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Current-Current Interaction and Calculation of the Nonleptonic Baryon Decays*

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Arguments are given why the nonleptonic weak interaction should, in a quark model with neutral-vector-boson strong interaction, be calculable in terms of low-energy contributions, which can be estimated from the knowledge of semileptonic processes. Fair agreement with experiments seems to support this possibility. The suggestion is also made that this model could be very helpful in understanding many properties of electromagnetic and weak interactions.

I. INTRODUCTION

THE universal current-current Hamiltonian for the weak interactions[1] has been extremely useful in explaining leptonic and semileptonic processes. An equally satisfactory understanding of the nonleptonic decays in this framework, however, has not yet been achieved.

Interesting results[2] have, on the other hand, been obtained by introducing a few low-lying intermediate states between the currents, in the current-current Hamiltonian, and using the information available from semileptonic processes. The picture that emerges from such a "saturation" scheme is, as we will review, consistent with experiments. This success is quite surprising. In fact, even if the current-current form is basically "correct," the local product of currents may be too singular to allow meaningful tests via a crude "saturation" approximation. Our experience with the calculation of electromagnetic mass-splittings may also serve as grounds for pessimism. It has been shown that one contribution of low-lying states to the Cottingham formula[3] fails to reproduce even the correct sign of the $\Delta I=1$ electromagnetic (e.m.) mass splittings.[4,5] Such a failure is relevant to the present discussion, because the S-wave decays in the soft-pion limit[6] are related to the parity-conserving (P.C.) matrix elements $\langle B'|H_W{}^{\text{P.C.}}|B\rangle$,[7] which are very similar (except for the missing photon propagator) to $\langle B'|H_{\text{e.m.}}|B\rangle$.

It has been recognized that additional "tadpole" terms,[8] reflecting high-energy contributions, must be present and account for most of the $\Delta I=1$ mass splittings[4,5]; and it has been suggested[8] that the $\Delta I=\frac{1}{2}$ rule in nonleptonic decays should emerge through a similar tadpole mechanism, thus casting severe doubts on low-energy saturation.

A possible interpretation of the "tadpoles" has been suggested by Bjorken.[9] By applying his method to the virtual "Compton-like" amplitudes, one finds in general divergent integrals, both in nonleptonic and e.m. amplitudes. We do not think that the occurrence of such divergencies is disastrous. Motivated by renormalization theory, we take the attitude that when these divergencies occur, they are going to supply us with *incalculable* "renormalization" constants. On the other hand, if such divergencies are not present, the possibility of calculating such amplitudes in terms of low-energy contributions seems to be likely. We would like to emphasize that this is the basic attitude taken in the present investigation.

In Sec. II we show that there exists at least one model of the strong interactions where the "divergent" terms have operator coefficients whose matrix elements vanish between the physical states of the weak decays. Such a privileged model is the quark model, where the interaction is mediated by a massive neutral vector meson coupled to the conserved baryon current, and the $SU_3\otimes SU_3$ chiral invariance of the theory is broken

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2 Y. Hara, Progr. Theoret. Phys. (Kyoto) **37**, 710 (1967); Y. T. Chiu and J. Schechter, Phys. Rev. Letters **16**, 1022 (1966); S. Biswas, A. Kumar and R. Saxena, *ibid.* **17**, 268 (1966); Y. T. Chiu, J. Schechter, and Y. Ueda, Phys. Rev. **150**, 1201 (1966).

3 W. N. Cottingham, Ann. Phys. (N. Y.) **25**, 424 (1963).

4 H. Harari, Phys. Rev. Letters **17**, 1303 (1967).

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6 H. Sugawara, Phys. Rev. Letters **15**, 870 (1965); **15**, 997 (1965); M. Suzuki, *ibid.* **15**, 986 (1965).

7 The Suzuki-Sugawara analysis assumes certain commutation relations between the weak Hamiltonian and the axial charges, which are true both in the JJ and in the intermediate-vector-boson pictures.

8 The term "tadpoles" was introduced by S. Coleman and S. Glashow, Phys. Rev. **134**, B671 (1964), indicating a general dynamical enhancement of a particular channel.

9 J. D. Bjorken, Phys. Rev. **148**, 1467 (1966).

only through mass terms in the Lagrangian.[10] It is worth noticing that this is the only renormalizable model of the strong interactions which guarantees either finiteness or universality of the radiative corrections to semileptonic processess.[11]

II. BJORKEN'S METHOD AND EVALUATION OF DIVERGENT PARTS

A. Intermediate-Vector-Boson Weak Interaction

We write the weak Lagrangian in the form

$$L_W = gJ^\mu(x)W_\mu(x)\,, \tag{1}$$

where J_μ is the Cabibbo current

$$J_\mu(x) = \cos\theta\,(V_\mu{}^{\pi^+}(x) + A_\mu{}^{\pi^+}(x)) + \sin\theta\,(V_\mu{}^{K^+}(x) + A_\mu{}^{K^+}(x))\,, \tag{2}$$

and $W_\mu(x)$ is the vector-boson field whose mass m_W relates the dimensionless coupling constant g to the Fermi coupling constant, via

$$g^2/m_W{}^2 = G/\sqrt{2}\,. \tag{3}$$

This Lagrangian leads to the nonleptonic amplitude

$$T(B\to B'\pi) = g^2\int\frac{d^4k}{(2\pi)^4}\left(g_{\mu\nu} - \frac{k_\mu k_\nu}{m_W{}^2}\right)\frac{T_{\mu\nu}}{-k^2+m_W{}^2}\,, \tag{4}$$

where

$$T_{\mu\nu} = -i\int d^4x\, e^{ikx}\langle B'\pi|T^*(J_\mu{}^\dagger(x)J_\nu(0))|B\rangle \tag{5}$$

and T^* denotes the covariant amplitude which represents the response of the S matrix to the second-order weak vector perturbation.[12]

We now apply Bjorken's analysis to $T_{\mu\nu}$. We analyze first the $k^\mu k^\nu T_{\mu\nu}$ part.[13] By using the chiral algebra we have

$$k^\mu k^\nu T_{\mu\nu} = k^\nu\langle B'\pi|\tilde{J}_\nu(0)|B\rangle + i\int d^3x\, e^{-i\mathbf{k}\cdot\mathbf{x}}\langle B'\pi|[J_0(x),D^\dagger(0)]_{x_0=0}|B\rangle + i\int d^4x\, e^{ikx}\langle B'\pi|T^*(D^\dagger(x),D(0))|B\rangle\,, \tag{6}$$

where $\tilde{J}_\nu(0)$ is a combination of neutral vector and axial currents, and $D(x) = \partial_\mu J^\mu(x)$. The first term integrates to zero by a symmetrical integration over k. The second term yields a quadratic divergence in (4):

$$\frac{g^2}{m_W{}^2}\int\frac{d^4k}{(2\pi)^4}\frac{1}{m_W{}^2-k^2} \times\int d^3x\, e^{-i\mathbf{k}\cdot\mathbf{x}}\langle B'\pi|[J_0(\mathbf{x},0),D^\dagger(0)]|B\rangle\,. \tag{7}$$

Logarithmic divergencies in (5) may arise from the third term in (6) and from $g_{\mu\nu}$ piece in (4), and according to the Bjorken analysis will be given by

$$\frac{(-ig^2)}{(2\pi)^4}\int\frac{d^4k}{m_W{}^2-k^2}\frac{1}{k^2} \times\left(\frac{1}{m_W{}^2}\int\langle B'\pi|[D^\dagger(x),[H,D(0)]]|B\rangle d^3x\, e^{-i\mathbf{k}\cdot\mathbf{x}} + \int d^3x\, e^{-i\mathbf{k}\cdot\mathbf{x}}\langle B'\pi|[J_\mu{}^\dagger(x),[H,J^\mu(0)]]|B\rangle\right), \tag{8}$$

where H is the Hamiltonian of the system.

We now evaluate (7) and (8) in the framework of the above-mentioned quark model, which is characterized by the Lagrangian

$$L = \bar{q}(-i\gamma\cdot\partial + g\mathbf{B} + M)q + L_B\,, \tag{9}$$

where L_B refers to the vector boson B_μ part, and M is a numerical quark mass matrix. In such a model the Cabibbo current has the form

$$J_\mu(x) = \bar{q}(x)\gamma_\mu(1+\gamma_5)\lambda^+ q(x)\,, \tag{10}$$

where

$$\lambda^+ = \begin{bmatrix} 0 & \cos\theta & \sin\theta \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{bmatrix}$$

and

$$D(x) = i\bar{q}(x)\mathfrak{M}q(x)\,, \tag{11}$$

where script letters here and in the following denote linear combinations of products of λ matrices with 1, and γ_5.

The equal-time commutator in (7) is now

$$[J_0(\mathbf{x},0),D^\dagger(\mathbf{0},0)] = \delta^3(\mathbf{x})i\bar{q}\mathfrak{M}'q\,. \tag{12}$$

In particular, the part relevant to nonleptonic decays ($\Delta S=1$) is

$$\bar{q}\mathfrak{M}'^{(\Delta S=1)}q = \bar{q}(\alpha\lambda_6 + \beta\lambda_7\gamma_5)q\,, \tag{13}$$

where α and β are constants.

We analyze next the logarithmically divergent part (8), and consider first the matrix element

$$\int d^3x\langle B'\pi|[D^\dagger(\mathbf{x},0),[H,D(0)]]|B\rangle\,.$$

Using the Hamiltonian H corresponding to (9), and the

[10] Some properties of this model have been considered by M. Gell-Mann, Phys. Rev. **125**, 1064 (1962) and by J. D. Bjorken, Ref. 9.

[11] C. G. Callan, Phys. Rev. **169**, 1175 (1968); G. Preparata and W. I. Weisberger, Phys. Rev., this issue, **175**, 1965 (1968).

[12] T^* consists in general of the time-ordered product of the currents and additional "Schwinger" terms. Here, and in the following, we assume that no $\Delta S=1$ "Schwinger" terms are present so that we can ignore them throughout our discussion.

[13] We have found that a similar analysis was carried out by V. S. Mathur and P. Olesen, Phys. Rev. Letters **20**, 1527 (1968).

expression (11) for $D(x)$, we find that

$$\int d^3x \langle B'\pi | [D^\dagger(x),[H,D(0)]]_{x_0=0} | B\rangle$$
$$= \int d^3x \langle B'\pi | [\bar{q}\mathfrak{M}\gamma^i(-i\overleftrightarrow{\nabla}_i+gB_i)q+\bar{q}\mathfrak{M}'q] | B\rangle . \quad (14)$$

The first term in (14) can be written in the form

$$\bar{q}\mathfrak{M}\gamma^i(-i\boldsymbol{\nabla}_i+gB_i)q = -\bar{q}\mathfrak{M}(-i\gamma\cdot\partial+g\mathbf{B})q + \eta_\mu\eta_\nu\bar{q}\mathfrak{M}\gamma^\mu(-i\partial^\nu+gB^\nu)q , \quad (15)$$

where $\eta_\mu \equiv (\mathbf{0},1)$.

The covariant form corresponding to (15) is obtained by the substitution[9] $\eta_\mu\eta_\nu \to k_\mu k_\nu/k^2$, yielding the following contribution to (8):

$$-\frac{ig^2}{(2\pi)^4}\int \frac{d^4k}{k^2(m_W{}^2-k^2)} \frac{3}{4m_W{}^2} \langle B'\pi | \bar{q}\mathfrak{M}''q | B\rangle . \quad (16)$$

A similar calculation applies to the second term in (8).

The crucial observation is that within the framework of this model the $S=1$ scalar and pseudoscalar densities can be expressed as four-divergences of the corresponding current operators.[13a]

Matrix elements of these densities therefore vanish between states of equal energy and momentum (provided such operators are, as they indeed are, nonsingular). As a consequence we find that the coefficients of both the quadratic and logarithmic divergencies [Eqs. (13) and (16)] vanish for the physical decay process.

This is different from what one finds, within this same approach, for the second-order e.m. mass shifts. There the coefficient of the leading logarithmic divergence[9] is

$$[J_\mu{}^{\text{e.m.}},[H,J_{\text{e.m.}}{}^\mu]] \propto \bar{q}Q^2q , \quad (16')$$

when Q is the 3×3 charge matrix. This density *cannot* be written as a four-divergence, and therefore its relevant matrix elements will in general be nonvanishing. Indeed, if we wish to attribute the prominent $\Delta I=1$ mass differences to such tadpole terms, these matrix elements should be quite large, as we will discuss later on.

B. Current-Current Interaction

We may obtain the current-current interaction formally from (4) by letting $m_W{}^2\to\infty$,[14] giving

$$T(B\to B'\pi) = \frac{G}{\sqrt{2}}\int \frac{d^4k}{(2\pi)^4} T_{\mu\mu} . \quad (17)$$

[13a] *Footnote added in proof*. After completion of this work, we learned that a similar observation has been made by C. Bouchiat, J. Iliopoulos, and J. Prentki, Nuovo Cimento **56A**, 1150 (1968).

[14] We must, however, warn that this procedure may be meaningless due to the possible bad behavior of the theory at small distances.

In addition to the quadratic divergence in (17) which by the above argument may be absent, there are logarithmic divergencies. If the Bjorken analysis can be pushed this far, the coefficient of this divergence is $[J_\mu{}^\dagger,(d^3/dt^3)J_\mu]$. Evaluation of this commutator gives in addition to quark densities (13) expressions of the form

$$g^2 m\bar{q}(\partial_\mu B_\nu - \partial_\nu B_u)\sigma_{\mu\nu}\mathcal{T}q , \quad (18)$$

which are quite different from a quark density, and their matrix elements may well be much smaller than those of the e.m. tadpole (16′). This, together with the fact that the leading quadratic divergence is absent leaves open the possibility that the unknown appropriately cut-off high-energy contribution to the nonleptonic amplitude is relatively small compared with the calculable low-energy contribution. This may serve as a motivation for the analysis of the low-energy part of the weak amplitudes to which we now proceed.

III. LOW-ENERGY CONTRIBUTIONS

Since we are interested in the region of small virtual momenta ($k^2 \ll$ experimental lower limit of $m_W{}^2$), Eq. (17) is an adequate starting point for the calculation. The kinematics of the amplitude $T_{\mu\nu}$, which appears in Eq. (17), is shown in Fig. 1.

In order to evaluate the contribution of the low-lying states in (17), it is most useful to write down a Cottingham-like formula,[3]

$$T_{\text{low}}(B\to B\pi) = \frac{G}{\sqrt{2}}\frac{1}{32\pi^3}\int_0^\infty d(-k^2)\int^{\nu_{\max}} (-k^2+\nu^2)^{1/2} \times \int_{-1}^{+1} dz\, \text{Im}T(k^2,\nu,z) , \quad (19)$$

where

$$\text{Im}T(k^2,\nu,z) = \tfrac{1}{2}(2\pi)^4 \sum_n \delta^4(p_n-p-k) \times \langle B'\pi | J_\mu{}^\dagger(0) | n\rangle\langle n | J_\mu(0) | B\rangle , \quad (20)$$

and $\nu = k\cdot p_B/m_B$, $z=\mathbf{k}\cdot\mathbf{q}/|\mathbf{k}|\,|\mathbf{q}|$, where $\mathbf{q}$ is the pion momentum in the rest frame of the decaying baryon B.

$\nu_{\max}$ is a cutoff energy defined by the condition

$$(k+p_B)^2 \leq M^2 \quad (21)$$

and M^2 will be specified below.

In (19) there are three different types of intermediate states n: n_s containing no disconnected parts (Fig. 2a),

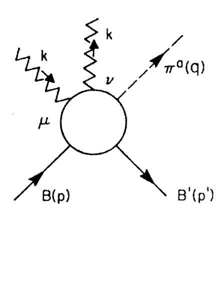


FIG. 1. Kinematics of the amplitude $T_{\mu\nu}$.

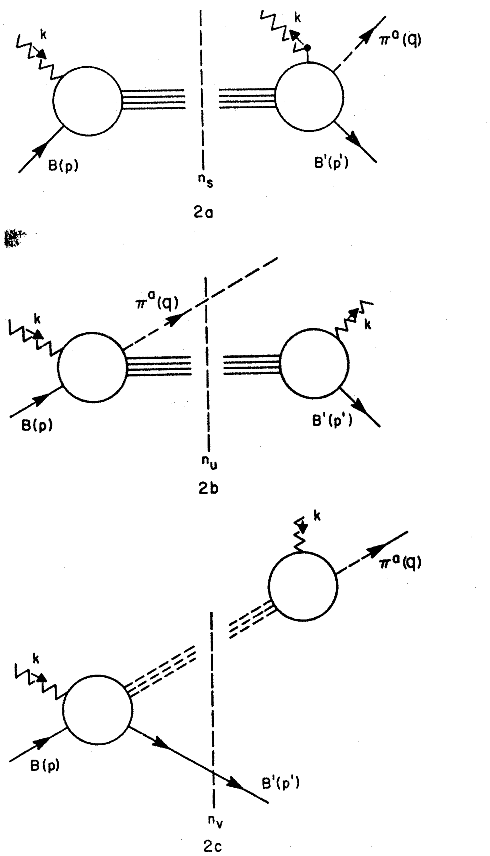


FIG. 2. Intermediate states contributing to formula (19).

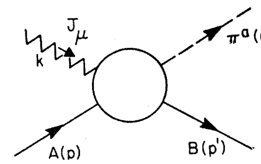
n_u containing a disconnected pion (Fig. 2b), and n_v containing a disconnected baryon (Fig. 2c).

In the following we will keep only the two lowest SU_3 multiplets in the s, u, and v channels. This is effectively achieved by choosing M of Eq. (21) to be slightly above $m_V+m_B \simeq 1800$ MeV. The final result will in fact not be particularly sensitive to the value of M^2.

Within this approximation the typical contributions to Im T of Eq. (20) involve weak-current form factors and weak meson production amplitudes. Following earlier calculations,[2] we take the weak-current baryon matrix elements from the fit to the Cabibbo theory and use universal dipole form factors. Lacking detailed experimental information about weak meson production amplitudes, we use the soft-pion limit.

In previous estimates of S-wave decays[2,6] partially conserved axial-vector current (PCAC) and the soft-pion limit were used at the outset, restricting the saturation procedure to the matrix elements of the weak Hamiltonian between single baryon states. In the framework of the present model, the PCAC extrapolation may be dangerous, since in the soft-pion limit the nonleptonic Hamiltonian carries effectively a momentum transfer q and the matrix elements of the quark densities [(13) and (16)] do not vanish any more. In practice, as we shall show later, the difference between this and our approach of using PCAC in the weak meson production amplitudes is relatively small.

FIG. 3. The general weak pion-production amplitude.



We consider the general weak meson production amplitudes $T_\mu{}^a(k,q)$ (see Fig. 3). If we assume that $T_\mu(k,q)-T_\mu{}^B(k,q)$ (where $T_\mu{}^B$ is the Born amplitude) is a smooth function of q when $q\to 0$,[15] we may write the amplitude

$$T_\mu{}^a(k,q)\simeq T_\mu{}^{aB}(k,q)+[T_\mu{}^a(k,0)-T_\mu{}^{aB}(k,0)]. \quad (22)$$

Making use of PCAC and the chiral $SU_3\otimes SU_3$ algebra, we have

$$T_\mu{}^a(k,0)-T_\mu{}^a(k,0)^B=-\frac{\sqrt{2}}{f_\pi}\langle B'(p')|\tilde{J}_\mu{}^a(0)|A(p)\rangle + \lim_{q\to 0}\left[\frac{\sqrt{2}}{f_\pi}iq^\alpha\int d^4x\, e^{iqx}\langle B|T^*[A_\alpha{}^a(x)J_\mu(0)]|B\rangle - T_\mu{}^a(k,q)^B\right], \quad (23)$$

where $\tilde{J}_\mu{}^a(0)$ is the result of the equal-time commutator $[A_0{}^a(x),J_\mu(0)]$. If we write $T_\mu{}^a(k,q)^B$ using a derivative coupling, the last term in (23) is identically zero when $q\to 0$. So in the soft-pion limit we have

$$T_\mu{}^a(k,q)=T_\mu{}^a(k,q)^B-(\sqrt{2}/f_\pi)\langle B|\tilde{J}_\mu{}^a(0)|A\rangle. \quad (24)$$

Equipped with Eq. (24) for the weak pion-production amplitude and with the usual weak-current form factors we now evaluate the contribution to the nonleptonic amplitude from $n_s=n_u=$ baryon octet and decuplet diagrams.

For the S-wave decays the dominant contribution comes from the equal-time commutator term in Eq. (24), so that we recover formally an S-wave amplitude identical with that obtained by direct application of PCAC[3] to the nonleptonic Hamiltonian,

$$S(B\to B'\pi^a)\cong(\sqrt{2}/f_\pi)\langle B'|\tilde{H}^a|B\rangle, \quad (25)$$

where $\tilde{H}_W$ is an effective nonleptonic Hamiltonian

TABLE I. We define $\mathfrak{M}=\bar{u}(p')(A-B\gamma_5)u(p)$ as the decay amplitude. The amplitudes satisfy the $\Delta I=\frac{1}{2}$ rule.

Decay	10^6A expt.[a]	Calculated	10^6B expt.[a]	Calculated
$\Lambda_0{}^-$	0.335±0.004	0.285	2.3±0.1	1.48
$\Sigma_+{}^+$	0.001±0.006	0	4.2±0.08	2.6
$\Sigma_0{}^+$	0.338±0.030	0.40	2.6±0.4	1.85
$\Sigma_-{}^-$	0.405±0.003	0.57	(−3.4±8.5)×10⁻²	0
$\Xi_-{}^-$	0.440±0.006	0.50	1.47±0.12	1.28

[a] The experimental figures have been taken from N. Cabibbo, in *Proceedings of the Thirteenth Annual International Conference on High-Energy Physics, Berkeley, 1966* (University of California Press, Berkeley, 1967), p. 29.

[15] We follow the procedure first used by L. S. Brown and C. M. Sommerfield, Phys. Rev. Letters **16**, 751 (1966).

evaluated by saturating a current-current Hamiltonian by the low-lying octet and decuplet states. In the SU_3-symmetry limit one can write

$$\langle B'|\tilde{H}_a|B\rangle = DD_{B'B}{}^a + FF_{B'B}{}^a + TT_{B'B}{}^a, \quad (26)$$

with D, F, T referring to the D and F octet, and **27** coupling, respectively. An estimate of $\langle B'|\tilde{H}_a|B\rangle$ was made by Hara[2], who obtained[16]

$$\begin{aligned} D &= -3.2\times10^{-5}\ \text{MeV},\\ F &= 3.8\times10^{-5}\ \text{MeV},\\ T &= -0.1\times10^{-5}\ \text{MeV}. \end{aligned} \quad (27)$$

Equations (25) and (28) yield a reasonable prediction of all S-wave decays. (See Table I).

In particular the $\Delta I=\frac{1}{2}$ selection rule seems to emerge in a dynamical way, because of mutual cancellation of octet and decuplet contributions.

The corrections to Eqs. (25) and (27) which arise from the Born term in Eq. (24), and the so far neglected n_ν diagrams, have been estimated. We find that such contributions give at most 20–30% corrections.

We turn now to consider the P-wave nonleptonic amplitudes. Neglecting again the n_ν-type diagrams, and the equal-time commutator term in the weak-production amplitude (24), we obtain from the Born diagrams effectively the results which have been previously obtained in Ref. 17, where we have to use for the "spurion" matrix elements the values of Eqs. (27). As is shown in Table I, this gives a substantially correct picture of the P-wave amplitudes.[18] We found that the neglected pieces (equal-time commutators and n_ν diagrams) give small corrections without altering the picture. It is however, interesting to notice that the possible effect of a P_{11} resonance in the P-wave weak production amplitude will add a contribution which is qualitatively of the right structure to improve agreement with experiment. Also here the $\Delta I=\frac{1}{2}$ rule is dynamically brought in through the matrix elements of the weak "Hamiltonian" between baryon states.

IV. CONCLUSIONS

We have shown that in a particular field-theoretical model, a justification can be given to the "saturation" approach to the nonleptonic interaction. A review of its implications has shown that, within the approximations made, it describes correctly the main features of the nonleptonic baryon decays.

That this situation is significantly different from what we have in the case of the e.m. mass difference, we think is supported by the following argument. Let us consider the S-wave amplitudes; if the deviations between the calculated and experimental values for the S-wave decays are interpreted in terms of additional "tadpole" contributions, we find for the magnitude of the tadpoles,

$$|F_W|+|D_W|\simeq 1.8\times10^{-5}\ \text{MeV}.$$

The analysis of the e.m. $\Delta I=1$ mass difference indicates[5] that the low-energy contributions need be augmented by a tadpole term with a magnitude

$$(F+D)_{\text{e.m.}}=-2.08\ \text{MeV},$$
$$F/D=-1.8,$$

giving

$$|F|+|D|=6.3\ \text{MeV}.$$

If we adopt Bjorken's interpretation[9] of the tadpole contributions, we would in general expect a ratio

$$r_{\text{th}}=\frac{(|D|+|F|)_W}{(|D|+|F|)_{\text{e.m.}}}=\frac{G\sin\theta\cos\theta}{\sqrt{2}} \times\int^{\Lambda_{\text{wk}}^2} dk^2(g_{VV}+g_{AA}) \bigg/ \int^{\Lambda_{\text{e.m.}}^2} dk^2\frac{g_{\text{e.m.}}}{k^2}$$

and $g_{VV}=g_{AA}=g_{\text{e.m.}}$ as a consequence of the "universality" of tadpoles. Using for the cutoffs the values $\Lambda_{\text{wk}}\simeq10$ BeV and $\Lambda_{\text{e.m.}}\simeq100$ BeV, we find that r_{th} is smaller than "r_{exp}" by almost an order of magnitiude and the situation is of course much worsened ,f we increase the value of Λ_{wk}^2. In spite of the crudeness of the argument we think that this is a fairly meaningful indication of the difference between the e.m. and the weak case, which seems to be incorporated in the model discussed previously.

The question now is: What have we learned from all this? We think optimistically that from the preceding discussion may emerge the basic adequacy of the current-current picture for low-energy nonleptonic interactions, and the interesting role played by this particular quark model in supplying us with information going beyond the realm of "current algebra." We think that investigating other features of such a model could be helpful in understanding the weak and e.m. interactions of the hadrons.

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[16] Uncertainties in F, D, and T of Eqs. (27) arise from the insufficient experimental information on the vertices $\langle B|J_\mu|B\rangle$ and $\langle B|J_\mu|\Delta\rangle$. The forms used by Hara for these vertices are rather simple and appealing. In particular the universal dipole form factor $(m_V^2/k^2-m_V^2)^2$ with $m_V^2\simeq0.71$ BeV^2 was used in all cases. We found only small variations ($\simeq$15%) when choosing different form factors, incorporating the correct static values (including the radii).

[17] C. Itzykson and M. Jacob, Nuovo Cimento **48A**, 655 (1967).

[18] The results for the P-wave decays are not as significant as those for the S-wave, due mainly to some subtle cancellations in the Born diagrams, which on the other hand are particularly sensitive to mass SU_3-breaking effects.