coupling constant, number of inelastic channels, phase space, statistical counting, etc. Dynamics may then be needed only for explaining a particular channel whose contribution tends to vanish as energy increases. Therefore, it would appear that the low-energy resonance region has greater dynamical content than the high-energy domain studied in this paper. Our model also suggests that the analysis of inelastic collisions will be simplified if the quasielastic part is removed. Concerning experiments in the future, it is desirable to measure the energy spectra for $\pi^{\pm}p$, $K^{\pm}p$, and $\bar{p}p$ scattering so that our calculated $R(s)$ can be checked, at least qualitatively, by experiments. Finally, we notice an interesting prediction on the width of $\pi\pi$ elastic diffraction scattering. According to our model, the quasielastic scattering for all processes is mainly due to the OPE diagram. However, such a diagram is forbidden in $\pi\pi$ scattering by G parity. We expect that the contribution due to other diagrams is very small. As a consequence, the quasielastic cross section for $\pi\pi$ will be very small. If the $\pi\pi$ total cross section behaves like all other total cross sections (i.e., does not increase with increasing energy), we predict that *shrinkage of the $\pi\pi$ diffraction peak is impossible.*

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Gauge Theory of Strong Interactions*

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A generalization of gauge invariance in strong interactions is given. This leads naturally to a structure of $[U(3)]^4$ with the gauge fields corresponding to the quantum numbers $J^{PC}=1^{++}, 1^{+-}, 1^{--}, 1^{-+}$. Currents of both first- and second-class types are present. A tentative assignment of the gauge fields to possible resonances is made.

1. INTRODUCTION

An important method in considering the symmetries of strong interactions has been to use gauge invariance of a Lagrangian to infer the existence of conserved currents. The gauge invariance is then broken by mass terms for example, resulting in partial-conservation laws. When these conserved or partially conserved currents are assumed to be the currents which describe the electromagnetic and weak interactions, several deep results follow.[1] The spin-1 gauge fields that are inferred by a local gauge invariance of a Lagrangian have been considered to be important because they must belong to the regular representation of the symmetry group. With the recent work of Kroll, Lee, and Zumino[2] and of others,[3] they too have been shown to be important in a dynamical sense in describing the electromagnetic and weak interactions of the hadrons. Since broken $SU(3)$ is well established as an algebraic symmetry, it would seem important to investigate the full consequences of an approach based on successive approximations to a completely exact local gauge invariance. Not only is such a method an important and fundamental way of arriving at the interaction of fields, but it has, in the past, also been used to infer the existence of pionic resonances which are identified with the gauge fields.[4]

In this paper, we examine a generalized application of gauge invariance,[5–7] and arrive at an underlying group structure of $[U(3)]^4$. A tentative assignment of the gauge fields to possible resonances is made.

* Supported in part by the National Research Council of Canada.

[1] S. Adler and R. Dashen, *Current Algebras* (W. A. Benjamin, Inc., New York, 1968).

[2] N. M. Kroll, T. D. Lee, and B. Zumino, Phys. Rev. **157**, 1376 (1967).

[3] T. D. Lee and B. Zumino, Phys. Rev. **163**, 1667 (1967); T. D. Lee, S. Weinberg, and B. Zumino, Phys. Rev. Letters **18**, 1029 (1967).

[4] J. J. Sakurai, Ann. Phys. (N.Y.) **11**, 1 (1960).

[5] The group $[U(3)]^4$ obtained here is not the same as the group obtained previously by the authors listed in Refs. 6 and 7. The transformations which these authors considered may still be applied along with the generalization given here and would result in a doubling of the number of gauge fields.

[6] P. G. O. Freund and Y. Nambu, Phys. Rev. Letters **12**, 714 (1964).

[7] A. Salam and J. C. Ward, Phys. Rev. **136**, B763 (1964). We shall follow the notation of these authors.

2. EXTENSION OF GAUGE PRINCIPLE

We shall work, for simplicity, with the free massless-fermion Lagrangian[7,8]

$$L_0=-\mathrm{tr}(\bar{\psi}\gamma_\mu\partial_\mu\psi)\,, \tag{2.1}$$

where

$$\psi=\sqrt{2}T^\alpha\psi^\alpha,\quad \alpha=0,1,\ldots,8$$

in the mixed-tensor representation of $U(3)$. The nine Hermitian matrices T^α are related to those of Gell-Mann[9] by

$$T^\alpha=\tfrac{1}{2}\lambda^\alpha.$$

We now rewrite Eq. (2.1) in a manifestly C-invariant way:

$$L_0=-\tfrac{1}{2}\,\mathrm{tr}\{\bar{\psi}\gamma_\mu\partial_\mu\psi+\bar{\chi}\gamma_\mu\partial_\mu\chi\}\,; \tag{2.2}$$

where, in the Dirac-Pauli representation of the γ matrices,

$$\chi\equiv\psi^C=\eta C\bar{\psi}^T\,,$$
$$\bar{\chi}=-\eta^*\psi^TC^{-1}\,,$$

with $C=\gamma_4\gamma_2$ and $|\eta|^2=1$. We may also split ψ and χ into left and right components,

$$\psi_{L,R}=\tfrac{1}{2}(1\pm\gamma_5)\psi\,,$$
$$\chi_{L,R}=\tfrac{1}{2}(1\mp\gamma_5)\chi\,.$$

The free massless Dirac Lagrangian then becomes

$$L_0=-\tfrac{1}{2}\,\mathrm{tr}(\bar{\psi}_L\gamma_\mu\partial_{1\mu}\psi_L+\bar{\psi}_R\gamma_\mu\partial_{2\mu}\psi_R+\bar{\chi}_L\gamma_\mu\partial_{3\mu}\chi_L+\bar{\chi}_R\gamma_\mu\partial_{4\mu}\chi_R)\,, \tag{2.3}$$

where the suffix i on $\partial_{i\mu}$ is formal, i.e., $\partial_{i\mu}\equiv\partial/\partial x_\mu$, $i=1,\ldots,4$.

The more usual form of local gauge invariance is now extended by demanding that *each* of the fields $\psi_{L,R}$ and $\chi_{L,R}$ may be transformed separately and independently, i.e.,

$$\psi_L'=P_0P\psi_LP^{-1}\,, \tag{2.4a}$$
$$\psi_R'=Q_0Q\psi_RQ^{-1}\,, \tag{2.4b}$$
$$\chi_L'=R_0R\chi_LR^{-1}\,, \tag{2.4c}$$
$$\chi_R'=S_0S\chi_RS^{-1}\,, \tag{2.4d}$$

where

$$P_0=\exp(i\epsilon_1^0)\,,$$
$$P=\exp(iT^j\epsilon_1^j)\,,$$
$$Q_0=\exp(i\epsilon_2^0)\,,$$
$$Q=\exp(iT^j\epsilon_2^j)\,,\quad\text{and so on}.$$

[8] Although all of our results could be derived with the simple quark representation of $U(3)$, we use the mixed-tensor representation, since it can be extended quite easily to include the spin-0 mesons. Since the pseudoscalar-meson nonet is self-conjugate, we need to introduce a new nonet of spin-0 mesons, and then apply our gauge principle to the non-Hermitian matrices $M=M_1+iM_2$ and $N=\mathcal{C}M\mathcal{C}^{-1}=M^\dagger$. The details are similar to the discussion in Ref. 7 except that only vector-gauge fields are introduced. This means that the tadpole mechanism used there will not give different contributions to the mass terms of the two sets of particles in this theory.

[9] M. Gell-Mann, Phys. Rev. **125**, 1067 (1962).

This is an extension of the method of arriving at chiral $SU(3)\times SU(3)$ symmetry by gauge invariance, in which case only the fields ψ_L and ψ_R would be gauged separately. It amounts to applying unitary transformations to all components of the Dirac four-spinor.

If we require that the Lagrangian L_0 be invariant under the four *independent* transformations (2.4) where the ϵ_n, $n=1,2,3,4$, depend on x_μ, then the derivatives $\partial_{1\mu},\cdots\partial_{4\mu}$ are replaced by the "covariant" derivatives

$$D_{1\mu}\psi_L=\partial_\mu\psi_L+ig_1[X_{1\mu},\psi_L]+ig_1^0X_{1\mu}^0\psi_L\,, \tag{2.5a}$$
$$D_{2\mu}\psi_R=\partial_\mu\psi_R+ig_2[X_{2\mu},\psi_R]+ig_2^0X_{2\mu}^0\psi_R\,, \tag{2.5b}$$
$$D_{3\mu}\chi_L=\partial_\mu\chi_L+ig_3[X_{3\mu},\chi_L]+ig_3^0X_{3\mu}^0\chi_L\,, \tag{2.5c}$$
$$D_{4\mu}\chi_R=\partial_\mu\chi_R+ig_4[X_{4\mu},\chi_R]+ig_4^0X_{4\mu}^0\chi_R\,, \tag{2.5d}$$

where $X_{n\mu}=\sqrt{2}T^iX_{n\mu}{}^i$. We have thus introduced 36 spin-1 fields.

For the Lagrangian (2.3) to be invariant under the transformations (2.4), we require

$$(D_{1\mu}\psi_L)'=P_0P(D_{1\mu}\psi_L)P^{-1}\,,$$

with analogous expressions for $D_{2\mu}\psi_R$, $D_{3\mu}\chi_L$, and $D_{4\mu}\chi_R$, in terms of the matrices Q, R, and S. This leads to the following transformation character for the fields $X_{1\mu}$, $X_{2\mu}$, $X_{3\mu}$, and $X_{4\mu}$:

$$X_{1\mu}'=PX_{1\mu}P^{-1}-(i/g_1)P\partial_\mu P^{-1}\,, \tag{2.6a}$$
$$X_{1\mu}^{0\prime}=X_{1\mu}^0-(1/g_1^0)\partial_\mu\epsilon_1^0\,, \tag{2.6b}$$
$$X_{2\mu}'=QX_{2\mu}Q^{-1}-(i/g_2)Q\partial_\mu Q^{-1}\,, \tag{2.6c}$$
$$X_{2\mu}^{0\prime}=X_{2\mu}^0-(1/g_2^0)\partial_\mu\epsilon_2^0\,, \tag{2.6d}$$
$$X_{3\mu}'=RX_{3\mu}R^{-1}-(i/g_3)R\partial_\mu R^{-1}\,, \tag{2.6e}$$
$$X_{3\mu}^{0\prime}=X_{3\mu}^0-(1/g_3^0)\partial_\mu\epsilon_3^0\,, \tag{2.6f}$$
$$X_{4\mu}'=SX_{4\mu}S^{-1}-(i/g_4)S\partial_\mu S^{-1}\,, \tag{2.6g}$$
$$X_{4\mu}^{0\prime}=X_{4\mu}^0-(1/g_4^0)\partial_\mu\epsilon_4^0\,. \tag{2.6h}$$

Note that each of the fields $X_{n\mu}$ transform independently as a representation of $SU(3)$.

We now define the field strength $Z_{n\mu\nu}$ of the field $X_{n\mu}$ by

$$Z_{n\mu\nu}=\partial_\mu X_{n\nu}-\partial_\nu X_{n\mu}+ig_n[X_{n\mu},X_{n\nu}]\,,$$
$$Z_{n\mu\nu}^0=\partial_\mu X_{n\nu}^0-\partial_\nu X_{n\mu}^0\,,\quad n=1,2,3,\text{ and }4\,,$$

and write

$$Z_{n\mu\nu}=\sqrt{2}T^iF_{n\mu\nu}{}^i\,,$$

where

$$F_{n\mu\nu}{}^i=\partial_\mu X_{n\nu}{}^i-\partial_\nu X_{n\mu}{}^i-\sqrt{2}f^{ijk}X_{n\mu}{}^jX_{n\nu}{}^k\,. \tag{2.7}$$

Then the Lagrangian

$$L=-\tfrac{1}{2}\,\mathrm{tr}\{\bar{\psi}_L\gamma_\mu D_{1\mu}\psi_L+\bar{\psi}_R\gamma_\mu D_{2\mu}\psi_R+\bar{\chi}_L\gamma_\mu D_{3\mu}\chi_L+\bar{\chi}_R\gamma_\mu D_{4\mu}\chi_R\}-\tfrac{1}{4}(F_{n\mu\nu}{}^iF_{n\mu\nu}{}^i+Z_{n\mu\nu}^0Z_{n\mu\nu}^0)\,, \tag{2.8}$$

where all indices are summed, is invariant under the local gauge transformations (2.4) and (2.6).

3. PARITY AND CHARGE-CONJUGATION INVARIANCE

By demanding local gauge invariance of the Lagrangian under the transformations (2.4) and (2.6), we have introduced interactions between the spin-1 fields and the baryons, and between the spin-1 fields themselves. The latter case is the well-known result of Yang and Mills[10] that the currents corresponding to the spin-1 fields are nonlinear in the spin-1 fields. In this section we focus our attention on the linear interaction terms in the Lagrangian (2.8), which may be written

$$\begin{aligned}L_{\text{int}}=&-\tfrac{1}{2}i\ \text{tr}\{g_1\bar{\psi}_L\gamma_\mu[X_{1\mu},\psi_L]+g_2\bar{\psi}_R\gamma_\mu[X_{2\mu},\psi_R]\\&+g_3\bar{\chi}_L\gamma_\mu[X_{3\mu},\chi_L]+g_4\bar{\chi}_R\gamma_\mu[X_{4\mu},\chi_R]\}\\&-\tfrac{1}{2}i\ \text{tr}\{g_1{}^0\bar{\psi}_L\gamma_\mu X_{1\mu}{}^0\psi_L+g_2{}^0\bar{\psi}_R\gamma_\mu X_{2\mu}{}^0\psi_R\\&+g_3{}^0\bar{\chi}_L\gamma_\mu X_{3\mu}{}^0\chi_L+g_4{}^0\bar{\chi}_R\gamma_\mu X_{4\mu}{}^0\chi_R\}.\end{aligned}\tag{3.1}$$

A. Charge-Conjugation Invariance

For the Lagrangian L_{int} to be invariant under charge conjugation, we must have

$$g_1=g_4 \quad\text{and}\quad \mathcal{C}X_{1\mu}\mathcal{C}^{-1}=X_{4\mu},$$
$$g_2=g_3 \quad\text{and}\quad \mathcal{C}X_{2\mu}\mathcal{C}^{-1}=X_{3\mu},$$

with similar results for the singlet terms. We can now define eigenstates of C by

$$Y_\mu{}^{(\pm)}=\tfrac{1}{2}(X_{1\mu}\pm X_{4\mu}),\tag{3.2a}$$

$$Z_\mu{}^{(\pm)}=\tfrac{1}{2}(X_{2\mu}\pm X_{3\mu}),\tag{3.2b}$$

where

$$\mathcal{C}Y_\mu{}^{(\pm)}\mathcal{C}^{-1}=\pm Y_\mu{}^{(\pm)},$$
$$\mathcal{C}Z_\mu{}^{(\pm)}\mathcal{C}^{-1}=\pm Z_\mu{}^{(\pm)}.$$

B. Invariance under Parity

For parity conservation, two cases are possible:

(i) $\quad g_1=g_2 \quad\text{and}\quad \mathcal{P}X_{1i}\mathcal{P}^{-1}=-X_{2i},\quad \mathcal{P}X_{3i}\mathcal{P}^{-1}=-X_{4i};$

(ii) $\quad g_1=-g_2 \quad\text{and}\quad \mathcal{P}X_{1i}\mathcal{P}^{-1}=X_{2i},\quad \mathcal{P}X_{3i}\mathcal{P}^{-1}=X_{4i}.$

Combinations of $Y_\mu{}^{(\pm)}$ and $Z_\mu{}^{(\pm)}$ transforming as vectors and axial vectors under P may now be made in each case.

Case (i). Let

$$Y_\mu{}^{(\pm)}=V_\mu{}^{(\pm)}+A_\mu{}^{(\pm)},\quad Z_\mu{}^{(\pm)}=V_\mu{}^{(\pm)}-A_\mu{}^{(\pm)}.$$

Then

$$\mathcal{P}V_i{}^{(\pm)}\mathcal{P}^{-1}=-V_i{}^{(\pm)},\quad \mathcal{P}A_i{}^{(\pm)}\mathcal{P}^{-1}=A_i{}^{(\pm)}.$$

Case (ii). Let

$$Y_\mu{}^{(\pm)}=V_\mu{}'^{(\pm)}+A_\mu{}'^{(\pm)},\quad Z^{(\pm)}=-V_\mu{}'^{(\pm)}+A_\mu{}'^{(\pm)}.$$

Then

$$\mathcal{P}V_i{}'^{(\pm)}\mathcal{P}^{-1}=-V_i{}'^{(\pm)},\quad \mathcal{P}A_i{}'^{(\pm)}\mathcal{P}^{-1}=A_i{}^{(\pm)}.$$

[10] C. N. Yang and R. L. Mills, Phys. Rev. **96**, 191 (1954).

The Lagrangian we are considering is also invariant under the chirality, or γ_5, transformation defined by

Case (i):

$$\psi_R\to\psi_L,\quad \chi_R\to\chi_L,\quad V^{(\pm)}\to V^{(\pm)},$$

and

$$A^{(\pm)}\to -A^{(\pm)};$$

Case (ii):

$$\psi_R\to\psi_L,\quad \chi_R\to\chi_L,\quad V'^{(\pm)}\to V'^{(\pm)},$$

and

$$A'^{(\pm)}\to -A'^{(\pm)},$$

so that we have a conserved quantum number $+1$ for processes involving V_μ fields, and -1 for processes involving A_μ fields. When we add mass terms to the Lagrangian, this quantum number will no longer be conserved.

The linear part of the interaction Lagrangian now reads [in case (i)]

$$\begin{aligned}L_{\text{int}}=\tfrac{1}{2}ig\ \text{tr}\{&\bar{\psi}\gamma_\mu[V_\mu{}^{(-)},\psi]-\bar{\chi}\gamma_\mu[V_\mu{}^{(-)},\chi]\\&+\bar{\psi}\gamma_\mu\gamma_5[A_\mu{}^{(+)},\psi]+\bar{\chi}\gamma_\mu\gamma_5[A_\mu{}^{(+)},\chi]\\&+(\bar{\psi}\gamma_\mu\psi-\bar{\chi}\gamma_\mu\chi)V_\mu{}^{0(-)}\\&+(\bar{\psi}\gamma_\mu\gamma_5\psi+\bar{\chi}\gamma_\mu\gamma_5\chi)A_\mu{}^{0(+)}\}.\end{aligned}\tag{3.3}$$

There are no linear fermion-interaction terms involving the fields $V_\mu{}^{(+)}$ and $A_\mu{}^{(-)}$ in this model. This is to be expected, for the currents corresponding to these fields are "second class," and are coupled to derivative terms.[11]

4. SYMMETRY BREAKING

The work of Kroll, Lee, and Zumino[2] has shown that although local gauge invariance may be a useful concept in arriving at the underlying group structure, it is unnecessary (and indeed is not wanted) in setting up a model field theory for low-energy processes involving the electromagnetic current. To arrive at an algebra of fields for the group structure $U(3)\times U(3)\times U(3)\times U(3)$ that we have obtained above, it is necessary to add a mass term to the Lagrangian. Adding the term

$$-m_0{}^2\ \text{tr}X_{n\mu}X_{n\mu}-\mu_0{}^2X_{n\mu}{}^0X_{n\mu}{}^0$$

to the Lagrangian reduces the gauge invariance to that of the "first kind," i.e., transformations with constant ϵ_n. For the spin-1 particles, there is both a dynamical and algebraic symmetry[12] of $[U(3)]^4$.

To reduce the symmetry further involves introducing *ad hoc* models for breaking the symmetry. For example, addition of a fermion mass term $m\bar{\psi}\psi$ reduces the symmetry under γ_5 transformations and leaves only $SU(3)$ symmetry. Alternatively, replacing the

[11] For example, the term arising from a second-class vector current in β decay contributes to the coefficient of k_μ.

[12] The concepts of dynamical and algebraic symmetry are discussed by S. Weinberg, *Proceedings of the Fourteenth International Conference on High-Energy Physics, Vienna, 1968* (CERN, Geneva, 1968), p. 253.

TABLE I. Possible nonets corresponding to the quantum numbers J^{PC}. The mixing angle is calculated using the mass-mixing (Ref. 2) model. The two missing mesons should be in the region ~1.2–1.6 GeV.

	J	P	C	$I=1$	$I=\frac{1}{2}$	$I=0$	$I=0$	θ
V	1	−	−	$\rho(765)$	$K^*(890)$	$\omega(783)$	$\phi(1020)$	40°
A	1	+	+	$A_1(1070)$	$K_A(1240)$	$D(1285)$	$E(1420)$	76°
$\tilde{V}$	1	−	+	$A_2{}^L(1270)$	$K^*(1400)$	$K_SK_S(1440)$	$\rho\rho(1410)$	90°
$\tilde{A}$	1	+	−	$B(1220)$	$K_A(1320)$			
					octet		singlet	

spin-one mass term above by a nondiagonal mass term $-\mathrm{tr}X_{m\mu}[m^2]_{mn}X_{n\mu}$ leads to terms like

$$-\mathrm{tr}\{m_1{}^2V_\mu{}^{(+)}V_\mu{}^{(+)}+m_2{}^2V_\mu{}^{(-)}V_\mu{}^{(-)}+m_3{}^2A_\mu{}^{(+)}A_\mu{}^{(+)}+m_4{}^2A_\mu{}^{(-)}A_\mu{}^{(-)}\}.$$

Similarly, the symmetry may be broken by off-diagonal terms in the nonlinear-interaction part of the Lagrangian.[13] The reduction of $SU(3)$ symmetry also involves extra assumptions beyond the gauge theory presented above, particularly since there will be mixing of the $\omega-\phi$ type for each set of nine fields.[14] In the following we shall assume that the Lagrangian has been rewritten completely in terms of the fields $V_\mu{}^{(\pm)i}$ and $A_\mu{}^{(\pm)i}$, and that the symmetry of the Lagrangian has been reduced to that of $SU(3)$.

If we assign a resonance to each of the vector fields, we see that there should be resonances with the quantum numbers $I^G=1^+$, $J^P=1^+$, and $I^G=1^-$, $J^P=1^-$, as well as the well-known quantum numbers of the ρ and A_1 mesons. These new quantum numbers would provide a natural explanation for the B meson at 1220 MeV, and for the $A_2{}^L(1270)$ if the spin of these mesons is confirmed to be one. In Table I, we show some tentative assignments of particles[15] to the 36 fields of $[U(3)]^4$. The mixing angle for the nine mesons $\tilde{V}$ is $\sim\frac{1}{2}\pi$, which means that there is essentially no mixing, or, that the Gell-Mann–Okubo mass formula holds *exactly* for the $A_2{}^L$, $K^*(1400)$, and $K_sK_s(1440)$. The two mesons needed to complete the identification of the $\tilde{A}$ nonet should be in the mass region around 1.35 GeV.

[13] Since local gauge invariance is now no longer applicable, the gauge-field part of the Lagrangian may be kept $SU(3)$-symmetric, and a symmetry-breaking term of the form $K^{ab}F_{\mu\nu}{}^aF_{\mu\nu}{}^b$ may be introduced into the matter part of the Lagrangian. See the first of Ref. 3 for details and a discussion of the general arbitrariness in the symmetry-violating term.

[14] There will also, in general, be mixing between the strange members of the $V^{(+)}$ $(A^{(+)})$ and $V^{(-)}$ $(A^{(-)})$ octets.

[15] Particle Data Group, Rev. Mod. Phys. **41**, 109 (1969); P. Antich *et al.*, Phys. Rev. Letters **21**, 1842 (1968).

5. CONCLUSIONS

We have examined the most general gauge theory of strong interactions, based on the existence of $SU(3)$ symmetry for the baryons, consistent with a four-component representation of the baryon fields, and have obtained a symmetry of $[U(3)]^4$. The currents corresponding to this symmetry have the form [when the Lagrangian is written in terms of the fields $V_\mu{}^{(\pm)}$ and $A_\mu{}^{(\pm)}$]

$$J_\mu{}^{V(-)i}=f^{ijk}(\bar{\psi}^j\gamma_\mu\psi^k+F_{\mu\nu}{}^{V(-)j}V_\nu{}^{(-)k}), \quad (5.1\text{a})$$

$$J_\mu{}^{V(+)i}=f^{ijk}F_{\mu\nu}{}^{V(+)j}V_\nu{}^{(+)k}, \quad (5.1\text{b})$$

$$J_\mu{}^{A(+)i}=f^{ijk}(\bar{\psi}^j\gamma_\mu\gamma_5\psi^k+F_{\mu\nu}{}^{A(+)j}A_\nu{}^{(+)k}), \quad (5.1\text{c})$$

$$J_\mu{}^{A(-)i}=f^{ijk}F_{\mu\nu}{}^{A(-)j}A_\nu{}^{(-)k}, \quad (5.1\text{d})$$

and the equations of motion are of the form

$$\partial_\mu F_{\mu\nu}{}^i-m^2\phi_\nu{}^i=gJ_\mu{}^i, \quad (5.2)$$

where ϕ refers to $V^{(\pm)}$, $A^{(\pm)}$. Using the techniques developed in Refs. 2 and 3, we can identify the spin mesons with the appropriate current operator and also arrive at an algebra similar to, but larger than, that obtained in Ref. 3.

In addition to providing a natural framework into which spin-1 resonances of abnormal C parity may be accommodated (such as the B meson and the $A_2{}^L$ if they are confirmed to be spin-1, we note that for the gauge fields, the dynamical and algebraic symmetries have the same group structure. The evidence against the existence of the new "second class" currents which enter here is not complete. Indeed there are indications that they may in fact exist.[16] In addition to the experimental processes listed[16] in which they may be observed, we note that, generally, we may expect to see effects due to their existence in comparing the scattering processes like $\nu+N\to N^*+\mu$ with $\bar{\nu}+N\to N^*+\mu$.[17,18]

ACKNOWLEDGMENTS

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[16] There are indications in μ capture in nuclei and nuclear β decay that second-class currents may be present with the same order of magnitude as the weak magnetism terms. See H. Ohtsubo [Phys. Letters **22**, 480 (1966)] for further references.

[17] M. Veltman, *Varenna School* (Academic Press Inc., New York, 1966), p. 187.

[18] N. Dombey [Phys. Rev. **174**, 2127 (1968)] has recently shown how second-class currents may affect weak η production.