PHENOMENOLOGICAL APPLICATIONS OF QCD FACTORIZATION TO SEMI-INCLUSIVE B DECAYS

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We have systematically investigated the semi-inclusive B decays $B \to MX$, which are manifestations of the quark decay $b \to Mq$, within a framework inspired by QCD-improved factorization. These decays are theoretically clean and have distinctive experimental signatures. We focus on a class of these that do not require any form factor information and therefore may be especially suitable for extracting information on the angles $\alpha$ and $\gamma$ of the unitarity triangle. The strong phase coming from final-state rescattering due to hard gluon exchange between the final states can induce large rate asymmetries for tree-dominated color-suppressed modes ($\pi^0, \rho^0, \omega, \phi, K^-X, K^0X, J/\psi X, \phi X, J/\psi X$) and the nonfactorizable hard spectator interactions in the 3-body decay $B \to Mq\bar{q}_2$, though phase-space suppressed, are extremely important for the tree-dominated modes ($\pi^0, \rho^0, \omega, \phi, K^-X, K^0X, J/\psi X$) and the penguin-dominated mode $\omega Xs$. Our result for $B(\bar{B} \to J/\psi Xs)$ is in agreement with experiment.

1 Why Semi-inclusive B Decays?

The semi-inclusive decays $B \to M + X$ that are of special interest originate from the quark level decay, $b \to Mq$. They are theoretically cleaner compared to exclusive decays and have distinctive experimental signatures. The theoretical advantages are: (i) A very important theoretical simplification occurs in the semi-inclusive decays over the exclusive decays if we focus on final states such that $M$ does not contain the spectator quark of the decaying $B(B_s)$ meson as then we completely by-pass the need for the transition form factor for $B(B_s) \to M$. Recall that for the exclusive case, in general, we need a knowledge of two such form factors if $M$ is a pseudoscalar meson or of four form factors if $M$ is a vector meson. (ii) There is no troublesome infrared divergent problem occurred at endpoints when working in QCD factorization, contrary to the exclusive decays where endpoint infrared divergences usually occur at twist-3 level. (iii) Some unknown strong (hard and soft) phases may arise from final-state interactions. However, these phases are mostly washed out in semi-inclusive decays. Consequently, the predictions of the branching ratios and partial rate asymmetries for $B \to MX$ are considerably clean and reliable. Since these semi-inclusive decays also tend to have appreciably larger...
branching ratios compared to their exclusive counterparts, they may therefore be better suited for extracting CKM-angles and for testing the Standard Model.

Earlier studies of semi-inclusive decays are based on generalized factorization, in which nonfactorizable effects are treated in a phenomenological way by assuming that the number of colors $N_{\text{eff}}^c$ is a free parameter to be fitted to the data or naive factorization, where $N_{\text{eff}}^c = 3$. Apart from the unknown nonfactorizable corrections, the factorization approach encounters another major theoretical uncertainty, namely the gluon’s virtuality $k^2$ in the penguin diagram is basically unknown, rendering the predictions of CP asymmetries not trustworthy.

The aforementioned difficulties with the conventional methods can be circumvented in the BBNS (Beneke, Buchalla, Neubert, Sachrajda) approach of QCD-improved factorization. Recently QCD factorization has been applied to charmless semi-inclusive decays $B \to K(K^*)X$ and $B \to \phi X_s$ in [4, 5]. Our goal is to extend the application of BBNS idea of QCD factorization to a certain class of semi-inclusive decays. In this regard our approach complements the recent works of He et al [4, 5].

2 QCD Factorization

We wish to suggest that the idea of QCD factorization can be extended to the case of semi-inclusive decays, $B \to M + X$, with rather energetic meson $M$, say $E_M \geq 2.1$ GeV. Recall that it has been shown explicitly that if the emitted meson $M_2$ is a light meson or a quarkonium in the two-body exclusive decay $B \rightarrow M_1 M_2$ with $M_1$ being a recoiled meson, the transition matrix element of an operator $O$, namely $\langle M_1 M_2|O|B\rangle$, is factorizable in the heavy quark limit. Schematically one has

$$\langle M_1 M_2|O_i|B\rangle = \langle M_1 M_2|O_i|B\rangle_{\text{fact}} \left[ 1 + \sum r_n \alpha_s^n + O(\frac{\Lambda_{\text{QCD}}}{m_b}) \right]$$

$$= \sum_j F_{ij}^{BM_1}(m_2^2) \int_0^1 du T^{ij}_i(u) \Phi_{M_2}(u)$$

$$+ \int_0^1 d\xi \, du \, dv \, T^{ij}_{ii}(\xi, u, v) \Phi_B(\xi) \Phi_{M_1}(u) \Phi_{M_2}(v), \quad (1)$$

where $F_{ij}^{BM_1}$ is a $B - M_1$ transition form factor, $\Phi_M$ is the light-cone distribution amplitude, and $T^{ij}, T^{ij}_{ii}$ are perturbatively calculable hard scattering kernels. The second hard scattering function $T^{ij}_{ii}$, which describes hard spectator interactions, survives in the heavy quark limit when both $M_1$ and $M_2$ are...
light or when $M_1$ is light and $M_2$ is a quarkonium. The factorization formula implies that naive factorization is recovered in the $m_b \to \infty$ limit and in the absence of QCD corrections. Nonfactorizable corrections are calculable since only hard interactions between the $(BM_1)$ system and $M_2$ survive in the heavy quark limit.

In order to have a reliable study of semi-inclusive decays both theoretically and experimentally, we will impose two cuts. First, a momentum cutoff imposed on the emitted light meson $M$, say $p_M > 2.1$ GeV, is necessary in order to reduce contamination from the unwanted background and ensure the relevance of the two-body quark decay $b \to Mq$. For example, an excess of $K(K^*)$ production in the high momentum region, $2.1 < p_{K(K^*)} < 2.7$ GeV, will ensure that the decay $B \to K(K^*)X$ is not contaminated by the background $b \to c$ transitions manifested as $B \to D(D^*)X \to K(K^*)X'$ and that it is dominated by the quasi two-body decay $b \to K^*q$ induced from the penguin process $b \to sg^* \to sq\bar{q}$ and the tree process $b \to u\bar{u}s$. Second, it is required that the meson $M$ does not contain the spectator quark in the initial $B$ meson and hence there us no $B-M$ transition form factors. Under these two cuts, we argue that the factorization formula (1) can be generalized to the semi-inclusive decay:

$$\langle MX|O|B\rangle = \langle MX|O|B\rangle_{\text{fact}} \left[ 1 + \sum r_n \alpha_s^n + O\left(\frac{\Lambda_{\text{QCD}}}{m_b}\right)\right]$$

$$= \int_0^1 du T^I(u) \Phi_M(u) + \int_0^1 d\xi du T^{II}(\xi,u) \Phi_B(\xi) \Phi_M(u). \quad (2)$$

However, this factorization formula is not as rigorous as the one (1) for the exclusive case, as we shall elucidate on below.

The factorizable hadronic matrix element $\langle MX|O|B\rangle$ in general consists of several terms:

$$\langle MX|O|B\rangle_{\text{fact}} = \langle M|j_1|0\rangle \langle X|j_2|B\rangle + \langle X|j'_1|0\rangle \langle M|j'_2|B\rangle$$

prohibited by 2nd cut

$$+ \langle X_1 M|j_1|0\rangle \langle X'_1|j_2|B\rangle + \langle X_2|j'_1|0\rangle \langle X'_2 M|j'_2|B\rangle$$

not favored by 1st cut, also $\alpha_s$ suppressed

$$+ \langle MX|j_1|0\rangle \langle 0|j_2|B\rangle.$$ power suppressed \quad (3)

However, several terms are prohibited or suppressed by aforementioned two cuts. The last term in Eq. (3) is the annihilation contribution and it is
suppressed by order $\Lambda_{\text{QCD}}/m_b$. In comparing Eq. (2) to the exclusive case Eq. (1), a crucial simplification that has occurred is that the semi-inclusive case does not involve any transition form factor(s). Since lack of knowledge of these form factors is often a serious limitation in quantitative applications, this adds to the appeal of the semi-inclusive case. Note also that when the emitted meson $M$ is a light meson or a quarkonium, the nonfactorizable corrections to naive factorization are infrared safe in the heavy quark limit and hence calculable.

In contrast to the exclusive case, the parton model implies that the semi-inclusive decay rate of the $B$ meson can be approximated by that of the free $b$ quark in the heavy quark limit, namely $\Gamma(B \rightarrow MX) \approx \Gamma(b \rightarrow Mq)$. Hence, the hard spectator interactions in semi-inclusive decays should be suppressed in the heavy quark limit. As we shall see later, they are suppressed by powers of $(\Lambda_{\text{QCD}}/m_b)$ at the decay rate level. However, these interactions will gain large enhancement for tree-dominated color-suppressed modes. Therefore, we will keep this term in Eq. (2).

To the order $O(\alpha_s)$, there are two additional contributions besides vertex corrections: the bremsstrahlung process $b \rightarrow Mqg$ ($g$ being a real gluon) and the process $b \rightarrow Mqq' \rightarrow Mqq'\bar{q}'$. The bremsstrahlung subprocess could potentially suffer from the infrared divergence. However, the vertex diagram in which a virtual gluon is attached to $b$ and $q$ quarks is also infrared divergent. This together with the above-mentioned bremsstrahlung process will lead to a finite and well-defined correction. Note that for exclusive hadronic decays, the infrared divergence occurring in the $BM_2$ system ($M_2$ being a recoiled meson) is absorbed into in the $B - M_2$ transition form factors. This finite correction is expected to be small as it is suppressed by a factor of $\alpha_s/\pi \approx 7\%$. Since $b \rightarrow Mqq$ does not interfere with $b \rightarrow Mq$, it can be counted as an order $O(\alpha_s)$ correction. In the presence of bremsstrahlung and the fragmentation of the quark-antiquark pair from the gluon, the factorizable configurations $\langle X_1 M | j_1 | 0 \rangle \langle X'_1 | j_2 | B \rangle$ and $\langle X_2 | j_1 | 0 \rangle \langle X'_2 M | j_2 | B \rangle$ with $X_1 + X'_1 = X$ and $X_2 + X'_2 = X$, that will break the factorization structure shown in the first term in Eq. (2), are allowed in Eq. (2). In general, one may argue that these configurations are suppressed since the momentum cut $p_{M} > 2.1$ GeV favors the two-body quark decay $b \rightarrow Mq$ and low multiplicity for $X$. However, it is not clear to us how rigorous this argument is. Therefore, in the present paper we will confine ourselves to vertex-type and penguin-type corrections as well as hard spectator interactions so that the factorization formula (2) is applicable to semi-inclusive decays at least as an approximation.

Note that QCD factorization is not applicable to the decay $B^0 \rightarrow \pi^0 D^0$ because the emitted meson $D^0$ is heavy so that it is neither small (with size of
order $1/\Lambda_{\text{QCD}}$ nor fast and cannot be decoupled from the $(B\pi)$ system. Hence, by the same token as the $\bar{B}^0 \rightarrow \pi^0 D^0$ decay, the above QCD factorization formula is also not applicable to $\bar{B}^0 \rightarrow D^0(\bar{D}^0) X$.

3 Two-body Decays of the $b$ Quark

A major advantage of studying the quasi-two-body decay of the $b$ quark is that it does not involve the unknown form factors and hence the theoretical uncertainty is considerably reduced. Hence, we first study the $CP$-averaged branching ratios and direct $CP$-violating partial rate asymmetries for some two-body hadronic $b$ decays of interest. The results are shown in Table I. Compared to the predictions of branching ratios based on naive factorization, there are three major modifications: (i) Decay modes $\pi^- u, \bar{K}^0 d$ and $K^- u$ are significantly enhanced owing to the large penguin coefficients $a_6$ and $a_4$. (ii) The modes $\pi^0 d, \rho^0 d, \omega d, J/\psi s, J/\psi d$ with neutral emitted mesons are suppressed relative to the naive factorization ones due to the smallness of $a_2$. (iii) The $\phi d$ mode has a smaller rate due to the large cancellation between $a_3$ and $a_5$. That is, while $\phi d$ is QCD-penguin dominated in naive factorization, it becomes electroweak-penguin dominated in QCD factorization.

For the prompt $\eta'$ production in semi-inclusive decays, we find the four-quark operator contributions to $b \rightarrow \eta' s$ can only account for about 10% of the measured result $B(B \rightarrow \eta' X_s) = (6.2 \pm 1.6 \pm 1.3^{+0.6}_{-1.5}(\text{bkg})) \times 10^{-4}$ for $2.0 < p_{\eta'} < 2.7$ GeV/c, where $X_s$ is the final state containing a strange quark. One important reason is that there is an anomaly effect in the matrix element $\langle \eta' | \bar{s} \gamma_5 s | 0 \rangle$ manifested by the decay constant $f_{\eta'}$. As a result, the decay rate of $b \rightarrow \eta' s$ induced by the $(S - P)(S + P)$ penguin interaction is suppressed by the QCD anomaly effect. If there were no QCD anomaly, one would have $B(b \rightarrow \eta' s) = 2.2 \times 10^{-4}$ from four-quark operator contributions which are about one third of the experimental value.

It is known that it proves to be useful to explicitly take into account the constraints from the CPT theorem when computing PRA’s for inclusive decays at the quark level (for a review, see [9]). The implication of the CPT theorem for partial rate asymmetries (PRA’s) at the hadron level in exclusive or semi-inclusive reactions is however more complicated. Consider the example $b \rightarrow du\bar{u}$. The corresponding semi-inclusive decays of the $b$ quark can be manifested as $b \rightarrow (\pi^-, \rho^-)u$ and $b \rightarrow (\pi^0, \rho^0, \omega) d$ at the two-body level and $(\pi^- \pi^0, K^0 K^-) u, (\pi^+ \pi^-, \pi^0 \pi^0, K^+ K^-) d$ at the three-body level and etc. The CPT theorem no longer constrains the absorptive cut from the $u$-loop penguin diagram to not to contribute separately to each aforementioned semi-inclusive $b$ decay, though the cancellation between $u\bar{u}$ and $c\bar{c}$ quarks will occur when all
semi-inclusive modes are summed over. In view of this observation, we shall keep all the strong phases in the calculation of direct CP violation in the individual semi-inclusive decay.

4 Semi-inclusive B Decays

Comparing to two-body decays of the b quark, there exist two more complications for semi-inclusive B decays. First, $B \rightarrow MX$ can be viewed as the two-body decay $b \rightarrow Mq$ in the heavy quark limit. For the finite b quark mass, it becomes necessary to consider the initial b quark bound state effect. Second, consider the 3-body decay $\bar{B} \rightarrow Mq_{1}\bar{q}_{2}$ with the quark content ($b\bar{q}_{2}$) for the $\bar{B}$ meson. One needs a hard gluon exchange between the spectator quark $\bar{q}_{2}$ and the meson $M$ in order to ensure that the outgoing $\bar{q}_{2}$ is hard. For the semi-inclusive case at hand, it has been argued that the hard spectator interaction is subject to a phase-space suppression since it involves three particles in the final states rather than the two-body one. However, we shall see below that it is not the case for color-suppressed decay modes, though hard spectator interactions are formally power suppressed in the heavy quark limit.

4.1 Initial bound state effect

The initial bound state effects on branching ratios and CP asymmetries have been studied recently using two different approaches: the light-cone expansion approach and the heavy quark effective theory approach. We will follow to employ the second approach. The nonperturbative HQET parameter $\mu_{i}^{2}$ is fixed from the $B^{*} - B$ mass splitting to be 0.36 GeV$^2$. Following we use $\mu_{\pi}^{2} = 0.5$ GeV$^2$, which is consistent with QCD sum rule and lattice QCD calculations. Compared to the two-body decays $b \rightarrow Mq$ shown in Table I, we see that the branching ratio of $B \rightarrow PX$ and $B \rightarrow VX$ owing to bound state effects is reduced by a factor of $(5 \sim 10)\%$ and $17\%$, respectively, while the CP asymmetry remains intact for $VX$ decays and for most of $PX$ modes.

4.2 Nonfactorizable hard spectator interactions

We now turn to the hard spectator interactions in the 3-body decay $B(p_{B}) \rightarrow M(p_{M}) + q_{1}(p_{1}) + \bar{q}_{2}(p_{2})$ with a hard gluon exchange between the spectator quark $\bar{q}_{2}$ and the meson $M$. As stressed in Sec.II, the validity of the free b quark approximation as implied by the parton model indicates that the hard spectator interaction in semi-inclusive decays is power suppressed in the heavy quark limit. Using the power counting $f_{B} \sim (\Lambda_{QCD})^{3/2}/m_{b}^{1/2}$, $f_{M} \sim \Lambda_{QCD}$, $\rho \sim \Lambda_{QCD}/m_{b}$ and taking into account the phase-space correction, it is easily
seen that the hard spectator interaction is of order $\Lambda_{\text{QCD}}/m_b$ in the heavy quark limit. However for the color-suppressed modes such as $\overline{B}_s \to \pi^0 X\bar{s}$ in which the factorizable contribution is color suppressed, the hard spectator interaction will become extremely important as it is color allowed.

It is known that the spectator interaction in $B \to K\pi$ decay, for example, is dominated by soft gluon exchange between the spectator quark and quarks that form the emitted kaon, indicating that QCD factorization breaks down at twist-3 order. This infrared divergent problem does not occur in the semi-inclusive decay, however.

4.3 Results and discussions

The results of calculations are shown in Table I. We see that the tree-dominated color-suppressed modes $(\pi^0, \rho^0, \omega) X\bar{s}$, $\phi X$, $J/\psi X\bar{s}$, $J/\psi X$ and the penguin-dominated mode $\omega X\bar{s} s$ are dominated by the hard spectator corrections. In particular, the prediction $\mathcal{B}(B \to J/\psi X\bar{s}) = 9.6 \times 10^{-3}$ is in agreement with the measurement of a direct inclusive $J/\psi$ production: $(8.0 \pm 0.8) \times 10^{-3}$ by CLEO \cite{12} and $(7.89 \pm 0.10 \pm 0.34) \times 10^{-3}$ by BaBar \cite{13}. This is because the relevant spectator interaction is color allowed, whereas the two-body semi-inclusive decays for these modes are color-suppressed. As a consequence, nonfactorizable hard spectator interactions amount to giving a large enhancement. In this work we found that it is the same spectator mechanism responsible for the enhancement observed in semi-inclusive decay $B \to J/\psi X\bar{s}$, and yet we do not encounter the same infrared problem as occurred in the exclusive case, and terms proportional to $m_B$ are not power suppressed, rendering the present prediction more reliable and trustworthy. It is conceivable that infrared divergences residing in exclusive decays will be washed out when all possible exclusive modes are summed over.

It is also interesting to notice that after including the spectator corrections, the branching ratios and PRA’s for the color-suppressed modes $\overline{B}_s^0 \to (\pi^0, \rho^0, \omega) X\bar{s}$, $B^- \to \phi X$ are numerically close to that predicted in \cite{4} based on naive factorization (see Table I). Note that the large CP asymmetries in $b \to (\pi^0, \rho^0, \omega) d$ decays (see Table I) are washed out to a large extent at hadron level by spectator interactions. By contrast, the nonfactorizable spectator interaction is in general negligible for penguin dominated (except for $\omega X\bar{s} s$) or color-allowed tree dominated decay modes. The channels $(B^-, \overline{B}^0) \to (\pi^0, \rho^0, \omega, \phi) X$ are not listed in Table II as they involve the unwanted form factors. For example, $B^- \to \pi^0 X$ contains a term $a_2 F^{B\pi}$ and $\overline{B}^0 \to \pi^0 X$ has a contribution like $a_4 F^{B\pi}$. Hence, the prediction of $(B^-, \overline{B}^0) \to \pi^0 X$ is not as
clean as $\overline{B}_s^0 \rightarrow \pi^0 X_\bar{s}$. Nevertheless, the former is also dominated by spectator interactions and is expected to have the same order of magnitude for branching ratios as the latter.

Owing to the presence of $B - \eta(\eta')$ form factors, the decays $B \rightarrow (\eta, \eta')X$ are also not listed in Table II. However, we find that the hard spectator corrections to the prompt $\eta'$ production in semi-inclusive decays are very small and hence the four-quark operator contributions to $b \rightarrow \eta' s$ can only account for about 10% of the measured result. Evidently this implies that one needs a new mechanism (but not necessarily new physics) specific to the $\eta'$. It has been advocated that the anomalous coupling of two gluons and $\eta'$ in the transitions $b \rightarrow s g^* \rightarrow \eta' g$ and $b \rightarrow s g^* g^*$ followed by $g^* g^* \rightarrow \eta'$ may explain the excess of the $\eta'$ production\textsuperscript{14,15}. An issue in this study is about the form-factor suppression in the $\eta' - g^* - g^*$ vertex and this has been studied recently in the perturbative QCD hard scattering approach\textsuperscript{16}. At the exclusive level, it is well known that the decays $B^\pm \rightarrow \eta' K^\pm$ and $B^0 \rightarrow \eta' K^0$ have abnormally large branching ratios\textsuperscript{17}. In spite of many theoretical uncertainties, it is safe to say that the four-quark operator contribution accounts for at most half of the experimental value and the new mechanism responsible for the prolific $\eta'$ production in semi-inclusive decay could also play an essential role in $B \rightarrow \eta' K$ decay.

From Table II it is clear that the semi-inclusive decay modes: $\overline{B}_s^0 \rightarrow (\pi^0, \rho^0, \omega)X_\bar{s}$, $\overline{B}_s^0 \rightarrow (K^- X, K^{*-} X)$ and $B^- \rightarrow (K^0 X_s, K^{*0} X_s)$ are the most promising ones in searching for direct CP violation; they have branching ratios of order $10^{-6} - 10^{-4}$ and CP rate asymmetries of order (10 – 40)%.

Note that a measurement of partial rate difference of $\overline{B}_s^0 \rightarrow (\pi^0, \rho^0, \omega)X_\bar{s}$ and $B^- \rightarrow (K^0 X_s, K^{*0} X_s)$ will provide useful information on the unitarity angle $\alpha$, while $\overline{B}_s^0 \rightarrow \rho^0 X_\bar{s}$ and $\overline{B}_s^0 \rightarrow (K^- X, K^{*-} X)$ on the angle $\gamma$. To have a rough estimate of the detectability of CP asymmetry, it is useful to calculate the number of $B - \overline{B}$ pairs needed to establish a signal for PRA to the level of three statistical standard deviations given by

$$N_B^{3\sigma} = \frac{9}{\Delta^2Br \epsilon_{\text{eff}}},$$

where $\Delta$ is the PRA, $Br$ is the branching ratio and $\epsilon_{\text{eff}}$ is the product of all of the efficiencies responsible for this signal. With about $1 \times 10^7$ $B\overline{B}$ pairs, the asymmetry in $K^{*0}$ channel starts to become accessible; and with about $7 \times 10^7$ $B\overline{B}$ events, the PRA’s in the other modes mentioned above will become feasible. Here we assumed, for definiteness, $\epsilon_{\text{eff}} = 1$ and a statistical significance of $3\sigma$ as in Eq. (4). Currently BaBar has collected 23 million $B\overline{B}$
It is conceivable that CP asymmetries in semi-inclusive B decays will begin to be accessible at these facilities. Likewise, PRA’s in semi-inclusive B_s decays may be measurable in the near future at the Fermilab’s Tevatron.

It is interesting to note that the decays $B^0_s \rightarrow (\pi^0, \rho^0, \omega) X_{s\bar{s}}$ and $B^- \rightarrow \phi X$ are electroweak-penguin dominated. Except for the last channel, they have sizable branching ratios and two of them have observable CP asymmetries. A measurement of these reactions will provide a good probe of electroweak penguins.

5 Conclusions

We have systematically investigated semi-inclusive B decays $B \rightarrow M X$ within a framework inspired by QCD-improved factorization. The nonfactorizable effects, such as vertex-type and penguin-type corrections to the two-body decay $b \rightarrow Mq$, and hard spectator corrections to the 3-body decay $B \rightarrow Mq_1 \bar{q}_2$ are calculable in the heavy quark limit. QCD factorization seems applicable when the emitted meson is a light meson or a charmonium.

There are two strong phases in the QCD factorization approach: one form final-state rescattering due to hard gluon exchange between $M$ and $X$, and the other from the penguin diagrams. The strong phase coming from final-state rescattering due to hard gluon exchange between the final states $M$ and $X$ can induce large rate asymmetries for tree-dominated color-suppressed modes $(\pi^0, \rho^0, \omega) X_{s\bar{s}}$. The predicted coefficient $a_2$ in QCD factorization is very small compared to naive factorization. Consequently, the color-suppressed modes $(\pi^0, \rho^0, \omega) X_{s\bar{s}}, \phi X$ and $J/\psi X, J/\psi X$ are very suppressed. Fortunately, the nonfactorizable hard spectator interactions in $B \rightarrow Mq_1 \bar{q}_2$, though phase-space suppressed, are extremely important for the aforementioned modes. Our prediction $B(B \rightarrow J/\psi X_s) = 9.6 \times 10^{-3}$ is in agreement with experiment. Contrary to the exclusive hadronic decay, the spectator quark corrections here are not subject to the infrared divergent problem, rendering the present prediction more clean and reliable.

$B^0 \rightarrow (\pi^0, \rho^0, \omega) X_{s\bar{s}}, \rho^0 X_{s\bar{s}}, B^0 \rightarrow (K^- X, K^+ X)$ and $B^- \rightarrow (K^0 X_s, K^{*0} X_s)$ are the most promising semi-inclusive decay modes in searching for direct CP violation; they have branching ratios of order $10^{-6} - 10^{-4}$ and CP rate asymmetries of order $(10 - 40\%)$. With about $7 \times 10^7 B\overline{B}$ pairs, CP asymmetries in these modes may be measurable in the near future at the BaBar, BELLE, CLEO and Tevatron experiments. The decays $B^+ \rightarrow (\pi^0, \rho^0, \omega) X_s$ and $B^- \rightarrow \phi X$ are electroweak-penguin dominated. Except for the last mode, they in general have sizable branching ratios and two of them have observable CP asymme-
tries. The above-mentioned reactions will provide good testing ground for the standard model and a good probe for electroweak penguins.

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| Mode      | BR          | PRA(%)  |
|-----------|-------------|---------|
| $b \rightarrow \pi^- u$ | $1.5 \times 10^{-4}$ ($1.3 \times 10^{-4}$) | -2 (7) |
| $b \rightarrow \rho^- u$ | $4.2 \times 10^{-4}$ ($3.5 \times 10^{-4}$) | -2 (7) |
| $b \rightarrow \pi^0 d$ | $5.3 \times 10^{-7}$ ($2.4 \times 10^{-6}$) | 93 (31) |
| $b \rightarrow \rho^0 d$ | $1.4 \times 10^{-6}$ ($5.9 \times 10^{-6}$) | 91 (33) |
| $b \rightarrow \omega d$ | $2.5 \times 10^{-6}$ ($5.8 \times 10^{-6}$) | -97 (34) |
| $b \rightarrow \phi d$ | $6.9 \times 10^{-8}$ ($2.3 \times 10^{-7}$) | -2 (0) |
| $b \rightarrow \pi^- c$ | $2.2 \times 10^{-2}$ | 0 |
| $b \rightarrow \rho^- c$ | $5.1 \times 10^{-2}$ | 0 |
| $b \rightarrow \eta d$ | $1.5 \times 10^{-6}$ | -59 |
| $b \rightarrow \eta' d$ | $1.0 \times 10^{-6}$ | 38 |
| $b \rightarrow K^0 s$ | $4.0 \times 10^{-6}$ ($2.5 \times 10^{-6}$) | -20 (4) |
| $b \rightarrow K^{+0} s$ | $2.6 \times 10^{-6}$ ($2.9 \times 10^{-6}$) | -24 (14) |
| $b \rightarrow K^- u$ | $9.2 \times 10^{-5}$ ($2.9 \times 10^{-5}$) | 5 (28) |
| $b \rightarrow K^{+u} u$ | $4.8 \times 10^{-5}$ ($5.1 \times 10^{-5}$) | 17 (44) |
| $b \rightarrow \bar{K}^0 d$ | $1.0 \times 10^{-4}$ ($2.0 \times 10^{-5}$) | 0.8 (1) |
| $b \rightarrow \bar{K}^{+0} d$ | $6.6 \times 10^{-5}$ ($2.6 \times 10^{-5}$) | 0.9 (3) |
| $b \rightarrow K^- c$ | $1.7 \times 10^{-3}$ | 0 |
| $b \rightarrow K^{+c} c$ | $2.7 \times 10^{-3}$ | 0 |
| $b \rightarrow \eta s$ | $1.9 \times 10^{-5}$ | -4 |
| $b \rightarrow \eta' s$ | $5.4 \times 10^{-5}$ | 1 |
| $b \rightarrow \pi^0 s$ | $1.8 \times 10^{-6}$ ($1.6 \times 10^{-6}$) | 19 (0) |
| $b \rightarrow \rho^0 s$ | $5.1 \times 10^{-6}$ ($4.3 \times 10^{-6}$) | 19 (0) |
| $b \rightarrow \omega s$ | $3.3 \times 10^{-7}$ ($1.3 \times 10^{-6}$) | 61 (0) |
| $b \rightarrow \phi s$ | $5.5 \times 10^{-5}$ ($6.3 \times 10^{-5}$) | 1 (0) |
| $b \rightarrow J/\psi s$ | $5.4 \times 10^{-4}$ | -0.5 |
| $b \rightarrow J/\psi d$ | $2.8 \times 10^{-5}$ | 10 |

Table 1: CP-averaged branching ratios and partial-rate asymmetries for some two-body hadronic $b$ decays. For comparison, the predicted branching ratios and rate asymmetries (in absolute values for $\gamma = 60^\circ$) based on naive factorization \cite{1} are given in parentheses.
Table 2: CP-averaged branching ratios and partial-rate asymmetries for some semi-inclusive hadronic B decays with $E_M > 2.1$ GeV for light mesons and $E_{J/\psi} > 3.3$ GeV for the $J/\psi$. Branching ratios due to hard spectator interactions in the 3-body decay $B \to M q \bar{q}_2$ are shown in parentheses. Here $X$ denotes a final state containing no (net) strange or charm particle, and $X_q$ the state containing the quark flavor $q$.

| Mode | BR     | PRA(%) |
|------|--------|--------|
| $B^0 \to \pi^- X$ ($\bar{B}^0 \to \pi^- X_s$) | $1.3 \times 10^{-4}$ ($5.1 \times 10^{-8}$) | -2 |
| $B^0 \to \rho^- X$ ($\bar{B}^0 \to \rho^- X_s$) | $3.4 \times 10^{-4}$ ($2.2 \times 10^{-7}$) | -2 |
| $\bar{B}^0_s \to \pi^0 X_s$ | $1.3 \times 10^{-6}$ ($8.7 \times 10^{-7}$) | 31 |
| $\bar{B}^0_s \to \rho^0 X_s$ | $4.8 \times 10^{-6}$ ($3.7 \times 10^{-6}$) | 22 |
| $\bar{B}^0_s \to \omega X_s$ | $5.5 \times 10^{-6}$ ($3.4 \times 10^{-6}$) | -37 |
| $B^- \to \phi X$ | $2.5 \times 10^{-7}$ ($1.9 \times 10^{-7}$) | -0.5 |
| $\bar{B}^0 \to \pi^- X_s$ ($\bar{B}^0_s \to \pi^- X_{cs}$) | $1.8 \times 10^{-2}$ ($8.4 \times 10^{-6}$) | 0 |
| $\bar{B}^0 \to \rho^- X_s$ ($\bar{B}^0_s \to \rho^- X_{cs}$) | $4.2 \times 10^{-2}$ ($1.1 \times 10^{-4}$) | 0 |
| $B^- \to K^0 X_s$ | $3.8 \times 10^{-6}$ ($2.9 \times 10^{-9}$) | -20 |
| $B^- \to K^{*0} X_s$ | $2.2 \times 10^{-6}$ ($1.1 \times 10^{-8}$) | -24 |
| $\bar{B}^0 \to K^- X$ ($\bar{B}^0_s \to K^- X_s$) | $8.7 \times 10^{-5}$ ($3.6 \times 10^{-9}$) | 5 |
| $\bar{B}^0 \to K^{*+} X$ ($\bar{B}^0_s \to K^{*+} X_s$) | $3.9 \times 10^{-5}$ ($1.4 \times 10^{-8}$) | 16 |
| $B^- \to K^{0*} X$ | $9.7 \times 10^{-5}$ ($7.5 \times 10^{-8}$) | 0.8 |
| $B^- \to \bar{K}^{0*} X$ | $5.4 \times 10^{-5}$ ($2.9 \times 10^{-7}$) | 0.9 |
| $\bar{B}^0 \to K^- X_c$ ($\bar{B}^0_s \to K^- X_{cs}$) | $1.4 \times 10^{-3}$ ($4.3 \times 10^{-7}$) | 0 |
| $\bar{B}^0 \to K^{*+} X_c$ ($\bar{B}^0_s \to K^{*+} X_{cs}$) | $2.3 \times 10^{-3}$ ($4.8 \times 10^{-6}$) | 0 |
| $\bar{B}^0 \to \pi^0 X_{ss}$ | $1.5 \times 10^{-6}$ ($5.0 \times 10^{-8}$) | 19 |
| $\bar{B}^0_s \to \rho^0 X_{ss}$ | $4.4 \times 10^{-6}$ ($2.2 \times 10^{-7}$) | 18 |
| $\bar{B}^0_s \to \omega X_{ss}$ | $7.4 \times 10^{-6}$ ($7.1 \times 10^{-6}$) | 2 |
| $B^- \to \phi X_s$ | $5.8 \times 10^{-5}$ ($2.8 \times 10^{-6}$) | 1 |
| $B \to J/\psi X_s$ | $9.6 \times 10^{-3}$ ($9.2 \times 10^{-3}$) | 0 |
| $B \to J/\psi X$ | $5.1 \times 10^{-4}$ ($4.9 \times 10^{-4}$) | 0.5 |