SLAC–PUB–8961
UK/TP 01–05
August 2001

Evading the CKM Hierarchy*

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Abstract

We show that the presence of intrinsic charm in the hadrons' light-cone wave functions, even at a few percent level, provides new, competitive decay mechanisms for B decays which are nominally CKM-suppressed. For example, the weak decays of the B-meson to two-body exclusive states consisting of strange plus light hadrons, such as $B \to \pi K$, are expected to be dominated by penguin contributions since the tree-level $b \to su\overline{u}$ decay is CKM suppressed. However, higher Fock states in the B wave function containing charm quark pairs can mediate the decay via a CKM-favored $b \to sc\overline{c}$ tree-level transition. Such intrinsic charm contributions can be phenomenologically significant. Since they mimic the amplitude structure of "charming" penguin contributions, charming penguins need not be penguins at all.

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*Research partially supported by the Department of Energy under contracts DE-AC03-76SF00515 and DE-FG02-96ER40989.

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I. INTRODUCTION

It is usually assumed in the analysis of B-meson decays that only the valence quarks of the initial and final-state hadrons participate in the weak interaction. Typical examples are the semi-leptonic decay $B^- \to \ell^- \bar{\nu} \pi^+$, which is based on the transition $b \to u \ell^- \bar{\nu}$; $B^- \to K^- \gamma$, which is based on the penguin amplitude $b \to s\gamma$; and $B^- \to K^- \pi^0$, which is based on $b \to s u \bar{u}$ and penguin $b \to s g^* \to s u \bar{u}$ transitions. In each case, it is assumed that the matrix elements of the operators of the effective weak Hamiltonian involve only the valence quarks of the incoming and outgoing hadrons. Any non-valence gluon or sea quarks present in the initial or final state wave functions appear only as spectators.

The wave functions of a bound state in a relativistic quantum field theory such as QCD necessarily contain Fock states of arbitrarily high particle number. For example, the B^--meson has a Fock state decomposition

$$\left| B^- \right\rangle = \psi_{b\bar{u}} \, | b\bar{u} \rangle + \psi_{b\bar{u}g} \, | b\bar{u}g \rangle + \psi_{b\bar{u}d\bar{d}} \left| b\bar{u}d\bar{d} \right\rangle + \psi_{b\bar{u}s\bar{s}} \, | b\bar{u}s\bar{s} \rangle + \psi_{b\bar{u}c\bar{c}} \, | b\bar{u}c\bar{c} \rangle + \cdots . \quad (1.1)$$

The Fock state decomposition is most conveniently done at equal light-cone time $\tau = t + z/c$ using light-cone quantization in the light-cone gauge $A^+ = 0$ [1,2]. The light-cone wave function $\psi_n(x_i, \vec{k}_{\perp i}, \lambda_i)$ depends on the momentum fraction of parton x_i, where $x_i = k_i^+/P^+$ and $\sum_i x_i = 1$, the transverse momentum $\vec{k}_{\perp i}$ where $\sum_i \vec{k}_{\perp i} = 0$, and the helicity λ_i. The light-cone wave functions are Lorentz invariant; i.e., they are independent of the total momentum $P^+ = P^0 + P^z$ and of $P_\perp$ of the bound state. The extra gluons and quark pairs in the higher Fock states arise from the QCD interactions. Contributions which are due to a single gluon splitting such as $g \to c\bar{c}$ are associated with DGLAP evolution, or they provide perturbative loop corrections to the operators; they are extrinsic to the bound-state nature of the hadron. In contrast, the $c\bar{c}$ pairs which are multiply-connected to the valence quarks cannot be attributed to the gluon substructure and are intrinsic to the hadron's structure. The intrinsic, heavy quarks are thus part of the non-perturbative bound state structure of the hadrons themselves [3], rather than part of the short-distance operators associated with the DGLAP evolution of structure functions or radiative corrections to the effective weak Hamiltonian.

Recently Franz, Polyakov, and Goeke have analyzed the properties of the intrinsic heavy-quark fluctuations in hadrons using the operator-product expansion [4]. For example, the light-cone momentum fraction carried by intrinsic heavy quarks in the proton $x_{Q\bar{Q}}$ as measured by the T^{++} component of the energy-momentum tensor is related in the heavy-quark limit to the forward matrix element $\langle p | \mathrm{tr}_c (G^{+\alpha} G^{+\beta} G_{\alpha\beta})/m_Q^2 | p \rangle$, where $G^{\mu\nu}$ is the gauge field strength tensor. Diagrammatically, this can be described as a heavy quark loop in the proton self-energy with four gluons attached to the light, valence quarks [5,6]. Since the non-Abelian commutator $[A_\alpha, A_\beta]$ is involved, the heavy quark pairs in the proton wavefunction are necessarily in a color-octet state. It follows from dimensional analysis that the momentum fraction carried by the $Q\bar{Q}$ pair scales as $k_\perp^2/m_Q^2$ where $k_\perp$ is the typical momentum in the hadron wave function. In contrast, in the case of Abelian theories, the contribution of an intrinsic, heavy lepton pair to the bound state's structure first appears in $\mathcal{O}(1/m_L^4)$. One relevant operator corresponds to the Born-Infeld $(F_{\mu\nu})^4$ light-by-light scattering insertion, and the momentum fraction of heavy leptons in an atom scales as $k_\perp^4/m_L^4$.

In the case of the proton, analyses [7–9] of the charm structure function measured by the EMC group indicates a significant charm quark excess beyond DGLAP or gluon-splitting predictions at large $x_{Bj} \sim 0.4$, and suggest that the intrinsic charm (IC) probability is $\lesssim 1\%$. Although these analyses are not conclusive [10], this value is consistent with the theoretical estimate of Franz et al. [4,11] An intrinsic charm component in the light hadrons of this scale has been invoked to explain the "$\rho\pi$" puzzle in J/ψ decay [12], leading charm production in $\pi\, N$ collisions [13–15], as well as the production of pairs of J/ψ at large x_F in these reactions [16].

The existence of intrinsic charm (IC) in the proton also implies the existence of IC in other hadrons, including the B-meson. In order to translate the estimate of IC probability in the proton to the IC of a B-meson, we are faced with two conflicting effects. The typical internal transverse momentum $k_\perp$ is larger in the B-meson, evidently favoring a larger IC probability in the B meson; on the other hand the proton's additional valence quark generates a larger combinatoric number of IC diagrams, favoring a larger IC probability in the proton. In evaluating the first effect, it is useful to compare positronium with H-atom: the kinematics of the heavy-light system make its ground-state radius a factor of two smaller than that of positronium, and thus its typical bound state momentum is a factor of two larger. This analogy should be applicable when comparing the internal scales of the B-meson to that of the light pseudoscalars; we note that the normalization of the light-cone wave function $\phi_i(x, \vec{k_\perp})$ with $i \in \pi, B$ is set by the decay constant f_i. Lattice calculations indicate $f_B \sim 191$ MeV [17], so that $f_B/f_\pi \sim 2$, suggesting that the momentum $k_\perp$ is significantly higher in the B meson than in light hadrons. Thus the IC component in the B-meson could be as large as four times that of the proton, that is, $\sim 4\%$. The IC component of the Λ_b baryon could be larger; in this case, the additional valence quark generates a larger combinatoric number of IC diagrams as well. The ways in which the decaying b quark interacts with its hadronic environment, particularly "spectator effects," are evidently important in explaining the lifetime difference in the B and Λ_b [18,19] hadrons; IC could play a role in this context as well.

The presence of intrinsic charm quarks in the B wave function provides new mechanisms for B decays. For example, Chang and Hou have considered the production of final states with three charmed quarks such as $\bar{B} \to J/\psi D \pi$ and $\bar{B} \to J/\psi D^*$ [20]; these final states are difficult to realize in the valence model, yet they occur naturally when the b quark of the intrinsic charm Fock state $|\, b\bar{u}c\bar{c}\rangle$ decays via $b \to c\bar{u}d$. In fact, the J/ψ spectrum for inclusive $B \to J/\psi X$ decays measured by CLEO and Belle shows a distinct enhancement at the low J/ψ momentum where such decays would kinematically occur [21,22]. Alternatively, this excess could reflect the opening of baryonic channels such as $B^- \to J/\psi \bar{p}\Lambda$ [23].

These ideas take on particular significance in view of the hierarchical structure of the CKM matrix — the weak transition $b \to sc\overline{c}$ is doubly Cabibbo enhanced with respect to a $b \to su\overline{u}$ transition. For example, intrinsic charm components in the initial and final hadron light-cone wave functions will allow $|\Delta S| = 1$ B-meson decays through processes such as that shown in Fig. 1; the small intrinsic charm probability is offset by the comparatively large CKM matrix elements associated with the $b \to sc\overline{c}$ transition, promoting their phenomenological impact.

As a specific illustration, consider the exclusive $|\Delta S| = 1$ decays, $B \to \pi K$ and $B \to \rho K$. The various πK final states from $\bar{B}^0$ and B^- decay are connected by isospin symmetry; the

same is true of the branching ratios to ρK. The presence of weak transitions involving intrinsic charm can alter the pattern of the predicted branching ratios. Since the same initial and final states are involved, the intrinsic charm contribution can interfere with the conventional amplitudes, and yield significant effects. We note that such intrinsic contributions function in a manner identical to the "charming penguin" contributions discussed by Ciuchini et al. [24–26]; thus charming penguins need not be penguins at all.

Halperin and Zhitnitsky have considered the role of IC in mediating the decays $B \to \eta' K$ [27], and $B \to \eta' X$ [28], arguing, as we have, that IC can be important when coupled with the Cabibbo-enhanced $b \to sc\bar{c}$ transition in decays to charmless final states [29]. They effect their numerical estimates in the factorization approximation, so that the importance of their IC mechanism is determined by the parameter $f_{\eta'}^{(c)}$, where

$$\langle 0|\bar{c}\gamma_\mu\gamma_5 c|\eta'(p)\rangle = i f_{\eta'}^{(c)} p_\mu \ . \tag{1.2}$$

Recent work has shown $f_{\eta'}^{(c)}$ to be ≈ -2 MeV [4], rather smaller [30,31,29,32] than $f_{\eta'}^{(c)} \approx 50 - 180$ MeV [27,28], so that efforts to reconcile the observed rate with Standard Model (SM) predictions continue [33]. Although other mechanisms could well be at work [34,35], we wish to point out that the factorization approximation does not capture the physics of IC. IC is produced in a higher Fock component of a hadron's light-cone wave function; it is naturally in a color octet state [30], so that the dynamical role it plays in mediating B-meson decay is intrinsically non-factorizable in nature.

Although we will specifically consider the role of IC in exclusive B-meson decays in this paper, the effect of IC can have a more general phenomenological impact on B physics. For example, it is well-known that the semileptonic branching fraction in inclusive B-meson decay, $B_{\rm sl}$, is smaller than SM predictions; however, the "natural" resolution of this puzzle – an increased $b \to sc\bar{c}$ rate – is at odds with the observed number of charm (and anti-charm) quarks per B-meson decay, n_c, as this is also too small compared to SM expectations [36]. IC in the B-meson can increase the charmless $b \to s$ rate, as illustrated in Fig. 1, thus reducing the semileptonic branching ratio. Earlier work ascribed a possible role to IC in resolving the $B_{\rm sl}/n_c$ puzzle [37], yet only IC in the light hadrons was considered. We believe that the role played by IC in the B meson in realizing strange, charmless final states to be potentially of greater importance. It should be recognized that IC does not significantly increase the inclusive yield of charmed hadrons, since the materialization of intrinsic charm is dynamically suppressed. For example, in hadron collisions the probability of materializing IC Fock states is of $\mathcal{O}(1/m_c^4)$ [5], save for an exceptional portion of phase space, at large x_F [13], for which it is of $\mathcal{O}(1/m_c^2)$ [13]. As noted by Chang and Hou, [20] the materialization of IC of the B-meson will lead to novel exclusive decays to charm final states; for example, IC can mediate $B^- \to J/\psi e^- \bar{\nu}_e$, as well as $B \to J/\psi\gamma$. If there were no IC in the B-meson, such final states could only be realized through OZI-violating processes.

II. IC IN $B \to \pi K$ AND $B \to \rho K$ DECAY

We will now consider the specific role of IC in mediating the exclusive decays $B \to \pi K$ and $B \to \rho K$. The operator-product expansion and renormalization group methods provide a systematic theoretical framework for analyzing exclusive hadronic B-meson decays. The

amplitude for the decay of a B meson to a hadronic final state f is given by $A(B \to f) = \langle f \mid \mathcal{H}_{\rm eff} \mid B\rangle$, and for $b \to sq\bar{q}$ decay $\mathcal{H}_{\rm eff}$ can be written as

$$\mathcal{H}_{\rm eff} = \frac{G_F}{\sqrt{2}} \left[\sum_{p=u,c} V_{pb}V_{ps}^*(C_1 O_1^s + C_2 O_2^s) - V_{tb}V_{ts}^* \sum_j C_j O_j \right] , \tag{2.1}$$

where the explicit form of the operators are given in Ref. [38]. The $O_{1,2}^p$ are the left-handed current-current operators arising from W-boson exchange; the sum over j contains the strong and electroweak penguin operators. The Wilson coefficients are computed at $\mu = M_W$ and are evolved, using renormalization group methods, to $\mu = \mathcal{O}(m_b)$; only C_2 is of $\mathcal{O}(1)$ at the scale $\mu \sim M_W$. The $C_i(\mu)$ are known; the operator matrix elements $\langle f \mid O_i(\mu) \mid B\rangle$, however, pose a continuing theoretical challenge. Various methods exist to estimate them [39–44]; however, all of these approaches assume that only the valence degrees of freedom participate in the decay process. Indeed, this viewpoint is shared by attempts to catalog all the possible contributions to the various exclusive decay amplitudes [45]. It is this assumption that we challenge.

It should be emphasized that Fock states of arbitrary particle number are necessary to describe a relativistic bound state. Furthermore, as shown in Ref. [46], the matrix element associated with a time-like form factor entering semi-leptonic decay has two distinct contributions: a contribution in which parton number is conserved, so that n partons $\to n$ partons, and one in which it is not. In the latter case, one of the anti-quarks in a non-valence Fock state fluctuation in the B-meson annihilates the b quark, yielding the transition $n+2$ partons $\to n$ partons. These conclusions emerge from Lorentz invariance in concert with the kinematic constraint of $q^+ > 0$. The omission of such contributions can become acute in the context of $b \to sq\bar{q}$ decays, since CKM factors favor $b \to st\bar{t}$ and $b \to sc\bar{c}$ transitions over those of $b \to su\bar{u}$ by a numerical factor of roughly 50. We shall show that when such decays are mediated by IC contributions, the CKM hierarchy can be evaded, impacting not only the branching ratios, but also the CP asymmetries associated with these decays.

Consider the following family of decay modes:

$$\mathcal{B}(\bar{B}^0 \to \pi^+ K^-) \quad \mathcal{B}(\bar{B}^- \to \pi^0 K^-) \quad \mathcal{B}(B^- \to \pi^- \bar{K}^0) \quad \mathcal{B}(\bar{B}^0 \to \pi^0 \bar{K}^0) \, , \tag{2.2}$$

as well as

$$\mathcal{B}(\bar{B}^0 \to \rho^+ K^-) \quad \mathcal{B}(\bar{B}^- \to \rho^0 K^-) \quad \mathcal{B}(B^- \to \rho^- \bar{K}^0) \quad \mathcal{B}(\bar{B}^0 \to \rho^0 \bar{K}^0) \, . \tag{2.3}$$

These decays are mediated by the short-distance weak transition $b \to sq\bar{q}$ where $q \in u, c, t$. The magnitude of the branching ratios themselves as well as the precise patterns of ratios of ratios are sensitive to the contributing decay topologies, the CKM factors accompanying a particular operator, and the numerical values of the accompanying Wilson coefficients. Intrinsic charm can play a particularly important role here since the operator with the largest Wilson coefficient combined with the largest combination of CKM matrix elements can now enter (the $p = c$ term in Eq. (2.1)) at tree level.

Let us begin by a constructing a simple numerical estimate for the above ratios. The transitions $b \to st\bar{t}$ and $b \to sc\bar{c}$ are favored over the $b \to su\bar{u}$ transition by a factor of λ^2;

however, the values of the CKM parameters can exacerbate the flavor hierarchy. That is, using $\sqrt{\rho^2+\eta^2}=0.38$ and $\lambda=0.2196$, which fall within the $\geq 5\%$ C.L. of recent fits [47], the $b\to su\bar{u}$ transition is suppressed with respect to $b\to st\bar{t}$ and $b\to sc\bar{c}$ transitions by a factor of roughly 50. If we neglect isospin-violating effects, in particular $|\Delta I|=1$ electroweak penguins, and assume that $b\to su\bar{u}$ transitions are negligible, just one decay topology exists for each decay in a particular "family," and the various decay rates are related by isospin symmetry. This simple model predicts the following pattern of ratios of branching ratios:

$$\mathcal{B}(\bar{B}^0\to\pi^+K^-):\mathcal{B}(\bar{B}^-\to\pi^0K^-):\mathcal{B}(B^-\to\pi^-\bar{K}^0):\mathcal{B}(\bar{B}^0\to\pi^0\bar{K}^0)=2:1:2:1\ , \quad (2.4)$$

as well as

$$\mathcal{B}(\bar{B}^0\to\rho^+K^-):\mathcal{B}(\bar{B}^-\to\rho^0K^-):\mathcal{B}(B^-\to\rho^-\bar{K}^0):\mathcal{B}(\bar{B}^0\to\rho^0\bar{K}^0)=2:1:2:1\ . \quad (2.5)$$

[The factors of 2 arise as the contributing components of the $\rho^\pm$, $\pi^\pm$ and ρ^0, π^0 wave functions differ by a normalization factor of $\sqrt{2}$.] The corresponding predictions for the branching ratios of the CP-conjugate decays are identical, and thus the associated CP asymmetries are zero in this limit. Measurements by the CLEO collaboration [48] of the CP-averaged modes are consistent with this prediction:

$$\begin{aligned}\mathcal{B}(\bar{B}^0\to\rho^\mp K^\pm)&=(16.0\ {}^{+7.6}_{-6.4}\pm2.8)\cdot10^{-6}\\ \mathcal{B}(B^-\to\rho^0K^-)&=(8.4\ {}^{+4.0}_{-3.4}\pm1.8)\cdot10^{-6}\ .\end{aligned} \quad (2.6)$$

The more extensive data for the $B\to\pi K$ modes [49–51] are summarized in Table I. In this case the simple 2:1:2:1 pattern is only roughly realized. No significant CP asymmetries have been observed thus far.

Before proceeding to detailed estimates, we discuss the contributions to $B\to\pi K$ and $B\to\rho K$ decays in terms of a parametrization based on the Wick contractions in the matrix elements of the operators of the effective Hamiltonian [45]. The individual parameters that appear are manifestly scale and scheme-independent, so that the deficiencies of analyses based on the factorization approximation [52] are avoided. We first assume, as in Ref. [45], that only the valence degrees of freedom of the mesons participate in the decay. The various $B\to\pi K$ amplitudes are then parameterized as

$$\begin{aligned}&\mathcal{A}(B^0\to K^+\pi^-)=V_{us}V_{ub}^*(E_1(s,u,u;B^0,K^+,\pi^-)-P_1^{\rm GIM}(s,u;B^0,K^+,\pi^-))\\ &-V_{ts}V_{tb}^*P_1(s,u;B^0,K^+,\pi^-)\end{aligned} \quad (2.7)$$

$$\begin{aligned}&\mathcal{A}(B^+\to K^+\pi^0)=\frac{V_{us}V_{ub}^*}{\sqrt{2}}(E_1(s,u,u;B^+,K^+,\pi^0)+E_2(u,u,s;B^+,\pi^0,K^+)\\ &-P_1^{\rm GIM}(s,u;B^+,K^+,\pi^0)+A_1(s,u,u;B^+,K^+,\pi^0))\\ &-\frac{V_{ts}V_{tb}^*}{\sqrt{2}}P_1(s,u;B^+,K^+,\pi^0)+\Delta A(B^+\to K^+\pi^0)\end{aligned} \quad (2.8)$$

$$\begin{aligned}&\mathcal{A}(B^+\to K^0\pi^+)=V_{us}V_{ub}^*(A_1(s,d,u;B^+,K^0,\pi^+)-P_1^{\rm GIM}(s,d;B^+,K^0,\pi^+))\\ &-V_{ts}V_{tb}^*P_1(s,d;B^+,K^0,\pi^+)\end{aligned} \quad (2.9)$$

$$\begin{aligned}&A(B^0\to K^0\pi^0)=\frac{V_{us}V_{ub}^*}{\sqrt{2}}(E_2(u,u,s;B^0,\pi^0,K^0)+P_1^{\rm GIM}(s,d;B^0,K^0,\pi^0)\\ &-\frac{V_{ts}V_{tb}^*}{\sqrt{2}}P_1(s,d;B^0,K^0,\pi^0)+\Delta A(B^0\to K^0\pi^0)\ .\end{aligned} \quad (2.10)$$

We have used the notation of Ref. [45] and the convention $\rho^-, \pi^- = \bar{u}d$ and $\rho^0, \pi^0 = (\bar{u}u - \bar{d}d)/\sqrt{2}$. The corresponding parametrization of the $B \to \rho K$ amplitudes is obtained from the replacement $(\pi^-, \pi^0, \pi^+) \to (\rho^-, \rho^0, \rho^+)$. The label "$E_i$" refers to operators with W^- emission topologies; "A_i" refers to annihilation topologies, whereas "P_i" contains penguin topologies. The term "$P_1^{\rm GIM}$" represents penguin contributions which vanish in the $m_c = m_u$ limit. The contribution labelled "ΔA" vanishes in the limit of isospin symmetry; electroweak penguin effects contribute to it, as do isospin-violating contributions in the matrix elements themselves [53].

In the effective theory, the effects of the heavy degrees of freedom, such as the $W^\pm, Z$ or t quark are replaced by effective coupling constants, the Wilson coefficients $C_i(\mu)$, multiplying effective vertices $O_i(\mu)$. Combinations of the products of $C_i(\mu)$ and $O_i(\mu)$ are individually scale and scheme invariant. Nevertheless, the physics of the diagrams of the full SM remains, and in Fig. 2 we illustrate the schematic diagrams in the full theory which underlie the effective vertices and parameters.

We now enlarge our considerations to include non-valence degrees of freedom in the meson wave functions. The form of the parametrization itself does not change, though additional terms arise from the decay processes which do not appear in valence approximation. Turning to Eq.(5) in Ref. [45], we see, adopting their conventions [The numerical values of the $C_i(\mu)$ can depend on the explicit form of the operators [54], but this impacts neither the identification of the effective parameters nor their numerical values [45].], that the terms of form

$$\frac{G_F}{\sqrt{2}} V_{cb}^* V_{cs} [C_1(\mu) Q_1^{scc}(\mu) + C_2(\mu) Q_2^{scc}(\mu)] \tag{2.11}$$

can contribute to $B \to \pi K$ and $B \to \rho K$ decay once IC in the hadron light-cone wave functions is considered. Such terms are Cabibbo-enhanced, and contain a Wilson coefficient of order unity, so that the phenomenological impact of these neglected terms can be substantial.

Let us define

$$A_1^{\rm IC}(s, q; B, M_1, M_2) = C_1 \langle M_1 M_2 | O_1^c | B \rangle_{\rm IC} + C_2 \langle M_1 M_2 | O_2^c | B \rangle_{\rm IC} \ , \tag{2.12}$$

where $q \in u, d$. The contributions to $A_1^{\rm IC}(s, q; B, M_1, M_2)$ arising from intrinsic charm are shown in a) and b) of Fig. 3 — we anticipate that contribution b), which is driven by the IC component of the Fock states of the light hadrons, is particularly significant for ρK final states. Thus, the parameter $A_1^{\rm IC}(s, q; B, M_1, M_2)$ will depend on whether ρK or πK final states are considered because the intrinsic charm content of the π and ρ are most likely different.

Fig. 3c) illustrates how intrinsic strangeness (IS) can modify the P_1 and $P_1^{\rm GIM}$ contributions — the decay topology indicated is realized from the "$CP(c, s, u; B, M_1, M_2)$ (connected penguin)" topology of Ref. [45] by pulling the strange quark line "backwards" into a $s\bar{s}$ pair. This pictorial description is not meant to trivialize the IS contribution: the latter is irrevocably associated with the bound state's structure, as it is entangled with the other quarks of its Fock component by two or more gluon attachments.

The essential point of the parametrization of Ref. [45] is that the parameters given therein are both scale and scheme independent. In particular, the emission topologies, E_i, are scale

and scheme-independent, irrespective of any penguin contributions, since we can consider transitions in which all four quark flavors are different: in such processes penguin contributions simply do not occur, and the E_i represent the physical amplitudes of the channels in question. Similarly, decay channels exist for which only annihilation topologies contribute; thus the annihilation contributions associated with the $O_1(\mu)$ and $O_2(\mu)$ are themselves scale and scheme independent. Here we consider an annihilation topology driven by the IC components of the hadrons' wave functions; thus, this contribution, too, is separately scale and scheme independent — and the additive parameter $A_1^{\rm IC}(s,q;B,M_1,M_2)$ parameterizes its contribution. In the limit of isospin symmetry, one value of $A_1^{\rm IC}(s,q;B,M_1,M_2)$ characterizes the πK final states and another characterizes those of ρK. We can now modify our earlier parametrization to include the presence of IC:

$$\mathcal{A}(B^0\to K^+\pi^-) = V_{us}V_{ub}^*(E_1(s,u,u;B^0,K^+,\pi^-) - P_1^{\rm GIM}(s,u;B^0,K^+,\pi^-))$$
$$-V_{ts}V_{tb}^*P_1(s,u;B^0,K^+,\pi^-) + V_{cs}V_{cb}^*A_1^{\rm IC}(s,u;B^0,K^+,\pi^-) \quad (2.13)$$
$$\mathcal{A}(B^+\to K^+\pi^0) = \frac{V_{us}V_{ub}^*}{\sqrt{2}}(E_1(s,u,u;B^+,K^+,\pi^0) + E_2(u,u,s;B^+,\pi^0,K^+)$$
$$-P_1^{\rm GIM}(s,u;B^+,K^+,\pi^0) + A_1(s,u,u;B^+,K^+,\pi^0))$$
$$-\frac{V_{ts}V_{tb}^*}{\sqrt{2}}P_1(s,u;B^+,K^+,\pi^0) + \frac{V_{cs}V_{cb}^*}{\sqrt{2}}A_1^{\rm IC}(s,u;B^+,K^+,\pi^0) + \Delta A(B^+\to K^+\pi^0) \quad (2.14)$$
$$\mathcal{A}(B^+\to K^0\pi^+) = V_{us}V_{ub}^*(A_1(s,d,u;B^+,K^0,\pi^+) - P_1^{\rm GIM}(s,d;B^+,K^0,\pi^+))$$
$$-V_{ts}V_{tb}^*P_1(s,d;B^+,K^0,\pi^+) + V_{cs}V_{cb}^*A_1^{\rm IC}(s,d;B^+,K^0,\pi^+) \quad (2.15)$$
$$A(B^0\to K^0\pi^0) = \frac{V_{us}V_{ub}^*}{\sqrt{2}}(E_2(u,u,s;B^0,\pi^0,K^0) + P_1^{\rm GIM}(s,d;B^0,K^0,\pi^0)$$
$$+\frac{V_{ts}V_{tb}^*}{\sqrt{2}}P_1(s,d;B^0,K^0,\pi^0) - \frac{V_{cs}V_{cb}^*}{\sqrt{2}}A_1^{\rm IC}(s,d;B^0,K^0,\pi^0) + \Delta A(B^0\to K^0\pi^0)\ , \quad (2.16)$$

as earlier the ρK final states are realized by the replacement $\pi\to\rho$. The structure of these relations show us that under the assumption of the unitarity of the CKM matrix IC enters in a manner *identical* to that of the penguin contributions P_1 and $P_1^{\rm GIM}$ — so that the contributions cannot be distinguished phenomenologically. The numerical size of the penguin contributions has been debated in the literature: in particular, Ciuchini et al have argued that the penguin contraction of the O_2^c operator, the "charming" penguin, entering P_1 and $P_1^{\rm GIM}$, can be subject to non-perturbative enhancement [24,25]. Recently, moreover, Ciuchini et al. have argued that the phenomenological deficiencies of the "QCD factorization" approach, without annihilation contributions, in describing the $B\to\pi K$ branching ratios can be rectified by introducing an additive phenomenological contribution to P_1 [26]. It is evident that such a phenomenological treatment will mimic the impact of IC! Thus we see that charming penguins need not be penguins at all. Ciuchini et al. have discussed that a variety of physical effects, such as annihilation contributions, e.g., could be covered by the "charming penguin" aegis; we have shown that IC is another effect, non-perturbative in nature, and potentially of substantial numerical size, which contributes in an identical manner.

The recent phenomenological analysis of Ciuchini et al. is of high interest [26], though it is unsatisfying: it is evident that annihilation terms, which they neglect, are numerically important [38,55]. Thus we need a theoretical framework in which the annihilation contri-

butions, including those associated with IC, can be estimated. To do this, we adopt the perturbative QCD treatment of Li, Yu, and collaborators [42,43,56,57], to which we now turn.

III. PERTURBATIVE QCD FRAMEWORK

In the usual perturbative QCD treatment of exclusive processes, the amplitude for a particular exclusive process is formed by the convolution of the nonperturbative distribution amplitudes, $\phi_H(x,Q)$, with the hard scattering amplitude T_H computed from the scattering of on-shell, collinear quarks [58,59]. To wit, for $B \to M_1 M_2$, we have

$$\mathcal{M}(B \to M_1 M_2) = \int_0^1 dz \int_0^1 dy \int_0^1 dx \phi_B(x,Q) T_H(x,y,z) \phi_{M_1}(y,Q) \phi_{M_2}(z,Q) \; , \tag{3.1}$$

where, e.g., $\phi_{M_2}(z,Q) = \int_0^Q d^2k_\perp \phi(x,k_\perp,\lambda_i)$. This formula is suitable if the distribution amplitudes vanish sufficiently rapidly at the endpoints, and if $\alpha_s(\mu)$ is sufficiently small for a perturbative treatment to be germane. However, in the case of some electromagnetic form factors [60], as here in B decay [61,62], the distribution amplitudes may not fall sufficiently quickly at the endpoints to permit both criteria to be satisfied. Eq. (3.1) itself emerges from an expansion of T_H in powers of $k_\perp^2/Q^2$; the solution [63,64] is to reorganize the contributions in $k_\perp$, so that the contributions to the hard scattering in the transverse configuration space (b, conjugate to $k_\perp$) are no longer of point-like size. The b dependence of the reorganized distribution amplitudes, the so-called "Sudakov exponent," suppresses the large b region, so that the resulting integrals are convergent and $\alpha_s(\mu)$ is more or less consistently small. As Q^2 increases, the Sudakov mechanism becomes more effective at screening the large b region [64]. The soft portion of the integrand can also be regulated by adopting a "frozen" running coupling for sufficiently low scales [65] or by incorporating transverse structure in the light-cone wave functions [66]. The procedure of Ref. [42,43] is appealing in its simplicity as a single parameter, Λ_{QCD}, suffices to regulate the nonperturbative dynamics.

In the approach of Refs. [42,43,56,57] the transition amplitude for $B \to M_1 M_2$ decay is estimated via a three-scale factorization formula, typified by the convolution

$$C(t) \otimes H(t) \otimes \phi(x,b) \otimes \exp\left[-s(P,b) - 2 \int_{1/b}^{t} \frac{d\bar{\mu}}{\bar{\mu}} \gamma_q[\alpha_s(\bar{\mu})] \right] \; , \tag{3.2}$$

where $\gamma_q = -\alpha_s/\pi$ is the quark anomalous dimension in axial gauge and $t \sim \mathcal{O}(m_B)$. The Wilson coefficient $C(t)$ reflects the RG evolution of a term of the effective Hamiltonian, such as given in Eq. (2.1), from the W mass scale, M_W, to the hard scale t. The light-cone wave function $\phi(x,b)$, on the other hand, parameterizes the nonperturbative dynamics manifest at scales below $1/b \sim \mathcal{O}(\Lambda_{QCD})$. The exponential factor organizes the large logarithms which occur when the system is evolved from the scale t to that of $1/b$; this contribution includes the Sudakov form factor, with exponent $s(P,b)$ and P the dominant light-cone component of a meson momentum, and suppresses the contribution of the large b region. The remaining part, the hard scattering amplitude $H(t)$, is channel-dependent and can be calculated perturbatively using the operators O_i of the effective Hamiltonian. The diagrams which comprise $H(t)$ for $B \to \pi K$ decay in $\mathcal{O}(\alpha_s)$ are shown in Fig. 4 — we consider

specifically $B \to \pi K$ decay in this section as this system has recently been treated by Li, Keum, and Sanda in the perturbative QCD approach [38]. These are the only diagrams in $\mathcal{O}(\alpha_s)$; contributions without a hard-gluon exchange between the spectator and other quarks are suppressed by wave functions and do not contribute [61,67].

Let us consider the leading-order contributions to $H(t)$ in $B \to \pi K$ decay. In diagrams a) and b) the K meson is produced in a color singlet state, so that the lower half of the process "factorizes" from the meson emission and represents a computation of the $B \to \pi$ form factor. In diagrams c) and d) the weak process produces the mesons in a color octet state; the exchanged, hard gluon ensures that the outgoing mesons emerge in color-singlet states. Finally, diagrams e)-h) describe annihilation processes; the gluon emission produces a $q\bar{q}$ pair, so that two color-singlet mesons can be produced. The limitations of the usual factorization formula, Eq. (3.1), become apparent upon the evaluation of the diagram in Fig. 4a) [61,62,68,42] for the $B \to \pi$ form factor. If the transverse momentum dependence of the quarks is neglected, the heavy b quark generates a singularity as x_3, the longitudinal momentum fraction of the spectator quark in the π-meson, goes to zero [69]. However, this singularity no longer occurs if the $k_\perp$ dependence in the heavy quark propagator is retained; thus the $B \to \pi$ form factor is regarded as perturbatively calculable, once proper resummation techniques are applied [42]. In the approach of Beneke *et al.* [44], however, the $B \to \pi$ form factor is treated as non-perturbative input. One consequence of this is that the perturbative corrections in the two schemes are organized differently: in Ref. [44], hard-scattering contributions in $\mathcal{O}(\alpha_s)$ are retained which would be regarded as non-leading order in the approach of Ref. [42]. However, infrared enhancements also plague the treatment of annihilation contributions in this approach [55]; thus the treatment of Ref. [42] is more systematic in that the endpoint singularities contributions from the contributions of all decay topologies are regulated in a consistent way. For further comparison of the two approaches, see Refs. [70,55,71].

Adopting the notation and conventions of Ref. [38] we find that the parameters of Eqs. (2.7-2.10) are given by [72]

$$\begin{aligned}
E_1(s,u,u;B,K,\pi) &= -f_K F_e - M_e\,, \\
E_2(u,u,s;B,\pi,K) &= -f_\pi F_{eK} - M_{eK}\,, \\
A_1(s,u,u;B,K,\pi) &= -f_B F_a - M_a\,, \\
P_1(s,q,B,K,\pi) &= -f_K F_e^P - M_e^P - f_B F_a^P - M_a^P\,, \\
P_1^{\rm GIM}(s,q,B,K,\pi) &= 0\,, \\
\Delta A(B\to\pi K) &= V_{ts}V_{tb}^*(f_\pi F_{eK}^P + M_{eK}^P)\,,
\end{aligned} \tag{3.3}$$

where "factorizable" and "non-factorizable" contributions are denoted by F and M, respectively. The subscripts e and a refers to emission and annihilation topologies, respectively, and the P superscript reflects the presence of penguin operators in the hard-scattering amplitude. Diagrams a) and b) in Fig. 4 give rise to the contribution $f_K F_e^{(P)}$, whereas Figs. 4 c) and d) give rise to $M_e^{(P)}$. "Factorizable" in this context means that the production of one meson is independent of that of the other, so that as illustrated in diagrams a) and b), e.g., the emission of the K-meson decouples from the dynamics of the $B \to \pi$ form factor. Switching the $\bar{s}$ and $\bar{u}$ labels in diagrams a)-d) of Fig. 4, one finds that the new diagrams a) and b) give rise to the contribution $f_\pi F_{eK}^{(P)}$, whereas the new diagrams c) and d) give rise

to $M_{eK}^{(P)}$. Note, moreover, diagrams e) and f) in Fig. 4 give rise to $f_B F_A$ whereas g) and h) give rise to M_A. Note that $q \in (u, d)$; the contributions F_e^P, M_e^P, F_a^P, and M_a^P subsume quark-charge-dependent pieces. Finally, electroweak penguin contributions give rise to the terms denoted by $\Delta A(B \to \pi K)$.

The explicit expressions for the various F and M are detailed in Ref. [38] and references therein. The best fit of their resulting branching ratios to the data of Ref. [49] suggests that $\gamma = \phi_3 \approx 90^\circ$. Combining the CLEO [49] results with the more recent measurements at BaBar [50] and Belle [51], noting Table I, suggest that γ can be smaller than 90°, a result in closer accord to the value of γ determined from global fits of the CKM matrix in the SM [47,73]. Significant CP asymmetries in the $B^\pm \to K^\pm \pi^0$ and $B^0(\bar{B}^0) \to K^\pm \pi^\mp$ modes are also predicted. The branching ratios in these modes are sensitive to γ. However, we shall show that the IC contribution may impact this picture in a significant way.

Significant CP asymmetries in the $B^\pm \to K^\pm \pi^0$ and $B^0(\bar{B}^0) \to K^\pm \pi^\mp$ decays are predicted in Ref. [74] since the factorizable annihilation diagrams generate significant strong phases through imaginary parts in the hard scattering amplitudes. Li, Keum, and Sanda argue that the Sudakov suppression mechanism required to regulate the endpoint regions of the amplitudes is reliable since some 90% of the contributions to the F_{e6}^P form factor comes from the region where $\alpha_s/\pi < 0.3$ [74,38]. This is an encouraging, if weak, self-consistency check, though it is crucial to recognize that this remark pertains specifically to the value of α_s associated with the gluon exchanged in the computation of the hard scattering kernel. The function $\alpha_s(\bar{\mu})$ also appears parametrically in the expression for the Sudakov exponents; here $\alpha_s(\bar{\mu})$ can be evaluated at the softest scale, namely $\alpha_s(1/b)$ with $b \sim 1/\Lambda_{\rm QCD}^{(4)}$. The sensitivity of the numerical results to soft physics can thus be codified in terms of the sensitivity of the results to the cut-off in the integration over the transverse coordinates. In specific, if we change the cut-off from $1/\Lambda_{\rm QCD}^{(4)}$ to $0.45/\Lambda_{\rm QCD}^{(4)}$, so that $\alpha_s(\bar{\mu})/\pi \leq 0.30$ throughout the entire integral, we find that the F_e^6 form factor changes by a factor of three. The numerical results depend sensitively on the manner in which the endpoint region is regulated. Hopefully the inclusion of "threshold resummation" will help mitigate this sensitivity [75]. It should also be emphasized that a more realistic regulator might comprise of the Sudakov suppression mechanism with a systematic scale setting procedure in concert with a "frozen" value of α_s at very low scales. In this manner, the sensitivity of the numerical results to the endpoint regions can be reduced, but we postpone this investigation to a later publication.

Let us now proceed to investigate the consequences of IC in $B \to \pi K$ decay. The hard scattering contribution to $B \to \pi K$ decay mediated by IC in the B-meson is shown in Fig. 5; the hard gluon must emerge from the spectator valence quark to generate a non-trivial contribution. This diagram gives rise to the parameter $A_1^{\rm IC}(s, q, B, K, \pi)$ in Eqs. (2.13-2.16). It is worth noting that IC in the K or the π mesons could also generate a contribution to $A_1^{\rm IC}(s, q, B, K, \pi)$, as illustrated schematically in Fig. 3b). We ignore this contribution, as we believe that IC in the B-meson to be of greater phenomenological significance. In order to estimate the contribution associated with Fig. 5, we must consider the structure of the B meson in the presence of IC. It is to this issue we now turn.

A. IC in the B Meson

The B-meson light-cone wave function $\mid \Psi_B \rangle$ as an eigenstate of its light-cone Hamiltonian H_{lc} satisfies the equation $(M_B^2 - H_{lc}) \mid \Psi_B \rangle = 0$. Expanding $\mid \Psi_B \rangle$ in terms of its Fock state decomposition yields an infinite set of integral equations for the Fock components $\psi_n(x_i, \vec{k}_{\perp i}, \lambda_i)$, so that

$$\psi_n(x_i, \vec{k}_{\perp i}, \lambda_i) = \frac{1}{M_B^2 - \sum_i m_{\perp i}^2/x_i} (\mathbf{V_{lc}} \Psi_B)_n \ , \tag{3.4}$$

where $\mathbf{V_{lc}}$ is the potential matrix in the Fock-state basis. The sum is over the partons in the Fock component n, namely $\sum_i k_i^-$, and $m_{\perp i}^2 \equiv \vec{k}_{\perp i}^2 + m_i^2$. The kinematics dictate $\sum_i x_i = 1$. Equation (3.4) suggests that minimizing the value of $\sum_i m_{\perp i}^2/x_i$ enhances the likelihood of a particular Fock-state component. This is equivalent to the statement that the kinematics of the particles in each Fock state tend to minimize the total light-cone energy. Thus the most favored configurations have zero relative rapidity and $x_i \propto m_{\perp i}$.

We enumerate the lowest-lying $\mid bc\bar{c}\bar{q} \rangle$ states consistent with the 0^- B-meson state in Table II, and we assess their relative likelihood, that is, their relative light-cone energy, by estimating the total invariant mass of each configuration of specific spin by summing the mass of each individual $(q_1\bar{q}_2)_{J^P}$ state. [For simplicity, we neglect the mass shift associated with the relative motion of the $q_1\bar{q}_2$ states.] To effect this estimate, we use the empirically determined masses [76], supplemented by model calculations [77,78] for those states which have not yet been observed [79]. The net result is that the enumerated states listed in Table II are of comparable likelihood.

The pattern of estimated masses can be understood by noting that in the heavy quark limit, the spin of the heavy quark decouples, and the ground state is a degenerate doublet of $(0^-, 1^-)$ states [80]. There are two excited P-wave doublets, comprised of $(0^+, 1^+)$ and $(1^+, 2^+)$ states; an estimate of spin-orbit effects suggests that the $(1^+, 2^+)$ doublet is more massive [81]. More recent investigations have suggested, however, that the sign of the spin-orbit force in heavy-light systems is inverted relative to atomic expectations [82], so that the $(0^+, 1^+)$ doublet is pushed to higher energy [82,77]; nevertheless, the configurations of two 1^+ mesons in a relative $l = 1$ state can be discounted, as they are more massive and thus less likely than the states enumerated in Table II. The combination of 0^- and 2^+ mesons in a relative $l = 2$ state can be similarly neglected, as the masses of the mesons themselves are comparable or slightly larger than those of the 0^- and 0^+ configurations — and the angular momentum penalty is two units. All the IC configurations involve at least one $l = 1$ configuration, either within or between the meson states. This implies that the numerator of the light-cone wavefunctions can involve terms which are linear in the relative transverse momenta $k_{\perp i}$ [83].

The recent observation of a significant branching ratio of B's decaying into the scalar P-wave state of charmonium [84]

$$\frac{\mathcal{B}(B^+ \to K^+ \chi_0)}{\mathcal{B}(B^+ \to K^+ J/\psi)} = 0.77 \pm 0.11 \pm 0.22 \tag{3.5}$$

is surprising from the standpoint of naive factorization since the $V-A$ structure of the weak vertex forbids the decay to χ_0 and χ_2 states [85,86]. However, QCD radiative corrections

to the inclusive decay rate are sizeable [85], so that once these corrections are included, the decay rates into the $\chi_{0,2}K$ and $J/\psi K$ channels may not be dissimilar [86].

On the other hand, the large observed decay rate for $B^+ \to K^+\chi_0$ could be evidence for a transition involving intrinsic charm of the B since the required orbital angular momentum is naturally present. The first four configurations in Table II are suggestive of a mechanism in which the $b\bar{q}$ annihilate to a K meson; the remaining charmonium state of either 1^-, 0^-, or 0^+ character is freed from the B-meson wave function. Helicity arguments favor annihilation in the 1^- ($b\bar{q}$) state, though the suppression associated with annihilation in 0^- state may not be severe [38]. Nevertheless, this mechanism populates the $\eta_c K$, $J/\psi K$, and $\chi_0 K$ final states, though the decay to the χ_2 state, however, is disfavored, offering a signal which can distinguish intrinsic charm from other mechanisms. Note that the weak annihilation of the $b\bar{q}$ quarks must be mediated by penguin operators to avoid suppression by CKM factors.

The IC mechanism can occur in inclusive decays as well. The likelihood of materializing IC Fock states is enhanced at the value of x_F which corresponds to the typical x of the IC configuration in the parent hadron [13]; in B-decay, this means that $x_{c\bar{c}} \sim \langle x_c \rangle + \langle x_{\bar{c}} \rangle \approx 0.44$ [20]. The kinematics, in turn, determine the invariant mass of X_s for which this value of $x_{c\bar{c}}$ can be realized, namely $M_B^2 = (M_{X_s}^2 + k_\perp^2)/(1 - x_{c\bar{c}}) + (M_{c\bar{c}}^2 + k_\perp^2)/x_{c\bar{c}}$, so that the IC mechanism is particularly favored in this region of X_s invariant mass. Note that the decay $B \to \chi_2 X_s$ continues to be disfavored. We can compare this scenario with the NRQCD prediction which emerges in leading order in the relative velocity of the $c\bar{c}$ pair; in this alternative picture, the B decays to the P-wave charmonia can be mediated by a decay to a color-octet charmonium state, which evolves non-perturbatively to a color-singlet χ_J state [87,85]. This mechanism, in concert with the color-singlet mechanism, computed in NLO, suggests that $\mathcal{B}(B \to \chi_2 X_s)/\mathcal{B}(B \to \chi_0 X_s) \sim 5$ [85]. Thus a comparison of the $B \to \chi_{0,2} X_s$ branching ratios as a function of X_s should serve to test and distinguish the IC and color-octet mechanisms.

Returning to Table II we see that the $\bar{D}$ and the 1^- and 0^+ $\bar{D}^*$ states can also be liberated from the B-meson via annihilation of the $\bar{b}c$ quarks. Indeed, this mechanism is illustrated for decay into the $\bar{D}^*$ (1^-) state in Fig.2e) of Ref. [20] — the annihilation can occur without CKM suppression into $\bar{d}u$ and $\bar{s}c$ quarks, or into $l^+\nu_l$. Moreover, the annihilation in this case can occur with a Wilson coefficient of order unity, so that the mechanism is numerically more significant in this case than in decays to charmonium final states. The helicity-favored process can yield either a $\bar{D}$ or $\bar{D}^*$ meson. IC has been argued to impact the $B \to \bar{D}^*$ form factor near the zero recoil region, and thus the extraction of $|V_{cb}|$ [20]. The error in $|V_{cb}|$ is of the order of 5% [76]; yet the presence of IC could become an important consideration as the form-factor studies gain in precision. The materialization of IC Fock states is more prominent at larger values of x_F; this argues for the materialization of IC in this case at small recoil. Potentially the presence of IC could impact the extrapolation of the empirically measured $B \to \bar{D}^*$ form factor to zero recoil. Interestingly, however, theoretical constraints exist on the slope and curvature of the $B \to \bar{D}^{(*)}$ form factors [88,89], and thus a comparison with high-precision experimental data may serve to identify the contribution from IC. It is worth noting that the decay to $\bar{D}_2^*$ is disfavored in the IC mechanism; comparing branching ratios to $\bar{D}_2^* h$ and $\bar{D}_0^* h$ final states could also be helpful in illuminating the mechanisms involved. These notions extend to the inclusive case as well, in a manner analogous to our earlier discussion.

Let us now proceed to specific numerical estimates of the IC mechanism in $B \to \pi K$.

B. Numerical Estimates

In constructing a first estimate of the impact of IC in $B \to \pi K$ decay we consider exclusively the role played by IC in the B-meson since we expect it to have larger phenomenological impact than the IC of the light mesons. The hard scattering diagram for $B \to \pi K$ decay in $\mathcal{O}(\alpha_s)$ is shown in Fig. 5. The color structure of the O_2^c operator which causes the weak decay also allows diagrams where the $g^* \to q\bar{q}$ originates from the $\bar{s}$ or $\bar{c}$ quark line. The light-cone-energy minimization argument suggests that $\langle x_q \rangle$ is pushed to very low x, as in Fig. 3b) of Ref. [20], so that the spectator quark must suffer a hard interaction to emerge in the pion final state with an appreciable likelihood. Thus these additional contributions are suppressed by the hadron wave functions, and we need not consider them further.

Although IC in the B meson can be in any of several low-lying configurations, as illustrated in Table II, we assume that our estimated lowest energy state, namely the $(b\bar{c})_{1^-}(c\bar{q})_{0^-}$ state, reflects the dominant arrangement of IC in the B meson wave function. Examining Fig. 5 we see that there are no hard gluon interactions across the weak vertex, so that the computation of the hard scattering amplitude factorizes. One portion of the weak vertex mediates the annihilation of the 1^- $\bar{b}c$ quarks, and the other describes the amplitude for the 0^- $(c\bar{q})$ state to emerge with the parton content of the πK final state, namely $\bar{s}q'q\bar{q}$. It should be realized, however, that the final contribution is non-factorizable since the two pieces of the hard scattering amplitude are tied together by the integration over the coordinates of the light-cone wave function of the $(b\bar{c})_{1^-}(c\bar{q})_{0^-}$ state. We know little about the form of this wave function, although the kinematics of the IC configuration suggests that it falls more steeply in the endpoint region than the wave function of the valence component. Consequently, reorganizing the perturbative contribution into integrals over the impact parameters may not be needed. Nevertheless, to effect an explicit, albeit simple, estimate we turn to the form factors computed by Ref. [38], as the amplitude for the 0^- $(c\bar{q})$ state to emerge as $\bar{s}q'q\bar{q}$ cosmetically resembles Fig. 4f). We thus estimate that

$$A_1^{\rm IC}(s,q,B,K,\pi) \sim f_{B_c^*} F_a^P \frac{a_1(m_b/2)}{a_6(m_b/2)} B \ , \tag{3.6}$$

where $f_{B_c^*} \sim 0.317$ GeV [90] reflects the annihilation of the 1^- $\bar{b}c$ quarks, and the remaining factors reflect an estimate of the lower half of the diagram in Fig. 5. The F_a^P form factor is dominated by the F_{a6}^P, engendered by the operators associated with $a_6(t)$, for which the contributions from Figs. 4 e) and f) sum, rather than approximately cancel. We note that $a_1(m_b/2)/a_6(m_b/2) \sim -20$. The parameter B reflects the probability amplitude to find the B meson in an IC configuration, as well as an adjustment for the $\sim 50\%$ penguin enhancement reported in Ref. [74], so that $B \sim 2(0.02)/3$. We take this as a rough estimate of the possible impact of IC in $B \to \pi K$ decay.

In evaluating the role of IC in $B \to \pi K$ decays, let us first consider what we might expect in its absence. We can relate the parameters of Eqs. (2.13-2.16) to the older, diagrammatic analyses of Ref. [91]. To wit, the parameter E_1 contains the factorizable, "color-allowed" tree contribution, whereas E_2 contains the factorizable, "color-suppressed" tree contribution — this is made manifest in the mapping to the amplitudes of the perturbative QCD

approach in Eq. (3.3). Typically the factorizable, "color-suppressed" tree and annihilation — specifically the $f_B F_a$ term in Eq. (3.3)) — contributions are thought to be smaller than the factorizable, "color-allowed" tree contribution [91]. The explicit numerical estimates of Ref. [38,55] support this in that the partial rate asymmetries in the modes with charged kaons, which contain the E_1 parameter, are much larger. For the purpose of our discussion, let us define the partial rate asymmetry in $K^\pm\pi^\mp$ decay as

$$A_{\rm CP}(K^+\pi^-) \equiv \frac{\mathcal{B}(B^+ \to K^+\pi^-) - \mathcal{B}(B^- \to K^-\pi^+)}{\mathcal{B}(B^+ \to K^+\pi^-) + \mathcal{B}(B^- \to K^-\pi^+)} \,, \tag{3.7}$$

where the extension to other modes is clear. We have argued (and explicit calculations suggest) that $A_{\rm CP}(K^+\pi^-)$ and $A_{\rm CP}(K^+\pi^0)$ should be larger than $A_{\rm CP}(K^0\pi^+)$ and $A_{\rm CP}(K^0\pi^0)$. Sizeable asymmetries are generated when the E_1 parameter is included; otherwise they are small. The IC contribution, $A_1^{\rm IC}$, can be significant relative to the penguin contribution, but it is small compared to E_1, so that $A_{\rm CP}(K^0\pi^+)$ and $A_{\rm CP}(K^0\pi^0)$ will continue to be small once the effects of IC are included, although they will be modified. Consequently we focus on the role of IC in $B \to K^\pm\pi^\mp$ decay. Using Eqs. (2.13,3.3) and the numerical results of Table I in Ref. [38] we find that

$$\begin{aligned} A_1^{\rm IC} &= 2.5 \cdot 10^{-4} + i\, 2.2 \cdot 10^{-3} \\ E_1 \quad &= -0.13 - i\, 4.1 \cdot 10^{-3} \\ P_1 &= 9.3 \cdot 10^{-3} - i\, 4.5 \cdot 10^{-3} \,, \end{aligned} \tag{3.8}$$

so that $|A_1^{\rm IC}|/|P_1| \sim 20\%$. The resulting branching ratios, $\mathcal{B}(B^0 \to K^+\pi^-)$ and its CP-conjugate, are plotted as a function of γ in Fig. 6. The inclusion of IC modifies the expected branching ratios, but its impact on the expected partial rate asymmetries is striking. IC modifies the penguin contribution, parameterized by P_1, substantially; in realizing $A_{\rm CP}$ this effect is amplified by the comparatively large value of E_1, generating the substantial variations seen. For similar reasons, one can also expect marked effects in $A_{\rm CP}(K^+\pi^0)$. The combinations of Wilson coefficients a_1 and a_6 differ in sign, as do the $\mathcal{O}(\lambda^2)$ pieces of $V_{cs}V_{cb}^*$ and $V_{ts}V_{tb}^*$, so that subtracting $A_1^{\rm IC}$ in Eq. (2.13) suppresses the strong phase relative to the amplitude parameterized by $V_{us}V_{ub}^* E_1$ and hence suppresses $A_{\rm CP}$. The numerical sign associated with $A_{\rm IC}$ is unclear; our remarks assume $B > 0$ in Eq. (3.6) for definiteness. A suppressed $A_{\rm CP}$, as can be realized through the inclusion of IC, is in better agreement with recent measurements, note $A_{\rm CP}(K^+\pi^-) = -0.07 \pm 0.08\,({\rm stat.}) \pm 0.02\,({\rm sys.})$ [92]. With the uncertainties from the IC contribution, we are compelled to echo the conclusion of Ref. [26]: with the present theoretical and experimental errors, it is not possible to extract the value of γ from these decays.

IV. SUMMARY

The role of non-valence components in the hadrons' light-cone wavefunctions, coupled with the hierarchical structure of the CKM matrix, offers new perspective on B decays. The effects can be striking in decays for which the tree-level contributions are Cabibbo-suppressed, so that we have focused on the role of IC in the B meson in mediating $B \to \pi K$ and $B \to \rho K$ decays. We have shown that such contributions cannot be distinguished from

the "charming" penguin contributions of Ciuchini et al. [24–26], demonstrating that charming penguins need not be penguins at all. The intrinsic charm effect, although quantitatively elusive, is sufficiently significant, on general grounds, to impact the $B \to \pi K$ and $B \to \rho K$ branching ratios, and the associated partial rate asymmetries, in a sizeable manner, though the largest variations are seen in $A_{\rm CP}(K^+\pi^-)$ and $A_{\rm CP}(K^+\pi^0)$. Thus we must agree with the discouraging conclusions of Ciuchini et al. [26]: the unestimated contributions to $B \to \pi K$ decay are of sufficient size that these modes cannot be used to estimate γ, or, better, that any discrepancy between γ determined in a fit of the $B \to \pi K$ branching ratios and that of an unitarity triangle fit is not yet meaningful.

The presence of IC in the B-meson touches many different aspects of B physics. It can mediate decays which are suppressed in valence approximation, such as $B \to J/\psi D\pi$ [20] and $B \to \chi_{0,2}K$. It offers a pathway by which the semi-leptonic branching ratio, $B_{\rm sl}$, can be suppressed without exacerbating the charm counting problem. It could potentially help explain the lifetime difference between the B meson and the Λ_b baryon.

At the present stage of knowledge of non-perturbative QCD wavefunctions, we only have a qualitative picture of the structure of non-valence components of the hadron's light-cone wave functions. Eventually lattice gauge theory or methods such as DLCQ may provide quantitative results. It is very important to re-measure the charm structure function of the proton at large x_{bj} to confirm or disprove the evidence for intrinsic charm indicated by the EMC measurements [7–9]. As emphasized by Chang and Hou [20], the decay of the B and Υ into charmed final states may provide the best phenomenological constraints on Fock states of the B containing intrinsic charm. We have emphasized here the impact of such Fock states on exclusive decays to light hadrons.

Additional insight into the magnitude of IC and/or IS contributions in the B may also be obtained from decays in which all four quarks in the weak transition differ, so that the penguin contributions which confound the clear identification of non-valence effects in $B \to (\pi, \rho)K$ decays are absent. For example, the decay $\bar{B}^0 \to K^- D^+$ proceeds at the quark-level via $b \to cs\bar{u}$. In the valence approximation W^- emission yields the only possible decay topology. The K^- meson forms from the decay of the W^- to a color-singlet $s\bar{u}$, so that the color transparency property of QCD suggests that a factorization picture should be an excellent description of this process. However additional decay topologies are possible once non-valence degrees of freedom are considered. For example, given intrinsic charm Fock states in the B-meson, the b and $\bar{c}$ can annihilate and decay to $s\bar{u}$ and form a K^- meson. The remaining c quark and spectator $\bar{d}$ form the D^+ meson. Additionally, using IS in the B-meson, the b and $\bar{s}$ can undergo W^- exchange to yield a $c\bar{u}$ final state. These quarks can rearrange with the remaining s quark and spectator $\bar{d}$ to form a D^+K^- final state. Thus the IC/IS-mediated contributions associated with the same CKM factors and short-distance operators as the conventional contributions could repair any quantitative deficiencies in the usual treatment of these decays. However, since the IC/IS contributions in such cases is neither Cabibbo nor Wilson coefficient-enhanced relative to the dominant contribution, their presence may be difficult to confirm.

ACKNOWLEDGMENTS

We are grateful to M. Beneke, G. Bonvicini, S. Dong, I. Dunietz, G. Hiller, W.-S. Hou, P. Hoyer, D.-S. Hwang, A. Kagan, M. Karliner, H. Li, H. Quinn, J. Tandean, and L. Wolfenstein for helpful comments and discussions. S.G. thanks the SLAC Theory Group for its hospitality during the completion of this work.

REFERENCES

[1] P. A. M. Dirac, Rev. Mod. Phys. **21**, 392 (1949).

[2] S. J. Brodsky and G. P. Lepage, in *Perturbative Quantum Chromodynamics,* edited by A. H. Mueller (World Scientific, Singapore, 1989) and references therein.

[3] S. J. Brodsky, P. Hoyer, C. Peterson, and N. Sakai, Phys. Lett. B **93**, 451 (1980); S. J. Brodsky, C. Peterson, and N. Sakai, Phys. Rev. D **23**, 2745 (1981).

[4] M. Franz, M. V. Polyakov, and K. Goeke, Phys. Rev. D **62**, 074024 (2000) [hep-ph/0002240].

[5] S. J. Brodsky, J. C. Collins, S. D. Ellis, J. F. Gunion, and A. H. Mueller, in *Proceedings of the 1984 Summer Study on the SSC, Snowmass, CO, June 23 - July 13, 1984*, edited by R. Donaldson and J. G. Morfin (AIP, New York, 1985).

[6] S. J. Brodsky, H. E. Haber, and J. F. Gunion, in *Proceedings of the DPF Workshop on $\bar{p}p$ Options for the Super Collider, Chicago, IL, Feb. 13-17, 1984*, edited by J. E. Pilcher and A. R. White (Chicago, 1984).

[7] E. Hoffmann and R. Moore, Z. Phys. C **20**, 71 (1983).

[8] B. W. Harris, J. Smith, and R. Vogt, Nucl. Phys. B **461**, 181 (1996) [hep-ph/9508403].

[9] F. M. Steffens, W. Melnitchouk, and A. W. Thomas, Eur. Phys. J. C **11**, 673 (1999) [hep-ph/9903441].

[10] J. Amundson, C. Schmidt, W. Tung, and X. Wang, JHEP **0010**, 031 (2000) [hep-ph/0005221].

[11] Franz et al. assume that the scale associated with the running coupling $\alpha_s(\mu)$ is that of the charm quark, $\mu \sim m_c$. However, the $c\bar{c}$ fluctuation exists within the proton bound state, and the appropriate momentum scale is that associated with the bound state itself. As the value of α_s is associated with soft scales, the other terms in the operator product expansion of $\mathcal{O}(1/m_c^2)$, which Franz et al. neglect in their numerical estimate, could well be of importance.

[12] S. J. Brodsky and M. Karliner, Phys. Rev. Lett. **78**, 4682 (1997) [hep-ph/9704379].

[13] S. J. Brodsky, P. Hoyer, A. H. Mueller, and W. Tang, Nucl. Phys. B **369**, 519 (1992).

[14] R. Vogt, S. J. Brodsky, and P. Hoyer, Nucl. Phys. B **383**, 643 (1992).

[15] R. Vogt and S. J. Brodsky, Nucl. Phys. B **438**, 261 (1995) [hep-ph/9405236].

[16] R. Vogt and S. J. Brodsky, Phys. Lett. B **349**, 569 (1995) [hep-ph/9503206].

[17] C. Bernard *et al.*, Nucl. Phys. Proc. Suppl. **94**, 346 (2001) [hep-lat/0011029].

[18] M. Neubert and C. T. Sachrajda, Nucl. Phys. B **483**, 339 (1997) [hep-ph/9603202].

[19] M. B. Voloshin, Phys. Rev. D **61**, 074026 (2000) [hep-ph/9908455]; Phys. Rept. **320**, 275 (1999) [hep-ph/9901445]; Phys. Lett. B **385**, 369 (1996) [hep-ph/9604335].

[20] C. V. Chang and W. Hou, hep-ph/0101162.

[21] R. Balest *et al.* [CLEO Collaboration], Phys. Rev. D **52**, 2661 (1995).

[22] S. E. Schrenk, presented at the 30$^{\text{th}}$ Int. Conf. on High Energy Physics, July 27 - August 2, 2000, Osaka, Japan.

[23] S. J. Brodsky and F. S. Navarra, Phys. Lett. B **411**, 152 (1997) [hep-ph/9704348].

[24] M. Ciuchini, E. Franco, G. Martinelli, and L. Silvestrini, Nucl. Phys. B **501**, 271 (1997) [hep-ph/9703353].

[25] M. Ciuchini, R. Contino, E. Franco, G. Martinelli, and L. Silvestrini, Nucl. Phys. B **512**, 3 (1998) [hep-ph/9708222]; Erratum-ibid.B531:656-660,1998.

[26] M. Ciuchini, E. Franco, G. Martinelli, M. Pierini, and L. Silvestrini, hep-ph/0104126.

[27] I. Halperin and A. Zhitnitsky, Phys. Rev. D **56**, 7247 (1997) [hep-ph/9704412].
[28] I. Halperin and A. Zhitnitsky, Phys. Rev. Lett. **80**, 438 (1998) [hep-ph/9705251].
[29] A. A. Petrov, Phys. Rev. D **58**, 054004 (1998) [hep-ph/9712497].
[30] F. Yuan and K. Chao, Phys. Rev. D **56**, 2495 (1997) [hep-ph/9706294].
[31] A. Ali, J. Chay, C. Greub, and P. Ko, Phys. Lett. B **424**, 161 (1998) [hep-ph/9712372].
[32] T. Feldmann, P. Kroll, and B. Stech, Phys. Rev. D **58**, 114006 (1998) [hep-ph/9802409].
[33] E. Kou and A. I. Sanda, hep-ph/0106159.
[34] W.-S. Hou and B. Tseng, Phys. Rev. Lett. **80**, 434 (1998) [hep-ph/9705304].
[35] A. L. Kagan and A. A. Petrov, hep-ph/9707354.
[36] A. F. Falk, M. B. Wise, and I. Dunietz, Phys. Rev. D **51**, 1183 (1995) [hep-ph/9405346].
[37] I. Dunietz, J. Incandela, F. D. Snider, and H. Yamamoto, Eur. Phys. J. C **1**, 211 (1998) [hep-ph/9612421].
[38] Y. Y. Keum, H. Li, and A. I. Sanda, Phys. Rev. D **63**, 054008 (2001) [hep-ph/0004173].
[39] M. Bauer, B. Stech, and M. Wirbel, Z. Phys. C **34**, 103 (1987); Z. Phys. C **29**, 637 (1985).
[40] A. Ali and C. Greub, Phys. Rev. D **57**, 2996 (1998) [hep-ph/9707251]; A. Ali, G. Kramer, and C. Lu, Phys. Rev. D **58**, 094009 (1998) [hep-ph/9804363]; Phys. Rev. D **59**, 014005 (1999) [hep-ph/9805403].
[41] Y. Chen, H. Cheng, B. Tseng, and K. Yang, Phys. Rev. D **60**, 094014 (1999) [hep-ph/9903453].
[42] H. Li and H. Yu, Phys. Rev. Lett. **74**, 4388 (1995) [hep-ph/9409313]; Phys. Lett. B **353**, 301 (1995); Phys. Rev. D **53**, 2480 (1996) [hep-ph/9411308].
[43] H. Li, Phys. Rev. D **52**, 3958 (1995) [hep-ph/9412340].
[44] M. Beneke, G. Buchalla, M. Neubert, and C. T. Sachrajda, Phys. Rev. Lett. **83**, 1914 (1999) [hep-ph/9905312]; M. Beneke, G. Buchalla, M. Neubert, and C. T. Sachrajda, Nucl. Phys. B **591**, 313 (2000) [hep-ph/0006124].
[45] A. J. Buras and L. Silvestrini, Nucl. Phys. B **569**, 3 (2000) [hep-ph/9812392].
[46] S. J. Brodsky and D. S. Hwang, Nucl. Phys. B **543**, 239 (1999) [hep-ph/9806358].
[47] A. Hocker, H. Lacker, S. Laplace, and F. Le Diberder, hep-ph/0104062.
[48] C. P. Jessop *et al.* [CLEO Collaboration], Phys. Rev. Lett. **85**, 2881 (2000) [hep-ex/0006008].
[49] D. Cronin-Hennessy *et al.* [CLEO Collaboration], Phys. Rev. Lett. **85**, 515 (2000).
[50] B. Aubert *et al.* [BABAR Collaboration], hep-ex/0105061.
[51] K. Abe *et al.* [BELLE Collaboration], hep-ex/0104030.
[52] A. J. Buras and L. Silvestrini, Nucl. Phys. B **548**, 293 (1999) [hep-ph/9806278].
[53] Were we interested in isospin-violating effects, we would need to include the effects of $\pi^0 - \eta, \eta'$ mixing, note S. Gardner, Phys. Rev. D **59**, 077502 (1999) [hep-ph/9806423].
[54] A. J. Buras, M. Jamin, E. Lautenbacher, and P. H. Weisz, Nucl. Phys. B **370**, 69 (1992) [B **375**, 501 (1992)]; Nucl. Phys. B **400**, 37 (1993) [hep-ph/9211304]; A. J. Buras, M. Jamin, and M. E. Lautenbacher, Nucl. Phys. B **400**, 75 (1993) [hep-ph/9211321]; improved hadronic matrix elements," Nucl. Phys. B **408**, 209 (1993) [hep-ph/9303284].
[55] M. Beneke, G. Buchalla, M. Neubert, and C. T. Sachrajda, hep-ph/0104110.
[56] C. V. Chang and H. Li, Phys. Rev. D **55**, 5577 (1997) [hep-ph/9607214].
[57] T. Yeh and H. Li, Phys. Rev. D **56**, 1615 (1997) [hep-ph/9701233].
[58] G. P. Lepage and S. J. Brodsky, Phys. Lett. B **87**, 359 (1979); G. P. Lepage and

S. J. Brodsky, Phys. Rev. Lett. **43**, 545 (1979) [Erratum-ibid. **43**, 1625 (1979)]; G. R. Farrar and D. R. Jackson, Phys. Rev. Lett. **43**, 246 (1979).
[59] G. P. Lepage and S. J. Brodsky, Phys. Rev. D **22**, 2157 (1980).
[60] N. Isgur and C. H. Llewellyn Smith, Phys. Rev. Lett. **52**, 1080 (1984); A. V. Radyushkin, Acta Phys. Polon. B **15**, 403 (1984); N. Isgur and C. H. Llewellyn Smith, Nucl. Phys. B **317**, 526 (1989); A. P. Bakulev and A. V. Radyushkin, Phys. Lett. B **271**, 223 (1991).
[61] A. Szczepaniak, E. M. Henley, and S. J. Brodsky, Phys. Lett. B **243**, 287 (1990).
[62] G. Burdman and J. F. Donoghue, Phys. Lett. B **270**, 55 (1991).
[63] J. Botts and G. Sterman, Nucl. Phys. B **325**, 62 (1989).
[64] H. Li and G. Sterman, Nucl. Phys. B **381**, 129 (1992).
[65] C. Ji, A. F. Sill, and R. M. Lombard, Phys. Rev. D **36**, 165 (1987); J. M. Cornwall, Phys. Rev. D **26**, 1453 (1982).
[66] Z. Dziembowski and L. Mankiewicz, Phys. Rev. Lett. **58**, 2175 (1987); T. Huang, Q. Shen, and P. Kroll, Z. Phys. C **50**, 139 (1991).
[67] C. Lu, K. Ukai, and M. Yang, Phys. Rev. D **63**, 074009 (2001) [hep-ph/0004213].
[68] R. Akhoury, G. Sterman, and Y. P. Yao, Phys. Rev. D **50**, 358 (1994).
[69] We define $p^{\pm} \equiv (p^0 \pm p^3)/\sqrt{2}$ as per Ref. [38] and write $p = (p^+, p^-, \mathbf{p}_\perp)$ where $\mathbf{p}_\perp = (\mathbf{p^1}, \mathbf{p^2})$. We work in the B meson rest frame, so that $p_B = (m_B/\sqrt{2})(1, 1, 0_\perp)$.
[70] Y. Keum and H. Li, Phys. Rev. D **63**, 074006 (2001) [hep-ph/0006001].
[71] S. J. Brodsky, hep-ph/0104153.
[72] The sign associated with the F_e^P, M_e^P, F_a^P, and M_a^P terms in $B \to \pi^0 K^0$ decay and its CP-conjugate, that is, in Eqs.(25) and (26) of Ref. [38] is in error.
[73] M. Ciuchini *et al.*, hep-ph/0012308.
[74] Y. Keum, H. Li, and A. I. Sanda, Phys. Lett. B **504**, 6 (2001) [hep-ph/0004004].
[75] H. Li, hep-ph/0102013; T. Kurimoto, H. Li and A. I. Sanda, hep-ph/0105003.
[76] D. E. Groom *et al.* [Particle Data Group Collaboration], Eur. Phys. J. C **15**, 1 (2000).
[77] D. Ebert, V. O. Galkin, and R. N. Faustov, Phys. Rev. D **57**, 5663 (1998) [Erratum-ibid. D **59**, 019902 (1998)] [hep-ph/9712318].
[78] S. Godfrey and N. Isgur, Phys. Rev. D **32**, 189 (1985); S. Godfrey and R. Kokoski, Phys. Rev. D **43**, 1679 (1991).
[79] The 0^+ D state, the 0^+ and 1^- B states, as well as the B_c states, have not yet been empirically identified, so that their masses are unknown. The masses of the unknown D and B states have been taken from the relativistic quark model calculations of Ref. [77], whereas the B_c masses are taken from Ref. [78].
[80] N. Isgur and M. B. Wise, Phys. Rev. Lett. **66**, 1130 (1991).
[81] A. De Rujula, H. Georgi, and S. L. Glashow, Phys. Rev. Lett. **37**, 785 (1976).
[82] N. Isgur, Phys. Rev. D **57**, 4041 (1998) and references therein.
[83] S. J. Brodsky, D. S. Hwang, B. Ma, and I. Schmidt, Nucl. Phys. B **593**, 311 (2001) [hep-th/0003082].
[84] K. Abe [Belle Collaboration], hep-ex/0107050.
[85] M. Beneke, F. Maltoni, and I. Z. Rothstein, Phys. Rev. D **59**, 054003 (1999) [hep-ph/9808360].
[86] M. Diehl and G. Hiller, JHEP **0106**, 067 (2001) [hep-ph/0105194].
[87] G. T. Bodwin, E. Braaten, T. C. Yuan, and G. P. Lepage, Phys. Rev. D **46**, 3703 (1992) [hep-ph/9208254]; G. T. Bodwin, E. Braaten, and G. P. Lepage, Phys. Rev. D **46**, 1914

(1992) [hep-lat/9205006].
[88] C. G. Boyd, B. Grinstein, and R. F. Lebed, Phys. Rev. D **56**, 6895 (1997) [hep-ph/9705252].
[89] I. Caprini, L. Lellouch, and M. Neubert, Nucl. Phys. B **530**, 153 (1998) [hep-ph/9712417].
[90] P. Colangelo, G. Nardulli, and N. Paver, Z. Phys. C **57**, 43 (1993).
[91] M. Gronau, J. L. Rosner, and D. London, Phys. Rev. Lett. **73**, 21 (1994) [hep-ph/9404282]; O. F. Hernandez, D. London, M. Gronau and J. L. Rosner, Phys. Lett. B **333**, 500 (1994) [hep-ph/9404281]; M. Gronau, O. F. Hernandez, D. London and J. L. Rosner, Phys. Rev. D **50**, 4529 (1994) [hep-ph/9404283].
[92] B. Aubert *et al.* [BaBar Collaboration], hep-ex/0107074.

FIGURES

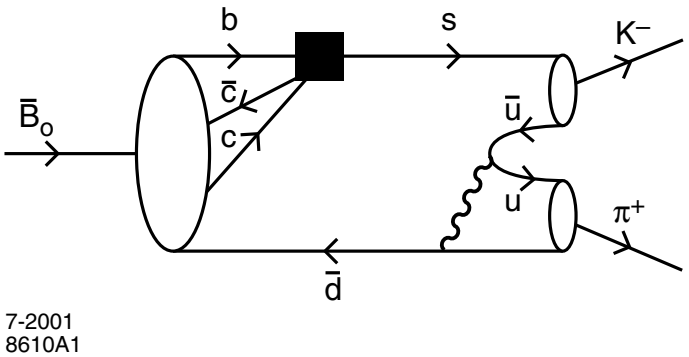

FIG. 1. Intrinsic charm in the B-meson can mediate the decay to a strange, charmless final state via the weak transition $b \to sc\overline{c}$. The square box denotes the weak transition operator.

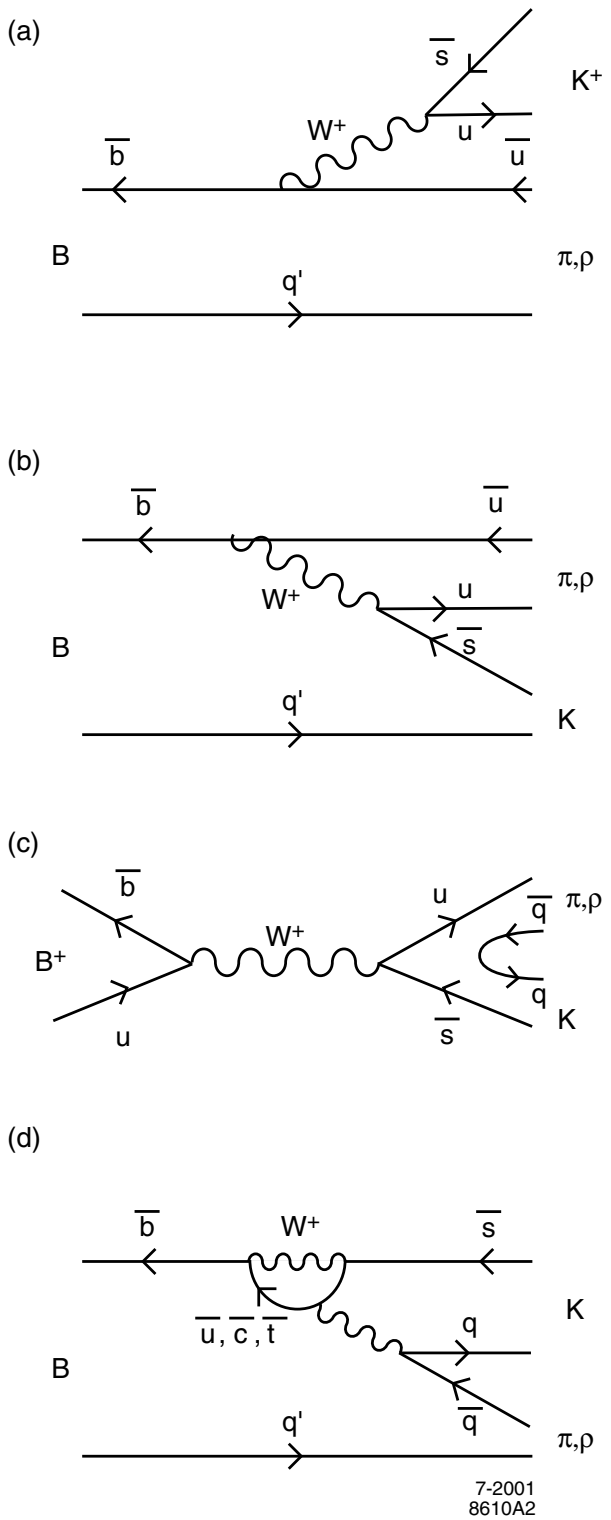

FIG. 2. Schematic illustrations of the full SM contributions to the amplitudes parametrized in Eqs. (2.7- 2.10). Diagram a) contributes to $E_1(s,u,u,B,K^+,\pi/\rho)$, diagram b) contributes to $E_2(u,u,s,B,\pi/\rho,K)$, and diagram c) contributes to $A_1(s,q,u,B^+,K,\pi/\rho)$. Diagram d) contributes to $P_1(s,q,B,K,\pi/\rho)$ and $P_1^{GIM}(s,q,B,K,\pi/\rho)$.

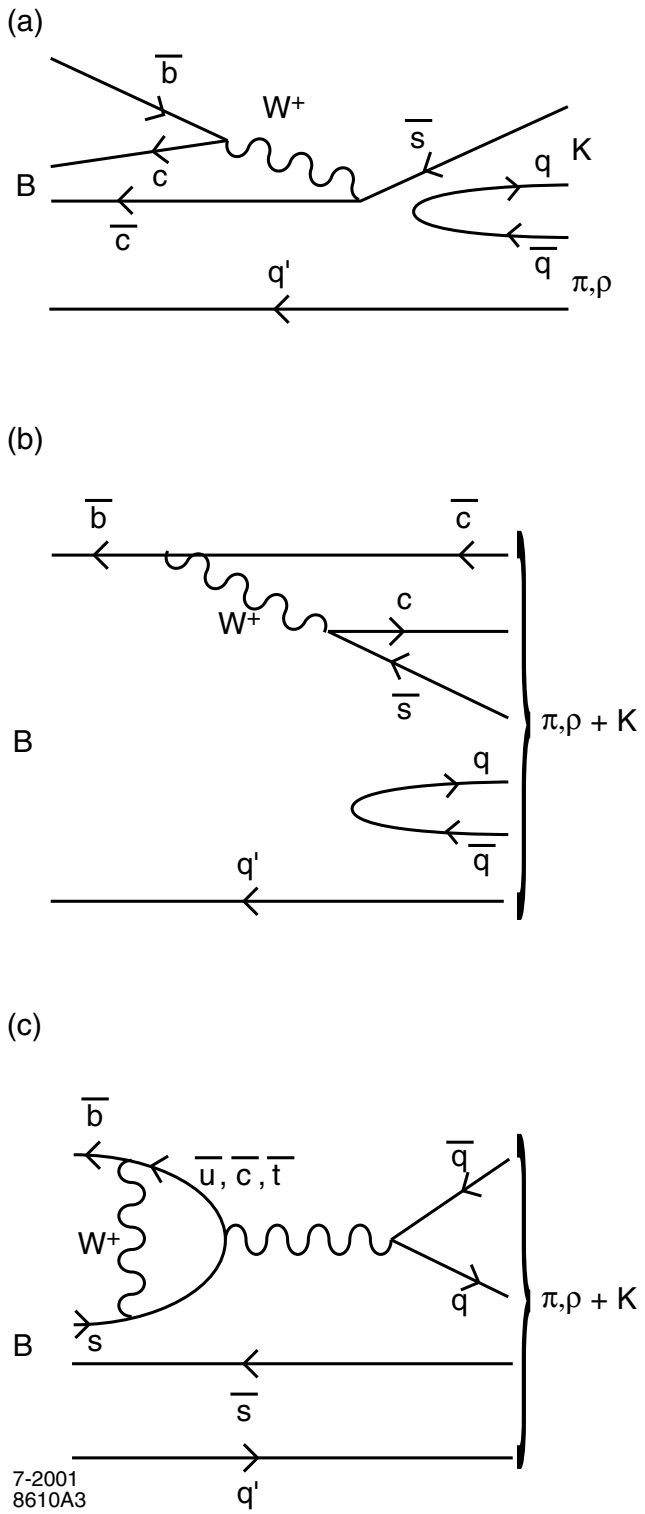

FIG. 3. Schematic illustrations of the full SM contributions to $B \to K\,\pi/\rho$ decay as mediated by intrinsic charm and strangeness in the hadron light-cone wave functions. Diagrams a) and b) contribute to $A_1^{IC}(s,q,B,K,\pi/\rho)$, whereas diagram c) modifies the value of $P_1(s,q,B,K,\pi/\rho)$ and $P_1^{GIM}(s,q,B,K,\pi/\rho)$ determined in the valence approximation.

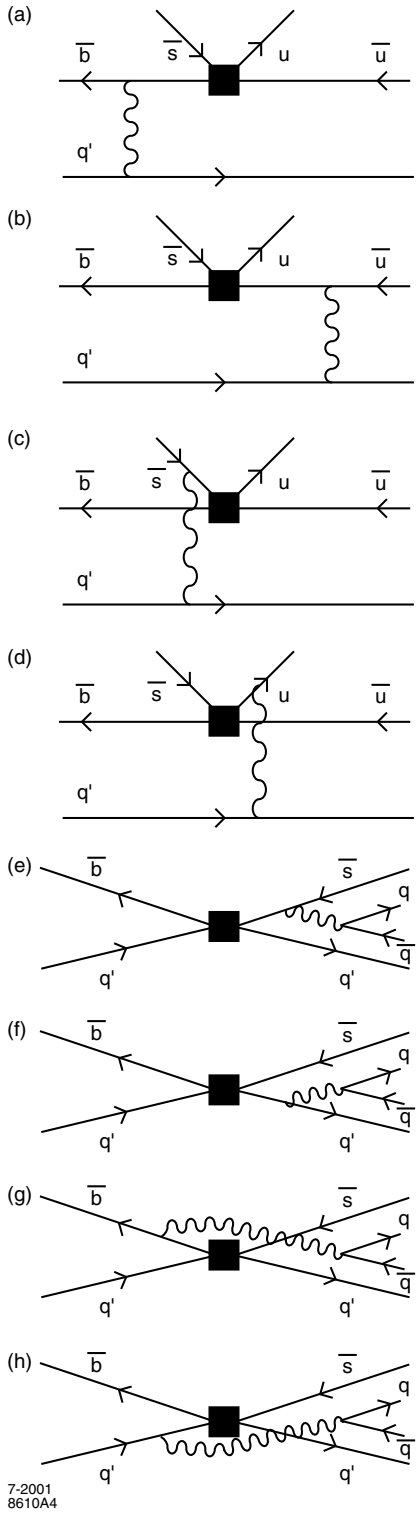

FIG. 4. Hard scattering diagrams in the valence approximation to $B \to K(\pi/\rho)$ decay.

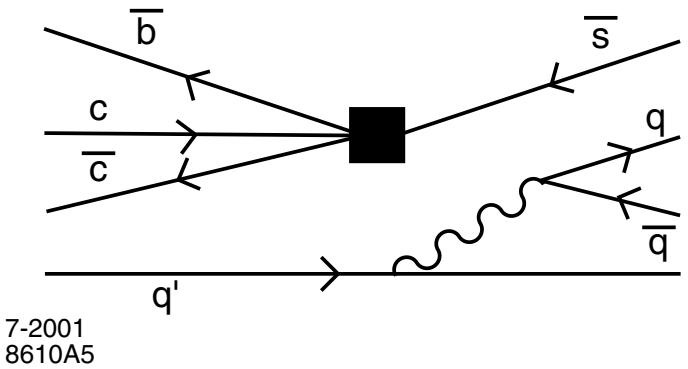

FIG. 5. Hard scattering diagram mediated by intrinsic charm in the B-meson wave function.

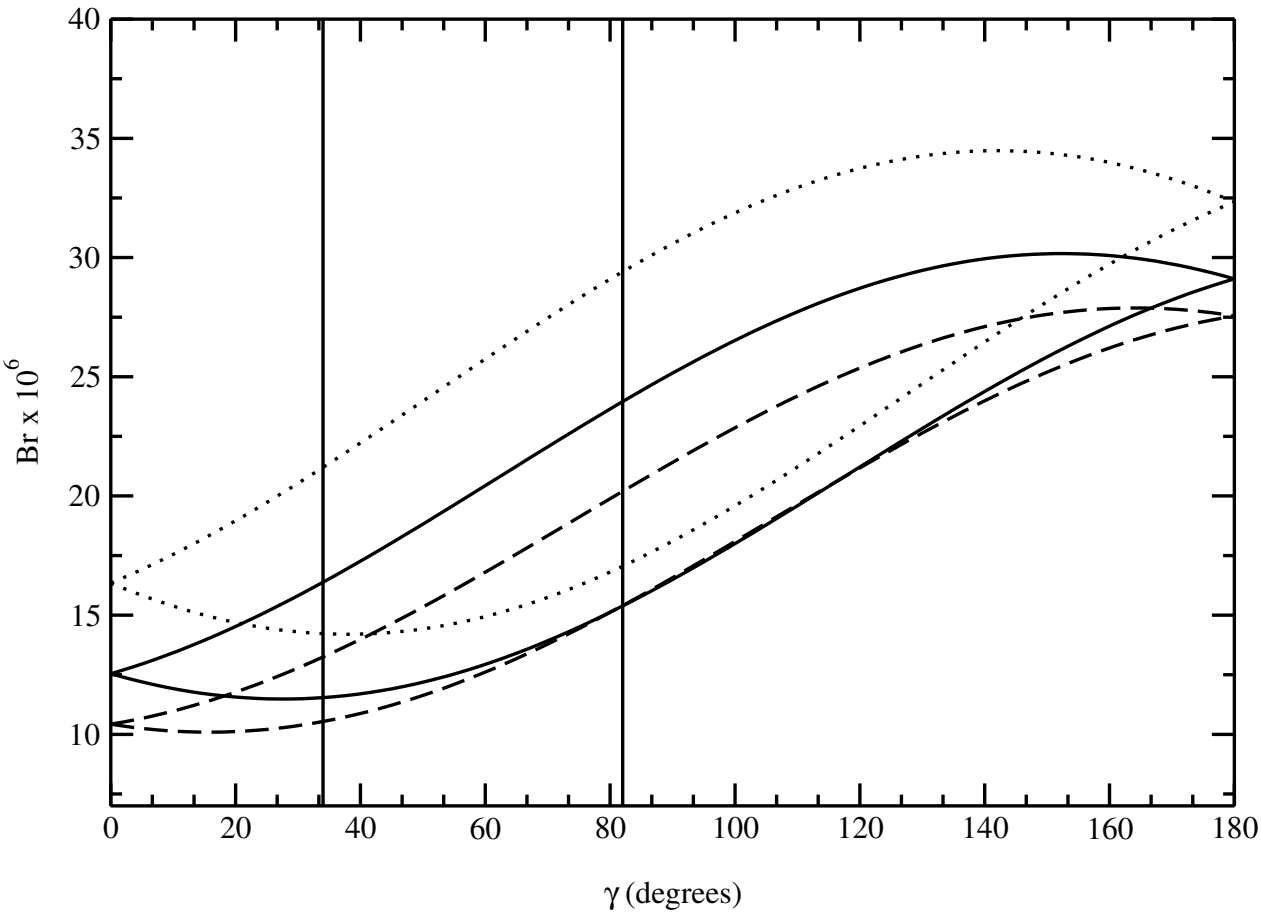

FIG. 6. The impact of IC in the B meson on the $B \to K^{\pm}\pi^{\mp}$ branching ratios as a function of γ. The *upper* curve for each type of line corresponds to $B^0 \to K^+\pi^-$ decay, whereas the *lower* curve for each type of line corresponds to $\bar{B}^0 \to K^-\pi^+$ decay. The solid line depicts the results of Ref. [38], realized from their Table I. The dashed line is the result once the IC contribution, as per Eqs. (2.13,3.6), is subtracted, and the dotted line is the result once the IC contribution, as per Eqs. (2.13,3.6), is added. The vertical solid lines enclose the $\geq 5\%$ C.L. fits to γ in the SM, $34^\circ \geq \gamma \geq 82^\circ$, of Ref. [47].

TABLES

TABLE I. Empirical branching ratios in $B \to \pi K$ decay in units of 10^{-6} from the CLEO [49], BaBar [50], and Belle [51] experiments. The symbol $\bar{\mathcal{B}}$ denotes the CP-averaged branching ratio. In constructing the averages, the statistical and systematic errors for each experiment were added in quadrature.

Mode	CLEO [49]	BaBar [50]	Belle [51]	Average
$\bar{\mathcal{B}}(\bar{B}^0 \to \pi^+ K^-)$	$17.2\,^{+2.5}_{-2.4} \pm 1.2$	$16.7 \pm 1.6 \pm 1.3$	$19.3\,^{+3.4\,+1.5}_{-3.2\,-0.6}$	17.3 ± 1.5
$\bar{\mathcal{B}}(B^- \to \pi^0 K^-)$	$11.6\,^{+3.0\,+1.4}_{-2.7\,-1.3}$	$10.8\,^{+2.1}_{-1.9} \pm 1.0$	$16.3\,^{+3.5\,+1.6}_{-3.3\,-1.8}$	$12.1\,^{+1.7}_{-1.6}$
$\bar{\mathcal{B}}(B^- \to \pi^- \bar{K}^0)$	$18.2\,^{+4.6}_{-4.0} \pm 1.6$	$18.2\,^{+3.3}_{-3.0} \pm 2.0$	$13.7^{+5.7+1.9}_{-4.8-1.8}$	$17.3\,^{+2.7}_{-2.4}$
$\bar{\mathcal{B}}(\bar{B}^0 \to \pi^0 \bar{K}^0)$	$14.6\,^{+5.9\,+2.4}_{-5.1\,-3.3}$	$8.2\,^{+3.1}_{-2.7} \pm 1.2$	$16.0^{+7.2+2.5}_{-5.9-2.7}$	$10.4\,^{+2.8}_{-2.5}$

TABLE II. Likely IC configurations in the $\bar{B}$ meson, where $q \in (u, d)$. The $|\, bc\bar{c}\bar{q}\rangle$ light-cone wave function will be maximized for states of minimal invariant mass, so that the configurations with the lowest masses are the most likely. The masses are assessed using the lowest-lying meson masses in each $(q_1\bar{q}_2)_{J^P}$ channel as described in the text. The mass increment associated with the relative motion of the two $q_1\bar{q}_2$ states has not been assessed, though it likely acts to increase the estimated mass in the $l = 1$ states.

Configuration	Estimated Mass
$(b\bar{q})_{1^-}$ $(c\bar{c})_{1^-}$ in a relative $l = 1$ state	$\gtrsim 8.4$
$(b\bar{q})_{0^-}$ $(c\bar{c})_{1^-}$ in a relative $l = 1$ state	$\gtrsim 8.4$
$(b\bar{q})_{0^-}$ $(c\bar{c})_{0^+}$ in a relative $l = 0$ state	8.7
$(b\bar{q})_{1^-}$ $(c\bar{c})_{0^-}$ in a relative $l = 1$ state	$\gtrsim 8.3$
$(b\bar{q})_{0^+}$ $(c\bar{c})_{0^-}$ in a relative $l = 0$ state	8.7
$(b\bar{c})_{1^-}$ $(c\bar{q})_{1^-}$ in a relative $l = 1$ state	$\gtrsim 8.4$
$(b\bar{c})_{0^-}$ $(c\bar{q})_{1^-}$ in a relative $l = 1$ state	$\gtrsim 8.3$
$(b\bar{c})_{0^-}$ $(c\bar{q})_{0^+}$ in a relative $l = 0$ state	8.7
$(b\bar{c})_{1^-}$ $(c\bar{q})_{0^-}$ in a relative $l = 1$ state	$\gtrsim 8.2$
$(b\bar{c})_{0^+}$ $(c\bar{q})_{0^-}$ in a relative $l = 0$ state	8.6