ISSN NO: 9726-001X

Volume 12 Issue 02 2024

JOURNAL OF CURRENT SCIENCE

A Study of Three-Body Nonleptonic B Decays within the Framework of Perturbative QCD

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Abstract

We create a formalism for three-body nonleptonic B meson decays using perturbative QCD. By developing power counting rules for various amplitude topologies, leading contributions are found. Two-meson distribution amplitudes are used to convert the analysis into the one for two-body decays. This concept, which is applicable to baryonic decays, predicts both no resonant and resonant contributions. Distinguishing between hard and soft dynamics in a QCD process is the underlying idea behind perturbative QCD (PQCD). While the latter is not calculable in perturbation theory, the former is, and is therefore treated as a universal input. The separation can be carried out within the context of kit factorization [2,3] or collinear factorization [1], in which an amplitude is defined as a convolution of a hard kernel H with a

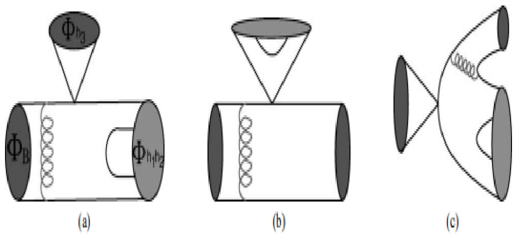

Fig. 1. Graphic definitions for topologies I, IIs, and III.

the other hand, the region with the two gluons being hard simultaneously is power-suppressed and not important. Therefore, a new input is necessary in order to catch dominant contributions to three-body decays in a simple manner. The idea is to introduce two-meson distribution amplitudes [13], by means of which a factorization formula for a B → h1h2h3 decay amplitude is written, in general, as M = ΦB ⊗ H ⊗ Ch (1) 1h2 ⊗ Φh3. It will be shown that both no resonant contributions and resonant contributions through two-body channels can be included through the parametrization of the two-meson distribution amplitude Φh1h2. Three-body decay amplitudes are classified into four topologies, depending on number of light mesons emitted from the four-fermion vertices. Topologies I and III, shown in Fig. 1(a) and (c), are associated with one light meson emission and three light meson emission, respectively. The bubbles denote the distribution amplitudes, which absorb nonperturbative dynamics. The hard kernel H contains only a single hard gluon exchange. The former involves transition of the B meson into two light mesons. In the latter case a B meson annihilates completely. For two light meson emission shown in Fig. 1(b), we assign IIs to the special amplitude corresponding to the scalar vertex, and II to the rest of the amplitudes. Both topologies II and IIs are expressed as a product of a heavy-to-light form factor and a time-like light–light form factor in the heavy-quark limit. The dominant kinematic region for three-body B meson decays is the one, where at least one pair of light mesons has the invariant mass of O (ΛM———— B) for no resonant contributions and of O(Λ——— 2) for resonant contributions, Λ———————— = MB − mb being the B meson and b quark mass difference. An example is the configuration, where all three mesons carry momenta of O(MB), but two of them move almost parallelly. In the above dominant region collinear factorization theorem applies to topology I, since it is free of end-point singularities as shown below. With the pair of mesons emitted with a small invariant mass, the evaluation of topologies II and IIs is the same as of two-body decays. The contribution from the region, where all three pairs have the invariant mass of O (M2 B), is powersuppressed. This contribution is the one, which can be calculated perturbatively in terms of the diagrams with two hard gluon exchanges. We define the power counting rules for the various topologies in the dominant kinematic region, and identify the leading ones. Consider first no resonant contributions. Topology I behaves like (ΛM—————————————————— B) −2, where one power of (ΛM———————— B) −1 comes from the hard gluon in Fig. 1(a), which kicks the soft spectator in the B meson into a fast one in a light meson [7], and another power is attributed to the invariant mass of the light meson pair. The overall product of the meson decay constants is not shown explicitly. Topology II exhibits the same power behaviour as topology I: the hard gluon in Fig. 1(b), i.e., the B meson transition form factor, gives a power of (ΛM——— B) −1, and the light–light form factor gives another power. The scalar vertex introduces an extra power m0/MB, m0 being the chiral symmetry breaking scale, to topology IIs. Topology III must involve

large energy release for producing at least a pair of fast mesons with the invariant mass of O (M2 B). That is, it behaves like (ΛM—
B) −1M−2 B. Hence, we have the relative importance of the decay amplitudes,

$$\mathcal{M}_{\rm I} : \mathcal{M}_{\rm II} : \mathcal{M}_{\rm IIs} : \mathcal{M}_{\rm III} = 1 : 1 : \frac{m_0}{M_B} : \frac{\bar{\Lambda}}{M_B},$$

indicating that topology III is negligible. For resonant contributions, we replace the power of (ΛM—
B) −1 associated with the light meson pair by Λ—
−2. Therefore, Eq. (2) still holds. Take topology I for the B+ → K+π+π− mode as an example, in which the B meson transit into a pair of peons. The π+ and π− mesons carry the momenta P1 and P2, respectively. The B meson momentum PB, the total momentum of the two peons, P = P1 + P2, and the kaon momentum P3 are chosen as

$$P_B = \frac{M_B}{\sqrt{2}}(1, 1, \mathbf{0}_T), \qquad P = \frac{M_B}{\sqrt{2}}(1, \eta, \mathbf{0}_T), \qquad P_3 = \frac{M_B}{\sqrt{2}}(0, 1-\eta, \mathbf{0}_T),$$

with the variable η = w2/M2 B, w2 = P2 being the invariant mass of the two-pion system. The light-cone coordinates have been adopted here. Define ζ = P + 1 /P + as the π+ meson momentum fraction, in terms of which, the other kinematic variables are expressed as

$$P_2^+ = (1-\zeta)P^+, \quad P_1^- = (1-\zeta)\eta P^+, \quad P_2^- = \zeta\eta P^+, \quad P_{1T}^2 = P_{2T}^2 = \zeta(1-\zeta)w^2.$$

The two peons from the B meson transition possess the invariant mass w2 ∼ O (ΛM—
B), implying the orders of magnitude P + ∼ O(MB), P − ∼ O(Λ) —
and PT ∼ O (ΛM—
B). In the heavy-quark limit, the hierarchy P + P1(2)T P − corresponds to a collinear configuration. Therefore, we introduce the two-pion distribution amplitudes [13],

$$\Phi_v(z,\zeta,w^2) = \frac{1}{2\sqrt{2N_c}}\int\frac{dy^-}{2\pi}e^{-izP^+y^-}\langle\pi^+(P_1)\pi^-(P_2)|\bar{\psi}(y^-)\not{n}\_T\psi(0)|0\rangle,$$

$$\Phi_s(z,\zeta,w^2) = \frac{1}{2\sqrt{2N_c}}\frac{P^+}{w}\int\frac{dy^-}{2\pi}e^{-izP^+y^-}\langle\pi^+(P_1)\pi^-(P_2)|\bar{\psi}(y^-)T\psi(0)|0\rangle,$$

$$\Phi_t(z,\zeta,w^2) = \frac{1}{2\sqrt{2N_c}}\frac{f_{2\pi}^\perp}{w^2}\int\frac{dy^-}{2\pi}e^{-izP^+y^-}\langle\pi^+(P_1)\pi^-(P_2)|\bar{\psi}(y^-)i\sigma_{\mu\nu}n_-^\mu P^\nu T\psi(0)|0\rangle,$$

with Nv being the twist-2 component, and is and at the twist-3 components. T = τ 3/2 is for the is vector I = 1 state, ψ the u–d doublet, z the momentum fraction carried by the spectator u quark, and n− = (0, 1, 0T) a dimensionless vector. The constant f ⊥ 2π of dimension of mass is defined via the local matrix element [13],

$$\lim_{w^2\to 0}\langle\pi^+(P_1)\pi^-(P_2)|\bar{\psi}(0)i\sigma_{\mu\nu}n_-^\mu P^\nu\frac{\tau^3}{2}\psi(0)|0\rangle = \frac{w^2}{f_{2\pi}^\perp}(2\zeta-1)P^+.$$

The matrix element with the structure γ5 n/− vanishes for topologies I and IIs, and contributes to topology II at twist 4. The one with the structure γ5 vanishes. For topologies II and IIs, a kaon–pion distribution amplitude is introduced in a similar way. For other two-pion systems, the distribution amplitudes can be defined with the appropriate choice of the matrix T. For instance, T = 1/2 is for the π0π0 is scalar (I = 0) state. A two-pion distribution amplitude can be related to the pion distribution amplitude through the calculation of the process γ γ ∗ → π+π− at large invariant mass w2 [14]. The extraction of the two-pion distribution amplitudes from the B → ππlb¯ decay has been discussed in [15]. Here we pick up the leading term in the complete Neugebauer expansion of I (z, ζ, w2) [13]:

$$\Phi_{v,t}(z,\zeta,w^2) = \frac{3F_{\pi,t}(w^2)}{\sqrt{2N_c}}z(1-z)(2\zeta-1), \qquad \Phi_s(z,\zeta,w^2) = \frac{3F_s(w^2)}{\sqrt{2N_c}}z(1-z),$$

where Fπ, set(w2) are the time-like pion electromagnetic, scalar and tensor form factors with Fπ, set (0) = 1. That is, the two-pion distribution amplitudes are normalized to the time-like form factors. For At in Eq. (9), we have adopted the parametrization,

$$\langle\pi^+(P_1)\pi^-(P_2)|\bar{\psi}(0)i\sigma_{\mu\nu}n_-^\mu P^\nu\frac{\tau^3}{2}\psi(0)|0\rangle = \frac{w^2}{f_{2\pi}^\perp}F_t(w^2)(2\zeta-1)P^+.$$

Note that the asymptotic functional form for the z dependence of Is is an assumption. For the B meson distribution amplitude, we employ the model [5],

$$\Phi_B(x) = N_B x^2(1-x)^2\exp\left[-\frac{1}{2}\left(\frac{xM_B}{\omega_B}\right)^2\right],$$

with the shape parameter dB = 0.4 GeV, and the normalization constant NB related to the decay constant fob = 190 MeV (in the convention fπ = 130 MeV) via
1 0 ΦB(x) dx = fob/ (2 √2Nc). The above ΦB is identified as Φ+ in the definition of the two leading-twist B meson distribution amplitudes Φ± given in [16,17]. Eq. (11), vanishing at x → 0, is consistent with the behaviour required by equations of motion [18]. Another distribution amplitude Φ—
B in our definition, identified as Φ—
B = (Φ− B − Φ+ B)/√2 with a zero normalization, contributes at the next-to-leading power Λ/M —
B [7]. It has been verified numerically [19] that the contribution to the B → π form factor from ΦB is much larger than from Φ—
B. The total decay rate is written as

ISSN NO: 9726-001X

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$$\Gamma = \frac{G_F^2 M_B^5}{512\pi^4} \int_0^1 d\eta (1-\eta) \int_0^1 d\zeta\, |\mathcal{M}|^2, \quad \mathcal{M} = \mathcal{M}_{\rm I} + \mathcal{M}_{\rm II} + \mathcal{M}_{\rm IIs},$$

with the amplitudes,

$$\mathcal{M}_{\rm I} = f_K \left(V_t^* \sum_{i=4,6} \mathcal{F}_{ei}^{P(u)} - V_u^* \mathcal{F}_{e2} \right), \qquad \mathcal{M}_{\rm IIs} = V_t^* F_s(\omega^2) F_{e6}^{P(d)},$$

$$\mathcal{M}_{\rm II} = (2\zeta - 1) F_\pi(\omega^2) \left[V_t^* \left(\sum_{i=3} F_{ei}^{P(d)} + \sum_{i=3,5} F_{ei}^{P(u)} \right) - V_u^* F_{e1} \right].$$

For a simpler presentation, we have assumed that the kaon–pion time-like form factor in topology IIs is equal to the pion time-like form factor multiplied by the ratio of the decay constants fake/fπ. This assumption is in fact not necessary, and the property of the kaon–pion form factor will be discussed elsewhere. The superscript P (q) stands for the amplitude from a penguin operator producing a pair of quarks q. Those without P (q) arise from tree operators. The subscript ea. stands for the emission topology (in contrast to the annihilation topology III) from the effective four-fermion operator Oi in the standard notation. We calculate the hard kernels by contracting the structures, which follow Ems. (5)– (7)

$$\frac{(\not{P}_B + M_B)\gamma_5}{\sqrt{2N_c}} \Phi_B(x), \qquad \frac{1}{\sqrt{2N_c}} \left[\not{P} \Phi_v(z,\zeta,w^2) + w \Phi_s(z,\zeta,w^2) + \frac{\not{P}_1 \not{P}_2 - \not{P}_2 \not{P}_1}{w(2\zeta - 1)} \Phi_t(z,\zeta,w^2) \right], \quad (14)$$

to Fig. 1. The factorization formulas for the $B \to 2\pi$ transition amplitudes are given by

$$\mathcal{F}_{e4}^{P(u)} = 8\pi C_F M_B^2 (1-\eta) \int_0^1 dx_1 dz \frac{\Phi_B(x_1)}{x_1 z M_B^2 + P_T^2}$$
$$\times \left\{ \left[(1+z)\Phi_v(z,\zeta,w^2) + \sqrt{\eta}(1-2z)\Phi_t(z,\zeta,w^2) + \sqrt{\eta}(1-2z)\Phi_s(z,\zeta,w^2) \right] \frac{\alpha_s(t_e^{(1)}) a_4^{(u)}(t_e^{(1)})}{z M_B^2 + P_T^2} \right.$$
$$\left. - \left[\eta \Phi_v(z,\zeta,w^2) - 2\sqrt{\eta}\Phi_s(z,\zeta,w^2) \right] \frac{\alpha_s(t_e^{(2)}) a_4^{(u)}(t_e^{(2)})}{x_1 M_B^2} \right\},$$

$$\mathcal{F}_{e6}^{P(u)} = -16\pi C_F M_B^2 r_0 \int_0^1 dx_1 dz \frac{\Phi_B(x_1)}{x_1 z M_B^2 + P_T^2}$$
$$\times \left\{ \left[(1+\eta-2z\eta)\Phi_v(z,\zeta,w^2) - \sqrt{\eta} z \Phi_t(z,\zeta,w^2) + \sqrt{\eta}(2+z)\Phi_s(z,\zeta,w^2) \right] \frac{\alpha_s(t_e^{(1)}) a_6^{(u)}(t_e^{(1)})}{z M_B^2 + P_T^2} \right.$$
$$\left. - \left[\eta \Phi_v(z,\zeta,w^2) - 2\sqrt{\eta}\Phi_s(z,\zeta,w^2) \right] \frac{\alpha_s(t_e^{(2)}) a_6^{(u)}(t_e^{(2)})}{x_1 M_B^2} \right\}, \quad (15)$$

with r0 = m0/MB. Fe2 is the same as FP (u) e4 but with a(u) 4 replaced by a2 (here a2 is close to unity). The definitions of the Wilson coefficients a(q)(t) are referred to [20]. The hard scales are defined by t (1) e = max [√Zeb, PT] and t (2) e = max [√x1 MB, PT]. The above collinear factorization formulas are well-defined, since the invariant mass of the two-pion system, proportional to PT, smears the end-point singularities from z → 0. The B meson transition form factors involved in topologies II and IIs are

$$F_{e4}^{P(d)} = 8\pi C_F M_B^2 \int_0^1 dx_1 dx_3 \int_0^\infty b_1 db_1 b_3 db_3 \Phi_B(x_1,b_1)$$
$$\times \left\{ \left[(1-\eta)(1+(1-\eta)x_3)\Phi_K(x_3) + r_0(1+\eta-2(1-\eta)x_3)\Phi_K^p(x_3) \right. \right.$$
$$\left. + r_0(1-\eta)(1-2x_3)\Phi_K^\sigma(x_3) \right] E_4^{(d)}(t_e^{(1)}) h_e(x_1,(1-\eta)x_3,b_1,b_3)$$
$$\left. + 2r_0(1-\eta)\Phi_K^p(x_3) E_4^{(d)}(t_e^{(2)}) h_e((1-\eta)x_3,x_1,b_3,b_1) \right\},$$

$$F_{e6}^{P(d)} = 16\pi C_F M_B^2 \sqrt{\eta} \int_0^1 dx_1 dx_3 \int_0^\infty b_1 db_1 b_3 db_3 \Phi_B(x_1,b_1)$$
$$\times \left\{ \left[(1-\eta)\Phi_K(x_3) + 2r_0\Phi_K^p(x_3) + r_0(1-\eta)x_3(\Phi_K^p(x_3) - \Phi_K^\sigma(x_3)) \right] \right.$$
$$\times E_6^{(d)}(t_e^{(1)}) h_e(x_1,(1-\eta)x_3,b_1,b_3)$$
$$\left. + 2r_0(1-\eta)\Phi_K^p(x_3) E_6^{(d)}(t_e^{(2)}) h_e((1-\eta)x_3,x_1,b_3,b_1) \right\}.$$

The definitions of the evolution factors E(q) I (t), which contain the Wilson coefficients a(q) I (t), of the hard functions he, and of the kaon distribution amplitudes ΦK, Up K and Φσ K are referred to [20]. FP (q) e3, FP (q) e5 and Fe1 are obtained from FP (d) e4 by substituting a(q) 3, a(q) 5 and a1 for a(d) 4, respectively. The PQCD evaluation of the form factors indicates the power behaviour in the asymptotic region, Fπ (w2) ∼ 1/w2, and their relative importance: Fest(w2)/Fπ (w2) ∼ m0/w. Therefore, the twist-3 contributions in Eq. (15) are down by a power of √η m0/w = m0/MB compared to the twist-2 ones, which is the accuracy considered here. To calculate the no resonant contribution, we propose the parametrization for the whole range of w2,

$$F_\pi^{(nr)}(w^2) = \frac{m^2}{w^2 + m^2}, \qquad F_{s,t}^{(nr)}(w^2) = \frac{m_0 m^2}{w^3 + m_0 m^2},$$

where the parameter m = 1 GeV is determined by the fit to the experimental data M2 J/ψ |Fπ (M2 J/ψ) | 2 ∼ 0.9 GeV2 [21], MJ/ψ being the J/ψ meson mass. These form factors can carry strong phases, which are assumed to be not very different, i.e., overall and negligible here.

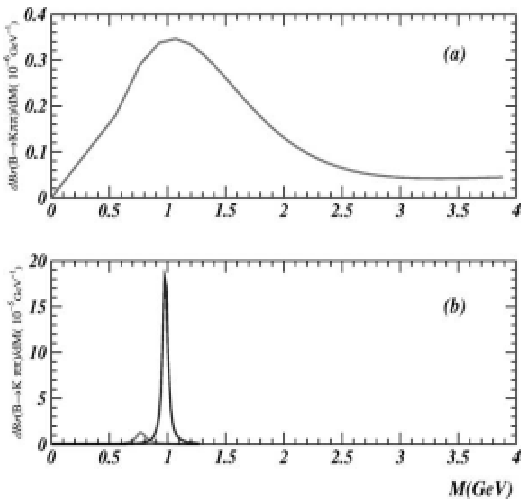

Fig. 2. (a) No resonant and (b) resonant contribution to the B+ → K+π+π− decay spectrum with respect to the two-pion invariant mass M(π+π−). The sharp peak corresponds to the f0 resonance with the width 50 MeV

To calculate the resonant contribution, we parametrize it into the time-like form factors,

$$F^{(r)}_{\pi,s,t}(w^2) = \frac{M_V^2}{\sqrt{(w^2 - M_V^2)^2 + \Gamma_V^2 w^2}} - \frac{M_V^2}{w^2 + M_V^2},$$

with ΓV being the width of the meson V. The subtraction term renders Eq. (18) exhibit the features of resonant contributions: the normalization F(r) π (0) = 0 and the asymptotic behaviour F(r) π (w2) ∼ 1/w4, which decreases at large w faster than the no resonant parametrization in Eq. (17). Eq. (18) is motivated by the pion time-like form factor measured at the ρ resonance [22]. It is likely that all F(r) π, set contain the similar resonant contributions. The relative phases among different resonances will be discussed elsewhere by employing the more sophisticated parametrization [23]. Here we assume the absence of the interference effect. We adopt m0 = 1.4 (1.7) GeV for the pion (kaon) and the unitarity angle φ3 = 90◦ [5]. For the B+ → ρ0(770) K+ and B+ → f0(980) K+ channels, we choose Γρ = 150 MeV and Γf0 = 50 MeV [24]. The no resonant contribution 0.61 × 10−6 to the B+ → K+π+π− branching ratio is obtained. Our results 1.8 × 10−6 and 13.2×10−6 are consistent with the measured three-body decay branching ratios through the B+ → ρ (770) K+ and B+ → f0(980) K+ channels, < 12 × 10−6 and (9.6+2.5+1.5+3.4 −2.3−1.5−0.8) × 10−6 [11], respectively. Since the f0 width has a large uncertainty, we also consider Γf0 = 60 MeV, and the branching ratio reduces to 10.5 × 10−6. The resonant contributions from the other channels can be analysed in a similar way. For example, the K∗ (892) resonance can be included into the K–π form factors by choosing the width ΓK∗ = 50 MeV. The no resonant and resonant contributions to the B+ → K+π+π− decay spectrum are displayed in Fig. 2. In the above formalism nonfactorizable contributions arise from the diagrams, in which a hard gluon attaches the spectator quark and the meson emitted from the weak vertex (topology I) or the meson pair (topologies

II and IIs). The nonfactorizable contributions, suppressed by ln−1(MB/Λ) ——————— [8], and topology III, being of O (Λ/M ————— B), can be evaluated systematically by means of the two-meson distribution amplitudes. The framework presented here is not only applicable to the study of three-body mesonic B meson decays, but also to baryonic decays [25], such as B → pp.¯. One simply introduces two-proton distribution amplitudes, and the calculation of the corresponding hard kernel is similar. In this Letter we have proposed a promising formalism for three-body nonleptonic B meson decays. This formalism, though at its early stage, is general enough for evaluating both no resonant and resonant contributions 264 C.-H. Chen, H.-N. Li / Physics Letters B 561 (2003) 258–265 to various modes, and as simple as that for two-body decays. In the future we shall discuss more delicate issues, such as CP asymmetries [26], phase shifts from meson–meson scattering [27], and interference effects among different resonances [28].

Acknowledgements

We thank S. Brodsky, H.Y. Cheng, M. Diehl and A. Garmash for useful discussions. This work was supported in part by the National Science Council of ROC under Grant No. NSC-91-2112-M-001-053, by the National Centre for Theoretical Sciences of ROC, and by Theory Group of KEK, Japan

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