TABLE I. 90% confidence limits for the Lorentz-invariant production cross section of heavy penetrating particles of mass 3 GeV/c^2. To apply these limits for higher masses they should be divided by that fraction of the bombarding-energy interval for which the mass in question is above threshold (see text). For short-lifetime heavy particles the additional factor given in the text must be applied.

$P_\perp$ (GeV/c)	$E\,d^3\sigma/d^3p$ [cm^2 GeV/(GeV/c)^3]	Lower limit on mass range (GeV/c^2)
1.0 −1.25	6.4×10^{-35}	0.6
1.25−1.75	2.7×10^{-35}	1.0
1.75−2.25	2.4×10^{-35}	1.4

cay of a particle of specific mass and lifetime, the limits should be multiplied by $\exp(35m/\tau p)$, where m is the mass in GeV/c^2, τ, the lifetime in nanoseconds, and p, the momentum in GeV/c. Figure 3 shows the limits in the case of no decays and in the case of a mass-2-GeV/c^2 particle with lifetimes of 3.5 and 0.5 nsec. Results of two other Fermilab experiments are shown for comparison.

We acknowledge the assistance of the internal target area group and especially D. D. Yovanovitch. At certain phases of the experiment significant contributions were made by R. Piccioni, D. Reeder, and J. Sidles.

*Work supported by the U. S. Atomic Energy Commission and the National Science Foundation. Experiment performed at Fermi National Accelerator Laboratory.

†Alfred P. Sloan Foundation Fellow.

‡Visitor from Département de Physique des Particules Elémentaires, Centre d'Etudes Nucléaires, Saclay, France.

[1]Preliminary results for this experiment have been reported: D. Bintinger *et al.*, Contribution No. 576 to Session A3 of the Seventeenth International Conference on High Energy Physics, London, England, 1-16 July 1974 (to be published).

[2]L. B. Leipuner *et al.*, Phys. Rev. Lett. 31, 1226 (1973).

[3]J. Appel *et al.*, Phys. Rev. Lett. 32, 428 (1974).

[4]J. W. Cronin *et al.*, Phys. Rev. D 10, 3093 (1974).

[5]B. Alper *et al.*, Phys. Lett. 46B, 265 (1973).

[6]D. C. Carey *et al.*, Phys. Rev. Lett. 33, 327 (1974).

[7]J. W. Cronin *et al.*, "Production of Hadrons with Large Transverse Momentum at 200, 300, and 400 GeV," paper submitted to the Seventeenth International Conference on High Energy Physics, London, England, 1–10 July 1974 (to be published).

Colorful ψ's*

I. Bars and R. D. Peccei

Institute of Theoretical Physics, Department of Physics, Stanford University, Stanford, California 94305

(Received 16 December 1974)

We suggest that the recently discovered long-lived particles $\psi(3105)$ and $\psi(3695)$ are members of a family of "colored" vector mesons. Our scheme is a natural generalization of SU(6)⊗SU(3)′. The narrow widths of these states are understood because of color conservation and the difficulty of radiating colored photons. The spectroscopy of these new colored states, which is very rich, is discussed along with predictions for photoproduction.

The discovery of the long-lived ψ particles[1-4] has had a profound effect in the world of elementary particle physics. Many theoretical schemes of various degrees of plausibility have been proposed which can account for the existence of these particles.[5] In this note we would like to suggest a novel and simple interpretation of the ψ's as "color" vector mesons. Our scheme will account in a natural way for their existence and for their long lifetime. Furthermore the scheme makes numerous testable predictions.

The long lifetime of the $\psi(3105)$ and $\psi(3695)$ suggests that these particles possess a quantum number that prevents their strong decay into ordinary hadrons. A natural candidate is color. The invariance group of the strong interactions is enlarged from SU(3) to SU(3)⊗$G_{\rm color}$. In the Han-Nambu scheme,[6] which we shall adopt, $G_{\rm color}$ is also SU(3). In this model, color is observable (in contrast to the red, white, and blue model of

Bardeen, Fritzsch, and Gell-Mann[7]). However, the lowest states in the model, the ordinary hadrons, are supposed to be color singlets because of dynamical properties of the quark interaction.

The quarks in the Han-Nambu model have integral charge and transform under SU(3)⊗ SU(3)′ according to the $(\underline{3},\underline{3}^*)$ representation. The electromagnetic current transforms according to the $(\underline{1},\underline{8})\oplus(\underline{8},\underline{1})$ representation and the electric charge is given by

$$Q=I_3+\tfrac{1}{2}Y+I_3{}'+\tfrac{1}{2}Y', \tag{1}$$

where $\vec{I}'$ and Y' are the isospin and hypercharge of the color SU(3)′ group.

We shall suppose that the ψ's are hadrons so that they do not have any *direct* coupling to the leptons. They can only be produced in e^+e^- collisions because they mix with the photon. Thus they must transform (at least in part) according to the $(\underline{1},\underline{8})$ representation of SU(3)⊗ SU(3)′. We shall suppose, in fact, that they do not transform according to a pure representation of SU(3)⊗SU(3)′ but rather that they are mixtures of the $(\underline{1},\underline{8})$ and the $(\underline{8},\underline{8})$ representations. This assumption is *natural* given our experience with ordinary hadrons. There too vector mesons are not pure SU(3) states but mixtures of octets and singlets. We remember that the mixing angle is naturally understood in SU(6) by classifying the mesons in the $\underline{35}$ representation. Generalizing this idea, we now expect that colored mesons classified as $(\underline{35},\underline{8})$ exist. This naturally gives a mixing of the $(\underline{1},\underline{8})$ and $(\underline{8},\underline{8})$ vector mesons.

In a quark model the $I=0$, $Y=0$, SU(3) singlet and octet states are mixed. The physical particles φ and ω are given by

$$\begin{aligned}\omega&=\cos\theta\varphi_{8,1}+\sin\theta\varphi_{1,1},\\ \varphi&=-\sin\theta\varphi_{8,1}+\cos\theta\varphi_{1,1},\end{aligned} \tag{2}$$

where $\varphi_{8,1}$ and $\varphi_{1,1}$ are, respectively, octets and singlets in ordinary SU(3) (and singlets in color). The mixing angle is

$$\tan\theta=1/\sqrt{2}, \tag{3}$$

a value in excellent agreement with experiment. Then, in a quark model, φ is pure $\lambda\bar{\lambda}$ while $\omega=(1/\sqrt{2})(\mathcal{P}\bar{\mathcal{P}}+\mathfrak{N}\bar{\mathfrak{N}})$. Our suggestion is that the $\psi(3105)$ and the $\psi(3695)$ are also mixtures of octet and singlet states, which are now, however, color octets. If we considered SU(6)⊗SU(3)′, a natural suggestion would be that they are precisely the same mixture as φ and ω. Then

$$\begin{aligned}\psi^a(3105)&=\cos\theta\varphi_{8,8}{}^a+\sin\theta\varphi_{1,8}{}^a\equiv\omega^a,\\ \psi^a(3695)&=-\sin\theta\varphi_{8,8}{}^a+\cos\theta\varphi_{1,8}{}^a\equiv\varphi^a.\end{aligned} \tag{4}$$

We have introduced the suggestive notation ω^a and φ^a, $a=1,\ldots 8$. The mixing of these states with the photon occurs only in the u'-spin singlet combination

$$\varphi_{1,8}{}^3+(\sqrt{3})^{-1}\varphi_{1,8}{}^8=\tfrac{2}{3}\sqrt{2}\left[\tfrac{1}{2}\sqrt{3}\,\omega^3+\tfrac{1}{2}\omega^8\right]+\tfrac{2}{3}\left[\tfrac{1}{2}\sqrt{3}\,\varphi^3+\tfrac{1}{2}\varphi^8\right]. \tag{5}$$

If color is broken only in the direction of electromagnetism then the u'-singlet states are directly identifiable with the observed ψ's. The other fourteen states in ω^a and φ^a, as well as the rest of unmixed $(\underline{8},\underline{8})$ mesons, are accessible in the e^+e^- channel only in pairs.

Besides the mixed states ω^a and φ^a there are 56 other states which belong to the $(\underline{8},\underline{8})$ representation: $\vec{\rho}^{\,a}$ with $I=1$, $Y=0$ (24 states); and $K_a{}^*$ and $\bar{K}_a{}^*$ with $I=\frac{1}{2}$, $Y=\pm1$ (32 states). If, as we still assume, color breaking is small, then one will see mainly only an ordinary SU(3)-breaking pattern. In this case we expect the 24 $\vec{\rho}^{\,a}$ states to lie at roughly the same mass as the $\psi(3105)$ and the strange-color vector mesons to have a mass given by

$$M_{K_a{}^*}{}^2=\tfrac{1}{2}[M_{\psi(3105)}{}^2+M_{\psi(3695)}{}^2]\cong11.6\ \mathrm{GeV}^2. \tag{6}$$

Because color is broken by electromagnetism we also expect mass breaking of perhaps 20–50 MeV among the charged and neutral members of each of the color multiplets. The spectroscopy of these states is very rich and varied, and the observation of the partners of the $\psi(3105)$ and $\psi(3695)$ in hadronic experiments will be a crucial test of the scheme. We should remark in this connection that some of these vector mesons have charges of ±2, and that we expected colored pseudoscalars also in this 3–4 GeV range.

The mixing scheme discussed above provides a *natural explanation for the presence of two narrow ψ's*. It makes many other predictions also. For example, in a quark model, we expect that the coupling of the photons to the vector mesons,

$$\langle0|J_\mu{}^{\mathrm{EM}}|V\rangle=c(m_v{}^2/f_v)\epsilon_\mu, \tag{7}$$

be given by the ratios

$$f_\rho{}^{-2}:f_\omega{}^{-2}:f_\varphi{}^{-2}:f_{\psi(3105)}{}^{-2}:f_{\psi(3695)}{}^{-2}=9:1:2:8:4. \tag{8}$$

These ratios yield leptonic widths for these vector mesons which, given the difficulty of including broken mass effects in these calculations, are in satisfactory agreement with experiment.

Up to now we have not discussed in detail why the width of the ψ's is so narrow. Since color is supposed to be conserved in our scheme, up to electromagnetic breaking, the process $\psi \to$ ordinary hadrons is forbidden to occur strongly. It can occur via electromagnetic breaking, and one expects a width of the order $\Gamma \sim \alpha^2\Gamma_s$, where Γ_s is a typical strong width ($\Gamma_s \sim 100-200$ MeV). However, radiative decays of the type $\psi \to \gamma +$ ordinary hadrons can occur, since the photon carries away the color. If one tries to estimate the width for these kinds of processes by comparing them to the radiative decays of ordinary hadrons one runs immediately into trouble. A typical radiative width (like $\omega \to \pi\gamma$) is of the order $\frac{1}{2}-1$ MeV. In the case of the ψ's, to make matters worse, phase space, and the availability of more channels certainly increase this estimate. The values extracted from the data[8] for the ψ's width are $\Gamma_{\psi(3105)} \cong 90$ keV, 200 keV $\leq \Gamma_{\psi(3695)} \lesssim 1$ MeV in gross contradiction with these estimates. *For any color scheme to be truly successful one has to understand why these estimates are misleading.* We shall provide a possible explanation shortly.

A second possible difficulty for color schemes arises from the observation of the decay[9] $\psi(3695) \to \psi(3105)\pi^+\pi^-$ with perhaps a branching ratio as high as 40%. The decay is also not forbidden by color and yet it has a small width. If one compares it, for instance, with the decay $\rho' \to \rho\pi^+\pi^-$, and corrects for phase space, one sees that the decay width for the $\psi(3695) \to \psi(3105)\pi^+\pi^-$ decay is a factor of 10^{-2} smaller than this strong decay. Our model can provide a "natural" explanation for such a suppression, as discussed below.

The decay of the $\varphi \to 3\pi$ has a very small width ($\Gamma \cong 600$ keV). The mechanism for the suppression of this ordinarily allowed decay is not clearly understood. However, in quark models, one has developed an empirical rule (sometimes known as Zweig's rule) by which processes which involve disconnected quark diagrams are suppressed by about 2 orders of magnitude in rates. Since the φ is purely $\lambda\bar{\lambda}$ its decay into three π's is suppressed by Zweig's rule. In our model the $\psi(3695)$ is also a pure $\lambda\bar{\lambda}$ state, while the $\psi(3105)$ does not contain any λ quarks. Hence, at least empirically, it is not disconcerting to find that the decay width for $\psi(3695) \to \psi(3105)\pi^+\pi^-$ is suppressed, after phase-space corrections, by an additional 10^{-2} relative to the $\rho' \to \rho\pi^+\pi^-$ width, where no Zweig-rule constraints apply.

It remains to understand why the ψ's do not have large widths. One needs to suppress the naive estimate made above for radiative decays by perhaps as much as 10^{-2}. Effectively one wants $\Gamma_{\rm rad} \sim \alpha^2\Gamma_s$ and not $\Gamma_{\rm rad} \sim \alpha\Gamma_s$. A possible explanation for this extra suppression follows, in a vector dominance model, if it were more difficult to convert high-mass vector mesons into photons than it is the case for ordinary vector mesons. Let us assume that we can write an effective Lagrangian for the interaction of two ($\underline{1}$, $\underline{8}$) vector mesons with an SU(3)$\otimes$SU(3)$'$ singlet as

$$\mathcal{L} = g\,\mathrm{Tr}(\varphi_{1,8}{}^{\mu}\varphi_{1,8\mu})S. \tag{9}$$

Radiative decays of the $\psi(3105)$ and $\psi(3695)$ can then be calculated by vector mesons dominating one of the $\varphi_{1,8}$. Unless one assumes some direct coupling between photons, colored vector mesons, and ordinary hadrons, the above term will account for *all* the radiative decays. Besides the strong coupling constant g, the matrix element will contain a propagation factor for the vector meson and a transition form factor for a vector meson to convert into an *on-mass-shell proton.* Hence the amplitude for radiative decays will be proportional to

$$A \propto g\frac{1}{M_v^2}e\frac{M_v^2}{f_v(0)} \sim \frac{ge}{f_v(0)}, \tag{10}$$

where $f_v(0)$ is the transition form factor. We remark, however, that what is measured by the decay width of vector mesons (including the ψ's) into lepton pairs is the transition form factor on the *vector-meson-mass shell;* $f_v(M_v^2)$, which we denoted simply by f_v in Eqs. (7) and (8). There is no *a priori* reason to suppose that $f_v(0) \cong f_v(M_v^2)$. In fact, since here $M_v^2 \cong 10$ GeV^2, as compared to the case of ordinary vector mesons, it might well be that here $1/f_v(0) \ll 1/f_v(M_v^2)$ and radiative decays are suppressed. We realize full well that this "explanation" for the paucity of radiative decays is obtained only within the context of an approximate dynamical model. It remains to be seen whether it works more generally in quark models. However, we feel that it is not physically implausible that high-mass states cannot convert as easily into photons as low-mass states do.

The above suppression mechanism has direct implications for photoproduction, since again one has $f_v(0)$ appearing instead of $f_v(M_v^2)$. From the point of view of our scheme it is logical to compare ω photoproduction to $\psi(3105)$ photoproduction, and φ photoproduction to $\psi(3695)$ photoproduction. One has then, according to standard lore,

$$\frac{\sigma_{\gamma N \to \psi(3105) N}}{\sigma_{\gamma N \to \omega N}} \cong \frac{f_\psi^{-2}(0)}{f_\omega^{-2}(0)} \left[\frac{\sigma_{\text{tot}}(\psi N)}{\sigma_{\text{tot}}(\omega N)}\right]^2 [t_{\min} \text{ effects}], \quad (11a)$$

$$\frac{\sigma_{\gamma N \to \psi(3695) N}}{\sigma_{\gamma N \to \varphi N}} = \frac{f_{\psi'}^{-2}(0)}{f_\varphi^{-2}(0)} \left[\frac{\sigma_{\text{tot}}(\psi' N)}{\sigma_{\text{tot}}(\varphi N)}\right]^2 [t_{\min} \text{ effects}]. \quad (11b)$$

Therefore, we get suppression factors (taking literally our rule $\Gamma_\psi \sim \alpha^2 \Gamma_{\text{strong}}$)

$$\frac{f_\psi^{-2}(0)}{f_\omega^{-2}(0)} = \frac{m_\omega \Gamma_{\psi \to e^+e^-}}{m_\psi \Gamma_{\omega \to e^+e^-}} \frac{f_\psi^{-2}(0)}{f_\psi^{-2}(10)}, \cong 2\times 10^{-2}, \quad (12a)$$

$$\frac{f_{\psi'}^{-2}(0)}{f_\varphi^{-2}(0)} = \frac{m_\omega \Gamma_{\psi' \to e^+e^-}}{m_{\psi'} \Gamma_{\varphi \to e^+e^-}} \frac{f_{\psi'}^{-2}(0)}{f_{\psi'}^{-2}(14)}. \cong 1\times 10^{-2}. \quad (12b)$$

The $t_{\min}$ effects are probably not very important for photon energies > 50 GeV, but they may become important at lower energies. The total cross sections for ψN and $\psi' N$ scattering are of course unknown, but one may expect them to be of the same order of magnitude (within factors such as 2 or 3) as ωN and φN total cross sections, respectively. If we take them to be equal to get the crudest estimate at high energies (neglecting $t_{\min}$ effects), then we simply obtain the suppression factors of Eqs. (12). Needless to say, our estimate is extremely crude, but we expect it to be correct within one order of magnitude.

Our scheme seems to be in rough agreement with the present data. It will not be difficult to test various aspects of it as more data become available.

We would like to thank all the members of the Stanford Linear Accelerator Center (SLAC) theory workshop for stimulating these ideas.

*Research sponsored by the National Science Foundation Grant No. MPS 073-08916.

[1]J.-E. Augustin *et al.*, Phys. Rev. Lett. 33, 1406 (1974).
[2]J. J. Aubert *et al.*, Phys. Rev. Lett. 33, 1404 (1974).
[3]C. Bacci *et al.*, Phys. Rev. Lett. 33, 1408 (1974).
[4]G. S. Abrams *et al.*, Phys. Rev. Lett. 33, 1453 (1974).
[5]See, for example, H. Harari, "ψ chology" (unpublished); I. Bars *et al.*, "Theoretical Speculations on the Nature of the ψ (3105) and ψ (3695)" (unpublished).
[6]M. Y. Han and Y. Nambu, Phys. Rev. 139, B1006 (1965).
[7]W. A. Bardeen, H. Fritzsch, and M. Gell-Mann, in *Scale and Conformal Symmetry in Hadron Physics*, edited by R. Gatto (Wiley, New York, 1973).
[8]Report of the Quantum Electrodynamic Subgroup of the Stanford Linear Accelerator Center workshop (unpublished.
[9]Stanford Linear Accelerator Center–Lawrence Berkeley Laboratory experimental group, private communication.
[10]J. F. Martin *et al.*, Phys. Rev. Lett. 34, 288 (1975).

Asymptotic SU(4) in the l^+l^- Annihilation of New Resonances

E. Takasugi and S. Oneda
Center for Theoretical Physics, Department of Physics and Astronomy, University of Maryland, College Park, Maryland 20742
(Received 19 December 1974)

The l^+l^- decays of the new narrow resonances are discussed as a good test of the hypothesis of asymptotic SU(4) and of the assignment of $\psi(3105)$ to an uncharmed member of the set of 1^{--} mesons belonging to the $15\oplus 1$ dimensional representation of SU(4).

We consider $\psi(3105) \to l + \bar{l}$ decays, assuming that $\psi(3105)$[1] is the uncharmed vector meson φ_c which belongs to the $15\oplus 1$ dimensional representation of SU(4) together with ρ, K^*, φ, ω, D^*, and F^*.[2] In the naive quark language the narrow width of φ_c may be explained if the φ_c is primarily a $c\bar{c}$ bound state.[2] However, one can also treat this problem in an entirely *algebraic* way without referring to quarks.[3]

Since SU(4) is expected to be *more* violated than