Gauge schemes of weak and electromagnetic interactions

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The modern development of unified gauge theories of the weak and electromagnetic interactions is reviewed.

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1. GAUGE INVARIANCE AND UNIFIED RENORMALIZABLE THEORIES OF WEAK AND ELECTROMAGNETIC INTERACTIONS

Cabibbo's Universal Theory of Weak Interactions. Until the discovery at CERN[1,2] in neutrino interactions with matter of events due to the interaction of weak neutral currents, the overwhelming majority of the experimentally known facts of the physics of weak processes could be well described in the lowest order in the Fermi weak coupling constant with an effective Lagrangian having the form of a product of weak charged currents:

$$\mathscr{L}_{\text{eff}}=(G/\sqrt{2})\, J_{\lambda}^{+}J^{\lambda}+\text{h.c.}, \tag{1}$$

where $G\approx 10^{-5}/m_p^2$ (m_p is the proton mass), and the weak current J^λ is a sum of the lepton and the hadron terms:

$$J^\lambda = l^\lambda + h^\lambda. \tag{2}$$

The leptonic current l^λ has $V-A$ space-time structure and can be represented in the form

$$l^\lambda=\bar{\nu}_\mu\gamma^\lambda(1+\gamma^5)\,\mu+\bar{\nu}_e\gamma^\lambda(1+\gamma^5)\,e\equiv l_V^\lambda-l_A^\lambda. \tag{3}$$

The current l^λ contains the fields of the four known leptons: the electron and muon and the two neutrinos with muonic and electronic quantum numbers, respectively. The hadronic part of the total weak current (2) can be expressed in the form

$$h^\mu=\cos\theta_c\,\bar{n}\gamma^\mu(1+\gamma^5)\,p+\sin\theta_c\bar{\lambda}\gamma^\mu(1+\gamma^5)\,p\equiv h^\mu_{\text{vect}}-h^\mu_{\text{ax}}\equiv h_V^\mu-h_A^\mu. \tag{4}$$

Here, n, p, and λ are the fields of the neutron, proton, and strange quarks, respectively, and θ_c, the Cabibbo angle, is approximately equal to 15°. The current h^μ also has $V-A$ space-time structure.

One can rewrite the expression (4) in a more general form without recourse to a concrete realization of the current in a quark or some other constituent model of the hadrons:

$$h^\mu=(\mathfrak{F}^\mu_{1V}+i\mathfrak{F}^\mu_{2V}-\mathfrak{F}^\mu_{1A}-i\mathfrak{F}^\mu_{2A})\cos\theta_c+(\mathfrak{F}^\mu_{4V}+i\mathfrak{F}^\mu_{5V}-\mathfrak{F}^\mu_{4A}-i\mathfrak{F}^\mu_{5A})\sin\theta_c, \tag{5}$$

where the vector weak hadronic current $\mathscr{F}^\mu_{iV}$ and the axial-vector weak hadronic current $\mathscr{F}^\mu_{jA}$ $(i,j=1,\ldots,8)$ transform in accordance with the octet representation of $SU(3)$.

Substituting the expression (2) in (1), we conclude that the effective Lagrangian of the weak interaction (1) can be represented as a sum of three types of term, each of which corresponds to a definite type of interaction, namely,

$$\mathscr{L}_{\text{eff}}=\mathscr{L}^l+\mathscr{L}^{lh}+\mathscr{L}^h, \tag{6}$$

where

$$\mathscr{L}^l=(G/\sqrt{2})\,l_\mu l^{+\mu}; \tag{7}$$

$$\mathscr{L}^{lh}=(G/\sqrt{2})\,l_\mu h^{+\mu}+\text{h.c.}; \tag{8}$$

$$\mathscr{L}^h=(G/\sqrt{2})\,h_\mu h^{+\mu}. \tag{9}$$

The Lagrangian (7) describes weak interactions with the participation of only leptons, for example, the decays $\mu^\pm\to e^\pm+\nu_e(\bar{\nu}_e)+\nu_\mu(\bar{\nu}_\mu)$, and also scattering processes induced by weak interaction: $\mu^-+\nu_e\to e^-+\nu_\mu$, $e^-+\bar{\nu}_e\to e^-+\bar{\nu}_e$, etc. Interactions of the type (8) are semileptonic. In particular, they can be used to describe β decay of the neutron, $n\to p+e+\bar{\nu}_e$, the decays $\pi\to e(\mu)+\nu_e(\nu_\mu)$, $\pi^\pm\to\pi^0+e^\pm+\nu_e(\bar{\nu}_e)$, and also the process of quasielastic scattering with charge exchange: $\nu_e+n\to e^-+p$, $\nu_\mu+n\to\mu^-+p$, etc. The third type of interaction (9) is of a nonleptonic nature. Interactions of this type are responsible for nonleptonic hadron decays, for example, $K\to 2\pi$, 3π, $\Lambda\to p+\pi$, etc.

Note that the semileptonic part of the effective Langrangian $\mathscr{L}^{lh}$ ensures fulfillment of the known selection rules for hadronic β decay. The strangeness-conserving hadronic current $(\mathscr{F}^\lambda_{1V}+i\mathscr{F}^\lambda_{2V}-\mathscr{F}^\lambda_{1A}-i\mathscr{F}^\lambda_{2A})\cos\theta_c$ and the corresponding Hermitian conjugate current are the $(1+i2)$ and $(1-i2)$ components of an isotopic triplet. Therefore, the Lagrangian (8) ensures fulfillment of the selection rules $\Delta I=1$ (I is the hadron isospin) for the semileptonic strangeness-conserving hadron decays. The selection rule $\Delta I=1/2$ for semileptonic strangeness-changing hadron decays is also contained in the Lagrangian (8), since the strangeness-changing hadronic currents $(\mathscr{F}^\lambda_{4V}\pm i\mathscr{F}^\lambda_{5V}-\mathscr{F}^\lambda_{4A}\pm i\mathscr{F}^\lambda_{5A})\sin\theta_c$ are the components of (different) isospin doublets. The selection rule $\Delta S=\Delta Q$ is ensured by using the $(4\pm i5)$ components of the currents h^μ and $h^{+\mu}$. The rates of hadron β decays predicted by means of (8) are in good agreement with experiments.[3]

As regards the part of the Lagrangian (9) responsible for the nonleptonic hadron decays, its structure of the (current ⊗ current) type has been less well confirmed experimentally. The approximate selection rule $|\Delta I|=1/2$ established experimentally in the study of weak nonleptonic hadron decays is not a direct consequence of the Lagrangian (9), which contains terms corresponding to transitions with both $|\Delta I|=1/2$ and $|\Delta I|=3/2$. However, there exist theoretical arguments, not very strong, for dynamical enhancement[4] of transitions with $|\Delta I|=1/2$ compared with $|\Delta I|=3/2$ transitions.

Discussing the Cabibbo weak interaction scheme

TABLE I.

Quark	Q	Y	B	I
p	2/3	1/3	1/3	1/2
n	−1/3	1/3	1/3	1/2
λ	−1/3	−2/3	1/3	0

(7)–(9), we have above all in mind Gell-Mann's three-quark model,[5] which is one of the simplest field-theory models for studying the commutation relations between currents. Rather than introduce fields for the observed particles, we assume that they are all composite states of three hypothetical objects (quarks) with fractional charges and spin $\frac{1}{2}$. The assumed quantum numbers of these objects are given in Table I. Here, Q is the charge of the quarks, Y is the hypercharge, B is the baryon number, and I is the isospin. The quarks transform in accordance with the lowest representation of $SU(3)$.

The three-quark model describes numerous static properties of hadrons manifested in strong, electromagnetic, and weak interactions. Using the three quarks, one can construct all the known $SU(3)$ multiplets of mesons and baryons. The model gives a good description of electromagnetic properties of hadrons such as the magnetic moments, the electromagnetic mass differences, the amplitudes of electromagnetic meson decays, etc. It describes well a large group of phenomena in weak interactions. The weak charged hadronic current in this model is represented by Eq. (4). The achievements of the Cabibbo theory, and also some of its difficulties, have been excellently presented in numerous papers. Without pretending to a complete list, we refer the interested reader to, for example, the monographs of Refs. 6, 7, 8, and 9.

The Cabibbo theory of weak interactions has serious shortcomings. One of them is related to the fact that the interaction Lagrangian, which has a current–current structure, does not ensure renormalizability of the theory, and one therefore lacks a justification for an important computational formalism-perturbation theory. On the other hand, the four-fermion point interaction leads to a linear growth of the total cross sections of weak interaction processes with the energy. This circumstance leads to a contradiction with unitarity at a sufficiently high energy ($E^* \sim 300$ GeV in the center-of-mass system). Another shortcoming of the Cabibbo scheme (independent of the problem of renormalizability of the weak interactions[10]) relates to the circumstance that in the framework of this approach one cannot describe processes due to weak neutral currents. For example, comparing the $SU(3)$ form of the weak charged quark current (4) with its neutral component

$$h_\mu^3 = \bar{p}\gamma_\mu (1+\gamma_5)\, p - (\bar{n}\cos\theta_c + \bar{\lambda}\sin\theta_c)\,\gamma_\mu (1+\gamma_5)\,(n\cos\theta_c + \lambda\sin\theta_c), \tag{10}$$

we see that the strangeness-changing neutral current $\bar{\lambda}n + \bar{n}\lambda$ has order of magnitude $\sin\theta_c$ compared with the charged current, which does not correspond to the experimental fact of a much stronger suppression of strangeness-changing neutral currents. To overcome this difficulty, it is necessary to enlarge the symmetry of the strong interactions.[11] This corresponds to increasing the number of fundamental fields (quarks). Simultaneously, a number of other problems of the theory is solved.

Weinberg's Model. The Higgs Phenomenon. Renormalizations. One of the main methods of field theory is the Lagrangian formalism. Proceeding from the Lagrangian of the system and using the variational principle, one obtains the equations of motion, while the dynamical quantities such as the energy and momentum, charge, etc., that are conserved in time are determined by means of Noether's theorem,[12,13] which asserts that invariances of the Lagrangian under different transformations of the coordinate system and the field functions correspond to definite conservation laws.

One of the most important requirements imposed on the Lagrangian is that of relativistic invariance or the condition of invariance under the complete inhomogeneous Lorentz group. The Lorentz group does not exhaust the transformations that lead to important conservation laws. For example, to the conservation law of electric charge there corresponds a gauge transformation of the second kind, which affects only the field functions and not the coordinates. It is well known that the requirement that the Lagrangian be invariant under a gauge transformation of the first kind

$$\Psi \to \exp(iB\lambda)\,\Psi \tag{11}$$

(where Ψ is any field, and $B=0$ for a meson field and $B=1$ for a baryon field) in the Lagrangian formalism leads to a conservation law for the total baryon (lepton) charge (hypercharge). Invariance of the Lagrangian under the isospin transformations $\Psi \to \exp(iT\lambda)\Psi$, where λ is a constant real vector in the isospace, corresponds to the isospin conservation law. The above results are rigorous consequences of Noether's theorem, which states that to every finite-parameter (depending on N constant parameters) continuous transformation of the field functions and coordinates that has zero variation of the action there correspond N dynamical invariants, i.e., combinations of the fields and their derivatives that are conserved in time. Note that transformations of the type (11) differ from Lorentz transformations in that they affect only the field functions and not the coordinates.

At the present time, a tendency is developing in theoretical physics to describe a huge range of properties and species of elementary particles on the basis of a unified theoretical approach. It is quite possible that many principles hitherto regarded as fundamental will need to be reviewed. The gulf between the internal (hypercharge, isospin, and so forth) and external (space-time) symmetries in the world of elementary particles is felt especially acutely. A new approach augmenting the old principles is the requirement of invariance under gauge transformations that depend on the space-time coordinates with the related idea of universal interactions and gauge fields.

Let us consider transformations of field functions of the type (11), for example, isospin rotations. Conservation of isospin is equivalent to the requirement that the interaction Lagrangian be invariant under rotations in the isospin space. This means that in cases when the electromagnetic interactions can be ignored the orientation of the isospin in the isospace has no physical significance. Then the distinction between the proton and the neutron becomes purely formal. However, it is usually assumed that this arbitrariness is restricted by the following circumstance: Once one has chosen what to call a proton and what a neutron at one point of space-time, there is no longer any freedom of choice at other space-time points. Yang and Mills noted that this is inconsistent with the concept of a localized field, which provides the basis of ordinary physical theories, and they put forward the principle of invariance of interactions under independent rotations of the isospin at all space-time points.[14]

An analogous situation obtains with regard to the ordinary gauge invariance of a charged field, which is described by a complex wave function Ψ. A change of the gauge results in a change of the phase factor,

$$\Psi \to \Psi' = \exp(i\alpha)\,\Psi\,, \tag{*}$$

i.e., a change that does not lead to any physical consequences. Since Ψ can depend on $x,\ y,\ z,\ t$, the relative phase factor of the function Ψ considered at two different space-time points is completely arbitrary. In other words, the arbitrariness in the choice of the phase factor has a local nature.

With regard to a method for ensuring gauge invariance, it should be noted that in electrodynamics, to compensate the change of the phase α with a change in $x,\ y,\ z,\ t$, it is necessary to introduce an electromagnetic (and vector) field A_μ, which transforms under gauge transformations in accordance with the law (gauge transformation of the second kind)

$$A'_\mu = A_\mu + (1/e)\,\partial\alpha/\partial x_\mu. \tag{12}$$

Indeed, the Lagrangian

$$\mathscr{L} = \overline{\Psi}\gamma_\mu\partial_\mu\Psi - ie\overline{\Psi}\gamma_\mu\Psi A_\mu = \overline{\Psi}\gamma_\mu\nabla_\mu A_\mu, \tag{13}$$

in which the derivative ∂_μ with respect to the four coordinates is replaced by the "extended" derivative

$$\partial_\mu \to \nabla_\mu = \partial_\mu - ieA_\mu, \tag{14}$$

is invariant under the simultaneous gauge transformations (*) and (12). It is precisely this class of symmetries that currently engages the interest of physicists. The presence of an exact dynamical symmetry of this kind, such as gauge symmetry, requires the introduction of definite universal interactions, which are realized by massless vector particles. The universality of the interactions means that the coupling constant characterizing of the interaction of a particular gauge field (for example, photons and also vector and tensor mesons) with other particles is the same for all processes in which they participate.

The invariance of the theory under the groups of dynamical gauge symmetry is ensured by introducing an interaction with a particular field by the replacement of the ordinary derivatives by the "extended" derivatives. Namely, interaction with the gravitational field corresponds to the substitution $\partial_\mu \to \delta^\lambda_\nu\partial_\mu + \Gamma^\lambda_{\mu\nu}$, where $\Gamma^\lambda_{\mu\nu}$ is the Christoffel symbol; interaction with the electromagnetic field, to the substitution (14); interaction with an arbitrary gauge four-vector field B_μ, to the substitution $\partial_\mu \to \partial_\mu - i\varepsilon_a T_a B_\mu$, where T_a is the generator of the representations of the gauge group with respect to which the operators of the fields interacting with the gauge field B_μ transform.

The idea of gauge fields is particularly important in the search for new schemes of weak interactions. The problem of the search for new approaches to the description of weak interactions is related above all to the nonrenormalizability of the old (current ⊗ current) Fermi scheme. Ignoring the higher perturbation orders in the framework of the (current ⊗ current) scheme, we inescapably encounter a further difficulty, namely, linear growth of the total cross sections with increasing energy, which at high energies leads unavoidably to a violation of unitarity, as we have already mentioned.

On the other hand, having in view the analogy between the vector variant of the theory of weak interaction and quantum electrodynamics, we can expect to be able to prove renormalizability of gauge theories of the weak interaction that employ gauge vector fields to construct the interaction. The vector nature of the electromagnetic and weak interactions and their unification into a single gauge scheme with a universal coupling constant makes it possible to approach once more the problem of a unified theory of interactions of different type that hitherto were in no way related to one another.

The first step in the construction of such theories is the choice of the gauge symmetry group. As an example, let us consider the Lagrangian that describes the electromagnetic and weak interaction of leptons proposed by Salam and Weinberg.[15,16]

Suppose that we have at our disposal only the four known leptons ν_μ, μ, ν_e, e. From the fields of the leptons, we can form the four-vector of the density of the electromagnetic current:

$$J^{\mathrm{e.m.}}_\mu = -(\bar{e}\gamma_\mu e + \bar{\mu}\gamma_\mu\mu) \tag{15}$$

and the weak $V-A$ current

$$J^W_\mu = \bar{e}\,[(1+\gamma_5)/2]\,\gamma_\mu\nu_e + \bar{\mu}\,[(1+\gamma_5)/2]\,\gamma_\mu\nu_\mu. \tag{16}$$

We consider only currents containing the electron field (the muon term can be treated similarly). We find the minimal group that includes these currents. We introduce a left-handed isodoublet of leptons L and a right-handed isosinglet R:

$$L = \frac{1+\gamma_5}{2}\begin{pmatrix}\nu_e\\ e\end{pmatrix};\quad R = \frac{1-\gamma_5}{2}\,e; \tag{17}$$

$$\frac{1+\gamma_5}{2}e = e_L;\quad \frac{1-\gamma_5}{2}e = e_R;\quad \frac{1+\gamma_5}{2}\nu = \nu_L. \tag{18}$$

We introduce "charges" and hypercharges of the left- and right-handed fields:

$$\left.\begin{array}{l}\overline{Q}_L = -1/2;\quad \overline{Q}_R = -1;\\ Y = 2\overline{Q};\quad Q = I_3 - Y/2;\\ Y_L = -1;\quad Y_R = -2.\end{array}\right| \tag{19}$$

We require that I and Y (the leptonic isospin and the hypercharge or charge) be conserved in weak and electromagnetic interactions. These (I, Y) conservation laws correspond to the gauge group $SU(2)\times U(1)$. The weak currents can be represented in the form

$$\left.\begin{aligned} &J^-_{\mu W} = \bar{L}\gamma_\mu\tau_- L; \quad J^+_{\mu W} = \bar{L}\gamma_\mu\tau_+ L;\\ &[\tau_+, \tau_-] = \tau = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}, \end{aligned}\right\} \tag{20}$$

where $\tau_\pm$ and τ_3 are the generators of $SU(2)$. The currents $J^\pm_{\mu W}$ can be expressed in terms of the fields by

$$\left.\begin{aligned} J^-_{\mu W} &= \bar{e}\gamma_\mu[(1+\gamma_5)/2]\,\nu_e;\\ J^+_{\mu W} &= \bar{\nu}_e\gamma_\mu[(1-\gamma_5)/2]\,e. \end{aligned}\right\} \tag{21}$$

We write the electromagnetic current in the form

$$\bar{e}\gamma_\mu e = -(1/2)\,\bar{L}\gamma_\mu\tau_3 L + (1/2)\,\bar{L}\gamma_\mu L + \bar{R}\gamma_\mu R. \tag{22}$$

It follows from (22) that to obtain the electromagnetic current it is sufficient to add to the currents $\bar{L}\gamma_\mu\tau L$ of the group $SU(2)$ the isoscalar current (hypercharge current) $\bar{L}\gamma_\mu L + \bar{R}\gamma_\mu R$. Thus, the minimal group containing the weak and electromagnetic interactions is $SU(2)\times U(1)$. We require invariance of the weak and electromagnetic interactions under the gauge group $SU(2)\times U(1)$, this being characterized by the following transformations of the fields:

$$L \xrightarrow{SU(2)} L' = \exp[-i(g/2)\,\tau\varepsilon(x)]\,L;\quad R \xrightarrow{SU(2)} R;$$
$$L \xrightarrow{U(1)} L' = \exp[i(g'/2)\,\Lambda(x)]\,L;\quad R \xrightarrow{U(1)} R' = \exp[ig'\Lambda(x)]\,R.$$

The original Lagrangian of the free massless leptonic fields has the form

$$\mathcal{L}_{\text{free}} = \bar{L}\gamma_\mu\partial_\mu L + \bar{R}\gamma_\mu\partial_\mu R. \tag{23}$$

If it is to be $SU(2)\times U(1)$ symmetric, it is necessary to introduce compensating fields $\mathbf{A}_\mu$ and B_μ, which generate interaction of the leptonic currents with vector fields:

$$\mathcal{L}_{\text{int}} = (g/2)\,\bar{L}\gamma_\mu\tau L\mathbf{A}_\mu - g'\,(\bar{L}\gamma_\mu L/2 + \bar{R}\gamma_\mu R)\,B_\mu, \tag{24}$$

where $\mathbf{A}_\mu$ and B_μ transform as follows under $SU(2)$:

$$\left.\begin{aligned} &\mathbf{A}_\mu \to \mathbf{A}_\mu - g\mathbf{A}_\mu\times\boldsymbol{\varepsilon}(x) - \partial_\mu\boldsymbol{\varepsilon}(x);\\ &B_\mu \to B_\mu; \end{aligned}\right\} \tag{25}$$

under the group $U(1)$:

$$\mathbf{A}_\mu \to \mathbf{A}_\mu;\quad B_\mu \to B_\mu - \partial_\mu\Lambda(x). \tag{26}$$

Then the total Lagrangian, invariant under $SU(2)\times U(1)$, is

$$\left.\begin{aligned} &\mathcal{L} = -(1/4)\,\mathbf{A}^2_{\mu\nu} - (1/4)\,B^2_{\mu\nu} + \bar{L}\gamma^\mu D^L_\mu L + \bar{R}\gamma^\mu D^R_\mu R;\\ &\mathbf{A}_{\mu\nu} = \partial_\mu\mathbf{A}_\nu - \partial_\nu\mathbf{A}_\mu + g\mathbf{A}_\mu\times\mathbf{A}_\nu;\\ &B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu;\quad D^L_\mu = \partial_\mu - i(g/2)\,\tau\mathbf{A}_\mu + i(g'/2)\,B_\mu;\\ &D^R_\mu = \partial_\mu + ig'B_\mu \end{aligned}\right\} \tag{27}$$

or

$$\begin{aligned} \mathcal{L} = &-(1/4)\,(\partial_\mu\mathbf{A}_\nu - \partial_\nu\mathbf{A}_\mu + g\mathbf{A}_\mu\times\mathbf{A}_\nu)^2\\ &-(1/4)\,(\partial_\mu B_\nu - \partial_\nu B_\mu)^2 + \bar{L}\gamma_\mu[\partial_\mu\\ &- i(g/2)\,\tau\mathbf{A}_\mu + i(g'/2)\,B_\mu]\,L + \bar{R}\gamma_\mu(\partial_\mu + ig'B_\mu)\,R. \end{aligned} \tag{28}$$

A serious shortcoming of the theory with this Lagrangian is that all the gauge fields $\mathbf{A}_\mu$ and B_μ are massless, whereas this has been established only for the photon. Introduction into the Lagrangian of a mass term $(m^2/2)\mathbf{A}_\mu\mathbf{A}_\mu$ obviously destroys the gauge invariance of the Lagrangian and leads to nonrenormalizability of the theory.

The possibility of gauge-invariant introduction of massive vector fields into the theory was pointed out by Higgs.[17] In his model, the compensating fields become massive as a result of spontaneous breaking of the gauge symmetry.

The essence of the phenomenon of spontaneous symmetry breaking, which is important for the construction of renormalizable theories, can be conveniently illustrated by the example of Goldstone's model in the classical field theory with Lagrangian of scalar particles of the form

$$\begin{aligned} \mathcal{L} &= \partial_\mu\varphi^*\partial_\mu\varphi - \mu_0^2\varphi^*\varphi - h\,(\varphi^*\varphi)^2 = \partial_\mu\varphi^*\partial_\mu\varphi - V(\varphi^*\varphi)\\ &= \frac{1}{2}(\partial_\mu\varphi_1)^2 + \frac{1}{2}(\partial_\mu\varphi_2)^2 - V(\varphi_1, \varphi_2), \end{aligned} \tag{29}$$

where the parameter $\mu_0^2 < 0$, and $\varphi = (\varphi_1 + i\varphi_2)/\sqrt{2}$ is the scalar field of the charged mesons. The Lagrangian (29) is invariant under gauge transformations of the form $\varphi \to \exp(i\alpha)\varphi$, $\varphi^* \to \exp(-i\alpha)\varphi^*$. Consider the term

$$V(\varphi^*\varphi) = \mu_0^2\varphi^*\varphi + h\,(\varphi^*\varphi)^2. \tag{30}$$

This potential has a minimum at $|\varphi| = \lambda = \sqrt{-\mu_0^2/2h}$ (Fig. 1).

We respresent the field φ in the form $\varphi(x) = \rho(x)\exp[i\theta(x)]$, where ρ and θ are real functions of x. The fact that V has a minimum at $|\varphi| = \lambda$ in no way fixes the phase of the field function $\varphi(x)$ with given modulus λ. Thus, to the ground state there corresponds the value $\varphi_{\text{g.s.}} = \lambda\exp(i\alpha)$, where the parameter α is completely arbitrary. The vacuum is infinitely degenerate. In classical theory, the ground state is determined, as is well known, by the conditions

$$\partial V/\partial\varphi = 0;\quad \partial^2 V/\partial\varphi^2 \geqslant 0.$$

In the example considered above, the vacuum with $\varphi = 0$ is unstable, and with $\varphi_{\text{g.s.}} = \lambda\exp(i\alpha)$ it is stable but degenerate and does not have the symmetry inherent in the original Lagrangian. We make the canonical transformation

$$\varphi_1 = \varphi_1' + \chi;\quad \varphi_2 = \varphi_2' \tag{31}$$

and the substitution (31) in the Lagrangian (29):

$$\begin{aligned} \mathcal{L} = &(\partial_\mu\varphi_1')^2/2 + (\partial_\mu\varphi_2')^2/2 - (\mu_0^2/2)\,[(\varphi_1')^2 + 2\varphi_1'\chi + \chi^2 + (\varphi_2')^2]\\ &-(h/4)\,[((\varphi_1')^2 + (\varphi_2')^2)^2 + 4\varphi_1'((\varphi_1')^2 + (\varphi_2')^2)\,\chi + 4\,(\varphi_1')^2\chi^2\\ &+ 2\chi^2((\varphi_1')^2 + (\varphi_2')^2)] - h\chi^4/4 - h\varphi_1'\chi^3\\ = &(\partial_\mu\varphi_1')^2/2 + (\partial_\mu\varphi_2')^2/2 - (\varphi_1')^2(\mu_0^2/2 + 3h\chi^2/2)\\ &-(\varphi_2')^2(\mu_0^2/2 + h\chi^2/2) - \varphi_1'(\mu_0^2\chi + h\chi^3)\\ &-(h/4)\,[(\varphi_1')^2 + (\varphi_2')^2] - h\chi\varphi_1'[(\varphi_1')^2 + (\varphi_2')^2] - \mu_0^2\chi^2/2 - h\chi^4/4. \end{aligned} \tag{32}$$

The equations of motion that follow from the Lagrangian

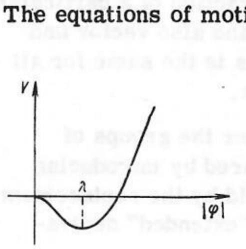

FIG. 1.

(32) redefined in this manner show that as a result of the spontaneous symmetry breaking there arise a field φ_1' with mass $\sqrt{2\mu_0}$ and a massless field φ_2' $(\mu_0^2+h\chi^2=0)$.

Goldstone's theorem. In a local, translationally invariant field theory with conserved current (local) and vacuum that is not invariant under the continuous symmetry group there must necessarily exist zero-mass particles.

Coleman's theorem. If the vacuum is not invariant under a particular transformation, then the Lagrangian is also not invariant under this transformation. The converse is not true.

We require invariance of the Goldstone Lagrangian under the group $U(1)$. Then the Lagrangian (32) takes the form

$$\mathcal{L}=-(1/4)F_{\mu\nu}^2+[(\partial_\mu-ieB_\mu)\varphi][(\partial_\mu+ieB_\mu)\varphi^*] -\mu_0^2\varphi^*\varphi-h(\varphi^*\varphi)^2;\ F_{\mu\nu}=\partial_\mu B_\nu-\partial_\nu B_\mu, \quad (33)$$

where $\mu_0^2<0$.

As was shown by Higgs[17] and Kibble,[18,19] one can in this case redefine the fields in such a way that the Goldstone field φ_2' disappears from the theory. To see this, we introduce the fields $\rho(x)$ and $\eta(x)$, make the gauge transformation

$$\left.\begin{aligned}&\varphi=\left(\frac{\chi+\rho}{\sqrt{2}}\right)\exp\left[i\,\frac{\eta(x)}{\chi}\right]\Rightarrow\left(\frac{\chi+\rho}{\sqrt{2}}\right)\exp\left[i\,\frac{\eta(x)}{\chi}+ie\Lambda(x)\right],\\&B_\mu\to B_\mu'=B_\mu-(1/e\chi)\,\partial_\mu\eta(x),\end{aligned}\right\} \quad (34)$$

and introduce $\Lambda(x)=-\eta(x)/e\chi$. Substituting (34) in the Lagrangian (33), we obtain the expression

$$\mathcal{L}=-(1/4)F_{\mu\nu}^2+[(\partial_\mu-ieB_\mu')(\chi+\rho)/\sqrt{2}][(\partial_\mu-ieB_\mu')(\chi+\rho)/\sqrt{2}]-V(\rho,\chi)=(\partial_\mu\rho)^2/2+e^2B_\mu B_\mu\rho^2/2 +e^2\chi^2B_\mu B_\mu/2-e^2B_\mu B_\mu\rho-(1/4)F_{\mu\nu}^2-V(\rho,\chi), \quad (35)$$

which does not contain the field $\eta(x)$.

From this result we obtain the following conclusions: Goldstone bosons can be eliminated by a special gauge transformation; at the same time, the gauge field B_μ acquires mass (the Euler–Lagrange equation for the vector field is $\Box B_\mu+m^2B_\mu=0$ with $m^2=e^2\chi^2$); the Higgs boson ρ also acquires mass.

We can now formulate a Higgs theorem in a generalized sense. Let G be a group of transformations that leaves the original Lagrangian invariant. The set of generators e_i of the group G that annihilates the vacuum state forms, by definition, the little subgroup of the vacuum. With each generator of G there is associated a gauge field. All fields corresponding to the generators of the little group of the vacuum remain massless. The remaining vector fields acquire mass.

It follows from this theorem that in gauge models with spontaneous breaking of local symmetry to the resultant symmetry $U(1)$ there remains one massive scalar Higgs boson and, accordingly, one massless vector field (the photon). As an example, we consider the Lagrangian

$$\mathcal{L}=-(1/4)F_{\mu\nu}^2+(1/2)(\partial_\mu\varphi-ig\mathbf{A}_\mu\tau\varphi)^*(\partial_\mu\varphi-ig\mathbf{A}_\mu\tau\varphi)+(m^2/2)\varphi\varphi^*-h(\varphi\varphi^*)^2,$$

where

$$\varphi=\begin{pmatrix}\varphi^+\\\varphi^0\end{pmatrix}$$

is an isodoublet of complex scalar fields (altogether four vector fields), and the vacuum expectation value is

$$\langle\varphi\rangle=\begin{pmatrix}0\\\lambda\end{pmatrix}.$$

In such a model, three of the four Goldstone bosons can be eliminated from the spectrum by an appropriate choice of the gauge. It is readily seen that the generators of $SU(2)$, applied to the vacuum state, do not form the little subgroup of the vacuum, as a result of which all three vector bosons in such a model acquire mass, and three new degrees of freedom—the longitudinal components of the fields of the vector bosons—arise as a result of the Goldstone bosons.

Let us consider briefly the renormalizability and unitarity of gauge theories. In the unitary gauge, the propagator of a vector field has the form $D_{\mu\nu}^c\sim(g_{\mu\nu}-k_\mu k_\nu/m_W^2)/(k^2-m_W^2)$ and the renormalizability of the theory is not obvious. In the gauge in which the propagator of the vector field is manifestly transverse in the four-dimensional sense (Landau gauge), $D_{\mu\nu}^c\sim(g_{\mu\nu}-k_\mu k_\nu/k^2)/(k^2-m_W^2)$, Goldstone bosons are absent.

The unitarity of the theory is not obvious. However, since the S matrix does not depend on the gauge, the total probability of transitions to unphysical states must vanish. In other words, by virtue of the gauge invariance the S matrix is simultaneously unitary and renormalizable.

Returning to the Weinberg–Salam model and using the Higgs mechanism, we introduce masses for the compensating fields, but in such a way that the photon field remains massless. For this, we consider the doublet

$$\phi=\begin{pmatrix}\phi^+\\\phi^0\end{pmatrix}. \quad (36)$$

For the field ϕ^0, we require a nonzero value of the vacuum expectation value as a consequence of the invariance of a self-interacting field of the type

$$-\mu^2(|\phi^+|^2+|\phi^0|^2)-h(|\phi^+|^2+|\phi^0|^2)^2.$$

As a result, the vector fields acquire mass through the appearance in the Lagrangian of the interaction

$$[\partial_\mu\phi^*+(ig/2)\phi^*\tau\vec{A}_\mu+(ig'/2)\phi^*B_\mu][\partial_\mu\phi-(ig/2)\tau\mathbf{A}_\mu-(ig'/2)B_\mu\phi];$$
$$\phi\xrightarrow{SU(2)}\phi'=\exp[-(ig/2)\tau\varepsilon(x)]\phi;\quad \phi\xrightarrow{U(1)}\phi' =\exp[-(ig'/2)\Lambda(x)]\phi.$$

Since the Lagrangian has the gauge invariance $SU(2)\times U(1)$, the Higgs transformations reduce in fact to the substitution

$$\phi=\begin{pmatrix}\phi^+\\\phi^0\end{pmatrix}\Rightarrow\begin{pmatrix}0\\\frac{\Lambda+\chi}{\sqrt{2}}\end{pmatrix}=\begin{pmatrix}0\\1\end{pmatrix}\frac{\Lambda+\chi}{\sqrt{2}}.$$

Rewriting the total Lagrangian

$$\mathcal{L}=-\mathbf{A}_{\mu\nu}^2/4-B_{\mu\nu}^2/4+\bar{L}\gamma_\mu[\partial_\mu-(ig/2)\tau\mathbf{A}_\mu-(ig'/2)B_\mu]L +\bar{R}\gamma_\mu(\partial_\mu+ig'B_\mu)R+[\partial_\mu\phi^*+(ig/2)\phi^*\tau\mathbf{A}_\mu+(ig'/2)\phi^*B_\mu] \times[\partial_\mu\phi-(ig/2)\tau\mathbf{A}_\mu\phi-(ig'/2)B_\mu\phi]-\sqrt{2}f(\bar{R}\phi^*L+\bar{L}\phi R),$$

we find that the last terms in the Lagrangian take the form

$$\partial_\mu\chi\partial_\mu\chi/4-\mu^2[(\Lambda+\chi)/\sqrt{2}]^2-h[(\Lambda-\chi)/\sqrt{2}]^4 +[(g^2+g'^2)/8](\Lambda^2+\chi^2)Z_\mu Z_\mu+(g^2/4)(\Lambda+\chi)^2 W^+_\mu W^-_\mu -(f/2)(\Lambda+\chi)\bar{e}e,$$

where

$$Z_\mu=(gA_{3\mu}-g'B_\mu)/\sqrt{g^2+g'^2} \qquad (37)$$

is a neutral vector field,

$$A_\mu=(g'A_{3\mu}+gB_\mu)/\sqrt{g^2+g'^2} \qquad (38)$$

is the electromagnetic field,

$$W^\pm_\mu=(A_{\mu1}\pm iA_{\mu2})/\sqrt{2} \qquad (39)$$

are charged vector fields.

We write out separately the Lagrangian responsible for the interaction of the weak leptonic currents and electromagnetic currents with the fields of the vector particles in the framework of the $SU(2)\times U(1)$ gauge-symmetric Weinberg–Salam model and discuss its consequences for some leptonic processes.

We consider the Lagrangian (24), rewriting it in the form

$$\mathscr{L}_{\text{int}}=(g/2)\bar{L}\gamma_\mu\tau_3LA_{\mu3}+(g/2)(\bar{L}\gamma_\mu\tau_1LA_{\mu1}+\bar{L}\gamma_\mu\tau_2LA_{\mu2}) -g'(\bar{L}\gamma_\mu L/2+\bar{R}\gamma_\mu R)B_\mu. \qquad (40)$$

By trivial calculations, we can reduce the second term of (40) to

$$(g\sqrt{2}/2)(J^+_{\mu W}W^-_\mu+J^-_{\mu W}W^+_\mu), \qquad (41)$$

where $J^\pm_W$ and W_μ are given by (20) and (39), respectively. We represent $A_{\mu3}$ and B_μ in the form

$$A_{\mu3}=A_\mu\sin\theta_W+Z_\mu\cos\theta_W;\quad B_\mu=A_\mu\cos\theta_W-Z_\mu\sin\theta_W. \qquad (42)$$

Then, taking

$$g\sin\theta_W=g'\cos\theta_W=e, \qquad (43)$$

we find

$$(g/2)\bar{L}\gamma_\mu\tau_3LA_{\mu3}-g'(\bar{L}\gamma_\mu L/2+\bar{R}\gamma_\mu R)B_\mu =(\sqrt{g^2+g'^2}/4)\{(4\sin^2\theta_W-1)\bar{e}\gamma_\mu e +\bar{\nu}_e\gamma_\mu(1+\gamma_5)\nu_e-\bar{e}\gamma_\mu\gamma_5 e\}Z_\mu-e\bar{e}\gamma_\mu eA_\mu.$$

In the final form, the Lagrangian of the interaction of the neutral and charged currents with the fields of the vector particles is written as

$$\mathscr{L}_{\text{int}}=-e\bar{e}\gamma_\mu eA_\mu+(\sqrt{g^2+g'^2}/4)\{(4\sin^2\theta_W-1)\bar{e}\gamma_\mu e +\bar{\nu}_e\gamma_\mu(1+\gamma_5)\nu_e-\bar{e}\gamma_\mu\gamma_5 e\}Z_\mu+(g/2\sqrt{2})\{\bar{\nu}_e\gamma_\mu(1+\gamma_5)eW^-_\mu +\bar{e}\gamma_\mu(1+\gamma_5)\nu_eW^+_\mu\}. \qquad (44)$$

In the Lagrangian (44), the first term describes the interaction of the electromagnetic current of the electron with the field of the massless photon, the second term the interaction of the weak neutral current with the field of the massive intermediate neutral boson Z, and the last term the interaction of the weak charged currents with the fields of the charged massive intermediate bosons $W^\pm_\mu$. To obtain the Lagrangian analogous to (44) for the interaction of the muonic currents with the fields of the vector particles it is sufficient in (44) to make the change of field variables $e\to\mu$, $\nu_e\to\nu_\mu$.

It follows from (43) that[1]

$$\text{tg}\,\theta_W=g'/g;\quad e=gg'/\sqrt{g^2+g'^2}. \qquad (45)$$

We can rewrite the second term of the expression (44) in the identical form

$$(\sqrt{g^2+g'^2}/2)(1/2)\{\nu_e\gamma_\mu(1+\gamma_5)\nu_e-\bar{e}\gamma_\mu(1+\gamma_5)e -4\sin^2\theta_W(-\bar{e}\gamma_\mu e)\}Z_\mu=(\sqrt{g^2+g'^2}/2)J^Z_\mu Z_\mu, \qquad (46)$$

where

$$J^Z_\mu=(1/2)[J^0_{\mu W}-4\sin^2\theta_W J^{\text{e.m.}}_\mu]. \qquad (47)$$

As a result, we arrive at the following "prescription" for constructing the interaction Lagrangian of the leptonic currents and the fields of the vector particles:

$$\mathscr{L}_{\text{int}}=eJ^{\text{e.m.}}_\mu A_\mu+(\sqrt{g^2+g'^2}/2)J^Z_\mu Z_\mu +(g/2\sqrt{2})(J^+_{\mu W}W^-_\mu+J^-_{\mu W}W^+_\mu). \qquad (48)$$

Comparing the theoretical and experimental values of the electromagnetic and weak coupling constants, we can readily estimate that

$$M_W=(g^2\sqrt{2}/2G)^{1/2}=(1/2\sqrt{2}\,|\sin\theta_W|)(e^2\sqrt{2}/G)^{1/2}\approx 60\text{ GeV}, \qquad (49)$$

$$M_Z=M_W/|\cos\theta_W|\approx 75\text{ GeV}. \qquad (50)$$

Here, we have used the experimental values of the Weinberg angle[2]:

$$\sin^2\theta_W\approx 0.25 \qquad (51)$$

and the Fermi coupling constant of the weak interaction $G=10^{-5}/m_p^2$, where m_p is the proton mass. As can be seen from the structure of the Lagrangian (48), it contains an interaction of weak neutral currents with the neutral intermediate boson. The ordinary Fermi theory does not contain an analog of such interactions directly in the first order in the Fermi coupling constant. The most direct verification of the theory can be made in experiments on the processes

$$\nu_e(\tilde{\nu}_e)+e\to\nu_e(\tilde{\nu}_e)+e; \qquad (52)$$

$$\nu_\mu(\tilde{\nu}_\mu)+e\to\nu_\mu(\tilde{\nu}_\mu)+e. \qquad (53)$$

The effective Lagrangian for these processes in the region of four-momentum transfers $|q^2|\lesssim m_Z^2$ and squares $s\leqslant m_Z^2$ of the energy in the center-of-mass system has the form

$$H_{\text{eff}}=(G/\sqrt{2})[\bar{\nu}_e\gamma_\mu(1+\gamma_5)\nu_e][\bar{e}\gamma_\mu(C_V+C_A\gamma_5)e], \qquad (54)$$

where

$$C_V=\pm 1/2\pm 2\sin^2\theta_W;\quad C_A=\pm 1/2. \qquad (55)$$

In the ordinary $V-A$ theory, $C_V=C_A=1$ for the first process and $C_V=C_A=0$ for the second. It follows from the available experimental data on the processes (52) and (53) that $\sin^2\theta_W\approx 0.2-0.35$. In the construction and use of the Weinberg–Salam $SU(2)\times U(1)$ gauge scheme, the hope was expressed that such a theory will be renormalizable. We emphasize, however, that some

[1] *Translator's note*: The Russian notation for the trigonometric, inverse trigonometric, hyperbolic trigonometric functions, etc., is retained here and throughout the article in the displayed equations.

[2] According to the latest data, $\sin^2\theta_W\approx 0.26\pm 0.03$ in $\nu_\mu N$ and $\tilde{\nu}_\mu N$ interactions.

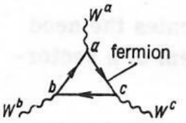

FIG. 2.

theories can contain so-called Adler anomalies.[20,21] The simplest of these anomalies are due to triangle diagrams of the form shown in Fig. 2. However, one can construct gauge theories in which the sum of the contributions of the anomalous diagrams vanishes. A necessary and sufficient condition for the existence of such theories is[22]

$$\mathrm{Sp}'\,\mathrm{Sp}\,[\gamma^5\{\Gamma^a,\ \Gamma^b\}\Gamma^c]=0 \tag{56}$$

for all a, b, and c, where the curly brackets denote the anticommutator. Note that the first trace is taken over the Dirac γ matrices, and the second is taken over the matrices L^a_{ij} and R^a_{ij} of the operator

$$\Gamma^a_{ij}=L^a_{ij}(1+\gamma_5)+R^a_{ij}(1-\gamma_5). \tag{57}$$

Thus, if we wish the theory to be renormalizable in all orders, the condition (56) must be satisfied. Note also that there can be cancellations between the leptonic and hadronic contributions to the anomalous Adler diagrams.

For models constructed on the basis of $SU(2)\times U(1)$ symmetry, the condition for the absence of anomalous contributions takes the very simple form

$$\sum_i (2T^i_3)\,Q_i=0, \tag{58}$$

where the summation is over all fermions that have charge Q_i and third component T^i_3 of the "weak isospin." Note that in the Weinberg–Salam model the sum (58) is equal to −2. Therefore, there must exist additional hadrons or (and) leptons, whose contribution to the sum (58) must be +2 if the sum is to vanish.

As for the purely leptonic case, one can also realize the Weinberg–Salam model in the hadronic sector. Below, we shall briefly consider some of the results of such a realization in the framework of a scheme with four quarks.

Weak Interactions of Hadrons. The GIM Quark Model. As was noted earlier, the weak strangeness-changing interactions of neutral and charged currents have the same order of magnitude in a model with three quarks (p, n, λ). This has the consequence that the decays $K_L\to\mu^+\mu^-$ and $K^+\to\mu^+\nu$ must have comparable probabilities,[23] which is in gross disagreement with experiment. This contradiction is an argument for the introduction of a fourth quark with unusual quantum number (charm or something else). The discovery of particles of the ψ family[24,25] was a serious argument in favor of schemes with four quarks. We shall discuss one of the most popular models, which includes a charmed quark alongside the three ordinary quarks.

In 1970, Glashow, Illiopoulos, and Maiani[11] proposed a model for the weak interactions (the GIM model), in which they introduced an additional "supercharged" quark, with a new quantum number "charm." Even earlier such a possibility had been considered in Ref. 26. The GIM model can be naturally generalized by the introduction of Yang–Mills fields; this removes the main difficulty of the theory of weak interactions—the problem of renormalizability.[3)] The Lagrangian of the weak interactions in the GIM model correctly reflects the empirical selection rules in weak interaction processes and the relations between the constants of the phenomenological theories of this interaction. The introduction of the fourth quark is also attractive on aesthetic grounds, for example, to achieve quark–lepton analogy. It is well known that in nature there exist four leptons (as yet we shall not mention the heavy leptons even if they do exist), namely, ν_e, ν_μ, e^-, μ^-, with charge differing from one another by at most unity. If we introduce a fourth quark with the same charge as the proton quark, it completes the symmetry between the quarks and the four leptons: Both quadruplets (leptonic and quark) have an as yet unexplained asymmetric mass spectrum and consist of two pairs of objects (leptons and quarks) whose charges differ from one another by unity.

The weak charged leptonic $V-A$ current can be represented in the form

$$J^l_\mu=\bar{l}C_l\gamma_\mu(1+\gamma_5)\,l, \tag{59}$$

where l is the column vector

$$\begin{pmatrix}\nu_e\\ \nu_\mu\\ e^-\\ \mu^-\end{pmatrix},$$

and the matrix C_l has the form

$$C_l=\begin{pmatrix}0&0&0&0\\0&0&0&1\\0&0&0&0\\0&0&0&0\end{pmatrix}. \tag{60}$$

By analogy with the expression (59), the weak charged hadronic $V-A$ current can be written as

$$J^H_\mu=\bar{q}C_H\gamma_\mu(1+\gamma_5)\,q, \tag{61}$$

where q is the column vector

$$\begin{pmatrix}c\\ p\\ n\\ \lambda\end{pmatrix}$$

of quarks. Here, p, n, and λ are the ordinary $SU(3)$ quarks, and the additional quark c has the same charge as p (i.e., +2/3), but has a new quantum number—charm $[C(c)=1]$. The matrix C_H has the form

$$C_H=\begin{pmatrix}0&0&-\sin\theta_c&\cos\theta_c\\0&0&\cos\theta_c&\sin\theta_c\\0&0&0&0\\0&0&0&0\end{pmatrix}. \tag{62}$$

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3) As we have already noted, in the framework of this model there exists a special type of divergent diagram associated with so-called Adler anomalies. Ways of eliminating such anomalies were discussed above.

Defining the matrices C_l^0 and C_H^0 as the commutators of the matrices C_l and C_l^+ and C_H and C_H^+, respectively,

$$C_l^0=[C_l,\ C_l^+]=\begin{pmatrix}1&0&0&0\\0&1&0&0\\0&0&-1&0\\0&0&0&-1\end{pmatrix}=[C_H,\ C_H^+]=C_H^0, \quad (63)$$

we obtain the following expression for the zeroth components with respect to the $SU(2)$ algebra of weak charges (both leptonic and hadronic):

$$J_{\mu l}^0=\bar{\nu}_\mu\gamma_\mu(1+\gamma_5)\nu_\mu+\bar{\nu}_e\gamma_\mu(1+\gamma_5)\nu_e - \bar{e}\gamma_\mu(1+\gamma_5)e-\bar{\mu}\gamma_\mu(1+\gamma_5)\mu; \quad (64)$$

$$J_{\mu H}^0=\bar{c}\gamma_\mu(1+\gamma_5)c+\bar{p}\gamma_\mu(1+\gamma_5)p-\bar{n}\gamma_\mu(1+\gamma_5)n-\bar{\lambda}\gamma_\mu(1+\gamma_5)\lambda. \quad (65)$$

To construct the neutral weak current interacting with the field of the neutral gauge vector boson Z, it is sufficient to use Eq. (47), which also holds in the hadron sector. Thus, in the GIM model the weak hadronic neutral current can be written as

$$J_\mu^{H(Z)}=\bar{c}\gamma_\mu\frac{1+\gamma_5}{2}c+\bar{p}\gamma_\mu\frac{1+\gamma_5}{2}p-\bar{n}\gamma_\mu\frac{1+\gamma_5}{2}n-\bar{\lambda}\gamma_\mu\frac{1+\gamma_5}{2}\lambda - 2\sin^2\theta_W\left(\frac{2}{3}\bar{c}\gamma_\mu c+\frac{2}{3}\bar{p}\gamma_\mu p-\frac{1}{3}\bar{n}\gamma_\mu n-\frac{1}{3}\bar{\lambda}\gamma_\mu\lambda\right), \quad (66)$$

where θ_W is the Weinberg angle and the fractional charges of the quarks have been taken into account in the term in the curly brackets. The interaction Lagrangian of the quark currents with the fields of the vector particles is given by Eq. (64). The charged quark currents interacting with the fields of the charged gauge bosons can be expressed in terms of (61), and the electromagnetic quark current has the form

$$J_\mu^{e.m.}=(2/3)(\bar{c}\gamma_\mu c+\bar{p}\gamma_\mu p)-(1/3)(\bar{n}\gamma_\mu n+\bar{\lambda}\gamma_\mu\lambda). \quad (67)$$

As can be seen from the expression (66), the neutral weak hadronic current does not contain nondiagonal transitions and the problem of the strangeness-changing neutral currents is automatically eliminated. The first experimental indication of a possible viability of the scheme of weak interactions with Lagrangian of the form (48) was the discovery with the Gargamelle chamber[1,2] of the neutral inclusive neutrino processes

$$\nu_\mu+\mathcal{N}\to\nu_\mu+\text{hadrons}, \quad (68)$$

$$\tilde{\nu}_\mu+\mathcal{N}\to\tilde{\nu}_\mu+\text{hadrons}. \quad (69)$$

The data on the reactions (68) and (69) (Ref. 27) agree with the Weinberg–Salam model of weak interactions and GIM if the Weinberg angle has the value (see also footnote[2)])

$$\sin^2\theta_W=0.33\pm0.05. \quad (70)$$

This same value is consistent with the data on neutral weak currents in the elastic scattering of neutrinos on protons, $\nu_\mu(\bar{\nu}_\mu)P\to\nu_\mu(\bar{\nu}_\mu)+P$, and in neutrino–nucleon interactions with production of a single secondary pion.[28] For a complete proof of the validity of the Weinberg–Salam and GIM model it will be necessary to continue the detailed experimental study of the spatial and isotopic structures of the neutral currents. At the least, the experimentally established effect of P parity violation in processes of interaction of weak neutral currents is an important criterion for selecting many-quark schemes ($n_q>4$) of weak interactions. In particular, this effect indicates the need to pay particular attention to the problem of a vector-like neutral current (see Sec. 3).

As can be seen from Eq. (61), the weak hadronic charged current is a linear combination of states that transform as a product of representations of $SU(4)$: $4^*\otimes4=15\oplus1$. The critical investigation of the GIM scheme is based above all on study of the semileptonic and nonleptonic hadron decays. When we speak of semileptonic processes, we have in mind the program of study of decays of charmed mesons and baryons of the form[118]

$$H_c\to H+(l\tilde{\nu}_l), \quad (71)$$

where H_c is the charmed hadron, H is an ordinary hadron, l is a lepton (μ or ν), and ν_l is the corresponding neutrino. Besides the processes (71), considerable interest attaches to the production of single charmed hadrons in interactions of neutrinos with matter.

2. FOUR-QUARK MODEL AND GENERAL PRINCIPLES FOR CONSTRUCTING GAUGE SCHEMES OF WEAK AND ELECTROMAGNETIC INTERACTIONS OF LEPTONS AND HADRONS

Gell-Mann's Hypothesis and the Universality of Weak Interactions. The SU(2) Algebra. Lepton–Hadron Analogy. The idea of lepton–hadron analogy was preceded by the hypothesis of universality of the weak interactions and Gell-Mann's related current algebra hypothesis.[29] This last hypothesis supposes that the time components of the vector and axial-vector currents satisfy the equal-time commutation relations

$$\left.\begin{aligned}&[\mathfrak{F}_k^0(x),\ \mathfrak{F}_l^0(y)]_{x^0=y^0}=i\delta(\mathbf{x}-\mathbf{y})f_{klm}\mathfrak{F}_m^0(x);\\&[\mathfrak{F}_k^0(x),\ \mathfrak{F}_l^{50}(y)]_{x^0=y^0}=i\delta(\mathbf{x}-\mathbf{y})f_{klm}\mathfrak{F}_m^{50}(x);\\&[\mathfrak{F}_k^{50}(x),\ \mathfrak{F}_l^{50}(y)]_{x^0=y^0}=i\delta(\mathbf{x}-\mathbf{y})f_{klm}\mathfrak{F}_m^{0}(x),\end{aligned}\right\} \quad (72)$$

where f_{klm} are the structure constants of the group $SU(3)$ [or $SU(4)$]. We define operators of "weak charges" corresponding to octets of weak vector and axial currents:

$$F_j(x^0)=\int d\mathbf{x}\mathfrak{F}_j^0(x) \quad (73)$$

and

$$F_j^5(x^0)=\int d\mathbf{x}\mathfrak{F}_j^{50}(x). \quad (74)$$

A more general formulation of Gell-Mann's hypothesis is the postulate that the following commutation relations hold [for the operators of the weak charges (73) and (74)]:

$$\left.\begin{aligned}&[F_k(x^0),\ F_l(x^0)]=if_{klm}F_m(x^0);\\&[F_k(x^0),\ F_l^5(x^0)]=if_{klm}F_m^5(x^0);\\&[F_k^5(x^0),\ F_l^5(x^0)]=if_{klm}F_m(x^0).\end{aligned}\right\} \quad (75)$$

These relations also have meaning when Schwinger terms are taken into account on the right-hand side of (72) (Ref. 30). The current algebra postulate (75) fixes the scale of the weak currents and permits a new formulation of the hypothesis of universality of the weak interactions. In its usual formulation, this asserts that the weak interactions responsible for β and μ decays have the same magnitude (i.e., they are characterized by the same coupling constant). This means that in the total weak current the leptonic and hadronic parts responsible for the strangeness-conserving transitions occur with the same coefficients, i.e., the currents $(\bar{p}n)$ and $(\bar{\nu}_\mu \mu + \bar{\nu}_e e)$ have the same coefficient. In general, allowance for strong-interaction renormalization effects must change the vector constant of the current $\bar{p}n$. However, as was shown in Refs. 31 and 32, it follows from the hypothesis of conservation of the vector current that this constant is not changed. Such a formulation of universality is somewhat inconvenient. It cannot be generalized to the case of decays of strange particles generated by the current $\bar{p}\lambda$, since one then obtains decay rates that are ten times greater than the experimentally observed values. In addition, one cannot explain the discrepancy of a few percent between the values of G_n and G_μ (G_n and G_μ are the Fermi constants of the weak interaction in the β decay of the neutron and the decay of the muon).

Following Gell-Mann, one can obtain a new formulation of the universality hypothesis that does not require one to express the weak current in terms of the fields of strongly interaction particles. We define the weak leptonic charges by

$$W_{l^+} = \frac{1}{2}\int d^3x J_l^0(x) = \frac{1}{2}\int d^3x\,[\nu_\mu^+(1+\gamma_5)\mu + \nu_e^+(1+\gamma_5)e]; \quad (76)$$

$$W_{l^-} = \frac{1}{2}\int d^3x\,[\mu^+(1+\gamma_5)\nu_\mu + e^+(1+\gamma_5)\nu_e] = W_{l^+}^+ \quad (77)$$

and introduce the equal-time commutators

$$[W_{l^+}, W_{l^-}] = 2W_{l_3}; \quad [W_{l_3}, W_{l\pm}] = \pm W_{l\pm}, \quad (78)$$

where

$$W_{l_3} = \frac{1}{4}\int d^3x\,[\nu_\mu^+(1+\gamma_5)\nu_\mu - \mu^+(1+\gamma_5)\mu + \nu_e^+(1+\gamma_5)\nu_e - e^+(1+\gamma_5)e]. \quad (79)$$

Thus, the charges $W_{l\pm}$ and W_{l_3} close the $SU(2)$ algebra. In the formulation of Gell-Mann universality, it is postulated that the hadronic weak charges

$$W_{h\pm} = \frac{1}{2}\int d^3x J^0_{h\pm}, \quad W_{h^-} = W_{h^+}^+ \quad (80)$$

satisfy the same commutation relations as the leptonic charges:

$$[W_{h^+}, W_{h^-}] = 2W_{h_3}, \quad [W_{h_3}, W_{h\pm}] = \pm W_{h\pm}. \quad (81)$$

If the hadronic current has the form [the case of $SU(3)$ symmetry]

$$J_h^\lambda = (\mathfrak{F}_1^\lambda + i\mathfrak{F}_2^\lambda - \mathfrak{F}_1^{5\lambda} - i\mathfrak{F}_2^{5\lambda})\cos\theta_c + (\mathfrak{F}_4^\lambda + i\mathfrak{F}_5^\lambda - \mathfrak{F}_4^{5\lambda} - i\mathfrak{F}_5^{5\lambda})\sin\theta_c, \quad (82)$$

then the relations (81) are satisfied if the hadronic weak charges have the form

$$W_{h\pm} = (1/2)(F_1 \pm iF_2 - F_1^5 \mp iF_2^5)\cos\theta_c + (1/2)(F_4 \pm iF_5 - F_4^5 \mp iF_5^5)\sin\theta_c; \quad (83)$$

$$W_{h_3} = (1/4)(1+\cos^2\theta_c)(F_3 - F_3^5) + (\sqrt{3}/4)\sin^2\theta_c(F_8 - F_8^5) - (1/2)\sin\theta_c\cos\theta_c(F_6 - F_6^5). \quad (84)$$

Thus, the definition of the weak current in the form (82) and the postulation of the commutation relations (75) are sufficient to ensure universality of the weak interactions in the sense of the relations (81). These relations coincide with the commutation relations for the generators of $SU(2)$. Note that whereas the existence of the analogous relations (76) can be directly verified for the leptonic weak charges, for the hadrons they are postulated. In other words, it is assumed that, despite the strong interaction effects, the weak hadronic charges form the same $SU(2)$ algebra as the leptonic charges. Frequently, the following generalization of the current algebra hypothesis is helpful:

$$\left.\begin{aligned}
&[F_k(x^0), \mathfrak{F}_l^\lambda(x)] = if_{klm}\mathfrak{F}_m^\lambda(x);\\
&[F_k(x^0), \mathfrak{F}_l^{5\lambda}(x)] = if_{klm}\mathfrak{F}_m^{5\lambda}(x);\\
&[F_k^5(x^0), \mathfrak{F}_l^\lambda(x)] = if_{klm}\mathfrak{F}_m^{5\lambda}(x);\\
&[F_k^5(x^0), \mathfrak{F}_l^{5\lambda}(x)] = if_{klm}\mathfrak{F}_m^\lambda(x).
\end{aligned}\right\} \quad (85)$$

These relations can be used to obtain the commutators of the charges F_k and F_k^5 with the effective Lagrangian of the nonleptonic weak decay. The Lagrangian of the nonleptonic weak decay constructed on the basis of the GIM scheme can be represented in the form

$$\mathscr{L}_{\rm eff}^{nl} = \mathscr{L}_{\rm eff}^{\rm pc} - \mathscr{L}_{\rm eff}^{\rm pv}, \quad (86)$$

where pc and pv stand for parity conserving and parity violating. Then, using (85), we obtain

$$\left.\begin{aligned}
&[F_k^5(x^0), \mathscr{L}_{\rm eff}^{\rm pc}(x)] = [F_k(x^0), \mathscr{L}_{\rm eff}^{\rm pv}(x)];\\
&[F_k^5(x^0), \mathscr{L}_{\rm eff}^{\rm pv}(x)] = [F_k(x^0), \mathscr{L}_{\rm eff}^{\rm pc}(x)];\\
&[F_k(x^0) + F_k^5(x^0), \mathscr{L}_{\rm eff}^{nl}(x)] = 0.
\end{aligned}\right\} \quad (87)$$

These commutation relations were used to study the nonleptonic decays of K mesons and hyperons. Note that the inclusion of $V+A$ currents in the weak interaction schemes somewhat modifies these relations and can violate some of the connections known in current algebra.

Consequences of the Commutation Relations for Weak Charges in a Scheme with Four Quarks. It is readily seen that the weak hadronic currents of the GIM model satisfy the charge algebra $SU(2)$. The next step is to find all possible schemes with quark currents in the four-quark model that realize the same algebra and in the three-quark limit go over into the ordinary Cabibbo scheme. The existence of schemes differing from the GIM scheme is perfectly possible, since the relations (81) are essentially nonlinear relations. The theoretical basis for the existence of different physically acceptable schemes of the weak quark currents is the very deep idea of universality of the weak interactions reformulated in the language of the commutation relations (81). The possible ambiguities that arise in virtually any theoretical model based on the least number of assumptions must be eliminated in what follows by comparing the predictions of the model with the experimental data.

Following Refs. 34–36, we formulate the main hypotheses used in the construction of the model:

1. The weak hadronic currents belong to the regular representation of $SU(4)$.

2. The $SU(2)$ algebra for the weak charges.

3. Universality of the weak interactions for leptons and hadrons.

More specifically, we shall assume that the most general form of the weak hadronic current for transitions with $\Delta Q=\pm 1$ has the form

$$\mathfrak{F}_{h\mu}=a(\mathfrak{F}_{\mu}^{1\pm i2}-\mathfrak{F}_{5\mu}^{1\pm i2})+b(\mathfrak{F}_{\mu}^{4\pm i5}-\mathfrak{F}_{5\mu}^{4\pm i5})$$
$$+c(\mathfrak{F}_{\mu}^{11\mp i12}-\mathfrak{F}_{5\mu}^{11\mp i12})+d(\mathfrak{F}_{\mu}^{13\mp i14}-\mathfrak{F}_{5\mu}^{13\mp i14}). \quad (88)$$

To determine the coefficients a, b, c, and d of the transition currents $p\leftrightarrow n$, $p\leftrightarrow\lambda$, $c\leftrightarrow n$, and $c\leftrightarrow\lambda$, respectively, we require that the weak hadronic charges (80) satisfy the commutation relations (81). The use of the relations (81) with allowance for the equal-time commutation relations (72) (the explicit form of the structure constants for the group $SU(4)$ is given in Ref. 120) leads to the following system of equations for the coefficients a, b, c, d:

$$\left.\begin{aligned}&a^3+ab^2+ac^2+bcd-a=0;\\&b^3+ba^2+bd^2+acd-b=0;\\&c^3+ca^2+cd^2+abd-c=0;\\&d^3+db^2+dc^2+abc-d=0.\end{aligned}\right\} \quad (89)$$

Solving this system, we find two classes of nontrivial solutions:

$$\text{I. } a^2+b^2=c^2+d^2=1;\quad ac=-bd; \quad (90)$$
$$\text{II. } a^2+b^2+c^2+d^2=1;\quad ad=bc. \quad (91)$$

Thus, the lepton–hadron symmetry (81) makes it possible to elucidate the structure of the hadronic current and, accordingly, fixes the scales of the leptonic and hadronic parts of the weak currents.

The solution (90) can be represented in the parametric form $a=d=\cos\theta_c$, $b=-c=\sin\theta_c$, where θ_c is the Cabibbo angle. Using the parametrization of the solution (90) in the expression for the neutral current, we can readily note that the strangeness-changing neutral currents $\mathscr{F}_{\mu}^{6}$ and $\mathscr{F}_{5\mu}^{6}$ and the charm-changing currents $\mathscr{F}_{\mu}^{9}$ and $\mathscr{F}_{5\mu}^{9}$ vanish. As a result, for the charged currents (88) we obtain an expression of the same form as for the corresponding quark currents in the GIM model.

The second solution (91) can be conveniently represented parametrically in the form

$$\left.\begin{aligned}&a=\cos\theta_c\cos\theta_r;\quad b=\sin\theta_c\cos\theta_r;\\&c=\cos\theta_c\sin\theta_r;\quad d=\sin\theta_c\sin\theta_r,\end{aligned}\right\} \quad (92)$$

where θ_r is a new parameter that can be conveniently called the charm angle.

Note that if this parameter vanishes, $\theta_r=0$, then, despite the existence of the quark c, charm-changing currents will be absent. It follows from data on the β decay of the neutron that the charm angle must be fairly small, at least $\theta_r\ll\theta_c$. The product $\cos\theta_c\cos\theta_r$ can be chosen in such a way that its value corresponds to the observed difference between the constants G_n and G_μ. As a result, the scheme with the solution (91) has a strong qualitative difference from the scheme with the solution (90), since the strangeness- and charm-changing currents $\mathscr{F}_{5\mu}^{13\mp i14}$ and $\mathscr{F}_{\mu}^{13\mp i14}$ are in this case additionally suppressed by the smallness of the charm angle θ_r. The production of charmed particles in weak interaction processes is then suppressed by a factor $\sin^2\theta_r$. The cross sections for the production of charmed strange particles is suppressed even more (by a factor $\sin^2\theta_r\sin^2\theta_c$).

Everywhere above we have assumed that the hadronic currents have a $V-A$ space-time structure. However, this restriction is not necessary and, as will be shown in what follows, is by no means dictated by the existing experimental data. We therefore go over to investigation of the more general formula

$$J_\mu=aV_\mu^{1\pm i2}+a'A_\mu^{1\pm i2}+bV_\mu^{4\pm i5}+b'A_\mu^{4\pm i5}$$
$$+cV_\mu^{11\mp i12}+c'A_\mu^{11\mp i12}+dV_\mu^{13\mp i14}+d'A_\mu^{13\mp i14} \quad (93)$$

for the current with as yet arbitrary relations between the vector and axial-vector parts of the currents. For convenience in the following calculations, we introduce half-sums and half-differences of the coefficients of the individual terms of the current (92):

$$A=(a+a')/2;\quad B=(b+b')/2;\quad C=(c+c')/2;\quad D=(d+d')/2;$$
$$A_1=(a-a')/2;\quad B_1=(b-b')/2;\quad C_1=(c-c')/2;\quad D_1=(d-d')/2. \quad (94)$$

The requirement of the lepton–hadron symmetry (81) leads to decoupled systems of equations for the sets of quantities $\{\mathscr{A}\}=\{A,B,C,D\}$ and $\{\mathscr{A}_1\}\equiv\{A_1,B_1,C_1,D_1\}$, which have exactly the same form as the system of equations (89). Choosing a solution of the type (90) for the coefficients $\{A_1,B_1,C_1,D_1\}$ and the trivial (zero) solutions for the set $\{A,B,C,D\}$, we obtain the GIM model, whereas a solution of the type (91) for the set $\{A_1,B_1,C_1,D_1\}$ and trivial solutions for the coefficients $\{A,B,C,D\}$ leads to the nontrivial result (92).

We now consider the following combination of solutions: of the type (91) for the set $\{\mathscr{A}\}$ and of the type (90) for the set $\{\mathscr{A}_1\}$. Then for the coefficients in the expression of the current (93) we obtain the parametric solution

$$\left.\begin{aligned}&a=\cos\alpha\cos\beta+\cos\theta_c;\quad c=\cos\alpha\sin\beta-\sin\theta_c;\\&a'=\cos\alpha\cos\beta-\cos\theta_c;\quad c'=\cos\alpha\sin\beta+\sin\theta_c;\\&b=\sin\alpha\cos\beta+\sin\theta_c;\quad d=\sin\alpha\sin\beta+\cos\theta_c;\\&b'=\sin\alpha\cos\beta-\sin\theta_c;\quad d'=\sin\alpha\sin\beta-\cos\theta_c,\end{aligned}\right\} \quad (95)$$

where α and β are free parameters.

One of the main requirements imposed on the structure of the current (93) is conservation of the Cabibbo form for the ordinary quark currents. Naturally, this requirement leads to definite restrictions on the currents of the charmed particles. The above requirement can be satisfied if, for example, $\beta\approx\pi/2$. The possibility that $\cos\beta\neq 0$ is an advantage of the above analysis. If $\cos\beta\neq 0$, it follows from the theory, for example, that one can reproduce the correct relation between the vector constants G_μ and G_n of μ and β decays.

The neutral component of the weak current [the third projection with respect to the $SU(2)$ algebra] in this scheme has the form

$$\begin{aligned}2J_\mu^3=&V_{3\mu}^3[2(A^2+A_1^2)+B^2+B_1^2+C^2+C_1^2]\\&+A_{3\mu}^3[2(A^2-A_1^2)+B^2-B_1^2+C^2-C_1^2]\\&-2V_{6\mu}^3[AB+A_1B_1+CD+C_1D_1]-2A_{6\mu}^3[AB-A_1B_1+CD-C_1D_1]\\&-\sqrt{3}V_{8\mu}^3[B^2+B_1^2+(2/3)(D^2+D_1^2)-(1/3)(C^2+C_1^2)]\\&+\sqrt{3}A_{8\mu}^3[B^2-B_1^2-(1/3)(C^2-C_1^2)+(2/3)(D^2-D_1^2)]\\&+2V_{9\mu}^3[AC+A_1C_1+BD+B_1D_1]+2A_{9\mu}^3[AC-A_1C_1+BD-B_1D_1]\\&-2\sqrt{2/3}V_{15\mu}^3[C^2+C_1^2+D^2+D_1^2]\\&-2\sqrt{2/3}A_{15\mu}^3[C^2-C_1^2+D^2-D_1^2].\end{aligned} \quad (96)$$

Experimentally, it is known that the processes due to the currents $J^3_{6\mu}(V^3_{6\mu}, A^3_{6\mu})$ (the transitions $n \leftrightarrow \lambda$) are strongly suppressed. In the model (96), the coefficient of the current $J^3_{6\mu}$ is $\sin 2\alpha$. It follows from the experiments that $\sin 2\alpha \ll \sin\theta_c$. The rigorous equality $\alpha = 0, \pi/2$ need not necessarily hold.

The coefficient of the current $J^3_{9\mu}$ is equal to $\sin 2\beta$. This circumstance automatically leads to a strong suppression of interactions involving neutral currents and a change of charm since $\beta \approx \pi/2$. As a result, in the current (93) the coefficients have the approximate values

$$\begin{aligned} &a \approx \cos\theta_c;\quad a' \approx -\cos\theta_c;\quad b \approx \sin\theta_c;\quad b' \approx -\sin\theta_c;\\ \text{a)}\ &c \approx 1 - \sin\theta_c;\quad c' \approx 1 + \sin\theta_c;\quad d \approx \cos\theta_c;\quad d' \approx -\cos\theta_c,\quad \alpha \approx 0;\\ \text{б)}\ &c \approx -\sin\theta_c;\quad c' \approx \sin\theta_c;\quad d \approx 1 + \cos\theta_c;\\ &d' \approx 1 - \cos\theta_c,\quad \alpha \approx \pi/2. \end{aligned} \tag{97}$$

The values of the parameters (97) do not contradict the existing experimental data and can be regarded as an alternative to the GIM model.

Right-Handed Currents[33-40] *$(\bar{c}n)$ and $(\bar{c}\lambda)$ and Some Problems in the Physics of Weak Decays of Hadrons.* In the physics of the weak nonleptonic decays of hyperons and K mesons there is a well-known empirical rule of enhancement of transition amplitudes with $\Delta I = 1/2$. Arguments based on current algebra or on asymptotic freedom in the framework of the GIM scheme are inadequate for the quantitative explanation of this enhancement, although an enhancement by a few times of the term in the Hamiltonian that transforms in accordance with the representation 8 of the group $SU(3)$ as compared with the term of the representation 27 can be justified. Experimentally, the amplitudes of transitions with $\Delta I = 1/2$ exceed the amplitudes of transitions with $\Delta I = 3/2$ by a factor of more than 20. This follows from comparison of the partial widths of the decays of K^0_S and K^+ mesons into two pions:

$$\Gamma(K^0_S \to \pi^+\pi^-)/\Gamma(K^+ \to \pi^+\pi^0) \sim 500 - 600. \tag{98}$$

Experimentally, it is also known that the amplitude of transitions with $\Delta I = 3/2$ in the decay $K^+ \to \pi^+\pi^0$ has the same order of magnitude as the amplitudes of the semileptonic channels of K^+-meson decay.

In the framework of the scheme with four quarks (p, n, λ, c), the inclusion of a right-handed charged weak quark current

$$J'_\mu = J_\mu + \bar{c}\gamma_\mu(1 - \gamma_5)\, n \tag{99}$$

makes it possible to explain[33] the enhancement of transitions with $\Delta I = 1/2$. In the present case, the effective Lagrangian for nonleptonic weak decays of quarks with $|\Delta S| = 1$ has this property:

$$\begin{aligned} \mathcal{L}_W = (G\cos\theta_c/\sqrt{2})\,\{&\bar{c}\gamma^\alpha(1+\gamma_5)\,\lambda\bar{n}\gamma_\alpha(1-\gamma_5)\,c\\ &+\sin\theta_c\,[\bar{p}\gamma^\alpha(1+\gamma_5)\,\lambda\bar{n}\gamma_\alpha(1+\gamma_5)\,p\\ &-\bar{c}\gamma^\alpha(1+\gamma_5)\,\lambda\bar{n}\gamma_\alpha(1+\gamma_5)\,c]\} + \text{h.c.} \end{aligned} \tag{100}$$

At the same time, there is no need to consider additionally dynamical mechanisms ensuring enhancement of transitions with $\Delta I = 1/2$.

The introduction into the theory of $V + A$ current of the type $(\bar{c}n)_R$ or $(\bar{c}\lambda)_R$ also makes it possible to eliminate the well-known difficulty associated with the existence of the decay $K^0_s \to 2\pi$. Let us recall the essence of the problem. In the naive quark model, the octet part of the effective nonleptonic Lagrangian responsible for transitions with $|\Delta S| = 1$ and $\Delta Q = 0$ transforms as the matrix λ_6 (Ref. 7). It is easy to show that in exact $SU(3)$ symmetry the decay $K^0_s \to 2\pi$ can be realized only through the 27-plet part of the Hamiltonian, and must therefore be strongly suppressed. But experimentally this decay is the main decay channel of the K^0_s meson. It seems remarkable that the main decay channel should be realized entirely through breaking of $SU(3)$ symmetry. Consideration of $V + A$ currents $(\bar{c}n)_R$ or $(\bar{c}\lambda)_R$ in the four-quark scheme removes this problem. Indeed, the P parity violating part of the interaction Lagrangian (100)

$$\bar{c}\gamma^\mu(1+\gamma_5)\,\lambda\bar{n}\gamma_\mu(1-\gamma_5)\,c + \text{h.c.} \tag{101}$$

transforms as the seventh component of the regular representation, as the matrix λ_7. Under these conditions, the decay $K^0_s \to 2\pi$ is allowed in exact $SU(3)$ symmetry.

On the other hand, the appearance in the Lagrangian of the weak nonleptonic interactions of a $V + A$ current of the type $(\bar{c}n)_R$ without suppression due to the smallness of the Cabibbo angle leads to some difficulties of the scheme. The most serious of them is the appreciable increase in the $K_L - K_S$ mass difference.[33] In what follows, we shall indicate a possible way of overcoming this difficulty based on the hypothesis that there exist virtual transitions that change the color quantum number (see also Refs. 41 and 81).

A further difficulty of the scheme relates to the restrictions that can be obtained from the chiral structure of the nonleptonic weak Hamiltonian responsible for $|\Delta S| = 1$ transitions in the study of the amplitudes of kaon and hyperon decays on the basis of current algebra and PCAC. We have already pointed out that the introduction of new currents with $V + A$ space-time structure can violate the commutation relations (87). Representing the nonleptonic Hamiltonian with $|\Delta S| = 1$ in the form

$$\mathcal{H}_W = \mathcal{H}^1_W + \mathcal{H}^3_W, \tag{102}$$

where the superscripts 1 and 3 refer to the components of the Hamiltonian responsible for the transitions with $\Delta I = 1/2$ and 3/2, respectively, we can, following Ref. 42, define four classes of models that differ by the form of the commutation relations between the Hamiltonian (102) and the operator $F^5_i, i = 1, 2, 3$:

$$[F^5_i, \mathcal{H}_W] = [F_i, \mathcal{H}^1_W] + [F_i, \mathcal{H}^3_W]; \tag{I}$$
$$[F^5_i, \mathcal{H}_W] = -[F_i, \mathcal{H}^1_W] + [F_i, \mathcal{H}^3_W]; \tag{II}$$
$$[F^5_i, \mathcal{H}_W] = [F_i, \mathcal{H}^1_W] - [F_i, \mathcal{H}^3_W]; \tag{III}$$
$$[F^5_i, \mathcal{H}_W] = -[F_i, \mathcal{H}^1_W] - [F_i, \mathcal{H}^3_W]. \tag{IV}$$

Note that the usual model with $V - A$ currents belongs to class I. The existence of the right-handed current $(\bar{c}n)_R$ leads to a commutator of type II. To understand the difference between these classes, we must return to the study of the nonleptonic decays of kaons and hyperons. Thus, a fairly successful prediction of the structure of the $K_S^0 \to 3\pi$ amplitude in terms of the experimental parameters of the $K_S^0 \to 2\pi$ decay using current algebra and PCAC requires that the Hamiltonian $H_W^{|\Delta S|=1}$ satisfy commutation relations of either class I or class II. At the same time, comparison of the s- and p-wave amplitudes of the hyperon decays leaves an arbitrariness in the choice between the commutation relations of the classes I and III. As a result, the nonleptonic Hamiltonian $H_W^{|\Delta S|=1}$ must belong to class I. Thus, the schemes of weak interactions with $SU(2)$ gauge symmetry in which one considers unsuppressed right-handed currents with the participation of the n quark encounter an additional difficulty. Introduction into the theory of a right-handed current of the type $(\bar{c}\lambda)_R$ keeps the theory in the class I, this being due to the fact that, except for the Cabibbo angle, the $(\bar{c}\lambda)$ current is a vector current. Such a structure of the current $(\bar{c}\lambda)$ does not violate the commutation relations (87).

It was shown in Ref. 81 that the introduction of an unsuppressed right-handed transition $(\bar{c}\lambda)_R$ in a theory in which there are color-changing transitions does not contradict the known value of the $K^0 \to \bar{K}^0$ transition amplitude. It is true that in this case the nonleptonic Hamiltonian responsible for transitions with $\Delta I = 1/2$ contains a small coefficient, namely, $\sin\theta_c$. As calculations in theories with asymptotic freedom show, introduction into the theory of right-handed currents makes it possible to obtain an appreciable dynamical enhancement of transition amplitudes with $\Delta I = 1/2$ by a factor of 10–20 compared with amplitudes of transitions with $\Delta I = 3/2$ (depending on the number of quarks in the considered model: $n_q = 4, 6, 8$).[43]

The GIM scheme, which fairly successfully describes a wide range of experimental data, may itself be part of a more complicated scheme containing, for example, six, eight, or more quarks. If such schemes do not at least contradict the experiments, it is worthwhile studying in their framework the mechanisms and values of the amplitudes of transitions of the ordinary quarks (p, n, λ) into the new quarks (c and others). An interesting way of verifying the existence of right-handed currents of the type $(\bar{c}n)_R$ or $(\bar{c}\lambda)_R$ could be by studying radiative decays of hyperons. It is known that the parity-violating parts of the amplitudes of the decays $\Sigma^+ \to P\gamma$ and $\Xi^- \to \Sigma^-\gamma$ vanish in the $SU(3)$-symmetric limit if the charged weak current has u-spin symmetry for strangeness-changing processes. This is true in both the Cabibbo and the GIM scheme. If the strong interactions are "asymptotically free," this result remains true for all weak interaction models containing only left-handed weak currents. If experimentally one were to observe significant parity violation in the decays $\Sigma^+ \to P\gamma$ and $\Xi^- \to \Sigma^-\gamma$, one could then conclude that there exist right-handed charged weak currents of the type $(\bar{c}n)_R$ or $(\bar{c}\lambda)_R$. Moreover, the available experimental data on the sign of the parameter of the angular asymmetry in these decays seems to favor the existence of the current $(\bar{c}\lambda)_R$; for more definite conclusions, it is necessary to raise considerably the accuracy in the measurement of the angular asymmetry in the radiative decays of Σ^+ and Ξ^-.

Summarizing our discussion of the possibility of introducing unsuppressed right-handed weak currents of the type $(\bar{c}n)_R$ or $(\bar{c}\lambda)_R$ into the theory, we can give preference to the second type of current. As an example of a scheme including right-handed currents we can mention the six-quark model (of the type proposed by Harari[44]) which will be discussed below.

On the Possibility of CP Parity Nonconservation in a Model with Four Quarks. We discuss briefly a possible generalization of the scheme. If it is assumed that the coefficients in the current (93) are complex, then the following relations for the coefficients that determine the weak currents can be obtained from the commutation relations:

$$|A|^2+|B|^2=|C|^2+|D|^2=1; \quad AC^*=-BD^*; \tag{103}$$
$$|A_1|^2+|B_1|^2=|C_1|^2+|D_1|^2=1; \quad A_1C_1^*=-B_1D_1^*. \tag{104}$$

The assumption that the coefficients are complex makes it possible to include in the treatment the neutral-current components J^7 and J^{10}, which have CP parity opposite to the sign of the CP parity of the currents J^6 and J^9. This example demonstrates the possibility of introducing CP parity violating interactions into the scheme. The possibility is realized by the existence of a solution of the commutation relations that is not identical with the widely adopted GIM model.[11]

Let us consider the example of the model of weak interactions with CP nonconservation[45-47] in the framework of the scheme with four quarks (p, n, λ, c). We assume that the gauge symmetry is $SU(2)_L \otimes SU(2)_R$. Using the complex solutions (103), we find the part of the Lagrangian of the weak interaction of the quarks with the W bosons that contains complex coefficients of the strangeness-changing transition currents:

$$\mathcal{L}_{\rm int} = g_L \{\bar{p}_L\gamma_\mu (n_L\cos\theta_c + \exp(i\delta_L)\lambda_L\sin\theta_c) W_{\mu L}^+ + \bar{c}_L\gamma_\mu [\exp(-i\delta_L) n_L \sin\theta_c + \lambda_L\cos\theta_c] W_{\mu L}^+\} + {\rm h.c.} + (L \to R). \tag{105}$$

The effective Hamiltonian of the nonleptonic weak interactions between the quarks with strangeness change $\Delta S = 1$ can be represented in the form[45,47]

$$\mathcal{H}_W^{|\Delta S|=1} = [f_L\exp(-i\delta_L) + f_R\exp(-i\delta_R)]\,\mathcal{L}^{\rm pc} + [f_L\exp(-i\delta_L) - f_R\exp(-i\delta_R)]\,\mathcal{L}^{\rm pv} + {\rm h.c.}, \tag{106}$$

where

$$\left.\begin{aligned} \mathcal{L}^{\rm pc} &= (V_{\Pi^-}V_{K^+} + A_{\Pi^-}A_{K^+})_{\rm sym} - (V'_{\Pi^-}V'_{K^+} + A'_{\Pi^-}A'_{K^+})_{\rm sym};\\ \mathcal{L}^{\rm pv} &= (V_{\Pi^-}A_{K^+} + A_{\Pi^-}V_{K^+})_{\rm sym} - (V'_{\Pi^-}A'_{K^+} + A'_{\Pi^-}V'_{K^+})_{\rm sym}. \end{aligned}\right\} \tag{107}$$

Here, we have introduced the following notation for the vector and axial quark currents: $V_{\Pi^-} = \bar{p}\gamma_\mu n$; $A_{\Pi^-} = \bar{p}\gamma_\mu\gamma_5 n$; $V'_{\Pi} = \bar{c}\gamma_\mu n$; $V_{K^+} = \bar{p}\gamma_\mu\lambda$; $V'_{K^+} = c\gamma_\mu\lambda$; $A_{K^+} = \bar{p}\gamma_\mu\gamma_5\lambda$; $A'_{K^+} = \bar{c}\gamma_\mu\gamma_5\lambda$; $(V_{\Pi^-}V_{K^+})_{\rm sym} = V_{\Pi^-}V_{K^+} + V_{K^+}V_{\Pi^-}$, etc. The effective coupling constants f_L and f_R are determined in the static case by the relation

$$f_{L,R} = g_{L,R}^2 \sin(2\theta_c)/16 m_{W_{L,R}}^2.$$

Redefining the phase of the field of the λ quarks, we can rewrite the Hamiltonian (106) as

$$\mathcal{H}_W^{|\Delta S|=1} = S^+ + P^+ + P^-, \tag{108}$$

where

$$\left.\begin{array}{l} S^{+}=G_S(\mathscr{L}^{\mathrm{pc}}+\mathscr{L}^{+\mathrm{pc}}); \\ P^{+}=G_P\cos\xi\,[\mathscr{L}^{\mathrm{pv}}+\mathscr{L}^{+\mathrm{pv}}];\; P^{-}=iG_P\sin\xi\,[\mathscr{L}^{\mathrm{pv}}-\mathscr{L}^{+\mathrm{pv}}]; \\ G_{\binom{S}{P}}=f_L\sqrt{1+\varepsilon^2\pm 2\varepsilon\cos\delta};\; \delta=\delta_R-\delta_L; \\ \mathrm{tg}\,\xi=2\varepsilon\sin\delta/(1-\varepsilon^2);\; \varepsilon=f_R/f_L. \end{array}\right\} \quad (109)$$

The Hamiltonian (108) contains explicitly CP parity violating terms. Note that CP parity will be conserved in nonleptonic hadron decays if one of the constants (f_L or f_R) is zero and (or) $\delta=\delta_L-\delta_R=0$. Note also that the CP-even and CP-odd operators P^+ and P^- in the Hamiltonian (108), which violate P parity, satisfy the relation

$$[T_3, P^-]=(i/2)(\mathrm{tg}\,\xi)\,P^+. \quad (110)$$

From this there follows automatically equality for the decay parameters of the K^0 mesons, $|\eta_{+-}|=|\eta_{00}|$, with relative phase Φ_{+-} approximately equal to $\Phi_{+-}=\tan^{-1}(2\Delta m/\Gamma_S)$, where $\Delta m=m_{K_L}-m_{K_S}$ is the mass difference of the K_L^0 and K_S^0 mesons and Γ_S is the total K_S^0 width.[47]

Thus, suppression of the CP violating effects can arise through a small value of the ratio f_R/f_L, while there remains the possibility of "maximal" CP violation $\delta=\pi/2$. It is readily seen that the CP parity violating part of the effective Hamiltonian (108) corresponds to transitions in which the strangeness changes by unity, $|\Delta S|=1$. Therefore, its contribution to the electric dipole moment d_n of the neutron is small. Under these conditions, the main contribution to the dipole moment of the neutron must be made by processes of interaction of quarks by the exchange of Higgs bosons (the terms corresponding to hypercharge-conserving transitions, $\Delta Y=0$, and terms that change sign under P and T transformations).

3. WEAK CURRENTS IN MANY-QUARK SCHEMES WITH COLOR DEGREE OF FREEDOM AND NEW QUANTUM NUMBERS

Color Degrees of Freedom in Gauge Theories. The investigation of weak and electromagnetic interactions in the framework of gauge theories appears incomplete if one does not use the available information about the internal structure of hadrons obtained in high-energy experiments. Study of the structure of hadrons revealed in lepton–hadron and hadron–hadron deep inelastic interactions, and also in processes of e^+e^- annihilation into hadrons at high energies, shows that the quarks behave as free point objects within the hadrons. At distances shorter than the radii of the hadrons, the effective coupling constant of the strong interaction between the quarks and the gluons is small ($g^2/4\pi \sim 1/5$). This circumstance can be used by assuming that the weak, electromagnetic, and strong interactions can be described in the framework of a unified theory[48-52] based on some gauge group characterized by a single universal coupling constant. The large differences between these forms of interaction established experimentally in the energy range of the contemporary accelerators must have a natural explanation in such theories. From this point of view, it is very important to study the properties of weak interactions due to neutral currents; indeed, it is important to study the conservation or nonconservation of P parity in neutrino–nucleon interactions, in weak muon–nucleon interactions at high energies, and in neutral transitions in atomic physics, and also to study neutral currents that change the strangeness and charm in lower and higher orders in the Fermi coupling constant G in perturbation theory.

In gauge theories there exists a further possibility for symmetry breaking; for example, P and CP nonconservation can be described by asymmetry of the vacuum.[53-55] Whatever the case, elucidation of the properties of neutral transitions could cast light on the possible interconnection between the strong, weak, and electromagnetic interactions.

The strong interaction can be included in a gauge-invariant manner in the considered schemes of weak and electromagnetic interactions by, for example, the introduction of a new degree of freedom of the quarks—the color quantum number.[56-60]

Note that the introduction into gauge theories of a fourth species of quark (with charm) made it possible to overcome serious difficulties of the old Cabibbo scheme, but a number of questions still remain open. Methods for overcoming them proposed in the literature differ; we have 1) the introduction into the theory of currents with new space-time structure, for example, $V+A$ currents; 2) an increase in the rank of the gauge group; 3) the introduction of new quantum numbers such as charm, beauty, etc.

Before we turn to the study of schemes with quarks that have new flavors [the groups $SU(6)$ and $SU(8)$], we must consider the possibility of eliminating these difficulties in simpler schemes with three or four quarks but taking into account the color degree of freedom. Quark color was introduced in order to reconcile the spin and statistics of quarks[56-59] when certain hadrons, for example, Δ^{++} and Ω^-, are constructed from quarks. Quark color is also attractive to the theoreticians to enable them to reconcile the theoretical predictions of the $\pi^0 \to 2\gamma$ decay width in the framework of the quark model and the experimental data. The color degree of freedom significantly improves the theoretical estimate for the ratio $R(s)$ of the total cross section of e^+e^- annihilation into hadrons to the cross section of the process $e^+e^- \to \mu^+\mu^-$ at high energies ($\sqrt{s}>5$ GeV), making it approach the experimentally observed value.

One of the most popular versions of colored quarks is the model with fractional quark charges. In this approach, the hadron symmetry is characterized by the group $SU(n)\times SU(3)^c$, which is the direct product of the ordinary $SU(n)$ group acting on the unitary space of flavors $(p, n, \lambda, c, \ldots)$ and the color group $SU(3)^c$. It is assumed that the color symmetry $SU(3)^c$ is exact. In this model, it is also assumed that all ordinary hadrons are singlets of the color group. At the same time, the quarks are distributed among color triplets and are therefore assumed to be unobservable in the free state. The quantum numbers of quarks that differ only by their color are assumed to be the same; the quark charges are fractional, as in the Gell-Mann–Zweig model:

$$\left.\begin{array}{l} Q_{p_i}=2/3;\ Q_{n_i}=-1/3; \\ Q_{\lambda_i}=-1/3;\ Q_{c_i}=2/3;\quad i=1,\ 2,\ 3. \end{array}\right\} \tag{111}$$

The physical currents (electromagnetic and weak) are regarded in the majority of investigations as singlets of the color group. For example, the electromagnetic current has the form

$$J_\mu^{\text{e.m.}}=\sum_{i=1,\,2,\,3}\left[\tfrac{2}{3}\,\bar{p}_i\gamma_\mu p_i-\tfrac{1}{3}\,\bar{n}_i\gamma_\mu n_i-\tfrac{1}{3}\,\bar{\lambda}_i\gamma_\mu\lambda_i+\tfrac{2}{3}\,\bar{c}_i\gamma_\mu c_i\right]. \tag{112}$$

A gauge model for the strong interactions is constructed in the framework of the gauge group $SU(3)^c$. As a result, the quanta of the strong interaction are eight colored massless Yang–Mills vector mesons (called gluons). In non-Abelian gauge theories of this kind, one has what is known as asymptotic freedom[61,13]—logarithmic tending to zero of the renormalized coupling constant α_{St} of the quark–gluon interaction at short distances ($r \lesssim 1$ F). This last circumstance makes it possible to use perturbation theory in the coupling constant α_{St} in this region of space (or at the corresponding large momentum transfers). This theory is constructed by analogy with quantum electrodynamics and is called quantum chromodynamics. The theory has infrared divergences in the region of small momentum transfers due to the fact that the gluons have no mass. Thus, quantum chromodynamics become unstable at large distances ($r \lesssim 1$ F), and perturbation theory breaks down in this region. It is assumed that the infrared instability of the theory ensures screening of the colored states (quarks, gluons, etc.) from the real world (quark confinement).

In the literature, a model with colored quarks that have integral charges has also been discussed. In the most popular version of such models—the Han–Nambu model[60]—the group of hadron symmetry is assumed to be a diagonal subgroup of the group $SU(n)'\times SU(3)''$, with $n=3, 4$. It is assumed that the quarks belong to the representation $(n, 3^*)$ of the group $SU(n)'\times SU(3)''$. The possibilities of constructing unified theories of the weak, electromagnetic, and strong interactions in the framework of such an approach were considered in Refs. 48–52. In this approach,[48] the generators of the color gauge group no longer commute with the generators of the gauge group $SU(2)\times U(1)$ of the weak and electromagnetic interactions. Therefore, the color gauge invariance of the theory is spontaneously broken and, as a consequence, the gluons not only acquire mass but are also mixed with the weak and electromagnetic gauge bosons. Despite this mixing, in such schemes one has Bjorken scaling in deep inelastic lepton–hadron interactions. It is also necessary to emphasize that in this class of unified gauge theories the property of asymptotic freedom is not satisfied in the rigorous sense.

We mention an interesting property of models with integrally charged quarks.[62] At sufficiently high energies, above the threshold for the production of colored states ($s \gg m_g^2$), quarks with integral charges can behave like the fractionally charged Gell-Mann–Zweig quarks. Indeed, let us consider the electromagnetic current[62,63] (m_g is the gluon mass, $s > m_g^2$):

$$J_{\mu\text{eff}}^{\text{e.m.}}=J_{\mu 0}+[m_g^2/(m_g^2+|q^2|)]\,J_{\mu 8}, \tag{113}$$

where $J_{\mu 0}$ is determined by (112), and we have the current

$$J_{\mu 8}=-\frac{2}{3}\sum_{\alpha=p,\,n,\,\lambda,\,c}\bar{q}_1^\alpha\gamma_\mu q_1^\alpha+\frac{1}{3}\sum_{\alpha=p,\,n,\,\lambda,\,c}\bar{q}_2^\alpha\gamma_\mu q_2^\alpha+\frac{1}{3}\sum_{\alpha=p,\,n,\,\lambda,\,c}\bar{q}_3^\alpha\gamma_\mu q_3^\alpha. \tag{114}$$

Below the threshold for the production of colored states ($s=|q^2|\ll m_g^2$),

$$J_{\mu\text{eff}}=J_{\mu 0}. \tag{115}$$

Above the threshold ($s \gg m_g^2$), we obtain the same result:

$$J_{\mu\text{eff}}\approx J_{\mu 0}. \tag{116}$$

The main objection to theories with integrally charged quarks arises from the fact that hitherto quarks and gluons have not been observed experimentally. The reader can become more acquainted with the physics of hadrons with color degrees of freedom in the reviews of Refs. 64–66.

We mention only that in models with color symmetry one can study some difficult problems of the weak interactions such as the problem of octet dominance in the nonleptonic decays of hadrons.[67] In such models, one constructs an effective Hamiltonian for nonleptonic strangeness-changing hadron decays ($|\Delta S|=1$), in which a term that transforms in accordance with the octet representation of $SU(3)$ is dominant. Therefore, the selection rule $\Delta I=1/2$ is satisfied. In the framework of such an approach, the problem of theoretical calculation of the rate of nonleptonic hadron decays remains unsolved.

Physical Restrictions on the Weak Currents in Many-Quark Schemes. In schemes with a large number of quarks (quarks with new quantum numbers: color, heaviness, etc.) it is necessary, in order to restrict the arbitrariness, to formulate basic requirements that the theory must successfully satisfy, namely:

1) the charged components of the currents $J_\mu^+, J_\mu^-, J_\mu^3$ (both left- and right-handed) must form the algebra $SU(2)_L$ or $SU(2)_L\times SU(2)_R$, and the truly neutral current may be a mixture of the components J_μ^3 and $J_\mu^{\text{e.m.}}$;

2) the charged currents containing the quarks p, n, λ must have the Cabibbo form;

3) the neutral current must be vectorlike;

4) there must be satisfied (completely or partly) the hypothesis of diagonality of the neutral transitions in both the lower and higher perturbation orders in the coupling constant of the weak interaction. It should be noted that among the neutral transitions it is helpful to distinguish genuinely neutral transitions, which are diagonal and take place without change of mass or other quantum numbers of the quarks (leptons), and nondiagonal neutral transitions, which are transitions

between quarks (leptons) that involve a change in the mass and some quantum number of the quark (lepton) besides the electric charge.

We have already mentioned above that the idea of a vectorlike neutral current is associated with attempts at unification of the known types of interaction of the elementary particles. In particular, it is to be expected that at high energies the differences between the strong and weak interactions disappear. For example, spatial parity in processes due to weak interactions will be conserved at asymptotically high energies (possibly in the region of the unitarity limit). This property is automatically satisfied in gauge schemes of the weak interactions based on left–right symmetry, for example, the group $SU(2)_L \times SU(2)_R \times U(1)$.[68-74] It is therefore of particular interest to study neutral weak interactions already at the energies of the contemporary accelerators with a view to elucidating the gauge symmetry corresponding to the weak interactions. Below, whenever we speak of the principle of a vectorlike neutral current, we have in mind the possibility of extending the gauge group to a left–right symmetric group, for example, $SU(2)_L \to SU(2)_L \times SU(2)_R$, $SU(3)_L \to SU(3)_L \times SU(3)_R$, etc. Therefore, one of the main requirements on a vectorlike theory is that all the quarks participating in the construction of the $V-A$ currents should also participate in the construction of the $V+A$ currents. Such vectorlike theories are free of "triangle" divergences[22] (Adler anomalies). In them, it is assumed that in the absence of masses of the elementary fermions (both the quarks and the leptons) the neutral gauge currents are vector currents[75] and that parity must be conserved in neutral transitions. When the interactions are switched on (fermion mass matrix) and the fermions acquire mass, scalar and pseudoscalar parts appear in the mass matrix. The unitary transformations that diagonalize the fermion mass matrix and eliminate its pseudoscalar part lead to the appearance in the Lagrangian of terms with an axial-vector current. This, in its turn, leads to P-odd effects in processes due to the interaction of neutral currents. At the present time, the situation in experimental physics[76] with regard to the study of neutral currents is not clear. For example, in neutrino interactions through the neutral current the P parity is violated in accordance with the Weinberg–Salam–GIM model. Study of neutral currents in atomic physics has not yet given a clear answer, since the experiments in this direction are contradictory and give large errors.[119]

Unification of the theories of the weak, electromagnetic, and strong interactions can be concretely realized in schemes in which the leptons are regarded as quarks with a new quantum number (fourth color)[48].

$$\begin{pmatrix} p_1, & p_2, & p_3; & \nu_e \\ n_1, & n_2, & n_3; & e^- \\ \lambda_1, & \lambda_2, & \lambda_3; & \mu^- \\ c_1, & c_2, & c_3; & \nu_\mu \end{pmatrix}. \tag{117}$$

In such an approach, it is necessary to distinguish two possibilities.

In Refs. 62 and 63 it was assumed that the quarks and leptons belong to the $(4.\overline{4})$ representation of $SU(4)' \times SU(4)''$. A unified theory of the interactions of elementary particles was then formulated in the language of the gauge symmetry $SU(2)_L \times SU(2)_R \times SU(4)$. With successive spontaneous breaking of the original symmetry the final invariant group is $U(1)$, corresponding to conservation of the electromagnetic current.

In a different approach,[70] based on Gell-Mann–Zweig quarks, the fermion fields transform in accordance with the following representations of the group $SU(2)_L \times SU(2)_R \times SU(4^c)$:

$$f = 2\{(2.\ 1.\ 4) + (1.\ 2.\ 4)\} = \begin{pmatrix} p_1, & p_2, & p_3; & \nu_e \\ n_1, & n_2, & n_3; & e^- \end{pmatrix} + \begin{pmatrix} c_1, & c_2, & c_3; & \nu_\mu \\ \lambda_1, & \lambda_2, & \lambda_3; & \mu^- \end{pmatrix}. \tag{118}$$

The group $SU(4)$ contains the subgroup $SU(3)^c \times U_1(Y)$, where Y is the $SU(4)$ hypercharge,

$$Y = \frac{1}{3}\begin{pmatrix} 1 & & & 0 \\ & 1 & & \\ & & 1 & \\ 0 & & & -3 \end{pmatrix}. \tag{119}$$

If we denote by $\mathbf{T}_L$ and $\mathbf{T}_R$ the generators of the groups $SU(2)_L$ and $SU(2)_R$, then

$$Q_{\text{e.m.}} = T_3 + \frac{1}{2}Y = \frac{1}{2}\begin{pmatrix} 1 & 1 & 1 & 1 \\ -1 & -1 & -1 & -1 \end{pmatrix} + \frac{1}{6}\begin{pmatrix} 1 & 1 & 1 & -3 \\ 1 & 1 & 1 & -3 \end{pmatrix}, \tag{120}$$

where $T_3 = T_3^L + T_3^R$.

As a result of successive spontaneous breaking of the symmetry $SU(2)_L \times SU(2)_R \times SU(4^c)$, the symmetry $SU(3^c) \times U(1)$ arises. Note that in the schemes considered above there are quark–lepton transitions that do not conserve the baryon and lepton numbers. The fermion quantum number is conserved. If the mass of the gauge vector meson which interacts with the lepton and quark $[\mathcal{L}_{\text{int}} \sim V_{\text{eq}}(\bar{q}l)]$ is $m_V \gtrsim 10^4$ to 10^5 GeV, then we have the proton lifetime $\tau_P \gtrsim 10^{30}$ years, which does not contradict the experimental data.[62] In the theory with Gell-Mann–Zweig quarks, the leptoquark bosons V_{eq} carry fractional charges. In contrast to quarks and gluons, such mesons can exist in the free state. Note that at the distances characterized by the masses of the leptoquark V_{eq} bosons, i.e., at $r \lesssim 10^{-18}$ cm $[m(V_{\text{eq}}) \gtrsim 10^5$ GeV$]$, all types of interaction are of the same order of magnitude.

As a result of the successive spontaneous symmetry breakings, the effective current–current Hamiltonian of the weak interaction takes the form

$$\begin{aligned} \mathcal{H}_W &= \frac{g^2}{4M_{W_L}^2} j_{\mu L} j_L^{\mu +} + \frac{g^2}{4M_{W_R}^2} j_{\mu R} j_R^{\mu +} \\ &+ \frac{g^2}{4M_{Z_A}^2} j_{\mu A} j_A^{\mu +} + \frac{g^2 + g^2}{4M_{Z_V}^2} j_{\mu V} j_V^{\mu +} + \text{h.c.} \\ &= \frac{G_F}{\sqrt{2}}\left[j_{\mu L} j_L^{\mu +} + \frac{M_{W_L}^2}{M_{W_R}^2} j_{\mu R} j_R^{\mu +} + j_{\mu A} j_A^{\mu +} + j_{\mu V} j_V^{\mu +} \right]. \end{aligned} \tag{121}$$

The charged weak currents in Eq. (121) are the same as in the Weinberg–Salam–GIM model; the right-handed currents $j_{\mu R}$ differ from the left-handed currents $j_{\mu L}$ only by their space-time structure, which has the form $V+A$. For significant suppression of the currents (in agreement with the experimental data) the spontaneous symmetry breaking must be realized in such a way that the mass of the W_R boson appreciably exceeds the mass of the W_L boson: $m(W_R^\pm) \gg m(W_L^\pm)$. In this scheme, the neutral axial-vector and vector currents are defined in accordance with

$$j_{\mu A} = \bar{p}\gamma_\mu\gamma_5 p - \bar{n}\gamma_\mu\gamma_5 n + \bar{\nu}_e\gamma_\mu\gamma_5\nu_e - \bar{e}\gamma_\mu\gamma_5 e + \ldots; \quad (122)$$

$$j_{\mu V} = (\bar{p}\gamma_\mu p - \bar{n}\gamma_\mu n + \bar{\nu}_e\gamma_\mu\nu_e - \bar{e}\gamma_\mu e) - 2\sin^2\theta_V J_\mu^{\text{e.m.}} + \ldots \quad (123)$$

To the currents (122)–(123), one must add terms similar to those written out but with the substitution

$$\begin{pmatrix} p \\ n \end{pmatrix} \to \begin{pmatrix} c \\ \lambda \end{pmatrix} \text{and} \begin{pmatrix} \nu_e \\ e \end{pmatrix} \to \begin{pmatrix} \nu_\mu \\ \mu \end{pmatrix}$$

of the field operators. The masses of the vector bosons in this scheme are determined by the relations

$$M_{Z_A} = M_{W_L^\pm}; \quad M_{Z_V} = M_{W_L^\pm}/\cos\theta_V;$$

$$4G_F/\sqrt{2} = g^2/M_{W_L^\pm}^2 = e^2/\sin^2\theta_V M_{W_L^\pm}^2,$$

where G_F is the Fermi coupling constant and $\tan\theta_V = \sqrt{3}g'/g$. If all the interactions are universal, i.e., if we have the equality $g[SU(2)] \equiv g = g[SU(4)] \equiv g'$, then the mixing angle θ_V and the masses of the vector bosons are completely determined: $\sin^2\theta_V = 3/4$, $M_{Z_A} \approx 62$ GeV, $M_{Z_V} \approx 2M_{Z_A} \approx 124$ GeV.[70] As a consequence of the equality of the coupling constants $g[SU(2)_L] = g[SU(2)_R] = g$, P parity is conserved in neutral-current interactions. Thus, in this scheme there need not in principle be P-violating effects in atomic transitions[119] due to the neutral currents and associated with weak-interaction processes of the form $l + N \to l + \ldots$ with lepton $l = \mu, e \ldots$. On the other hand, since only left-handed neutrinos (right-handed antineutrinos) participate in experiments with neutrinos, the neutral-current interaction containing the neutral current $\bar{\nu}\gamma_\mu(1+\gamma_5)\nu$ has the same form in this scheme as in the Weinberg–Salam–GIM model under the condition $\sin^2\theta_V = 2\sin^2\theta_W$.[70] Note that in neutral-current interactions exchange of the axial vector boson Z_A is dominant, since $M_{Z_V}^2 = 4M_{Z_A}^2$. If exchange of only the Z_A boson were important, then the cross sections of $\bar{\nu}_\mu N$ and $\nu_\mu N$ interactions through the neutral current would be the same.

Another requirement used in the construction of gauge schemes based on the symmetry group $SU(2) \times U(1)$ is the hypothesis that the neutral weak current is diagonal.[77] In accordance with this hypothesis, all nondiagonal neutral transitions between quarks (leptons) are suppressed. In the scheme with four quarks, this requirement reduces to partial suppression of the transitions $n \leftrightarrow \lambda$ and $p \leftrightarrow c$. In schemes with color symmetry of the type $SU(4)' \times SU(3)''$ of the Han–Nambu model, the diagonality principle consists of the requirement that the transitions $p_2 \leftrightarrow p_3$, $n_1 \leftrightarrow \lambda_1$, $n_2 \leftrightarrow n_3$, etc., be suppressed. This hypothesis is justified by the experimental data, which indicate suppression of the strangeness-changing neutral currents: The effective coupling constant of the transition $n \leftrightarrow \lambda$ satisfies the relation $G(|\Delta S| = 1) \leq 10^{-3}G(\Delta S = 0)$ (at least, at low energies). Similarly, there are restrictions on the existence of nondiagonal neutral currents in the lepton world as well [the $(\mu \leftrightarrow e)$ probability of the decay $\mu \to e + \gamma$ is $\lesssim 10^{-8}$]. The hypothesis of diagonality of the neutral currents is used naturally in the construction of the GIM model.[11] We emphasize that in the framework of the four-quark scheme the requirement that the charged and neutral currents satisfy the algebra $SU(2)$ leads to a different interesting possibility. Essentially, one can allow the existence of neutral transitions of the type $n \leftrightarrow \lambda$ and $p \leftrightarrow c$.[33–36,45] Moreover, a reasonable choice of the parameters of the theory in accordance with the existing experimental data makes it possible to suppress such transitions as much as one wishes. Therefore, experimental study of the nondiagonal neutral currents ($n \leftrightarrow \lambda$, $p \leftrightarrow c$, etc.) at high energies is of decisive importance for the choice of the variant of the theory.[78]

Note that if the charged current in any many-quark scheme is represented in the form $J_\mu^+ = \bar{Q}(q')MO_\mu Q(q)$, where Q is a column of quarks with the same charge, the operator O_μ denotes $V-A$ or $V+A$ space-time structure, and the matrix M determines all possible charged transitions between quarks, then the principle of diagonality of the neutral current will be satisfied if $M^+M = 1 (MM^T = 1)$, i.e., if the matrix M is unitary (orthogonal).[77]

The next important step is the suppression of transitions of the type $K^0 \leftrightarrow \bar{K}^0$ and $K_L^0 \to \mu^+\mu^-$ in the higher orders.[79–80] This condition, like the diagonality principle, can be generalized for all systems that are neutral with respect to the charge but charged with respect to the remaining quantum numbers (for example, $c_1\bar{p}_1 \leftrightarrow \bar{c}_1 p_1$; see Refs. 74 and 81). It is known that if in the simple four-quark model one introduces the right-handed current $(c\bar{n})_R$, it is impossible to suppress the transition without additional assumptions.

Schemes using color of the Han–Nambu type take into account all possible virtual transitions associated with exchange of states that are not necessarily color singlets. Therefore, in such schemes it is possible to require suppression of these neutral transitions.

As a result, it is found that the above requirement of suppression[74,81] of $K^0 \leftrightarrow \bar{K}^0$, $D^0 \leftrightarrow \bar{D}^0$ and all other transitions between systems which are neutral with respect to the charge but charged with respect to the other quantum numbers leads, if all the considered quarks have equal masses, to a simple relationship between the right- and left-handed current matrices carrying charge of the same kind:

$$M_R^+ = M_L^+ A_1; \quad M_R^- = M_L^- A_2, \quad (124)$$

where the minors of the matrices A_1 and A_2 are orthogonal and symmetric.[74,81]

Further particularization of the scheme entails a choice of the left-handed current. For example, it can be chosen by analogy with the GIM model. In this case, one can introduce into the theory without suppression only the right-handed current $(\bar{c}\lambda)_R$. One can ensure the necessary suppression of transitions of the type $\bar{K}^0 \leftrightarrow K^0$, $K_L^0 \to \mu^+\mu^-$. In quark schemes of the Han–Nambu type, one can also include an unsuppressed current $(\bar{c}n)_R$ in accordance with Eq. (124), but only in the case when the weak left- and right-handed currents

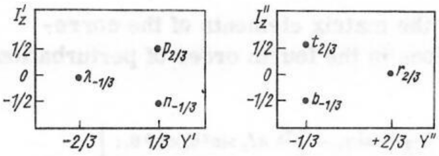

FIG. 3.

are not singlets of the color group.[81]

Construction of Six- and Eight-Quark Gauge Schemes and their Physical Consequences. The usual Weinberg–Salam–GIM model of weak interactions with four quark species p_i, n_i, λ_i, c_i (colored states) and four leptons $(\nu_\mu, \mu^-, \nu_e, e^-)$, leaves a number of difficulties unresolved, as we have already mentioned in the previous sections.

Recently obtained experimental data indicate the possible existence of "heavy" leptons and quarks.[82-89] This forces us to consider a scheme more complicated than the Weinberg–Salam–GIM model, a theory containing a greater number of quarks and leptons $(n_q > 4, n_l > 4)$.

One of the simplest six-quark schemes, which has all the positive features of the GIM scheme, is the scheme of Ref. 44, and also Refs. 121–123.

The six-quark model of Harari[44] contains the ordinary $SU(3)$ triplet of quarks (p, n, λ) and an $SU(3)$ antitriplet of heavy quarks. The latter contains an isodoublet (t, b) with charge content $(2/3, -1/3)$ and an isosinglet r with charge $+2/3$. The new quarks are ascribed an additive quantum number "heaviness" ($H = 1$ for heavy quarks and $H = 0$ for ordinary quarks). In addition, all quarks can be ascribed three colors, it being assumed that there are no colored composite hadrons in the free state. The quantum numbers of the quarks are conveniently represented in the form of the diagram in Fig. 3. Harari's model gives the ratio $R(s) = \sigma(e^+e^- \to \text{hadrons})/\sigma(e^+e^- \to \mu^+\mu^-) = 2$ below the production threshold of the heavy mesons, while above it $R = 5$. The meson spectrum contains 36 components for each J^P value: the ordinary octet and singlet of light ($H = 0$) mesons, nine heavy mesons ($H = -1$) contained in the $SU(3)$ multiplets $\overline{6}$ and 3, nine mesons with $H = -1$ from the $SU(3)$ multiplets 6 and $\overline{3}$, and also an octet and singlet of $H = 0$ mesons constructed from heavy quark–antiquark pairs.

The charged currents in Harari's model have the form[44,90]

$$J^+ = (p, t, r)\, M_{L,R} \begin{pmatrix} \overline{n} \\ \overline{\lambda} \\ \overline{b} \end{pmatrix}, \qquad (125)$$

where

$$M_L = \begin{pmatrix} \cos\theta_c; & -\sin\theta_c; & 0 \\ \cos\Phi\sin\theta_c; & \cos\Phi\cos\theta_c; & -\sin\Phi \\ \sin\Phi\sin\theta_c; & \sin\Phi\cos\theta_c; & \cos\Phi \end{pmatrix};$$

$$M_R = \begin{pmatrix} 0 & 0 & 1 \\ -\sin\varphi; & \cos\varphi; & 0 \\ \cos\varphi; & \sin\varphi; & 0 \end{pmatrix},$$

i.e.,

$$J(V+A) = (p\overline{b})_R + [t\overline{\lambda}(\varphi)]_R + [r\overline{n}(\varphi)]_R;$$
$$J(V-A) = [p\overline{n}(\theta_c)]_L$$
$$+ t(-\cos\Phi\sin\theta_c\overline{n} + \cos\Phi\cos\theta_c\overline{\lambda} - \sin\Phi\overline{b})_L$$
$$+ r(-\sin\Phi\sin\theta_c\overline{n} + \cos\Phi\overline{b} + \sin\Phi\cos\theta_c\lambda)_L.$$

It is readily noted that the condition of diagonality of the neutral current is satisfied if the matrices M_L and M_R are orthogonal. This ensures the absence of neutral transition currents with change $|\Delta S| = 1$ of the strangeness and $|\Delta H| = 1$ of the heaviness, and the particular form (125) leads to the selection rule $\Delta H = \Delta Q$ for heaviness-changing semileptonic decays. The angle θ_c in Eq. (125) is the Cabibbo angle. If the second parameter Φ is also assumed to be small, then the dominant weak left-handed transitions will be the transitions $p \leftrightarrow n, t \leftrightarrow \lambda, r \leftrightarrow b$.

The symmetry of Harari's model is the algebra $U(6)$. Subalgebras are, respectively, the algebras of operators that couple the light quarks to one another and the heavy quarks to one another. The symmetry here ensures conservation of the baryon number and the heaviness. The electric charge is equal to

$$Q = (Y_L + Y_H)/2 + (I_L^z + I_H^z) + H/3. \qquad (126)$$

The indices L and H here refer to the light and heavy quarks.

An interesting six-quark model was proposed by Barnett.[91] In the framework of this model, one introduces the usual triplet of quarks p, n, λ and a charmed triplet $\mathscr{P}, \mathfrak{N}, \Lambda$. These last quarks have the same quantum numbers as p, n, and λ, but in addition each of them has one further quantum number (charm), which is equal to unity. In addition, it is assumed that each quark species has three colors. As regards hadron spectroscopy in this model, there is in addition to the ordinary meson nonet ($C = 0$) a nonet with charm $C = +1$ (for example, $\mathscr{P} - \overline{n} = \Pi_1^+, \ldots$), a new nonet with $C = 0$ composed of charmed quark–antiquark pairs (for example, $\mathscr{P}, \overline{\mathfrak{N}} = \Pi_2^+$, etc.) and a nonet with $C = -1$ ($p\overline{\mathfrak{N}} = \pi_3^+$, etc.). The weak charged current, which in the framework of the Weinberg–Salam gauge symmetry leads to the absence of strangeness-changing terms in the corresponding neutral current, has in Barnett's model the form

$$\left.\begin{aligned} J_{\mu L} &= \overline{p}\,\gamma_\mu(1+\gamma_5)(n\cos\theta_c + \lambda\sin\theta_c) \\ &+ \overline{\mathscr{P}}\gamma_\mu(1+\gamma_5)(\lambda\cos\theta_c - n\sin\theta_c); \\ J_{\mu R} &= \overline{p}\gamma_\mu(1-\gamma_5)(\mathfrak{N}\cos\theta_a + \Lambda\sin\theta_a) \\ &+ \overline{\mathscr{P}}\gamma_\mu(1-\gamma_5)(\Lambda\cos\theta_a - \mathfrak{N}\sin\theta_a). \end{aligned}\right\} \qquad (127)$$

Below the threshold for the production of charmed particles, the current (127) leads to the standard results of the ordinary $SU(3)$ theory. The structure of the second term ensures cancellation of the strangeness-changing neutral currents.

Barnett's model in its fractionally charged variant leads to the value $R = 4$ for the process of e^+e^- annihilation into hadrons, while its variant with integral charges of the quarks gives $R = 8$.

In Refs. 71 and 74, analogous currents were considered in the framework of an eight-quark (eight-lepton) scheme with a light quartet of quarks $(p, n, \lambda, \varkappa)$

and a heavy quartet (t, b, d, r) (the additional quarks have the charges $Q_\varkappa = 2/3$ and $Q_d = -1/3$). These currents have the form

$$J_\mu\binom{L}{R} = (\bar{n}, \bar{\lambda}, \bar{b}, \bar{d})\, O_\mu^{L,R} M^{L,R} \begin{pmatrix} p \\ t \\ r \\ \varkappa \end{pmatrix}. \tag{128}$$

In Eq. (128), the matrices M^L and M^R are orthogonal four-dimensional matrices, each of which depends on six parameters.

As we have already seen, the many-quark models contain a certain arbitrariness because of the large number of free parameters. The usual requirements of the theory-conservation of the Cabibbo form for the currents involving the ordinary $SU(3)$ quarks, diagonality of the neutral current, and suppression of the $K^0 \leftrightarrow \bar{K}^0$ transition in the first and higher perturbation orders—still leave an appreciable arbitrariness in the choice.

In the listed schemes, one can construct systems that are neutral with respect to the electric charge but charged with respect to the other quantum numbers. These systems are analogs of the $SU(3)$-symmetric model of hadrons. For example, in the model with six quarks, such schemes have the following quark content:

$$b\bar{n}\,(\bar{b}n),\quad b\bar{\lambda}\,(\bar{b}\lambda),\quad t\bar{p}\,(\bar{t}p),\quad r\bar{p}\,(\bar{r}p),\quad t\bar{r}\,(\bar{t}r).$$

In what follows, we shall denote these systems by $D_i^0(\bar{D}_i^0)$. As in the case of the system $K^0(\bar{K}^0)$, the weak interaction will result in mixing of the states D_i^0 and $\bar{D}_i^0$.

To particularize the weak-interaction scheme determined by the currents (125), we require that for equality of the masses of all the quarks [in the limit of exact $SU(6)$ symmetry] the transitions $D_i \leftrightarrow \bar{D}_i$ be suppressed in higher orders. Allowance for breaking of the original exact symmetry leads to significant transitions of some of the systems D_i^0 and $\bar{D}_i^0$ into one another (the probabilities of such transitions are comparable with the decay probabilities of the corresponding states). The processes $D_i^0 \leftrightarrow \bar{D}_i^0$ can in particular be a source of production of pairs of charged leptons with the same charges in interactions of neutrinos with matter. As before, the transitions $K^0 \leftrightarrow \bar{K}^0$ remain suppressed.

As is shown in Refs. 71 and 74, simultaneous suppression of the transitions $K^0 \leftrightarrow \bar{K}^0$ and $D_i^0 \leftrightarrow \bar{D}_i^0$ is possible only if

$$\varphi - \Phi - \theta_c = 0. \tag{129}$$

The requirement of simultaneous suppression of mixings of the neutral K^0 and $\bar{K}^0$ and also D_i^0 and $\bar{D}_i^0$ mesons significantly restricts the values of the parameters φ, Φ, and θ_c and reduces to the condition of symmetry of the matrix A, where $M_L = AM_R$.

We consider the case of broken $SU(6)$ symmetry, when the quark masses differ from one another if the quarks belong to different quark triplets,[74] namely,

$$m_p \approx m_n \approx m_\lambda = m_1;$$
$$m_t \approx m_b \approx m_r = m_2,\quad m_2 \gg m_1.$$

In this case, for the matrix elements of the corresponding transitions in the fourth order of perturbation theory we find

$$\left.\begin{aligned}
&1.\ K^0 \leftrightarrow \bar{K}^0{:}\ S_4 = \bar{\lambda}\gamma^\alpha \frac{1+\gamma_5}{2} n\bar{\lambda}\gamma_\alpha \frac{1+\gamma_5}{2} nI_2 \sin^2\theta_c \cos^2\theta_c;\\
&2.\ D_1^0 \leftrightarrow \bar{D}_1^0{:}\ S_4 = \bar{n}\gamma^\mu\gamma^\nu \frac{\gamma^5}{2} \bar{b}n\gamma_\nu\gamma_\mu \frac{\gamma_5}{2} bI_1 \cos^2\theta_c;\\
&3.\ D_2^0 \leftrightarrow \bar{D}_2^0{:}\ S_4 = \bar{\lambda}\gamma^\mu\gamma^\nu \frac{\gamma^5}{2} b\bar{\lambda}\gamma_\nu\gamma_\mu \frac{\gamma_5}{2} bI_1 \sin^2\theta_c;\\
&4.\ D_3^0 \leftrightarrow \bar{D}_3^0{:}\ S_4 = \bar{p}\gamma^\mu\gamma^\nu \frac{\gamma^5}{2} t\bar{p}\gamma_\nu\gamma_\mu \frac{\gamma^5}{2} tI_1 \sin^2\Phi;\\
&5.\ D_4^0 \leftrightarrow \bar{D}_4^0{:}\ S_4 = \bar{p}\gamma^\mu\gamma^\nu \frac{\gamma^5}{2} r\bar{p}\gamma_\nu\gamma_\mu \frac{\gamma^5}{2} rI_1 \cos^2\varphi;\\
&6.\ D_5^0 \leftrightarrow \bar{D}_5^0{:}\ S_4 = r\gamma^\alpha \frac{1+\gamma_5}{2} \bar{t}r\gamma_\alpha \frac{1+\gamma_5}{2} tI_2 \sin^2\Phi \cos^2\Phi,
\end{aligned}\right\} \tag{130}$$

where we have introduced the notation

$$\left.\begin{aligned}
&I_1 \approx -\mathrm{i}\,\frac{g^4 m_2^2}{8m_W^4 (4\pi)^2}\left[\ln\frac{m_2^2}{m_W^2} + \frac{\sqrt{a-4B}}{\sqrt{a}} \ln\frac{\sqrt{a}+\sqrt{a-4B}}{\sqrt{a}-\sqrt{a-4B}}\right];\\
&a = (k_2 - k_4)^2/m_W^2;\quad B = m_2^2/m_W^2.
\end{aligned}\right\} \tag{131}$$

For $(k_2 - k_4)^2 = 0$ [$k_i (i = 1, 2, 3, 4)$ are the quark momenta (see Ref. 74)],

$$I_1 \approx -\mathrm{i}\,(g^4 m_2^2/8m_W^4 (4\pi)^2)\,[2 + \ln(m_2^2/m_W^2)]; \tag{132}$$
$$I_2 \approx -\mathrm{i}\,(g^4 m_2^2/8m_W^4)\,(4\pi)^2. \tag{133}$$

Thus, if allowance is made for the mass difference of the quark triplets, the $K^0 \leftrightarrow \bar{K}^0$ transition remains suppressed as before, namely

$$(m_L - m_S)/m_{K^0} = (g^4 m_2^2/8m_W^4 (4\pi)^2)\cos^2\theta_c \sin^2\theta_c f_K^2, \tag{134}$$

where m_L and m_S are the masses of the K_L^0 and K_S^0 mesons. At the same time, irrespective of the choice of the parameter Φ, there always exists, as can be seen from Eqs. (130), a system D_i^0 whose transition into $\bar{D}_i^0$ will take place without additional suppression. For example, $D_1^0 = (b\bar{n})$,

$$(m_{D_1L} - m_{D_1S})/m_{D_1} \approx \mathrm{i} I_1 \cos^2\theta_c f_{D_1}^2. \tag{135}$$

The mixing of such D_1^0 and $\bar{D}_1^0$ states may be the mechanism that could lead to dimuonic events in neutrino experiments with the same sign of the charges of the muons. Assuming that the parameter Φ is zero, $\Phi = 0$, one can follow the interesting interconnection between the left- and right-handed currents [the equation $\Phi = 0$ leads to correspondence of this scheme to the GIM model ($t \equiv c$)]:

$$J_L = \bar{p}n\,(\theta_c) + \bar{r}b + \bar{t}\lambda\,(\theta_c); \tag{136}$$
$$J_R = \bar{p}b + \bar{r}n\,(\theta_c) + \bar{t}\lambda\,(\theta_c). \tag{137}$$

If the theory of weak and electromagnetic interactions is constructed on the basis of the gauge symmetry $SU(2)_L \times SU(2)_R \times U(1)$ (see Refs. 71 and 72), then, considering the mixing of the charged $W_L^\pm$ and $W_R^\pm$ bosons,

$$\left.\begin{aligned}
W_L^\pm &= W_1^\pm \cos\alpha + W_2^\pm \sin\alpha;\\
W_R^\pm &= -W_1^\pm \sin\alpha + W_2^\pm \cos\alpha,
\end{aligned}\right\} \tag{138}$$

one can obtain the following quark Lagrangian for the interaction of the currents with the gauge bosons [cf. Eq. (121)]:

$$\begin{aligned}
\mathcal{L} = e\Big(&\tfrac{2}{3}\bar{p}\gamma_\mu p - \tfrac{1}{3}\bar{n}\gamma_\mu n + \tfrac{2}{3}\bar{t}\gamma_\mu t - \tfrac{1}{3}\bar{\lambda}\gamma_\mu\lambda\\
&+ \tfrac{2}{3}\bar{r}\gamma_\mu r - \tfrac{1}{3}\bar{b}\gamma_\mu b\Big) A_\mu + \frac{g}{2}(\bar{p}\gamma_\mu\gamma_5 p - \bar{n}\gamma_\mu\gamma_5 n + \ldots)\, Z_\mu^A\\
&+ \frac{\sqrt{g^2+\tilde{g}^2}}{2}(\bar{p}\gamma_\mu p - \bar{n}\gamma_\mu n + \ldots - 2\sin^2\theta_V J_\mu^{\text{e.m.}})\, Z_\mu^V\\
&+ \frac{g\cos\alpha}{2} J_{\mu L}^+ W_1^- - \frac{g\sin\alpha}{2} J_{\mu R}^+ W_1^-\\
&+ \frac{g\sin\alpha}{2} J_{\mu L}^+ W_2^- + \frac{g\cos\alpha}{2} J_{\mu R}^+ W_2^- + \text{h.c.}
\end{aligned} \tag{139}$$

The restrictions on the masses of the intermediate vector bosons in the case of the Lagrangian (139) differ somewhat from those given earlier [see Eq. (121)]:

$$M^2_{W_2} \gg M^2_{W_1} \frac{\sqrt{2}\, e^2 \cos^2\alpha}{4G \sin^2\theta_V}\,;\; M_{Z_V} = \frac{M_{W_1}}{\cos\alpha \cos\theta_V}\,;\; M_{Z_A} = M_{Z_V}\cos\theta_V, \tag{140}$$

where G is the Fermi coupling constant of the weak interaction, and the angle θ_V is as before related to the Weinberg angle by $\sin^2\theta_V = 2\sin^2\theta_W$. Substituting in Eq. (140) the numerical value $\sin^2\theta_W = 3/8$ of the Weinberg angle, we obtain the following bounds for the possible masses of the intermediate bosons, these depending on the mixing angle of the left- and right-handed worlds:

$$45\ \text{GeV} \leqslant M_{W_1} \leqslant 62\ \text{GeV},\ M_{Z_A} \approx 62\ \text{GeV}, M_{Z_V} \approx 124\ \text{GeV}\,. \tag{141}$$

Note that the lower bound for the mass of the charged W_1 boson, $M_{W_1} \sim 45$ GeV, corresponds to maximal mixing ($\alpha = \pi/4$; see Ref. 72).

The given mass values of the intermediate gauge bosons can be obtained by spontaneously breaking the original symmetry of the Lagrangian by the Higgs mechanism. Thus, mixing of the intermediate bosons $W^\pm_L$ and $W^\pm_R$ can be achieved by using the adjoint representation $(3, \overline{3})$ of scalar fields of the group $SU(2)_L \times SU(2)_R$.[72] Note that at very high energies, where it is necessary to take into account the exchange of $W^\pm_2$ bosons, mixing between the right- and left-handed worlds disappears. It is of interest to write down the Hamiltonian of the weak nonleptonic interaction due to exchanges of charged $W^\pm_1$ bosons that follows from the form of the Lagrangian (139):

$$\mathscr{H}^{\text{eff}} = (G/\sqrt{2})\{J_L J_L^+ + \text{tg}\,\alpha J_L J_R^+ + J_R J_L^+ + \text{tg}^2\alpha J_R J_R^+\} + \text{h.c.}, \tag{142}$$

where J_L and J_R are determined by (136) and (137). In the presented approach, there is a space-time $V+A$ coupling of the "ordinary" λ quark to the heavy $t(c)$ quark. As was noted above, there is a number of indications that in a theory with four quarks (p, n, λ, c) one can take into account precisely this current. Of course, it is understood that the magnitude $\tan\alpha$ of the mixing cannot appreciably differ from unity. On the other hand, the interpretation of the resonance recently observed by Lederman in the $\mu^+\mu^-$ system with mass 9.5 GeV (Ref. 88) as a state analogous to the ψ meson consisting of heavy quarks b and $\overline{b}$ ($Q_b = -1/3$ and $m_b = 5$ GeV) indicates that the mixing angle α is somewhat less than $\pi/4$, since otherwise we should observe an appreciable raising of the total cross sections for the interactions of antineutrinos with matter.

Note that the suppression of $D^0_i \leftrightarrow \overline{D}^0_i$ transitions in the higher perturbation orders in the case when all the considered quarks have equal masses is valid for any mixing angle α,[74] although the effect is of interest for not very small values of α. For example, the presence in the theory of right-handed currents leading to transitions between the heavy and light quarks, and also appreciable mixing α, has the consequence that weak decays of hadronic systems formed from the heavy quarks b, t, r occur with an appreciable effective coupling constant, which corresponds to short lifetimes of these systems ($10^{-14} - 10^{-16}$ sec) if the b quarks have masses ≈ 5 GeV. On the other hand, without mixing between the left- and right-handed worlds ($\alpha = 0$) some of these systems may be long-lived, i.e., they may decay in the higher perturbation orders.

It can be seen from the relations (136) and (137) that one species of current goes over into the other if the light quark (p) is replaced by the heavy quark (r), and vice versa. The p and r quarks are in the isodoublets (p, n) and (r, b), respectively, whereas the components of the currents (136) and (137) composed of the isosinglets t and λ remain unchanged. Transition from this scheme to the eight-quark scheme with four light (p, n, λ, c) and four heavy $(\mathscr{P}, \mathfrak{N}, \Lambda, C)$ quarks leads in the framework of the indicated connection to the currents[71,74]

$$J_L = p\overline{n}\,(\theta_c) + \mathscr{P}\overline{\mathfrak{N}}\,(\theta_c) + c\overline{\lambda}\,(\theta_c) + C\overline{\Lambda}\,(\theta_c); \tag{143}$$

$$J_R = \mathscr{P}\overline{n}\,(\theta_c) + p\overline{\mathfrak{N}}\,(\theta_c) + C\overline{\lambda}\,(\theta_c) + c\overline{\Lambda}\,(\theta_c). \tag{144}$$

The substitutions $p \leftrightarrow \mathscr{P}$ and $c \leftrightarrow C$ realize the connection between the left- and right-handed currents, $J_R \leftrightarrow J_L$. The scheme with the currents (143) and (144) satisfies all the requirements discussed above. In particular, $K^0 \leftrightarrow \overline{K}^0$ transitions are suppressed, whereas $D^0_i \leftrightarrow \overline{D}^0_i$ meson mixing occurs very strongly for some systems D^0_i.

In the eight-quark scheme, the c quark is assumed to be heavier than its partners in the fundamental $SU(4)$ representation, but it can be assumed to be lighter than the quarks $\mathscr{P}, \mathfrak{N}, \Lambda, C$. The production threshold of heavy particles with $H = 1$ may therefore lie far above the threshold for the production of charmed particles containing the c quark with $H = 0$; for example, it may lie above the $\overline{c}p$ meson production threshold.

In the framework of the quark–lepton analogy, it is convenient to consider also the eight-lepton variant of the scheme (143)–(144) with four heavy leptons M^0, M^-, E^0, E^-. The leptonic currents in this case have the form[71-74]

$$J^l_L = e^-\overline{\nu}_e\,(\theta_l) + \mu^-\overline{\nu}_\mu\,(\theta_l) + E^-\overline{E}^0\,(\theta_l) + M^-\overline{M}^0\,(\theta_l); \tag{145}$$

$$J^l_R = E^-\overline{\nu}_e\,(\theta_l) + M^-\overline{\nu}_\mu\,(\theta_l) + e^-\overline{E}^0\,(\theta_l) + \mu^-\overline{M}^0\,(\theta_l) \tag{146}$$

and under the substitution $e^- \leftrightarrow E^-, \mu^- \leftrightarrow M^-$ we arrive at symmetry between the left- and right-handed leptonic currents: $J^l_L \leftrightarrow J^l_R$. The quark–lepton analogy can be taken further in view of the electron and muon quantum numbers of the leptons. We shall call $p, n(e, \nu_e)$ and $c, \lambda(\mu, \nu_\mu)$ light quarks (respectively, leptons) and $\mathscr{P}, \mathfrak{N}\,(E^-, E^0)$ and $C, \Lambda(M^-, M^0)$ heavy quarks (respectively, leptons). We ascribe the quarks (respectively, leptons) $p, n, \mathscr{P}, \mathfrak{N}, (e, \nu_e, E^-, E_0)$ the electron quantum number, and the quarks (respectively, leptons) $c, \lambda, C, \Lambda(\mu, \nu_\mu, M^-, M^0)$ the muon (strange) quantum number. It was noted in Ref. 92 that the angle θ_l that mixes the electron and muon quantum numbers plays a role in the leptonic currents analogous to that of the Cabibbo angle θ_c in the hadronic current, which leads to strangeness- and charm-changing transitions. In the eight-quark (eight-lepton) scheme of weak interactions considered here, one can assume the existence

TABLE II.

Electron		Muon		H
quarks	leptons	quarks	leptons	
p n	e ν_e	c λ	μ ν_μ	0, light quarks and leptons
$\mathcal{P}$ $\mathfrak{N}$	E^- E^0	C Λ	M^- M^0	1, heavy quarks and leptons

of such an analogy between the angles θ_l and θ_c. The suppression of the weak currents responsible for transitions between quarks (leptons) from different columns of Table II is ensured by the factors $\sin^2\theta_c$ ($\sin^2\theta_l$). The neutral nondiagonal transitions of the type $n \leftrightarrow \lambda$ or $e \leftrightarrow \mu$ are also suppressed naturally in this approach.

Weak interactions of the type (143)–(146) can be considered in the framework of the gauge symmetry $SU(2)_L \times SU(2)_R \times U(1)$. Breaking this symmetry by means of the Higgs mechanism to the symmetry $U(1)$, we conclude that interactions between the neutral currents of the quarks and leptons will be realized by two heavy bosons Z_A and Z_V, the masses of these bosons being $m_{Z_A} = m_{W^\pm_{L,R}}$ and $m_{Z_V} = (\sqrt{g^2+g'^2}/g')m_{Z_A}$. By means of the constant g' corresponding to the symmetry $U(1)$, and also by virtue of the fact that neutrino beams are obtained from the decays $\pi \to \mu\nu$ and $K \to \mu\nu$, in which the neutrino (antineutrino) has left-handed (respectively, right-handed) helicity, one can satisfactorily describe[70] the existing data on the $\nu_\mu(\bar{\nu}_\mu) + \mathcal{N} \to \nu_\mu(\bar{\nu}_\mu) + X$ reaction.[76]

Thus, in the framework of $SU(2)_L \times SU(2)_R \times U(1)$ gauge symmetry and the eight-quark (eight-lepton) scheme one can construct a model of the weak interactions that satisfies the following conditions.

1. Below the threshold for the production of heavy particles with $|H| = 1$ it coincides for charged currents with the GIM mechanism.

2. It enables one to include the right-handed currents $(p\bar{q}_1)_R$ and $(n\bar{q}_2)_R$, where q_1 and q_2 are new quarks with charges $-1/3$ and $2/3$, respectively, and a current of the type $(p\bar{q}_1)_R$ can lead to a growth with the energy of the cross section ratio $\sigma(\bar{\nu} + \mathcal{N} \to \mu^+ + X)/\sigma(\nu + \mathcal{N} \to \mu^- + X)$. The degree of growth depends on the extent of mixing of the charged states of the W bosons corresponding to the symmetry groups $SU(2)_L$ and $SU(2)_R$.

Right-handed currents involving heavy leptons can lead to the production of "direct" leptons with unusual polarization in hadron–hadron collisions.[90,93-96]

3. It enables one to explain why the neutral currents in the $\nu\mathcal{N}$ interactions can coincide above the heavy-particle production threshold with the predictions of the GIM mechanism. Note that in the weak hadron–hadron transitions due to neutral currents, parity may still be violated when allowance is made for the quark masses. The viability of the scheme with the two neutral Z_A and Z_V bosons could be tested by investigating effects of nonconserving of P parity in muon–nucleon experiments with change in the polarization of the primary muons from the decays $\pi \to \mu\nu$ due to their kinematic depolarization in the laboratory system.[97]

4. It contains the possibility of mixing of the electron and muon quantum numbers determined by the angle θ_l. For weak interactions that include the light leptons $(\nu_\mu, \mu^-, \nu_e, e^-)$, the mixing angle leads to interesting consequences, these concerning above all $\nu_\mu \leftrightarrow \nu_e$ neutrino oscillations.[92]

Inclusion of the mixing angle θ_l in the right-handed leptonic currents containing transitions $M^0(\theta) \to \mu$ and $E^0(\theta) \to e$ of the heavy leptons is of interest from the point of view of the search for $\mu \to e\gamma$ and $\mu \to 3e$ decays. For example, the fraction of $\mu \to e\gamma$ decays relative to the basic decay $\mu \to e\nu\bar{\nu}$ is, as one can show,[92]

$$\frac{\Gamma(\mu \to e\gamma)}{\Gamma(\mu \to e\nu\tilde{\nu})} = \frac{3}{32}\frac{\alpha}{\pi}\left(\frac{M_1^2 - M_2^2}{M_{W_1}^2}\right)^2 \sin^2\theta_l \cos^2\theta_l \,\mathrm{tg}^4\alpha. \tag{147}$$

The mass values of the heavy leptons M^0 and E^0, respectively, are equal to M_2 and M_1. The expression (147) contains the quantity $\tan\alpha$, which characterizes the mixing between the $W^\pm_L$ and $W^\pm_R$ bosons [see Eq. (138)].[72] Other possibilities for the existence of a nonzero mixing angle θ_l can appear in decays of hadrons with a μe pair in the final state (for example, the decay $K^0_L \to \mu e$ or

$$D^- \to \bar{E}^0 e^- + \cdots,\quad \bar{E}^0 \to e^+ + \cdots,\quad \to \mu^+ + \cdots).$$

But if the lepton mixing angles are different for the left- and right-handed currents $(\theta_l^L \neq \theta_l^R)$, then if we have, say, a partly or fully polarized muon beam (to be specific with right-handed polarization) we arrive at the possible existence of processes of the form

$$\mu^-_R + \mathcal{N} \to M^0(\theta_e^R) + \cdots,\quad M^0 \to \mu^- + \cdots,\quad \to e^- + \cdots,$$

i.e., experimentally one can observe an "apparent" violation of the law of conservation of the muon and electron quantum numbers associated with the difference $(\theta_l^L - \theta_l^R) \neq 0$ of the mixing angles.[72]

In the schemes of weak interactions considered above, the diagonality hypothesis was used in the most rigorous sense. However, the currently known experimental confirmation of this hypothesis is insufficient. The suppression of the neutral nondiagonal transitions $n \leftrightarrow \lambda(\mu \leftrightarrow e)$ may have deep physical reasons, a reflection of which is the existence of the hadronic Cabibbo angle and its leptonic analog θ_l.

The existence of a mixing angle of hadrons (leptons) with different hypercharges may essentially mean that there can exist several worlds in nature[72,74,87]:

1) a nonstrange world (electronic) of quarks (leptons):

$$\begin{pmatrix} p, n, \mathcal{P}, \overline{\mathfrak{N}} \ldots \\ e^-, \nu_e, E^-, E^0 \ldots \end{pmatrix}$$

with strangeness (respectively, muon quantum number) equal to zero;

2) a strange–charmed (muon) world of quarks (leptons) with electron number equal to zero for the leptons and nonzero values of the strangeness or charm for the quarks:

$$\begin{pmatrix} c, & \lambda, & C, & \Lambda \ldots \\ \mu, & \nu_\mu, & M^0, & M^- \ldots \end{pmatrix};$$

3) a world characterized by a new quantun number, a possible manifestation of which is the detection of the τ lepton:

$$\begin{pmatrix} \ldots & ? \\ \nu_\tau, & \tau \ldots \end{pmatrix}.$$

Transitions between quarks (leptons) from the different worlds, even through charged weak currents, are then due to the Cabibbo angle θ_c and its leptonic analog θ_l. Neutral nondiagonal transitions between quarks (leptons) from different worlds are suppressed much more strongly. The question of the nature of this suppression is evidently related to the currently unknown internal structure of the quarks (leptons). On the other hand, neutral transitions between quarks (leptons) within the same world may be observed at a given stage in the development of the experiments. For example, this applies to the currents $(\overline{\nu}_\mu M^0)$, $(\overline{\nu}_e E^0)$, $(\overline{\mu}^- M^-)$, $(\overline{e}^- E^-)$, $(\overline{p}\mathcal{P})$, $(\overline{c}C)$, $(\overline{n}\mathfrak{N})$, $(\overline{\lambda}\Lambda)$. The mixing angle of two different worlds may be appreciable precisely for the leptons. It is helpful to give an example of weak interactions that takes into account these possibilities. We put the leptons in doublets of the group $SU(2)_L \times SU(2)_R$, in accordance with

$$\begin{pmatrix} \nu_e & E^0 \\ e^- & E^- \end{pmatrix}_{L,R}, \quad \begin{pmatrix} \nu_\mu & M^0 \\ \mu^- & M^- \end{pmatrix}_{L,R}, \quad \begin{pmatrix} \nu_\tau & \mathcal{T}^0 \\ \tau^- & \mathcal{T}^- \end{pmatrix}_{L,R}. \tag{148}$$

In constructing the scheme (148), we have used a "weakened" diagonality principle, i.e., we have assumed the suppression of only those nondiagonal transitions in which the lepton quantum number is not additionally conserved.

The further development of the theoretical schemes in the framework of six, eight, etc., quarks is to increase the rank of the gauge group. For example, in Harari's six-quark model one can construct a scheme of weak interactions based on the gauge group $SU(3)_L \times SU(3)_R$. For this, we must place the considered quarks (leptons) in representations of the group; for example, the quarks in the model of Ref. 44 can be put in triplets as follows:

$$\begin{pmatrix} p \\ n \\ \lambda \end{pmatrix}_L \begin{pmatrix} t \\ b \\ r \end{pmatrix}_L \begin{pmatrix} t \\ \lambda \\ n \end{pmatrix}_R \begin{pmatrix} p \\ r \\ b \end{pmatrix}_R.$$

A natural generalization of the weak interaction scheme in the framework of eight quarks is to consider the gauge group $SU(4)_L \times SU(4)_R$. In this case, the quarks can be placed in quadruplets, for example,

$$\begin{pmatrix} p \\ n \\ \mathfrak{N} \\ \mathcal{P} \end{pmatrix}_{L,R} \begin{pmatrix} c \\ \lambda \\ \Lambda \\ C \end{pmatrix}_{L,R}.$$

Note that the appearance of several left-handed (right-handed) multiplets in both the six- and eight-quark models is obviously not fortuitous. Indeed, the discovery of the τ lepton may well indicate the existence of multiplets with a new quantum number $(\tau^-, \nu_\tau, \ldots)$. Of course, this fact calls for careful experimental verification; it is necessary to study decay channels of the τ lepton such as $\tau^- \to e^-\gamma$, $\tau \to \mu\gamma$, $\tau \to 3e$, $\tau \to 3\mu$.

4. SOME EXPERIMENTAL CONSEQUENCES OF GAUGE SCHEMES OF THE WEAK INTERACTION

Heavy Leptons (Quarks)[101,102] *and Interpretation of Many-Lepton Events in Neutrino Experiments.* One of the first indications of the possible existence of particles with a new hidden quantum number (charm, beauty, etc.) was provided by the results of neutrino experiments on the production of pairs of charged leptons (muons).[103-106] In the meanwhile, the following restrictions on the cross sections for the production of lepton pairs in neutrino and antineutrino experiments on nucleon targets have been established:

$$\left.\begin{aligned} &\frac{\sigma^\nu(\mu^+\mu^-)}{\sigma^\nu(\mu^-)} \sim 10^{-2}; \quad \frac{\sigma^\nu(\mu^-\mu^-)}{\sigma^\nu(\mu^-)} \sim 10^{-3}; \quad \frac{\sigma^\nu(\mu^-e^+)}{\sigma^\nu(\mu^-)} \sim 5\cdot 10^{-3}; \\ &\frac{\sigma^{\bar\nu}(\mu^+\mu^-)}{\sigma^{\bar\nu}(\mu^+)} \sim 8\cdot 10^{-3}; \quad \frac{\sigma^{\bar\nu}(\mu^+\mu^+)}{\sigma^{\bar\nu}(\mu^+)} \leqslant 10^{-3}; \quad \frac{\sigma^{\bar\nu}(\mu^+e^-)}{\sigma^{\bar\nu}(\mu^+)} \sim 5\cdot 10^{-3}. \end{aligned}\right\} \tag{149}$$

According to the latest estimates, the cross section for the production of charmed (beautiful) hadrons in $\nu_\mu N$ interactions at energies of order 100 GeV of the original neutrino beam is 15–20% of the total cross section of the $\nu_\mu N$ interaction. Therefore, if the semileptonic decay mode of the charmed (beautiful) mesons (baryons) is 5–10%, this assumption agrees with the experimental results (149).

The appearance of dimuonic events with muons of one sign in the framework of the four-quark scheme (p, n, λ, c) is hard to explain through the $D_0 \leftrightarrow \overline{D}{}^0$ regeneration process, since this regeneration is strongly suppressed in the GIM scheme. However, the associated production of charmed particles in the neutrino interaction with cross section ~1% of the total cross section can completely explain the appearance of dileptons with leptons of one sign. The same hypothesis does not contradict the results of experiments in which there has been observed the production of pairs and triplets of muons in the final state[107] in μN interaction processes:

$$\frac{N(3\mu)}{N(\mu)} \sim 5\cdot 10^{-4}; \quad \frac{N(2\mu)}{N(\mu^-)} \sim 10^{-3}. \tag{150}$$

Note that in many-quark schemes with right-handed currents one can ensure the regeneration of several neutral systems (D_i^0 mesons) at a fairly high level (see Sec. 3). It is also necessary to point out that the observation of dilepton pairs with different signs of the charges of the leptons can be interpreted as the production and subsequent decay of heavy leptons, for example,

$$\begin{array}{l} \nu_\mu + \mathcal{N} \longrightarrow M^0 + \cdots \\ \qquad\qquad\quad \hookrightarrow \mu^- \mu^+ \nu \\ \qquad\qquad\quad \hookrightarrow \mu^- e^+ \nu \\ \qquad\qquad\quad \hookrightarrow \mu^- E^0 e^+ \\ \qquad\qquad\qquad\quad\;\; \hookrightarrow e^- \mu^+ \nu . \end{array} \tag{151}$$

There is also another possible mechanism for the appearance of many-lepton events associated with decays of hadrons with new quantum numbers through channels with the production of a neutral heavy lepton[90]:

$$\begin{array}{l} \nu_\mu + \mathcal{N} \longrightarrow \mu^- + X^+ \\ \qquad\qquad\qquad\quad \hookrightarrow M^0 \; \nu_\mu + \text{hadrons} \\ \qquad\qquad\qquad\quad \hookrightarrow M^0 \mu^+ \\ \qquad\qquad\qquad\quad \hookrightarrow \mu^- + \cdots \\ \qquad\qquad\qquad\quad \hookrightarrow \mu^- + \mu^+ + \nu . \end{array} \tag{152}$$

The production of a heavy lepton in the leptonic block of the $\nu_\mu N$ interaction with the simultaneous production of a hadron with new quantum number in the hadronic block may lead to many-lepton events through the above cascade of decays. Thus, the muon triplets $(\mu^- \mu^+ \mu^-)$ and $(\mu^- \mu^+ \mu^+)$ observed in the experiment of Refs. 103 and 108 can be interpreted as the production and decay of charged and neutral heavy leptons.[109] In the case of production of the triplet $(\mu^- \mu^- \mu^+)$,

$$\begin{array}{l} \nu_\mu + \mathcal{N} \longrightarrow M^- + \text{hadrons} \\ \qquad\qquad\quad \hookrightarrow M^0 \mu^- \; \nu_\mu \\ \qquad\qquad\qquad \hookrightarrow \mu^- \mu^+ \nu_\mu . \end{array} \tag{153}$$

and in the case of production of the triplet $(\mu^- \mu^+ \mu^+)$,

$$\begin{array}{l} \nu_\mu + \mathcal{N} \longrightarrow M^0 \quad + \quad X^+ + \text{hadrons} \\ \qquad\qquad\quad \hookrightarrow \mu^- \mu^+ \nu_\mu \;\; \hookrightarrow M^0 \mu^+ \\ \qquad\qquad\qquad\qquad\qquad\quad \hookrightarrow \nu_\mu + \text{hadrons} . \end{array} \tag{154}$$

The collection of mechanisms (153) and (154) can lead to the production of a large number of leptons $(n>3)$ in the final state.

Note that it is most natural to look for heavy leptons precisely in neutrino experiments. In the gauge schemes considered above, there have also been considered heavy charged leptons and heavy neutral leptons with both muon and electron quantum numbers (M^-, E^-, M^0, E^0). There exists one further interesting possibility associated with the inclusion in the theory of a larger number of leptons [six leptons with muon (electron) quantum number], which presupposes the existence of two species of charged lepton. The $SU(2)$ gauge schemes in this case contain lepton doublets of the type

$$\begin{pmatrix} \nu_\mu(\theta) \\ \mu^- \end{pmatrix}_L ; \quad \begin{pmatrix} M^0(\theta) \\ M^- \end{pmatrix}_L ; \quad \begin{pmatrix} \tilde{M}^0(\theta) \\ \mu^- \end{pmatrix}_R ; \begin{pmatrix} N^0_\mu(\theta) \\ M^- \end{pmatrix}_R , \tag{155}$$

where θ is the angle of the mixing of the electron and muon lepton worlds. Leptons with electron quantum number are put in doublets similarly. One can also have an $SU(2)$ gauge scheme that contains nondiagonal neutral transitions between the leptons listed above:

$$\begin{pmatrix} (\cos\alpha\nu_\mu + \sin\alpha M^0)(\theta) \\ \cos\beta\mu^- + \sin\beta M^- \end{pmatrix}_L ; \quad \begin{pmatrix} (\cos\tilde{\alpha}\tilde{M}^0 + \sin\tilde{\alpha} N^0_\mu)(\theta) \\ \cos\tilde{\beta}\mu^- + \sin\tilde{\beta} M^- \end{pmatrix}_R \tag{156}$$

and similar doublets of leptons with electron quantum numbers.

An indication of the existence of nondiagonal leptonic transitions was obtained, as we have noted, in the neutrino experiment with the SKAT chamber,[85-87] in which there may have been an observation of the reaction

$$\begin{array}{l} \nu_\mu + P \to M^0 + 2\pi^0 + P . \\ \qquad\qquad\; \hookrightarrow \mu^- e^- \nu \end{array} \tag{157}$$

The fundamental importance of these many-lepton schemes arises because they contain at least three species of neutral heavy leptons $M^0, \tilde{M}^0, N^0_\mu$.

It is quite possible that the existence of heavy neutral leptons has been indicated by experiments on the detection of "direct" leptons in hadron–hadron collisions and the measurement of their polarization (μ^+) at large transverse momenta: $p_\perp = 2-2.8$ GeV/c.[93,98,99] The production of new particles (for example, charmed mesons) in hadron–hadron collisions with subsequent decays into heavy leptons:

$$\begin{array}{l} X \to M^0\mu^+ + \ldots;\; E^0 e^+ + \ldots;\; M^0 \to \mu^- + \ldots \\ \;\hookrightarrow M^0 \tilde{\nu}_\mu + \ldots;\; E^0 \tilde{\nu}_e + \ldots;\; E^0 \to e^- + \ldots, \end{array} \tag{158}$$

where the currents $(\bar{M}^0 \mu^-)_R$ and $(\bar{E}^0 e^-)_R$ are right-handed, could explain the observed longitudinal polarization of the μ^+.[90,94-96] According to the latest results[100] [the preliminary result[98,99] was $p_\parallel(\mu^+) = -(0.85 \pm 0.35)$],

$$p_\parallel(\mu^+) = -(0.41 \pm 0.17). \tag{159}$$

If the X mesons are identified with charmed mesons whose mass does not exceed 2 GeV, then if the decays (158) are to occur the mass of M^0 must not exceed 2 GeV.

Analysis of the event in the SKAT chamber made it possible to establish restrictions on the mass of the putative M^0 lepton[85-87]: $1.4 \leq m(M^0) \leq 2.1$ GeV. An estimate of the M^0 lifetime from the observed range of its decay products gives $\tau \sim (5-7) \times 10^{-12}$ sec. If the M^0 mass is near the upper limit, the decays of the charmed particles through the channel (158) may be suppressed and therefore cannot make an appreciable contribution to the yield of "direct" leptons produced in $P\mathcal{N}$ collisions. Then the situation with regard to the μ^+ polarization in $P\mathcal{N}$ collisions could be understood if one assumes the existence of a further heavy neutral lepton $\tilde{M}^0$ as in the scheme (156), whose mass must be appreciably less than the mass of the M^0. At the same time, as follows from the scheme (156), the $\tilde{M}^0(\tilde{E}^0)$ leptons must not be produced in neutrino experiments. The observation of such leptons is then possible in $\mu(e)\mathcal{N}$ interactions. Accordingly, the lifetime of the $\tilde{M}^0(\tilde{E}^0)$ leptons must be appreciably greater than τ_{M0} (τ_{E0}).

The production in $\nu_\mu \mathcal{N}(\bar{\nu}_\mu \mathcal{N})$ interactions of heavy leptons, and also the existence of nondiagonal neutral leptonic currents, leads to a certain change in the expected values of the cross sections of $\nu_\mu \mathcal{N}(\bar{\nu}_\mu \mathcal{N})$ interactions in reactions with both charged and neutral currents. For example, the production of M^0 in the reaction $\nu_\mu + \mathcal{N} \to M^0 + \ldots$ with the subsequent decays $M^0 \to \mu^- + \ldots$ and $M^0 \to \nu_\mu + \ldots$ will give formally contributions to the total cross section of reactions with both charged and neutral currents, and in the case of the $V-A$ variant the contribution to the cross section of reactions with neutral currents at energies far from the M^0 production threshold will be three times as large as the contribution to the cross section of reactions induced by antineutrinos. However, measurement of the total cross sections of $\nu_\mu \mathcal{N}(\overline{\nu}_\mu \mathcal{N})$ interactions through charged currents in the energy range 100–250 GeV indicates that these contributions are themselves slight, which then indicates either small coupling constants of the transitions $\nu_\mu \leftrightarrow M^-$ and $\nu_\mu \leftrightarrow M^0 (G_{\text{eff}} \ll G_F)$, or an appreciable mass of the heavy lepton $M^-[m(M^-) \geq 8 \text{ GeV}]$.

Since the heavy leptons with electron quantum number may be appreciably lighter than the heavy muon leptons, it would be interesting to make corresponding searches in beams of high-energy electron neutrinos, an admixture of which is always present in a ν_μ beam.[4)] That this suggestion is reasonable is also indicated by experiments in colliding e^+e^- beams, in which μe pair production processes are explained by the production of the heavy charged leptons τ with mass ≈ 1.9 GeV. This lepton may either be ascribed a new quantum number or a muon or electron quantum number. In the latter case, the $\tau^\pm$ in the scheme with eight leptons may occupy the position of $E^\pm$.[71,74]

Behavior of Total Cross Sections of $\nu_\mu \mathcal{N}(\overline{\nu}_\mu \mathcal{N})$ Interactions and Production of Heavy Quarks. The possible existence of new heavy quarks and heavy leptons[82,85,86,88] (besides the charmed quarks) poses the experimentalists a number of problems associated with their detection. In the considered gauge schemes, there exist both left- and right-handed transition currents of the ordinary valence quarks (p, n) of the nucleon into heavy quarks. It is natural to seek processes due to these transitions, above all in neutrino experiments. On the basis of the quark–parton picture, the differential distributions of the secondary muons with allowance for right-handed quark transitions will be

$$\left.\begin{aligned} &\frac{d^2\sigma^{\nu\mathcal{N}}}{dx\,dy} = \frac{G^2 ME}{\pi}\{[q_L(x) + \overline{q}_L(x)(1-y)^2] \\ &+ \varkappa[\overline{q}_R(x) + q_R(x)(1-y)^2]\}; \\ &\frac{d^2\sigma^{\tilde{\nu}\mathcal{N}}}{dx\,dy} = \frac{G^2 ME}{\pi}\{[q_L(x)(1-y)^2 + \overline{q}_L(x)] \\ &+ \varkappa[q_R(x) + \overline{q}_R(x)(1-y)^2]\}; \\ &x = \frac{Q^2}{2M\nu};\quad y = \frac{\nu}{E_\nu};\quad \varkappa = \mathrm{tg}^2\alpha\left(\frac{1-\eta}{1+\eta}\right)^2;\quad \eta = \frac{M^2_{W_1^\pm}}{M^2_{W_2^\pm}}.\ ^{5)} \end{aligned}\right\} \qquad (160)$$

In these expressions, the functions $q_L(x)$ and $\overline{q}_L(x)$ are the densities, multiplied by x, of the distributions of the primary quarks (or partons) and antiquarks in the nucleon, these quarks being coupled to the secondary quarks through the $V-A$ current, and $q_R(x)$ and $\overline{q}_R(x)$ are the corresponding quantities for the distributions of the quarks in the nucleon coupled to the other quarks through the $V+A$ current.

The variable x characterizes the fraction of the nucleon longitudinal momentum carried by the quark–parton that interacts with the lepton in the system in which the target nucleon moves with relativistic velocity. The parameter $0 \leq \varkappa \leq 1$ in (160) characterizes the intensity with which the right-handed currents are involved, and it depends on the particular model. For example, in the scheme with six or eight quarks based on the $SU(2)_L \times SU(2)_R \times U(1)$ gauge symmetry (see Sec. 3), the parameter $\varkappa$ depends on the degree of mixing of the charged $W_L^\pm$ and $W_R^\pm$ bosons.[71,72] The maximal increase in the cross section in the case of the $\nu_\mu \mathcal{N}$ interaction through the charged currents is by 1/3 due to right-handed transitions of the type $(n \leftrightarrow q)_R (Q_q = 2/3)$[110]; in the antineutrino interaction, the cross section can increase because of the transition $(p \leftrightarrow b)_R (Q_b = -1/3)$ by four times.[91] The existence of right-handed currents of the type $(p \leftrightarrow b)_R$ would lead to the appearance of a y anomaly[103] in antineutrino experiments and to a growth of the ratio of the cross section of the reaction $\overline{\nu}_\mu + \mathcal{N} \to \mu^+ + \ldots$ to the cross section of the process $\nu_\mu + \mathcal{N} \to \mu^- + \ldots$ with increasing primary energy. The new data[111–114] do not yet confirm the existence of the previously observed y anomaly, i.e., the behavior of $d\sigma(\overline{\nu}_\mu \mathcal{N})/dy$ inconsistent with the quark–parton model with three quark species. On the other hand, if mixing between the left- and right-handed worlds does not occur ($\varkappa = 0$), then the new currents will not have any influence on the behavior of the cross sections for the interaction of neutrinos and antineutrinos with nucleons, and heavy quarks (of the b and r type) will not be produced in neutrino experiments. In Harari's scheme, the left-handed currents can lead to the production of the r quark with its subsequent decay and production of the b quark (see Sec. 3). In the $\overline{\nu}_\mu \mathcal{N}$ interaction, such processes can be realized only on quarks in the sea of $q\overline{q}$ pairs with the production of $\overline{r}$ quarks and the decay of these into $\overline{b}$ quarks. In this case, searches for heavy quarks in neutrino experiments may be fruitless, but they will give a bound on $\varkappa$. The masses of the heavy quarks are expected to be in the range 2–100 GeV. Attempts to observe unusual hadrons should be attempted either in the process of e^+e^- annihilation or in hadron–hadron collisions. Also, as will be shown below, the observed polarization of the "direct" μ^+ in $P\mathcal{N}$ collisions is in practice very hard to explain in the framework of the standard GIM scheme. The only possibility remaining is to assume that already at energies $\sqrt{s} = 11.6$ GeV there occurs production of particles with the new quantum number "heaviness," the mass of these particles lying in the range 2–4 GeV. This question requires a careful experimental investigation.

By analogy with the lepton scheme that includes nondiagonal neutral transitions, it is helpful to consider a quark scheme based on either the gauge symmetry $SU(2) \times U(1)$ or $SU(2)_L \times SU(2)_R \times U(1)$ with four quark doublets:

4)The admixture of $\nu_e(\overline{\nu}_e)$ in $\nu_\mu(\overline{\nu}_\mu)$ beams is $\lesssim 0.5\%$. Therefore, searches for E^- and E^+ require a considerable time.
5) For $\alpha = \pi/4$, $\eta = M^2_{W_A}/M^2_{W_V}$.

$$\left.\begin{array}{l}\begin{pmatrix}\cos\alpha p+\sin\alpha\mathcal{P}\\ \cos\beta n+\sin\beta\mathfrak{N}\end{pmatrix}_{(\theta)_L};\quad \begin{pmatrix}\cos\alpha c+\sin\alpha C\\ \cos\beta\lambda+\sin\beta\Lambda\end{pmatrix}_{(\theta)_L};\\ \begin{pmatrix}\cos\tilde{\alpha}p+\sin\tilde{\alpha}\mathcal{P}\\ \cos\tilde{\beta}n+\sin\tilde{\beta}\mathfrak{N}\end{pmatrix}_{(\theta)_R};\quad \begin{pmatrix}\cos\tilde{\alpha}c+\sin\tilde{\alpha}C\\ \cos\tilde{\beta}\lambda+\sin\tilde{\beta}\Lambda\end{pmatrix}_{(\theta)_R}.\end{array}\right\} \quad (161)$$

In this scheme, the production of new particles with the heaviness quantum number can also occur through $V-A$ transitions.

Polarization Experiments and Right-Handed Leptonic Currents. In connection with the possible existence of right-handed leptonic currents discussed in Secs. 3 and 4,[94,115] it would be interesting to make an experimental search for leptons with right-handed polarization (and antileptons with left-handed polarization) in neutrino, muon, and electron interactions with nucleons. For example, one could look for the right-handed currents $(\mu^- M^0)_R$ or $(\bar{e}^- E^0)_R$ in the reactions $\mu^\pm + \mathcal{N} \to \overline{M}^0(M^0)$ + hadrons or $e^\pm + \mathcal{N} \to \overline{E}^0(E^0)$ + hadrons.[94,115] It then becomes very important to measure the polarization of positive muons in the decays

$$M^0 \Rightarrow \begin{cases}\mu^+ + \mu^- + \bar{\nu}_\mu\\ \mu^+ + e^- + \nu_e\\ \mu^+ + \text{hadrons}.\end{cases} \quad (162)$$

From the measured sign and magnitude of the μ^+ polarization one could establish the structure of the current $(\bar{\mu}^- M^0)$.

It is somewhat harder to establish the structure of the current $(\bar{\mu}^- M^0)$ in the decays

$$M^0 \Rightarrow \begin{cases}\mu^- + \mu^+ + \nu_\mu\\ \mu^- + e^+ + \nu_e\\ \mu^- + \text{hadrons}\end{cases} \quad (163)$$

because of the depolarization of the negative muons in the process of production of mesic atoms when μ^- mesons are stopped and captured in atomic levels and there are then cascade transitions of the atom to the ground state. However, even in this case the muon retains 7–15% of the original polarization, and to establish the nature of the interaction it is sufficient to measure the sign of the angular asymmetry of the electrons in the decay $\mu^- \to e^- + \bar{\nu}_e + \nu_\mu$. In addition, if the $(\bar{\mu}^- M^0)$ current has $V+A$ form, then in a beam of μ^- with left-handed polarization the production of M^0 will be impossible (and in a beam of μ^+ with right-handed polarization the production of $\overline{M}^0$ will be impossible). It is perfectly possible to produce $\mu^\pm$ beams with given polarization by separating different ranges of angles of emission of the muons from the decays of pions and kaons.[97] Changing the polarization of the muon beam, one could follow the change in the yield of events with the production of $\overline{M}^0$ and M^0.

One could look for neutrinos with right-handed polarization[90,115] (or antineutrinos with left-handed polarization), which can exist in the six- and eight-quark gauge schemes, by the beam dump method. In this method, one directs the proton beam directly onto a nuclear shield in order to reduce drastically the secondary flux of neutrinos with left-handed polarization (or antineutrinos with right-handed polarization) that result from the decays of the K and π mesons produced in proton–nuclear collisions. The relative fraction of unusual neutrinos (or antineutrinos) produced by the decay of the short-lived hypothetical hadrons (charmed bosons or fermions) must increase, so that the search for such processes is considerably facilitated. Neutrinos with right-handed polarization (or antineutrinos with left-handed polarization) will, when they reach a detector set up directly behind the nuclear shield, induce the reactions

$$\nu_{\mu R}(\nu_{eR}) + \mathcal{N} \to M^-(E^-) + \text{hadrons}; \quad (164)$$
$$\bar{\nu}_{\mu L}(\bar{\nu}_{eL}) + \mathcal{N} \to M^+(E^+) + \text{hadrons}. \quad (165)$$

The heavy leptons $M^\pm$ and $E^\pm$ produced in the reactions (164) and (165) can be identified by their leptonic decay modes:

$$M^- \to \mu^-(e^-) + \bar{\nu}_\mu(\bar{\nu}_e) + \nu_\mu;\quad E^- \to e^-(\mu^-) + \bar{\nu}_e(\bar{\nu}_\mu) + \nu_e; \quad (166)$$
$$M^+ \to \mu^+(e^+) + \nu_\mu(\nu_e) + \bar{\nu}_\mu;\quad E^+ \to e^+(\mu^+) + \nu_e(\nu_\mu) + \bar{\nu}_e. \quad (167)$$

The most probable value of the transverse momentum of the $\mu^\pm$ and $e^\pm$ products of the decays (166) and (167) will be $p_\perp \approx m_M/4$, which will make it possible to subtract the background from ordinary neutrino interactions and identify the processes (164) and (165). Right-handed polarization of the neutrinos in the processes (164) and (165) can be established from the energy and angular distributions of the charged leptons from the decays (166) and (167). Detailed calculations of these distributions are presented in Ref. 116.

CONCLUSIONS

Study of the interaction of elementary particles on the basis of the gauge approach has made it possible not only to eliminate a number of fundamental difficulties of the old Cabibbo theory (its nonrenormalizability, the difficulty associated with the unitarity limit, etc.), but also to predict new interesting phenomena in weak-interaction physics in a wide range of high and superhigh energies. The gauge approach first demonstrated the possibility of constructing a unified theory of the weak and electromagnetic interactions and indicated how the strong interactions of quarks could be included in this scheme. It appears very probable that the observed considerable difference between the effective coupling constants of these interactions is due to the fact that the experimental data at present at our disposal have been obtained at not very high energies. The transition to the region of superhigh energies (physics at distances comparable with a quantum of length 10^{-17} cm; see Ref. 117) must tell us whether we are dealing with a unified interaction or whether there exist in nature fundamentally different forms of interaction.

The main effects expected in the different gauge theories concern above all the production of new hadrons and leptons and the interaction quanta—the vector bosons—in different collisions of elementary particles (lepton–hadron collisions, e^+e^- annihilation, hadron–hadron interactions).

The search for and detailed study of the new particles and processes may lead in the not too distant future to the solution of a number of questions relating to the establishment of the lepton spectrum and, possibly, the elucidation of the nature of the analogy between quarks and leptons and in what their difference resides.

Let us now briefly list the topical problems for which the modern theory of elementary particles awaits an answer: 1) the μe problem and the possibly related λ/n problem (the origin of the Cabibbo angle); 2) the quark–lepton problem (the number of quarks and leptons and the nature of their analogy); 3) the problem of the search for quarks, gluons, and W bosons; 4) the type of gauge symmetries of the weak interactions and the number of guage bosons; 5) the energies at which the weak interactions become comparable with the strong interactions; 6) the problem of CP nonconservation in decays of the new hadrons.

The final choice of the unified gauge scheme of interactions of the elementary particles will largely depend on the solution of these problems.

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