Chiral Perturbation Theory for Vector Mesons

Elizabeth Jenkins and Aneesh V. Manohar
Department of Physics, University of California at San Diego, 9500 Gilman Drive, La Jolla, California 92093-0319

Mark B. Wise
California Institute of Technology, Pasadena, California 91125
(Received 19 June 1995)

We derive a heavy vector-meson chiral Lagrangian in which the vector mesons are treated as heavy static matter fields. The unknown couplings of the chiral Lagrangian are further related using the $1/N_c$ expansion. Chiral perturbation theory is applied to the vector-meson mass matrix. At one loop there are large corrections to the individual vector-meson masses, but the singlet-octet mixing angle remains almost unchanged. The parity-violating s-wave $\phi \to \rho\pi$ weak decay amplitude is derived in the combined chiral and large N_c limits. Rare ϕ decays provide a sensitive test of nonleptonic neutral current structure.

PACS numbers: 12.39.Fe, 11.15.Pg, 13.25.Gv, 14.40.Cs

An important application of chiral perturbation theory is to describe the interactions of matter fields (such as nucleons [1] or hadrons containing a heavy quark [2]) with low-momentum pseudo Goldstone bosons—the pions, kaons and, eta. In this paper we use chiral perturbation theory to describe the interactions of the ρ, K^*, ϕ, and ω vector mesons with low-momentum pseudo Goldstone bosons. The results of this work are relevant for experiments at the ϕ factory being built at Frascati [3]. We apply chiral perturbation theory to transitions of the form $V \to V'X$, where V and V' are vector mesons. The mass differences between the nine lowest-lying vector mesons are small compared with the chiral symmetry breaking scale of $\sim$1 GeV, so chiral perturbation theory is applicable as a systematic expansion procedure for such decays. Chiral perturbation theory has previously been used to study processes such as $\rho \to \pi\pi$, which do not have a vector meson in the final state. Decays such as $\rho \to \pi\pi$ do not have soft pions in the final state, so the application of the chiral Lagrangian to such processes is not justified and should be considered as a phenomenological model.

The pseudo Goldstone boson fields can be written as a 3×3 special unitary matrix

$$\Sigma = \exp\frac{2i\mathbf{\Pi}}{f}, \tag{1}$$

where

$$\mathbf{\Pi} = \begin{bmatrix} \pi^0/\sqrt{2} + \eta/\sqrt{6} & \pi^+ & K^+ \\ \pi^- & -\pi^0/\sqrt{2} + \eta/\sqrt{6} & K^0 \\ K^- & \overline{K}^0 & -2\eta/\sqrt{6} \end{bmatrix}. \tag{2}$$

Under chiral $SU(3)_L \times SU(3)_R$, $\Sigma \to L\Sigma R^\dagger$, where $L \in SU(3)_L$ and $R \in SU(3)_R$. At leading order in chiral perturbation theory, f can be identified with the pion or kaon decay constant ($f_\pi \simeq 132$ MeV, $f_K \simeq 160$ MeV). It is convenient when describing the interactions of the pseudo Goldstone bosons with other fields to introduce

$$\xi = \exp\frac{i\mathbf{\Pi}}{f} = \sqrt{\Sigma}. \tag{3}$$

Under chiral $SU(3)_L \times SU(3)_R$,

$$\xi \to L\xi U^\dagger = U\xi R^\dagger, \tag{4}$$

where, in general, U is a complicated function of L, R, and the meson fields $\mathbf{\Pi}$. For transformations $V = L = R$ in the unbroken $SU(3)_V$ subgroup, $U = V$.

The vector-meson fields are introduced as a 3×3 octet matrix

$$\mathcal{O}_\mu = \begin{bmatrix} \rho^0_\mu/\sqrt{2} + \phi^{(8)}_\mu/\sqrt{6} & \rho^+_\mu & K^{*+}_\mu \\ \rho^-_\mu & -\rho^0_\mu/\sqrt{2} + \phi^{(8)}_\mu/\sqrt{6} & K^{*0}_\mu \\ K^{*-}_\mu & \overline{K}^{*0}_\mu & -2\phi^{(8)}_\mu/\sqrt{6} \end{bmatrix}, \tag{5}$$

and as a singlet

$$S_\mu = \phi^{(0)}_\mu. \tag{6}$$

Under chiral $SU(3)_L \times SU(3)_R$,

$$\mathcal{O}_\mu \to U\mathcal{O}_\mu U^\dagger, \quad S_\mu \to S_\mu, \tag{7}$$

and under charge conjugation,

$$C\mathcal{O}_\mu C^{-1} = -\mathcal{O}_\mu^T, \quad CS_\mu C^{-1} = -S_\mu, \quad C\xi C^{-1} = \xi^T. \tag{8}$$

We construct a chiral Lagrangian for vector mesons by treating the vector mesons as heavy static fields

 0031-9007/95/75(12)/2272(4)$06.00

[4,5] with fixed four-velocity v^μ, $v^2 = 1$. The three polarization states of vector mesons with velocity v^μ satisfy $v \cdot S = v \cdot \mathcal{O} = 0$. The chiral Lagrange density which describes the interactions of the vector mesons with the low-momentum π, K, and η mesons has the general structure

$$\mathcal{L} = \mathcal{L}_{\rm kin} + \mathcal{L}_{\rm int} + \mathcal{L}_{\rm mass}\,. \tag{9}$$

Processes such as $\phi \to K^+K^-$ cause a net disappearance of vector mesons and can be taken into account by including anti-Hermitian terms in $\mathcal{L}$. The vector-meson widths are small compared with M_K and will be neglected in our analysis.

At leading order in the derivative and quark mass expansions

$$\mathcal{L}_{\rm kin} = -i\,S_\mu^\dagger(v\cdot\partial)S^\mu - i\,{\rm Tr}\,\mathcal{O}_\mu^\dagger(v\cdot\mathcal{D})\mathcal{O}^\mu \tag{10}$$

and

$$\mathcal{L}_{\rm int} = ig_1\,S_\mu^\dagger\,{\rm Tr}\,(\mathcal{O}_\nu A_\lambda)v_\sigma\epsilon^{\mu\nu\lambda\sigma} + {\rm H.c.} + ig_2\,{\rm Tr}\,(\{\mathcal{O}_\mu^\dagger,\mathcal{O}_\nu\}A_\lambda)v_\sigma\epsilon^{\mu\nu\lambda\sigma}, \tag{11}$$

where

$$\mathcal{D}^\nu\mathcal{O}^\mu = \partial^\nu\mathcal{O}^\mu + [V^\nu,\mathcal{O}^\mu] \tag{12}$$

and

$$V^\mu = \frac{1}{2}(\xi\partial^\mu\xi^\dagger + \xi^\dagger\partial^\mu\xi),\qquad A^\mu = \frac{i}{2}(\xi\partial^\mu\xi^\dagger - \xi^\dagger\partial^\mu\xi)\,. \tag{13}$$

The terms in $\mathcal{L}_{\rm kin}$ appear with minus signs because the polarization vector is spacelike. Charge conjugation invariance requires that the product of $\mathcal{O}_\mu^\dagger$ and $\mathcal{O}_\nu$ in the second term of Eq. (11) be an anticommutator. (It is important to remember that in the matrix $\mathcal{O}_\mu^\dagger$ the field $\rho_\mu^{-\dagger}$ is not equal to ρ_μ^+, etc. In heavy vector-meson chiral perturbation theory, ρ_μ^+ destroys a ρ^+ but it does not create the corresponding antiparticle. A separate field $\rho_\mu^{-\dagger}$ is introduced to create a ρ^-.) Finally, to linear order in the quark mass expansion,

$$\begin{aligned}\mathcal{L}_{\rm mass} = {}& \mu_0\,S_\mu^\dagger S^\mu + \mu_8\,{\rm Tr}\,\mathcal{O}_\mu^\dagger\mathcal{O}^\mu \\ &+ \lambda_1\left[{\rm Tr}\,(\mathcal{O}_\mu^\dagger\mathcal{M}_\xi)S^\mu + {\rm H.c.}\right] \\ &+ \lambda_2\,{\rm Tr}\,(\{\mathcal{O}_\mu^\dagger,\mathcal{O}^\mu\}\mathcal{M}_\xi) \\ &+ \sigma_0\,{\rm Tr}\,\mathcal{M}_\xi S_\mu^\dagger S^\mu + \sigma_8\,{\rm Tr}\,\mathcal{M}_\xi{\rm Tr}\,\mathcal{O}_\mu^\dagger\mathcal{O}^\mu,\end{aligned} \tag{14}$$

where $\mathcal{M}$ is the quark mass matrix $\mathcal{M} = {\rm diag}(m_u, m_d, m_s)$ and

$$\mathcal{M}_\xi = \frac{1}{2}(\xi\mathcal{M}\xi + \xi^\dagger\mathcal{M}\xi^\dagger)\,. \tag{15}$$

Note that the fields S and $\mathcal{O}$ appearing in Eqs. (10)–(12) are understood to be velocity-dependent fields which are rescaled by a common phase factor (either $e^{-i\mu_0 v\cdot x}$ or $e^{-i\mu_8 v\cdot x}$). (The velocity-dependent vector-meson fields are related to the vector-meson fields by $\phi_v^\mu = \sqrt{2m}\,e^{imv\cdot x}\phi^\mu$ and have dimension 3/2.) This rescaling removes either μ_0 or μ_8 from Eq. (14), so only the singlet-octet mass difference $\Delta\mu \equiv \mu_0 - \mu_8$ is relevant. Phenomenologically, the parameter $\Delta\mu < 200$ MeV is comparable to mass splittings of order m_s, so in our power counting we treat $\Delta\mu$ as a quantity of order m_q. $\Delta\mu$ is of order $1/N_c$, and so vanishes in the large N_c limit.

We begin by considering the spectrum of vector mesons produced at leading order in chiral perturbation theory. The analysis is identical to the well-known SU(3) analysis [6]. Neglecting isospin breaking due to the up and down quark mass difference, i.e., $m_u = m_d = \hat{m}$, we find that

$$m_\rho = \bar{\mu}_8 + 2\lambda_2\hat{m},\quad m_{K^*} = \bar{\mu}_8 + \lambda_2(\hat{m} + m_s), \tag{16a}$$

and the $\phi^{(0)} - \phi^{(8)}$ mass matrix is

$$M^{(08)} = \begin{bmatrix} \bar{\mu}_0 & -\frac{2}{\sqrt{6}}\lambda_1(m_s - \hat{m}) \\ -\frac{2}{\sqrt{6}}\lambda_1(m_s - \hat{m}) & \bar{\mu}_8 + \frac{2}{3}\lambda_2(\hat{m} + 2m_s)\end{bmatrix}, \tag{16b}$$

where

$$\bar{\mu}_0 = \mu_0 + \sigma_0{\rm Tr}\mathcal{M},\quad \bar{\mu}_8 = \mu_8 + \sigma_8{\rm Tr}\mathcal{M}\,. \tag{17}$$

(Isospin breaking effects have been studied in Ref. [7].) Using Eqs. (16a) and (16b), it is possible to express the elements of $M^{(08)}$ in terms of the measured vector-meson masses (up to a sign ambiguity for $M_{12}^{(08)}$)

$$M_{11}^{(08)} = m_\omega + m_\phi - \frac{4}{3}m_{K^*} + \frac{1}{3}m_\rho\,, \tag{18a}$$

$$M_{22}^{(08)} = \frac{4}{3}m_{K^*} - \frac{1}{3}m_\rho\,, \tag{18b}$$

$$M_{12}^{(08)} = M_{21}^{(08)} = \pm\left[\left(\frac{4}{3}m_{K^*} - \frac{1}{3}m_\rho - m_\omega\right)\times\left(m_\phi - \frac{4}{3}m_{K^*} + \frac{1}{3}m_\rho\right)\right]^{1/2}. \tag{18c}$$

The eigenstates of $M^{(08)}$ are parametrized by a mixing angle Θ_V

$$|\phi\rangle = \sin\Theta_V|\phi^{(0)}\rangle - \cos\Theta_V|\phi^{(8)}\rangle, \tag{19a}$$

$$|\omega\rangle = \cos\Theta_V|\phi^{(0)}\rangle + \sin\Theta_V|\phi^{(8)}\rangle, \tag{19b}$$

where Eqs. (18) imply the usual SU(3)$_V$ prediction for the tangent of the mixing angle

$$\tan\Theta_V = \mp\sqrt{\frac{m_\phi - \frac{4}{3}m_{K^*} + \frac{1}{3}m_\rho}{\frac{4}{3}m_{K^*} - \frac{1}{3}m_\rho - m_\omega}} \simeq \mp 0.76\,. \tag{20}$$

In the large N_c limit [8,9], quark loops are suppressed so that the leading diagrams in the meson sector contain a single quark loop. As a result, the octet and singlet mesons can be combined into a single "nonet" matrix

$$N_\mu = \mathcal{O}_\mu + \frac{I}{\sqrt{3}}S_\mu\,, \tag{21}$$

which enters the chiral Lagrangian. The kinetic, interaction, and mass terms at leading order in $1/N_c$ are

$$\mathcal{L}_{\rm kin} \to -i\, \mathrm{Tr}\, N_\mu^\dagger (v \cdot \mathcal{D}) N^\mu, \tag{22}$$

$$\mathcal{L}_{\rm int} \to i g_2\, \mathrm{Tr}\, (\{N_\mu^\dagger, N_\nu\} A_\lambda) v_\sigma \epsilon^{\mu\nu\lambda\sigma}, \tag{23}$$

and

$$\mathcal{L}_{\rm mass} \to \mu\, \mathrm{Tr}\, N_\mu^\dagger N^\mu + \lambda_2\, \mathrm{Tr}\, (\{N_\mu^\dagger, N^\mu\} \mathcal{M}_\xi). \tag{24}$$

Comparing with Eqs. (10)–(14), one finds that, in the $N_c \to \infty$ limit,

$$\Delta\mu \to 0, \quad \sigma_0 \to \frac{2\lambda_2}{3}, \quad \sigma_8 \to 0, \tag{25}$$

$$g_1 \to \frac{2 g_2}{\sqrt{3}}, \quad \lambda_1 \to \frac{2\lambda_2}{\sqrt{3}}, \quad \tan\Theta_V \to \frac{1}{\sqrt{2}}, \tag{26}$$

the $|\phi\rangle$ state becomes "pure" $|s\bar{s}\rangle$, and the nonet matrix is

$$N_\mu = \begin{bmatrix} \rho_\mu^0/\sqrt{2} + \omega_\mu/\sqrt{2} & \rho_\mu^+ & K_\mu^{*+} \\ \rho_\mu^- & -\rho_\mu^0/\sqrt{2} + \omega_\mu/\sqrt{2} & K_\mu^{*0} \\ K_\mu^{*-} & \overline{K}_\mu^{*0} & \phi_\mu \end{bmatrix}. \tag{27}$$

If the minus sign is chosen in Eq. (18c), the prediction for the mixing angle at leading order in chiral perturbation theory, Eq. (20), is close to its value for large N_c. At leading order in chiral perturbation theory, the partial width for the Zweig forbidden decay $\phi \to \rho\pi$ summed over all three modes is

$$\Gamma(\phi \to \rho\pi) = \frac{2h^2 |\boldsymbol{p}_\pi|^3}{\pi f^2}. \tag{28}$$

The coupling, h, which vanishes as $N_c \to \infty$, is

$$h = \frac{g_1}{\sqrt{2}} \sin\Theta_V - \frac{g_2}{\sqrt{3}} \cos\Theta_V. \tag{29}$$

The measured branching ratio gives $h \approx 0.05$, which also suggests that the couplings are close to the $N_c \to \infty$ values.

In the nonrelativistic constituent quark model, assuming the $|\phi\rangle$ is pure $|s\bar{s}\rangle$, $g_1 = 2/\sqrt{3}$ and $g_2 = 1$. In the nonrelativistic chiral quark model [10], g_1 and g_2 are reduced by a factor of 0.75 from their values in the nonrelativistic constituent quark model.

In chiral perturbation theory the leading corrections to the expressions for the vector-meson masses in Eqs. (16) are of order $m_q^{3/2}$ (recall we are treating $\Delta\mu$ as of order m_q) and arise from one-loop self-energy diagrams giving

$$\begin{aligned}
\delta m_\rho &= -\frac{1}{12\pi f^2}\left[g_2^2\left(\frac{2}{3}m_\pi^3 + 2m_K^3 + \frac{2}{3}m_\eta^3\right) + g_1^2 m_\pi^3\right], \\
\delta m_{K^*} &= -\frac{1}{12\pi f^2}\left[g_2^2\left(\frac{3}{2}m_\pi^3 + \frac{5}{3}m_K^3 + \frac{1}{6}m_\eta^3\right) + g_1^2 m_K^3\right], \\
\delta M_{11}^{(08)} &= -\frac{1}{12\pi f^2} g_1^2 (3m_\pi^3 + 4m_K^3 + m_\eta^3), \\
\delta M_{22}^{(08)} &= -\frac{1}{12\pi f^2}\left[g_2^2\left(2m_\pi^3 + \frac{2}{3}m_K^3 + \frac{2}{3}m_\eta^3\right) + g_1^2 m_\eta^3\right], \\
\delta M_{12}^{(08)} &= \delta M_{21}^{(08)} = \frac{1}{12\pi f^2}\sqrt{\frac{2}{3}}\, g_1 g_2 (-3m_\pi^3 + 2m_K^3 + m_\eta^3).
\end{aligned} \tag{30}$$

The singlet-octet mixing angle Θ_V including these corrections is

$$\tan\Theta_V = \mp\sqrt{\frac{m_\phi - \frac{4}{3}m_{K^*} + \frac{1}{3}m_\rho - \delta m}{\frac{4}{3}m_{K^*} - \frac{1}{3}m_\rho - m_\omega + \delta m}}, \tag{31}$$

where

$$\begin{aligned}
\delta m &= -\frac{4}{3}\delta m_{K^*} + \frac{1}{3}\delta m_\rho + \delta M_{22}^{(08)} \\
&= -\frac{1}{12\pi f^2}\left(g_1^2 + \frac{2}{3}g_2^2\right)\left(\frac{1}{3}m_\pi^3 - \frac{4}{3}m_K^3 + m_\eta^3\right).
\end{aligned} \tag{32}$$

Using the relation between g_2 and g_1 in Eq. (26), we find that

$$\delta m \to -\frac{2g_2^2}{12\pi f^2}\left(\frac{1}{3}m_\pi^3 - \frac{4}{3}m_K^3 + m_\eta^3\right). \tag{33}$$

With $g_2 = 0.75$, Eq. (33) yields $\delta m \simeq -4$ MeV. The combination of mass shifts, δm, that affects the mixing angle Θ_V is very small even though the corrections to the individual masses are substantial (e.g., $\delta m_\rho \simeq -300$ MeV). δm, which is of order $1/N_c$, must transform like a **27** of flavor SU(3). The linear combination of the cubed pseudoscalar meson mass in Eq. (33) transforms like a **27**, and is numerically small. This same linear combination occurs in the violation of the Gell-Mann–Okubo formula for baryon masses [11].

For N_c large, the $\phi \to \rho\pi$ decay amplitude is of order $N_c^{-3/2}$, since the leading order $1/\sqrt{N_c}$ amplitude is forbidden by Zweig's rule. At leading order in chiral perturbation theory it occurs at tree level because of order $1/N_c$ deviations from the relations $\tan\Theta_V = 1/\sqrt{2}$ and $g_1/g_2 = 2/\sqrt{3}$. At order $m_s \ln m_s$ in the chiral expansion the order $N_c^{-3/2}$ contribution arises from one-loop vertex and wave function corrections calculated with vertices from the nonet Lagrange density. The π and η loops do not contribute when one uses the nonet Lagrangian. The resulting decay amplitude for each of the three $\rho\pi$ modes is

$$\begin{aligned}
\mathcal{A}(\phi \to \rho\pi) = &\frac{i}{f}\epsilon^{\mu\nu\lambda\sigma}\epsilon_\mu(\phi)\epsilon_\nu^*(\rho) p_{\pi\lambda} v_\sigma (2\sqrt{m_\phi m_\rho}) \\
&\times\left[\sqrt{2}\,h - g_2^3\left(\frac{m_K^2}{8\pi^2 f^2}\right)\ln\left(\frac{m_K^2}{\mu^2}\right) + \cdots\right],
\end{aligned} \tag{34}$$

where the ellipsis denotes terms of higher order in the chiral and $1/N_c$ expansions. The terms of order m_s have a dependence on the subtraction point μ which cancels that of the logarithm in Eq. (34). With $\mu = 1$ GeV, $g_2 = 0.75$, the magnitude of the term of order $m_s \ln m_s$ in Eq. (34) is about 1.5 times as large as the measured $\phi \to \rho\pi$ decay amplitude. This suggests either that g_2 is smaller than the chiral quark model value or that there is a partial cancellation between order $m_s \ln m_s$ and order m_s contributions to the decay amplitude.

The Frascati ϕ factory is expected to produce of order 10^{10} ϕ's, allowing even very rare ϕ decay processes to be experimentally accessible. The $\phi \to \rho\pi$ decay amplitude has a small parity-violating s-wave amplitude that is induced by the weak interactions. This amplitude can be predicted in the combined limits of chiral $SU(2)_L \times SU(2)_R$ symmetry and large N_c. In these limits the part of the weak Hamiltonian that dominates the s-wave $\phi \to \rho\pi$ amplitude is due to Z^0 exchange,

$$\mathcal{H}_W = \eta \frac{G_F}{2\sqrt{2}}\left(1 - \frac{4}{3}\sin^2\theta_W\right) \times (\bar{s}_\alpha \gamma_\mu s_\alpha)[\bar{u}_\beta \gamma^\mu \gamma_5 u_\beta - \bar{d}_\beta \gamma^\mu \gamma_5 d_\beta], \tag{35}$$

where $\eta \sim 1.56$ [12] arises from QCD scaling between the weak scale and low energies. In the large N_c limit, $\log\eta$ is of order $1/N_c$ times logarithms of the form $\log M_W/\Lambda$, and we have chosen to include corrections of this form [13]. The $\phi \to \rho\pi$ matrix element takes the form

$$\langle\rho\pi|\mathcal{H}_W|\phi\rangle = \eta \frac{G_F}{2\sqrt{2}}\left(1 - \frac{4}{3}\sin^2\theta_W\right) f_\phi\, \epsilon_\mu(\phi) \times \langle\rho\pi|\bar{u}\gamma^\mu\gamma_5 u - \bar{d}\gamma^\mu\gamma_5 d|0\rangle, \tag{36}$$

where the ϕ decay constant f_ϕ is defined by

$$\langle 0|\bar{s}\gamma_\mu s|\phi\rangle = f_\phi \epsilon_\mu(\phi). \tag{37}$$

The measured $\phi \to e^+e^-$ decay width implies that $f_\phi \simeq (492\ \mathrm{MeV})^2$. The left-handed isovector current $\bar{u}\gamma^\mu P_L u - \bar{d}\gamma^\mu P_L d$ transforms as $(3_L, 1_R)$ under chiral $SU(2)_L \times SU(2)_R$. For matrix elements between the vacuum and a ρ plus soft pions, this current is represented by the operator

$$\bar{u}\gamma^\mu P_L u - \bar{d}\gamma^\mu P_L d = \frac{f_\rho}{2\sqrt{2m_\rho}} \mathrm{Tr}(\xi \mathcal{O}_\mu^\dagger \xi^\dagger \tau^3), \tag{38}$$

where f_ρ is defined analogous to f_ϕ and has the value $f_\rho \simeq (407\ \mathrm{MeV})^2$ from the $\rho \to e^+e^-$ partial width. In Eq. (38), ξ and $\mathcal{O}_\mu$ are the 2×2 matrix analogs of the corresponding 3×3 matrices used in the case of chiral $SU(3)_L \times SU(3)_R$. The right-handed isovector current is given by exchanging ξ and $\xi^\dagger$ in Eq. (38). Using Eq. (38) to evaluate the matrix element in Eq. (36) we find that in the combined chiral and large N_c limits

$$\begin{aligned}\langle\rho^+\pi^-|\mathcal{H}_W|\phi\rangle_{s\text{-wave}} &= -\langle\rho^-\pi^+|\mathcal{H}_W|\phi\rangle_{s\text{-wave}} \\ &= -i\eta\frac{G_F}{\sqrt{2}}\left(1 - \frac{4}{3}\sin^2\theta_W\right) \\ &\quad \times \left(\frac{f_\phi f_\rho}{f}\right)\epsilon^*(\rho)\cdot\epsilon(\phi)\end{aligned} \tag{39}$$

and

$$\langle\rho^0\pi^0|\mathcal{H}_W|\phi\rangle_{s\text{-wave}} = 0. \tag{40}$$

Interference between the s-wave and p-wave amplitudes is possible for aligned ϕ's, but it requires a final state interaction phase. The s-wave $\rho^+\pi^-$ branching ratio is 10^{-11}, which is too small to be measured at the Frascati ϕ factory. However, an enhancement of the parity-violating decay rate could make the signal observable. This provides a very interesting test of new physics, because it probes nonleptonic neutral currents involving strange quarks. We will consider the application of chiral perturbation theory to other processes such as $\phi \to \rho\gamma\gamma$ elsewhere.

This work was supported in part by the Department of Energy under Grant No. DOE-FG03-90ER40546 and No. DE-FG03-92-ER40701. E. J. was supported in part by the NYI program, through Grant No. PHY-9457911 from the National Science Foundation, and by the Alfred P. Sloan Foundation. A. M. was supported in part by the PYI program, through Grant No. PHY-8958081 from the National Science Foundation.

[1] J. Gasser, M. E. Sainio, and A. Švarc, Nucl. Phys. **B307**, 779 (1988); E. Jenkins and A. V. Manohar, Phys. Lett. B **255**, 558 (1991); **259**, 353 (1991).
[2] M. B. Wise, Phys. Rev. D **45**, 2188 (1992); G. Burdman and J. F. Donoghue, Phys. Lett. B **280**, 287 (1992); T.-M. Yan *et al.,* Phys. Rev. D **46**, 1148 (1992); P. Cho, Nucl. Phys. **B396**, 183 (1993).
[3] *The DAΦNE Physics Handbook,* edited by L. Maiani, G. Pancheri, and N. Paver (INFN, Frascati, Italy, 1992).
[4] E. Eichten and B. Hill, Phys. Lett. B **234**, 511 (1990).
[5] H. Georgi, Phys. Lett. B **240**, 447 (1990).
[6] J. J. J. Kokkedee, *The Quark Model* (Benjamin, New York, 1969).
[7] R. Urech, Report No. hep-ph/9504238.
[8] G. 't Hooft, Nucl. Phys. **B72**, 461 (1974).
[9] G. Veneziano, Nucl. Phys. **B117**, 519 (1976).
[10] A. Manohar and H. Georgi, Nucl. Phys. **B234**, 189 (1984).
[11] E. Jenkins, Nucl. Phys. **B368**, 190 (1992).
[12] J. Dai, M. J. Savage, J. Liu, and R. Springer, Phys. Lett. B **271**, 403 (1991).
[13] W. Bardeen, A. Buras, and J. M. Gerard, Nucl. Phys. **B293**, 787 (1987).