to the same friction process leading to an excitation of ^{16}O close to an effective α threshold in the composite system. The mass-12 fragments are unlikely to originate from breakup, since this process would yield less γ rays compared to ^{16}O de-exciting by γ emission—in contradiction to the higher M_γ observed. On the other hand, *transfer* of an α particle in the course of the friction process will transform more angular momentum of the relative to the intrinsic motion, resulting in a higher M_γ—in agreement with our data. Since considerably higher excitation of ^{16}O will be needed to transfer one or two nucleons, the low yield of masses between 12 and 16 is consistent with this picture. The larger value of $|Q|$ for ^{12}C in comparison to ^{16}O is approximately ascribed to the difference in rotational and Coulomb energies. It may also be noted that semiclassical matching of Coulomb trajectories is incapable of explaining the large negative Q value of the friction-induced transfer.

(iii) In a classical description of friction, Tsang[4] distinguishes between sliding, rolling, and sticking of the fragments to each other, leading to a value of l_f/l_i being larger than, equal to, or smaller than $\frac{5}{7}$, respectively. With all the reservations in mind concerning the implicit assumption of rigid-body moments of inertia even for the light fragments, it is tempting to visualize a process leading to the formation of a rotating nuclear molecule, generated by friction. With $l_i = 42\hbar$ and $l_i - l_f = 12\hbar$ for ^{16}O we find $l_f/l_i = 0.7$.

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Does the Fourth ψ Particle Exist at About 4.9 GeV?*

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We show that, in a generalized three-triplet quark model, there can be four new fundamental vector-meson states with the photon quantum numbers. We propose that the broad peak in the total cross section for e^+-e^- annihilation into hadrons at 4.1 GeV and the narrow resonances $\psi(3.1)$ and $\psi(3.7)$ account for three of these. We predict, among other things, that there exists a fourth resonance ($\Gamma \sim 300$ MeV) at about 4.9 GeV with a partial width to electrons of approximately 1 keV.

The Stanford Linear Accelerator Center–Lawrence Berkeley Laboratory e^+-e^- colliding-beam experiment[1] at SPEAR has recently revealed not only very narrow resonances $\psi(3.1)$ and $\psi(3.7)$ but also a broad enhancement of the total cross section at the center-of-mass energy $\sqrt{s} \approx 4.1$ GeV. If it is a single vector-meson resonance [which we call $\psi(4.1)$], its total width (Γ) and partial decay width to an electron pair ($\Gamma_{\psi(4.1)\to e^+e^-}$) are reported to be approximately 250–300 MeV and 4 keV, respectively. In this paper we propose that $\psi(3.1)$, $\psi(3.7)$, $\psi(4.1)$, and a yet unobserved resonance $\psi(4.9)$ make up a set of four new states with the photon quantum numbers allowed by a generalized three-triplet quark model (TTQM). (We always exclude daughters and excited states.) We introduce a single working hypothesis which reflects our ignorance about the dynamics of color-nonsinglet states. It predicts and explains various properties of $\psi(3.1)$, $\psi(3.7)$, $\psi(4.1)$, and $\psi(4.9)$. The existence of other members of the multiplet and some of their properties are also discussed.

We start with n SU(3) triplets of quarks[2,3] $q_{i,j}$ ($i=\mathcal{P}, \mathfrak{N}, \lambda$, and $j=1,2,3,\ldots,n$) with arbitrary charges $Q_{i,j}$:

$$\begin{pmatrix} Q_{\mathcal{P},j} \\ Q_{\mathfrak{N},j} \\ Q_{\lambda,j} \end{pmatrix} = \begin{pmatrix} \frac{2}{3}+\alpha_j \\ -\frac{1}{3}+\alpha_j \\ -\frac{1}{3}+\alpha_j \end{pmatrix}. \tag{1}$$

The approximate symmetry obeyed by these quarks is SU(3)$\otimes G$. The color symmetry group G is not yet specified. We assign the ordinary hadrons to G-singlet states. A mixing between G-singlet and -nonsinglet states is possible but assumed to be extremely small. This is consistent with the apparent large mass splitting between these states. The condition[2]

$$\sum_{j=1}^{n} \alpha_j = 0 \tag{2}$$

guarantees that the usual hadron spectroscopy is preserved among color-singlet states. In this scheme, the partial conservation of axial-vector current anomalous constant S for the π^0 decay into two photons becomes independent of α_j: $S=n/6$. Since the experimental value for S is very close to $\frac{1}{2}$, we must choose $n=3$.[2,3]

This model is nothing but a simple generalization of already existing TTQM's. In fact, results of other models can be recovered by specific choices of $\{G;\alpha_j\}$; for example, in the Fritzsch-Gell-Mann model,[3] $\{\mathrm{SU}(3)';\ \alpha_j=0\}$; in the Han-Nambu model,[2] $\{\mathrm{SU}(3)'';\ \alpha_j=\frac{1}{3},\frac{1}{3},-\frac{2}{3}\}$; and in the Tati model,[4] $\{\mathrm{SO}(3);\ \alpha_j=1,0,-1\}$.

In general, the electromagnetic current J_μ transforms as follows:

$$J_\mu \sim \sum_{j=1,2,3} \left(\tfrac{2}{3}\bar{q}_{\mathcal{P},j}q_{\mathcal{P},j} - \tfrac{1}{3}\bar{q}_{\mathfrak{N},j}q_{\mathfrak{N},j} - \tfrac{1}{3}\bar{q}_{\lambda,j}q_{\lambda,j}\right) + \sum_{i=\mathcal{P},\mathfrak{N},\lambda} \left(\alpha_1\bar{q}_{i,1}q_{i,1} + \alpha_2\bar{q}_{i,2}q_{i,2} + \alpha_3\bar{q}_{i,3}q_{i,3}\right). \tag{3}$$

With use of the naive quark-parton model, the ratio $R=\sigma(e^+e^- \to \text{hadrons})/\sigma(e^+e^- \to \mu^+\mu^-)$ is expected to increase from $R=2$ in the quasiasymptotic region, where only color-singlet hadrons are produced, to

$$R = 2 + 3\sum_{j=1}^{3}\alpha_j^2 \tag{4}$$

in the real asymptotic region, where all states are produced.

How many vector mesons with the photon quantum numbers are accommodated by a TTQM? The first term in (3) transforms as a linear combination of the ordinary ρ^0, ω, and φ states. In discussing the maximum number of states accommodated by the second term in (3), it is convenient to consider the following three steps: When SU(3) and G are both exact, the number is one. When SU(3) is exact but G broken, this number becomes two. (Note that a possible G-singlet combination, by definition, can be written in terms of a linear combination of the ordinary vector mesons.) Finally, when SU(3) and G are both broken, this number grows to four. This doubling of states can be caused by, for example, ω-φ mixing.[5] In order to obtain more than two nondegenerate new vector-meson states, we must assume that the color symmetry G as well as the ordinary SU(3) is broken. Since the rank of SO(3) is one, the maximum number of states obtained for this choice of G is two. We thus assume an SU(3) for the color symmetry: $G=\mathrm{SU}(3)_c$. We further assume, in close analogy with the ordinary SU(3) breaking, that SU(3)$_c$ is broken to SU(2)$_c$ in such a way that both the color isospin (I_c) and the color hypercharge (Y_c) are conserved. At this stage, we have not yet introduced any color-symmetry breaking which mixes color-singlet states with color-nonsinglet ones. An extremely small mixing between these states as well as that between color-nonsinglet states with different I_c is, however, necessary for explaining some decays of ψ's. As will be seen below, this symmetry-breaking scheme will be incorporated in our hypothesis.

Let (a, b) denote a state which belongs to a certain irreducible representation of $SU(3)\otimes SU(3)_c$. Here a and b specify members belonging to certain irreducible representations of SU(3) and $SU(3)_c$, respectively. For example, the ρ^0 meson is represented by $(\rho^0, 1)$ where 1 stands for a singlet representation. We shall introduce the following hypothesis on the dynamics of color-nonsinglet states:

Hypothesis.—A transition matrix element between two states with color quantum numbers satisfies

$$\langle (a,b)|T|(a',b')\rangle \cong C_{bb'}(T)\langle (a,1)|T|(a',1)\rangle$$

with

$$C_{bb}(T)=O(1) \text{ and } |C_{bb'}(T)| \ll 1 \text{ for } b\neq b'.$$

This hypothesis may follow from an interesting possibility that nature chooses to decouple the dynamics in the ordinary SU(3) space and the dynamics in the color space and furthermore, to conserve color quantum numbers to a very good accuracy.

An immediate consequence of the hypothesis is, when T is taken to be the Hamiltonian, equality of the ω-φ mixing angles for $b=1$ and for $b\neq 1$. With an arbitrary constant β for the excess charges $\alpha_j = \beta(\frac{2}{3}, -\frac{1}{3}, -\frac{1}{3})$, the decomposition of the electromagnetic current can be completed as

$$J_\mu \sim \sqrt{\tfrac{3}{2}}[(\rho^0,1)+\tfrac{1}{3}(\omega,1)-\tfrac{1}{3}\sqrt{2}(\varphi,1)]+\beta[(\omega,\rho^0)+\tfrac{1}{3}\sqrt{3}(\omega,\omega_8)+\tfrac{1}{2}\sqrt{2}(\varphi,\rho^0)+\tfrac{1}{6}\sqrt{6}(\varphi,\omega_8)]. \quad (5)$$

The ratio of photon couplings for the color-nonsinglet vector mesons can now be predicted:

$$g^2(\omega,\rho^0):g^2(\omega,\omega_8):g^2(\varphi,\rho^0):g^2(\varphi,\omega_8)=1:\tfrac{1}{3}:\tfrac{1}{2}:\tfrac{1}{6}. \quad (6)$$

Let us now proceed to the most crucial point in this paper, namely, the assignment of $\psi(3.1)$, $\psi(3.7)$, and $\psi(4.1)$ to various states given in (5). To this end, note that the hypothesis predicts relationships among the masses:

$$m(\varphi,\rho^0)/m(\omega,\rho^0)\cong m(\varphi,\omega_8)/m(\omega,\omega_8)\cong m(\varphi,1)/m(\omega,1)=1.30. \quad (7)$$

Since $m_{\psi(4.1)}/m_{\psi(3.1)}=1.32$, it is tempting to assign the pair $[\psi(3.1), \psi(4.1)]$ to either $[(\omega,\rho^0), (\varphi,\rho^0)]$ or to $[(\omega,\omega_8), (\varphi,\omega_8)]$. In order to decide between these choices, we use the experimental fact that the partial decay width of $\psi(3.1)$ into an electron pair (~ 5.2 keV)[6] is definitely larger than that of $\psi(3.7)$ (~ 2 keV), as well as the prediction (6) that $g^2(\omega,\rho^0), g^2(\varphi,\rho^0) > g^2(\omega,\omega_8)$. Therefore, we shall *take $\psi(3.1)$ and $\psi(4.1)$ as (ω,ρ^0) and (φ,ρ^0), respectively*. Next, there are two choices for the assignment of $\psi(3.7)$. The partial decay width of $\psi(3.7)$ into an electron pair favors the choice of *identifying $\psi(3.7)$ with (ω,ω_8)*. In fact, rough estimates of the photon-vector-meson coupling constants [defined by $\langle 0|J_\mu(0)|\psi\rangle=(m_\psi{}^2/f_\psi)\epsilon_\mu$, where ϵ_μ is the polarization vector of ψ] from the experimental data[1] satisfy the relation (6) remarkably well:

$$(f_{\psi(3.1)}{}^2/4\pi):(f_{\psi(3.7)}{}^2/4\pi):(f_{\psi(4.1)}{}^2/4\pi)\cong(10\pm1):\sim27:\sim18\cong1:3:2. \quad (8)$$

What has not yet been assigned is (φ,ω_8). One of the most unambiguous predictions in this model is the following: *There exists a fourth new vector-meson resonance (φ,ω_8), which we call $\psi(4.9)$, with the photon quantum numbers. The mass is*

$$m(\varphi,\omega_8)\cong(m_\varphi/m_\omega)m_{\psi(3.7)}\cong 4.9 \text{ GeV}, \quad (9)$$

and the partial decay width to an electron pair is

$$\Gamma_{\psi(4.9)\to e^+e^-}\cong\tfrac{1}{6}(m_{\psi(4.9)}/m_{\psi(3.1)})\Gamma_{\psi(3.1)\to e^+e^-} \cong 1 \text{ keV}. \quad (10)$$

We summarize the assignment of the new vector-meson resonances:

$$\psi(3.1)=(\omega,\rho^0),\quad \psi(3.7)=(\omega,\omega_8)$$
$$\psi(4.1)=(\varphi,\rho^0),\quad \psi(4.9)=(\varphi,\omega_8). \quad (11)$$

Next we shall briefly discuss possible decays of ψ's. We find that the decays are completely dictated by this hypothesis: (a) The decays $\psi(4.1)\to\psi(3.1)$ + ordinary hadrons and $\psi(4.9)\to\psi(3.7)$ + ordinary hadrons are allowed by strong interactions since $C_{\rho^0,\rho^0}, C_{\omega_8,\omega_8}=O(1)$. Both $\psi(4.1)$ and $\psi(4.9)$ are therefore predicted to be broad resonances. Experimentally $\psi(4.1)$ has been found to be a possible broad resonance. *We ex-*

pect $\psi(4.9)$ *to be as broad as* $\psi(4.1)$ ($\Gamma \sim 300$ *MeV*). (b) The decays $\psi(3.1) \to$ ordinary hadrons, $\psi(3.7) \to$ ordinary hadrons, and $\psi(3.7) \to \psi(3.1)$ + ordinary hadrons[7] are all suppressed because they are not transitions between states with the same color quantum numbers ($|C_{1,\rho^0}|$, $|C_{1,\omega_8}|$, $|C_{\rho^0,\omega_8}| \ll 1$).[8] (c) The hypothesis implies that *the selection rules for transitions between color-singlet states, e.g.,* $\Delta \vec{I} = \Delta Y = 0$ *and G-parity conservation, are strictly preserved for those between color-nonsinglet states.* In particular, among the suppressed decays in (b), $\psi(3.1) \to$ even number of pions is further suppressed because of G-parity nonconservation. On the other hand, for the same reason $\psi(3.7) \to \psi(3.1) + 2\pi$ is preferred to $\psi(3.7) \to \psi(3.1) + \pi^0$. (d) Radiative decays of ψ's can be discussed by taking T to be J_μ. The decays $\psi \to \gamma$ + ordinary hadrons [e.g., $\psi(3.1) \to \gamma + \eta$] are suppressed compared to similar decays of the corresponding ordinary vector meson (e.g., $\varphi \to \gamma + \eta$) since $|C_{1,\rho^0}|$, $|C_{1,\omega_8}| \ll 1$. Also, the radiative decay of any one of the ψ's into $\gamma + \pi^0$ is forbidden by isospin conservation to the lowest order in electromagnetic interactions. (e) The matrix elements for $\psi(4.1) \to \psi(3.1) + K + \bar{K}$ and $\psi(4.9) \to \psi(3.7) + K + \bar{K}$ are comparable to that for $\varphi \to K + \bar{K}$. However, because of the small three-body phase space, the former two decays are even more strongly suppressed than the latter.

As another consequence of the hypothesis, we can predict at least fourteen more vector-meson states with color quantum numbers. They are an isotriplet of (ρ, ρ^0) at 3.1 GeV, isodoublets of (K^*, ρ^0) and $(\bar{K}^*, \rho^0)$ at 3.5 GeV, an isotriplet of (ρ, ω_8) at 3.6 GeV, and isodoublets of (K^*, ω_8) and $(\bar{K}^*, \omega_8)$ at 4.2 GeV. According to our hypothesis, the decays $\psi(3.7) \to (\rho, \rho^0) + \pi$ and $\psi(3.7) \to \psi(3.1) + 2\pi$ are both suppressed.[7] Since, however, the latter decay has been seen as one of the major decay modes, we should expect the former to be seen as well.

We conclude our discussion of ψ decays by pointing out that our model raises the following major problem: It is widely known that naive estimates of radiative decay widths for $\psi(3.1)$ and $\psi(3.7)$, based on symmetries, are comparable to or even larger than their total widths observed by experiments. However, we note that by symmetry arguments alone radiative decays of ψ's cannot rigorously be related to the corresponding decays of ordinary vector mesons even in the exact-symmetry limit. It is plausible that some dynamical difference between binding mechanisms of a color-singlet state and of a color-nonsinglet state may invalidate the naive estimates. In our picture, suppressions of these radiative decays are consequences of our hypothesis, whose dynamical justification remains to be understood.

It is amusing to see that the mass ratio of $\psi(3.7)$ to $\psi(3.1)$ ($\cong 1.19$) is very close to that of ω_8 to ρ^0 ($\cong 1.21$). This coincidence tempts us to speculate that the symmetry breaking in the color SU(3) space may be similar to that in the ordinary SU(3) space. If this is the case, assuming the Gell-Mann–Okubo mass formula in the color space, we expect that there exist four additional sets of nonets, i.e., (ρ, K^*), (ω, K^*), (K^*, K^*), and (φ, K^*) at 3.5, 3.5, 4.0, and 4.7 GeV, respectively.

In conclusion, we stress the importance of searching for $\psi(4.9)$ in e^+-e^- colliding-beam experiments. Because of its predicted large total width (~300 MeV) and small partial width to an electron pair (~1 keV), a careful study of the total cross section for hadron production in the vicinity of $\sqrt{s} = 4.9$ GeV is highly desirable. The area under $\psi(4.9)$ is expected to be approximately one third of the area under $\psi(4.1)$ when integrated over $\sqrt{s}$. To check this, the background under each of these resonances must be carefully subtracted. Also, for searching for the other color-nonsinglet vector-meson states predicted in this paper, neutrino-production as well as hadron-hadron collisions at Brookhaven National Laboratory, CERN, and Fermi National Accelerator Laboratory is very promising.

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[8]$|C_{1,\rho^0}|$, $|C_{1,\omega_8}|\ll 1$, and $|C_{1,\rho^0}|$, $|C_{\rho^0,\omega_8}|\ll 1$ because of the smallness of color-singlet-octet mixing and of color-isospin violation, respectively. Because of these simultaneous suppressions of C_{1,ρ^0}, it may be that $|C_{1,\rho^0}|\ll|C_{1,\omega_8}|$, $|C_{\rho^0,\omega_8}|$. This has been pointed out to the authors by V. Rittenberg.

SU(4) Symmetry and Nonleptonic Decays*

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I derive the new relation $\Lambda_-^0:\Sigma_0^+:\Xi_-^-=1:-\sqrt{3}:-2$ for the S-wave amplitudes of the hyperon nonleptonic decays, assuming SU(4) 20-plet dominance for the weak interaction of current-current type. I also discuss the nonleptonic decays of charmed mesons.

The purpose of this Letter is to discuss the nonleptonic decays of the hyperons and charmed mesons, by assuming the validity of the mechanism proposed by Glashow, Iliopoulos, and Maiani[1] (GIM mechanism), which can suppress the strangeness-changing neutral current in semileptonic decays. The weak current is of the form[1,2]

$$J_\mu=\bar{u}\gamma_\mu(1+\gamma_5)(\cos\theta d+\sin\theta s)+\bar{c}\gamma_\mu(1+\gamma_5)(\cos\theta s-\sin\theta d),\qquad(1)$$

where θ is the Cabibbo angle and u, d, s, and c represent quark fields of fractional charges. The strong interaction possesses approximate SU(4) symmetry. I assign meson multiplets to $\underline{15}\oplus\underline{1}$ and the baryon multiplet $(\frac{1}{2}^+)$ to $\underline{20}$,[3] as usual.

Nonleptonic decays occur through the current-current interactions which may be mediated by weak bosons. The bilinear form contains two parts[4] belonging to a 20-dimensional[9] and an 84-dimensional SU(4) representations; denote them by traceless tensors $T_{[cd]}{}^{[ab]}$ and $T_{\{cd\}}{}^{\{ab\}}$. A generalization of the SU(3) octet dominance is the SU(4) 20-plet dominance,[4] since the SU(3) decompositions are as follows: $\underline{20}\supset\underline{8}$ and $\underline{84}\supset\underline{27}\oplus\underline{8}\oplus\underline{1}$ for $\Delta C=0$. Thus I assume the 20-plet dominance throughout this Letter.

Now it is straightforward to discuss the nonleptonic decays of the baryons. We can easily derive[10]

$$\Lambda_-^0:\Sigma_0^+:\Xi_-^-=1:-\sqrt{3}:-2,\qquad(2)$$

for the S-wave amplitudes. [I only consider CP-conserving interactions.] The derivation is analogous to that of the Lee-Sugawara relation for the S-wave amplitudes in the SU(3) symmetry. I shall sketch the derivation. The parity-nonconserving part of the matrix element of the nonleptonic decay has the following form[11]:

$$\begin{aligned}&a(T_{[kl]}{}^{[ij]}\bar{B}_{[nm]}{}^kB_i{}^{[nm]}M_j{}^l-T_{[ij]}{}^{[kl]}\bar{B}_{[nl]}{}^iB_k{}^{[nm]}M_m{}^j)\\&+b(T_{[kl]}{}^{[ij]}\bar{B}_{[im]}{}^nB_j{}^{[km]}M_n{}^l-T_{[ij]}{}^{[kl]}\bar{B}_{[km]}{}^jB_n{}^{[im]}M_l{}^n)\\&+c(T_{[kl]}{}^{[ij]}\bar{B}_{[ij]}{}^lB_m{}^{[kn]}M_n{}^m-T_{[ij]}{}^{[kl]}\bar{B}_{[kn]}{}^mB_l{}^{[ij]}M_m{}^n)\\&+d(T_{[kl]}{}^{[ij]}\bar{B}_{[ij]}{}^mB_m{}^{[kn]}M_n{}^l-T_{[ij]}{}^{[kl]}\bar{B}_{[kn]}{}^mB_m{}^{[ij]}M_l{}^n),\end{aligned}\qquad(3)$$