Eur. Phys. J. C 23, 691–699 (2002)
Digital Object Identifier (DOI) 10.1007/s100520200894

THE EUROPEAN
PHYSICAL JOURNAL C

B meson decays to baryons in the diquark model

C.-H.V. Chang[1], W.-S. Hou[2]

[1] Department of Physics, National Taiwan Normal University, Taipei, Taiwan, R.O.C.
[2] Department of Physics, National Taiwan University, Taipei, Taiwan, R.O.C.

Received: 27 December 2001 /
Published online: 5 April 2002 –

Abstract. We study B meson decays to two charmless baryons in the diquark model, including strong and electroweak penguins as well as the tree operators. It is shown that penguin operators can enhance $\bar{B} \to \boldsymbol{B}_s\bar{\boldsymbol{B}}$ considerably, but affect $\bar{B} \to \boldsymbol{B}_1\bar{\boldsymbol{B}}_2$ only slightly, where $\boldsymbol{B}_{(1,2)}$ and $\boldsymbol{B}_s$ are non-strange and strange baryons, respectively. The γ dependence of the decay rates due to tree–penguin interference is illustrated. In principle, some of the $\boldsymbol{B}_s\bar{\boldsymbol{B}}$ modes could dominate over $\boldsymbol{B}_1\bar{\boldsymbol{B}}_2$ for $\gamma > 90°$, but in general the effect is milder than their mesonic counterparts. This is because the O_6 operator can only produce vector but not scalar diquarks, while the opposite is true for O_1 and O_4. Predictions from the diquark model are compared to those from the sum rule calculation. The decays $\bar{B} \to \boldsymbol{B}_s\bar{\boldsymbol{B}}_s$ and inclusive baryonic decays are also discussed.

1 Introduction

B meson decays provide a unique setting for baryon pair production, since this is impossible for the D system. Once observed, these decays could shed light on our understanding of baryon production, and may offer further probes [1] of underlying weak decay dynamics such as CP violating phases.

Many rare mesonic B decays have been observed at the 10^{-5} level in recent years, heralding the start of the "B Factory" era. However, rare baryonic decays have yet to be discovered. The most recent published limits come from the CLEO collaboration [2]. Based on 5.8 million $B\bar{B}$ events, CLEO finds $B \to \bar{\Lambda}p$, $\bar{\Lambda}p\pi^-$ and $p\bar{p} < 0.26, 1.3$, and 0.7×10^{-5}, respectively. There was some 2.8σ excess in the $\bar{B}^0 \to p\bar{p}$ channel, but this was insufficient to claim discovery. The B factories, i.e. the Belle and BaBar Collaborations, have now each accumulated an order of magnitude more data. A preliminary result from Belle [3] has pushed the $\bar{B}^0 \to p\bar{p}$ limit down to the 10^{-6} level. This rules out the CLEO hint, and puts two-body baryonic modes in strong contrast to the corresponding mesonic modes. However, though still elusive, it is quite possible that charmless baryonic modes are just around the corner.

Theoretical work on rare baryonic decays is sparse. Most of it was stimulated by the surprising (and false [4]) 1987 results [5] of $\mathrm{Br}(B^- \to p\bar{p}\pi^-) = (3.7 \pm 1.3 \pm 1.4) \times 10^{-4}$ and $\mathrm{Br}(\bar{B}^0 \to p\bar{p}\pi^+\pi^-) = (6.0 \pm 1.3 \pm 1.4) \times 10^{-4}$ from the ARGUS Collaboration. Pole models [6,7] and the sum rule approach [8] have been proposed for calculating the two-body decay widths, while Ball and Dosch [9] pursued the diquark model appoach. All of these works are now a decade old. Not until very recently, sensing that experiment is about to move forward, did theorists start to pay attention again. Hou and Soni [1] pointed out the need for reduced energy release on the baryon side, e.g. charmless baryonic B decays may be more prominent in association with η' or γ. Chua, Hou and Tsai studied $\bar{B} \to D^{*-}N\bar{N}$ [10] and $B^0 \to \rho^- p\bar{n}$, $\pi^- p\bar{n}$ [11] using a factorization approach of current produced baryons.

For two-body baryonic decays, the sum rule and diquark model approaches are the most relevant. The sum rule calculation [8] predicts that the branching ratio of $\bar{B} \to \boldsymbol{B}_1\bar{\boldsymbol{B}}_2$ is typically of order 10^{-6} while $\mathrm{Br}(\bar{B} \to \boldsymbol{B}_s\bar{\boldsymbol{B}}) \approx (0.3 \sim 1.0) \times 10^{-5}$ and $\mathrm{Br}(\bar{B} \to B_c\bar{B}_{(c)}) \approx \mathrm{O}(10^{-3})$. Here $\boldsymbol{B}_{(1,2)}$, $\boldsymbol{B}_s$ and $\boldsymbol{B}_c$ denote non-strange, strange and charmed baryons, respectively. The CLEO limit of $\mathrm{Br}(\bar{B}^0 \to p\bar{\Lambda}) < 0.26 \times 10^{-5}$ [2] is already at odds with the sum rule prediction of $\mathrm{Br}(B^+ \to p\bar{\Lambda}) \sim 1.0 \times 10^{-5}$. Furthermore the predicted $\mathrm{Br}(\bar{B}^0 \to p\bar{p}) \sim 1.0 \times 10^{-6}$ is an order of magnitude smaller than $B^- \to p\bar{\Lambda}$. This is in contrast with the diquark model [9] which typically gives a larger rate for $\bar{B} \to \boldsymbol{B}_1\bar{\boldsymbol{B}}_2$ over $\boldsymbol{B}_s\bar{\boldsymbol{B}}$, although it can predict really only relative rates. For example, the $\bar{B}^0 \to p\bar{p}$ mode has larger rate than all the $\bar{B} \to \boldsymbol{B}_s\bar{\boldsymbol{B}}$ decays. Unfortunately, [9] did not include penguin operators. Judging from the role played by penguins in mesonic B decays like $B \to \pi\pi$, $K\pi$, an enhancement by penguins of $\mathrm{Br}(\bar{B} \to \boldsymbol{B}_s\bar{\boldsymbol{B}})$ is to be expected. Furthermore, with penguins ignored, $\mathrm{Br}(B^- \to \Lambda\bar{p}) = 0$ in the diquark model, preventing us from comparing with Belle and CLEO search directly.

The purpose of this paper is to complete the diquark model treatment of two-body charmless baryonic B decays by including penguin diagrams, and to assess the importance of these penguin effects. Whether the approach of using diquarks to describe baryon formation is correct or not is still an open question. Even if one assumes that the idea is reasonable, to calculate the relative baryonic decay rates, we still need to make many other dynamical assumptions, which introduce further uncertainties. Our goal in this paper is simply to clarify the actual predictions of the diquark model by expanding on the work of [9]. The experimental measurements in the future will decide whether the diquark model or the sum rule approach is more relevant in the description of baryonic decays, or how they might be improved upon. Thus, we do not attempt to improve the diquark model towards absolute rate calculations.

We find that penguin operators indeed could enhance $\bar{B} \to \boldsymbol{B}_s \bar{\boldsymbol{B}}$ decay rate by a factor of ~ 5 for $\gamma = 90°$. Due to tree–penguin interference, the decay widths now depend on the unitarity phase angle γ ($\equiv \arg V_{ub}^*$, in the convention of PDG [12]). The enhancement in baryonic B decays is milder than in the mesonic decays because the operator O_6 cannot generate scalar diquarks. Penguins also affect non-strange decays, $\bar{B} \to \boldsymbol{B}_1 \bar{\boldsymbol{B}}_2$, introducing also a γ dependence, but the effect here is small. In general, $\bar{B} \to \boldsymbol{B}_s \bar{\boldsymbol{B}}$ is still smaller than $\bar{B} \to \boldsymbol{B}_1 \bar{\boldsymbol{B}}_2$, and $\bar{B}^0 \to p\bar{p}$ typically has the largest rate. But for large $\gamma > 90°$, the $\bar{B} \to \Sigma^+ \bar{p}$, $\Sigma^+ \overline{\Delta^{++}}$ rate could become larger than $\bar{B}^0 \to p\bar{p}$. We find that the pattern of decay widths calculated using diquark model is quite different from the sum rule results. More experimental data should shed light on the two models.

This paper is organized as follows. We first review the diquark model and baryonic B decays. The connection between penguin and diquark operators is discussed. In Sect. 3, we study inclusive baryonic decays, with two-body exclusive decays discussed in Sect. 4, in both cases including the effect of penguins. The conclusion is given in the last section.

2 Diquark model and penguin operators

It is well known that the strong force between two quarks in a color-antitriplet combination is attractive, hence it has been speculated for a long time that they will form a bound or correlated state, called the diquark. The flavor antisymmetric combinations form scalar diquarks, while flavor symmetric combinations form vector diquarks. The diquark picture is useful in the description of baryons. The spin 1/2 octet and spin 3/2 decuplet baryons can be understood as bound states of a quark and a scalar or vector diquark, respectively [13].

Diquarks can be generated in weak decays [14–16]. The tree level weak decay effective Hamiltonian is

$$\mathcal{H}_{\text{eff}} = \frac{G_F}{\sqrt{2}} \sum_{i=u,c} \left\{ V_{iq}^* V_{ib} [c_1(\mu) O_1^i(\mu) + c_2(\mu) O_2^i(\mu)] \right\} + \text{h.c.}, \quad (1)$$

where

$$O_1^u = (\bar{q}_\beta u_\alpha)_{\text{V-A}} (\bar{u}_\alpha b_\beta)_{\text{V-A}},$$
$$O_2^u = (\bar{q} u)_{\text{V-A}} (\bar{u} b)_{\text{V-A}},$$

and likewise for $O_{1,2}^c$. The quark q could be d or s. The current–current operators are defined by $(\bar{q}_1 q_2)_{\text{V-A}} \equiv \bar{q}_1 \gamma_\mu (1-\gamma_5) q_2$. The operators could be rewritten, after a Fierz transformation, in terms of scalar diquark field operators. For example

$$[c_1 O_1^u + c_2 O_2^u] = -(c_2 - c_1)(ub)_{\text{L}k} (ud)_{\text{L}k}^* + \text{color sextet current}, \quad (2)$$

with the scalar diquark field operator defined by

$$(qQ)_{\text{L}k} = \epsilon_{klm} (\bar{q}_l^C (1-\gamma_5) Q_m), \quad (3)$$

where k, l, m are color indices. Here $\bar{q}^C \equiv q^{\text{T}} C$. The scalar diquark field operator can create a scalar diquark from the vacuum with the strength g_{qQ}, usually called the "diquark decay constant":

$$\langle 0 | \epsilon_{klm} (\bar{q}_l^C \gamma_5 Q_m) | (qQ)_k^{0+} \rangle \equiv \sqrt{\frac{2}{3}} \delta_{il} g_{qQ}. \quad (4)$$

We see from (2) that a b quark can decay into a scalar diquark plus an antiquark. Since the baryons are bound states of a diquark and a quark, it is natural to expect that in decays to baryonic final states the diquark operators will dominate and the sextet current operators can be ignored. Following this reasoning, [9] gives a picture of two-body baryonic B decays. The antiquark produced in the decay combines with the spectator antiquark to form a scalar or vector antidiquark. As the diquark and antidiquark fly apart, they pull a quark–antiquark pair from the vacuum, resulting finally in a baryon–antibaryon pair. Of course, this may not be the only mechanism, but it is assumed to be the dominant one in the diquark model [9]. It is interesting to note that $O_{1,2}$ can only generate scalar diquarks, hence $\bar{B}$ decaying to a decuplet baryon plus either an octet or decuplet antibaryon, $\bar{B} \to \boldsymbol{B}^* \overline{\boldsymbol{B}}^{(*)}$, are predicted to have small rates [9]. These decays can only arise from the penguin operators $O_{5,6}$, as will be discussed below.

In addition to the tree level effective Hamiltonian, it is well known that penguin diagrams are important for the charmless decays of B mesons. In mesonic decays like $B \to K\pi$ and $\pi\pi$, penguin diagrams are crucial in the calculation of decay rates and CP asymmetries. Their effects can be described by the effective penguin operators O_3 through O_{10},

$$\mathcal{H}_{\text{penguin}} = -\frac{G_F}{\sqrt{2}} \left\{ V_{tq}^* V_{tb} \sum_{i=3}^{10} c_i(\mu) O_i(\mu) \right\} + \text{h.c.}, \quad (5)$$

where

$$O_{3(5)} = \sum_{q'} (\bar{q}' q')_{\text{V-A(V+A)}} (\bar{q} b)_{\text{V-A}},$$
$$O_{4(6)} = \sum_{q'} (\bar{q}'_\beta q'_\alpha)_{\text{V-A(V+A)}} (\bar{q}_\alpha b_\beta)_{\text{V-A}},$$
$$O_{7(9)} = \frac{3}{2} \sum_{q'} e_{q'} (\bar{q}' q')_{\text{V+A(V-A)}} (\bar{q} b)_{\text{V-A}},$$
$$O_{8(10)} = \frac{3}{2} \sum_{q'} e_{q'} (\bar{q}'_\beta q'_\alpha)_{\text{V+A(V-A)}} (\bar{q}_\alpha b_\beta)_{\text{V-A}}, \quad (6)$$

with O_{3-6}, O_{7-10} the QCD and electroweak penguin operators, respectively, and $q = d, s$, $(\bar{q}_1 q_2)_{\mathrm{V\pm A}} \equiv \bar{q}_1\gamma_\mu(1 \pm \gamma_5)q_2$. The sum over q' runs over all quark flavors that exist in the effective field theory.

The penguin operators can also be written in terms of diquark operators. Let us consider $b \to s$ penguins (a similar discussion follows for the $q = d$ case). The $(\mathrm{V-A})\times(\mathrm{V-A})$ type penguin operators $O_{3,4}$ and $O_{9,10}$ are similar in form to $O_{1,2}$. They can be written, after a Fierz transformation, as the sum of diquark operators $\pm\sum_{q'}(q'b)_{\mathrm{L}}(q's)^\dagger_{\mathrm{L}}$ plus color sextet terms. Again the latter will be ignored. The sum includes the operator $(ub)_{\mathrm{L}}(us)^\dagger_{\mathrm{L}}$, which is of exactly the same form as that obtained from $O_{1,2}$, as well as $(db)_{\mathrm{L}}(ds)^\dagger_{\mathrm{L}}$, which is not present in the tree operators. Note that the $q' = s$ piece $(sb)_{\mathrm{L}}(ss)^\dagger_{\mathrm{L}}$ vanishes since the scalar diquark is antisymmetric in their flavor constituents.

On the other hand, the $(\mathrm{V-A})\times(\mathrm{V+A})$ type penguin operators O_5 and O_6 do not give rise to scalar diquark operators. The reason is that a scalar diquark is a quark–quark correlation. The operators $O_{5,6}$ contains one left-handed and one right-handed quark field, which cannot form a scalar combination under the Lorentz transformation. In other words, the scalar diquark content of octet baryons implies that the operators $O_{5,6}$ (and the corresponding electroweak penguin operators $O_{7,8}$) do not contribute to the $\bar{B}$ decays to octet baryon plus either an octet or decuplet antibaryon: $B \to \boldsymbol{B}\overline{\boldsymbol{B}}^{(*)}$. This is contrary to the significant role that they play in the mesonic decays such as $B \to K\pi$, $\pi\pi$. However, $O_{5,6}$ could generate operators that consist of vector diquarks which O_1 through O_4 could not produce. By Fierz transformation, for example,

$$\begin{aligned} &c_5(\bar{s}\gamma^\mu b)_{\mathrm{V-A}}(\bar{q'}\gamma_\mu q')_{\mathrm{V+A}} \\ &\quad + c_6(\bar{s}_\alpha\gamma^\mu b_\beta)_{\mathrm{V-A}}(\bar{q'}_\beta\gamma_\mu q'_\alpha)_{\mathrm{V+A}} \\ &\quad = -\frac{1}{2}(c_6 - c_5)(q'_{\mathrm{R}}b_{\mathrm{L}})^*_k(q'_{\mathrm{R}}s_{\mathrm{L}})^{*\dagger}_k + \text{color sextet}, \end{aligned} \tag{7}$$

with the vector diquark field operator defined by

$$(q_{\mathrm{R}}Q_{\mathrm{L}})^*_k \equiv \epsilon_{klm}(\bar{q}^C_l\gamma^\mu(1-\gamma_5)Q_m). \tag{8}$$

Since decuplet baryons are bound states of a vector diquark and a quark, $O_{5,6}$ will produce $\bar{B}$ decays to a decuplet baryon plus either an octet or decuplet antibaryon: $\bar{B} \to \boldsymbol{B}^*\overline{\boldsymbol{B}}^{(*)}$. As mentioned above, these decays cannot be generated by the tree $O_{1,2}$ and the penguin $O_{3,4}$ operators. For example, the novel channel $B^- \to \Omega^-\overline{\Xi}^0$ (five strange quarks in the final state) could arise from $B^- \to (ss)^*(\overline{su})$.

To sum up, the effective Hamiltonian that generates scalar diquarks can now be collected as

$$\begin{aligned} \mathcal{H}_{\text{diquark}} \sim -\frac{G_{\mathrm{F}}}{\sqrt{2}}\Big\{&\mathcal{A}_1(ub)_{\mathrm{L}}(ud)^\dagger_{\mathrm{L}} + \mathcal{A}_2(ub)_{\mathrm{L}}(us)^\dagger_{\mathrm{L}} \\ &+\mathcal{A}_3(sb)_{\mathrm{L}}(sd)^\dagger_{\mathrm{L}} + \mathcal{A}_4(db)_{\mathrm{L}}(ds)^\dagger_{\mathrm{L}} + \text{h.c.}\Big\}, \end{aligned}$$

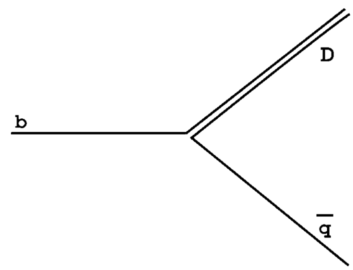

Fig. 1. The inclusive decay $b \to \boldsymbol{D}\bar{q}$

with the coefficients

$$\begin{aligned} \mathcal{A}_1 &\equiv V^*_{ud}V_{ub}(c_2 - c_1) - V^*_{td}V_{tb}(c_4 - c_3 + c_9 - c_{10}), \\ \mathcal{A}_2 &\equiv V^*_{us}V_{ub}(c_2 - c_1) - V^*_{ts}V_{tb}(c_4 - c_3 + c_9 - c_{10}), \\ \mathcal{A}_3 &\equiv -V^*_{td}V_{tb}\left(c_4 - c_3 - \frac{1}{2}c_9 + \frac{1}{2}c_{10}\right), \\ \mathcal{A}_4 &\equiv -V^*_{ts}V_{tb}\left(c_4 - c_3 - \frac{1}{2}c_9 + \frac{1}{2}c_{10}\right). \end{aligned} \tag{9}$$

For vector diquarks, we have

$$\begin{aligned} \mathcal{H}_{\text{diquark}*} \sim -\frac{G_{\mathrm{F}}}{\sqrt{2}}\sum_{q'=u,d,s}\Big\{&\mathcal{B}_1(q'_{\mathrm{R}}b_{\mathrm{L}})^*(q'_{\mathrm{R}}d_{\mathrm{L}})^{*\dagger} \\ &+\mathcal{B}_2(q'_{\mathrm{R}}b_{\mathrm{L}})^*(q'_{\mathrm{R}}s_{\mathrm{L}})^{*\dagger} + \text{h.c.}\Big\}, \end{aligned} \tag{10}$$

with the coefficients

$$\begin{aligned} \mathcal{B}_1 &\equiv -\frac{1}{2}V^*_{td}V_{tb}\left(c_5 - c_6 + \frac{3}{2}e_{q'}c_7 - \frac{3}{2}e_{q'}c_8\right), \\ \mathcal{B}_2 &\equiv -\frac{1}{2}V^*_{ts}V_{tb}\left(c_5 - c_6 + \frac{3}{2}e_{q'}c_7 - \frac{3}{2}e_{q'}c_8\right). \end{aligned} \tag{11}$$

3 Inclusive baryonic decays

Before we discuss the more difficult exclusive baryonic decays, the diquark picture could actually give us useful insight in the inclusive baryonic B decays.

In the diquark model, one postulates that baryons are bound states of a (scalar or vector) diquark and an antiquark. It is natural to expect that baryonic B decays proceed dominantly via the process of b quark decaying into a scalar diquark plus an antiquark. Since the subsequent hadronization process always generates at least one baryon, with the other antibaryon guaranteed by baryon number conservation, the inclusive baryonic decay rates can be approximated by the rates of $b \to \boldsymbol{D}\bar{q}$ (Fig. 1), where $\boldsymbol{D}$ denotes a scalar diquark such as (ud) or (cd).

This is the approach mentioned by Neubert and Stech in [15,17]. Though citing the result from the above two papers, [9] adopted a different method to calculate the inclusive rates. It computes the rates of the B meson decaying into a diquark and an antidiquark, i.e. $B \to \boldsymbol{D}\overline{\boldsymbol{D}'}$. The two-body channel $B \to \boldsymbol{D}\overline{\boldsymbol{D}'}$, in which both diquarks are fast moving, implies that at least two fast moving baryons will be generated. Though baryons always appear in pairs in baryonic decays, this could be too restrictive for inclusive baryonic decays, since it ignores the possibility of slowly moving antibaryons.

The $b \to \boldsymbol{D}\bar{q}$ rate turns out to be rather close to the observed inclusive rate. Let us take a closer look. After the $b \to \boldsymbol{D}\bar{q}$ decay, the diquark $\boldsymbol{D}$, the antiquark $\bar{q}$ and the spectator antiquark jointly form a color singlet, just like the three antiquarks in an antibaryon. The fast moving diquark $\boldsymbol{D}$ pulls a quark q' from the vacuum to form a baryon, leaving behind a slow antiquark $\bar{q}'$. The color configuration of $\bar{q}'$, $\bar{q}$ and the spectator antiquark is again just like that of an antibaryon. Since $\bar{q}$ is moving fast while the other two are slow, the system breaks up into hadrons through fragmentation. This generates all kinds of possible final products, but at least one antibaryon has to be generated due to baryon number conservation. One possible scenario is for $\bar{q}$ to form a fast moving antibaryon by pulling the two antiquarks with it. Two body baryonic decays are just such a case. Another scenario is that the fast moving $\bar{q}$ captures one quark to form a meson, leaving behind a slower antiquark. The final products of the decay then consist of a fast baryon, a fast meson and a slow antibaryon plus possible soft mesons. The baryon–antibaryon pair mass would then be far below m_B. One could also break two strings and capture two new antiquarks to form a fast antibaryon with the remaining quarks and antiquarks combining into mesons. The final products would then be one fast moving baryon, one fast antibaryon, plus two (or more) soft mesons. In all the above scenarios, one baryon and one antibaryon are generated. But the second scenario clearly is not included in the $B \to \boldsymbol{D}\overline{\boldsymbol{D}'}$ picture of [9].

We list the branching ratios of $b \to \boldsymbol{D}\bar{q}$ decays in Table 1. The transition amplitude of $b \to \boldsymbol{D}\bar{q}$ is assumed to factorize into the product of the diquark decay constant as defined in (4), and the quark level amplitude $\langle q|(\bar{q}b)|b\rangle$. For comparison, we also list the corresponding numbers for $\bar{B}^0$ decay from [9] by adding up appropriate diquark–antidiquark decay rates. For example, the rate of $b \to (cd)\bar{u}$ would correspond to the sum of the rates of $\bar{B}^0 \to (cd)(\overline{ud})$, $(cd)(\overline{ud})^*$. Unlike the $b \to \boldsymbol{D}\bar{q}$ case, calculating the latter not only involves diquark decay constants and masses, it also depends on $\bar{B}$ meson to diquark form factors. The diquark decay constants we use are [9, 18]

$$g_{ud}, g_{us} = 0.179, 0.215\,\mathrm{GeV}^2,$$
$$g_{cd} \cong g_{cs} \cong 0.35\,\mathrm{GeV}^2, \tag{12}$$

and the diquark masses are

$$m_{ud}, m_{us}, m_{cd}, m_{cs} = 0.5, 0.7, 1.7, 2.0\,\mathrm{GeV}. \tag{13}$$

As expected, the inclusive baryonic decays are dominated by the two charmed modes: $b \to (cd)\bar{u}$ and $b \to (cs)\bar{c}$. The combined branching ratio is about 4.6%. Adding in the rates of the smaller modes $b \to (cs)\bar{u}$ and $b \to (cd)\bar{c}$ gives a prediction for the total inclusive baryonic B decay branching ratio of

$$\mathrm{Br}(B \to \mathrm{baryon} + X) = 4.8\%. \tag{14}$$

Table 1. Estimate of inclusive baryonic branching ratios. The line separates charmed vs. charmless final states

	$b \to \boldsymbol{D}\bar{q}$	$B \to \boldsymbol{D}\overline{\boldsymbol{D}'}$ [9]
$b \to (cd)\bar{u}$	2.2×10^{-2}	2.0×10^{-3}
$b \to (cs)\bar{c}$	2.4×10^{-2}	5.8×10^{-3}
$b \to (cs)\bar{u}$	1.1×10^{-3}	2.0×10^{-4}
$b \to (cd)\bar{c}$	1.3×10^{-3}	3.0×10^{-4}
$b \to (ud)\bar{u}$	4.8×10^{-5}	5.2×10^{-6}
$b \to (us)\bar{u}$	2.1×10^{-5}	3.9×10^{-7}
$b \to (ds)\bar{d}$	2.0×10^{-5}	0

This is in reasonably good agreement with the experimental result [12]:

$$\mathrm{Br}\,(B \to \mathrm{baryon} + X) = 6.8 \pm 0.6\%. \tag{15}$$

The minor deficit is to be expected in consideration of the possibility of decaying into vector diquarks and other excited states as well as other mechanisms such as current produced baryons [10, 11]. Though [9] quotes a reasonable prediction from [15], their approach of simply adding up the $B \to \boldsymbol{D}\overline{\boldsymbol{D}'}$ rates would have given a branching ratio that is too small, $\sim 0.8\%$. This is an indication that $B \to \boldsymbol{D}\overline{\boldsymbol{D}'}$ is not inclusive enough. In fact, our discussion shows that this is rather an estimate of the fraction of the baryonic events where both baryons are energetic.

We make some observations before turning to exclusive modes. For baryonic decays, the single charm channel $b \to (cd)\bar{u}$ has roughly the same rate as the double charm channel $b \to (cs)\bar{c}$ since the decay constants g_{cd} and g_{cs} are equal and both have two-body phase space. This is different from the quark level picture for inclusive b decays, where $b \to c\bar{c}s$ is suppressed by a factor of 3–5 compared to $b \to c\bar{u}d$ because of having two massive final quarks in three body phase space. In [9], i.e. for $B \to \boldsymbol{D}\overline{\boldsymbol{D}'}$, the difference is even more dramatic: $b \to (cs)\bar{c}$ is more than twice $b \to (cd)\bar{u}$ because of $B \to \boldsymbol{D}$ form factors. This feature of the diquark model can be tested by experiment. For example, one can study the inclusive decays of $\bar{B} \to \Xi_c^0 + X$ and $\bar{B} \to \Sigma_c^0 + X$. These decays arise dominantly from $b \to (cs)\bar{c}$ and $b \to (cd)\bar{u}$, with the diquark (cs) or (cd) picking up a d quark. The diquark model would predict $\mathrm{Br}(\bar{B} \to \Xi_c^0 + X) \sim \mathrm{Br}(\bar{B} \to \Sigma_c^0 + X)$ in strong contrast to the expectation that $b \to c\bar{c}s$ is less than $b \to c\bar{u}d$ by a factor of three or more. $\mathrm{Br}\,(\bar{B} \to \Xi_c^0 X) \times \mathrm{Br}(\Xi_c^0 \to \Xi^- \pi^+)$ has been measured by CLEO [19]. While Σ_c^0 decays into $\Lambda_c \pi^-$ with 100% branching ratio, one would need absolute measurements of $\mathrm{Br}\,(\Xi_c^0 \to \Xi^- \pi^+)$ to perform the test.

For mesonic final states, the charmed rates are 40 to 50 times larger than the corresponding charmless rates. In baryonic decays, however, $b \to (cd)\bar{u}$ is 400 times larger than $b \to (ud)\bar{u}$. Part of the reason is that the charmed diquark decay constant is larger: $g_{cd} \sim 2g_{ud}$. On the other hand, while penguin operators enhance charmless mesonic decays, a similar enhancement is much weaker in the baryonic modes, as we will discuss in the next section.

We can calculate from Table 1 the total inclusive charmless baryonic decay branching ratio from the diquark picture:

$$\mathrm{Br}\,(B \to \text{charmless baryons} + X) = 8.9 \times 10^{-5}, \quad (16)$$

which is relatively small. Although this estimate is probably less reliable than (14), considering the numerous possible modes to be discussed in detail in the next section, the largest two-body decay $\bar{B}^0 \to p\bar{p}$ *is likely below* 10^{-6}, *considerably smaller than charmless mesonic decays that are typically of order* 10^{-5}.

In view of the small two-body branching ratios, it is possible that three body decays could be larger. In a calculation analogous to that of $B^0 \to D^{*-}p\bar{n}$ [10], it was estimated that $B^0 \to \rho p\bar{n}$ should be of order 10^{-5} [11], hence considerably larger than two-body modes. We note that the mechanism advocated in [11], that of current produced $p\bar{n}$ pair, is not contained in the diquark model discussed here.

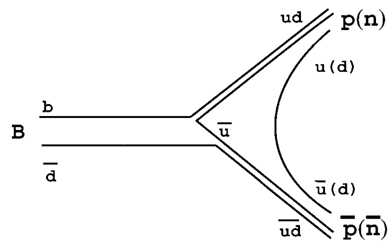

Fig. 2. $B \to p\bar{p}$, $n\bar{n}$ through $(ud)\overline{(ud)}$

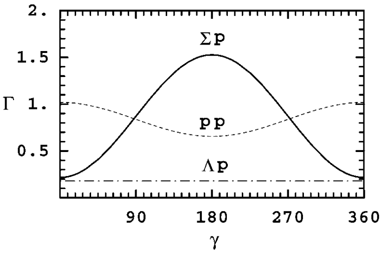

Fig. 3. The γ dependence of $B \to p\bar{p}$ (dashed), $\Sigma^+\bar{p}$ (solid) and $\Lambda\bar{p}$ (dot-dashed) rates in arbitrary units

4 Exclusive decays

As described above, two-body baryonic decays proceed via $b \to \boldsymbol{D}^{(*)}\bar{q}$ through diquark operators. The diquark $\boldsymbol{D}^{(*)}$ captures a quark from vacuum quark pair creation to form a baryon. The antiquark $\bar{q}$ pairs up with the spectator antiquark to form an antidiquark $\overline{\boldsymbol{D}}'^{(*)}$, which then captures the antiquark from pair creation and becomes the antibaryon. Admittedly, this is not a simple process compared to meson pair formation. We have seen that $\mathcal{H}_{\mathrm{diquark}}$ generates only scalar diquarks from b decay, and hence octet baryons, while $\mathcal{H}_{\mathrm{diquark}^*}$ generates only vector diquarks and hence decuplet baryons. Octet and decuplet antibaryons can result from b decay mediated by either $\mathcal{H}_{\mathrm{diquark}}$ or $\mathcal{H}_{\mathrm{diquark}^*}$.

We find that most decay channels involve only one diquark operator. We shall follow [9] which calculates the matrix element of the operators by a decomposition into four components: the diquark decay constant g_D, the $\bar{B}$ meson to antidiquark form factor $\langle \overline{\boldsymbol{D}}' | (qb) | \bar{B} \rangle$, the quark pair creation wavefunction, and the baryon wavefunction (a diquark and a quark form a baryon). The authors adopt a pole model to calculate the form factor, take a harmonic oscillator wavefunction in the ground state as the baryonic wavefunction and use a non-local wavefunction for pair creation. Since the latter consists of an undetermined normalization factor, together with other uncertainties of the four steps, the diquark model cannot be expected to predict absolute exclusive rates. But the model may give a reasonable estimate of ratios of decay rates.

It will become clear that the penguin contributions usually involve the same or similar matrix elements as the tree contributions. As a result, the matrix elements calculated in [9] can be used directly in our evaluation of the penguin effects. Since the purpose of this paper is to investigate the effect of penguin operators in the diquark model, we do not attempt to improve the calculation of the amplitudes in [9], expecting most of our conclusions to be insensitive to details. We refer readers to [9] for a discussion of the methods of calculating the various factors.

4.1 Non-strange decays $\bar{B} \to \boldsymbol{B}_1\bar{\boldsymbol{B}}_2$

Let us start with the non-strange decays $\bar{B} \to \boldsymbol{B}_1\bar{\boldsymbol{B}}_2$, such as $\bar{B}^0 \to p\bar{p}$, $n\bar{n}$. Taking $\bar{B}^0 \to p\bar{p}$ as an example, the decay occurs only through the $\bar{B}^0 \to (ud)(\bar{u}\bar{d})$ diagram, as shown in Fig. 2. The other decay $\bar{B}^0 \to n\bar{n}$ is obtained by replacing the $u\bar{u}$ pair by $d\bar{d}$. Note that the decay through $\bar{B}^0 \to (dd)(\bar{d}\bar{d})$ is impossible due to the antisymmetry of the constituent quark flavors in a scalar diquark. The decay rate can be written as

$$\Gamma(\bar{B}^0 \to p\bar{p}) = |\mathcal{A}_1|^2 \times |\langle p\bar{p}|\, (ub)_k (ud)_k^\dagger \,|\bar{B}\rangle|^2. \quad (17)$$

The constant $\mathcal{A}_1$, as defined in (9) and evaluated at the scale $\mu = m_b$, is equal to $4.1 \times 10^{-3}\mathrm{e}^{-\mathrm{i}\gamma} + 4.5 \times 10^{-4}$. The rate without penguins, as cited in [9], is

$$\begin{aligned} &\left|V_{ud}^* V_{ub}(c_2 - c_1)\right|^2 \times |\langle p\bar{p}|\, (ub)_k (ud)_k^\dagger \,|\bar{B}\rangle|^2 \\ &= \left|4.5 \times 10^{-3}\mathrm{e}^{-\mathrm{i}\gamma}\right|^2 \times |\langle p\bar{p}|\, (ub)_k (ud)_k^\dagger \,|\bar{B}\rangle|^2. \quad (18) \end{aligned}$$

Note that the two expressions differ only in the short distance coefficients and they share the same matrix element. The matrix element $\langle p\bar{p}|\, (ub)_k (ud)_k^\dagger \,|\bar{B}\rangle$ will be taken from [9]. Some interesting observations can be made even without obtaining the absolute value.

The rate now depends on the unitarity phase angle γ due to the interference of penguins with trees, analogous to the mesonic decays [20]. This dependence is shown in Fig. 3, together with the γ dependence of strange decays discussed later. The penguin contribution is roughly

Table 2. Relative rates of $\bar{B} \to \boldsymbol{B}_1\bar{\boldsymbol{B}}_2$, normalized to $\Gamma(p\bar{p})_{\text{no penguin}} = 1$ for diquark approach [9], and $\Gamma(p\bar{p}) = 1$ in sum rule approach [8]. The angle γ is taken to be 90°

	This work	[9]	[8]
$p\bar{p}$	0.84	1	1
$n\bar{n}$	0.84	1	0.3
$n\bar{p}$	0	0	0.6
$p\overline{\Delta^+}$	0.28	0.33	0.1
$n\overline{\Delta^0}$	0.28	0.33	
$p\overline{\Delta^{++}}$	0.63	0.75	0.25
$n\overline{\Delta^+}$	0.70	0.83	

one tenth of tree in amplitude, and its effect is milder compared to $B \to \pi^+\pi^-$. The reason is because, unlike $B \to \pi^+\pi^-$, which receives sizable O_6 contribution through chiral enhancement, the O_6 operator does not contribute in the diquark picture, as discussed earlier.

Other $\bar{B} \to \boldsymbol{B}_1\bar{\boldsymbol{B}}_2$ modes exhibit similar properties. Their rates are listed in Table 2, where all rates are normalized to the $p\bar{p}$ mode. Decays involving a decuplet antibaryon, $\bar{B} \to \boldsymbol{B}_1\bar{\boldsymbol{B}}_2^*$, is also possible, as mentioned in Sect. 2. An interesting channel to search for is the mode $B^- \to n\bar{p}$. In [9] it is pointed out that this decay is impossible in the diquark picture. This prediction is still true even after the penguin contribution is taken into account. None of the candidate decay diagrams such as the tree $B^- \to (ud)(\bar{u}\bar{u})$ and the penguin $B^- \to (dd)(\bar{u}\bar{d})$ survive due to the antisymmetry of the constituent quark in a scalar diquark. This is very different from the sum rule calculation [8], in which $\mathrm{Br}(B^- \to n\bar{p})$ is about as large as $\mathrm{Br}(\bar{B}^0 \to p\bar{p})$. Thus in the diquark model, one has the dynamical result of $\mathrm{Br}(\bar{B}^0 \to p\bar{p}) \sim \mathrm{Br}(\bar{B}^0 \to n\bar{n})$, but $\mathrm{Br}(B^+ \to p\bar{n})$ and $\mathrm{Br}(B^- \to n\bar{p})$ vanish, as seen from Table 2.

The tree operators $O_{1,2}$ and penguin operators $O_{3,4}$ generate scalar diquarks and produce only the $\bar{B}$ decays into an octet baryon plus an octet or decuplet antibaryon $\bar{B} \to \boldsymbol{B}\overline{\boldsymbol{B}}^{(*)}$. However, $O_{5,6}$ could generate a vector diquark and hence $\bar{B} \to \boldsymbol{B}^*\overline{\boldsymbol{B}}^{(*)}$ is possible. Decays $\bar{B}^0 \to \Delta^+\bar{p}$, $\Delta^0\bar{n}$, $\Delta^+\overline{\Delta^+}$, $\Delta^0\overline{\Delta^0}$ would arise from the vector diquark operator $(u_\mathrm{R}b_\mathrm{L})^*(u_\mathrm{R}d_\mathrm{L})^{*\dagger}$, while $\bar{B}^0 \to \Delta^-\overline{\Delta^-}$ and $B^- \to \Delta^0\overline{\Delta^+}$, $\Delta^-\overline{\Delta^0}$ from $(d_\mathrm{R}b_\mathrm{L})^*(d_\mathrm{R}d_\mathrm{L})^{*\dagger}$. The amplitudes of these decays are proportional to $\mathcal{B}_1 \sim 1.8 \times 10^{-4}$, which is about one twentieth of $\mathcal{A}_1 \sim 4.5 \times 10^{-3}$. The vector diquark decay constants are roughly the same as scalar diquark decay constants [18]:

$$g_{ud*} = 0.216\,\mathrm{GeV}^2,$$
$$g_{us*} = 0.245\,\mathrm{GeV}^2. \quad (19)$$

Assuming that the respective form factors are also of the same order as that for $\bar{B}^0 \to p\bar{p}$, we expect the branching ratio of $\bar{B} \to \boldsymbol{B}^*\overline{\boldsymbol{B}}^{(*)}$ to be about $0.0025 \times \mathrm{Br}(\bar{B}^0 \to p\bar{p})$. The small rates of $\bar{B} \to \boldsymbol{B}^*\overline{\boldsymbol{B}}^{(*)}$ is a testable feature of the diquark model. Because of further numerical uncertainties, this type of modes are not listed in Table 2.

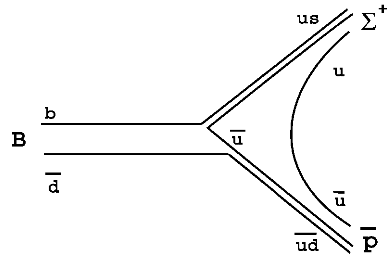

Fig. 4. $\bar{B}^0 \to \Sigma^+\bar{p}$

4.2 Strange decays $\bar{B} \to \boldsymbol{B}_s\overline{\boldsymbol{B}}$

In mesonic B decays, strange decays like $B \to K\pi$ have larger rates than non-strange decays like $B \to \pi\pi$ because of penguin contributions, with $\mathrm{Br}(B \to K^-\pi^+) \approx 1.88 \times 10^{-5}$ compared to $\mathrm{Br}(B \to \pi^-\pi^+) \approx 4.7 \times 10^{-6}$. One may wonder if the same could happen in the baryonic decays between $\bar{B} \to \boldsymbol{B}_s\bar{\boldsymbol{B}}$, such as $\bar{B}^0 \to \Lambda\bar{n}$, $\Sigma^+\bar{p}$, $\Sigma^0\bar{n}$, versus $\bar{B} \to \boldsymbol{B}_1\bar{\boldsymbol{B}}_2$ such as $\bar{B}^0 \to p\bar{p}$. This is indeed the case in the sum rule calculation, which gives $\mathrm{Br}(\bar{B}^0 \to p\bar{p}) = 1.6 \times 10^{-6}$ while $\mathrm{Br}(\bar{B}^0 \to \Sigma^+\bar{p}) = 6 \times 10^{-6}$. In the diquark calculation of [9], $\mathrm{Br}(\bar{B}^0 \to \Sigma^+\bar{p})$ is only 0.15 times $\mathrm{Br}(\bar{B}^0 \to p\bar{p})$. However, since the penguin operators are not included in [9], it is of interest to include the penguins to find the actual prediction of the diquark picture.

To include penguins, we proceed just like in the discussion of non-strange decays. Taking $\bar{B}^0 \to \Sigma^+\bar{p}$ as an example, only one diquark diagram, Fig. 4, through $\bar{B}^0 \to (us)(\bar{u}\bar{d})$ will contribute,

$$\Gamma(\bar{B}^0 \to \Sigma^+\bar{p}) = |\mathcal{A}_2|^2 \times |\langle\Sigma\bar{p}|(ub)_k(us)_k^\dagger|\bar{B}\rangle|^2, \quad (20)$$

The constant $\mathcal{A}_2$ as defined in (9), evaluated at the scale m_b, is equal to $1.0\times10^{-3}\mathrm{e}^{-\mathrm{i}\gamma}-2.2\times10^{-3}$. The rate without penguins is

$$|V_{us}^*V_{ub}(c_2-c_1)|^2 \times |\langle\Sigma\bar{p}|(ub)_k(us)_k^\dagger|B\rangle|^2$$
$$= |1.0\times10^{-3}\mathrm{e}^{-\mathrm{i}\gamma}|^2 \times |\langle\Sigma\bar{p}|(ub)_k(us)_k^\dagger|B\rangle|^2. \quad (21)$$

Again the two expressions share the same matrix element, which we take from [9] and find that

$$|\langle\Sigma\bar{p}|(ub)_k(us)_k^\dagger|B\rangle|^2 \simeq 3.0 \times |\langle p\bar{p}|(ub)_k(ud)_k^\dagger|B\rangle|^2. \quad (22)$$

The expression for $\mathcal{A}_2$ indicates that the contribution from penguins is almost twice as large as the tree contribution in amplitude and cannot be ignored. The actual branching ratio of $\bar{B}^0 \to \Sigma^+\bar{p}$ will depend on the angle γ (see Fig. 3). The penguin operators do enhance the rate significantly and $\mathrm{Br}(\bar{B}^0 \to \Sigma^+\bar{p})$ is larger than $\mathrm{Br}(\bar{B}^0 \to p\bar{p})$ for $\gamma > 90°$. However, the effects are milder than in the mesonic decays. $\mathrm{Br}(\bar{B}^0 \to \Sigma^+\bar{p})$ is at most twice $\mathrm{Br}(\bar{B}^0 \to p\bar{p})$, when $\gamma = 180°$. This is largely because the operators $O_{5,6}$ do not contribute to the same final state, unlike the role of chiral enhancement and constructive interference (between O_4 and O_6) in $B \to K\pi$. Analogous to $B \to K\pi$ and $B \to \pi\pi$, a large γ will enhance $\mathrm{Br}(\bar{B}^0 \to \Sigma^+\bar{p})$ but decrease $\mathrm{Br}(\bar{B}^0 \to p\bar{p})$. Similar results hold for the decays $\bar{B}^0 \to \Lambda\bar{n}$, $\Sigma^0\bar{n}$.

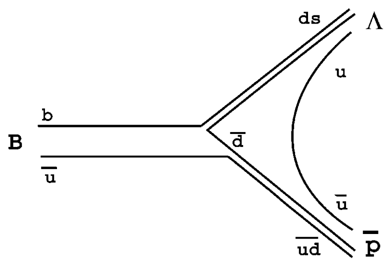

Fig. 5. $B^- \to \Lambda\bar{p}$ through $(ds)\overline{(ud)}$

Table 3. The relative rates of $\bar{B} \to \boldsymbol{B}_s\bar{\boldsymbol{B}}$ normalized as in Table 2. The angle γ is taken to be 90°

	This work	[9]	[8]
$p\bar{p}$	0.84	1	1
$\Sigma^+\bar{p}$	0.88	0.15	7.5
$\Sigma^0\bar{n}$	0.21	0.037	
$\Sigma^-\bar{n}$	0.72	0	
$\Sigma^0\bar{p}$	0.18	0	3.8
$\Lambda\bar{n}$	0.21	0.037	
$\Lambda\bar{p}$	0.18	0	< 3.8
$\Sigma^+\overline{\Delta^+}$	0.57	0.10	7.5
$\Sigma^0\overline{\Delta^0}$	0.17	0.030	
$\Sigma^+\overline{\Delta^{++}}$	1.1	0.2	7.5
$\Sigma^0\overline{\Delta^+}$	0.080	0.014	
$\Lambda\overline{\Delta^0}$	0.086	0.015	
$\Lambda\overline{\Delta^+}$	0.023	0.004	

There are pure penguin contributions that were not given in [9]. The modes $B^- \to \Lambda\bar{p}$, $\Sigma^0\bar{p}$ are two examples. Nonzero tree contribution would require the diquark decay channel $B^- \to (us)(\bar{u}\bar{u})$, which is impossible due to the antisymmetry of the constituent. However, these modes can be generated through the diquark operator $(db)_k(ds)_k^\dagger$, which arises only from penguin operators, as shown in Fig. 5. This diagram is not calculated in [9], but is identical to $\bar{B}^0 \to \Lambda\bar{n}$, $\Sigma^0\bar{n}$, respectively, after an isospin transformation $u \leftrightarrow d$. Hence the rate for $B^- \to \Lambda\bar{p}$ is given by

$$\Gamma(B \to \Lambda\bar{p}) = |\mathcal{A}_4|^2 \times |\langle\Lambda\bar{n}|(ub)_k(us)_k^\dagger|\bar{B}\rangle|^2. \quad (23)$$

The coefficient A_4 is equal to -2.2×10^{-3}. Since there is no tree–penguin interference, the rates for $B^- \to \Lambda\bar{p}$, $\Sigma^0\bar{p}$ are independent of the angle γ, just like $B \to K^0\pi^0$

The branching ratios of the modes $\bar{B} \to \boldsymbol{B}_s\bar{\boldsymbol{B}}$ are listed in Table 3. For comparison, the sum rule [8] results are also listed, which gives a strikingly different pattern. $B^- \to \Sigma^+\overline{\Delta^{++}}$ and $\bar{B}^0 \to \Sigma^+\overline{\Delta^+}$ have rates close to or larger than the one of $\bar{B}^0 \to \Sigma^+\bar{p}$. The γ dependences of their rates are identical to that of $\Sigma^+\bar{p}$.

The $O_{5,6}$ operators could generate $\bar{B} \to \boldsymbol{B}_s^*\overline{\boldsymbol{B}}^{(*)}$ via vector diquarks. The decays $\bar{B}^0 \to \Sigma^{*+}\bar{p}$, $\Sigma^{*0}\bar{n}$, $\Sigma^{*+}\overline{\Delta^+}$, $\Sigma^{*0}\overline{\Delta^0}$, $\Sigma^{*+}\overline{\Delta^{++}}$ and $B^- \to \Sigma^{*0}\bar{p}$, $\Sigma^{*+}\overline{\Delta^{++}}$ can arise from the vector diquark operator $(u_\mathrm{R}b_\mathrm{L})^*(u_\mathrm{R}s_\mathrm{L})^{*\dagger}$, while $\bar{B}^0 \to \Sigma^{*0}\overline{\Delta^0}$, $\Sigma^{*-}\overline{\Delta^-}$ and $B^- \to \Sigma^{*-}\bar{n}$, $\Sigma^{*-}\overline{\Delta^0}$, $\Sigma^{*0}\overline{\Delta^+}$ from $(d_\mathrm{R}b_\mathrm{L})^*(d_\mathrm{R}s_\mathrm{L})^{*\dagger}$. The amplitudes of these decays are proportional to $\mathcal{B}_2 \sim 9.0 \times 10^{-4}$, which is about one half of $\mathcal{A}_4$. As described above, $\mathrm{Br}(\bar{B}^0 \to \Lambda\bar{p})$ is proportional to $\mathcal{A}_4^2$. Assuming that the form factors are of the same order, we expect the rates of $\bar{B} \to \boldsymbol{B}_s^*\overline{\boldsymbol{B}}^{(*)}$ to be roughly $1/4 \times \mathrm{Br}(\bar{B}^0 \to \Lambda\bar{p}) \sim 0.03 \times \mathrm{Br}(\bar{B}^0 \to p\bar{p})$. This is still an order of magnitude smaller than the typical $\bar{B}$ decays to an octet baryon plus an octet or decuplet antibaryon. Again, since further uncertainties are involved, the $\bar{B} \to \boldsymbol{B}_s^*\overline{\boldsymbol{B}}^{(*)}$ modes are not listed in Table 3.

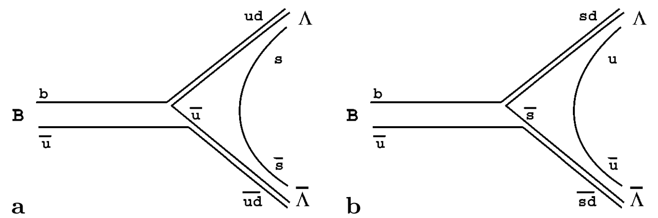

Fig. 6a,b. $\bar{B}^0 \to \Lambda\bar{\Lambda}$ through **a** $(ud)\overline{(ud)}$ and **b** $(sd)\overline{(sd)}$

Table 4. The relative rate of $B \to B_s\bar{B}_s$. Diquark rates are normalized so that $\Gamma(p\bar{p})_\text{no penguin} = 1$. Sum rule rates are normalized so that $\Gamma(p\bar{p})_\text{sum rule} = 1$. The angle γ is assumed to be 90°

	Rate	[9]	[8]
$p\bar{p}$	0.84	1	1
$\Lambda\bar{\Lambda}$	0.38	0.39	
$\Xi^0\bar{\Lambda}$	0.34	0.059	
$\Lambda\overline{\Sigma^*}^0$	0.068	0.082	
$\Xi^0\overline{\Sigma^*}^0$	0.11	0.02	
$\Lambda\overline{\Sigma^{*+}}$	0.068	0.082	
$\Xi^0\overline{\Sigma^{*+}}$	0.11	0.02	

4.3 $\bar{b} \to \boldsymbol{B}_s\bar{\boldsymbol{B}}_s$

This category includes decays like $\bar{B}^0 \to \Lambda\bar{\Lambda}$, $\Xi^0\bar{\Lambda}$, etc. The relative rates of these modes are listed in Table 4. The mode $\bar{B}^0 \to \Xi^0\bar{\Lambda}$ is more straightforward since only one diquark diagram, $\bar{B}^0 \to (us)(\bar{u}\bar{d})$, is involved. The enhancement effects from the penguin is very similar to $\bar{B}^0 \to \Sigma^+\bar{p}$.

$\bar{B}^0 \to \Lambda\bar{\Lambda}$ decay can arise from two diquark diagrams and hence is more complicated. The tree contribution is through $\bar{B}^0 \to (ud)(\bar{u}\bar{d})$ decay plus an $s\bar{s}$ pair creation (Fig. 6a). The penguin operators will enhance the Wilson coefficient in this diagram just like in the case of $\bar{B}^0 \to p\bar{p}$. However, there is one more contribution from the penguin operator: through $\bar{B}^0 \to (sd)(\bar{s}\bar{d})$ with $u\bar{u}$ pair creation (Fig. 6b). We shall argue that the penguin contribution is smaller than the tree.

The short distance coefficient for penguins, $\mathcal{A}_3 = 3.5 \times 10^{-4}$, is smaller than one tenth of the tree $V_{ud}^*V_{ub}(c_2 - c_1) = 4.5 \times 10^{-3}$. The matrix element for Fig. 6b is basically the same as in Fig. 6a, except replacing $s\bar{s}$ pair creation by $u\bar{u}$. To estimate the effect of this replacement, we can compare relative rates of $(\bar{B}^0 \to p\bar{p})$ versus $\bar{B}^0 \to \Lambda\bar{\Lambda}$ calculated in [9], which is about 2.6 : 1. Since the only difference between them at tree level is just in the pair creation, this fixes the relative weight of $s\bar{s}$ pair creation versus $u\bar{u}$. From this, we expect the penguin contribution overall to be one fifth of the tree. We therefore ignore the penguin contribution in our reporting of $\Lambda\bar{\Lambda}$ rates in Table 4. To obtain the actual numerical value and γ dependence in the future, one would have to evaluate Fig. 6b.

5 Discussion and conclusion

In this paper, we have discussed the effects of penguin operators on two-body baryonic B decays in the diquark picture. We point out that the penguin operators $O_{3,4}$ can also be transformed into operators with diquark fields and hence the calculation of their contributions is very similar to that of the tree operators. On the other hand, $O_{5,6}$ do not generate scalar diquark operators, indicating that their contribution to $\bar{B}$ decays into an octet baryon is small. As a result, penguin operators will significantly enhance $\bar{B} \to \boldsymbol{B}_s\bar{\boldsymbol{B}}$. Though the effects may not be large enough to reverse the general relative order of $\bar{B} \to \boldsymbol{B}\bar{\boldsymbol{B}}$ and $\bar{B} \to \boldsymbol{B}_s\bar{\boldsymbol{B}}$ as in the mesonic decays, some modes do have comparable rates. For example, $\bar{B}^0 \to \Sigma^+\bar{p}$, after penguin enhancement, is larger than $\bar{B}^0 \to p\bar{p}$ for $\gamma > 90°$. The $B^- \to \Sigma^+\overline{\Delta^{++}}$ mode is even of order 30% larger than $\bar{B}^0 \to \Sigma^+\bar{p}$.

The penguin operators $O_{5,6}$ could generate $\bar{B}$ decays to a decuplet baryon. These channels were predicted to vanish in [9] since the tree operators $O_{1,2}$ only generate scalar diquarks. However, the $(\mathrm{V}+\mathrm{A})\times(\mathrm{V}-\mathrm{A})$ type penguin operators $O_{5,6,7,8}$ could generate vector diquarks after a Fierz transformation, which could form decuplet baryons as the final product. Their rates are nevertheless small. We estimate the non-strange and strange decays $\mathrm{Br}(\bar{B} \to \boldsymbol{B}^*\bar{\boldsymbol{B}}^{(*)})$, $\mathrm{Br}(\bar{B} \to \boldsymbol{B}_s^*\bar{\boldsymbol{B}}^{(*)})$ are about 0.25%, 3% of $\mathrm{Br}(\bar{B}^0 \to p\bar{p})$, respectively.

The diquark model calculation of the exclusive rates depends on the pair creation model with an undetermined normalization factor. Hence, absolute rates cannot be obtained. However, as a result of duality, the inclusive rates are independent of the pair creation model. We estimate the inclusive rate for baryonic decays by computing the rate of $b \to \boldsymbol{D}\bar{q}$. This calculation relies only on the assumption of the diquark model and the values of the diquark decay constants, without a further dynamical assumption about the form factors. The total rate we get is very close to the experiment result, indicating that the diquark model is a reasonable picture for baryon production. Actually, the theoretical value is somewhat smaller, leaving some room for other mechanisms.

Since the inclusive prediction relies only on the decay constants in the diquark model, the agreement also indicates that the values of g_{cd} and g_{cs} used are reasonable. The ratio of exclusive decay rates can further check the values of g_{ud} and g_{us}. For example, the modes $\bar{B}^0 \to p\bar{p}$ and $\bar{B}^0 \to \Sigma_c^+\bar{p}$ have an identical $\bar{B}$ to antidiquark form factor, and they differ only in the diquark decay constants and CKM factors. Assuming that the transition form factor of B to the antidiquark $\overline{(ud)}$ is not very sensitive to the momentum transfer, it will cancel in the ratio of their rates. A similar argument applies for the pair creation wavefunction. The ratio can be written as

$$\frac{\Gamma(\bar{B}^0 \to p\bar{p})}{\Gamma(\bar{B}^0 \to \Sigma_c^+\bar{p})} = \left|\frac{V_{ub}}{V_{cb}}\right|^2 \times \left(\frac{g_{ud}}{g_{cd}}\right)^2, \tag{24}$$

allowing one in principle to test diquark decay constant ratios. Likewise, the ratio of $\bar{B}^0 \to \Sigma^+\overline{\Lambda_c^+}$ to $\bar{B}^0 \to \Lambda_c^+\overline{\Lambda_c^+}$ could test g_{us}/g_{cd}.

It should be clear from the previous sections that the diquark picture gives rather different predictions from the sum rule calculation. Most significantly, we note that $\bar{B} \to \boldsymbol{B}_1\bar{\boldsymbol{B}}_2$ is suppressed compared to $\bar{B} \to \boldsymbol{B}_c\bar{\boldsymbol{B}}$, $\boldsymbol{B}_c\bar{\boldsymbol{B}}_c$ and $\boldsymbol{B}_s\bar{\boldsymbol{B}}$ in a sum rule treatment [8]. The reason is that, as the sum rule authors claim, quark pair creation is mainly a soft process instead of a hard one like in the diquark non-local pair creation model. A soft process would favor heavier quarks carrying a larger momentum in the final product to pick up soft quarks from the vacuum. In the sum rule calculation, therefore, the amplitude for producing an additional quark from the vacuum is of order 1 in $\bar{B} \to \boldsymbol{B}_c\bar{\boldsymbol{B}}$ and $\boldsymbol{B}_c\bar{\boldsymbol{B}}_c$ but suppesed in $\bar{B} \to \boldsymbol{B}_1\bar{\boldsymbol{B}}_2$. Such effects are much less pronounced in the diquark model. As a result, $\bar{B} \to \boldsymbol{B}_s\bar{\boldsymbol{B}}$ typically is still smaller than $\bar{B} \to \boldsymbol{B}_1\bar{\boldsymbol{B}}_2$, even after penguins are taken into account.

Another feature of the diquark model is that several decay modes are missing due to the antisymmetry of the constituent quark flavor in a scalar diquark. For example, there is no $B^- \to n\bar{p}$ while $\bar{B}^0 \to p\bar{p}$ and $n\bar{n}$ have the same rates. The modes $B^- \to \Lambda\bar{p}$, $\Sigma^0\bar{p}$ are pure penguins and are smaller than $\bar{B}^0 \to p\bar{p}$. The sum rule approach predicts $B^- \to n\bar{p}$, $\Lambda\bar{p}$, $\Sigma^0\bar{p}$ are of the same order as $\bar{B}^0 \to p\bar{p}$. The decays arising from the penguin operator $(\bar{s}s)_{\mathrm{V}\pm\mathrm{A}}(\bar{s}b)_{\mathrm{V}-\mathrm{A}}$ such as the novel one $\bar{B}^0 \to \Omega\bar{\Xi}^-$ are supposedly possible in a sum rule calculation (though this is not mentioned in [8]). However the above penguin operator $(\bar{s}s)_{\mathrm{V}\pm\mathrm{A}}(\bar{s}b)_{\mathrm{V}-\mathrm{A}}$ does not have a scalar diquark component and thus such decays should be suppressed. The authors of [9] estimate that these decays, forbidden by the scalar diquark model, should be suppressed by at least a factor of 3. The predictions emerging from the two pictures are, anyway, different enough to be tested in the near future by experimental observation of B meson baryonic decays.

The path to observation and especially understanding charmless baryonic B decays was bound to be a long and winding one.

Acknowledgements. This work is supported in part by the National Science Council of R.O.C. under Grants NSC-89-2112-M-002-063 and NSC-90-2112-M-003-010, the BCP Topical Program of NCTS, and the MOE CosPA Project.

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