarrow \eta K$ rate is barely within the experimental limit, and $B \rightarrow PK$, PX_s could be at the 0.1% and 1% level, respectively. Smaller but less definite results are found for the mechanism of $g^* \rightarrow \eta' g$ via the gluon anomaly. © 1998 Published by Elsevier Science B.V. All rights reserved.

Large branching ratios for exclusive $B \rightarrow \eta' K$ and semi-inclusive $B \rightarrow \eta' X_s$ decays have been reported recently, giving [1]

$$\begin{aligned} Br(B^\pm \rightarrow \eta' K^\pm) &= (7.1_{-2.1}^{+2.5} \pm 0.9) \times 10^{-5}, \\ Br(B^0 \rightarrow \eta' K^0) &= (5.3_{-2.2}^{+2.8} \pm 1.2) \times 10^{-5}, \\ Br(B \rightarrow \eta' X_s) &= (6.2 \pm 1.6 \pm 1.3) \times 10^{-4} (2.0 < p_{\eta'} < 2.7 \text{ GeV}). \end{aligned} \quad (1)$$

Factorization calculations of four quark operators in the Standard Model (SM) indicate that exclusive branching ratios could be accounted for by choice of form factors [2–4]. However, the four quark operators do not seem sufficient for the semi-inclusive decay [2]. Several mechanisms have been proposed [5–10] to explain the latter, and ways to distinguish some of the mechanisms have also been proposed [11]. We wish to study here the implications for gluonic origins [5,6,8] of such enhancements. In particular, we concentrate on the mechanism proposed in Ref. [8], which postulates an ad hoc effective Hamiltonian of the form

$$H_{\text{eff}} = a \alpha_s G_F \bar{s}_L b_R G_{\mu\nu} \tilde{G}^{\mu\nu}. \quad (2)$$

In the notation of Ref. [8], a contains a factor of V_{ts} . Similar effective Hamiltonians can be generated perturbatively at one loop level in the SM [12]. However, the resulting $Br(B \rightarrow \eta' X_s)$ would be smaller by more than two orders of magnitude. It was therefore argued in Ref. [8] that Eq. (2) may arise from non-perturbative effects. It may also arise from new physics. It turns out that a single parameter a could account

for both inclusive and exclusive η' rates of Eq. (1). Although the semi-inclusive m_{X_s} spectrum seems to favor 3-body decay over 2-body, which supports the anomaly induced $g^* \rightarrow \eta' g$ mechanism [5,6], some smearing effect cannot yet be ruled out [8]. Fitting Eq. (2) to Eq. (1), we find that η and pseudoscalar glueball P modes could then be quite large: $Br(B \rightarrow \eta K)$ is barely within the experimental limit, while $B \rightarrow PK$, PX_s could be at the 0.1% and 1% level, respectively.

To calculate $Br(B \rightarrow \eta' X_s)$ from Eq. (2), one has to project out the $G_{\mu\nu} \tilde{G}^{\mu\nu}$ content in η' , i.e. extract the matrix element $\langle 0 | G_{\mu\nu} \tilde{G}^{\mu\nu} | \eta' \rangle$. Using the definitions

$$\begin{aligned} \eta &= \eta_8 \cos \theta - \eta_0 \sin \theta, & \eta' &= \eta_8 \sin \theta + \eta_0 \cos \theta, \\ \langle 0 | j_\mu^8 | \eta_8 \rangle &= \langle 0 | \frac{\bar{u} \gamma_\mu \gamma_5 u + \bar{d} \gamma_\mu \gamma_5 d - 2 \bar{s} \gamma_\mu \gamma_5 s}{\sqrt{6}} | \eta_8 \rangle = i f_8 p_\mu, \\ \langle 0 | j_\mu^0 | \eta_0 \rangle &= \langle 0 | \frac{\bar{u} \gamma_\mu \gamma_5 u + \bar{d} \gamma_\mu \gamma_5 d + \bar{s} \gamma_\mu \gamma_5 s}{\sqrt{3}} | \eta_0 \rangle = i f_0 p_\mu, \end{aligned} \quad (3)$$

we obtain [13]

$$\begin{aligned} \langle 0 | \partial^\mu j_\mu^8 | \eta \rangle &= f_8 \cos \theta m_\eta^2 = \frac{1}{\sqrt{6}} \langle 0 | i 2 m_u \bar{u} \gamma_5 u + i 2 m_d \bar{d} \gamma_5 d - i 4 m_s \bar{s} \gamma_5 s | \eta \rangle, \\ \langle 0 | \partial^\mu j_\mu^8 | \eta' \rangle &= f_8 \sin \theta m_{\eta'}^2 = \frac{1}{\sqrt{6}} \langle 0 | i 2 m_u \bar{u} \gamma_5 u + i 2 m_d \bar{d} \gamma_5 d - i 4 m_s \bar{s} \gamma_5 s | \eta' \rangle, \\ \langle 0 | \partial^\mu j_\mu^0 | \eta \rangle &= -f_0 \sin \theta m_\eta^2 = \frac{1}{\sqrt{3}} \langle 0 | i 2 m_u \bar{u} \gamma_5 u + i 2 m_d \bar{d} \gamma_5 d + i 2 m_s \bar{s} \gamma_5 s + \frac{3 \alpha_s}{4 \pi} G_{\mu\nu} \tilde{G}^{\mu\nu} | \eta \rangle, \\ \langle 0 | \partial^\mu j_\mu^0 | \eta' \rangle &= f_0 \cos \theta m_{\eta'}^2 = \frac{1}{\sqrt{3}} \langle 0 | i 2 m_u \bar{u} \gamma_5 u + i 2 m_d \bar{d} \gamma_5 d + i 2 m_s \bar{s} \gamma_5 s + \frac{3 \alpha_s}{4 \pi} G_{\mu\nu} \tilde{G}^{\mu\nu} | \eta' \rangle. \end{aligned} \quad (4)$$

Neglecting the small up and down quark masses, we arrive at

$$\begin{aligned} \langle 0 | \frac{3 \alpha_s}{4 \pi} G_{\mu\nu} \tilde{G}^{\mu\nu} | \eta \rangle &= \sqrt{\frac{3}{2}} (\cos \theta f_8 - \sqrt{2} \sin \theta f_0) m_\eta^2, \\ \langle 0 | \frac{3 \alpha_s}{4 \pi} G_{\mu\nu} \tilde{G}^{\mu\nu} | \eta' \rangle &= \sqrt{\frac{3}{2}} (\sin \theta f_8 + \sqrt{2} \cos \theta f_0) m_{\eta'}^2. \end{aligned} \quad (5)$$

We emphasize that Eqs. (3)–(5) do not make explicit assumptions of the quark and gluon content of η and η' mesons. The parameters θ , f_8 and f_0 , however, are still not quite certain. In the following we shall use (A) $\theta = -17.0^\circ$, $f_8/f_\pi \simeq f_0/f_\pi = 1.06$ [13], and (B) $\theta = -22.0^\circ$, $f_8/f_\pi = 1.38$ and $f_0/f_\pi = 1.06$ [14,15], to illustrate the sensitivity to η - η' parameter values. To account for the observed central value of $Br(B \rightarrow \eta' X_s) = 6.2 \times 10^{-4}$, one needs $a_{\text{in}} \approx 0.012(0.015) \text{ GeV}^{-1}$, where the subscript stands for inclusive. The numbers are given for Set A (Set B) respectively, a notation which we shall employ from now on.

Assuming that Eq. (2) also dominates the large exclusive $Br(B \rightarrow \eta' K)$, one finds

$$A(B^- \rightarrow \eta' K^-) = a G_F \frac{4 \pi}{3} \sqrt{\frac{3}{2}} (\sin \theta f_8 + \sqrt{2} \cos \theta f_0) m_{\eta'}^2 \frac{m_B^2 - m_K^2}{2(m_b + m_s)} F_0^K(m_{\eta'}^2), \quad (6)$$

where the form factor $F_0^K(m_{\eta'}^2)$ is estimated to be 0.33 [16]. With the a_{in} values obtained earlier, we find

$Br(B \rightarrow \eta' K) \simeq 6.6 \times 10^{-5}$, which is consistent with experimental data. In turn, we could extract $a_{ex} = 0.012(0.016) \text{ GeV}^{-1}$ from $Br(B^\pm \rightarrow \eta' K^\pm)$, and $a_{ex} = 0.011(0.014) \text{ GeV}^{-1}$ from $Br(B_0 \rightarrow \eta' K^0)$, which are consistent with a_{in} .

We can easily obtain $Br(B \rightarrow \eta K(\eta X_s))$ by using Eq. (5). If the mechanism of Ref. [8] is the sole source for the large $B \rightarrow \eta' K(\eta' X_s)$ branching ratios, then $Br(B \rightarrow \eta K(\eta X_s))$ is simply related to $Br(B \rightarrow \eta' K(\eta' X_s))$ by

$$Br(B \rightarrow \eta K(\eta X_s)) = R^2 \Phi Br(B \rightarrow \eta' K(\eta' X_s)), \quad (7)$$

where

$$R = \frac{A(B \rightarrow \eta K(\eta X_s))}{A(B \rightarrow \eta' K(\eta' X_s))} = \frac{\cos \theta f_8 - \sqrt{2} \sin \theta f_0}{\sin \theta f_8 + \sqrt{2} \cos \theta f_0} \frac{m_\eta^2}{m_{\eta'}^2}, \quad (8)$$

and Φ is a phase space correction factor. For inclusive decays,

$$\Phi = \frac{\sqrt{(1 - (m_\eta + m_s)^2/m_b^2)(1 - (m_\eta - m_s)^2/m_b^2)}}{\sqrt{(1 - (m_{\eta'} + m_s)^2/m_b^2)(1 - (m_{\eta'} - m_s)^2/m_b^2)}}, \quad (9)$$

while for exclusive decays

$$\Phi = \frac{\sqrt{(1 - (m_\eta + m_K)^2/m_B^2)(1 - (m_\eta - m_K)^2/m_B^2)}}{\sqrt{(1 - (m_{\eta'} + m_K)^2/m_B^2)(1 - (m_{\eta'} - m_K)^2/m_B^2)}}. \quad (10)$$

Note that Eqs. (5) and (8) give $R \simeq 0.42(0.69)$. This means that the slight numerical change from Set A to Set B results in a factor of 2.7 difference in predictions for η modes. The branching ratio for inclusive $B \rightarrow \eta X_s$ is therefore

$$Br(B \rightarrow \eta X_s) = 1.1(3.0) \times 10^{-4} \quad (p_\eta > 2.1 \text{ GeV}), \quad (11)$$

which is still consistent with the experimental upper limit of 4.4×10^{-4} at the 90% C.L. for $2.1 \text{ GeV} < p_\eta < 2.7 \text{ GeV}$. For exclusive decays, using a_{in} from above gives

$$Br(B^- \rightarrow \eta K^-) = 1.2(3.2) \times 10^{-5}, \quad (12)$$

which is higher than the experimental limit of 0.8×10^{-5} (90% C.L.). If Set B is the case, then the Fritzsche scenario of Eq. (2) would be in grave difficulty. Scaling from exclusive $B \rightarrow \eta' K$ decay gives similar results. It is somewhat surprising that the one parameter model of Fritzsche, Eq. (2), could account for both $B \rightarrow \eta'$ and $B \rightarrow \eta$ data. Our results underpin the importance of improved experimental measurements in probing this mechanism.

It is interesting to speculate that, since this mechanism invokes the large gluon content of η' , there should be a large branching ratio for $B \rightarrow PK$, where P is the pseudoscalar glueball. To obtain the branching ratio for $B \rightarrow PK$, we need to know the matrix element $\langle 0 | \alpha_s G_{\mu\nu} \tilde{G}^{\mu\nu} | P \rangle$. This can be obtained by using QCD sum rules, as in Refs. [17,18]. Early calculations give $(3/4\pi) \langle 0 | \alpha_s G_{\mu\nu}^a \tilde{G}_{\mu\nu}^a | P \rangle \approx 0.87 f_\pi m_p^2$ for glueball mass around 1.4 GeV (i.e. taking $\eta(1440)$ as glueball candidate). Recent lattice calculations suggest, however, that glueball masses are heavier, i.e. $m_p = 2.3 \pm 0.2 \text{ GeV}$ [19], and the sum rule result could be quite different. We will follow the analysis in Ref. [18], but include the η contribution to the low energy spectral density which was previously neglected.

The basic idea of QCD sum rule analysis in the present case is to match the dispersion relation for the vacuum topological susceptibility $T(s)$ with the result found by using the operator product expansion. The definitions of T and the spectral density ρ are given by

$$\begin{aligned} T(s) &= i \int d^4x e^{iqx} \langle 0 | T(j_{\text{ps}}(x) j_{\text{ps}}(0)) | 0 \rangle, \quad j_{\text{ps}} = \frac{3\alpha_s}{4\pi} G_{\mu\nu}^a \tilde{G}_{\mu\nu}^a, \\ \rho(s) &= \rho^{\text{pole}}(s) + \rho^{\text{cont}}(s) \theta(s - s_1), \\ \rho^{\text{pole}}(s) &= \sum_i f_i^2 m_i^4 \delta(s - m_i^2), \quad \rho^{\text{cont}}(s) \simeq \left(\frac{3\alpha_s}{4\pi} \right)^2 \frac{2}{\pi} \left(1 + 5 \frac{\alpha_s}{\pi} \right) s^2 \equiv bs^2, \\ f_i m_i^2 &= \langle 0 | \frac{3\alpha_s}{4\pi} G_a^{\mu\nu} \tilde{G}_a^{\mu\nu} | i \rangle, \end{aligned} \quad (13)$$

where $i = P, \eta', \eta$, and s_1 is the continuum threshold. To improve series convergence, and to remove subtraction dependence, one makes the following Borel transformations on $T(s)$: $\int_0^\infty \text{Im } T(s) e^{-s/M^2} ds/s$ and $\int_0^\infty \text{Im } T(s) e^{-s/M^2} ds$, with M^2 a free parameter. One gets

$$\begin{aligned} & \int_0^\infty \frac{e^{-s/M^2}}{s} \rho^{\text{pole}}(s) ds \\ &= \int_0^{s_1} \frac{e^{-s/M^2}}{s} \rho^{\text{cont}}(s) ds + \pi \left(\frac{3\alpha_s}{4\pi} \right)^2 \left(D_4 + \frac{D_6}{M^2} + \dots \right) \\ & \times \int_0^\infty e^{-s/M^2} \rho^{\text{pole}}(s) ds = \int_0^{s_1} e^{-s/M^2} \rho^{\text{cont}}(s) ds + \pi \left(\frac{3\alpha_s}{4\pi} \right)^2 \left(D_6 + O\left(\frac{1}{M^2}\right) + \dots \right), \end{aligned} \quad (14)$$

where $D_4 = 4 \langle 0 | G_{\mu\nu}^a G_{\mu\nu}^a | 0 \rangle$ and $D_6 = 8 g_s f^{abc} \langle 0 | G_{\mu\alpha}^a G_{\alpha\beta}^b G_{\beta\mu}^c | 0 \rangle$. Numerically we use $\tilde{D}_{4,6} \equiv \pi(3\alpha_s/4\pi)^2 D_{4,6} = (1.44 \pm 0.61) \times 10^{-2} \text{ GeV}^4$, $(0.25 \pm 0.11) \times 10^{-2} \text{ GeV}^6$ [18] respectively. To further simplify, we follow Ref. [18] and take M^2 to infinity, leading to

$$f_P^2 m_P^2 + f_\eta^2 m_\eta^2 + f_{\eta'}^2 m_{\eta'}^2 = \frac{bs_1^2}{2} + \tilde{D}_4, \quad f_P^2 m_P^4 + f_\eta^2 m_\eta^4 + f_{\eta'}^2 m_{\eta'}^4 = \frac{bs_1^3}{3} + \tilde{D}_6. \quad (15)$$

Since $f_{\eta, \eta'}$ and $m_{\eta, \eta'}$ are known, f_P and m_P can be determined from knowing what to take for the continuum threshold s_1 .

Some information on s_1 can be obtained from comparing with lattice studies, which give $m_P = 2.3 \pm 0.2 \text{ GeV}$ in the quenched approximation. The quenched approximation is equivalent to QCD without quarks (but lattice studies incorporate running coupling with $N_f = 3$). Applying QCD sum rules, we obtain similar results as in Eq. (15), except that there are no η and η' mesons. Also, in the absence of quarks, there are no light meson states for glueball to decay into, hence the continuum threshold s_0 in this case would be larger than s_1 . Using the lattice glueball mass of 2.3 GeV , s_0 is determined to be 7.4 GeV^2 [18]. For the lower bound on s_1 , we require it to be larger than the glueball mass, since otherwise the sum rule approach breaks down. The glueball mass in QCD need not be the same as the lattice result in quenched approximation. To get an idea on how m_P and f_P can be determined, we take s_1 to range from 7.4 GeV^2 to 6 GeV^2 . Using the two sets of θ and $f_{0,8}$ discussed earlier and $\alpha_s = 0.33$ (at m_P scale), we get

$$\begin{aligned} m_P &= 2.42 \text{ GeV}, f_P = 1.15 (1.16) f_\pi, \quad \text{for } s_1 = 7.4 \text{ GeV}^2, \\ m_P &= 2.29 \text{ GeV}, f_P = 0.92 (0.94) f_\pi, \quad \text{for } s_1 = 6.0 \text{ GeV}^2. \end{aligned} \quad (16)$$

The difference between Set A and B is small in the given range for s_1 . In Ref. [18] the lower limit for s_1 was taken to be as low as 3 GeV^2 . We consider this to be too low for the following reason. The effect of including η in the sum rule is not significant for s_1 between 7.4 to 6.0 GeV^2 , but cannot be ignored for $s_1 \simeq 3 \text{ GeV}^2$. The inclusion of η increases the value for s_1 such that s_1 is larger than 3.24 GeV^2 (3.15 GeV^2) for Set A (Set B), and below these values there is no physical solution to the sum rule equations. When s_1 is close to the above lower bounds, the solution is in any case very sensitive to small changes.

We are now ready to calculate the branching ratio for $B \rightarrow PK (PX_s)$. Using the glueball masses and decay constants in Eq. (16), we obtain

$$\frac{Br(B \rightarrow PK (PX_s))}{Br(B \rightarrow \eta' K (\eta' X_s))} \simeq \begin{cases} 1.9 (3.2) & \text{for } m_P = 1.44 \text{ GeV} \\ 11 (20) & \text{for } m_P = 2.29 \text{ GeV} \\ 21 (36) & \text{for } m_P = 2.42 \text{ GeV} \end{cases} \quad (17)$$

As argued, the low glueball mass value of 1.44 GeV is not plausible. The numerical value for this case is taken from Ref. [17] for sake of illustration. The branching ratio is very sensitive to m_P because $\langle 0 | (3\alpha_s/4\pi) G_{\mu\nu} \tilde{G}^{\mu\nu} | P \rangle \propto m_P^2$. Hence, $Br(B \rightarrow PX_s)$ is at least twice as large as $Br(B \rightarrow \eta' K (\eta' X_s))$, and is likely as large as 1% , which is truly huge. Since the perturbative estimate for the inclusive $b \rightarrow sg^*$ penguin is itself $\simeq 1\%$ [20], this indicates once again that the Fritzsch mechanism is probably not the sole source for the large branching ratios for $B \rightarrow \eta' K (\eta' X_s)$. However, it does suggest that one should search for the pseudoscalar glueball in B decays, since the scenario is not yet completely ruled out. For exclusive decay, one finds $Br(B^\pm \rightarrow PK^\pm) \sim 0.1\%$ or even higher for $m_P \simeq 2.3 \text{ GeV}$. Taking $Br(\eta_c \rightarrow K\bar{K}\pi) \simeq 6.6\%$ as a guide, one should search for a fast $K^\pm$ ($p_K \simeq 2.1 \text{ GeV}$) recoiling against a $K\bar{K}\pi$ system with mass $\sim 2.2\text{--}2.3 \text{ GeV}$. Background should be small since both $B \rightarrow D_s K + X$ and usual $D \rightarrow K\bar{K} + X$ decays are suppressed. Such a large rate for $B \rightarrow PK$ should be readily observable if Γ_P is not too broad, which is likely the case since the 2^{++} glueball candidate, the $\xi(2220)$, has a width of only 20 MeV [21].

Before we conclude, let us comment on the predictions for $B \rightarrow \eta X_s$ and PX_s for the mechanism proposed in Refs. [5,6]. It was proposed in Ref. [5] that the coupling of η' to two gluons via the gluon anomaly may be able to explain the semi-inclusive decay by the process $b \rightarrow sg^* \rightarrow sg\eta'$ within the SM. However, it was subsequently pointed out that, because the effective η' - g - g coupling arises from the Wess-Zumino term [13,22],

$$L_{\text{eff}} = \alpha_s \frac{\sqrt{N_F}}{4\pi f_0} \eta_0 G_{\mu\nu} \tilde{G}^{\mu\nu}, \quad (18)$$

it is proportional to α_s . If this runs with the q^2 of the virtual gluon g^* , the branching ratio is reduced by a factor of 3 [6]. Enhanced bsg coupling due to new physics is then needed [6], with the tantalizing prospect of potentially large CP violation effects. A critical question is the form factor behavior of Eq. (18). Analogy with the photon anomaly suggests $1/k^2$ suppression [9], where k is the difference between the two gluon momenta. However, the analogy fails precisely [6] because of binding effects in the two gluon channel, with the normalization of Eq. (18) fixed by the η - η' system at low energy. The study of the off-shell behavior of the gluon anomaly has yet to begin.

Turning back to phenomenology, since one of the gluons appear in the final state, one can no longer use Eq. (5). Using naive η - η' mixing as an estimate of the ratio $A(b \rightarrow sg\eta)/A(b \rightarrow sg\eta') \sim \tan\theta$, the inclusive $Br(B \rightarrow \eta X_s)$ could well be below the experimental upper limit. For the exclusive process, as the actual transition is $b\bar{q} \rightarrow \eta' + sg\bar{q}$ and there are no IR singularities, it was noted [6] that if the $sg\bar{q}$ system evolves into the K meson when the gluon is soft, it could account for the exclusive $B \rightarrow \eta' K$ rate. However, since it is

difficult to make quantitative predictions, the large exclusive $B \rightarrow \eta' K$ branching ratio could well be coming from other mechanisms, such as from four quark operators in the SM with appropriate form factors.

To evaluate the branching ratio for $B \rightarrow PX_s$, we need to estimate the Pgg coupling. For sake of illustration we use an Ansatz given in Ref. [23],

$$L_{\text{eff}} = -\frac{2c\Lambda}{k^2} P G_{\mu\nu} \tilde{G}^{\mu\nu}, \quad (19)$$

where k is the difference of the two gluon momenta. For $g^* \rightarrow Pg$ where P and g are both on-shell, one has $k^2 = 2q^2 - m_p^2$ where q is the virtual gluon momentum. But in the small k^2 limit, Eq. (19) could generate finite glueball mass [23] at one loop level, giving

$$m_p^2 = \frac{16c^2\Lambda^2}{\pi^2} (4\ln 2 - 1). \quad (20)$$

Of course there should be other contributions to m_p^2 , but saturating m_p^2 with the above relation gives the upper bound $(c_1\Lambda)_{\text{max}} = 0.59 m_p$, and Eq. (19) becomes

$$L_{\text{eff}} \simeq -\frac{1.2 m_p^2}{2q^2 - m_p^2} \frac{P}{m_p} G_{\mu\nu} \tilde{G}^{\mu\nu}. \quad (21)$$

Comparing with Eq. (18), the most salient feature is the absence of an α_s factor. It also highlights our ignorance of the form factor behavior of Eq. (18). We emphasize that $m_{\eta'}$ does not arise from $\eta'-P$ mixing, nor from some Ansatz analogous to Eq. (19), but from topological effects through the gluon anomaly [22]. Since the effective P - g - g coupling of Eq. (21) is only slightly larger than the η' - g - g coupling of Eq. (18), we find $Br(B \rightarrow PX_s) \sim > Br(B \rightarrow \eta' X_s)$ is likely. Though sizable, the branching ratios for $B \rightarrow PX_s$ and PK in the scenario of Refs. [5,6] are not as large as in the case of Ref. [8]. However, more quantitative predictions in the present case are difficult, especially for the exclusive case.

To conclude, we have shown that the mechanism proposed by Fritzsche for explaining large $Br(B \rightarrow \eta' K (\eta' X_s))$ gives large branching ratios for B decays with an η or a glueball in the final state. The predicted branching ratio for $B \rightarrow \eta K$ is barely within the experimental upper limit, while $Br(B \rightarrow PK)$ could be as large as 0.1%, with $Br(B \rightarrow PX_s)$ ten times larger. If this is the case, perhaps the pseudoscalar glueball P could be first discovered in B decays. Improved experiments will provide decisive information about this mechanism.

Acknowledgements

This work was supported in part by the ARC, the NSC, and the NNSF. X.G.H. acknowledges the support of K.C. Wong Education Foundation, Hong Kong.

References

- [1] CLEO Collaboration, R. Godang et al., CLNS 97-1522, hep-ex/9711010; P. Kim, talk presented at FCNC97, Santa Monica, California, February 1997; J. Alexander, B. Behrens at BCONF97, Honolulu, Hawaii, March 1997; D.H. Miller, talk presented at EPS-HEP97, Jerusalem, Israel, August 1997; A. Anderson et al., CLEO CONF 97-221, 1997; D.M. Asner et al., CLEO CONF 97-13; J. Smith, talk presented at the 1997 Aspen Winter Physics Conference on Particle Physics, January 1997.
- [2] A. Datta, X.-G. He, S. Pakvasa, hep-ph/9707259.
- [3] A. Ali, C. Greub, hep-ph/9707251.
- [4] H.-Y. Cheng, B. Tseng, hep-ph/9707316.
- [5] D. Atwood, A. Soni, Phys. Lett. B 405 (1997) 150.
- [6] W.-S. Hou, B. Tseng, hep-ph/9705304, to appear in Phys. Rev. Lett.

- [7] I. Halperin, A. Zhitnitsky, hep-ph/9705251.
- [8] H. Fritzsch, hep-ph/9708348.
- [9] A. Kagan, A. Petrov, hep-ph/9707354.
- [10] F. Yuan, K.-T. Chao, *Phys. Rev. D* 56 (1997) 2495; D.S. Du, C.S. Kim, Y. Yang, hep-ph/9711428; M.R. Ahmady, E. Kou, A. Sugomoto, hep-ph/9710509.
- [11] D.S. Du, M.Z. Yang, hep-ph/9711272.
- [12] X.-G. He et al., *Phys. Rev. Lett.* 64 (1990) 1003; Jiang Liu, Y.-P. Yao, *Phys. Rev. D* 41 (1990) 2147.
- [13] R. Akhoury, J.-M. Frère, *Phys. Lett. B* 220 (1989) 258; P. Ball, J.-M. Frère, M. Tytgat, *Phys. Lett. B* 365 (1996) 367.
- [14] E. Venugopal, B. Holstein, hep-ph/9710382.
- [15] Particle Data Group, *Phys. Rev. D* 54 (1996) 1.
- [16] M. Bauer, B. Stech, *Phys. Lett. B* 152 (1985) 380; M. Bauer, B. Stech, M. Wirbel, *Z. Phys. C* 34 (1987) 103.
- [17] V. Novikov et al., *Phys. Lett. B* 86 (1979) 347; K. Senba, M. Tanimoto, *Phys. Lett. B* 105 (1981) 297.
- [18] G. Gabadadze, Preprint, hep-ph/9711380.
- [19] G.S. Bali et al., UKQCD Collaboration, *Phys. Lett. B* 309 (1993) 378.
- [20] W.S. Hou, A. Soni, H. Steger, *Phys. Rev. Lett.* 59 (1987) 1521; W.S. Hou, *Nucl. Phys. B* 308 (1988) 561.
- [21] J.Z. Bai et al., BES Collaboration, *Phys. Rev. Lett.* 76 (1996) 3502.
- [22] G.M. Shore, G. Veneziano, *Nucl. Phys. B* 381 (1992) 2; P.H. Damgaard, H.B. Nielsen, R. Sollacher, *Nucl. Phys. B* 414 (1994) 541, and references therein.
- [23] G.J. Gounaris, *Nucl. Phys. B* 320 (1989) 253.