Gluon enhancements in charmed-meson decays

V. Barger, J. P. Leveille, and P. M. Stevenson
Physics Department, University of Wisconsin—Madison, Madison, Wisconsin 53706
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We address tests of nonspectator diagrams in charmed-meson decays. In particular, we explore the consequences of a model in which such diagrams are enhanced by the emission of a single gluon. For a lifetime ratio $\tau(D^+)/\tau(D^0)=5$, we predict (i) $\tau(F^+)/\tau(D^+)\simeq 0.6$, (ii) an enhanced $\Delta S=0$ rate for D^+, (iii) a substantial multipion branching fraction for the F^+, (iv) an $F^+\to\tau^+\nu$ branching fraction of order 10–20%, and (v) restrictive bounds on the $D^0\to\bar{K}^0\pi^0$ mode.

I. INTRODUCTION

Recent experimental results[1-3] on charm-particle lifetimes and semileptonic branching fractions require a revision of the usual theoretical model for weak decays. In the conventional model the heavy quark decays, with the light-quark constituents acting as spectators.[4] This predicts equal lifetimes and equal semileptonic branching fractions for all charmed particles, whereas preliminary data[1-3] indicate that the ratios $\tau(D^+)/\tau(D^0)$, $B(D^+\to eX)/B(D^0\to eX)$, and $\tau(D^+)/\tau(\Lambda_c^+)$ may be of order three or larger.

A reasonable explanation is that contributions from diagrams with interactions between initial quarks are larger than previously suspected.[5] For Λ_c^+ decay we have recently found that the W-exchange quark diagram is indeed comparable to the spectator diagrams.[6] For mesons it is usually argued that such contributions are very small due to the helicity suppression inherent in the $V-A$ decay of a pseudoscalar to a light-quark–antiquark pair.[4] However, it has been proposed recently that this suppression can be circumvented by an accompanying gluon, which allows the quarks to interact in a spin-1 state.[7,8] This gluon can be regarded either as being radiated perturbatively by the initial quarks,[7] or as a constituent in the initial meson wave function.[8] Quantitative estimates of nonspectator transitions in a one-gluon model give rates comparable to those of spectator diagrams, but are subject to considerable theoretical uncertainty. It is also possible that multiple-gluon emission may be important.

In this paper we examine tests of the gluon mechanism which are independent of any absolute rate calculations. These tests involve (i) a relation among D^+, D^0, and F^+ lifetimes, (ii) inclusive Cabibbo-suppressed ($\Delta S=0$) decays, (iii) final states in F^+ decay without s or $\bar{s}$ quarks, and (iv) $D\to\bar{K}$ two-body decay modes. Predictions for these quantities are made in terms of the $\tau(D^+)/\tau(D^0)$ lifetime ratio. Our point is that if the D^0-D^+ lifetime difference is due to nonspectator interactions, a definite pattern of modifications to the conventional-model predictions[4] is implied. In some cases the usual expectation is hardly modified (which in itself is an important prediction), while in other cases significantly different results emerge. Most of our results depend only on the presence of nonspectator interactions at a significant, but not necessarily dominant, level. However, in certain cases we examine predictions specific to the one-gluon model.

II. SPECTATOR AND NONSPECTATOR CONTRIBUTIONS

Our analysis is based on the conventional effective weak-interaction Lagrangian[9]

$$\mathcal{L}_{\rm eff}=(G_F/\sqrt{2})[f_1(\bar{u}d')(\bar{s}'c)+f_2(\bar{s}'d')(\bar{u}c)]\,, \qquad (1)$$

where $(q\bar{Q})$ denotes a color-singlet $V-A$ current and $f_1\equiv(f_++f_-)/2$, $f_2\equiv(f_+-f_-)/2$. The short-distance enhancement factors f_+, f_- are due to hard-gluon renormalization effects (in the absence of strong interactions $f_+=f_-=1$). They are given by[9]

$$f_-=[\alpha_s(m_c^2)/\alpha_s(m_W^2)]^\gamma\,,\quad f_+=(f_-)^{-1/2}\,,$$

where $\gamma=12/(33-2F)$, with F the effective number of flavors; m_c and m_W are the charm-quark and W-boson masses, respectively. Numerically[10]

$$f_1\simeq 1.39\,,\quad f_2\simeq -0.70\,,$$

where we have taken $m_c=1.5$ GeV, $m_W=84$ GeV, $F=4$ and used $\alpha_s(m^2)=\pi\gamma/\ln(m^2/\Lambda^2)$ with $\Lambda=0.5$ GeV.

In Eq. (1), $d'=d\cos\theta+s\sin\theta$ and $s'=s\cos\theta-d\sin\theta$, where θ is the Cabibbo angle; we neglect small differences between Cabibbo angles in the charm and noncharm sectors that arise in the six-quark model[11] and take $\sin^2\theta\simeq 0.05$.

We parametrize the spectator contributions (see Fig. 1) as

$$\Gamma(c\to su\bar{d};\, su\bar{s};\, du\bar{d};\, du\bar{s})=(C^4;\, C^2S^2;\, C^2S^2;\, S^4)N,$$

$$\Gamma(c\to sl\nu;\, dl\nu)=(C^2;\, S^2)\tfrac{1}{2}L\,, \qquad (2)$$

where $C\equiv\cos\theta$, $S\equiv\sin\theta$, and $l=e$ or μ. N and L are

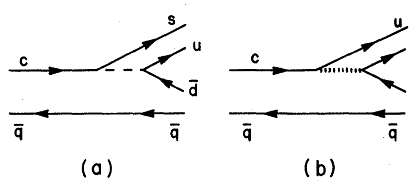

FIG. 1. A typical spectator contribution: the Cabibbo-favored $c \to su\bar{d}$ decay. The diagrams (a) and (b) represent the first and second terms of the effective Lagrangian, Eq. (1). (The diagrams do not show gluon radiative corrections which are implicitly included in our analysis.)

reduced rates for c quark decay with gluon corrections [including the short-distance factor[12] $(f_-^2+2f_+^2)/3$ in N as well as radiative gluon corrections in both L (Ref. 13) and N].

The contributions of nonspectator diagrams (see Fig. 2) can be parametrized as

$$\Gamma_g(D^0(c\bar{u}) \to s\bar{d}; d\bar{d}; s\bar{s}; d\bar{s}) = (C^4; C^2S^2; C^2S^2; S^4)f_1^{\,2}G(D^0),$$
$$\Gamma_g(D^+(c\bar{d}) \to u\bar{d}; u\bar{s}) = (C^2S^2, S^4)f_2^{\,2}G(D^+)\,, \qquad (3)$$
$$\Gamma_g(F^+(c\bar{s}) \to u\bar{d}, u\bar{s}) = (C^4, C^2S^2)f_2^{\,2}G(F^+)\,.$$

These partial widths represent exchange or annihilation contributions including the possible radiation of gluons which removes the helicity suppression. At this stage the G parameters are unrelated. In the one-gluon model,[7] the ratios of these parameters are fixed, so we have chosen to write Eq. (3) with explicit f_1, f_2 factors for later convenience. We have assumed that the gluon radiation enhances only nonleptonic decay modes. This is obviously true in the one-gluon model because of color conservation, but may not be true if multiple-gluon emission is important.

We can neglect purely leptonic decays which are Cabibbo-suppressed and/or helicity-suppressed. However, we need to include the $F^+ \to \tau^+\nu_\tau$ contribution which may be significant. This rate is given by

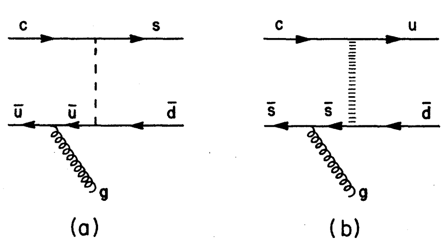

FIG. 2. Typical nonspectator contributions, enhanced by gluon emission: (a) the $c\bar{u} \to s\bar{d}$ exchange contribution to D^0 decay, and (b) the $c\bar{s} \to u\bar{d}$ contribution to F^+ decay. Note that in the one-gluon model these involve different operators. Multiple-gluon emission is not shown but is also possible.

$$\Gamma_{F\to\tau\nu} = C^2G_F^{\,2}f_F^{\,2}m_F m_\tau^{\,2}(1-m_\tau^{\,2}/m_F^{\,2})^2/(8\pi), \qquad (4)$$

where $f_F^{\,2} = 12|\psi(0)|_F^{\,2}/m_F$ in the nonrelativistic approximation. Theoretical estimates of the decay constant f_F vary widely. To make numerical estimates we assume $f_F = f_D$ and take $f_D \sim 430$ MeV, as determined by Suzuki[7] from experimental $D^{*+,0}$ and $D^{+,0}$ electromagnetic mass splitting. This value of f_F is about a factor of 3 larger than f_K or f_π, in qualitative accord with the notion that nonspectator interactions are important. This gives $\Gamma_{F\to\tau\nu} \sim 4.7\times10^{11}\ \mathrm{s}^{-1}$ which implies a substantial $F^+ \to \tau^+\nu_\tau$ branching fraction, in the range 10–20%, for an F^+ lifetime of 2×10^{-13} to 4×10^{-13} s.

From Eqs. (2) and (3), and including the $F\to\tau\nu$ contribution, we find the total decay rates

$$\Gamma(D^0) = L+N+f_1^{\,2}G(D^0)\,,$$
$$\Gamma(D^+) = L+N+f_2^{\,2}S^2G(D^+)\,, \qquad (5)$$
$$\Gamma(F^+) = L+N+f_2^{\,2}C^2G(F^+) + \Gamma_{F\to\tau\nu}\,.$$

III. TESTS OF NONSPECTATOR INTERACTIONS

A. Lifetimes in the one-gluon model

If nonspectator interactions proceed dominantly through the emission of a single gluon, as in Ref. 7, then we can relate the reduced rates G appearing in Eq. (3). The initial quarks, having emitted a single color-octet gluon, can interact only through that term of Eq. (1) which represents an exchange (and not an annihilation) diagram. This leads to the specific pattern of short-distance enhancement factors explicitly displayed in Eq. (3).[14] The remaining differences between the G's in the one-gluon model are due to SU(3)-breaking effects:

$$G(D^0) : G(D^+) : G(F^+) = 1:1:r\,, \qquad (6)$$

where

$$r = \frac{|\psi(0)|_F^{\,2}}{|\psi(0)|_D^{\,2}}\left(\frac{m_F}{m_D}\right)^4\left(\frac{m_u}{m_s}\right)^2 \simeq 0.55\,.$$

This numerical estimate of r is based on $m_u = 336$ MeV, $m_s = 540$ MeV, and on the assumption that the wave function at the origin is comparable for D and F. In this model the nonspectator contributions [see Eq. (3)] to the total inclusive rate are in the ratios

$$\Gamma_g(D^0):\Gamma_g(D^+):\Gamma_g(F^+) = f_1^{\,2}:f_2^{\,2}S^2:f_2^{\,2}C^2r \qquad (7a)$$
$$\simeq 1.9:0.03:0.26\,. \qquad (7b)$$

From Eqs. (5) and (6) we find, in the one-gluon model, the lifetime relation

$$\frac{\tau(D^+)}{\tau(F^+)} = 1+\Delta(rC^2-S^2)+\tau(D^+)\Gamma_{F\to\tau\nu}\,, \qquad (8a)$$

where

$$\Delta \equiv \frac{f_2^{\,2}}{(f_1^{\,2}-S^2 f_2^{\,2})}\left(\frac{\tau(D^+)}{\tau(D^0)}-1\right). \tag{8b}$$

For example, if $\tau(D^+)/\tau(D^0)=5$, then we get $\Delta \simeq 1.0$ and

$$\tau(F^+) \simeq 2\tau(D^+)/[3+2\tau(D^+)\Gamma_{F\to\tau\nu}]$$
$$\simeq (0.5\text{–}0.6)\tau(D^+)$$

for $\tau(D^+)$ in the range 8×10^{-13} to 2×10^{-13} s. Figure 3 shows the variation of $\tau(F^+)/\tau(D^+)$ as a function of $\tau(D^+)/\tau(D^0)$ for various values of $B(F\to\tau\nu)$. Qualitatively this model predicts an F^+ lifetime that is always somewhere between the D^0 and D^+ lifetimes.

In this model the semileptonic widths are unaffected by exchange interactions, so the inclusive semileptonic branching fractions into electrons are in the same ratio as the lifetimes:

$$B_e(D^0) : B_e(D^+) : B_e(F^+) = \tau(D^0) : \tau(D^+) : \tau(F^+). \tag{9}$$

However, if multiple-gluon emission is important, then the semileptonic decays of the F^+ can be enhanced by the annihilation interaction $c\bar{s} \to e^+\nu_e$ + gluons, where the gluons form a color-singlet state which metamorphoses into hadrons. In this case $B_e(F^+)$ would be larger than that expected from Eq. (9).

B. Inclusive $\Delta S=0$ decays

We now turn to inclusive Cabibbo-suppressed $\Delta S=0$ decays; i.e., transitions in which the final state has the same net strangeness as the initial meson. In the pure spectator picture the D^0, D^+, and F^+ branching fractions are each given by

$$B(\Delta S=0) = S^2\left|\frac{2C^2N+L}{N+L}\right| \tag{10a}$$
$$= S^2[2C^2 - 2(1-2S^2)B_e(D^+)] \simeq 8\%, \tag{10b}$$

where the numerical value is based on[1] $B_e(D^+)$

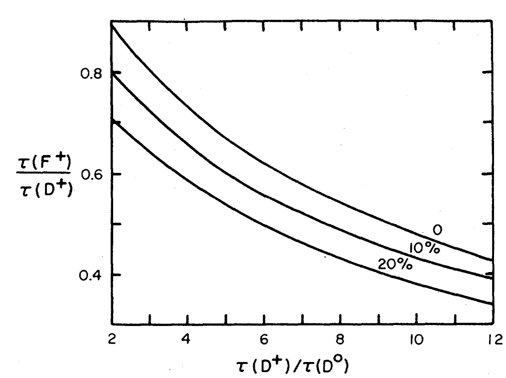

FIG. 3. The lifetime ratio $\tau(F^+)/\tau(D^+)$ as a function of $\tau(D^+)/\tau(D^0)$ in the one-gluon model [Eq. (8)] for various values of $B(F\to\tau\nu)$.

$\simeq 0.16$. Including exchange interactions we find that $\Delta S=0$ branching ratios for D^0 and F^+ are essentially unchanged; the former is very slightly increased, the latter slightly decreased. Specifically, we obtain from Eqs. (3) and (5)

$$B(D^0;\Delta S=0) = S^2[2C^2 - 2(1-2S^2)B_e(D^0)], \tag{11}$$

which differs from the spectator result only in that $B_e(D^0) < B_e(D^+)$. Since the right-hand side of Eq. (11) is less than $2C^2S^2$, we conclude that $B(D^0;\Delta S=0) = 8\text{–}10\%$. The corresponding result for the F^+ is

$$B(F^+;\Delta S=0) = S^2\{2C^2(1-B_{F\to\tau\nu}) - (1-2S^2)[C^2\Delta_F + 2B_e(F^+)]\}, \tag{12}$$

where

$$\Delta_F = f_2^{\,2}G(F^+)/\Gamma(F^+).$$

Whatever the value of Δ_F in the entire range $0 < \Delta_F < (1-B_{F\to\tau\nu})/C^2$, the right-hand side of Eq. (12) is less than the spectator result, but greater than $(1-E_{F\to\tau\nu})S^2$. Hence, $B(F^+;\Delta S=0)$ must lie in the range 4–8%.

However, for the D^+ the situation is very different. Here nonspectator interactions contribute only to Cabibbo-suppressed decays, so the $\Delta S=0$ branching fraction can be noticeably enhanced. It is given by

$$B(D^+;\Delta S=0) = S^2[2C^2 + (1-2S^2)(C^2\Delta_+ - 2B_e(D^+))], \tag{13}$$

where

$$\Delta_+ = f_2^{\,2}G(D^+)/\Gamma(D^+) \quad (0<\Delta_+<1/S^2).$$

So far our analysis has been general, but if we now specialize to the one-gluon-model ratios of Eq. (6), we can identify Δ_+ with the Δ defined in Eq. (8b). This leads to a prediction for $B(D^+;\Delta S=0)$ in terms of $\tau(D^+)/\tau(D^0)$. For example, if $\tau(D^+)/\tau(D^0) = 5$, we obtain $B(D^+;\Delta S=0) \simeq 12\%$. There is some indication from measurements[1] of the inclusive branching fractions $B(D\to\bar{K}X)$ that the $\Delta S=0$ branching fraction for the D^+ is indeed larger than that for D_0.

C. F^+ multipion decay modes

In F^+ decays the spectator transition $c\bar{s}\to s\bar{s}ud$ will give final states with either overt or hidden strangeness (e.g., $F^+\to K^-K^+X^+$ or $F^+\to\eta X^+$). Nonresonant multipion final states are not expected at a significant level. The situation is markedly different with the nonspectator transition $c\bar{s}\to u\bar{d}$ [+ gluon(s)], where states of three or more pions are anticipated. The branching fraction for this transition with no s or $\bar{s}$ quarks originating from the weak interaction can be estimated in the one-gluon model. We obtain

$$B(F^+ \to \text{no } s \text{ or } \bar{s}) = C^4 r \Delta \tau(F^+)/\tau(D^+) \simeq (25-30)\% , \quad (14)$$

where Δ and $\tau(F^+)/\tau(D^+)$ are given in Eq. (8). The numerical value is based on $\tau(D^+)/\tau(D^0) = 5$ and $\tau(F^+)/\tau(D^+) \simeq 0.5-0.6$. Since $\bar{u}u$ or $\bar{d}d$ pair creation is favored over $\bar{s}s$, the $u\bar{d}g$ final state evolves mainly into multipion or η plus pions modes. Although we are unable to separate the relative proportion, phase-space considerations suggest that the multipion states will predominate. In this connection it is interesting to note that an emulsion event $F^- \to \pi^-\pi^+\pi^-\pi^0$ has been observed as one of six charged charm-meson events.[2]

D. $D \to \bar{K}\pi$ decays

In the spectator model[15] the transition $D^0 \to \bar{K}^0\pi^0$ is color-suppressed; this prediction is at variance[16] with recent experimental measurements.[1] The $D^0 \to K^-\pi^+$, $D^0 \to \bar{K}^0\pi^0$, and $D^+ \to \bar{K}^0\pi^+$ amplitudes are related in the spectator model as[10]

$$A_s^{-+} = a, \quad A_s^{00} = \frac{a}{\sqrt{2}} \frac{\epsilon\chi_-}{\chi_+}, \quad A_s^{0+} = a\left(1 + \frac{\epsilon\chi_-}{\chi_+}\right), \quad (15)$$

where $\chi_+ = (2f_+ + f_-)/3 \simeq 1.16$ and $\chi_- = (2f_+ - f_-)/3 \simeq -0.24$. In the SU(3) limit without form factors $\epsilon = 1$. The $\bar{K}^0\pi^0$ amplitude is particularly sensitive to the value of this factor, given by

$$\epsilon = (f_K/f_\pi)(1 - K^2/D^2)^{-1}(1 - K^2/D_S{}^2) \simeq 1.5 , \quad (16)$$

where $f_K/f_\pi = 1.28$ is the ratio of leptonic decay constants, $K = 0.495$ GeV, $D = 1.867$ GeV, and D_S is the 0^+ charmed-meson mass. With the above estimate for ϵ the spectator-model prediction is $B(\bar{K}^0\pi^0) \simeq 0.05\ B(K^-\pi^+)$ whereas comparable D^0 branching fractions are observed[1] for these modes ($B^{00} = 2.0 \pm 0.9\%$, $B^{-+} = 2.8 \pm 0.6\%$).

Again the exchange diagrams offer a way out of the discrepancy. This is because the nonspectator interactions give amplitudes in which there is no color-suppression factor between the $\bar{K}^0\pi^0$ and $K^-\pi^+$ modes:

$$A_g^{-+} = b, \quad A_g^{00} = -b/\sqrt{2}, \quad A_g^{0+} = 0 . \quad (17)$$

In the extreme limit where the nonspectator contributions completely dominate, $B^{00} = 0.5 B^{-+}$ (Refs. 5 and 8); with the present experimental uncertainties this is not yet excluded. Adding the spectator and exchange diagram contributions, we are left with an amplitude sum rule

$$A^{00} = [A^{0+} - A^{-+}]/\sqrt{2} . \quad (18)$$

In fact, this sum rule follows directly from the $I = 1$, $I_3 = 1$ transformation properties of the effective Lagrangian.[17] Equation (18) implies the following triangle inequality on the $\bar{K}^0\pi^0$ branching fraction

$$\tfrac{1}{2}\left[(B^{-+})^{1/2} - \left(B^{0+}\frac{\tau^0}{\tau^+}\right)^{1/2}\right]^2 \leq B^{00} \leq \tfrac{1}{2}\left[(B^{-+})^{1/2} + \left(B^{0+}\frac{\tau^0}{\tau^+}\right)^{1/2}\right]^2 , \quad (19)$$

where τ^0, τ^+ denote the D^0, D^+ lifetimes. Using the experimental values[1] $B^{-+} = 2.8 \pm 0.6\%$ and $B^{0+} = 2.1 \pm 0.4\%$, and taking $\tau^+/\tau^0 = 5$, we find the numerical bounds

$$0.5 \pm 0.2\% \leq B^{00} \leq 2.7 \pm 0.5\% .$$

These bracket the measured value[1] $B^{00} = 2.0 \pm 0.9\%$. The bounds become more restrictive with decreasing τ^0/τ^+.

The inclusion of nonspectator contributions does not change the usual $\tan^2\theta_C$ predictions for the ratios $(D \to \pi\pi)/(D \to \bar{K}\pi)$ and $(D \to \bar{K}K)/(D \to \bar{K}\pi)$. The observed discrepancies from these predictions can, however, be reasonably explained in terms of six-quark-model mixing angles and SU(3) breaking.[10]

E. Distintictive decay modes

Some special decay modes exist in which the so-called spectator quark is absent in the final state. These modes are of special interest since they can occur only through the nonspectator interaction. An example is the transition $c\bar{u} \to s\bar{d}g$ with an $s\bar{s}$ pair created by the gluon. This would give rise to $D^0 \to \bar{K}^0 K^0 \bar{K}^0$ or $D^0 \to \bar{K}^0\phi$ decays. A more exotic example is $F^+ \to p\bar{n}$; although certain to be rare, such a decay has a spectacular signature.

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[1]G. Goldhaber, in *Proceedings of the Topical Workshop on the Production of New Particles in Super High Energy Collisions*, edited by V. Barger and F. Halzen (University of Wisconsin, Madison, 1979).

[2]J. Prentice, in *Proceedings of the Topical Workshop on the Production of New Particles in Super High Energy Collisions*, edited by V. Barger and F. Halzen (University of Wisconsin, Madison, 1979).

[3]J. Kirkby, in *Proceedings of the 1979 International Symposium on Lepton and Photon Interactions at High Energies, Fermilab*, edited by T. B. W. Kirk and H. D. I. Abarbanel (Fermilab, Batavia, Illinois, 1980).

[4]See, for example, M. K. Gaillard, in *Weak Interactions—Present and Future*, proceedings of SLAC Summer Institute on Particle Physics, 1978, edited by M. C. Zipf (SLAC, Stanford, 1978); J. Ellis, M. K. Gaillard, and D. V. Nanopoulos, Nucl. Phys. B100, 313 (1975); L. Maiani and T. Walsh, in *Proceedings of International Symposium on Lepton and Photon Interactions at High Energies, Hamburg, 1977*, edited by F. Gutbrod (DESY, Hamburg, 1977); B. W. Lee, M. K. Gaillard, and J. Rosner, Rev. Mod. Phys. 47, 277 (1975).

[5]Isospin tests for the dominance of nonspectator diagrams in D^0 decay have been proposed by S. P. Rosen, Phys. Rev. Lett. 44, 4 (1980); Phys. Lett. 89B, 246 (1980); Phys. Rev. D 21, 2631 (1980).

[6]V. Barger, J. P. Leveille, and P. M. Stevenson, Phys. Rev. Lett. 44, 226 (1980).

[7]M. Bander, D. Silverman, and A. Soni, Phys. Rev. Lett. 44, 7 (1980); M. Suzuki (unpublished). We calculate a $D^0 \to s\bar{d}g$ decay rate in agreement with Suzuki's result, which is a factor of 2 smaller than that of Bander *et al.*

[8]H. Fritzsch and P. Minkowski, Phys. Lett. 90B, 455 (1980).

[9]B. W. Lee and M. K. Gaillard, Phys. Rev. Lett. 33, 108 (1974); G. Altarelli and L. Maiani, Phys. Lett. 52B, 351 (1974).

[10]V. Barger and S. Pakvasa, Phys. Rev. Lett. 43, 812 (1979).

[11]M. Kobayashi and K. Maskawa, Prog. Theor. Phys. 49, 652 (1973); V. Barger, W. F. Long, and S. Pakvasa, Phys. Rev. Lett. 42, 1585 (1979); R. E. Shrock, S. Treiman, and L.-L. Wang, *ibid.* 42, 1589 (1979).

[12]G. Altarelli, N. Cabibbo, and L. Maiani, Phys. Rev. Lett. 35, 635 (1975).

[13]M. Suzuki, Nucl. Phys. B145, 420 (1978); N. Cabibbo and L. Maiani, Phys. Lett. 79B, 109 (1978).

[14]Here we are assuming that single-gluon emission does not change the short-distance factors of Eq. (1).

[15]D. Fakirov and B. Stech, Nucl. Phys. B133, 315 (1978); N. Cabibbo and L. Maiani, Phys. Lett. 73B, 418 (1978).

[16]This problem has been discussed by N. Deshpande, M. Gronau, and D. Sutherland, Phys. Lett. 90B, 431 (1980); J. F. Donoghue and B. R. Holstein, Phys. Rev. D 21, 1334 (1980); H. Fritzsch, Phys. Lett. 86B, 343 (1979).

[17]J. F. Donoghue and B. R. Holstein, Phys. Rev. D 12, 1454 (1975); R. Kingsley, S. Treiman, F. Wilczek, and A. Zee, *ibid.* 11, 1919 (1975); L.-L. Wang and F. Wilczek, Phys. Rev. Lett. 43, 816 (1979).