paper it is noted that the strong W-pair model leads to an effective cutoff Λ on higher-order weak interactions of order $\Lambda \sim \mathfrak{M}(W)$. If the J particles are indeed W's, the value $\Lambda \simeq 4$ GeV deduced from the $(K_L{}^0 - K_S{}^0)$ mass difference [R. N. Mohapatra, J. Subba Rao, and R. E. Marshak, Phys. Rev. 171, 1502 (1968)] would find a natural explanation. See also R. E. Marshak, Riazuddin, and C. P. Ryan, *Theory of Weak Interactions in Particle Physics* (Wiley, New York, 1969), Chap. IX.

[9]S. V. Pepper, C. Ryan, S. Okubo, and R. E. Marshak, Phys. Rev. 137, B1259 (1965).

[10]S. Okubo, V. S. Mathur, and J. E. Kim, University of Rochester Report No. UR-506 (unpublished).

[11]Cf. A. Rousset, in *Neutrinos—1974*, AIP Conference Proceedings No. 22, edited by C. Baltay (American Institute of Physics, New York, 1974).

[12]On the contrary, these small decay rates pose a serious problem for the unified gauge theories and require the postulation of a charmed quark [cf. B. W. Lee, in *Proceedings of the Sixteenth International Conference on High Energy Physics, The University of Chicago and National Accelerator Laboratory, 1972* edited by J. D. Jackson and A. Roberts (National Accelerator Laboratory, Batavia, Ill., 1973).

[13]There are other diagrams wherein the W's and h's are essentially interchanged.

[14]It should be noted that there are two final states, $W^{0\prime} K_2{}^0$ and $W^{0\prime} K_1{}^0$, and that only $W_1{}^0$, (with $CP = +1$) can be identified with the observed $J(3108)$ particle.

[15]This prediction depends on t conservation in the strong decays $W^0 \to W^0 + K^0$ and $\overline{W}^0 \to \overline{W}^{0\prime} + \overline{K}^0$ and the fact that only one third of the decays of K^0 and $\overline{K}^0$ yield $\pi^+\pi^-$.

[16]R. Holleebeek (private communication) gave the range 125 keV to 1 MeV for the decay width of $J(3695)$ and a branching ratio into $\pi^+\pi^-$ of the order of 30%.

[17]Even allowing for the fact that W possesses spin —implying a necessarily stronger interaction with the electromagnetic field than for a spin 0 or $\frac{1}{2}$ particle —one could hardly expect $\mathfrak{M}(W^+) - \mathfrak{M}(W^0)$ [W^+, W^0 in the same isodoublet] to exceed several hundred MeV.

[18]Cf. F. J. Sciulli, in *Neutrinos—1974*, AIP Conference Proceedings No. 22, edited by C. Baltay (American Institute of Physics, New York, 1974).

[19]Cf. I. T. Drummond, Nucl. Phys. B35, 269 (1971); C. Rebbi, Lett. Nuovo Cimento 1, 967 (1971).

Quark-Model Comparison on the Narrow e^+e^- Resonances

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The narrow e^+e^- resonances $\psi(3.1\text{ GeV}/c^2)$ and $\psi(3.7\text{ GeV}/c^2)$ are used to test quark models: $SU(3)\otimes SU(3)'$ encompasses them readily but also predicts four more such states; SU(4) includes more than one state only by extension to $(q\overline{q}\,q\overline{q})$ representations; the two-triplet model encounters severe difficulties; and scalar colored quarks seem practically excluded. The importance of photoproduction, especially of charged counterparts, is emphasized.

Recent measurements[1,2] have indicated the presence of two sharp resonances in the e^+e^- channel, at 3.1 and 3.7 GeV/c^2. These resonances are most likely very heavy vector mesons, highly stable against strong decay with an inhibition factor of order 10^3 to 10^4.

Such a large factor suggests a new quantum number for hadronic matter, and we ascribe it to quarks.[3] Several schemes already exist in the literature for the extension of quark parameters beyond simple SU(3); the object of the present note is to compare the most popular schemes in their application to the new discoveries.[1,2] It will turn out that their implications are sufficiently different to allow ultimate experimental distinction.

(I) SU(3) ⊗ SU(3)′, the Han-Nambu model.[4] —The nine fundamental objects in this three-triplet model are denoted by $\mathcal{P}_i$, $\mathfrak{N}_i$, and λ_i, where $i = 1, 2, 3$ is the SU(3)′ index and $\mathcal{P}$, $\mathfrak{N}$, and λ refer to the SU(3) index of conventional strong interactions. Their electric charges are given by $Q(\mathcal{P}_2) = Q(\mathcal{P}_3) = +1$, $Q(\mathfrak{N}_1) = Q(\lambda_1) = -1$, all other $Q = 0$; the electromagnetic current operator is[5]

$$\begin{aligned} J_\mu &= \overline{\mathcal{P}}_2\gamma_\mu\mathcal{P}_2 + \overline{\mathcal{P}}_3\gamma_\mu\mathcal{P}_3 - \overline{\mathfrak{N}}_1\gamma_\mu\mathfrak{N}_1 - \overline{\lambda}_1\gamma_\mu\lambda_1 \\ &= \tfrac{1}{2}J_\mu(3,0) + (2\sqrt{3})^{-1}J_\mu(8,0) \\ &\qquad - \tfrac{1}{2}J_\mu(0,3) - (2\sqrt{3})^{-1}J_\mu(0,8). \end{aligned} \tag{1}$$

Since Eq. (1) is symmetric in SU(3) and SU(3)′, the electromagnetic current will couple directly to $\psi(0,3)$ and $\psi(0,8)$ just as to $\psi(3,0)$ and $\psi(8,0)$. Recalling again the medium strong interactions $H(8,0)$, we see that the electromagnetic current will hence generate $\psi(8,3)$ and $\psi(8,8)$, or, properly orthogonalized,[6] $\psi(\omega,3)$, $\psi(\varphi,3)$, $\psi(\omega,8)$, and $\psi(\varphi,8)$. One thus expects the process $e^+ + e^- \rightarrow (\gamma) \rightarrow \psi$ to involve four neutral vector mesons that are octet[7] with respect to SU(3)′.

Because of the wide mass separation of SU(3)′ octet and singlet (ordinary) states, we assume that any $H(0,8)$ interaction is negligible. The lowest SU(3)′ octet states would thus be stable against strong decay by emission of ordinary hadrons; the higher SU(3)′ octet states can decay strongly into the lowest[8] ones, however, if the mass difference is sufficient (greater than about 1–3 pion masses). Not all of the four neutral mesons listed[6] will decay strongly into each other. In particular, conservation of I' spin for strong interactions ensures that $\psi(i,3) \leftrightarrow (i,8)$ is allowed only electromagnetically. This suggests that if we assign one of these to 3.1 GeV/c^2, the other should be assigned to 3.7 GeV/c^2. Suppose that $\psi(3.1) = \psi(\omega,3)$ and $\psi(3.7) = \psi(\omega,8)$. What can we say about the other two? If $\Delta m = |m(\varphi,j) - m(\omega,j)| \gtrsim 1$ GeV/c^2, the transition between $\psi(\varphi,j)$ and $\psi(\omega,j)$ will be rapid, with the emission of a $K\bar{K}$ pair. This means that $\psi(\varphi,3)$ and $\psi(\varphi,8)$ will hardly be observable by e^+e^- if their masses much exceed 4.1 and 4.7 GeV/c^2, respectively; and that their masses must lie above 2.1 and 2.7 GeV/c^2, or the observed states would not be so narrow. The same arguments apply with less force for $\Delta m \gtrsim 0.6$ GeV/c^2, where the emitted particle is an η meson. The ultimate limit is $\Delta m \gtrsim 2m_\pi \simeq 0.3$ GeV/c^2, where the emission is greatly inhibited by the usual λ-quark argument for eliminating $\varphi(1020) \rightarrow n\pi$.

One other possibility is that both $\psi(3.1)$ and $\psi(3.7)$ have the same SU(3)′ index (3 or 8) and are therefore related as an ordinary φ and ω. By arguments like the above, both should be observable in the e^+e^- channel; but the $\psi(3.7)$ should be broadened by a component of $\psi(3.7) \rightarrow \psi(3.1)$ + (ordinary mesons). This question should be accessible to experimental study.

There are two more neutral vector nonstrange mesons in the same mass region, $\psi(3,3)$ and $\psi(3,8)$, which only couple to the electromagnetic current in second order. In electron-positron collisions they would be produced with sharp low-energy but broad upper-energy edge, as $e^+ + e^- \rightarrow \psi(3,3) + \{\gamma \text{ or } \pi^0\}$. Their decay will be by electromagnetic transition to ordinary mesons; e.g., $\psi(3,8)$ or $\psi(3,3) \rightarrow \pi^0 + \gamma$.

These considerations suggest a classification of the 81 SU(3)⊗SU(3)′ mesons into three groups: nine "ordinary" $\psi(i,0)$, eighteen "semicharmed" $\psi(i,3)$ and $\psi(i,8)$; and fifty-four "charmed" $\psi(i,j)$ with $j = 1, 2, 4, 5, 6, 7$. The charmed mesons can decay only through weak interactions: $\psi(i,j) \rightarrow \psi(k,0) + l + \bar{\nu}$. The mean lifetime for decay can be variously estimated, but is probably $\gtrsim 10^{-15}$ sec. The semicharmed ones are capable of electromagnetic decay into ordinary states, including the vacuum, and have lifetimes generally on the order of 10^{-20} sec. The semicharmed bosons could be photoproduced on ordinary targets and must include some charged counterparts to the neutral $\psi(3.1)$ and $\psi(3.7)$; the establishment of these charged counterparts would provide an important experimental test of the model.

(II) SU(4), the original "charm".[9]—It is simplest to list the individual SU(4) quarks as $\mathcal{P}$, $\mathcal{N}$, λ, and $\mathcal{P}'$; or to represent the basic tetrahedron of SU(4) as $[\underline{4}] = (\underline{3}) + \mathcal{P}'$, where $(\underline{3})$ is the SU(3) triangle. The charmed quark $\mathcal{P}'$ decays into the others only through weak interactions. The quark charges are now $Q(\mathcal{P}) = Q(\mathcal{P}') = \frac{2}{3}$, $Q(\lambda) = Q(\mathcal{N}) = -\frac{1}{3}$.

The immediate representations for mesons $(q\bar{q})$ are given by

$$[\underline{4}] \otimes [\underline{4}^*] = [\underline{1}] + [\underline{15}] = (\underline{8}) + (\underline{1}) + \mathcal{P}'(\underline{3}^*) + (\underline{3})\bar{\mathcal{P}}' + \mathcal{P}'\bar{\mathcal{P}}'. \qquad (2)$$

If the $\mathcal{P}'$ mass is much greater than that of the $(\underline{3})$ quarks, the representations in Eq. (2) will be split in such a way as to make the $\mathcal{P}'\bar{\mathcal{P}}'$ an essentially pure component, as with the ordinary $\varphi = \lambda\bar{\lambda}$. Then a narrow vector meson that decays only electromagnetically would be $\psi(\mathcal{P}'\bar{\mathcal{P}}')$. However, this representation provides only a single meson where at least two have been observed. The $\mathcal{P}'(\underline{3}^*)$ and $(\underline{3})\bar{\mathcal{P}}'$ each contain one neutral, nonstrange meson; but they have charm quantum number $C = \pm 1$, and by analogy with the K^0 and $\bar{K}^0$ are subject only to weak and not to first-order electromagnetic decay.

It is necessary to seek higher (self-conjugate) representations. The next lowest to $[\underline{15}]$ is $[\underline{20}]_{sc}$, which can be obtained as follows:

$$\begin{aligned} &[\underline{4}] \otimes [\underline{4}] = [\underline{6}] + [\underline{10}], \\ &[\underline{6}] \otimes [\underline{6}^*] = [\underline{1}] + [\underline{15}] + [\underline{20}], \qquad (3) \\ &[\underline{20}]_{sc} = (\underline{6})\mathcal{P}' + (\underline{6}^*)\mathcal{P}' + (\underline{8})\mathcal{P}'\bar{\mathcal{P}}'. \end{aligned}$$

All of the representations in Eq. (3) involve $(q\bar{q}$-

$q\bar{q}$); however, in the presumably massive quarks $\mathcal{P}'$ and $\bar{\mathcal{P}}'$ they are of the same order as $\psi(\mathcal{P}'\bar{\mathcal{P}}')$ and may therefore be found in the same mass range. Equation (3) contains two more neutral, noncharmed mesons in the $(\underset{\sim}{8})\mathcal{P}'\bar{\mathcal{P}}'$ part: $\psi(8, \mathcal{P}'\bar{\mathcal{P}}')$, $\psi(3, \mathcal{P}'\bar{\mathcal{P}}')$, using a notation $\psi(i, \mathcal{P}'\bar{\mathcal{P}}')$ similar to that for section I above, except that the second index does not run $j = 0$ to 8 but only $j = 0, 1, \cdots$ for the number of $\mathcal{P}'\bar{\mathcal{P}}'$ pairs. The previous $\psi(\mathcal{P}'\bar{\mathcal{P}}')$ is in this notation $\psi(0, \mathcal{P}'\bar{\mathcal{P}}')$. As before, the presence of $H(8, 0)$ induces mixing of $\psi(0, \mathcal{P}'\bar{\mathcal{P}}')$ and $\psi(8, \mathcal{P}'\bar{\mathcal{P}}')$ to the observed mesons $\psi(\omega, \mathcal{P}'\bar{\mathcal{P}}')$ and $\psi(\varphi, \mathcal{P}'\bar{\mathcal{P}}')$. These are then candidates for $\psi(3.1)$ and $\psi(3.7)$ (with the same remarks about strong decay between them for $\Delta m \gtrsim 0.6$–1.0 GeV/c^2). As before, the $\psi(3, \mathcal{P}'\bar{\mathcal{P}}')$ couples to the electromagnetic field only in second order; it will be unstable against decay by π^0 emission into $\psi(0, \mathcal{P}'\bar{\mathcal{P}}')$ unless their mass separation is very small.

It thus appears that SU(4) is also adequate to fit the observations to date, though it cannot extend to very many more neutral vector mesons like $\psi(3.1)$ and $\psi(3.7)$. The device of adding higher representations seems to carry no further; the self-conjugate representations of $(q\bar{q}\ q\bar{q}\ q\bar{q})$ repeat those found above or add a few with $(\mathcal{P}'\bar{\mathcal{P}}'\ \mathcal{P}'\bar{\mathcal{P}}')$, which will presumably be in another mass range. So far as narrow neutral vector mesons go, the SU(4) and SU(3) $\otimes$ SU(3)$'$ schemes predict the same variations, but with a factor of 2 in the multiplicity:

$$\psi(\omega, \mathcal{P}'\bar{\mathcal{P}}'),\quad \psi(\varphi, \mathcal{P}'\bar{\mathcal{P}}'),\quad \psi(\rho^0, \mathcal{P}'\bar{\mathcal{P}}') \quad \text{for SU(4)},$$

$$\left\{\begin{matrix} \psi(\omega, 8), & \psi(\varphi, 8), & \psi(\rho^0, 8) \\ \psi(\omega, 3), & \psi(\varphi, 3), & \psi(\rho^0, 3) \end{matrix}\right\} \quad \text{for SU(3)} \otimes \text{SU(3)}'. \tag{4}$$

Again the SU(4) scheme predicts additional semicharmed mesons, which can be photoproduced from ordinary targets: $\psi(\rho^{\pm}, \mathcal{P}'\bar{\mathcal{P}}')$ and $\psi(K, \mathcal{P}'\bar{\mathcal{P}}')$. The charged counterparts among these are reduced in multiplicity by a factor of 3 from SU(3) $\otimes$ SU(3)$'$, being four instead of twelve. The distinction of these multiplicities provides an additional reason for experimental study of photoproduction. In this SU(4) scheme the eighteen fully charmed mesons that decay only by weak interactions are $\psi((\underset{\sim}{6})\mathcal{P}')$, $\psi((\underset{\sim}{3}^*)\mathcal{P}')$, $\psi((\underset{\sim}{6}^*)\mathcal{P}')$, $\psi((\underset{\sim}{3})\bar{\mathcal{P}}')$; they can be produced by strong interactions only in pairs or with a charmed baryon.

(III) Two triplets and color.—In the model[10] with two fundamental fermion triplets t_1 and t_2 having charm $C = 1$ and $C = -2$, respectively, mesons with $C = 0$ can be constructed from structures $t_i\bar{t}_i$. The low-lying mesons will consist of a linear combination of these two structures, $(a_1t_1\bar{t}_1 + a_2t_2\bar{t}_2)$, and it is tempting to identify the newly observed mesons with the other linear combination $(b_1t_1\bar{t}_1 + b_2t_2\bar{t}_2)$ orthogonal to the low-lying mesons. However, the structure of a baryon with $C = 0$ is $t_1t_1t_2$, so both structures $t_1\bar{t}_1$ and $t_2\bar{t}_2$ have finite transition amplitudes to a baryon-antibaryon pair. Thus, the great stability of the new mesons causes difficulty for this model.

The only obvious way out seems rather artificial: cancelation of the amplitudes from $t_1\bar{t}_1$ and $t_2\bar{t}_2$ to baryon-antibaryon; i.e., $f(t_1\bar{t}_1 \to B\bar{B}) = (b_2/b_1) \times f(t_2\bar{t}_2 \to B\bar{B})$. Such a situation occurs if we assume that all the strong interactions of t_1 and t_2 are identical and $M_1 = M_2$. Then

$$f(t_1\bar{t}_1 \to \Omega) = f(t_2\bar{t}_2 \to \Omega), \tag{5}$$

where Ω is any particular final state of hadrons. Since t_1 and t_2 differ at least electromagnetically, there will be two classes of mesons: $2^{-1/2}(t_1\bar{t}_1 + t_2\bar{t}_2)$ and $2^{-1/2}(t_1\bar{t}_1 - t_2\bar{t}_2)$, symmetric and antisymmetric under $t_1 \leftrightarrow t_2$. These two classes are connected by electromagnetic interactions; but the antisymmetric mesons cannot decay into symmetric mesons by strong interactions invariant under $t_1 \leftrightarrow t_2$. The baryon attains $C = 0$ by breaking the symmetry of interchange of t_1 and t_2, but by Eq. (5) an antisymmetric meson cannot decay into a baryon-antibaryon pair. The antisymmetric meson states that can occur are the usual $(\underset{\sim}{8})$ and $(\underset{\sim}{3})$ of SU(3).

Such a scheme does not seem as attractive as the SU(3) $\otimes$ SU(3)$'$ scheme. One difficulty is to envisage the perturbation causing such a large mass difference between the symmetric and antisymmetric mesons while still keeping Eq. (5).

The situation in the three-triplet color scheme of red, white, and blue quarks[11] is similar. The ordinary mesons are symmetric in the color variables: $\bar{q}_Rq_R + \bar{q}_Wq_W + \bar{q}_Bq_B$. There are also $\bar{q}q$ states of type D (the two-dimensional representation of the three-element perturbation group) in the color indices, and with the usual representations $(\underset{\sim}{8})$ and $(\underset{\sim}{1})$ of SU(3). These are presumably the candidates for the lowest-lying charmed bosons; there is also a higher set of $(q\bar{q}\ q\bar{q})$ states that is totally antisymmetric (A) in the color indices. In this scheme the baryon is antisymmetric in color: $q_Rq_Wq_B + q_Wq_Bq_R + q_Bq_Rq_W$

$-q_W q_R q_B - q_B q_W q_R - q_R q_B q_W$, and the electromagnetic interaction is totally symmetric, being independent of color. Thus, mesons of symmetry D or A cannot decay strongly to ordinary mesons or to baryon-antibaryon pairs; nor can they decay or be produced electromagnetically.

[1]J. J. Aubert *et al.*, Phys. Rev. Lett. 33, 1404 (1974); J.-E. Augustin *et al.*, Phys. Rev. Lett. 33, 1406 (1974); C. Bacci *et al.*, Phys. Rev. Lett. 33, 1408 (1974).

[2]G. S. Abrams *et al.*, Phys. Rev. Lett. 33, 1453 (1974).

[3]We do not consider the possibility that $\psi(3.1)$ or $\psi(3.7)$ is the weak-interaction meson Z^0, as the masses appear an order of magnitude too low. In any case, this question is subject to fairly immediate experimental check by observation of $\psi \rightarrow l + \bar{\nu} +$ hadrons as a major decay component.

[4]M. Y. Han and Y. Nambu, Phys. Rev. 139, B1006 (1965).

[5]Here, with summation over $q = \mathcal{P}, \mathcal{N}, \lambda$,

$$J_\mu(3,0) = \textstyle\sum_i [\bar{\mathcal{P}}_i \gamma_\mu \mathcal{P}_i - \bar{\mathcal{N}}_i \gamma_\mu \mathcal{N}_i],$$

$$\sqrt{3}\, J_\mu(8,0) = \textstyle\sum_i [\bar{\mathcal{P}}_i \gamma_\mu \mathcal{P}_i + \bar{\mathcal{N}}_i \gamma_\mu \mathcal{N}_i - 2\bar{\lambda}_i \gamma_\mu \lambda_i],$$

$$J_\mu(0,3) = \textstyle\sum_q [\bar{q}_1 \gamma_\mu q_1 - \bar{q}_2 \gamma_\mu q_2],$$

$$\sqrt{3}\, J_\mu(0,8) = \textstyle\sum_q [\bar{q}_1 \gamma_\mu q_1 + \bar{q}_2 \gamma_\mu q_2 - 2\bar{q}_3 \gamma_\mu q_3].$$

The notation in Eq. (1) will be used throughout; namely, an object $F(i,j)$ refers to the ith index of the representation $(\underline{8})$ in SU(3), with $i=0$ indicating the representation $(\underline{1})$, and likewise for j and SU(3)$'$. In this notation the medium-strong interaction responsible for singlet-octet mixing is $H(8,0)$, and we assume it to be always present. To represent this mixing in observed particles, we write $F(\omega,j)$ and $F(\varphi,j)$ as distinct from the unmixed ideal states $F(0,j)$ and $F(8,j)$. It will be convenient to extend this device to $F(\rho,j)$ for $i=1,2,3$; to $F(i,K)$ for $j=4,5,6,7$; even to $F((\underline{8}),j)$ for $i=1$ to 8, etc. Denote by ψ the lowest vector particle configuration $^3S_1$ on the $\bar{q}q$ model. Then in the present notation, $\psi(\omega,0)=\omega$, $\psi(\varphi,0)=\varphi$, $\psi(3,0)=\rho^0$, $\psi(K,0)=K^*$, etc.

[6]By analogy with the ordinary ω-φ situation, we assume for these the combinations suggested by the simple quark model; *viz.* (suppressing the $^3S_1$ notation),

$$\psi(\omega,3) = \tfrac{1}{2}[(\bar{\mathcal{P}}_1\mathcal{P}_1 + \bar{\mathcal{N}}_1\mathcal{N}_1) - (\bar{\mathcal{P}}_2\mathcal{P}_2 + \bar{\mathcal{N}}_2\mathcal{N}_2)],$$

$$\psi(\varphi,3) = (\sqrt{2})^{-1}(\bar{\lambda}_1\lambda_1 - \bar{\lambda}_2\lambda_2),$$

$$\psi(\omega,8) = (2\sqrt{3})^{-1}[(\bar{\mathcal{P}}_1\mathcal{P}_1 + \bar{\mathcal{N}}_1\mathcal{N}_1) + (\bar{\mathcal{P}}_2\mathcal{P}_2 + \bar{\mathcal{N}}_2\mathcal{N}_2) - 2(\bar{\mathcal{P}}_3\mathcal{P}_3 + \bar{\mathcal{N}}_3\mathcal{N}_3)],$$

$$\psi(\varphi,8) = (\sqrt{6})^{-1}[(\bar{\lambda}_1\lambda_1 + \bar{\lambda}_2\lambda_2) - 2\bar{\lambda}_3\lambda_3].$$

[7]There may, of course, be further SU(3)$'$ multiplets like $\psi(i,(\underline{27}))$, but we assume them to be at much higher energies.

[8]The narrowness of the observed and presumably vector bosons implies that the lowest configuration in the SU(3)$'$ boson octet is $^3S_1$ instead of $^1S_0$, a minor variation from ordinary particles.

[9]B. J. Bjørken and S. L. Glashow, Phys. Lett. 11, 255 (1964).

[10]Y. Nambu, in *Preludes in Theoretical Physics*, edited by A. de-Shalit *et al.* (North-Holland, Amsterdam, 1966); H. Bacry, J. Nuyts, and L. van Hove, Phys. Lett. 9, 279 (1964).

[11]M. Gell-Mann, in *Proceedings of the Sixteenth International Conference on High Energy Physics, University of Chicago and National Accelerator Laboratory, 1972,* edited by J. D. Jackson and A. Roberts (National Accelerator Laboratory, Batavia, Ill., 1973), Vol. 4, p. 333.

Higgs-Particle Interpretation of Narrow e^+e^- Resonances

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(Received 16 December 1974)

We present a gauge model in which parity and $[\mathrm{SU}(2)\otimes \mathrm{Y}]_W$ are spontaneously broken. Three neutral Higgs particles couple directly to e^+e^- ($\mu^+\mu^-$) with coupling strengths proportional to a heavy-electron (muon) mass. Identification of two Higgs particles with e^+e^- resonances at 3.105 and 3.695 GeV limits the two heavy-lepton masses to 1–3 GeV. The heavy leptons are not produced by ν_L beams, but have sequential signature in e^+e^- production.

In this note we propose that the two narrow resonances observed in the e^+e^- system[1-4] are Higgs particles characteristic of an underlying, spontaneously broken, gauge symmetry of the weak and electromagnetic interactions. To support this interpretation for the newly discovered resonances we need a specific model with two or more neutral Higgs particles. Their couplings to e^+e^- must allow narrow widths for the e^+e^- decay channel, perhaps of the order of a few keV,