Perturbative QCD study of $B_s$ decays to a pseudoscalar meson and a tensor meson

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We study two-body hadronic $B_s \to P T$ decays, with $P(T)$ being a light pseudoscalar (tensor) meson, in the perturbative QCD approach. The CP-averaged branching ratios and the direct CP asymmetries of the $\Delta S = 0$ modes are predicted, where $\Delta S$ is the difference between the strange numbers of final and initial states. We also define and calculate experimental observables for the $\Delta S = 1$ modes under the $B_0^0 - \bar{B}_0^0$ mixing, including CP averaged branching ratios, time-integrated CP asymmetries, and the CP observables $C_1$, $D_f$ and $S_f$. Results are compared to the $B_s \to P V$ ones in the literature, and to the $B \to P T$ ones, which indicate considerable $U$-spin symmetry breaking. In our work provides theoretical predictions for the $B_s \to P T$ decays for the first time, some of which will be potentially measurable at future experiments.

Two-body hadronic $B$ meson decays have attracted a lot of attentions, because of their importance for studies of CP violation, CKM angle determination, and both weak and strong dynamics. The two $B$ factories have measured hadronic $B$ decays into light tensor ($T$) mesons recently $^1$-$^3$, which were also intensively investigated in several theoretical methods, such as the naive factorization hypothesis $^4$-$^6$, the perturbative QCD (PQCD) approach $^7$, and the QCD factorization approach $^8$. With much higher production efficiency of $B_s$ mesons at the LHCb than at the $B$ factories, many data for two-body hadronic $B_s$ decays have been published $^9$, $^{10}$, but no decays into tensor mesons were observed so far.

The $B_s$ decays into tensor mesons have not been analyzed theoretically either to our knowledge. The naive factorization hypothesis does not apply to modes involving only the annihilation amplitudes, and only the amplitudes with tensor mesons being emitted from the weak vertex. Besides, branching ratios for color-suppressed decays estimated in the naive factorization are usually too small. As for the QCD factorization $^{11}$, owing to lack of data for $B_s \to P T$ branching ratios, $P$ being a light pseudoscalar meson, the penguin-annihilation parameters cannot be determined through global fits. If the parameters associated with the $B_s \to P T$ modes were approximated by the $B_s \to P V$ ones $^8$, large theoretical uncertainties would be introduced. Both the annihilation amplitudes and the nonfactorizable tensor-emission amplitudes are calculable in the PQCD approach without inputs of free parameters. Encouraged by successful applications of the PQCD approach to many two-body hadronic $B$ meson decays $^7$, $^{12}$-$^{14}$, we will make predictions for the $B_s \to P T$ branching ratios and CP-violation observables in this letter, which can provide useful hints to relevant experiments.

The effective electroweak Hamiltonian relevant to the $B_s \to P T$ decays is written as

$$H_{eff} = \frac{G_F}{\sqrt{2}} \left[ \sum_{i=1}^{2} V_{ub}^* V_{uD} C_{i}(\mu) \mathcal{O}_{1}^{u}(\mu) - \sum_{j=3}^{10} V_{tb}^* V_{tD} C_{j}(\mu) \mathcal{O}_{j}^{u}(\mu) \right],$$

where $V$’s are the CKM matrix elements with $D$ denoting a down-type quark $d$ or $s$, $O_{i,j}(\mu)$ are the tree and penguin four-quark operators $^{15}$, and $C_{i,j}(\mu)$ are the corresponding Wilson coefficients, which evolve from the $W$ boson mass down to the renormalization scale $\mu$. In the PQCD approach a hadronic transition matrix element of a four-quark operator is further factorized into two pieces $^{16}$: the kernel with hard gluon exchanges characterized by the $b$ quark mass, and the nonperturbative hadron wave functions characterized by the QCD scale $\Lambda_{QCD}$.

The leading-order diagrams contributing to the $B_s \to P T$ decays are displayed in Fig. 1 where (a) and (b) are factorizable emission-type diagrams, (c) and (d) are nonfactorizable emission-type diagrams, (e) and (f) are factorizable annihilation-type diagrams, and (g) and (h) are nonfactorizable annihilation-type diagrams. As indicated in Fig. 1 the factorizable tensor-emission amplitudes do not exist, since a tensor meson cannot be produced via a $V$ or $A$ current. The PQCD results for the $B \to P T$ (without $B_s$) decays $^7$ are basically in agreement with the experimental data $^{17}$-$^{18}$ and those from the QCD factorization $^8$. The extension of the PQCD formalism to the $B_s \to P T$ decays is straightforward because of the similarity between $B$ and $B_s$ decays in SU(3) symmetry: the factorization formula for every diagram can be obtained by substituting the quantities in the $B_s \to P T$ decays for...
FIG. 1: Leading-order diagrams for $B_s \to PT$ decays.

the corresponding ones in the $B \to PT$ decays [7]. The confrontation of the $B \to PT$ calculations to the data has restricted the parameters involved in the $P$ and $T$ meson wave functions to some extent. In this work we will adopt the $B_s$ meson wave function in [14], and the $P$ and $T$ meson wave functions in [7].

A neutral meson and its charge conjugate partner, including the $K^0_0 - \bar{K}^0_0$, $D^0 - \bar{D}^0$, $B^0 - \bar{B}^0$, and $B_s^0 - \bar{B}_s^0$ systems, mix through the weak interaction. The $B_s^0 - \bar{B}_s^0$ mixing is the strongest, since the mass difference $\Delta M$ between the mass eigenstates is much larger than the decay width $\Gamma$ of the $B_s$ meson. The frequent oscillation between the $B_s^0$ and $\bar{B}_s^0$ mesons due to the strong mixing has rendered difficult measurements of $B_s$ decay observables at the $B$ factories, such as measurements of time-dependent CP-violation parameters. However, these measurements become feasible in LHCb experiments, because of the time dilation caused by energetic $B_s$ mesons. The mass eigenstates of the $B_s$ mesons are superpositions of the flavor eigenstates,

$$|B_{sL,H}\rangle = p|B_s^0\rangle \pm q|\bar{B}_s^0\rangle,$$

where $p$ and $q$ are complex coefficients. We neglect the difference between the mass eigenstates and the CP eigenstates, and assume that $B_{sL,H}$ is CP even (odd) as suggested in [19]. The time-dependent $B_s \to PT$ differential branching
ratios are then expressed as \cite{20}

\[
\frac{d}{dt} Br(B_s^0(t) \to f) = \Phi(B_s \to f) e^{-\Gamma t} |A_f|^2 \left[ \frac{1 + |\lambda_f|^2}{2} \right]
\times \left[ \cosh \left( \frac{\Delta \Gamma}{2} t \right) + \cos(\Delta M t) C_f - \sin(\Delta M t) S_f - \sinh \left( \frac{\Delta \Gamma}{2} t \right) D_f \right],
\]

\[
\frac{d}{dt} Br(\bar{B}_s^0(t) \to f) = \Phi(B_s \to f) e^{-\Gamma t} |\bar{A}_f|^2 \left[ \frac{1 + |\lambda_f|^2}{2} \right]
\times \left[ \cosh \left( \frac{\Delta \Gamma}{2} t \right) - \cos(\Delta M t) C_f + \sin(\Delta M t) S_f - \sinh \left( \frac{\Delta \Gamma}{2} t \right) D_f \right],
\]

\[
\frac{d}{dt} Br(B_s^0(t) \to \bar{f}) = \Phi(B_s \to f) e^{-\Gamma t} |\bar{A}_f|^2 \left[ \frac{1 + |\lambda_f|^2}{2} \right]
\times \left[ \cosh \left( \frac{\Delta \Gamma}{2} t \right) + \cos(\Delta M t) C_{\bar{f}} - \sin(\Delta M t) S_{\bar{f}} - \sinh \left( \frac{\Delta \Gamma}{2} t \right) D_{\bar{f}} \right],
\]

\[
\frac{d}{dt} Br(B_s^0(t) \to \bar{f}) = \Phi(B_s \to f) e^{-\Gamma t} |\bar{A}_f|^2 \left[ \frac{1 + |\lambda_f|^2}{2} \right]
\times \left[ \cosh \left( \frac{\Delta \Gamma}{2} t \right) - \cos(\Delta M t) C_{\bar{f}} + \sin(\Delta M t) S_{\bar{f}} - \sinh \left( \frac{\Delta \Gamma}{2} t \right) D_{\bar{f}} \right],
\]

with the mass difference \( \Delta M = (116.4 \pm 0.5) \times 10^{-10} \) MeV, the decay width difference \( \Delta \Gamma = (0.100 \pm 0.013) \times 10^{12} \) s\(^{-1}\) \cite{17}, \( \Phi(B_s \to f) \) being the phase space of the corresponding mode, and \( A_f \) being the \( B_s^0 \) \( \to f \) decay amplitude. We have employed the definitions of the amplitude ratios \( \lambda_f \) and \( \lambda_{\bar{f}} \), and the CP asymmetry observables \( C_{f,j}, D_{f,j} \), and \( S_{f,j} \) used in \cite{20}.

Since the oscillation period is much shorter than the lifetime of the \( B_s \) meson, Eq. (3) can be integrated over \( t \), and lead to the time-integrated branching ratios

\[
Br(B_s^0(\infty) \to f) = \Phi(B_s \to f) |A_f|^2 \left[ \frac{1 + |\lambda_f|^2}{2} \right] \left[ \frac{\Gamma - D_f \frac{\Delta \Gamma}{\Gamma^2}}{\Gamma^2 + \Delta M^2} + C_f \frac{\Gamma + S_f \Delta M}{\Gamma^2 + \Delta M^2} \right],
\]

\[
Br(\bar{B}_s^0(\infty) \to f) = \Phi(B_s \to f) |\bar{A}_f|^2 \left[ \frac{1 + |\lambda_f|^2}{2} \right] \left[ \frac{\Gamma - D_{\bar{f}} \frac{\Delta \Gamma}{\Gamma^2}}{\Gamma^2 + \Delta M^2} - C_{\bar{f}} \frac{\Gamma + S_{\bar{f}} \Delta M}{\Gamma^2 + \Delta M^2} \right],
\]

\[
Br(B_s^0(\infty) \to \bar{f}) = \Phi(B_s \to f) |\bar{A}_f|^2 \left[ \frac{1 + |\lambda_f|^2}{2} \right] \left[ \frac{\Gamma - D_{\bar{f}} \frac{\Delta \Gamma}{\Gamma^2}}{\Gamma^2 + \Delta M^2} - C_{\bar{f}} \frac{\Gamma + S_{\bar{f}} \Delta M}{\Gamma^2 + \Delta M^2} \right],
\]

\[
Br(\bar{B}_s^0(\infty) \to \bar{f}) = \Phi(B_s \to f) |\bar{A}_f|^2 \left[ \frac{1 + |\lambda_f|^2}{2} \right] \left[ \frac{\Gamma - D_{\bar{f}} \frac{\Delta \Gamma}{\Gamma^2}}{\Gamma^2 + \Delta M^2} + C_{\bar{f}} \frac{\Gamma + S_{\bar{f}} \Delta M}{\Gamma^2 + \Delta M^2} \right].
\]

The terms proportional to \( (\Delta \Gamma/\Gamma)^2 \approx 0.006 \) have been dropped, and the approximation \( |p/q|^2 = 1 \) has been made in the above expressions. If it happens that the \( B_{sH} \) state is CP odd while \( B_{sH} \) is CP even, the substitutions \( \Delta M \to -\Delta M \) and \( \Delta \Gamma \to -\Delta \Gamma \), or equivalently, \( D_{f,j} \to -D_{f,j} \) and \( S_{f,j} \to -S_{f,j} \) need to be done.

For \( \Delta S = 0 \) modes, a \( B_s^0 \) (\( \bar{B}_s^0 \)) meson decays to the final state \( f \) (\( \bar{f} \)), but not to \( \bar{f} \) (\( f \)) with \( f \neq \bar{f} \). In this case one can determine the initial \( B_s^0 \) or \( \bar{B}_s^0 \) meson through the final state even under the frequent \( B_s^0 \) \( - \bar{B}_s^0 \) oscillation. The ordinary definitions of CP-averaged branching ratios and direct CP asymmetries then apply directly. The predictions for the CP-averaged branching ratios and the direct CP asymmetries of these \( \Delta S = 0 \) modes are listed in Table I. The dominant topological amplitudes for each decay channel are also listed, including the color-favored (\( T \)), color-suppressed (\( C \)), and annihilation-type (\( A \)) tree amplitudes, and the corresponding penguin amplitudes PT, PC, and PA. Two types of theoretical uncertainties are estimated here: the first type comes from the variation of the nonperturbative parameters in the meson wave functions (see \cite{14,15}), except that we have adopted the recent lattice QCD result for the \( B_s \) meson decay constant, 0.228(10) GeV \cite{21}; the second type reflects the unknown next-to-leading-order QCD corrections characterized by the variations of the QCD scale \( \Lambda_{QCD} = (0.25 \pm 0.05) \) GeV and of the hard scales. It is observed that both types of uncertainties are roughly of the same order for most channels.

As shown in Table I, only the \( B_s^0 \to \pi^+ K_s^0 \) decay has a sizable branching ratio arising from the dominant amplitude \( T \), and the branching ratios of the other modes are of order \( 10^{-7} \). For color-suppressed modes such as \( B_s^0 \to \bar{K}_s^0 a_2^0, K^0 f_2 \) and \( K^0 f_2^\prime \), there is no significance difference between their branching ratios and those of their \( PV \) partners.
significant interferences in the $\Delta S_R$ implies that the above ratios are equal to each other in the U-spin symmetry limit \cite{22}. Combining our predictions with the superposition of the flavor states $\bar{\eta}$ the strange (down) quark mass. The physical U-spin conjugate processes of the other modes do not exist due to the $(\bar{\eta} S_B$ the dominant amplitudes require the $\bar{s}s$ \cite{14}, because the factorizable emission contributions are less important. For the color-favored $B^0 \rightarrow K^- a_2^\pm$ decay, whose factorizable tensor-emission amplitude is forbidden, its branching ratio $1.50 \times 10^{-7}$ is much smaller than the $B^0 \rightarrow K^- \rho^\pm$ one, $1.78 \times 10^{-5}$. Most modes in Table \[\text{II}\] exhibit large direct CP asymmetries caused by the interference between the tree and penguin amplitudes. The direct CP asymmetry in the $B^0 \rightarrow K^0 f_2^\prime$ decay would vanish if $f_2^\prime$ was a pure $\bar{s}s$ state. After receiving a tree contribution from the mixing of the isospin-1 states, this mode gets a small CP asymmetry.

To examine whether the U-spin symmetry holds in the $B_{(s)} \rightarrow PT$ decays, we define the following ratios

$$R_{CP}(B^0_s \rightarrow f) = \frac{-A_{CP}(B^0_s \rightarrow f)}{A_{CP}(B^0_s \rightarrow UF)};$$

$$R_T(B^0_s \rightarrow f) = \frac{\tau(B^0) Br(B^0_s \rightarrow f)}{\tau(B^0) Br(B^0_s \rightarrow f)};$$

where $U$ stands for the U-spin transformation, $d \leftrightarrow s$. The relation between two decay modes in a U-spin pair implies that the above ratios are equal to each other in the U-spin symmetry limit \cite{22}. Combing our predictions with the $B \rightarrow PT$ ones \cite{3}, we obtain $R_{CP}(B^0_s \rightarrow \pi^+ K^0_s^-) = 0.29^{+0.10}_{-0.08}$ and $R_T(B^0_s \rightarrow \pi^+ K^0_s^-) = 0.74^{+0.24}_{-0.19}$. $R_{CP}(B^0_s \rightarrow K^- a_2^\pm) = 1.9^{+0.5}_{-0.5}$ and $R_T(B^0_s \rightarrow K^- a_2^\pm) = 5.2^{+0.6}_{-1.0}$. The central values indicate that the U-spin symmetry is considerably broken in the $B_{(s)} \rightarrow PT$ decays by hadronic effects at order $(m_s - m_d)/\Lambda_{QCD}$ \cite{22}, $m_s$ ($m_d$) being the strange (down) quark mass. The physical U-spin conjugate processes of the other modes do not exist due to the superposition of the flavor states $\bar{q}q$ in final-state mesons.

**Table I:** Branching ratios (in units of $10^{-7}$) and direct CP asymmetries of the $\Delta S=0$ $B^0_s \rightarrow PT$ decays.

| Modes | Amplitudes | $Br$ | $A_{CP}$ (%) |
|-------|------------|------|--------------|
| $B^0_s \rightarrow \pi^+ K^- a_1^-$ | $T$ | $90^{+40}_{-12}$ | $13^{+2}_{-2}$ |
| $B^0_s \rightarrow \pi^0 K^0 a_1^-$ | $C,PA$ | $1.3^{+0.6}_{-0.5}$ | $47^{+6}_{-6}$ |
| $B^0_s \rightarrow K^- a_2^0$ | $C,PA$ | $2.0^{+0.4}_{-0.3}$ | $35^{+7}_{-7}$ |
| $B^0_s \rightarrow K^0 a_2^0$ | $C,PA$ | $3.4^{+0.7}_{-0.6}$ | $24^{+5}_{-6}$ |
| $B^0_s \rightarrow K^0 f_2$ | $PA$ | $2.0^{+0.5}_{-0.4}$ | $4.8^{+2}_{-1}$ |
| $B^0_s \rightarrow K^0 a_2^0$ | $T,PA$ | $1.5^{+0.3}_{-0.2}$ | $39^{+1}_{-1}$ |
| $B^0_s \rightarrow \eta K^0$ | $C,PA$ | $0.55^{+0.29}_{-0.19}$ | $74^{+13}_{-12}$ |
| $B^0_s \rightarrow \eta' K^0$ | $C,PT$ | $3.5^{+1.2}_{-1.0}$ | $-30^{+1}_{-1}$ |

For $\Delta S = 1$ $B^0_s$ ($B^0_s$) meson decays, we first consider those modes, whose final states are CP eigenstates, i.e. $f = \bar{f}$. In this case the four equations in Eq. \[\text{III}\] reduce to two, and one has to measure the CP observables $C_f$, $D_f$ and $S_f$ through time-dependent branching ratios, which require a lot of data accumulation. Alternatively, we define the time-integrated CP asymmetries for these decays

$$A_{CP}(B_s(\infty) \rightarrow f) \equiv \frac{Br(B^0_s(\infty) \rightarrow f) - Br(B^0_s(\infty) \rightarrow f)}{Br(B^0_s(\infty) \rightarrow f) + Br(B^0_s(\infty) \rightarrow f)} = \frac{C_f T + S_f \Delta M}{\Gamma^2 + \Delta M^2} \Gamma - D_f \Delta M;$$

and assess if there is a chance to measure it at the early stage of data accumulation.

The PQCD predictions for all the experimental observables, together with the dominant topological amplitudes and uncertainties, are shown in Table \[\text{IV}\]. It is observed that the $\eta'$-involved modes $B^0_s \rightarrow \eta a_2^0(f_2, f_2')$ have branching ratio larger than those of the corresponding $\eta$-involved modes $B^0_s \rightarrow \eta a_2^0(f_2, f_2')$. This pattern is understood, since the dominant amplitudes require the $\bar{s}s$ constituent, which is more in $\eta'$ than in $\eta$. The branching ratios of the $\Delta I = 1$ modes, like $B^0_s \rightarrow \eta a_2^0$ and $\eta' a_2^0$, are highly suppressed, compared to those of the corresponding $\Delta I = 0$ modes, $B^0_s \rightarrow \eta f_2$ and $\eta' f_2$. This suppression can be explained as follows. Neglecting the $f_2 - f_2'$ mixing effect, both $B^0_s \rightarrow \eta a_2^0$ and $\eta' f_2$ are dominated by the amplitudes $PC$ naively. However, the minus sign in the flavor constituent $(\bar{u}u - \bar{d}d)/\sqrt{2}$ renders $PC(u)$ and $PC(d)$ cancel in the former mode, while they become constructive in the latter. The source of the discrepancy between the $B^0_s \rightarrow \eta a_2^0$ and $\eta f_2$ branching ratios is the same.

Contrary to the $\Delta S = 0$ decays, the tree and penguin contributions are never simultaneously sizable to form significant interferences in the $\Delta S = 1$ decays listed in Table \[\text{III}\] so the direct CP violation $C_f$’s are tiny. One
seemingly exceptional mode is $B^0 \to \pi^0 f_2$, which has the tree and penguin contributions of the same order, but still a small direct CP asymmetry. A careful analysis reveals that the strong phases of the tree and penguin amplitudes are almost equal, $\phi^f_{\pi} \approx \phi^p_{\pi}$, and the direct CP asymmetry is proportional to $\sin(\phi^f_{\pi} - \phi^p_{\pi})$. Besides, the time-integrated CP asymmetries in Table III differ dramatically from the corresponding direct CP asymmetries $-C_f$'s. According to Eq. (6), the differences mainly come from the large mixing parameter $\Delta M$.

### TABLE II: Branching ratios (in units of $10^{-7}$) and CP observables for the $\Delta S = 1 B^0 \to PT$ decays, whose final states are CP eigenstates.

| Modes          | $B_f$ | $C_f$ | $D_f$ | $S_f$ | CP time-integrated $AC_{CP}$ (%) |
|----------------|-------|-------|-------|-------|----------------------------------|
| $\pi^0 a_2$ PA | 0.90  | 0.19  | 0.31  | 0.08  | 1.03                             |
| $\pi^0 f_2$ A,PC | 0.048 | 0.012 | 0.002 | 0.006 | 0.02                             |
| $\pi^0 f_2$ PC  | 1.2   | 0.5   | 0.1   | 0.05  | 0.1                             |
| $\eta^0_s C, A$ | 0.047 | 0.013 | 0.010 | 0.02  | 0.01                             |
| $\eta^0_s f_2$ PC | 0.2     | 0.02  | 0.005 | 0.02  | 0.04                             |
| $f'_2$ PC       | 0.004 | 0.004 | 0.004 | 0.003 | 0.01                             |
| $f'_2$ PA, PT   | 0.01  | 0.001 | 0.001 | 0.001 | 0.01                             |

There exist more complicated $\Delta S = 1$ modes, in which either a $B^0_s$ or $\bar{B}^0_s$ meson can decay into $f$ and $\bar{f}$ with $f \neq \bar{f}$. Even though a final state is identified in this case, there is no way to determine whether the initial state is a $B^0_s$ or $\bar{B}^0_s$ meson directly. It is then difficult to distinguish the four channels in Eq. (6), and time-dependent measurements are also required. For experimental access, we define the CP asymmetry parameter only by charge-tag of final states

$$AC_{CP} = \frac{Br(B^0_s/\bar{B}^0_s(\infty) \to f) - Br(B^0_s/\bar{B}^0_s(\infty) \to f)}{Br(B^0_s/\bar{B}^0_s(\infty) \to f) + Br(B^0_s/\bar{B}^0_s(\infty) \to f)}$$

(7)

All the CP observables, and the sum of the branching ratios of a pair of channels defined by

$$Br \equiv \frac{1}{2} \left[Br(B^0_s(\infty) \to f) + Br(B^0_s(\infty) \to f) + Br(B^0_s(\infty) \to f) + Br(\bar{B}^0_s(\infty) \to f)\right]$$

(8)

are presented in Table III. For the $B^0_s \to K^0 K^*_{0}$ set, all the $f$-related CP observables are equal to the $\bar{f}$-related ones, and the CP asymmetry parameter $AC_{CP}$ is exactly zero. There are no tree contributions, and the penguin amplitudes share one common weak phase in these decays. It is then straightforward to arrive at $\lambda_f = \lambda_{\bar{f}}$, and thus $C(D, S)_f = C(D, S)_{\bar{f}}$ and $AC_{CP} = 0$.

### TABLE III: Branching ratios (in units of $10^{-7}$) and CP observables for the rest $\Delta S = 1$ decays.

| Modes           | $C_f$ | $D_f$ | $S_f$ | $C_f$ | $D_f$ | $S_f$ | $Br$ | $AC_{CP}$ (%) |
|-----------------|-------|-------|-------|-------|-------|-------|-------|----------------|
| $\pi^0 a_2$     | -0.15 | 0.04  | 0.02  | -0.08 | 0.06  | 0.05  | 0.05  | 1.03    |
| $K^+ K^-_{\pi}$ | 0.95  | 0.06  | 0.02  | -0.08 | 0.06  | 0.05  | 0.05  | 1.03    |

In this letter we have investigated the $B_s \to PT$ decays in the PQCD approach, whose branching ratios and CP asymmetry parameters were predicted. It was noticed that the absence of the factorizable tensor-emission amplitudes in these decays leads to differences from the $B_s \to PV$ ones. Owing to the significant $B^0_s - \bar{B}^0_s$ mixing effect, the time-integrated CP asymmetries have been redefined and calculated for the $\Delta S = 1$ modes. The U-spin symmetry was found to be considerably broken, when the $B^0_s \to \pi^+ K^0_{-\pi}$ and $K^- a_2^+$ branching ratios are compared to the corresponding $B^0 \to K^+ a_2^+$ and $\pi^- K_{-\pi}$ ones. The branching ratios of some modes reach $O(10^{-6})$ or even $O(10^{-5})$, including $B^0 \to \eta f_2$, $\eta f_2$, $\bar{K}^+ K^-_{-\pi}$, $K^0K^0_{-\pi}$, and $\pi^+ K_{-\pi}$, which are expected to be measured at LHCb experiments. There is also potential to observe CP violation effects in the $B^0 \to \pi^+ K^0_{-\pi}$, $K^+ K^0_{-\pi}$, and $K^0K^0_{-\pi}$ decays in the near future.

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