in good agreement with the experimental result[4] of $\sim -0.8\hbar/\tau_1$.

For the sake of completeness, we estimate the mass difference arising due to $K_2 \to \pi \to \sigma + \pi \to 3\pi$. We follow the same procedure as for the mass difference arising from $K_1 \to \sigma \to 2\pi$. With the normalization of $\Gamma_{K_2 \to 3\pi}(m_K^2) = \hbar/\tau_2$ we get a mass difference $\delta m = 1.0\hbar/\tau_2$, justifying the neglect of the contribution to the mass difference coming from the process $K_2 \to 3\pi$.

Direct search for the σ and σ'.—The direct search for σ and σ' is easier if we have a reaction which does not include the production of the ρ^0. Such a reaction is

$$\pi^+ + d \to p + p + \pi^0 + \pi^0. \tag{8}$$

One may then look for bumps in the missing-mass plot, corresponding to the σ and σ'. Such an experiment has already been done by Gelfand et al.[7] In their plot of the missing mass, we make a theoretical estimate of about 10 events for the σ ($m_\sigma = 400$ MeV and $\Gamma_\sigma = 70$ MeV) and about 30 events for the σ' ($m_{\sigma'} = 700$ MeV and $\Gamma_{\sigma'} = 150$ MeV). There are some candidates for the σ but the statistics are not good enough to decide the question on σ yet; however, the σ' with the width of 150 MeV seems to be ruled out.

Discussion.—(1) From the above analysis we conclude that if the σ exists, then a large constant s-wave ($I = 0$) phase shift in the region of the ρ is the only alternative which is consistent with the asymmetry in the $\pi^+\pi^-$ center-of-mass distribution in the ρ^0 production, the K_1K_2 mass difference, and the experiment of Gelfand et al.[7] This conclusion, however, is based on the sign of $\delta m = m_{K_1} - m_{K_2}$ being negative, the experimental evidence for which is rather weak.[4]

(2) One can explain the K_1K_2 mass difference in terms of the strong s-wave interaction around the ρ-meson position, and hence one must take into account the large s wave around the ρ in calculating the mass difference.

(3) The existence of the σ implies two zeros in the asymmetry plot, one at about the σ position and the other between the σ- and the ρ-meson positions, providing a sensitive test for the existence of the σ meson. The presently existing data on the asymmetry[6] then force the position of the possible σ to be less than 400 MeV and the width ≤ 100 MeV.

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NONRENORMALIZATION THEOREM FOR THE STRANGENESS-VIOLATING VECTOR CURRENTS

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On the assumptions that (i) the vector currents and the electromagnetic current belong to the same unitary octet,[1-3] and that (ii) the breaking of unitary symmetry is due to a term behaving like the eighth component of an octet,[1] we prove the following result: To first order in the symmetry-breaking interaction all the vector coupling constants are not renormalized. A useful application of the result is to strangeness-violating leptonic decays of baryons and mesons: The vector coupling constants (i.e., the limits of the vector amplitudes for vanishing momentum transfer) are uniquely predicted up to first order in symmetry breaking.

To prove the theorem we first remark that, as noted by Gell-Mann,[4] the vector octet we are considering (which includes the electromagnetic current) must have $\mathcal{C} = -1$ ($\mathcal{C}$ is the charge-conjugation quantum number of the components 1, 3, 4, 6, 8; the charge conjugation of the components 2, 5, 7 is $-\mathcal{C}$). It follows that, for first-class covariants[5] (γ_μ and $\sigma_{\mu\nu}k_\nu$ where k_ν is the momentum transfer), to first order in the symmetry breaking the ith component of the current can be expanded as

$$a_0\mathrm{Tr}(\bar{B}B\lambda_i)+b_0\mathrm{Tr}(\bar{B}\lambda_i B)+a\,\mathrm{Tr}(\bar{B}B\{\lambda_i,\lambda_8\})+b\,\mathrm{Tr}(\bar{B}\{\lambda_i,\lambda_8\}B)+c[\mathrm{Tr}(\bar{B}\lambda_i B\lambda_8)-\mathrm{Tr}(\bar{B}\lambda_8 B\lambda_i)]$$
$$+g\,\mathrm{Tr}(\bar{B}B)\,\mathrm{Tr}(\lambda_i\lambda_8)+h[\mathrm{Tr}(\bar{B}\lambda_i)\,\mathrm{Tr}(B\lambda_8)+\mathrm{Tr}(\bar{B}\lambda_8)\,\mathrm{Tr}(B\lambda_i)],\tag{1}$$

where B represents the baryons, the matrices λ_i are defined as in reference 1, $a_0, b_0, \cdots, h$ are first-class amplitudes.

For the electromagnetic current, defined as $j_{\mathrm{em}}=j_3+j_8/\sqrt{3}$, the expansion (1) becomes

$$(a_0+2a/\sqrt{3})\,\mathrm{Tr}(\bar{B}B\lambda_3)+(b_0+2b/\sqrt{3})\,\mathrm{Tr}(\bar{B}\lambda_3 B)+(a_0-2a/\sqrt{3})(1/\sqrt{3})\,\mathrm{Tr}(\bar{B}B\lambda_8)+(b_0-2b/\sqrt{3})(1/\sqrt{3})\,\mathrm{Tr}(\bar{B}\lambda_8 B)$$
$$+(2/\sqrt{3})(\tfrac{2}{3}a+\tfrac{2}{3}b+g)\,\mathrm{Tr}(\bar{B}B)+c[\mathrm{Tr}(\bar{B}\lambda_3 B\lambda_8)-\mathrm{Tr}(\bar{B}\lambda_8 B\lambda_3)]+h[\mathrm{Tr}(\bar{B}\lambda_3)\,\mathrm{Tr}(B\lambda_8)+\mathrm{Tr}(\bar{B}\lambda_8)\,\mathrm{Tr}(B\lambda_3)$$
$$+(2/\sqrt{3})\,\mathrm{Tr}(\bar{B}\lambda_8)\,\mathrm{Tr}(B\lambda_8)].\tag{2}$$

We know that this current is nonrenormalized; in the limit of vanishing momentum transfer we thus have

$$(2)=-\tfrac{1}{2}[\mathrm{Tr}(\bar{B}B\lambda_3)-\mathrm{Tr}(\bar{B}\lambda_3 B)+(1/\sqrt{3})\,\mathrm{Tr}(\bar{B}B\lambda_8)-(1/\sqrt{3})\,\mathrm{Tr}(\bar{B}\lambda_8 B)].\tag{3}$$

Equation (3) has the unique solution $a_0=-\frac{1}{2}$, $b_0=\frac{1}{2}$, $a=b=c=g=h=0$. We conclude that in the limit of vanishing momentum transfer, the amplitudes $a_0, b_0, \cdots, h$ in the general expansion (1) conserve the same values as in the unitary-symmetric limit.[6]

We can illustrate in more physical terms the meaning of Eq. (3). We have expressed, with Eq. (2), the (first-class) electromagnetic amplitudes for $\gamma\to B+\bar{B}$, with first-order symmetry breaking, in terms of the seven coupling constants a_0, b_0, a, b, c, g, h. We have required, with Eq. (3), that, for vanishing momentum transfer, the amplitudes must be proportional to the charge of B. This requirement is equivalent to a set of seven[7] inhomogeneous linear relations that must be satisfied by the seven coupling constants. The determinant formed with the coefficients of this linear set of equations is different from zero. The (unique) solution must thus coincide with the obvious solution of the unitary-symmetric limit.

We add a few remarks on vector and axial currents at first order in the symmetry-breaking interaction. First-class axial amplitudes (those proportional to the covariants $\gamma_\mu\gamma_5$ and $k_\mu\gamma_5$) can be expanded in exactly the same form (1). The expansion implies the known relations among first-class amplitudes (vector and axial): $A(\bar{\Lambda}\Sigma^-)=A(\bar{\Sigma}^+\Lambda)$ and $A(\bar{\Sigma}^+\Sigma^0)=-A(\bar{\Sigma}^0\Sigma^-)$. For second-class amplitudes (vector covariant k_μ and axial covariant $\sigma_{\mu\nu}k_\nu\gamma_5$) instead of (1) one has the expansion

$$a\,\mathrm{Tr}(\bar{B}B[\lambda_i,\lambda_8])+b\,\mathrm{Tr}(\bar{B}[\lambda_i,\lambda_8]B)$$
$$+h[\mathrm{Tr}(\bar{B}\lambda_i)\,\mathrm{Tr}(B\lambda_8)-\mathrm{Tr}(\bar{B}\lambda_8)\,\mathrm{Tr}(B\lambda_i)].\tag{4}$$

The expansion (4) implies (i) the absence of second-class amplitudes in the unitary-symmetric limit[8]; (ii) the known relation $A(\bar{\Lambda}\Sigma^-)=-A(\bar{\Sigma}^+\Lambda)$ among second-class amplitudes (all the other $\Delta S=0$ second-class amplitudes must vanish); (iii) the relation

$$-\sqrt{6}[A(\bar{p}\Lambda)+A(\bar{\Lambda}\Xi^-)]=A(\bar{n}\Sigma^-)+A(\bar{\Sigma}^+\Xi^0)$$

among $\Delta S=1$ second-class amplitudes.

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[3]N. Cabibbo, Phys. Rev. Letters 10, 531 (1962).

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[6]The vector couplings for hyperon β decay are thus uniquely predicted (for small momentum transfers). For a coupling of the form[3]

$$1/\sqrt{2}G\sin\theta[\bar{\psi}_f\gamma_\mu(f_V+f_A\gamma_5)\psi_i][\bar{\psi}_e\gamma_\mu(1+\gamma_5)\psi_\nu],$$

we have:

$$f_V(\Sigma^-n)=-1,\ f_V(\Sigma^0p)=-1/\sqrt{2},\ f_V(\Lambda p)=-(\tfrac{3}{2})^{1/2},$$

$f_V(\Xi^-\Lambda)=(\frac{3}{2})^{1/2}$, $f_V(\Xi^0\Sigma^+)=1$, $f_V(\Xi^-\Sigma^0)=1/\sqrt{2}$;

where $f_V(\Sigma^- n)$ is the vector coupling for $\Sigma^- \to n$ + leptons. From β decay of O^{14} and Al^{26} we obtain $\cos\theta = 0.980$ or $\sin\theta = 0.20$ [see J. Sakurai, Phys. Rev. Letters 12, 79 (1964)]. From the recent values $R=(1.07\pm0.13)\times10^{-3}$ for the branching ratio and $f_A/f_V=1.03$ for Λ β decay (V. G. Lind, T. O. Binford, M. L. Good, and D. Stern, to be published), we find $|f_V(\Lambda p)|=1.29\pm0.13$ in excellent agreement with $f_V=-(\frac{3}{2})^{1/2}=-1.22$. For $K^0\to\pi^-+e^+ +\nu$ we write a matrix element $(1/\sqrt{2})(G\sin\theta)f(p+q)_\mu \bar{u}_e\gamma_\mu \times(1+\gamma_5)u_\nu$. The prediction is $f=1$. From data on K_2^0 [D. Luers, I. S. Mittra, W. J. Willis, and S. S. Yamamoto, Phys. Rev. 133, B1276 (1964); Proceedings of the Sienna International Conference on Elementary Particles (Società Italiana di Fisica, Bologna, Italy, 1963), Vol. 1, p. 23; for the branching ratio we take a weighted average of 0.56 ± 0.03], we obtain $|f|=0.96 \pm 0.20$. From K^+ data [B. Roe et al., Phys. Rev. Letters 7, 346 (1961)] using $\Delta T=\frac{1}{2}$ rule, we obtain instead $|f|=1.18\pm0.06$, in apparent disagreement with the predicted value and with the K_2^0 data.

[7]The amplitude for $\gamma\to\bar{\Sigma}^0+\Sigma^0$ can be expressed by charge independence in terms of the other $\bar{\Sigma}\Sigma$ amplitudes; the (first-class) amplitude for $\gamma\to\bar{\Sigma}^0+\Lambda$ is equal to the amplitude for $\gamma\to\bar{\Lambda}+\Sigma^0$ and is expressed as a linear combination of the other amplitudes for neutral baryons [Okubo's relation: S. Okubo, Phys. Letters 4, 14 (1963)].

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ERRATUM

NUCLEAR SPIN ORDERING IN ADSORBED He^3. M. H. Lambert [Phys. Rev. Letters 12, 67 (1964)].

Further experiments have shown that the observed specific heat anomaly is not due to spin ordering. A complete report is in preparation.