not.

We thank A. Luther for suggesting this study and P. Moussa for an illuminating discussion. Laboratoire de Physique des Solides is a Laboratoire associé au Centre National de la Recherche Scientifique.

[a]Permanent address: Institute of Physics of the University of Zagreb, Croatia, Yugoslavia.

[1]C. N. Yang and T. D. Lee, Phys. Rev. 87, 404 (1952).

[2]T. D. Lee and C. N. Yang, Phys. Rev. 87, 410 (1952).

[3]M. Suzuki, Prog. Theor. Phys. 38, 1225 (1967).

[4]J. D. Bessis, J. M. Drouffe, and P. Moussa, J. Phys. A 9, 2105 (1976).

[5]D. A. Kurtze and M. E. Fisher, J. Stat. Phys. 19, 205 (1978).

[6]G. A. Baker, M. E. Fisher, and P. Moussa, Phys. Rev. Lett. 42, 615 (1979).

[7]P. J. Kortman and R. B. Griffiths, Phys. Rev. Lett. 27, 1439 (1971).

[8]M. E. Fisher, Phys. Rev. Lett. 40, 1610 (1978).

[9]R. Abe, Prog. Theor. Phys. 38, 72 (1967); G. A. Baker, Phys. Rev. Lett. 20, 990 (1968); D. S. Gaunt and G. A. Baker, Phys. Rev. B 1, 1184 (1970); G. A. Baker and P. Moussa, J. Appl. Phys. 49, 1360 (1978); M. Barnsley, D. Bessis, and P. Moussa, J. Math. Phys. 535 (1979).

[10]There is an exact correspondence [M. Suzuki, Prog. Theor. Phys. 56, 1454 (1976)] between the ground-state energy of the ITF in D dimensions and the free energy of classical Ising model in $D+1$ dimensions, with an interaction going to zero in D dimensions and to infinity in $D+1$. This leads to identical critical behaviors as shown from the exact solution for $D=1$ in zero longitudinal field [P. Pfeuty, Ann. Phys. (N.Y.) 57, (1970); B. M. McCoy and T. T. Wu, *The two dimensional Ising Model* (Harvard University Press, Cambridge, Mass. 1973)].

[11]S. D. Drell, M. Weinstein, and S. Yankielowicz, Phys. Rev. D 14, 487 (1976).

[12]R. Jullien, P. Pfeuty, J. N. Fields, and S. Doniach, Phys. Rev. B 18, 3568 (1978).

[13]This method is equivalent to a renormalization-group study of the maximum eigenvalues of the transfer matrix of the classical Ising model, which supposes that the thermodynamic limit is partially taken before the method is applied. This real-space renormalization method is thus supposed to be more correct than other real-space renormalization methods from recent arguments by R. Griffiths and P. A. Pearce, Phys. Rev. Lett. 41, 917 (1978).

[14]K. Uzelac, R. Jullien, and P. Pfeuty, unpublished.

[15]J. L. Cardy, Nucl. Phys. B115, 141 (1976).

[16]K. A. Penson, R. Jullien, and P. Pfeuty, Phys. Rev. B 19, 4653 (1979).

Euclidean Group as a Dynamical Symmetry of Surface Fluctuations: The Planar Interface and Critical Behavior

D. J. Wallace[a]
Physics Department, The University, Southampton SO9 5NH, England

and

R. K. P. Zia
Physics Department, Virginia Polytechnic Institute & State University, Blacksburg, Virginia 24061

(Received 26 June 1979)

The statistical mechanics of the interface between two discrete thermodynamic phases is studied in terms of a field f which describes the deviation of the interface from planar. Exploiting the dynamical Euclidean invariance of the Hamiltonian of the field f, we construct ϵ expansions for Ising-like critical behavior in $1+\epsilon$ dimensions, with a critical temperature of order ϵ. Scaling functions for the interface profile and width in a pinning potential are obtained.

Recent interest in the interface between two discrete thermodynamic phases has focused on two main characteristics. One is the divergence of the interfacial width with the correlation length[1] as the temperature T is raised to the critical temperature T_c; this behavior can be described for example by ϵ expansions in $4-\epsilon$ dimensions[2] starting from the conventional Ginzburg-Landau-Wilson Hamiltonian. The second is the collective wandering of the interface as a whole,[3] which arises from the inability of surface tension to maintain a planar interface. In particular, the lowest-order calculations[4] in the capillary-wave approach imply that in the absence of a stabilizing potential, such as gravity, the mean square width of the interface diverges for all $T<T_c$ in the thermodynamic limit for dimension of space

$d \leq 3$. This statement must be modified for lattice Ising models where for $d=3$ the finite lattice spacing succeeds in limiting the wandering of the interface to a finite region for all temperatures less than a roughening temperature[5] T_R ($<T_c$). Above T_R the wandering is unbounded, as in the continuum models using capillary waves.

The capillary-wave approach is basically a low-temperature description of the interface. It incorporates into the statistical mechanics apparently only the low-energy configurations—the large-distance deviations from planar of an essentially sharp interface, with surface tension as the restoring potential. In this paper we describe how the renormalization group enables one to control the short-distance behavior in the interface in such a model. Hence we can construct, from this model of a sharp interface, a renormalized interface, whose local width diverges as the correlation length when T approaches T_c. In order to control such calculations in a perturbation theory in T it is necessary that T_c be small in some sense. In fact, we show that perturbative calculations are limited to an ϵ expansion in $1+\epsilon$ dimensions,[6] with a T_c of order ϵ. The fact that the capillary waves transform as a nonlinear realization of the Euclidean group of d dimensions[7] plays an important role in these calculations, and conversely these calculations establish that the capillary waves can be interpreted as the Goldstone modes whose fluctuations lower to zero the critical temperature as $d \to 1^+$. They therefore play a role for systems with discrete internal symmetry which is equivalent to that played by spin waves in a continuous internal symmetry.[8]

Our starting point is the reduced Hamiltonian for a field f which describes the deviation from planar of an essentially sharp interface (see for example Bull, Lovett, and Stillinger[4]),

$$\mathcal{H}(f) = T^{-1} \int d^{d-1}x \{[1+(\nabla f)^2]^{1/2} + \tfrac{1}{2} m^2 f^2\}. \tag{1}$$

The first term is the surface area of the interface; the coefficient $1/T$ represents $\sigma/k_B T$ where σ is the interfacial energy per unit area at zero temperature. The mass term in $\mathcal{H}$ represents a pinning potential such as gravity, or step-function magnetic field $H(z) = H\epsilon(z)$. Starting from the Ginzburg-Landau-Wilson description of the interface,[2] it is straightforward to see that m^2 represents $HM\xi^{-1}$ for such a magnetic field, where M is spontaneous magnetization and ξ is the correlation length. The loop expansion for the partition function $Z = \int Df e^{-H}$ and correlation functions is an expansion in powers of T; it is obtained by expansion of the square root to give a free propagator $T(q^2+m^2)^{-1}$ and vertices $-\frac{1}{8}T^{-1}[(\nabla f)^2]^2$, etc.

The ultraviolet and infrared properties of the model are controlled by the dimension of T. Since $\mathcal{H}$ is dimensionless, $T = \kappa^{-(d-1)}$ where κ is an inverse length, or momentum. We therefore expect that for $d>1$ the loop integrals which accompany powers of T in perturbation theory do not change the low-momentum behavior of the free propagator and vertices. This property ensures that the description of the large-distance interface wandering by the evaluation of $\langle f^2 \rangle$ *at lowest order in perturbation theory* is valid. Correspondingly $\mathcal{H}$ is naively nonrenormalizable for $d>1$, i.e., its high-momentum behavior is not controllable in straightforward perturbation in T. Now, it is intuitively clear that if we are to have an interface which is "fuzzy" over a length scale of the correlation length ξ, then this will be obtained only by incorporating fluctuations of a sharp interface over *all* momenta $q > \xi^{-1}$. The description of critical behavior in the model (1) is therefore equivalent to controlling the high-momentum behavior of the model.

This problem can be solved perturbatively in $1+\epsilon$ dimensions using renormalization-group techniques. Our approach follows closely the calculations of Brézin and co-workers[8] for Heisenberg models in $2+\epsilon$ dimensions. Using dimensional regularization,[9] we have calculated the two- and four-point vertex functions up to and including two loops.[10] The renormalization required to remove poles in $(d-1)$ are a coupling-constant renormalization

$$T \equiv \kappa^{-(d-1)} t_0 = \kappa^{-(d-1)}\{t + t^2(d-1)^{-1} + t^3[(d-1)^{-2} + \tfrac{1}{4}(d-1)^{-1}] + O(t^4)\},$$

and a mass renormalization $m^2/T = m_R^2 \kappa^{d-1}/t$. Here t is the dimensionless renormalized coupling constant; it is defined to absorb a factor

$$I = -\int d^{d-1}q\,(2\pi)^{-d+1} q^2 (q^2+1)^{-1},$$

which can be extracted from each loop integral. No wave-function renormalization is required. We have shown that this is true to all orders in perturbation theory as a consequence of the Ward identities resulting from the nonlinear Euclidean-group transformations on f in (1). This can also be understood because f represents a length, the deviation from planar, and hence develops no anomalous dimension.

The resulting renormalization-group equation for all vertex functions Γ_R is

$$\left(\kappa\frac{\partial}{\partial\kappa}+\beta(t)\frac{\partial}{\partial t}+\gamma_1(t)m_R^2\frac{\partial}{\partial m_R^2}\right)\Gamma_R=0, \tag{2}$$

where $\beta(t)=(d-1)t-t^2-\frac{1}{2}t^3+\ldots$ and $\gamma_1(t)=-(d-1)+(\beta/t)$. The fixed points of $\dot{t}=\beta(t)$ are (i) $t=0$, which is infrared stable $[\beta'(0)>0]$ and controls the mean-field low-momentum behavior of $\mathcal{H}$, and (ii) $t_c=(d-1)-\frac{1}{2}(d-1)^2+O((d-1)^3)$, which is ultraviolet stable $[\beta'(t_c)=-(d-1)-\frac{1}{2}(d-1)^2+O((d-1)^3)]$, and is the effective coupling for the high-momentum behavior. The interpretation of t_c as the critical temperature is supported by the following arguments.

(i) It is convenient to introduce a physical length scale, the correlation length $\xi\equiv\kappa^{-1}\zeta(t)$, which is independent of the pinning potential m_R, and invariant under a renormalization-group transformation, i.e., $[\kappa\,\partial/\partial\kappa+\beta(t)\,\partial/\partial t]\xi=0$. The solution of this equation is

$$\xi=\kappa^{-1}t^{1/(d-1)}\exp\int_0^t dt'[1/\beta(t')-t'^{-1}(d-1)^{-1}].$$

For small t, we see $\xi=\kappa^{-1}t^{1/(d-1)}[1+O(t)]$; this correlation length can be identified with the phase coherence length introduced for Heisenberg and Ising systems.[11] As t is increased to t_c, ξ diverges according to $\xi\propto(t_c-t)^{-\nu}$ where $\nu=-1/\beta'(t_c)=(d-1)^{-1}-\frac{1}{2}+O(d-1)$. This agrees with the lowest-order results from lattice models.[6]

(ii) The interfacial free energy per unit area contains a contribution $1/T$ from the Hamiltonian (1). As $t\to t_c$, the bare temperature $T\to\infty$ and hence the interfacial energy vanishes. The way in which it vanishes is obtained by solving $dt_0/dt=(d-1)t_0/\beta(t)$ for t_0, leading to

$$1/T=\kappa^{d-1}\exp[-(d-1)\int^t dt'/\beta(t')]\propto\xi^{-(d-1)}.$$

Thus the interfacial energy vanishes according to the expected scaling law (see, e.g., Widom[1] and also Ref. 12).

(iii) The solution of the renormalization-group equation (2) for the mean square width of the interface $\langle f^2\rangle$ has the scaling form $\langle f^2\rangle=\xi^2P_2(y)$ where $y=m_R^2\xi^{d+1}\kappa^{d-1}t^{-1}$ is the scaling variable associated with the pinning potential m_R^2. We see that this is compatible with the identification $m^2\sim HM\xi^{-1}$, i.e., HM scales like ξ^{-d} as expected. The loop expansion gives the function $P_2(y)$ as a power series in $y^{(d-1)/2}$. The explicit form of the first two terms is $P_2(y)=y^{-(3-d)/2}[1-y^{(d-1)/2}+O(y^{d-1})]$; we have discarded numerical coefficients (of order 1) which are only physically significant when $O(y^{d-1})$ is calculated. The behavior of $P_2(y)$ for y large, i.e., large pinning potential, can be obtained from a simple model similar to that used by Weeks[5] to study roughening. If we take cells of base area $\sim\xi^{d-1}$, the interface height over a single cell may be modeled by a single variable f, with a probability distribution proportional to $\exp(-m^2\xi^{d-1}f^2/T)$; we neglect intercell interactions. This model for large y gives $\langle f^2\rangle=Tm^{-2}\xi^{-(d-1)}\propto\xi^2y^{-1}$. Thus we see that for $y\ll1$ the divergence of $P_2(y)$ describes the divergence of interface width due to wandering for $d<3$, for $y\sim1$, $\langle f^2\rangle\sim\xi^2$, and finally $\langle f^2\rangle$ tends to zero for large pinning potential ($y\to\infty$). The $\ln y$ term expected in three dimensions[4] is also indicated.

(iv) Loop expansions can be similarly obtained for the expectation values $\langle f^{2n}\rangle=\xi^{2n}P_{2n}(y)$. These functions can be regarded as the moments $P_{2n}=\int_{-\infty}^{\infty}P(z)z^{2n}dz$ of a probability distribution $P(z)$ for the position of the center of the interface. Given the explicit forms of the first two orders in $\langle f^{2n}\rangle$, one can perform an inverse Mellin transform to obtain the probability distribution $P(z)$. To first order, a Gaussian distribution is obtained, a result reported by Huang and Webb[13] for the model of Buff, Lovett, and Stillinger.[4] In second order, a new type of graph appears for $n\geq2$. If we denote by R the ratio of the integral associated with this new type to the $n=1$ type, corrections to the Gaussian distribution can be written explicitly. The scaling form of $P(z)$ produced by this kind of renormalization-group treatment is

$$P(z)\propto\exp\{-\tfrac{1}{2}y^{(3-d)/2}(z/\xi)^2[1-(R-1)y^{(d-1)/2}+\tfrac{1}{6}Ry(z/\xi)^2+\ldots]\}. \tag{3}$$

In principle, systematic corrections to higher order can also be obtained. The function $P(z)$ can also be associated with a mean interface profile $\varphi(z)$ according to $d\varphi/dz=P(z)$. For small z, φ is the error function, and we may anticipate good fits to data on optical reflectivity[14] with this distribution.

In conclusion we see that the renormalization group enables one to control the high-wave-vector capillary waves. These waves can be identified as the Goldstone modes whose fluctuations reduce the critical temperature of Ising-like models to zero as the dimension of space is lowered to 1. Perturbative calculations yield ϵ expansions in $1+\epsilon$ dimensions. Although these are numerically rather poor for the exponent ν, and the role of nonperturbative effects such as overhangs[15] requires elucidation, these calculations do appear to support a useful interface phenomenology.

One of us (D.J.W.) thanks the Physics Departments of Virginia Polytechnic Institute & State University, Harvard University, and Massachusetts Institute of Technology for support and hospitality and one of us (R.K.P.Z.) thanks W. Klein and Boston University for hospitality during part of this work.

This work was supported in part through funds provided by the U. S. Department of Energy under Contract No. EY-76-C-02-3069 and by the National Science Foundation under Grant No. DMR 77-10210.

[a]On leave of absence at Virginia Polytechnic Institute & State University, Blacksburg, Va. 24061, Physics Department, Harvard University, Cambridge, Mass. 02138, and Center for Theoretical Physics, Massachusetts Institute of Technology, Cambridge, Mass. 02139. Permanent address after 1 October, 1979: Physics Department, University of Edinburgh, James Clerk Maxwell Building, Mayfield Road, Edinburgh EH9 3JZ, Scotland.

[1]J. D. van der Waals, Z. Phys. Chem. 13, 657 (1894); J. W. Cahn and J. E. Hilliard, J. Chem. Phys. 28, 258 (1958); B. Widom, in *Phase Transitions and Critical Phenomena,* edited by C. Domb and M. S. Green (Academic, New York, 1972), Vol. 2, p. 79; S. Fisk and B. Widom, J. Chem. Phys. 50, 3219 (1969); M. E. Fisher, M. N. Barber, and D. Jasnow, Phys. Rev. A 8, 1111 (1973).

[2]T. Ohta and K. Kawasaki, Prog. Theor. Phys. 58, 467 (1977); J. Rudnick and D. Jasnow, Phys. Rev. B 17, 1351 (1978). See also M. J. de Oliveira, D. Furman, and R. B. Griffiths, Phys. Rev. Lett. 40, 977 (1978).

[3]M. V. Smoluchowski, Ann. Physik 25, 205 (1908).

[4]F. P. Buff, R. A. Lovett, and F. H. Stillinger, Phys. Rev. Lett. 15, 621 (1965); J. Zittartz, Phys. Rev. 154, 529 (1967); R. A. Lovett, P. W. De Haven J. J. Vieceli, and F. P. Buff, J. Chem. Phys. 58, 1880 (1973); D. Jasnow and J. Rudnick, Phys. Rev. Lett. 41, 698 (1978).

[5]J. D. Weeks, G. H. Gilmer, and H. K. Leamy, Phys. Rev. Lett. 31, 549 (1973); S. T. Chui and J. D. Weeks, Phys. Rev. B 14, 4978 (1976); J. M. Kosterlitz, J. Phys. C 10, 3735 (1977); H. J. F. Knops, Phys. Rev. Lett. 39, 766 (1977); R. H. Swendsen, Phys. Rev. B 15, 689 (1977); J. D. Weeks, J. Chem. Phys. 67, 3106 (1977); H. T. Davis, J. Chem. Phys. 67, 3636 (1977); T. Ohta and K. Kawasaki, Prog. Theor. Phys. 60, 365 (1978).

[6]Results to first order in $\epsilon = d-1$ may be obtained using lattice renormalization-group methods: A. A. Migdal, Zh. Eksp. Teor. Fiz. 69, 1457 (1975) [Sov. Phys. JETP 42, 743 (1975)]; W. Klein, D. J. Wallace, and R. K. P. Zia, Phys. Rev. Lett. 37, 639 (1976). The field-theory approach we describe permits a systematic expansion to all orders in ϵ.

[7]N. J. Günther, D. A. Nicole, and D. J. Wallace, to be published. The interface is discussed in broken-symmetry terms also by M. S. Jhon, R. C. Desai, and J. S. Dahler, J. Chem. Phys. 68, 5615 (1978), and references therein. The dual-string model also involves a nonlinear realization (of the Poincaré group) on a two-component field representing the position of the string; see, e.g., P. Goddard, J. Goldstone, C. Rebbi, and C. B. Thorn, Nucl. Phys. B56, 104 (1973). We acknowledge conversations with J. Goldstone and C. B. Thorn on this topic.

[8]N. D. Mermin and H. Wagner, Phys. Rev. Lett. 17, 1133 (1966); P. C. Hohenberg, Phys. Rev. 158, 383 (1967); S. Coleman, Commun. Math. Phys. 31, 259 (1973). The spin-wave approach also gives ϵ expansions in $2+\epsilon$ dimensions: A. M. Polyakov, Phys. Lett. B59, 79 (1975), A. A. Migdal, Ref. 6; E. Brézin and J. Zinn-Justin, Phys. Rev. Lett. 36, 691 (1976), and Phys. Rev. B 14, 3110 (1976); E. Brézin, J. Zinn-Justin, and J. C. Le Guillou, Phys. Rev. D 14, 2615 (1976), and Phys. Rev. B 14, 4976 (1976); D. R. Nelson and R. A. Pelcovits, Phys. Lett. A57, 23 (1976); D. R. Nelson and R. A. Pelcovits, Phys. Rev. B 16, 2191 (1977); S. Hikami and E. Brézin, J. Phys. A 11, 1141 (1978).

[9]G. 't Hooft and M. Veltman, Nucl. Phys. B44, 185 (1973); C. G. Bollini and J. J. Giambiagi, Phys. Lett. B40, 566 (1972); D. J. Amit, *Field Theory, The Renormalization Group, and Critical Phenomena* (McGraw-Hill, New York, 1978).

[10]No additional vertices are required in the field theory; the naive measure Df is normally invariant under the Euclidean transformations on f because these involve a total derivative $\partial_i(x^2+f^2)$ and give Jacobian 1. Even if we had chosen parametrization of the displacement by a different field $\tilde{f} = F(f) = f + O(f^2)$ which did require a nontrivial measure, the use of dimensional regularization enables us to ignore the problem (cf. Brézin, Zinn-Justin, and Le Guillou, Ref. 8).

[11]The phase coherence length and its relationship to the bulk correlation length are discussed in Fisher, Barber, and Jasnow, Ref. 1.

[12]M. E. Fisher and A. E. Ferdinand, Phys. Rev. Lett. 19, 169 (1967); P. C. Watson, in *Phase Transitions and Critical Phenomena,* edited by C. Domb and M. S. Green (Academic, New York, 1972), Vol. 2, p. 101; R. K. P. Zia, Phys. Lett. A64, 345 (1978). A collection of exact results on surface tension in Ising models is also given in J. Bric-

mont, J. L. Lebowitz, and C. E. Pfister, to be published.
[13]J. S. Huang and W. W. Webb, Bull. Am. Phys. Soc. 12, 333 (1967).
[14]J. S. Huang and W. W. Webb, J. Chem. Phys. 50, 3677 (1969); E. S. Wu and W. W. Webb, Phys. Rev. A 8, 2065, 2077 (1973).
[15]In a $d=2$ antiferromagnetic model, E. Muller-Hartmann and J. Zittartz [Z. Physik B27, 261 (1977)] use a "no-overhang" method to calculate the surface energy successfully. See also H. N. V. Temperley, Proc. Cambridge Philos. Soc. 48, 683 (1952), and K. Y. Lin and F. Y. Wu, Z. Physik B33, 181 (1979). We are grateful to M. E. Fisher, D. R. Nelson, and F. Y. Wu for discussions on this topic. In a similar spirit, D. S. Fisher, R. H. Morf, and B. I. Halperin (to be published) show that the entropy associated with the wandering of a grain boundary in two dimensions determines a melting transition temperature identical to that of the dislocation unbinding of the bulk discussed by J. M. Kosterlitz and D. J. Thouless, J. Phys. C 6, 1181 (1973).

Two-Body Hadronic Decays of D Mesons

V. Barger and S. Pakvasa
Department of Physics and Astronomy, University of Hawaii at Manoa, Honolulu, Hawaii 96822
and Department of Physics, University of Wisconsin, Madison, Wisconsin 53706
(Received 25 May 1979)

This Letter accounts for the observed (i) enhancement of $\Gamma(D^0\to K^-K^+)/\Gamma(D^0\to K^-\pi^+)$ and (ii) suppression of $\Gamma(D^0\to\pi^-\pi^+)/\Gamma(D^0\to K^-\pi^+)$ relative to $\tan^2\theta_C$ in terms of (i) the ratio f_K/f_π of leptonic decay constants and (ii) different effective Cabibbo angles in non-charm and charm sectors. Strong new constraints on mixing angles in the six-quark model are implied.

In the four-quark model of the weak interaction,[1] relations between two-body channels in weak decays of charmed hadrons have been derived from SU(3) symmetry.[2] In particular, the ratios of Cabibbo-suppressed to allowed decay modes for D^0 decay were predicted to be

$$\frac{\Gamma(D^0\to K^-K^+)}{\Gamma(D^0\to K^-\pi^+)}=\frac{\Gamma(D^0\to\pi^-\pi^+)}{\Gamma(D^0\to K^-\pi^+)}=\tan^2\theta_C\simeq 0.05.$$

Taking into account phase-space corrections based on physical π and K masses, the SU(3) predictions become

$$\Gamma(D^0\to K^-K^+)/\Gamma(D^0\to K^-\pi^+)=4.7\%,$$
$$\Gamma(D^0\to \pi^-\pi^+)/\Gamma(D^0\to K^-\pi^+)=5.5\%.$$

Recent measurements at SPEAR are markedly different from these original theoretical expectations. The reported experimental ratios are[3]

$$\Gamma(D^0\to K^-K^+)/\Gamma(D^0\to K^-\pi^+)=(11.3\pm3.0)\%,$$
$$\Gamma(D^0\to \pi^-\pi^+)/\Gamma(D^0\to K^-\pi^+)=(3.3\pm1.5)\%. \quad (1)$$

In this Letter, we show that the experimental results can be understood in a model based on charmed-quark decay with hard-gluon corrections. Essential ingredients of the explanation are (i) known symmetry breaking in the ratio of K to π leptonic decay constants and (ii) weak-current angles in the six-quark model[4] allowed by analyses of the K^0-$\bar{K}^0$ mass matrix.[5,6]

Weak decays of mesons that have new flavor are nominally presumed to proceed through the decay of the heavy quark, with the light-quark constituent acting as a spectator. For two-body decays, estimates suggest that diagrams in which a quark-antiquark pair is created from the vacuum are negligible. With use of this model, charm-decay rates into two-body hadronic channels were first examined in asymptotically free gauge theory by Ellis, Gaillard, and Nanopoulos.[7] Further calculations of these modes were made by Fakirov and Stech[8] and then by Cabibbo and Maiani.[9] In the present work we extend the predictions to Cabibbo-suppressed two-body decays, with the mixing angles of the six-quark model.

The effective nonleptonic Lagrangian for $c\to\alpha u\bar{\beta}$ including the short-distance renormalization effects of hard gluons is[7,10]

$$\mathcal{L}_{\rm eff}=\frac{G}{\sqrt{2}}U_{u\beta}U_{c\alpha}{}^*[\tfrac{1}{2}(c_++c_-)(\bar{u}\beta)(\bar{\alpha}c)+\tfrac{1}{2}(c_+-c_-)(\bar{\alpha}\beta)(\bar{u}c)], \quad (2)$$

where $(q_1\bar{q}_2)$ denotes a color-symmetric $V-A$ current and U is the Kobayashi-Maskawa[4] mixing matrix of the six-quark model (hereafter called KM matrix). The renormalization constants are $c_-=[\alpha(m_c{}^2)/$

