PERTURBATIVE QCD AND NON-LEPTONIC B DECAYS[a]

Chia-Hung V. Chang
Department of Physics, National Tsing-Hua University, Hsin-Chu, Taiwan, R.O.C.

Hsiang-nan Li
Department of Physics, National Chung-Cheng University, Chia-Yi, Taiwan, R.O.C.

We develop a perturbative QCD factorization theorem for the non-leptonic heavy meson decays. The factorization involves three scales: an infrared cutoff of order Λ_{QCD}, a hard scale of order the heavy meson mass, and an ultraviolet cutoff of order the W boson mass. Unlike the conventional Bauer-Stech-Wirbel factorization approximation, our factorization is scale independent. Applying this formalism, we clarify the controversy over the parameters a_2/a_1 for charm and bottom decays. It is found that the nonfactorizable contribution plays an important role in the explanation of the sign and magnitude of a_2/a_1.

1 Non-leptonic decays

Since the birth of Quantum Chromodynamics (QCD), large scales in the processes involving this strong interaction have been exploited to *simplify* the calculation. Because of its property of asymptotic freedom, the strong coupling constant at the large scale is weak and a perturbative expansion is possible. Actually this simplification is essential in most of QCD calculations, such as deep inelastic scattering. The physics of heavy quarks also provides us with such an opportunity. Heavy quark effective field theory (HQEFT) is a beautiful example of expanding in powers of the heavy quark mass. While HQEFT is spectacularly successful in dealing with the semileptonic decays of heavy mesons, its application to their non-leptonic decays is less impressive. Nonleptonic heavy meson decays are more difficult mainly because of the complicated QCD corrections. Gluon exchanges among quarks usually generated large logarithmic corrections, which will invalidate the perturbation theory. The QCD corrections are usually taken care of by integrating out the W meson at the scale of M_W. Effective Hamiltonians are generated to reproduce the weak decays. Logarithmic corrections are calculated using renormalization group running. While the semileptonic decays involve only the form factor of conserved currents, which does not run under RG, the nonleptonic decays are

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described by four-quark current-current operators. For example, the relevant operator for the $B \to D\pi$ decays is

$$H = \frac{4G_F}{\sqrt{2}} V_{cb} V_{ud}^* (\bar{c}_L \gamma_\mu b_L)(\bar{d}_L \gamma^\mu u_L) \ . \tag{1}$$

This operator not only run but also mix with other four fermion operators under RG. The resultant effective Hamiltonian related to Eq. (1) is written as

$$H_{\rm eff} = \frac{4G_F}{\sqrt{2}} V_{cb} V_{ud}^* [\, c_1(\mu) O_1 + c_2(\mu) O_2 \,] \ , \tag{2}$$

with $O_1 = (\bar{c}_L \gamma_\mu b_L)(\bar{d}_L \gamma^\mu u_L)$ and $O_2 = (\bar{d}_L \gamma_\mu b_L)(\bar{c}_L \gamma^\mu u_L)$. This operator mixing indicates the gluon exchange will change the color indices of the quarks emerged from the weak decay before their hadronization. c_1 and c_2 are the Wilson coefficients, whose evolution from the W boson mass M_W down to a lower scale μ is determined by renormalization-group (RG) running [1]. Though Wilson coefficients are μ-dependent, physical quantities such as decay amplitudes are not. In principle, the matrix elements of the four-fermion operators contain a μ dependence, which exactly cancels that of the Wilson coefficients.

This beautiful scale independent effective field theory scheme, however, is handicapped by our inability to deal with matrix elements of four fermion operator. Unlike form factors, for which the conservation of currents usually provide us with some model independent information, various schemes and models are needed to estimate the hadronic matrix elements. The estimates are usually μ-independnt. Hence, the decay amplitudes turn out to be scale-dependent. The most popular approach is the Bauer-Stech-Wirbel (BSW) factorization approximation [2]. It is assumed that nonleptonic matrix elements can be factored into two matrix elements of (axial) vector currents. Since the currents are conserved, the matrix elements have no anomalous scale dependence. Presumably μ should be set to the dominant scale of the matrix elements. However, the matrix elements involve both the heavy quark scale and the small hadronic scale. Naively setting μ to the heavy quark mass will lose large logarithms associated with the hadronic scale. It is then quite natural that theoretical predictions by the BSW approach are sensitive to the scale we choose [3].

To circumvent this problem, a phenomenological approach is adopted to bypass the strong scale dependence. The Wilson coefficients $c_{1,2}$ are regarded as free parameters, and determined by experimental data [2]. In this model two equivalent parameters $a_1 = c_1 + c_2/N_c$ and $a_2 = c_2 + c_1/N_c$ describe the external and internal W-emission amplitudes, respectively. These two parameters, especially a_2, are quite sensitive to the scale. Despite limited success of this approach, it was found that the ratio a_2/a_1 from an individual fit to the CLEO data of $B \to D^{(*)}\pi(\rho)$ [4] varies significantly [5]. An allowed domain (a_1, a_2) exists for the three classes of decays $\bar{B}^0 \to D^{(*)+}$, $\bar{B}^0 \to D^{(*)0}$ and $B^- \to D^{(*)0}$ only when the experimental errors are expanded to a large extent[6]. Moreover, it has been very difficult to explain the two ratios associated with the $B \to J/\psi K^{(*)}$ decays [7],

$$R = \frac{\mathcal{B}(B \to J/\psi K^*)}{\mathcal{B}(B \to J/\psi K)} , \quad R_L = \frac{\mathcal{B}(B \to J/\psi K_L^*)}{\mathcal{B}(B \to J/\psi K^*)} , \tag{3}$$

simultaneously in the BSW framework, where $\mathcal{B}(B \to J/\psi K)$ is the branching ratio of the decay $B \to J/\psi K$. For all the qualified success of this approach, the problem remains to understand the parameters a_1, a_2. From the data of charm and bottom decays a negative a_2/a_1 and a positive a_2/a_1 are needed [2,8], respectively. If the scale is set at the heavy quark mass, to explain them using the formula $a_1 = c_1 + c_2/N_c$ and $a_2 = c_2 + c_1/N_c$ would require N_c to be infinity and 2.3 for charm and bottom respectively.

In our opinion, the problem stems from the scale sensitivity generated by the factorization approximation. A scale independent scheme is called for to deal with the physics below the heavy quark mass. It was shown recently that the perturbative QCD (PQCD) calculation based on the Brodsky-Lepage approach to the exclusive processes is applicable to heavy meson decays at large recoil [9,10]. The breakthrough is due to the all-order Sudakov resummation of large radiative corrections, which suppresses contributions from the long-distance region. This formalism, taking into account the evolution from the typical scale of hard subprocesses characterized by the heavy meson mass to a lower hadronic scale, is μ-independent for semileptonic decays. However, it can not be an appropriate tool for nonleptonic decays, because it does not involve the scale M_W. In this contribution we shall develop a PQCD formalism combined with the effective Hamiltonian in Eq. (2), which further incorporates the evolution from M_W down to the hard scale. This three-scale factorization theorem, as demonstrated below, moderates the scale-setting ambiguity.

2 Three scale PQCD factorization theorem

We first illustrate the main idea of PQCD factorization theorems. The key is that the amplitude can be factored into short distance parts and long distance parts. The long distance parts are characterized by their infrared divergences, while the short distance part is infrared safe and can be calculated perturbatively. In other words, by factoring out the infrared divergences order by order, we get an infrared finite and hence calculable quantity. This is possible because the infrared divergence and long distance physics happens in certain configurations that the Eikonal approximations is applicable, which simplifies the structure tremendously. We explain this idea by considering one-loop corrections to a generic decay process through a current. These corrections are ultraviolet finite, since the conserved current is not renormalized. However, they are also infrared divergent. The divergence arises in the configuration when the gluons are soft or collinear to light partons. In these configurations, Eikonal approximation can be applied. Due to the approximation these infrared divergences can be factored and isolated. Radiative corrections are classified into the reducible and irreducible types. Irreducible corrections contain only single soft logarithms and are absorbed into a soft function U. Reducible corrections, containing double logarithms from the combination of soft and collinear divergences, are absorbed into a wave function $\phi(P, b, \mu)$ and explicitly resummed into a Sudakov factor e^{-s}:

$$\phi(P, b, \mu) = \exp[-s(P, b)]\phi(b, \mu) . \tag{4}$$

b is the conjugate variable of the transverse momentum, and $1/b$ can be regarded as an infrared cutoff of order the hadonic scale. The Sudakov factor strongly suppresses the contributions from the large b region. With Sudakov suppression, the irreducible soft corrections, appearing in the form $1 - e^{i\mathbf{l}\cdot\mathbf{b}}$, l being the loop momentum, cancel asymptotically $(b \to 0)$ [10]. Hence, they will be neglected below (*ie* $U = 1$).

To factorize a one-loop correction, we divide it into two terms. In the first term, the fermion propagators are replaced by their eikonal approximation. This term has the same infrared structure as the full diagram. Being infrared sensitive, it is absorbed into U or ϕ, depending on which type the correction is. The second term, equal the full diagram with the first term subtracted , is infrared finite and has the same ultraviolet structure as the full diagram. It can be absorbed into a hard scattering amplitude $H(t, \mu)$, where t denotes the typical scale of the hard decay process. We then get the $O(\alpha_s)$ factorization formula.

$$H(t, \mu)\phi(b, \mu) \tag{5}$$

The presence of μ implies that both ϕ and H need renormalization. Let γ_ϕ be the anomalous dimension of ϕ. Then the anomalous dimension of H must be $-\gamma_\phi$, because the full diagram is ultraviolet finit. Using RG, the convolution of H with ϕ is μ-independent as indicated by

$$H(t, \mu)\phi(b, \mu) = H(t, t)\phi(b, 1/b) \exp\left[-\int_{1/b}^{t} \frac{d\bar{\mu}}{\bar{\mu}} \gamma_\phi(\alpha_s(\bar{\mu}))\right] . \tag{6}$$

The contribution characterized by momenta smaller than $1/b$, *ie* the infrared divergence, is absorbed into the initial condition $\phi(b, 1/b)$, which is of nonperturbative origin.

Indeed the effective Hamiltonian in Eq. (2) can be constructed in a similar way. Consider a typical one-loop QCD correction to the W-exchange diagram. Again we express the full diagram, which is ultraviolet finite after coupling constant and wavefunction renormalization, into two terms. The first term, obtained by shrinking the W boson line into a vertex, corresponds to the local four-fermion operators O_i. It is absorbed into a hard scattering amplitude $H(t, \mu)$, with a typical scale $t \ll M_W$, since gluons involved in this term do not "see" the W boson. The second term is equal to the full diagram with the first term subtracted. The physics below the scale M_W disappears. This term only involves momenta of order M_W, and can be absorbed into a "harder" function $H_r(M_W, \mu)$ (not a scattering amplitude), in which gluons do "see" the W boson. We obtain the $O(\alpha_s)$ factorization formula.

$$H_r(M_W, \mu) H(t, \mu) \tag{7}$$

This formula in fact represents a matrix relation because of the mixing between O_1 and O_2. Solving their RG equations, we derive

$$H_r(M_W, \mu) H(t, \mu) = H_r(M_W, M_W) H(t, t) \exp\left[\int_t^{M_W} \frac{d\bar{\mu}}{\bar{\mu}} \gamma_{H_r}(\alpha_s(\bar{\mu}))\right], \tag{8}$$

where the anomalous dimension γ_{H_r} of H_r is also a matrix. We emphasize that the factorization in Eq. (8) is not complete because of the presence of infrared divergences in H. Without large logarithms, $H_r(M_W, M_W)$ can now be safely approximated by its lowest-order expression $H_r^{(0)} = 1$.

We are now ready to contruct a three-scale factorization theorem by combining Eqs. (6) and (8). Consider the decay amplitude up to $O(\alpha_s)$ without integrating out the W boson. We first factorize out the infrared sensitive wave functions as described above. Though devoid of infrared divergences, the hard part still invloves two scales t and M_W. The second factorization is then employed to separate these two scales, and H_r can be moved out of the hard part, a step valid up to $O(\alpha_s)$. We identify the remaining contribution, including the four-fermion amplitude and the associated soft subtraction, as H. The anomalous dimension of H is given by $\gamma_H = -(\gamma_\phi + \gamma_{H_r})$. We thus get the three-scale factorization formula

$$H_r(M_W, \mu) H(t, \mu) \phi(b, \mu) = c(t) H(t, t) \phi(b, 1/b) \exp\left[-\int_{1/b}^{t} \frac{d\bar{\mu}}{\bar{\mu}} \gamma_\phi(\alpha_s(\bar{\mu}))\right] \tag{9}$$

where the exponential factor in Eq. (8) has been identified as the Wilson coefficient $c(t)$. It implies that μ in $c(\mu)$ should be set to the hard scale t. The two-stage evolutions from $1/b$ to t and from t to M_W are both included, and the final expression is μ-independent. Note that this μ independence is a direct consequence of the current conservation as stated before.

The above conclusion is quite natural from the effective field theory approach [11]. An effective field theory is constructed for a scale $\mu < M_W$ by integrating out the W boson at $\mu = M_W$. Matching corrections are determined by the matching condition requiring that the low-energy light-particle Green functions of the two theories be equal. The effective theory is then evolved by RG running from $\mu = M_W$ to a lower scale, insuring that the amplitudes are μ-independent. The scale μ in a continuum effective field theory is actually a scale to separate the long-distance from the short-distance physics with the physics above the scale μ absorbed into the coefficients in the effective Hamiltonian, such as the Wilson coefficients $c_{1,2}(\mu)$. The effective field theory constructed this way has exactly the same low-energy behaviour as the full theory, including infrared divergences, physical cuts, and etc. Thus the infrared divergences in the decay amplitudes calculated using the effective field theory can be factorized in the same way as the full theory. The factorization formula for the μ-independent amplitude, $c(M_W, \mu)H(t, \mu)\phi(b, \mu)$, is identical to Eq. (9) with the Wilson coefficient c identified as H_r.

3 Numerical results

We apply the above formalism to the nonleptonic $B(P_1) \to D(P_2)\pi(P_3)$ decays. The decay rate is given by [10]

$$\Gamma = \frac{1}{128\pi} G_F^2 |V_{cb}|^2 |V_{ud}|^2 M_B^3 \frac{(1-r^2)^3}{r} |\mathcal{M}|^2 , \tag{10}$$

with $r = M_D/M_B$, M_B (M_D) being the B (D) meson mass. In the rest frame of the B meson, P_1 has the components $P_1 = (M_B/\sqrt{2})(1, 1, \mathbf{0}_T)$. The nonvanishing components of P_2 and P_3 are respectively $P_2^+ = M_B/\sqrt{2}$, $P_2^- = rM_D/\sqrt{2}$, $P_3^+ = 0$, and $P_3^- = (1-r^2)M_B/\sqrt{2}$. Let $k_1(k_2)$ be the momentum of the light valence quark in the B (D) meson and k_3 be the momentum of a valence quark in the pion. These k's may be off-shell by the amount of their transverse components k_T of order $\Lambda_{\rm QCD}$. We define the momentum fractions x as $x_1 = k_1^-/P_1^-$, $x_2 = k_2^+/P_2^+$, and $x_3 = k_3^-/P_3^-$. To leading power in $1/M_B$, the factorization formula for $\mathcal{M}$ in the transverse configuration space is written as

$$\mathcal{M} = \int_0^1 [dx] \int_0^\infty [d^2\mathbf{b}] \phi_B(x_1, b_1, 1/b_1) \phi_D(x_2, b_2, 1/b_2) \phi_\pi(x_3, b_3, 1/b_3) \times c(t) H(x_i, b_i, t) \exp[-S(x_i, b_i)] \tag{11}$$

with $[dx] = dx_1 dx_2 dx_3$ and $[d^2\mathbf{b}] = d^2\mathbf{b}_1 d^2\mathbf{b}_2 d^2\mathbf{b}_3$. The Sudakov factor e^{-S} is the product of e^{-s} in Eq. (4) and the exponential in Eq. (9) from each wave function. Below we shall neglect the b dependence of the wave functions [9].

Without large logarithms, H can be reliably treated by perturbation theory. To leading order in α_s, the hard part for the decay $B^- \to D^0\pi^-$ consists of four sets of diagrams[12]. The first set of diagrams correspond to the external W emission [2,3], while the second to the internal W emission. They have been calculated using the PQCD formalism in [9,10] without including the Wilson coefficients. Denote their contributions to the amplitude $\mathcal{M}$ as $\mathcal{M}_a$ and $\mathcal{M}_b$. It is easy to find that the Wilson coefficients associated with $\mathcal{M}_a$ and $\mathcal{M}_b$ are respectively a_1 and a_2. Readers are referred to [10] for the complete formulas of $\mathcal{M}_a$ and $\mathcal{M}_b$. In addition, there are two sets of diagrams that do not have similar parts in the factorization approximation and will be called the nonfactorizable diagrams. Their contributions are

$$\begin{aligned}\mathcal{M}_c &= 32\sqrt{2N_c}\pi \mathcal{C}_F \sqrt{r} M_B^2 G_F \int_0^1 [dx] \int_0^\infty b_1 db_1 b_2 db_2 \phi_B(x_1)\phi_D(x_2)\phi_\pi(x_3) \\ &\times \left[\alpha_s(t_1) \frac{c_1(t_1)}{N_c} e^{-S_c^{(1)}(x_i,b_i)} (x_1 - x_2 - x_3(1-r^2)) h_c^{(1)}(x_i, b_i) \right. \\ &\left. + \alpha_s(t_2) \frac{c_1(t_2)}{N_c} e^{-S_c^{(2)}(x_i,b_i)} (1 - (x_1+x_2)(1-r^2)) h_c^{(2)}(x_i, b_i) \right] .\end{aligned} \tag{12}$$

The functions $h_c^{(j)}$, $j = 1$ and 2, are given by

$$\begin{aligned} h_c^{(j)} &= [\theta(b_1 - b_2) K_0(AM_B b_1) I_0(AM_B b_2) \\ &+ \theta(b_2 - b_1) K_0(AM_B b_2) I_0(AM_B b_1)] \\ &\times \begin{pmatrix} K_0(B_j M_B b_2) & \text{for } B_j \geq 0 \\ \frac{i\pi}{2} H_0^{(1)}(|B_j| M_B b_2) & \text{for } B_j \leq 0 \end{pmatrix} , \end{aligned} \tag{13}$$

with $A^2 = x_1 x_3(1-r^2)$, $B_1^2 = (x_1+x_2)r^2 - (1-x_1-x_2)x_3(1-r^2)$, and $B_2^2 = (x_1 - x_2)x_3(1-r^2)$. The Sudakov exponent $S_c^{(j)}$ is written as

$$\begin{aligned} S_c^{(j)} &= s(x_1 P_1^+, b_1) + s(x_2 P_2^+, b_2) + s((1-x_2)P_2^+, b_2) + s(x_3 P_3^-, b_3) \\ &+ s((1-x_3)P_3^-, b_3) - \frac{1}{\beta_1} \sum_{i=1}^{3} \ln \frac{\ln(t_j/\Lambda)}{-\ln(b_i \Lambda)} , \end{aligned} \tag{14}$$

with $b_3 = b_2$, $\beta_1 = (33 - 2n_f)/12$ and n_f the number of flavors. The scale t_j is chosen as $t_j = \max(AM_B, |B_j|M_B, 1/b_1, 1/b_2)$. The amplitude $\mathcal{M}_d$ is obtained from Fig. 2(d) accordingly. The amplitudes for the decay $\bar{B}^0 \to D^+\pi^-$ can be derived in a similar way. However, it is found that only the external W-emission contribution, the same as $\mathcal{M}_a$, is important.

The wave functions are chosen as [10],

$$\phi_\pi(x) = \frac{5\sqrt{6}}{2} f_\pi x(1-x)(1-2x)^2, \quad \phi_{B,D}(x) = \frac{N_{B,D}}{16\pi^2} \frac{x(1-x)^2}{M_{B,D}^2 + C_{B,D}(1-x)}, \tag{15}$$

with $f_\pi = 132$ MeV the pion decay constant. $N_B = 650.212$ GeV^3 and $C_B = -27.1051$ GeV^2 correspond to $f_B = 200$ MeV. N_D is determined by $f_D = 220$ MeV, and C_D is fixed by data for the decay $\bar{B}^0 \to D^+\pi^-$. All other parameters are referred to [10].

The experimental data of the branching ratios are $\mathcal{B}_0 = \mathcal{B}(\bar{B}^0 \to D^+\pi^-) = (3.08 \pm 0.85) \times 10^{-3}$ and $\mathcal{B}_- = \mathcal{B}(B^- \to D^0\pi^-) = (5.34 \pm 1.05) \times 10^{-3}$ [4]. Our predictions using the full Hamiltonian in Eq. (1), *ie* $c_1 = 1$ and $c_2 = 0$, are $\mathcal{B}_0 = 3.08 \times 10^{-3}$ and $\mathcal{B}_- = 5.10 \times 10^{-3}$. If the three-scale factorization formulas is employed, we obtain $\mathcal{B}_0 = 3.08 \times 10^{-3}$ and $\mathcal{B}_- = 5.00 \times 10^{-3}$, differing from the previous results only by 2%. We claim that the main theoretical uncertainty comes from higher corrections to the hard part. They are estimated to be 15-20% using the value of $\alpha_s(t)/\pi$ evaluated at the scale t below which half of contributions have been accumulated. A careful observation reveals that when the evolution of the Wilson coefficients is included, the amplitude $\mathcal{M}_b$, proportional to $a_2(t)$, becomes smaller, while $\mathcal{M}_c$, proportional to $c_1(t)/N_c$, becomes larger. The two changes cancel each other, and the total decay rate remains almost the same. $\mathcal{M}_d$ is less important because of pair cancellation. Our calculation indicates that the nonfactorizable contribution $\mathcal{M}_c$ is substantial, and the limit of the BSW factorization approximation. That is why the naive choice of $a_{1,2} = a_{1,2}(M_B)$ in the BSW model fails to explain the data.

Applying the three-scale factorization theorem to the modes $D \to K\pi$, we obtain the predictions $\mathcal{B}(\bar{D}^0 \to K^+\pi^-) = 4.05\%$ and $\mathcal{B}(D^- \to K^0\pi^-) = 2.67\%$, consistent with the data $(4.01 \pm 0.14)\%$ and $(2.74 \pm 0.29)\%$, respectively. With the running scale t reaching below the c quark mass, $\mathcal{M}_b$ becomes more negative and overcomes the positive contribution of $\mathcal{M}_c$. This explains the observed destructive interference of the external and internal W emissions absent in the B meson decays. Hence, the nonfactorizable diagrams play an important role in the explanation of the heavy meson decay data. It is clear that a PQCD formalism based on the original Hamiltonian without Wilson coefficients [10,13] can not account for this change of sign in the charm decays.

From the above analysis, we suggest that $c_{1,2}$ could be regarded as Wilson coefficients as they originally are, instead of as fitting parameters in the BSW model. That is, the controversy over the extraction of a_2/a_1 from the bottom and charm decays does not exist in our theory.

The scale dependence of our formalism can be tested by substituting $2t$ for t in the factorization formula. It is found that the prediction decreases by only 5%. In the conventional effective field theory, the substitution of M_b by $2M_b$ for the argument of $c_{1,2}$ results in a 10-20% difference [1]. Hence, the scale-setting ambiguity is moderated in the three-scale factorization theorems. Our formalism provides a more sophisticated choice of the scale, and is expected to give more definitive predictions, when it is applied to inclusive nonleptonic B meson decays. This subject will be discussed elsewhere.

This formalism has been extended to $B \to J/\psi K^{(*)}$ modes [14], and a thorough phenomenology investigation has been performed in the above work. It turns out that both R and R_L can be accommodated. Recently the irreducible soft corrections are also taken into account [15]. The non-factorizable contributions are clearly identified and are shown to be positive for bottom decays and negative for charm decays, consistent with experiment data.

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References

1. G. Buchalla, A. Buras and M. Lautenbacher, hep-ph/9512380.
2. M. Bauer, B. Stech and M. Wirbel, Z. Phys. C29, 637 (1985); 34, 103 (1987).
3. M. Neubert, V. Rieckert, B. Stech and Q.P. Xu in *Heavy Flavours*, edt. by A.J. Buras and M. Lindner (World Scientific, Singapore, 1992).
4. J.L. Rodriguez, hep-ph/9604011; CLEO Collaboration, M.S. Alam *et al.*, Phys. Rev. D**50**, 43 (1994).
5. H.Y. Cheng and B. Tseng, Phys. Rev. D**51**, 6259 (1995).
6. M. Gourdin, A.N. Kamal, Y.Y. Keum, and X.Y. Pham, Phys. Lett. B**333**, 507 (1994).
7. M. Gourdin, A.N. Kamal, and X.Y. Pham, Phys. Rev. Lett. **73**, 3355 (1994).

8. A.N. Kamal and T.N. Pham, Phys. Rev. D50, 395 (1994).
9. H-n. Li, Phys. Rev. D52, 3958 (1995); C.Y. Wu, T.W. Yeh, and H-n. Li, Phys. Rev. D53, 4982 (1996).
10. C.Y. Wu, T.W. Yeh, and H-n. Li, Phys. Rev. D55, 237 (1997).
11. See, for example, H. Georgi, Ann. Rev. Nucl. Part. Sci, 43, 209 (1993).
12. C.-H. V. Chang and H-n. Li, Phys. Rev. D55, 5577 (1997).
13. C.E. Carlson and J. Milana, Phys. Rev. D49, 5908 (1994).
14. T.W. Yeh and H.-n. Li, hep-ph/9701233, Phys. Rev. D56 (1997).
15. H.-n. Li and B. Tseng, hep-ph/9706441.