New elementary particles

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Institute of Theoretical and Experimental Physics
Usp. Fiz. Nauk **117**, 227–312 (October 1975)

The first reports of the discovery of new heavy neutral bosons appeared in the autumn of 1974. The purpose of this review is to describe the properties of these particles and their interpretation within the framework of various theoretical models. Problems connected with the search for so-called charmed particles are also discussed.

PACS numbers: 14.40.–n, 12.40.Dd

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I. INTRODUCTION

In November 1974 a remarkable discovery was made in elementary-particle physics: a new particle was discovered—a narrow resonance of very large mass—3.3 times the mass of the proton (3.1 GeV). The width of the new resonance was found to be very small—70 keV; i.e., its lifetime is 10^{-20} sec—a large value by nuclear standards. The new particle was discovered almost simultaneously in two laboratories in the U.S.A. It was detected at Brookhaven National Laboratory by S. Ting's group, who analyzed the effective-mass spectrum of electron-positron pairs produced in collisions of 30-GeV protons with beryllium nuclei[1] 1). It was also found using the colliding electron-positron beams of the Stanford accelerator by anlyzing the cross section of electron-positron annihilation[2].

The history of how this resonance was discovered is of some interest. Ting's group began their experiment in the spring of 1974. The experiment which they carried out was an exceedingly difficult one: it is sufficient to say that the probability for producing an electron-positron pair in the appropriate effective-mass region in collisions of protons with the Be nucleus was $\sim 10^{-8}$ of the total probability for inelastic collisions. Similar experiments on the reaction p + nucleus $\to \mu^+\mu^-$ + anything had been carried out previously by Lederman's group[3] and yielded a number of interesting results: the cross section for muon pair production was determined as a function of the effective mass of the pair, revealing a "shoulder" in the mass spectrum of the muon pairs in the region 3–4 GeV. Lederman's group did not see any resonances in the mass spectrum of the muon pairs. However, the experimental arrangement of Ting's group had a major advantage over that of Lederman: its resolution in the energy of the detected pairs was much better, and this facilitated the search for narrow resonances. By autumn, Ting's group accumulated a sufficient number of events (about 500 by the end of October), which fell mainly within a narrow band of the effective-mass spectrum of the e^+e^- pairs in the vicinity of 3 GeV.

In 1973–74, B. Richter, M. Perl, G. Goldhaber, M. Chinowsky and their coworkers at Stanford (SLAC) used the colliding electron-positron beams of the SPEAR accelerator to measure the cross section for electron-positron pair annihilation into strongly interacting particles—hadrons, and these measurements yielded a number of valuable results2). Until June 1974, SPEAR was operating with each of the beams at an energy 2.5 GeV; in June, it was closed down for modifications, with an increase in the energy to 4–5 GeV; it came into operation again in September. Richter's group decided to use the recently closed-down accelerator to repeat their experiments which were discontinued in June, but with more closely spaced steps in energy. Starting the measurements on November 9th, it was not long before they observed an enhancement in the cross section for the annihilation $e^+e^- \to$ hadrons in the region 3.1–3.2 GeV; then, raising the accuracy with which the energy of the colliding beams was measured to 0.1%, on November 10th they made the first conclusive observation of a resonance—a more than tenfold (and, in several hours, after raising the accuracy to 0.01%, even a hundredfold) enhancement in the cross section.

Two papers concerning the discovery of the new particle were submitted to Physical Review Letters, one after the other, on November 12 and 13. Ting's group designated the new particle the J, while Richter's group designated it the ψ. When the news of the discovery of the J-ψ particle became known in Europe by telephone, a group of Italian physicists who were working with the colliding e^+e^- beams of the ADONE accelerator at Frascati (Italy) raised the energy of the accelerator to $2 \times 1.55 = 3.1$ GeV—the value required to produce the resonance (the accelerator was previously operating at an energy $2 \times 1.5 = 3.0$ GeV)—and, within several days,

1)The references to the cited literature in this paper are numbered separately in each chapter.

2)See the reviews [4,5] for a discussion of the results of these experiments.

they also observed this resonance in the same reaction, $e^+e^- \to$ hadrons, in which it was seen at Stanford. The Italian physicists dictated their paper by telephone to the editorial office of Physical Review Letters. As a result, the issue of Physical Review Letters dated December 2, 1974 contained at once three papers[1,2,6] reporting the discovery of the new particle. Before long, one more particle was discovered using the intersecting electron-positron rings at Stanford, this particle being similar to the first one, with a mass equal to 3.9 times the mass of the proton (3.7 GeV) and a width of the order of several hundred kiloelectron volts[7]. Since they were seen as resonances in the e^+e^- system, both new particles must have integral spin, i.e., they are bosons. We shall henceforth call these new particles ψ bosons (or ψ mesons); we shall refer to the boson with mass 3.1 GeV as simply the ψ, and the boson with mass 3.7 GeV as the ψ'.

The properties of the new particles were surprising mainly because of their small widths for such large masses. The previously known particles—the heavy mesons and baryons—with masses in the region 2—3 GeV, had widths of the order of several hundred MeV, i.e., three orders of magnitude greater than those of the new particles. The small widths of the new mesons therefore indicated the discovery of some new type of elementary particle, whose decay into the ordinary strongly interacting particles—pions, kaons and nucleons—is strongly suppressed.

The possible existence of new types of particles had been widely discussed by physicists in recent years, and in this sense their discovery was not surprising (moreover, there had even been theoretical proposals[8-10] to search for new particles in experiments on e^+e^- annihilation, where they were indeed discovered, as well as in photoproduction[11], where they were found later). However, these particles were expected[10,11] to have larger widths than those that are actually observed or to have completely different properties, so that the discovery was, as it were, a fortuitous bonus.

The small widths for the decays of the new mesons into ordinary hadrons might mean either that the new mesons have no strong interaction at all with the ordinary hadrons or that their decays are strongly suppressed by some selection rules.

The discovery brought about an avalanche of theoretical papers providing interpretations of the new particles. First of all, it was necessary to establish whether these particles are hadrons, having a strong interaction, or whether they belong to the class of particles which have only weak and electromagnetic interactions, i.e., whether they are so-called intermediate bosons.

The possible existence of bosons which have only weak and electromagnetic interactions—intermediate bosons—has been discussed for many years in connection with the problem of weak interactions (see, e.g.,[12]). In analogy with the photon, which carries electromagnetic interactions, it was assumed that there exist charged intermediate bosons which are carriers of weak interactions involving the transfer of charge from hadrons to leptons, such as the process of β decay. A number of theoretical arguments related to the construction of a unified theory of weak and electromagnetic interactions on the one hand, and the more recent discovery of weak neutral currents in a neutrino experiment on the other hand, indicate that there may also exist neutral intermediate bosons Z^0. One might have tried to identify such neutral intermediate bosons Z^0 with the new mesons, since theoretical estimates show that the Z^0 widths should be similar to those that are observed experimentally. However, on the basis of the unified theory of weak and electromagnetic interactions, the Z^0 bosons were expected to have masses of the order of several tens of GeV, so that their discovery at 3 GeV was unexpected. Further investigations of the properties of the ψ bosons established that parity is conserved in the interactions of the ψ with leptons and that, although the decays $\psi \to$ hadrons indicate a semi-weak interaction of the ψ with hadrons (as in the case of intermediate bosons), the interaction of the ψ with hadrons is strong in other processes, such as the decays $\psi' \to \psi$ + hadrons. This shows that the ψ cannot be the neutral intermediate bosons of the weak interactions. (Certain possible ways of describing the ψ as intermediate bosons which have so far not been excluded experimentally are discussed here in Chap. V.)

Thus, the only remaining plausible description of the ψ mesons is the possibility that they are hadrons but that, because of some selection rules, such as the presence of new quantum numbers for these mesons, their decay into ordinary hadrons is strongly suppressed (although the decay $\psi' \to \psi$ + hadrons is allowed).

During the past few years, theoretical predictions have been made of the possible, and in certain cases even necessary, existence of hadrons with new, hitherto experimentally unknown quantum numbers, which are conserved in the strong interactions. Theoretically, on the basis of two different arguments, two hypotheses have been put forward regarding the existence of hadrons with two new quantum numbers. One of these quantum numbers has become known as "charm" and the other as "color."

The need to introduce the new quantum number "charm" in the theory originated from the physics of weak interactions. By analyzing virtual processes at high energies, the theory of weak interactions which assumes that there exist only ordinary and strange particles (i.e., that the symmetry of the hadrons is SU(3)) was used to infer that there is a relatively large probability for decays of strange particles into lepton pairs with zero charge (such as $K^0_L \to \mu^+\mu^-$; see the review[13]). Experimentally, such decays were not seen with the expected probability. To overcome this difficulty, it was conjectured[14] that there exist hadrons which possess a new quantum number—charm. The strong interactions are approximately symmetric with respect to all types of hadrons, including charmed hadrons (if there exists a single type of charmed particle, i.e., a single new quantum number, then the symmetry of the strong interactions is SU(4); if there are two types, then the symmetry is SU(5), etc), while the structure of the weak interactions is such that the contribution of charmed particles completely cancels the contribution of the ordinary particles in the amplitudes for decays of the type $K^0_L \to \mu^+\mu^-$ in virtual processes at high energies. In the hypothesis of charmed particles, the "charm" quantum number C is conserved in the strong and electromagnetic interactions, but not in the weak interactions.

An important difference between charmed and ordinary particles is that, whereas the average charge within any SU(3) multiplet of ordinary particles is $\langle Q \rangle = 0$, an SU(3) multiplet of charmed particles has $\langle Q \rangle \neq 0$, so that each such multiplet can be characterized by the value $\sigma = 3\langle Q \rangle$, which has been designated "supercharge"[15].

In several respects, the physics of charmed particles should be analogous to that of strange particles. Like strange particles, charmed particles should be produced in pairs. The charmed mesons and baryons of lowest mass can decay into ordinary or strange particles only via the weak interactions and, by nuclear standards, should have rather long lifetimes $\sim 10^{-13}$–10^{-14} sec.

In terms of the hypothesis of charmed particles, the ψ bosons cannot be charmed themselves and must have the quantum number $C = 0$. This statement follows from the fact that the ψ are produced in electron-positron collisions, which in the conjectured scheme must proceed via a virtual photon, $e^+e^- \to \gamma \to \psi$, as well as from the values of the widths and lifetimes of the ψ mesons. The small widths of the ψ mesons must therefore be explained by invoking the additional hypothesis that the ψ and ψ' are, as it were, weakly bound systems of two charmed particles with $C = \pm 1$, which have a very small probability of coming together at a single point and annihilating into ordinary hadrons.

The introduction of the new quantum number "color," which can take three values, say red, blue and yellow, is due to a number of theoretical advantages of models involving this quantum number (see, e.g.,[16]). It is assumed in such models that, in addition to the usual SU(3) symmetry, the strong interactions are SU(3)'-symmetric with respect to transformations in the three-dimensional space of the three colors and that the carriers of the colored interactions are 8 vector mesons.

As the interaction in "color" theories is carried by an octet of vector particles, the effective charge in this case does not grow at small distances, as in quantum electrodynamics or in meson theories, but decreases; i.e., the theory does not involve the so-called "zero-charge" problem—the vanishing of the physical charge for any value of the bare charge. The decrease of the effective charge at small distances means that the particles appear as free particles asymptotically, at large momentum transfers ("asymptotic freedom"). This conclusion is confirmed by experimental data on deep inelastic electroproduction and neutrino scattering.

Another virtue of color theories is that they provide a consistent description of the data on baryon spectroscopy in terms of model-dependent concepts (the quark model; see Chap. IV). Theories involving color symmetry also have other advantages.

According to the hypothesis of color, the ordinary hadrons are mixtures of all three colors with various weights, i.e., they are "white." In the framework of the color hypothesis, there may appear a large number of new hadrons, which are distinguished by the new quantum numbers. In this case, however, the method of incorporating electromagnetic and weak interactions in the theory is not so well determined by the properties of the scheme as in the case of the hypothesis of charm. In particular, it may be assumed that the electromagnetic interaction violates the color symmetry. It is this variant of the color hypothesis which is preferred for the description of the ψ bosons. In this case, the ψ bosons are colored particles. It is possible to explain their small width for the decay into hadrons and their production in the e^+e^- reaction: both processes take place because of the violation of color by the electromagnetic interaction. One of the main consequences for the ψ mesons is that the principal decay mode is $\psi \to \gamma$ + hadrons, i.e., ψ decays must involve photons in the majority of cases.

With either hypothesis—charm or color—it is possible to make a number of predictions regarding the types and masses of new particles and their decays and interactions, which can be compared with experiment, both in processes involving the ψ mesons which have already been studied and in possible experimental searches for new particles. One of the major problems here is the question of whether these hypotheses can account for the results of experiments on e^+e^- annihilation into hadrons[5], which were not incorporated in the theoretical scheme before the discovery of the ψ mesons. In comparing the two hypotheses with the existing experimental data, it cannot be said that everything goes smoothly: each of the hypotheses has its difficulties. (This is particularly true of the color hypothesis.)

In addition to the hypotheses of intermediate bosons, charmed particles and colored particles, other theoretical schemes for describing the ψ mesons have been proposed. However, they have not been developed in such great detail as the foregoing schemes.

Regardless of whether any one of the theoretical schemes which we are considering here proves to be correct, now that the ψ mesons have been discovered we can expect a vigorous development of elementary-particle physics in the course of the next few years—the discovery of new particles and the study of their properties, which may be the key to a clearer understanding of the laws of nature.

In Chapters II and III of this review we shall outline the experimental facts about the ψ mesons and give their theoretical interpretation on the basis of the phenomenological approach, as far as possible without invoking any theoretical hypotheses. Then, in Chap. IV, we shall consider the basic ideas of the above-mentioned theoretical hypotheses—charm and color—in the language of the quark model. A discussion of alternative interpretations of the ψ mesons appears in Chap. V. Finally, Chap. VI contains a theoretical analysis of the production mechanisms of the new particles.

II. EXPERIMENTS ON THE PRODUCTION OF NEW PARTICLES IN COLLIDING ELECTRON-POSITRON BEAMS

The main physical characteristics of the ψ mesons (their widths, spins, parities, etc.) have been determined in experiments on the electron-positron pair annihilation process $e^+e^- \to \psi$ carried out with the colliding-beam accelerator SPEAR (Stanford, U.S.A.)[1-5]. Data on the annihilation process $e^+e^- \to \psi$, in agreement with the results of SPEAR, have also been obtained using the accelerators ADONE (Frascati, Italy)[6] and DORIS (DESY, Hamburg, West Germany)[7].

The experimental arrangement of SPEAR[5,8]—a magnetic detector (Fig. 1)—consisted of a solenoid of diameter 3 m and length 3 m, with a uniform magnetic field of 4 kG, whose axis was parallel to the beam axis. The particles which appeared in the interaction region

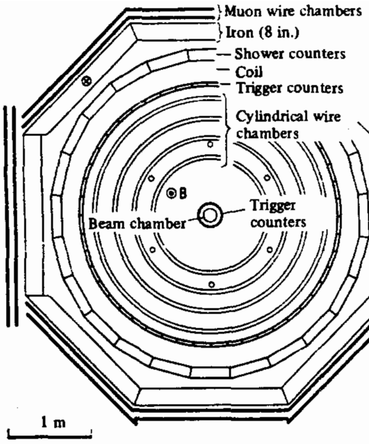


FIG. 1. The experimental arrangement of the SPEAR magnetic detector.

traversed in turn a vacuum chamber, an inner trigger (3 mm), a scintillator to suppress the cosmic-ray background, a magnetostriction wire spark chamber, a cylindrical layer of 48 plastic scintillators of thickness 2.8 cm as outer triggers, a magnetic coil (with a thickness of one radiation length), a cylindrical array of 24 scintillation counters (with lead) to identify electrons, a 20-cm iron yoke of the magnet which also served as a hadron filter, and finally two gaps of planar spark chambers to detect muons. The apparatus subtended a solid angle $0.65 \times 4\pi$ sr, which was symmetric with respect to 90° to the beam axis. The SPEAR data[3)] on the cross sections for the processes $e^+e^- \to$ hadrons, $e^+e^- \to e^+e^-$ and $e^+e^- \to \mu^+\mu^-$ in the vicinity of the resonance ψ (3.1 GeV) and the process $e^+e^- \to$ hadrons in the vicinity of the resonance ψ' (3.7 GeV) are shown in Figs. 2 and 3. These data demonstrate that there is not the slightest doubt about the existence and small widths of the resonances ψ and ψ'. The experimental resolution in the total energy of the beams was approximately 2.8 MeV in the SPEAR experiments, so that only an upper bound $\Gamma < 2.8$ MeV can be obtained on the widths of the resonances from the direct experimental data.

In the course of further, more accurate measurements of the energy dependence of the cross section for e^+e^- annihilation into hadrons at SPEAR, a maximum was seen in the cross section at a total e^+e^- c.m. energy equal to 4.15 GeV, which can be interpreted as a resonance. The experimental data on the total cross section for the annihilation process $e^+e^- \to$ hadrons are shown in Fig. 4. We shall henceforth denote this resonance by ψ (4.15). As can be seen from Fig. 4, the ψ (4.15), unlike the ψ and ψ', is a broad resonance, with a width of about 200—300 MeV.

1. Determination of the resonance widths. Allowance for radiative effects

The values of the partial widths of the resonances can be determined from the experimental curves of Figs. 2 and 3 by making use of the Breit-Wigner formula

$$\sigma_f(E) = \pi\lambdabar^2 \frac{2J+1}{4} \frac{\Gamma_{ee}\Gamma_f}{(E-M)^2+(\Gamma^2/4)}, \qquad (2.1)$$

[3)]By the middle of February 1975, about 100 000 events had been seen in the SPEAR experiments in the region of the ψ resonance and about 30 000 events in the region of ψ'.

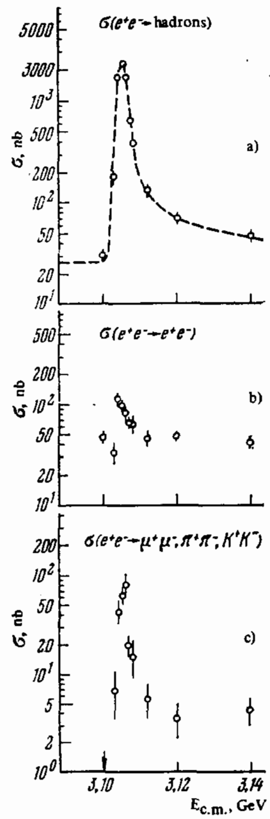


FIG. 2. The e^+e^- annihilation cross section in the region of energies near the ψ (3.1) resonance. The total annihilation cross section $\sigma(e^+e^- \to \text{hadrons})$ (a), $\sigma(e^+e^- \to e^+e^-)$ (b), and $\sigma(e^+e^- \to \mu^+\mu^-)$ (c). It was not possible to distinguish muons from pions and kaons in this experiment, but subsequent experiments showed that the contribution of pions and kaons to the cross section $\sigma(e^+e^- \to \mu^+\mu^-, \pi^+\pi^-, K^+K^-)$ is small.

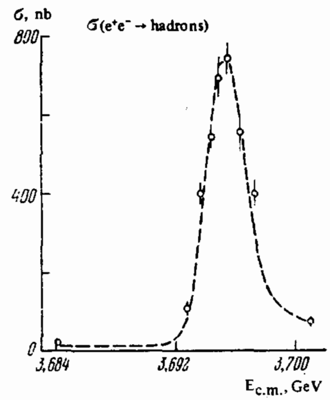


FIG. 3. The cross section for the annihilation process $e^+e^- \to$ hadrons near the ψ' (3.7) resonance.

where σ_f is the cross section for the partial process $e^+e^- \to \psi \to f$ (for example, f = h (hadrons), e^+e^- or $\mu^+\mu^-$), $\lambdabar \approx 2/M$ is the wavelength of the e^+ or e^- in the c.m.s., M is the mass of the resonance, J is its spin, Γ_{ee} is the width of the decay into e^+e^-, Γ_f and Γ are the partial and full widths, and E is the total e^+e^- energy in the c.m.s. If the experimental resolution ΔE is much greater than the width of the resonance, $\Delta E \gg \Gamma$, then the value of $\sigma_f(E)$ within the peak is determined by an integral of Eq. (2.1) with respect to the energy:

$$\sigma_f^{\text{peak}} \approx \frac{1}{\Delta E}\int \sigma_f(E)\,dE = \frac{2\pi^2}{\Delta E}\frac{2J+1}{M^2}\Gamma_{ee}\frac{\Gamma_f}{\Gamma}. \qquad (2.2)$$

In the case of a Gaussian distribution in the energies of the e^+e^- beams, which holds in the experiments under consideration, we have

$$w(E, E')\,dE = \frac{1}{\sqrt{2\pi}\,\sigma} e^{-(E'-E)^2/2\sigma^2},$$

giving

$$\Delta E = \sqrt{2\pi}\sigma \qquad (2.3)$$

Using the experimental values within the peak (or, as is more often more convenient, $\int \sigma_f^{\exp}(E)\,dE$), Eq. (2.2)

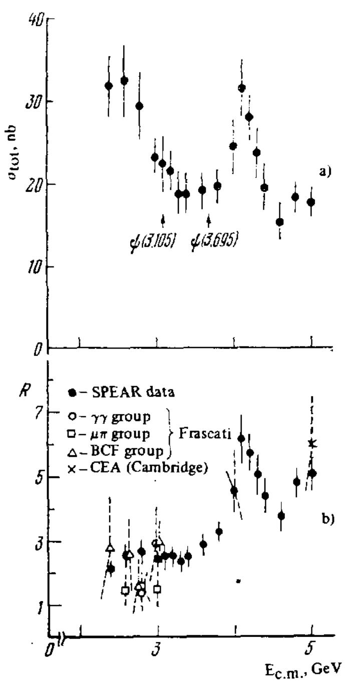


FIG. 4. The cross section for the annihilation process $e^+e^- \rightarrow$ hadrons, $\sigma_{tot}(e^+e^- \rightarrow \text{hadrons})$, as a function of the c.m. energy (a) and the ratio $R = \sigma(e^+e^- \rightarrow \text{hadrons})/\sigma(e^+e^- \rightarrow \mu^+\mu^-)$ as a function of the energy (b).

makes it possible to determine the value of $\Gamma_e\Gamma_f/\Gamma$ and, if the ratio Γ_e/Γ is known (i.e., if the ratio of the $\psi \rightarrow e^+e^-$ decay probability to the total decay probability is known), to determine Γ_e and the full width Γ. Although Eqs. (2.1) and (2.2) in principle provide a solution to the problem of determining the widths on the basis of the experimental data on the cross sections, they require an important correction due to radiative effects.

The radiative effects are associated with the emission of virtual and real photons in the process $e^+e^- \rightarrow \psi$. To first order in $e^2 = \alpha = 1/137$, the contribution of virtual photons to the $e^+e^- \rightarrow \psi$ interaction vertex is described by the diagram of Fig. 5, while the emission of real photons, i.e., the process $e^+e^- \rightarrow \psi + \gamma$, corresponds to the diagram of Fig. 6.

If the total energy E of the e^+ and e^- in the center-of-mass system is somewhat greater than M, there is a preference for the emission of a soft photon with energy $E_\gamma \approx E - M$ by the electron or positron in the annihilation process, since the e^+e^- system is then at the resonance after the emission of the photon. For this reason, there is an appreciable enhancement in all the partial cross sections for e^+e^- annihilation in the region above the resonance, i.e., there is a "radiative tail," which is clearly seen in the curves of Figs. 2 and 3.[4)] The partial cross sections for e^+e^- annihilation $e^+e^- \rightarrow \psi + \gamma \rightarrow f + \gamma$ involving the emission of a soft

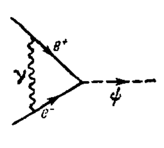


FIG. 5

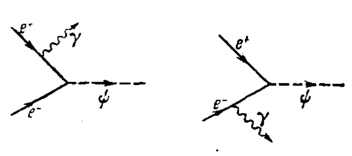


FIG. 6

[4)]An analogous effect should be seen in ψ production in hadronic collisions (or photoproduction) followed by the decay $\psi \rightarrow e^+e^-$. In this case, the photon is emitted by the final e^+e^-: the "radiative tail" appears to the left of the resonance, and its form is determined by (2.4) with the substitution $E - M \rightarrow M - E_{e^+e^-}$.

photon ($E_\gamma \ll M$) in the region above the resonance, $E > M$, can readily be calculated to first order in α on the basis of the diagram of Fig. 6[9,10]. There is obviously no need to allow here for photon emission by the final particles, whose contribution should be small (of order Γ/E_γ with respect to the contribution of the diagrams of Fig. 6 when $E_\gamma \gg \Gamma$). However, it is important to allow for radiative corrections of higher order, since, in addition to α, their contribution is proportional to a product of large logarithms, $\ln(M/m_e) \times \ln[M/(E-M)]$ (where m_e is the mass of the electron). The partial cross section for e^+e^- annihilation into channel f of the ψ decay, with the emission of an undetected soft photon in the region $E > M$, $E - M \gg \Gamma$, $E - M \gg \Delta E$, calculated in the doubly logarithmic approximation (i.e., with allowance for the terms $\sim[\alpha \ln(M/m_e)\ln(M/E_\gamma)]^n$), has the form[9-14]

$$\sigma_f(E) = \frac{4\alpha}{\pi}\int \sigma_f(E')\,dE' \cdot \frac{1}{E-M}\left(\ln\frac{M}{m_e} - \frac{1}{2}\right)\gamma(E - M), \quad (2.4)$$

where $\int \sigma_f(E')dE'$ is determined in (2.2), and

$$\gamma(\varepsilon) = \exp\left[-\frac{4\alpha}{\pi}\left(\ln\frac{M}{m_e} - \frac{1}{2}\right)\ln\frac{M}{2\varepsilon}\right]. \quad (2.5)$$

Terms of order $\alpha\ln(M/m_e)$ are neglected in Eq. (2.4), so that the accuracy of this expression is of order 10% in the cases in which we are interested. A more accurate expression, which takes into account the terms linear in $\alpha\ln(M/m_e)$ and which therefore has an accuracy ~1%, is given in[12-14].

In the case of a narrow resonance, when the experimental resolution is such that $\Delta E \gg \Gamma$, the partial cross sections for annihilation within the peak, averaged with respect to the experimental resolution, with allowance for radiative corrections calculated in the doubly logarithmic approximation, differ from (2.2) by the factor $\gamma(\Delta E)$[9-14]:

$$\sigma_{f\,\mathrm{rad}}^{\mathrm{peak}} = \sigma_f^{\mathrm{peak}}\gamma(\Delta E), \quad (2.6)$$

where σ_f^{peak} is defined by Eq. (2.2), and $\gamma(\Delta E)$ is defined by Eq. (2.5). Substituting the numerical values $M = 3.1$ GeV and $\Delta E = 2.8$ MeV in (2.5), we find $\gamma(\Delta E) \approx 0.5$, i.e., the radiative corrections are very important.

We quote another frequently used formula, which expresses the integral of the experimentally measured cross section over the region of the resonance and the "radiative tail," with allowance for both the experimental resolution and the emission of undetected soft photons, as a function of the upper limit of integration[11]:

$$A(E, E_{\min}) = \int_{E_{\min}}^{E} dE'' \int_{-\infty}^{\infty} dE'\sigma_{f\,\mathrm{rad}}(E')\,w(E'', E') \quad (2.7)$$

$$= \int \sigma_f(E')\,dE' \cdot \gamma(E - M).$$

The same assumptions are used in deriving (2.7) as in (2.4).

It follows from (2.4), (2.6) and (2.7) that the resonance widths can be determined from the experimental data on the cross sections by considering both the region of the peak and the region of the radiative tail. The best procedure seems to be to compare the experimental data with the general theoretical formula which takes the radiative effects into account and which is applicable in both the resonance region and the region of the radiative tail. However, such a formula depends on the specific form of the experimental resolution and can be used only for numerical calculations.

The values of the widths of the resonances ψ and ψ' obtained by analyzing the experimental data (subtracting the background and allowing for radiative effects) are listed in Table I[3-5]. To complete the picture, this table also includes the data on the resonance ψ (4.15).

In calculating Γ_{ee} and Γ in Table I (with the exception of Γ for the ψ (4.15)), the spins of the resonances were taken as 1 (see below) and $\sigma_h(E)$ was determined experimentally by the number of events with hadron tracks. In the case of the ψ resonance, B_{ee} (or $B_{\mu\mu}$) was determined experimentally as the ratio of the number of events involving e^+e^- (or $\mu^+\mu^-$) production to the total number of events in the resonance region, where hadronic events were detected only if the hadrons included charged hadrons. (It is natural to assume that decays of the ψ or ψ' into only neutral hadrons have a small probability.) It was assumed in determining B_{ee} and Γ that there are no unobserved neutral types of ψ decays such as $\psi \to n\gamma$, $\psi \to \gamma\pi^0$ or $\psi \to \bar{\nu}\nu$. The experimental bounds $\Gamma(\psi \to 2\gamma)/\Gamma_{ee} < 0.05$ and $\Gamma(\psi \to \pi^0\gamma)/\Gamma_{ee} < 0.13$ (90% confidence) have been obtained on the relative decay probabilities[7b].

Using the data of a neutrino experiment, the estimate $\Gamma_{\bar{\nu}\nu} < (10^{-2}-10^{-3})\Gamma_{ee}$ can be obtained for the $\psi \to \bar{\nu}\nu$ decay probability (see Sec. 1 of Chap. V).

2. The Spin and Parity of the ψ Meson

The currently available experimental data indicate that the spin and parity of the ψ meson are $J^P = 1^-$ and that parity is conserved in the decays $\psi \to e^+e^-$ and $\psi \to \mu^+\mu^-$ (and apparently also in the decays $\psi \to$ hadrons).

The conclusions about the spin and parity of the ψ meson follow from measurements of the angular distributions of muons and electrons in the reactions $e^+e^- \to \mu^+\mu^-$ and $e^+e^- \to e^+e^-$ in the resonance region and measurements of the total cross section for the reaction $e^+e^- \to \mu^+\mu^-$ in the region of interference between the resonance and background.

In the resonance region, the process $e^+e^- \to \mu^+\mu^-$ is described by the diagram of Fig. 7. In the case of zero spin of the ψ meson, the angular distribution of the emitted muons is isotropic:

$$J=0, \quad \frac{d\sigma}{d\Omega} \sim \text{const.} \tag{2.8}$$

If the spin of the ψ meson is 1 and if parity is conserved in the ψe^+e^- or $\psi\mu^+\mu^-$ interaction, then the angular distribution has the following form:

$$J=1,\ P \text{ is conserved},\ d\sigma/d\Omega \sim 1+\cos^2\theta, \tag{2.9}$$

where θ is the angle between the e^+ and μ^+ momenta. If parity is not conserved (but CP is conserved), we have, instead of (2.9):

$$J=1,\ P \text{ is not conserved},$$
$$\frac{d\sigma}{d\Omega} \sim (1+\cos^2\theta) + \frac{8\lambda^2}{(1+\lambda^2)^2}\cos\theta, \tag{2.10}$$

where, in the case of $\mu - e$ universality, $\lambda = g_A/g_V$ is the ratio of the axial-vector and vector coupling constants for the interaction $\psi_\nu \bar{e}\gamma_\nu(1+\lambda\gamma_5)e$ or $\psi_\nu\bar{\mu}\gamma_\nu(1+\lambda\gamma_5)\mu$. (In the absence of $\mu - e$ universality $\lambda^2 = g_A(e)g_A(\mu)/g_V(e)g_V(\mu)$.) In particular, for the two-component interaction, $\lambda = 1$ and

$$J=1, \quad V-A, \quad \frac{d\sigma}{d\Omega} \sim (1+\cos\theta)^2. \tag{2.11}$$

We see from (2.10) and (2.11) that nonconservation of parity produces a forward-backward asymmetry. Figure 8 shows the experimental data on the angular distribution in the reaction $e^+e^- \to \mu^+\mu^-$ in the region of the ψ meson. These data force us to reject spin 0 for the ψ meson (as well as spin 2, for which the angular distribution has the form $d\sigma/d\Omega \sim 1 - 3\cos^2\theta + 4\cos^4\theta$, and higher spins). It follows from Fig. 8 that the experimental data are in satisfactory agreement with Eq. (2.9), i.e., with the hypothesis that the ψ has spin 1 and parity is conserved. These same data can be used to obtain the bound $\lambda^2 \leq 0.05$ on the possible magnitude of parity nonconservation.

By studying the total cross section for the reaction $e^+e^- \to \mu^+\mu^-$ in the region of interference between the resonance and background below the resonance, it is possible to determine not only the spin of the resonance, but also its parity. The background annihilation process $e^+e^- \to \mu^+\mu^-$ is described by the diagram of Fig. 9, which involves the exchange of a virtual photon. The interference between the diagrams of Figs. 7 and 9 is obviously non-zero only in the case of a vector ψ meson (or when the state with $J^P = 1^-$ contributes strongly to the interactions of the ψ meson with both the electron and the muon). Consequently, the observation of an appreciable interference between the resonance and background, in conjunction with the above-mentioned fact that parity is conserved in the decays $\psi \to e^+e^-$ and $\psi \to \mu^+\mu^-$, would constitute a proof that the ψ is a vector meson[5)]. By summing the diagrams of Figs. 7 and 9, taking into account both resonance and background effects and assuming $\mu - e$ universality and parity conservation, but omitting for simplicity the radiative corrections, which are actually very important in the interference region, the following expression is obtained for the total cross section of the process $e^+e^- \to \mu^+\mu^-$[13]:

TABLE I

Resonance	M, MeV	$\int\sigma_h(E)\,dE$, nb·MeV	Γ_{ee}, keV	B_{ee}	Γ
ψ	3095 ± 5	9900 ± 1500	4.8 ± 0.6	0.069 ± 0.009	69 ± 15 keV
ψ'	3684 ± 5	3700 ± 600	2.2 ± 0.3	$(0.97\pm0.16)\cdot10^{-2}$	225 ± 56 keV
ψ (4.15)	~ 4150	~ 5650	4.2 ± 0.6	$1.7\cdot10^{-5}$	250 ± 50 MeV

The integral $\int\sigma_h(E)dE$ over the resonance region (without radiative corrections) is determined by Eq. (2.2) with f = h, Γ_{ee} is the $\psi \to e^+e^-$ decay width, and $B_{ee} = \Gamma_{ee}/\Gamma$. For the ψ meson, $\Gamma_{\mu\mu}/\Gamma_{ee} = 1.00 \pm 0.05$; for the ψ' meson, $\Gamma_{\mu\mu}/\Gamma_{ee} = 0.89 \pm 0.16$. The mass difference between the ψ and ψ' is determined with better accuracy than the masses themselves: $M_{\psi'} - M_\psi = 590 \pm 1$ MeV.

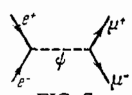


FIG. 7

5)The remaining possibility here that parity is not conserved only in the interaction of the ψ with electrons is ruled out by the experimental data on the e^+e^- angular distribution in the reaction $e^+e^- \to e^+e^-$, which are in agreement with the theoretical result for $J^P = 1^-$ in the resonance region if allowance is made for the contribution from the diagram of Fig. 10.

$$\sigma_{e^+e^-\to\mu^+\mu^-}(E)=\frac{3\pi}{M^2}\left[\frac{\Gamma_{ee}^2}{(E-M)^2+(\Gamma^2/4)}+\frac{4\alpha}{3}\,\frac{\Gamma_{ee}(E-M)}{(E-M)^2+(\Gamma^2/4)}+\frac{4\alpha^2}{9}\right]. \quad (2.12)$$

The expression (2.12) has a minimum at

$$E-M\approx-\frac{3\Gamma_{ee}}{2\alpha}\approx-1\ \mathrm{MeV}. \quad (2.13)$$

We note that there is no such minimum below the resonance (for $J_\psi = 0$ or 1) in the case of the process $e^+e^- \to e^+e^-$ and that the interference term is positive here; this is so because, in addition to the diagrams analogous to Figs. 7 and 9, there is a contribution to this process from the diagram of Fig. 10 involving the exchange of a space-like photon, whose interference with the diagram of Fig. 7 for $E - M < 0$ in the angular region studied in the SPEAR experiments is positive and dominant over the interference term in the total cross section from the diagrams of Figs. 7 and 9.

The experimental values and theoretical curves for the ratio of cross sections $\sigma(e^+e^- \to \mu^+\mu^-)/\sigma(e^+e^- \to e^+e^-)$ in the interference region, shown in Fig. 11, definitely favor a destructive interference in the region below the resonance and thus demonstrate that the spin and parity of the ψ meson are $J^P = 1^-$, i.e., the same as for the photon; consequently, the charge parity of the ψ meson is negative.

We note that the foregoing conclusions about the P and C parities of the ψ meson refer to its interactions with leptons. These conclusions might not apply to the

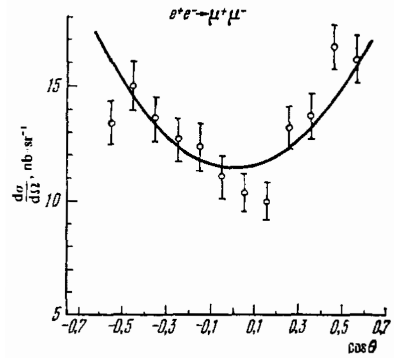


FIG. 8. The angular distribution of muons in the reaction $e^+e^- \to \mu^+\mu^-$.

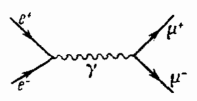


FIG. 9

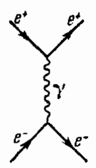


FIG. 10

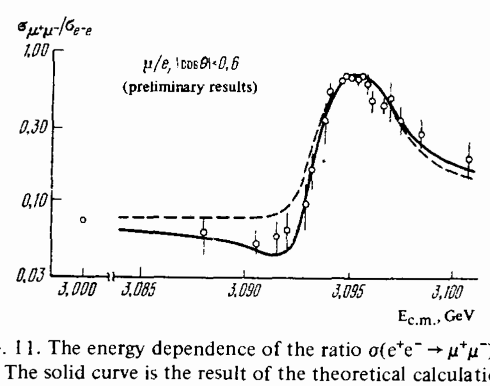


FIG. 11. The energy dependence of the ratio $\sigma(e^+e^- \to \mu^+\mu^-)/\sigma(e^+e^- \to e^+e^-)$. The solid curve is the result of the theoretical calculation assuming that the resonance interferes with the background, and the dashed curve is the same with no interference.

interactions of the ψ with hadrons if these interactions did not conserve the P and C parities. Since the P and C parities are conserved in the interactions of the ψ with leptons, it is natural to assume that this is also the case for the interactions of the ψ with hadrons. We shall henceforth adopt this hypothesis. (Certain additional arguments in favor of C-parity conservation in decays of the ψ into hadrons will be given in the next section.)

3. Non-leptonic decays of the ψ mesons

The study of non-leptonic decays of the ψ meson is of interest in several respects: a) by studying hadronic decays of the ψ meson, it is possible to determine the isotopic spin T and G-parity of the produced hadronic states, i.e., if the ψ meson is a hadron and if isospin is conserved in hadronic decays of the ψ, its T and G can be determined; b) searches for the decays $\psi \to \gamma$ + hadrons enable us to test various hypotheses about the nature of the ψ meson; c) if the ψ meson is a hadron, we can hope to find new hadronic states among the products of the decay $\psi \to \gamma$ + hadrons; d) it is important to study the products of the decay $\psi \to$ hadrons (the spectrum of states, multiplicity, etc.) in order to understand the nature of the ψ meson itself, as well as other processes at high energies (the annihilation processes $\bar{p}p$, $e^+e^- \to$ hadrons, etc.).

Let us consider the decays $\psi \to$ hadrons. We shall suppose that C-parity is conserved in ψ decays. We shall assume at first that the hadrons are produced in a state having a definite value of the isotopic spin $T = 0$ or 1, without considering larger values of T. Since the C-parity of the ψ meson is negative, for $T = 0$ or 1 the G-parity of the produced hadrons is -1 or $+1$, respectively. By virtue of the fact that the G-parity of the pion is negative, this implies that the decays $\psi \to$ pions are allowed into an odd number of pions for $T = 0$ and into an even number of pions for $T = 1$. Let us suppose that the ψ meson is a hadron, characterized by a definite value of the isotopic spin $T = 0$ or $T = 1$, and that the decays $\psi \to$ hadrons occur via the strong (or medium-strong) interaction, which conserves isotopic spin. Even in this case, however, we cannot say that the hadrons in the decay of the ψ are produced in a pure state with definite T. The occurrence of the decay $\psi \to e^+e^-$, which in this case must proceed via a virtual photon, implies that a certain fraction of the ψ decays into hadrons must also proceed via a virtual photon, i.e., that there may be a mixture of states with $T = 0$ and $T = 1$. The partial width for the decay $\psi \to H$ via a virtual photon (where H is some hadronic state) is given by

$$\Gamma_{\psi\to\gamma\to H}=\Gamma_{ee}R_H, \quad (2.14)$$

where R_H is the ratio of the cross section for the annihilation $e^+e^- \to H$ to the cross section for the annihilation $e^+e^- \to \mu^+\mu^-$ outside the resonance:

$$R_H=\frac{\sigma(e^+e^-\to H)}{\sigma(e^+e^-\to\mu^+\mu^-)}. \quad (2.15)$$

Thus, in the case which we are considering here, the fraction of hadrons having isotopic spin different from that of the ψ is uniquely determined by the value of R_H, which can be measured independently. In other words, the ratio of the cross section for e^+e^- annihilation into hadrons having isospin different from that of the ψ to the cross section for e^+e^- annihilation into muons must be the same in the resonance and in the background. The isotopic spin of the ψ meson can be deduced from the experimental data on the cross sections for e^+e^- annihilation into $2\pi^+2\pi^-$, $2\pi^+2\pi^-\pi^0$, $3\pi^+3\pi^-$ and $3\pi^+3\pi^-\pi^0$. According to these data[5],

$$\frac{[\sigma(2\pi^+2\pi^-)/\sigma(\mu^+\mu^-)]_{\rm res}}{[\sigma(2\pi^+2\pi^-)/\sigma(\mu^+\mu^-)]_{\rm back}}=0.82\pm0.22,$$

$$\frac{[\sigma(3\pi^+3\pi^-)/\sigma(\mu^+\mu^-)]_{\rm res}}{[\sigma(3\pi^+3\pi^-)/\sigma(\mu^+\mu^-)]_{\rm back}}=1.10\pm0.54,$$

whereas the analogous ratios for $2\pi^+2\pi^-\pi^0$ and $3\pi^+3\pi^-\pi^0$ production are greater than 7. (We note that, since $R=\sigma(e^+e^-\to \text{hadrons})/\sigma(e^+e^-\to\mu^+\mu^-)\approx 2.5$ outside the resonance, the ratio of the probability for ψ decay into hadrons via a virtual photon to the total probability for the decays $\psi\to$ hadrons is $B_{ee}R\approx 0.068\times 2.5=0.17$.) These facts together imply that the isotopic spin of the ψ is zero and that its G-parity is negative[6]. These same facts also favor the conservation of C-parity in the decay of the ψ into hadrons. In fact, if C-parity were not conserved in the decay of the ψ into hadrons, there would be no reason for all the decays $\psi\to 2\pi^+2\pi^-$ to proceed via a virtual photon.

The fact that the isotopic spin of the ψ boson is zero has a number of consequences for hadronic decays. In particular, the decay $\psi\to\pi^+\pi^-$ can proceed only via a virtual photon, and the width of this decay satisfies the relation

$$\frac{\Gamma_{\pi^+\pi^-}}{\Gamma_{ee}}=\frac{1}{4}|F_\pi(M^2)|^2\approx 10^{-3}, \qquad (2.16)$$

where $F_\pi(q^2)$ is the electric form factor of the pion, with $|F_\pi(3.1^2\ \text{GeV}^2)|^2\approx 4\times 10^{-3}$ [15]. Unlike the $\pi^+\pi^-$ system, the K^+K^- system with $J^P=1^-$ can exist in a state with $T=0$, so that (since $|F_\pi|\approx|F_K|$[15]) $\Gamma_{K^+K^-}$ must be much greater than $\Gamma_{\pi^+\pi^-}$. However, it must be borne in mind that the direct decay $\psi\to K^+K^-$ (not via a virtual photon) is also forbidden in exact SU(3) symmetry if the ψ is an SU(3) singlet. Therefore measurements of the $\psi\to\pi^+\pi^-$ and $\psi\to K^+K^-$ decay probabilities may provide a test of SU(3) symmetry in the decays of the ψ mesons. At the present time, there are only the experimental upper limits $\Gamma_{\pi^+\pi^-}/\Gamma_{ee}<2.5\times10^{-2}$ and $\Gamma_{K^+K^-}/\Gamma_{ee}<2.5\times10^{-2}$ on $\Gamma_{\pi^+\pi^-}$ and $\Gamma_{K^+K^-}$.

Similarly, the width of the decay $\psi\to\bar{p}p$ must be much greater than the width obtained under the assumption that the decay occurs via a virtual photon. The latter is equal to

$$\frac{(\Gamma_{pp})_{\text{virt}\,\gamma}}{\Gamma_{ee}}=v_p[|G_M(M^2)|+b|G_E(M^2)|^2], \qquad (2.17)$$

where $G_E(q^2)$ and $G_M(q^2)$ are the electric and magnetic form factors of the proton, v_p is its velocity, and $b\approx 0.2$ is a kinematic factor. With $|G_M|=|G_E|=0.12$ (an extrapolation according to the $1/q^2$ law from the experimentally measured value $|G_M|=|G_E|=0.27$ at $q^2=4.4\ \text{GeV}^2$; see[16]), we have

$$\frac{\Gamma(\psi\to\gamma\to\bar{p}p)}{\Gamma_{ee}}=1.5\cdot10^{-2}. \qquad (2.17')$$

[6]Strictly speaking, this is true only if the ψ meson is a hadron and interacts with leptons via a virtual photon. If the ψ meson were not a hadron and had a direct decay interaction with leptons and hadrons, the formulation given in the text would have to be replaced by the following (see Sec. 1 of Chap. V): the main contribution to the direct decay interaction $\psi\to$ hadrons comes from hadronic states with $T=0$ and $G=-1$. The conclusions which we have indicated for $T=0$ remain valid here.

If the form factor of the proton falls off in the time-like region in the same way as in the space-like region, i.e., according to the q^{-4} law, we obtain instead of (2.17′)

$$\frac{\Gamma(\psi\to\gamma\to\bar{p}p)}{\Gamma_{ee}}\approx 0.3\cdot10^{-2}. \qquad (2.17'')$$

The values (2.17′) and (2.17″) are to be compared with the experimental value $\Gamma_{pp}/\Gamma_{ee}\approx 3\times10^{-2}$.

As we have already mentioned, it would be of great interest to search for the decays $\psi\to\gamma+$ hadrons. If the ψ meson is a colored particle and if color is violated by the electromagnetic interaction (see Chap. IV), these would be the principal decay modes of the ψ which involve hadron emission. The hadrons should be produced in this case in a state with $T=0$ and $G=+1$, i.e., the number of pions in the decays $\psi\to n\pi+\gamma$ must be even. Experimental information on this question is obtained by studying the missing-mass spectrum in the decay $\psi\to 2\pi^+2\pi^-+X$. The spectrum of m_X in the region $m_X\approx m_{\pi^0}$ exhibits a large peak which is displaced to the right of the point $m_X=0$, but the experimental resolution in the mass is such that we cannot be sure that this decay does not involve the emission of photons in large numbers. The statement that the peak at $m_X\approx 0-m_{\pi^0}$ is due to the pion and not the photon is also supported by the fact that the $\pi^+\pi^-\pi^0$ mass spectrum (with the π^0 in the region of the peak) exhibits a pronounced maximum at the mass of the ω meson[5].

At the present time, there exists only an experimental upper bound on the probability of decays $\psi\to\gamma\eta_c$ followed by the decay $\eta_c\to2\gamma$ (where η_c is a pseudoscalar meson with hidden charm; see Chap. IV)[7b]:

$$\frac{\Gamma(\psi\to\gamma\eta_c)}{\Gamma_{ee}}\frac{\Gamma(\eta_c\to2\gamma)}{\Gamma(\eta_c\to\text{anything})}<0.2\ (90\%\ \text{confidence}) \qquad (2.18)$$

4. Decays of the ψ' Meson

An important experimental fact regarding the decays of the ψ' meson is the observation of the decay $\psi'\to\psi\pi^+\pi^-$ with a probability[4,5]

$$\frac{\Gamma(\psi'\to\psi\pi^+\pi^-)}{\Gamma(\psi'\to\text{anything})}=0.31\pm0.04. \qquad (2.19)$$

It has also been found experimentally that

$$\frac{\Gamma(\psi'\to\psi+\text{anything})}{(\Gamma\psi'\to\text{anything})}=0.54\pm0.08, \qquad (2.19')$$

and hence

$$\frac{\Gamma(\psi'\to\psi+\text{anything})}{\Gamma(\psi'\to\psi\pi^+\pi^-)}=1.8\pm0.1. \qquad (2.19'')$$

Definite conclusions about the quantum numbers of the ψ' meson can be drawn from these data. The two-pion system produced in the decay $\psi'\to\psi2\pi$ can exist in states with isotopic spin $T=0,1,2$. For these three values, the corresponding values of the ratio of the sum of the $\psi'\to\psi\pi^+\pi^-$ and $\psi'\to\psi\pi^0\pi^0$ decay probabilities to the $\psi'\to\psi\pi^+\pi^-$ decay probability are given by

$$\frac{\Gamma(\psi'\to\psi\pi^+\pi^-)+\Gamma(\psi'\to\psi\pi^0\pi^0)}{\Gamma(\psi'\to\psi\pi^+\pi^-)}=1.5;\ 1.0;\ 3.0. \qquad (2.20)$$

By comparing (2.20) with (2.19″), we can infer that the pure two-pion state with $T=2$ is excluded by the experimental data.

Now the two-pion state with $T=1$ is C-odd. If we assume that the ψ' is a hadron (arguments in favor of this hypothesis will be given later) and is therefore produced in e^+e^- collisions via a virtual photon, then $J^P_{\psi'} = 1^-$ and the C-parity of the ψ' is negative. Since the C-parity of the ψ is also negative, conservation of C-parity in the decay $\psi' \to \psi 2\pi$ implies that the decay $\psi' \to \psi 2\pi$ with pion production in the state with $T = 1$ (i.e., in a p-wave state) is forbidden. Thus, only the two-pion state with $T = 0$ (of the pure states) is consistent with the experimental value (2.19). It follows then from Eqs. (2.19″) and (2.20) that the decays $\psi' \to \psi$ + anything go predominantly into the channel $\psi' \to \psi\pi\pi$. Assuming that isotopic spin is conserved in the decay $\psi' \to \psi\pi\pi$ and taking $T_\psi = 0$ in accordance with our earlier discussion, we conclude that $T_{\psi'} = 0$, i.e., that all the quantum numbers J, P, C and T of the ψ' meson are the same as for the ψ.

The absolute value of the partial width for the decay $\psi' \to \psi\pi^+\pi^-$ can be determined from the data of Table I and the ratio (2.19):

$$\Gamma(\psi' \to \psi\pi^+\pi^-) = 70 \pm 20 \text{ keV}. \qquad (2.21)$$

We shall use (2.21) to estimate the effective strength of the interaction in the decay $\psi' \to \psi\pi\pi$. Let us consider the simplest model, in which the effective interaction Hamiltonian governing this decay has the form

$$H = \sqrt{4\pi}\,\lambda\psi'_\mu\psi_\mu\pi^+\pi^-, \qquad (2.22)$$

where λ is a coupling constant. In Eq. (2.22) we are taking into account only the s-wave contribution in the two-pion system and neglecting the $\pi\pi$ interaction in the final state. The d wave in the $\pi\pi$ system may be expected to give a small contribution to the amplitude. Estimates of the effect of the $\pi\pi$ interaction will be made later. The interaction Hamiltonian (2.22) corresponds to a $\psi' \to \psi\pi^+\pi^-$ decay width

$$\Gamma(\psi' \to \psi\pi^+\pi^-) = \frac{\lambda^2}{48\pi^2}\frac{\Delta^3}{MM'}f, \qquad (2.23)$$

where M′ is the ψ' mass, $\Delta = M' - M$, and $f \approx 1/3$ is a correction factor which takes into account the finite value of the pion mass. Comparing (2.23) with (2.21), we find

$$\lambda^2 \approx 6, \qquad (2.24)$$

i.e., the decay $\psi' \to \psi\pi\pi$ is governed by a strong interaction. This fact is of paramount importance for the theoretical description of the ψ mesons.

It should be noted that, although the values of λ^2 were found to be rather large, their corresponding effective interaction is much weaker than the ordinary strong interactions, i.e., it is, as it were, a medium-strong interaction. To see this, let us use the Hamiltonian (2.22) to make a perturbation-theory calculation of the cross section for the process $\pi^+ + \psi' \to \pi^+ + \psi$ at relatively low energies $E_\pi \lesssim 1$ GeV in the ψ' rest system. (In this approximation, scattering of the π by the ψ' occurs in the s wave.) We find

$$\sigma(\pi^+ + \psi' \to \pi^+ + \psi) = \frac{\lambda^2}{4}\frac{1}{MM'}. \qquad (2.25)$$

We see from Eq. (2.25) that, with λ^2 determined by (2.24), $\sigma(\pi^+ + \psi' \to \pi^+ + \psi)$ is small in comparison with the ordinary hadronic cross sections $\sigma \sim 1/m_\pi^2$, and for $E_\pi < 1$ GeV is also small in comparison with the unitarity limit $\sigma_{unit} = 4\pi/E_\pi^2$ for the s wave.

Let us now estimate the strength of the $\psi'\psi\pi\pi$ interaction in another limiting case, assuming that the decay $\psi' \to \psi\pi\pi$ occurs via a scalar σ meson with mass $m_\sigma \approx 550$ MeV and $T = 0$, according to the scheme $\psi' \to \psi\sigma$, $\sigma \to \pi\pi$. (This possibility is not excluded by the pion spectrum in the decay $\psi' \to \psi\pi^+\pi^-$ [5] (Fig. 12).)

On the basis of the data on the pion spectrum in this decay, we are assuming that the σ is an ordinary strongly interacting meson with a $\sigma \to \pi\pi$ decay width of order 100—200 MeV. Writing the $\psi' \to \psi\sigma$ decay interaction in the form

$$H = \sqrt{4\pi}\,gm_\sigma\psi'_\mu\psi_\mu\sigma, \qquad (2.26)$$

we find for this case

$$\Gamma(\psi' \to \psi\sigma) = \frac{g^2}{2}\sqrt{2m_\sigma(\Delta - m_\sigma)}\frac{m_\sigma^2}{MM'}, \qquad (2.27)$$

so that the width (2.21) corresponds to a coupling constant

$$g^2 = 3\cdot 10^{-2}. \qquad (2.28)$$

The cross section for the scattering process $\pi^+\psi' \to \pi^+\psi$ is described in this model by the diagram involving exchange of the σ meson and at low pion energies $E_\pi \le m_\sigma$ is given by

$$\sigma(\pi^+\psi' \to \pi^+\psi) \approx \frac{8\pi}{3}g^2\frac{\Gamma_\sigma}{m_\sigma}\frac{1}{M'M}, \qquad (2.29)$$

where Γ_σ is the width of the σ meson. For g^2 given by (2.28) and $\Gamma_\sigma/m_\sigma \lesssim 1/2$, the value (2.29) is approximately an order of magnitude smaller than that given by Eq. (2.25). Thus, in either model, we conclude that the decay $\psi' \to \psi\pi\pi$ is governed by a strong interaction, but that this interaction is much weaker (by 2—3 orders of magnitude in the cross section) than the ordinary strong interactions.

At the present time, there is no unique interpretation of the relative weakness of the $\pi\pi\psi'\psi$ interaction: it may be due to a generally medium-strong interaction of the ψ with the ordinary hadrons, a specifically weak nondiagonal interaction $\psi'\pi \to \psi\pi$[7)], or finally a weakness of only the interaction of the π and ψ at low energies. This last possibility may, for example, be a consequence of Adler's PCAC condition (the hypothesis of partial conservation of the axial-vector current). In this case, the small value of the matrix element is determined by the parameter $(p_{\pi^+}p_{\pi^-})/m_{char}^2 \sim 1/5$, and it is possible that it is by itself insufficient to explain the weakness of the $\pi\pi\psi'\psi$ interaction. Moreover, in discussing the applicability of the PCAC hypothesis to the decay $\psi' \to \psi\pi\pi$, it must be borne in mind that the expression for

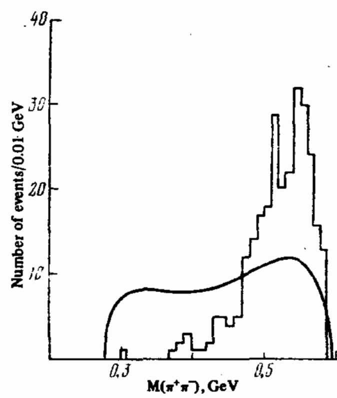


FIG. 12. The spectrum in the effective mass $M(\pi^+\pi^-)$ of the two-pion system in the decay $\psi' \to \psi\pi^+\pi^-$ (see [5]).

[7)]This weakness might, for example, be due to a suppression of the decay $\psi' \to \psi\pi\pi$ by some selection rule (conservation of isospin if $T_{\psi'} = 1$).

the matrix element may contain the invariant $p_\psi p_\pi$, which tends to zero as $p_\pi \to 0$ and whose ratio to m^2_{char} is of order 1 in the real decay $\psi' \to \psi\pi\pi$. The spectrum in $M(\pi^+\pi^-)$ in the decay $\psi' \to \psi\pi\pi$, was calculated in[17] on the basis of the PCAC hypothesis.

Since the decay $\psi' \to \psi$ + anything occurs with 50% probability while the probability for the decay $\psi \to \mu^+\mu^-$ or $\psi \to e^+e^-$ is about 7%, approximately 3.5% of the events involving muons (or electrons) which are seeen in the decays of the ψ' arise from the casdade process

$$\psi' \to \psi + \text{anything}, \quad \psi \to \mu^+\mu^- \ (e^+e^-),$$

and 1.5% of the events arise from decays of the type

$$\psi' \to \psi + 2\pi^0, \quad \psi \to \mu^+\mu^-,$$

which is much greater than the fraction of muons (or electrons) from the direct decay $\psi' \to \mu^+\mu^-$ (e^+e^-). This circumstance hinders the experimental determination of the probability of the direct decays $\psi' \to \mu^+\mu^-$ (e^+e^-) and hence the full width of the ψ', since the experimental arrangement of SPEAR does not span the entire solid angle, π^0 mesons are not detected, and the energy of the μ or e is determined with poor accuracy.

III. EXPERIMENTS ON THE PRODUCTION OF NEW PARTICLES IN BEAMS OF HADRONS, PHOTONS AND NEUTRINOS

1. The discovery of a narrow resonance in the reaction $pBe \to e^+e^-X$

At the same time that the ψ mesons were observed in e^+e^- collisions, a narrow resonance in the e^+e^- system was discovered by Ting's group at Brookhaven[1]. They studied the inclusive production of electron-positron pairs in collisions of protons with beryllium nuclei:

$$p\,\text{Be} \to e^+e^-X \tag{3.1}$$

(where X is an arbitrary hadronic state) at an initial proton energy 28.5 GeV. The distribution which was found in the effective mass $m_{e^+e^-}$ of the electron-positron pair is shown in Fig. 13. The experimental resolution was 25 MeV, and the observed form of the peak is consistent with a resonance width which is negligibly small in comparison with the experimental resolution. We note that some of the events to the left of the resonance, i.e., with smaller mass, are due to radiative decay of the resonance into e^+e^- and a photon (see the footnote in Sec. 1 of Chap. II).

The experimenters who discovered this resonance designated it the J particle. If the decay $J \to e^+e^-$ is due to the electromagnetic interaction, the spin and parity of the resonance are given by $J^P = 1^-$. We shall henceforth identify the J with the ψ (3.1).

The total cross section for producing ψ mesons was not measured in the experiment[1], since only particles in a very limited momentum range were detected. To estimate it, the authors of[1] assumed that the cross section for ψ production has the same dependence on the longitudinal and transverse momenta as the cross section for ρ-meson production:

$$E\frac{d^3\sigma}{dp^3} \sim \exp(-6p_\perp), \tag{3.2}$$

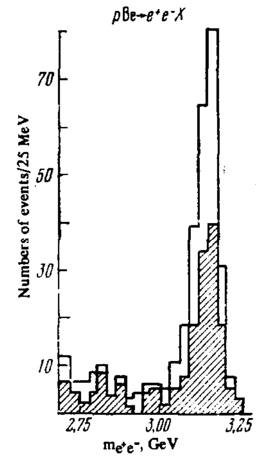


FIG. 13. The distribution in the invariant mass of the e^+e^- pair in the reaction $pBe \to e^+e^-X$ according to the data of [1]. The crosshatched and non-crosshatched events refer to somewhat different operational regimes of the accelerator.

where $p_\perp$ is the transverse momentum of the ψ meson in GeV, and E is the energy of the resonance in the center-of-mass system. The product of the cross section for producing the ψ and the probability of its decay into the e^+e^- channel then has the value

$$\sigma(pN \to \psi X)\frac{\Gamma(\psi \to e^+e^-)}{\Gamma(\psi \to \text{anything})} \sim 10^{-34}\,\text{cm}^2, E_p = 28.5 \text{ GeV}. \tag{3.3}$$

Resonances of large mass, and in particular the ψ' and ψ (4.15), were not observed at the level of the yield of e^+e^- pairs equal to 1% of the yield of pairs of mass m_ψ. Actually, the ψ (4.15) resonance should not be seen in such an experiment, since the probability of its decay into e^+e^- is very small, $\sim 10^{-5}$. For the ψ', the relatively small value of the partial width for the decay into e^+e^- can only partially account for the negative result of the search for this resonance in the experiment (Γ_{ee}/Γ_{tot} is smaller for the ψ' than for the ψ by a factor ~ 10[2]). The additional suppression of the cross section for ψ' production by about an order of magnitude is often explained by the thermodynamic model and the large mass of the ψ' in comparison with the mass of the ψ. However, it is not clear to what extent this model can be applied to the production of ψ mesons (see the remarks at the end of Sec. 2b of Chap. VI). It is possible that the experiment[1] is exhibiting threshold effects for the production of heavy particles.

2. The production of ψ mesons in nucleon-nuclear and nucleon nucleon collisions

More recently, the production of ψ mesons in hadronic collisions has also been observed at Fermilab (Batavia, U.S.A.) and at CERN, using the colliding-beam accelerator. The accelerator at Batavia was used to study the reaction[3]

$$n\,\text{Be} \to \mu^+\mu^-X \tag{3.4}$$

at an initial neutron energy $E_n = 250 \pm 50$ GeV. Figure 14 shows the distribution of events in the mass $m_{\mu\mu}$ of the muon pair in the region $m_{\mu\mu} > 1.4$ GeV. The peak in the mass region ~ 3 GeV includes 43 events.

The existing experimental arrangement can be used to study the production cross section as a function of the transverse momentum $p_\perp$ and the longitudinal momentum $p_\parallel$ of the ψ mesons (the scaling variable $x = p_\parallel/p_N$ is frequently adopted instead of $p_\parallel$). The results indicate, in particular, that the distribution in the transverse momentum is much broader than in the case of ρ-meson production (we note that, apart from dynami-

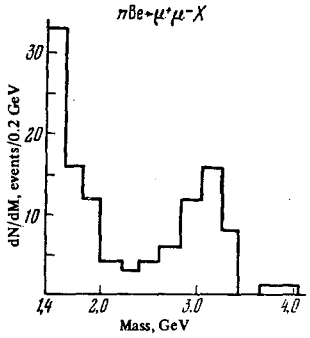


FIG. 14. The distribution in the $\mu^+\mu^-$ invariant mass in the reaction $nBe \to \mu^+\mu^- X$ according to the data of [3].

cal reasons associated with the large mass of the ψ mesons, the broader distribution in $p_\perp$ may be due to the cascade process $pBe \to \psi' X \to \psi X'$).

If we assume, in accordance with the data which have been obtained, that the cross section depends on $p_\perp$ and x in the form

$$E\frac{d^3\sigma}{dp^3} \sim e^{-10x}e^{-2p_\perp^2}, \tag{3.5}$$

then

$$\sigma(n\,\mathrm{Be} \to \psi X)\frac{\Gamma(\psi \to \mu^+\mu^-)}{\Gamma(\psi \to \text{anything})} = \begin{cases} 3.6\cdot 10^{-33}\,\mathrm{cm^2/nucleon}, & x > 0.24, \\ 1.7\cdot 10^{-33}\,\mathrm{cm^2/nucleon}, & x > 0.32. \end{cases} \tag{3.6}$$

If the parameters of the distribution in x and $p_\perp$ are varied within the range allowed by the experimental data, the cross section varies by about a factor of two.

Using the CERN intersecting storage rings, the production of ψ mesons has been observed in pp collisions[4]

$$pp \to \psi X, \quad \psi \to e^+e^-, \tag{3.7}$$

the initial energy of each of the beams being ~30 GeV.

According to preliminary data obtained with statistics of 8 events of the type (3.7), the cross section for producing ψ mesons has a value of order

$$\sigma(pp \to \psi X)\frac{\Gamma(\psi \to e^+e^-)}{\Gamma(\psi \to \text{anything})} \sim 4\cdot 10^{-32}\,\mathrm{cm^2},\ E_p \sim 10^3\,\mathrm{GeV}. \tag{3.8}$$

It is difficult to compare the results of the experiments on ψ-meson production which have been carried out at various initial energies and at different laboratories, since the experimental arrangements have different resolutions in the energy and emission angle of the leptons, and the determination of the cross section always involves a certain extrapolation procedure. However, it is clear that the value of the cross section rises with energy.

3. Experiments on the Photoproduction of ψ Mesons

The group of experimenters who carried out the experiment on ψ-meson production in a neutron beam also observed the photoproduction of ψ mesons on beryllium nuclei[5]:

$$\gamma\,\mathrm{Be} \to \psi X, \quad \psi \to \mu^+\mu^-. \tag{3.9}$$

The initial energy of the photons was greater than 80 GeV and had an average value ~150 GeV.

The experimental yield of ψ mesons was

$$\sigma(\gamma\,\mathrm{Be} \to \psi X)\frac{\Gamma(\psi \to \mu^+\mu^-)}{\Gamma(\psi \to \text{anything})} = (16 \pm 5)\,10^{-33}\ \mathrm{cm^2/nucleus}. \tag{3.10}$$

The conversion of this number into a value of the cross section for the photoproduction of ψ mesons by nucleons requires, firstly, a separation of the mechanisms of diffraction production by the nucleus as a whole and by the individual nucleons and, secondly, allowance for the partial width of the decay.

It was assumed in[5] that the differential cross section has the form

$$\frac{d\sigma}{dt}(\gamma\,\mathrm{Be} \to \psi X) \sim A^2e^{40t} + Ae^{bt}, \tag{3.11}$$

where A is the atomic number, and t is the square of the momentum transfer (in GeV^2). The first and second terms in Eq. (3.11) correspond to scattering by the nucleus as a whole and by the individual nucleons, respectively. The value of the parameter b was not determined precisely, but it was found that the value b = 4 GeV^{-2} provides a good description of the data that were obtained.

Adopting the assumption (3.11) and the slope b = 4 GeV^{-2}, we find from (2.10) a cross section for ψ-meson photoproduction on nucleons having the value

$$\sigma(\gamma N \to \psi X) = (13 \pm 5)\cdot 10^{-33}\,\mathrm{cm^2},\ E_\gamma \sim 150\,\mathrm{GeV}. \tag{3.12}$$

We recall in this connection that the photoproduction cross section is 14 000 nb for ρ mesons, 2000 nb for ω mesons, and 600 nb for φ mesons (1 nb = 10^{-33} cm^2).

The photoproduction of ψ mesons has also been observed at lower energies at Stanford (SLAC, U.S.A.)[6]. It amounts to approximately 10 nb at 21 GeV:

$$\sigma(\gamma N \to \psi X) \approx 10\cdot 10^{-33}\,\mathrm{cm^2},\ E_\gamma = 21\,\mathrm{GeV}. \tag{3.13}$$

It has also been found that the cross section has a strong energy dependence: as the initial energy varies from 14 to 21 GeV, the cross section rises by about a factor 3. The SLAC accelerator was used to measure not only the total cross section, but also the differential cross section for producing ψ mesons at small angles, which was found to have the value[6]

$$\left.\frac{d\sigma(\gamma N \to \psi N)}{dt}\right|_{t=0} = (25 \pm 5)\cdot 10^{-33}\,\mathrm{cm^2/GeV^2}, \tag{3.14}$$

with a slope equal to

$$b = 2.5 - 3.0\ \mathrm{GeV^{-2}}.$$

The first estimates of the cross section for ψ'-meson photoproduction were also obtained:

$$\sigma(\gamma N \to \psi' X) \approx 0.5\sigma(\gamma N \to \psi X). \tag{3.15}$$

However, this estimate was made using very poor statistics of only a few events.

4. Qualitative Effects of the Production of Charmed Particles in Neutrino Reactions

In all recent neutrino experiments, searches have been made for reactions involving the production of charmed particles:

$$\nu_\mu N \to \mu^- h_c X_s \tag{3.16}$$

where h_c is a hadron containing the c quark (see Chap. IV).

In most cases, it is not feasible to make a direct search for a resonance h_c in the effective-mass distribution; usually, only the momentum of the muon and/or the total energy of the hadrons are measured. It is therefore more common to discuss qualitative effects: the production of muon pairs, violation of the $\Delta Q = \Delta S$ rule (where ΔQ and ΔS are the changes of charge and strangeness of the hadrons), and violations of scaling. We shall describe these effects briefly.

a) Muon pair production

$$\nu_\mu N \to \mu^+\mu^- X \tag{3.17}$$

may be due either to the production of ψ mesons or pairs of charmed particles

$$\nu_\mu N \to \psi X, \quad \psi \to \mu^+\mu^-, \qquad \nu_\mu N \to h_c \bar{h}_c X \tag{3.18}$$

by neutral currents or to the production of a charmed particle by charged currents, which subsequently decays into a lepton pair and hadrons:

$$\nu_\mu N \to \mu^- h_c X, \quad h_c \to \mu^+ \bar{\nu} X', \tag{3.19}$$

where X and X′ are systems of hadrons. The production of a μe pair is also possible, but the identification of the leptons in most experimental arrangements involves transmission through a large length of iron, and electrons are not detected. Searches for μe pairs are possible in bubble chambers.

Two groups of experimenters have now reported the observation of muon pairs in the neutrino experiment at Batavia[7,8] (see below). In addition, one event involving the production of a μe pair has been observed in the "Gargamelle" chamber at CERN. The authors estimated the background of ordinary weak interactions to be $\sim 10^{-3}$ events[9].

b) An effective violation of the $\Delta Q = \Delta S$ rule, such as the observation of the reaction

$$\nu_\mu p \to \mu^-\pi^-\pi^+\pi^+\pi^+\Lambda \tag{3.20}$$

(here $\Delta S = -1$, $\Delta Q = +1$), might also serve as circumstantial evidence for the production of charmed particles. In fact, suppose that a charmed quark is produced in the reaction

$$\nu_\mu n \to \mu^- c, \tag{3.21}$$

and undergoes non-leptonic decay

$$c \to p\lambda\bar{n}, \tag{3.22}$$

so that the reaction $\nu_\mu n \to \mu^- p\lambda\bar{n}$ effectively takes place. It is clear that $\Delta Q = -\Delta S$ in this case. At sufficiently high energy, the amplitudes for such processes are not small in comparison with the amplitudes for the single production of strange particles in the ordinary weak interaction, and there could be a large violation of the $\Delta Q = \Delta S$ rule (the validity of this rule in strange-particle decays has been established with an accuracy of order 1%).

One event of the reaction (3.20) has been identified in the neutrino experiment at Brookhaven. The authors estimated the background to be 3×10^{-5} events[10].

c) Scaling in neutrino reactions may be violated as a result of threshold effects associated with the production of new particles. Heavy particles would be produced mainly with large energy transfers ν and small values of Q^2. We would therefore expect an abundance of events with $y \equiv \nu/E_\nu \sim 1$ and $x \equiv Q^2/2m\nu \ll 1$.

In addition, specific effects may be expected, for example, in the distribution with respect to the invariant mass of the hadronic system. Experimental searches for a characteristic violation of scaling in neutrino reactions have also led to positive results[11], which we shall discuss later in greater detail (see Sec. 4f of Chap. VI). We note that, since the entire discussion concerns only the indirect effects of the new short-living particles, we would not be able to reach a definitive conclusion that charmed particles exist even if these effects were observed.

5. Experiments on the Production of Muon Pairs in a Neutrino Beam

As we have already mentioned, the reaction (3.17) involving muon pair production has been observed by two groups of experimenters at Batavia. In the experiment of the HPW (Harvard-Pennsylvania-Wisconsin) group[7], 14 muon pairs were detected, with overall statistics of 8×10^4 neutrino and antineutrino events and an initial energy $E_{\nu,\bar{\nu}} > 40$ GeV. If allowance is made for the efficiency of muon detection, the probability of muon pair production comprises about 1% of the total cross section:

$$\frac{\sigma(\nu_\mu(\bar{\nu}_\mu) N \to \mu^+\mu^- X)}{\sigma(\nu_\mu(\bar{\nu}_\mu) N \to \mu X)} = (9 \pm 3)\cdot 10^{-3}, \tag{3.23}$$

where only the statistical error is given, the systematic error having been estimated by the authors of[7] as a factor 2.

As the distribution in the mass of the produced muons does not have a peak at the mass of the ψ meson, we consider the reaction (3.19) involving the production of charmed particles. The second muon can also be produced as a result of ordinary leptonic decays of pions or kaons. The possibility of separating this mechanism from the reaction involving the production and decay of charmed particles is connected with the fact that charmed particles should have a much shorter lifetime.

The background conditions have been discussed in detail by the California Institute of Technology group[8]. Of 19 observed events which were candidates for muon pair production, it was found that 15 could be attributed to the background. The background for the remaining 4 events (in the neutrino exposure) was estimated as 0.4 events. The magnitude of the observed effect is consistent with the ratio (3.23).

We note that the experimental papers also contain analyses of other possible mechanisms of muon pair production (the production of intermediate bosons or heavy leptons) and demonstrate that the characteristics of the observed events are evidently in agreement only with the hypothesis of the production and leptonic decay of short-lived hadrons.

IV. QUARK MODELS OF THE NEW PARTICLES

The discovery of the ψ mesons has evoked intense interest in quark models of hadrons. In most of the published theoretical papers devoted to the interpretation of the ψ mesons as hadrons, these particles are considered in the framework of a model involving 12 quarks, in which there exist four types of quarks of three different colors. In addition, the existence of particles whose properties are similar to those of the ψ mesons was predicted several years ago on the basis of the quark model.

The experimental situation is rather uncertain at the present time, and it is possible that the most popular hypotheses regarding the quark structure of the ψ mesons will be ruled out experimentally. Nevertheless, it seems appropriate to discuss these hypotheses. This is so for a number of reasons. First, these hypotheses may prove to be correct. Second, if they do not turn out to be correct in themselves, their future modifications may be correct. Third, without these hypotheses it would be difficult to see the point of the program of experiments to search for new particles which is now being carried out using all the world's greatest accelerators. A fourth and final reason is that the quark model now involves a number of beautiful physical ideas, mechanisms and effects. The new mesons provide a good reason for a broad section of the physics community to become acquainted with them.

In what follows, we shall consider the status of the quark model on the eve of the discovery of the ψ mesons. In particular, we shall discuss the arguments for going over from three to twelve quarks. We shall then see how the ψ mesons are described in the framework of the quark model. Finally, we shall discuss the properties of the so-called charmed particles, which are predicted by the 12-quark model and which have so far not been discovered.

1. Quark Models

a) **Three quarks**. The three-quark model reproduces many properties of hadrons which show up in the strong, electromagnetic and weak interactions. With three quarks p, n and λ, having electric charges $2/3$, $-1/3$, and $-1/3$, respectively, baryonic charge $1/3$ and spin $1/2$, it is possible to construct all known SU(3) multiplets of mesons and baryons. In particular, this includes the nonets of pseudoscalar and vector mesons, as well as baryon multiplets: an octet with $J^P = 1/2^+$ and a decuplet with $J^P = 3/2^+$. These multiplets correspond to the lowest orbital states of the quarks—the states with $l = 0$. Many SU(3) multiplets with $l \geq 1$ are also known, such as the meson multiplets with $l = 1$ ($J^P = 2^+, 1^+, 0^+$) or the baryon multiplets with negative parity. However, these multiplets have generally been studied much more poorly: not all of the particles contained in them have been discovered experimentally, and in certain cases the values of the spin and parity of particles that have already been discovered have not been established.

The three-quark model provides a good description of the electromagnetic properties of hadrons: their magnetic moments, electromagnetic mass differences, the amplitudes for electromagnetic decays of mesons, etc. Finally, the model provides an excellent description of a wide class of weak processes. In this model, the weak charged hadronic current has the form

$$j_W = \bar{p}n_\theta = \bar{p}\,(n\cos\theta + \lambda\sin\theta),$$

where $\bar{p}$ is the creation operator of the p quark, n and λ are the destruction operators of the n and λ quarks, and θ is the Cabibbo angle. The angle is found experimentally to be $\theta \approx 15^\circ$.

The interactions of the current $\bar{p}n_\theta$ and of the charged leptonic currents $\bar{\nu}_e e$ and $\bar{\nu}_\mu\mu$ show up in β decay of nuclei, in numerous weak decays of mesons and baryons, and in neutrino-induced reactions such as $\nu + n \to \mu^- + p$.

However, the three-quark model suffers from a number of serious defects, which can be overcome only by increasing the number of quarks. One serious difficulty for the three-quark model is the experimental absence of neutral strangeness-changing currents.

b) **Neutral currents**. First of all, the three-quark model cannot provide a description of processes governed by weak neutral currents. This refers to neutral leptonic currents like $\bar{\nu}_\mu\nu_\mu$, $\bar{\nu}_e\nu_e$, $\bar{e}e$ and $\bar{\mu}\mu$, the strangeness-conserving hadronic currents $\bar{p}p$, $\bar{n}n$, and $\bar{\lambda}\lambda$, and finally the strangeness-changing neutral hadronic currents $\bar{n}\lambda$ and $\bar{\lambda}n$.

Theoretical arguments for the existence of neutral currents have been discussed for several years. The point is that, if such currents are introduced, it is possible to construct a consistent renormalizable theory of the weak interaction and to avoid the divergences which occur in the higher orders of the weak interaction. Two elements which are inherent in a large class of renormalizable theories of the weak interaction are: 1) symmetry between the charged and neutral weak currents, and 2) intermediate bosons—the charged $W^\pm$ and neutral Z^0, which interact with these currents. An example of such a theory is the simplest variant of the Weinberg model, in which the Z boson interacts with the lepton currents $\bar{\nu}_e\nu_e$, $\bar{\nu}_\mu\nu_\mu$, $\bar{e}e$ and $\bar{\mu}\mu$ and the hadronic currents $\bar{p}p$ and $\bar{n}_\theta n_\theta$, where $n_\theta = n\cos\theta + \lambda\sin\theta$, the constant for this interaction being approximately the same as for the interactions of charged W bosons with the currents $\bar{e}\nu_e$, $\bar{\mu}\nu_\mu$ and $\bar{n}_\theta p$. An interaction of the neutral neutrino current $\bar{\nu}_\mu\nu_\mu$ with the strangeness-conserving hadronic currents ($\bar{p}p$ and $\bar{n}n$) was detected experimentally at CERN in 1973 and subsequently confirmed in similar experiments at other laboratories. This refers to reactions of the type ν_μ + nucleon $\to \nu_\mu$ + hadrons or $\bar{\nu}_\mu$ + nucleon $\to \bar{\nu}_\mu$ + hadrons. The cross sections for these reactions are, respectively, ~ 0.2 and ~ 0.4 of the cross sections for the analogous reactions governed by the interactions of charged currents, namely ν_μ + nucleon $\to \mu^-$ + hadrons and ν_μ + nucleon $\to \mu_\mu$ + hadrons. Several events have also been observed which can be interpreted as the scattering process $\nu_\mu + e^- \to \nu_\mu + e^-$, which is governed by the interaction of the neutral currents $\bar{\nu}_\mu\nu_\mu$ and $\bar{e}e$. Thus, the first experimental confirmations of the theory have been obtained.

In addition, this theory predicts a number of processes involving neutral strangeness-changing currents: $K_L \to \bar{\mu}\mu$, $K \to \bar{e}e\pi$, $K \to \bar{\nu}\nu\pi$, etc.

Experimentally, these processes are many orders of magnitude weaker (for example, the decay $K_L \to \bar{\mu}\mu$ is 8 orders of magnitude weaker) than the corresponding decays which proceed via charged currents, namely $K^+ \to \mu^+\nu$, $K \to e\nu\pi$ and $K \to \mu\nu\pi$. However, the simplest variant of the theory (the three-quark model) yields similar probabilities for the "neutral" decays (such as $K_L \to \mu^-\mu$) and the "charged" decays (such as $K^+ \to \mu^+\nu$). This follows from the fact that

$$\bar{n}_\theta n_\theta = \bar{n}n\cos^2\theta + \bar{\lambda}\lambda\sin^2\theta + (\bar{n}\lambda + \bar{\lambda}n)\cos\theta\sin\theta,$$

so that the interaction of the currents $\bar{n}\lambda$ and $\bar{\lambda}n$ is of the same order of magnitude as that of the current $\bar{n}n$, while the interaction of the latter is in turn of the same order of magnitude as that of the charged currents $\bar{n}p$ and $\bar{\lambda}p$.

A similar difficulty occurs in connection with the mass difference between the K_L and K_S mesons. The product $(\bar{n}\lambda)(\bar{n}\lambda)$ gives an excessively rapid transition $K^0 \to n\bar{\lambda} \to \lambda\bar{n} \to \bar{K}^0$ and hence an inadmissibly large mass difference between the K_L and K_S mesons, which is many orders of magnitude greater than the experimental value of this quantity.

It should be stressed that these difficulties are in general present independently of the Weinberg model and the existence of Z^0 bosons and were found before this model was proposed. The point is that effective interactions of the type $(\bar{\mu}\mu)(\bar{\lambda}n)$ or $(\bar{n}\lambda)(\bar{n}\lambda)$ occur in second-order perturbation theory even in a theory involving only charged currents.

The simplest diagram of second order in the weak interaction, corresponding to transitions $\bar{\lambda}n \to \bar{\mu}\mu$, has the form shown in Fig. 15a.

From calculations of the $K_L \to \bar{\mu}\mu$ decay amplitude and the mass difference between the K_L and K_S mesons, with allowance for the strong interactions, assuming only that the latter are SU(3)-symmetric, it has been inferred that the large virtual momenta lead to an inadmissibly large contribution to these quantities.

Thus, in theories involving the three quarks p, n and λ, it is not possible to obtain the small probability for decays of the type $K_L \to \bar{\mu}\mu$ and the small value of the mass difference between the K_L and K_S mesons which are required experimentally.

c) **A fourth quark.** This difficulty is eliminated by introducing a fourth quark having the same charge as the p quark. If it is assumed that the two quark doublets (p, n_θ) and (c, λ_θ), where $\lambda_\theta = -n \sin\theta + \lambda \cos\theta$, appear in a symmetric way in the weak interaction, the neutral strangeness-conserving currents remain, while the strangeness-changing currents vanish:

$$\bar{n}_\theta n_\theta + \bar{\lambda}_\theta \lambda_\theta = \bar{n}n + \bar{\lambda}\lambda.$$

The c quark here plays a very peculiar role in the neutral current itself: this quark is required "only" in order to introduce the combinations n_θ and λ_θ in a symmetric way in the current. However, the c quark plays a direct role in the second and higher orders of the weak-interaction theory, cancelling the unwanted diagrams involving p quarks. For example, the diagram of Fig. 15b cancels the above-mentioned diagram of Fig. 15a.

Taking into account the fact that the charged currents have the form $\bar{p}n_\theta$ and $\bar{c}\lambda_\theta$, it is readily verified that the sum of the diagrams of Figs. 15a and 15b vanishes in the limit $m_c = m_p$ and that the unwanted interactions of the type $(\bar{\mu}\mu)(\bar{\lambda}n)$ do not occur, even in second order in the weak interaction. When $m_c \neq m_p$, the contribution from the region of large virtual momenta $p^2 \gg m_c^2, m_p^2$ again vanishes in the sum of the diagrams of Figs. 15a and 15b, so that the resulting contribution is small—of order $G^2(m_c^2 - m_p^2)\sin\theta\cos\theta/\sqrt{2}\pi^2$, and

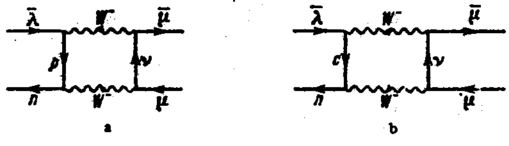


FIG. 15

$$\frac{\Gamma(K_L \to \mu^+\mu^-)}{\Gamma(K^+ \to \mu^-\nu)} \approx \frac{G^2 m_c^4 \cos^2\theta}{2\pi^4} \quad \text{for} \quad m_c \gg m_p.$$

As the limits on the decay $K_L \to \bar{\mu}\mu$ are very strong, $\Gamma(K_L \to \bar{\mu}\mu) \lesssim 10^{-8}\Gamma_{K_L}$, it follows that m_c cannot be much greater than 10 GeV.

d) **SU(4) symmetry.** If the foregoing cancellations are not to be violated by the strong interaction, the latter must be at least approximately symmetric with respect to the two quarks c and p. In conjunction with SU(3) symmetry of the three quarks (p, n, λ), this leads to the condition of SU(4) symmetry. In nature, SU(4) symmetry is broken much more strongly than SU(3) symmetry. However, it cannot be violated completely, since in that case the cancellations which we require would disappear. The violation of SU(4) symmetry must tend to zero at large virtual momenta much greater than the mass difference between the p and c quarks. The existence of broken SU(4) symmetry implies, for example, that the well-known meson nonets must appear in 16-plets, while the well-known baryon octet and decuplet must appear in appropriate 20-plets.

We note that, if the strong interactions played a negligibly small role for large virtual momenta ($p^2 \gtrsim m_c^2$) and if the quarks were practically bare for such momenta, we would not require SU(4) symmetry of the strong interaction in order to have a cancellation between the diagrams of types a and b of Fig. 15. However, the strong-interaction models discussed below do not possess such a "superconvergence" property.

e) **Charm and supercharge.** In considering the particles belonging to SU(4) multiplets, it is convenient to introduce an additive quantum number "supercharge" given by $\sigma = 3\langle Q\rangle$, where $\langle Q\rangle$ is the average charge of an SU(3) multiplet. For the triplet of quarks p, n and λ, we have $\sigma = 0$, so that all particles constructed from the quarks p, n and λ are superneutral. The c quark has $\sigma = 2$, and particles which contain the c quark or any other quarks with $\sigma \neq 0$ may be called supercharged.

To distinguish between the n and λ quarks, we say that the λ quark, unlike the n quark, carries the quantum number "strangeness" S ($S(\lambda) = -1$). Similarly, to distinguish between the p and c quarks, we say that the c quark carries the quantum number "charm" C ($C(c) = +1$). It follows from the foregoing remarks that $C = \sigma/2$. Particles which contain the c quark are called charmed or supercharged. This terminology seems unfortunate, but there are as yet no better alternatives. In the literature, particles containing the pair $\bar{c}c$ are called particles with hidden charm.

f) The analogy with leptons. It is interesting to compare the four quarks with the four leptons:

$$p,\ n_\theta,\ c,\ \lambda_\theta,$$
$$\nu_e,\ e,\ \nu_\mu,\ \mu.$$

These particles appear in an analogous way in the weak currents. In particular, they are represented by their left-handed components $\psi_L = (1/2)(1+\gamma_5)\psi$ in the charged weak current. The four left-handed isotopic doublets, which appear in the weak current in a symmetric way, have the form

$$\begin{pmatrix} p_L \\ n_{\theta L}\end{pmatrix},\quad \begin{pmatrix} c_L \\ \lambda_{\theta L}\end{pmatrix},\quad \begin{pmatrix} \nu_{eL} \\ e_L\end{pmatrix},\quad \begin{pmatrix} \nu_{\mu L} \\ \mu_L\end{pmatrix}.$$

The approach to lepton-hadron symmetry was in fact the original basis for the hypothesis that c quarks exist. However, we cannot rule out the possibility that both lepton and quark series will be required (see Sec. 3j of Chap. VI). Would lepton-quark symmetry be preserved in this case? What form would it take?

g) **Colored quarks and white hadrons.** To cope with the difficulties which arise in connection with the weak interactions, it seems sufficient for the present to go over from three to four quarks. However, to overcome the difficulties in understanding the strong and electromagnetic interactions of hadrons, it appears that we must triple the number of quarks of each type: p, n, λ and c.

This triplication eliminates the inconsistency between the spin and statistics of the quarks which arises in the quark model of the baryons. For example, the Ω^- hyperon consists of three λ quarks with parallel spins in an s-wave state, which is forbidden by the Pauli principle. If all three λ quarks belonging to the Ω^- hyperon are different, say λ_1, λ_2 and λ_3, then this inconsistency disappears. In this case, the wave function of the three λ quarks in the Ω^- hyperon is completely antisymmetric and has the form $\epsilon^{ikl}\lambda_i\lambda_k\lambda_l$ (i, k, l = 1, 2, 3). This is also true of the wave functions of the quarks in the other baryons. For instance, the Δ^{++} baryon corresponds to a wave function $\epsilon^{ikl}p_ip_kp_l$.

The 12 quarks can be tabulated as follows:

$$\begin{pmatrix} p_1 & p_2 & p_3 \\ n_1 & n_2 & n_3 \\ \lambda_1 & \lambda_2 & \lambda_3 \\ c_1 & c_2 & c_3 \end{pmatrix}.$$

It is conventional to say that the columns of this table are distinguished by color. For example, 1 is yellow, 2 is blue, and 3 is red. If the strong interactions are completely degenerate with respect to color, we are dealing with colored SU(3) symmetry, which is usually designated SU(3)$'$. In that case, the ordinary baryons, constructed by means of the SU(3)$'$-invariant antisymmetric tensor ϵ^{ikl}, belong to the singlet representation of the group SU(3) and are, as it were, white. The ordinary mesons are also white. For example,

$$\pi^+ = \frac{1}{\sqrt{3}}\, p_i\tilde{n}^i = \frac{1}{\sqrt{3}}\, p_i\tilde{n}^k\delta_k^i = \frac{1}{\sqrt{3}}\,(p_1\tilde{n}^1 + p_2\tilde{n}^2 + p_3\tilde{n}^3).$$

(The tensor δ_k^i is also an invariant tensor of the color group SU(3)$'$.)

h) **The alternatives.** In introducing the concept of color, we are faced with the following choice.

1) We may assume that color symmetry is a strict symmetry, i.e., that not only the strong interaction, but also the electromagnetic and weak interactions, are SU(3)$'$-invariant. In this case, the charges of the quarks do not depend on their color and must have fractional values.

2) We may assume that the electromagnetic interaction (and possibly the weak interaction) violates color symmetry. In this case, the charges of the quarks may have integral values.

Searches in cosmic rays, in accelerator experiments and in the matter of our environment have not revealed any fractionally charged quarks. (The absence of so-called relic quarks in the matter of our environment has been established with particularly high accuracy. The experimental upper limt in this case is about 15 orders of magnitude below the value obtained from calculations based on the hot-universe model, which yield 1 quark per 10^{10} protons). In the case of strict color SU(3)$'$ symmetry and fractionally charged quarks, it is therefore natural to attempt to construct a theory in which quarks (and, in general, colored particles) cannot exist in a free state.

If, on the other hand, quarks have integral charges, then the electromagnetic current is not a color singlet. In this case, it is natural to expect other color charges as well, so that colored particles are physically observable.

We shall bear in mind both possibilities in the discussion which follows.

i) **Quarks with integral charges.** The electric charges of colored quarks are equal to the charges of the ordinary superneutral fractionally charged quarks only after averaging with respect to color:

$$Q_{p_i} = Q_{c_i} = \frac{2}{3} + \alpha_i,$$
$$Q_{n_i} = Q_{\lambda_i} = -\frac{1}{3} + \alpha_i,$$
$$\sum_{i=1}^{3} \alpha_i = 0.$$

It is most common to discuss colored quarks having the following set of integral charges:

$$\begin{array}{cc} & \begin{array}{ccc} y & b & r \end{array} \\ Q = & \begin{pmatrix} 1 & 1 & 0 \\ 0 & 0 & -1 \\ 0 & 0 & -1 \\ 1 & 1 & 0 \end{pmatrix} \begin{array}{c} p \\ n \\ \lambda \\ c \end{array}. \end{array}$$

These values of the charges correspond to $\alpha_1 = \alpha_2 = 1/3$ and $\alpha_3 = -2/3$. Stable quarks with integral charges are not excluded experimentally, provided that their masses are sufficiently high. Mass spectrometric searches for heavy singly charged particles rule out a concentration of such particles greater than 10^{-18} for masses between 6 and 16 GeV. It would be useful to extend the range of these searches.

j) **Unstable quarks.** Colored quarks with integral charges may be unstable. A number of variants of unstable quarks have been discussed in the literature. In one of them, it is conjectured that the baryonic charges of the quarks are integral and are described, for example, by the matrix

$$\begin{array}{cc} & \begin{array}{ccc} y & b & r \end{array} \\ B = & \begin{pmatrix} 1 & 1 & -1 \\ 1 & 1 & -1 \\ 1 & 1 & -1 \\ 1 & 1 & -1 \end{pmatrix} \begin{array}{c} p \\ n \\ \lambda \\ c \end{array}. \end{array}$$

In this case, the quarks, violating the color symmetry, could decay into ordinary white mesons and nucleons (the red quarks would decay into antinucleons).

According to another variant, the baryonic charges of the quarks have the fractional value 1/3, but the quarks decay into leptons without conserving baryonic and leptonic charges. In this case, it is possible to reconcile the quark lifetime $\lesssim 10^{-10}$ sec with the proton lifetime $\gtrsim 10^{37}$ sec by assuming that the quark is about 10 times as heavy as the proton and taking into account the fact that proton decay would take place only in third-order perturbation theory in the interaction responsible for quark decay, since all three quarks in the proton would have to decay simultaneously.

It should be noted, however, that nonconservation of baryonic charge would necessarily have led to a much smaller baryon asymmetry of the universe, even in the earliest stages of the hot universe, then that which is observed at the present time.

The following scheme summarizes the variants of quarks which we have discussed:

QUARKS

- With fractional electric charge (apparently physically unobservable)
- With integral electric charge (physically observable)

Unstable — Stable

- With integral baryonic charge B (quarks decay into mesons and nucleons, conserving B. The proton is stable)
- With fractional baryonic charge B (quarks decay into leptons without conserving B. The proton is unstable)

k) Gluons. What are the strong interactions between the quarks like? At first sight, it is impossible to answer this question, especially in view of the fact that we are completely ignorant of whether or not quarks exist. However, there are a number of requirements which must be imposed on the strong interaction of the quarks in order to give the experimentally observed regularities in the properties of hadrons. It is usually assumed that the strong interactions between quarks are due to the exchange of vector particles known as gluons (from the English word "glue").

In Weinberg's theory, gluons must be SU(4) singlets, i.e., they must not carry any quantum numbers such as isospin, strangeness or supercharge. In general, the symmetry group G(W) of the weak and electromagnetic interactions on the one hand, and the symmetry group G(S) of the strong interaction on the other hand, must be independent. The symmetry group of the full Lagrangian must factorize and have the form G(W) × G(S). Otherwise, the Feynman diagrams containing both exchanges of gluons and exchanges of intermediate bosons would give excessively large effects of nonconservation of parity and strangeness in hadronic processes.

l) The gluon cannot be white. It might appear to be most natural to assume that there exists only a single gluon, which is a singlet with respect to each of the groups SU(4) and SU(3)′. The theory of the strong interaction would then be analogous to the theory of the electromagnetic interaction. However, such a gluon, which is a singlet with respect to SU(3)′, is unsatisfactory for at least three reasons.

Firstly, although we could account for the existence of mesons in this case (since a quark and antiquark attract one another), it would be difficult to explain the existence of baryons, since two quarks repel one another. This conclusion is based on an analysis of a certain set of diagrams, and it seems highly improbable that the diagrams which are not taken into account would change the repulsion to an attraction.

Secondly, a theory involving a single white gluon would have too high a degree of symmetry: instead of SU(4) × SU(3)′, it would have the symmetry SU(12), since all 12 quarks would interact identically with the gluon. The subgroup SU(9) corresponding to the nine colored quarks p_i, n_i and λ_i would be particularly hazardous in this case. The known mesons, together with their colored analogues, would have to belong to 81-plets. If color symmetry were a strict symmetry, the mass splitting in such 81-plets would occur only because of the violation of ordinary SU(3) symmetry and would have to be no greater than in the ordinary nonets, which is excluded experimentally.

If color SU(3)′ symmetry were violated much more strongly than ordinary SU(3) symmetry, the masses of the mesons belonging to a single SU(9) multiplet would be strongly split. At the present time, we cannot exclude the existence of a large number of relatively heavy colored mesons. But this possibility does not seem particularly attractive.

Finally, a third objection against a white gluon is the fact that, if white gluons are exchanged, the effective gluon "charges" of the quarks grow as the quarks approach one another and as the momentum transfers increase. This enhancement of the interaction at small distances is a result of the total contribution of a certain series of diagrams, and it leads to an asymptotic ultraviolet instability of the theory. This theoretical phenomenon was discovered in ordinary electrodynamics as early as the mid-50s. The charge at distance r is related to the charge at distance R by the equation

$$e^2(R) = \frac{e^2(r)}{1 + [e^2(r)/3\pi] \ln (R^2/r^2)}.$$

This equation implies that the physical renormalized charge e(R) at large distances R tends not to the constant $\sqrt{\alpha}$, but to zero, like $[\ln(R/r)]^{-1/2}$. This difficulty was called the "zero-charge problem" in the 1950s Landau and Pomeranchuk). It follows from the same formula that a finite e(R) corresponds to e(r) which increases with decreasing r.

Experiments on deep inelastic interactions of electrons and neutrinos with nucleons indicate that the strong interaction between quarks does not grow, but falls off, at large momentum transfers. Apparently, the phenomenon known as "asymptotic freedom" occurs. The white-gluon theory cannot reproduce the property of asymptotic freedom, giving instead an asymptotic growth of the strong interaction.

m) Colored gluons and the masses of white hadrons. The foregoing difficulties are absent in a model in which the strong interaction is governed by the exchange of an octet of colored gluons interacting with eight colored charges—the generators of the group SU(3)′. The corresponding strong current has the form

$$j_S = \bar{p}\gamma_\mu\lambda_i p + \bar{n}\gamma_\mu\lambda_i n + \bar{\lambda}\gamma_\mu\lambda_i\lambda + \bar{c}\gamma_\mu\lambda_i c.$$

where γ_μ are the four Dirac matrices (μ = 1, 2, 3, 4), and λ_i are the eight well-known matrices which generate the group SU(3)′:

$$\lambda_1=\begin{pmatrix}0&1&0\\1&0&0\\0&0&0\end{pmatrix},\quad \lambda_2=\begin{pmatrix}0&-i&0\\i&0&0\\0&0&0\end{pmatrix},\quad \lambda_3=\begin{pmatrix}1&0&0\\0&-1&0\\0&0&0\end{pmatrix},\quad \lambda_4=\begin{pmatrix}0&0&1\\0&0&0\\1&0&0\end{pmatrix},$$

$$\lambda_5=\begin{pmatrix}0&0&-i\\0&0&0\\i&0&0\end{pmatrix},\quad \lambda_6=\begin{pmatrix}0&0&0\\0&0&1\\0&1&0\end{pmatrix},\quad \lambda_7=\begin{pmatrix}0&0&0\\0&0&-i\\0&i&0\end{pmatrix},\quad \lambda_8=\frac{1}{\sqrt{3}}\begin{pmatrix}1&0&0\\0&1&0\\0&0&-2\end{pmatrix}.$$

It is easy to see that SU(12) degeneracy does not occur in this case, since the strong interaction is not SU(12)-invariant: its symmetry is SU(4) × SU(3)′. Moreover, within the limitations of the potential approximation, it is readily shown that the masses of white particles must be less than those of colored particles, and there is an attraction not only between a quark and antiquark, but also between quarks.

Let us consider a potential interaction between two quarks or between a quark and an antiquark. This interaction is governed by the exchange of a gluon which carries color charge, so that it depends on the color state in which the interacting quarks occur. Such an exchange potential is proportional to

$$\lambda^1\lambda^2 \equiv \sum_{i=1}^{8} \lambda_i^1\lambda_i^2,$$

where the upper index 1 (2) indicates that a matrix operates on the color degrees of freedom of the first (second) quark. Obviously,

$$2\lambda^1\lambda^2 = (\lambda^1 + \lambda^2)^2 - (\lambda^1)^2 - (\lambda^2)^2.$$

Adding the squares of the eight matrices λ_i, we find that $(\lambda)^2 = 16/3$. As to the quadratic Casimir operator $(1/4)(\lambda^1 + \lambda^2)^2$, its value is determined by the dimensionality of the state in which the two quarks occur. If this state is the singlet state (in the case $q + \overline{q}$), the eigenvalue of $(\lambda^1 + \lambda^2)^2$ is equal to zero. If it is the triplet state, this eigenvalue is equal to 16/3. Thus, a quark and an antiquark in the singlet state attract one another twice as strongly as two quarks in the triplet state. (We denote the corresponding potentials by 2U and U, respectively.)

We now take into account the fact that a white baryon, say a nucleon, consists of three quarks which are in a state $\epsilon^{ijk}q_iq_jq_k$ that is completely antisymmetric with respect to color, so that each pair of quarks in the nucleon is in a triplet state of the type $\epsilon^{ijk}q_iq_j$. As there are three such pairs in the nucleon, the total binding energy of the quarks in the nucleon is 3U, which is 1.5 times as large as the binding of a quark and an antiquark in a white meson. Thus, the binding energy of the quarks in white particles is proportional to the number of quarks of which these particles are composed.

n) **"Realistic" colored quarks.** According to the hypothesis of realistic quarks with integral charges, free quarks are rather heavy (for example, $m_q \sim 10$ GeV), but their masses in white hadrons are almost completely "eaten up" by their mutual binding. The above-mentioned saturation property of gluon forces in the case of an exchange gluon potential suggests that such a possibility might be realized if U is close to m_q. In this case, the remaining effective quark mass is small: $\overline{m}_q = m_q - U$ is less than or of the order of several hundred MeV. (It should be emphasized that, with such a large binding energy, discussions based on the analysis of a potential can only be expected to serve as a guide. It suffices to note that a small increase in the depth of the potential may give a negative effective quark mass $\overline{m}_q$ in the naive approximation of a nonrelativistic potential, or an imaginary one in the case of an equation of the Bethe-Salpeter type.)

If $\overline{m}_q \ll m_q$, the masses of colored particles will be much greater than those of the white hadrons. Thus, the masses of the triplet diquarks will be of order m_q. In the sextet diquarks and the octet $q\overline{q}$ pairs, the particles repel one another, the repulsion energy being equal to U/4 in the first case and U/8 in the second.

We note that for an arbitrary representation $T^{[a]}_{[b]}$ of the group SU(3) we have

$$C(a, b) = \frac{1}{3}(a^2 + b^2 + ab + 3a + 3b),$$

where C(a, b) is the value of the quadratic Casimir operator $(1/4)(\Sigma_m\lambda^m)^2$. The dimensionality of the representation $T^{[a]}_{[b]}$ is given by

$$N = \frac{1}{2}(a+1)(b+1)(a+1+b+1).$$

For N = 1, we have a = b = 0 and C = 0; for N = 3, a = 1, b = 0 and C = 3/4; for N = 6, a = 2, b = 0 and C = 10/3; for N = 8, a = b = 1 and C = 3; for N = 10, a = 3, b = 0 and C = 4.

In the case of "realistic" heavy quarks, their charges must naturally be integral, and the fact that they are not observed experimentally must be attributed to the large mass of the quarks and/or their instability. In this case, only two of the eight gluons are necessarily electrically neutral, while six of them have charges $\pm(2\alpha_1 + \alpha_2)$, $\pm(2\alpha_2 + \alpha_1)$ and $\pm(\alpha_1 - \alpha_2)$ (see the expression given earlier for the charges of colored quarks in Sec. 1i).

If we take the limit as $m_q \to \infty$ and $U \to \infty$ in such a way that $\overline{m}_q = m_q - U$ remains of order 1/3 GeV, then the model of "realistic" quarks goes over into a model of unphysical confined quarks. This is the simplest model of confinement. Other models are more refined.

o) **Confined quarks.** Many theoretical papers in recent years have been devoted to the problem of quark confinement. The chief hopes here are connected with the properties of the octet of colored gluons.

It is assumed that the eight gluons are massless and are generated by eight conserved color currents. As the gluons themselves carry color charges, the gluon sources are not only the quarks, but also the gluons themselves. To cite an analogy with photons, the gluons represent "luminous light" (ordinary photons do not "emit light"). It is obvious that this luminosity must become particularly large at small gluon momenta and that an infrared instability of the theory may then occur. The theoretical expectations are that, owing to the infrared instability, the interaction potential between color changes will not fall off as the distance between them increases and that two colored charges, say a quark and an antiquark which form a meson, cannot therefore become free from one another, even for arbitrarily high excitation energy of the meson.

So far, physics has always dealt with potentials which fall off to zero with increasing distance, either slowly, as in the case of the Coulomb potential (Fig. 16a), or rapidly, as in the case of the Yukawa potential (Fig. 16b). If colored gluons actually lead to potentials which increase with distance, we are faced with the situation shown in Fig. 16c. No matter how high the energy transferred to a quark, it cannot escape from the peak shown in Fig. 16c. The quark may be excited into one of the higher levels, but it then goes down into its original state, radiating its energy in the form of particles (photons, leptons and white quarks).

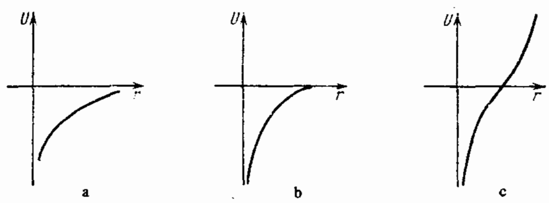


FIG. 16

Confinement must also be extended to other colored particles, including gluons, so that no free colored charges can exist and all observable particles are white.

In this approach, in contrast with the model of realistic colored quarks, color symmetry must be a strict symmetry, and the electric charges of the quarks must be independent of their color and hence must have fractional values.

p) Non-Abelian local symmetry and asymptotic freedom. A theory of massless vector particles interacting with conserved isotopic currents was first proposed by Yang and Mills in 1954. Such a theory is said to possess local SU(2) symmetry. This means that it is invariant under isotopic rotations whose magnitude is a function of the space-time point. Unlike the theory of photons, which is invariant under local Abelian (mutually commuting) gauge transformations, the Yang-Mills theory is invariant under non-Abelian gauge transformations. Colored gluons correspond to local SU(3) symmetry. An important property of such non-Abelian gauge theories is that the effective charges become weaker in magnitude at small distances, tending to zero as the distance tends to zero, corresponding to asymptotically large momentum transfers. Such theories are called asymptotically free. In the case of non-Abelian theories, unlike Abelian theories, an initial charge "placed" in a vacuum experiences anti-screening. Thus, in the case of $SU(3)'$ symmetry, the color charges at distances R and r are related to one another by the equation

$$g^2(R)=\frac{g^2(r)}{1-(25/12\pi)\,g^2(r)\ln(R^2/r^2)}.$$

Owing to the minus sign in the denominator, the charge is greater at large distances than at small distances. At sufficiently large R, the denominator vanishes and $g^2(R)$ becomes infinite. We see that asymptotic freedom in the ultraviolet region (at small r) is intimately related to the instability in the infrared region (at large R). It is this infrared instability of the theory of colored gluons on which those theorists who are trying to construct a theory of absolute confinement of colored quarks inside white hadrons are pinning their hopes.

It should be emphasized, however, that between the infinitely rising effective charge at large distances which follows from the expression for $g^2(R)$ and the hump-like potential shown in Fig. 16c there is a gap which is filled at present only with optimistic hopes and studies of one-dimensional models or models involving a discrete lattice-like space-time. The goal of studying such models is to obtain something like a tube of lines of force of the gluon field between two colored charges; this is usually called a string. The energy of such a string would grow in proportion to its length, so that two color charges joined by a string could not become free from one another.

The property of asymptotic freedom implies, in particular, that it is possible to apply perturbation theory in studying strongly bound quarks situated at the lower part of the hump in Fig. 16c.

Asymptotic freedom is the theoretical basis of the parton model, which provides a successful explanation of the phenomenon of scaling in deep inelastic electroproduction, in neutrino reactions at high energies and in hadron-hadron collisions. Owing to the weak interaction between quarks at small distances, they may be considered as partons—hypothetical point-like constituents of hadrons.

Now that we have examined the main arguments in favor of color, we shall consider briefly two other effects which are interpreted by invoking colored quarks.

q) The annihilation $e^+e^- \to$ hadrons. According to the parton model, e^+e^- annihilation into hadrons at high energies proceeds via annihilation into a quark-antiquark pair. It is assumed here that the cross section for this process is what it would be for point-like bare quarks and that the subsequent transformation of the $q\bar{q}$ pair into hadrons does not alter the value of this cross section. In that case, the ratio of the cross section for annihilation into hadrons to the cross section for annihilation into a $\mu^+\mu^-$ pair is given by

$$R\equiv\frac{\sigma(e^+e^-\to \text{hadrons})}{\sigma(e^+e^-\to\mu^+\mu^-)}=\sum_i Q_i^2,$$

where Q_i is the charge of the quark of type i, and the summation is carried out over all types of quark pairs which are produced. Below the threshold for producing charmed particles, allowance must be made for the production of $p\bar{p}$, $n\bar{n}$ and $\lambda\bar{\lambda}$ pairs. The quantity R then has the value $4/9 + 1/9 + 1/9 = 2/3$ in the model of three colorless quarks, or 2 in the model of three colored quarks. Experimentally, $R \approx 2.5$ in the range 1.5–3.5 GeV, and $R \approx 5.5$ in the range 4.5–7.5 GeV. We note that, at energies much higher than the threshold for producing the $c\bar{c}$ pair, we must have $R = 2 + (4/9)3 = 10/3$. In the parton model, a large asymptotic value of R would imply the existence of other types of quarks in addition to the p, n, λ and c. The possibility of p, n and λ quarks with large charges and supercharges ($\alpha_1 \neq 0$ and/or $\alpha_2 \neq 0$), which would also lead to large values of R, is apparently inconsistent with the data on deep inelastic interactions of electrons and neutrinos with nucleons.

r) The decay $\pi^0 \to 2\gamma$. The calculation of the $\pi^0 \to 2\gamma$ decay width in the framework of the quark model on the basis of the partially conserved axial-vector current has a remarkable property: the corrections to it due to the virtual strong interactions should be small. This calculation yields

$$\Gamma=\frac{F^2m_\pi^3}{64\pi},\quad \text{where}\quad F=\frac{\sqrt{2}}{\pi}\frac{\alpha S}{f_\pi},$$

$\alpha = 1/137$, $f_\pi = 0.96m_{\pi^+}$ is the constant which determines the $\pi^+ \to \mu^+\nu$ decay width, $S = Q_p^2 - Q_n^2 = (Q_p - Q_n)(Q_p + Q_n) = 1/3$ in the model of singly-colored quarks, and

$$S=\sum_{\alpha=1}^{3}(Q_{p\alpha}^2-Q_{n\alpha}^2)=\sum_{\alpha=1}^{3}Q_{p\alpha}+\sum_{\alpha=1}^{3}Q_{n\alpha}=1$$

in the model of colored quarks. Experiment is inconsistent with the value $S = 1/3$, but is in good agreement with $S = 1$.

Other merits of colored quarks and colored gluons will be mentioned later, when we discuss the properties of ψ mesons and charmed hadrons.

2. Quarks and ψ Mesons

A theoretical model of the ψ mesons must first of all provide answers to the following two questions: 1) What hadronic degrees of freedom correspond to the ψ mesons? 2) What selection rule is reponsible for the small widths of these particles?

Within the scope of quark models, three possibilities have been considered: 1) the ψ mesons are colored mesons (colored analogous to the φ, ω or ρ); 2) the ψ mesons are $^3S_1$ orthostates of the $c\bar{c}$ system, known as orthocharmonium (in analogy with orthopositronium); and 3) the ψ mesons are orthostates ofso me other quarks, and not the c quark. (Situations in which one thing was sought and another found are not so rare in physics: it suffices to recall the history of how muons were discovered.)

The discussion which follows is concerned mainly with the properties of orthocharmonium. But we shall begin with a few words about the hypothesis of colored ψ mesons.

a) **Are the ψ mesons colored?** It is easy to see why this possibility has atrracted attention. If the strong interaction conserves color, strong hadronic decays of the ψ mesons would be forbidden and it would be possible to account for their small widths. There is no unique answer to the question of what color states of the vector 81-plet to associate with the ψ and ψ' mesons. If color symmetry is broken only by the electromagnetic interactions and if the quark charges have the values (1, 1, 0) for the p and (0, 0, −1) for the n and λ, then the electromagnetic current transforms as follows:

$$j_{\text{e.m.}}=\bar{p}_1p_1+\bar{p}_2p_2-\bar{n}_3n_3-\bar{\lambda}_3\lambda_3=(\bar{p}_1p_1+\bar{p}_2p_2+\bar{p}_3p_3)-(\bar{p}_3p_3+\bar{n}_3n_3+\bar{\lambda}_3\lambda_3)$$
$$\sim\begin{pmatrix}1&&\\&0&\\&&0\end{pmatrix}\begin{pmatrix}1&&\\&1&\\&&1\end{pmatrix}'-\begin{pmatrix}1&&\\&1&\\&&1\end{pmatrix}\begin{pmatrix}0&&\\&0&\\&&1\end{pmatrix}',$$

where the unprimed matrices operate in the space SU(3), while the primed matrices operate in the color space SU(3)′. The first term conserves color, since it is proportional to the matrix $\begin{pmatrix}1&&\\&1&\\&&1\end{pmatrix}'$, but the second term does not. Now since the matrix $\begin{pmatrix}0&&\\&0&\\&&1\end{pmatrix}'$ is diagonal, the second term conserves the projection of color. (In this sense, it is analogous to the mass term $\Delta m\bar{\lambda}\lambda$, which violates ordinary SU(3) symmetry.) Moreover, this term is invariant with respect to the SU(2)′ subgroup of transformations in the space of the yellow and blue colors, and also with respect to the ordinary SU(3) group and its isotopic subgroup. Therefore the transition $\psi\to\gamma\to e^+e^-$ is possible only for the mesons ω_8 and φ_8, which are isotopic scalars with respect to both SU(2) and SU(2)′. The index 8 indicates that the appropriate meson states are constructed in terms of the matrix λ_8 in the space SU(3)′:

$$\varphi_8=\frac{1}{\sqrt{6}}(\lambda_1\bar{\lambda}_1+\lambda_2\bar{\lambda}_2-2\lambda_3\bar{\lambda}_3),$$
$$\omega_8=\frac{1}{\sqrt{6}}\left[\frac{1}{\sqrt{2}}(p_1\bar{p}_1+n_1\bar{n}_1)+\frac{1}{\sqrt{2}}(p_2\bar{p}_2+n_2\bar{n}_2)-\frac{2}{\sqrt{2}}(p_3\bar{p}_3+n_3\bar{n}_3)\right].$$

The transition $\psi\to\gamma\to e^+e^-$ is not possible for the remaining colored mesons, such as φ_3 or ρ_8. It is therefore natural to assume that $\psi=\omega_8$ and $\psi'=\varphi_8$. Such colored ψ and ψ' mesons must decay via the electromagnetic interaction (in general, this does not apply to the decays $\psi'\to\psi$ + hadrons). In particular, hadronic decays of the ψ and ψ' must be accompanied by photon emission. The isoscalar character of the photons which carry color in this case implies that their G-parity is negative, so that the pion system in the decay $\psi\to n\pi+\gamma$ must have positive G-parity, i.e., n must be even. An experimental test of this prediction might decide the fate of this scheme. It would also be of interest to search for the decays $\psi\to$ hadrons + e^+e^-. We note that the decays $\psi\to$ hadrons + γ must have the highest probability according to this scheme; the decays $\psi\to$ hadrons + e^+e^- should comprise a fraction $\sim\alpha$ of those of the first type, and the decay $\psi\to e^+e^-$ should have a still smaller probability. The fact that experimentally the decay $\psi\to e^+e^-$ accounts for about 6% of the full width of the ψ meson constitute an argument against the interpretation of the ψ particles as colored mesons.

From the purely theoretical point of view, this interpretation seems unattractive. It would entail the prediction of enormous SU(9) multiplets: in particular, the ψ mesons would belong to an 81-plet. There is no such approximate SU(9) degeneracy of white and colored mesons in the colored-gluon model. We recall that, in the approximation of an exchange potential, a quark and an antiquark in the color octet repel one another and do not form a bound state, while a pair of quarks in the color triplet attract one another. A color octet of mesons should therefore not exist in the potential approximation; but if such mesons do actually exist, there is all the more reason for the existence and much lighter masses of diquarks which form a color triplet. It should be stressed that these would have to be physically observable.

There have also been discussions of more complex color interpretations of the ψ mesons based on the idea that color is violated by not only the electromagnetic interaction, but also the strong interaction. In this case, φ_8 and ω_8 may decay strongly, and ω_3, φ_3 and ρ_8 must be narrow, where

$$\omega_3=\frac{1}{\sqrt{2}}\left[\frac{1}{\sqrt{2}}(p_1\bar{p}_1+n_1\bar{n}_1)-\frac{1}{\sqrt{2}}(p_2\bar{p}_2+n_2\bar{n}_2)\right],$$
$$\varphi_3=\frac{1}{\sqrt{2}}(\lambda_1\bar{\lambda}_1-\lambda_2\bar{\lambda}_2),$$
$$\rho_8=\frac{1}{\sqrt{6}}\left[\frac{1}{\sqrt{2}}(p_1\bar{p}_1-n_1\bar{n}_1)+\frac{1}{\sqrt{2}}(p_2\bar{p}_2-n_2\bar{n}_2)-\frac{2}{\sqrt{2}}(p_3\bar{p}_3-n_3\bar{n}_3)\right].$$

b) **Charmonium.** The assumption that the ψ meson is the $^3S_1$ ground state of the quarks $c\bar{c}$ immediately raises a number of questions:

1) Can we account for the observed width of the decays $\psi\to e^+e^-$?

2) Can we account for the observed width of the decays $\psi\to$ hadrons?

3) What state of the $c\bar{c}$ system corresponds to the ψ' meson? How should this meson decay?

4) What properties should mesons corresponding to the $^1S_0$ state of the $c\bar{c}$ pair have?

5) What properties should mesons corresponding to states of the $c\bar{c}$ pair with non-zero orbital angular momentum (P and D states) have?

6) What state of the $c\bar{c}$ system corresponds to the 4.15-GeV resonance if it belongs to the same family as the ψ and ψ' mesons, and why does this resonance have such a large width?

c) **The decays $\psi\to e^+e^-$ and $\psi\to\mu^+\mu^-$.** It is reasonable to compare the decay $\psi\to\gamma\to e^+e^-$ with the decays $\rho\to e^+e^-$, $\omega\to e^+e^-$ and $\varphi\to e^+e^-$. If all these mesons had the same mass, their widths according to the quark model would be in the same ratio as the squares of the corresponding charges:

$$\Gamma_\psi:\Gamma_\rho:\Gamma_\omega:\Gamma_\varphi=Q_c^2:\frac{(Q_p-Q_n)^2}{2}:\frac{(Q_p+Q_n)^2}{2}:Q_\lambda^2=8:9:1:2.$$

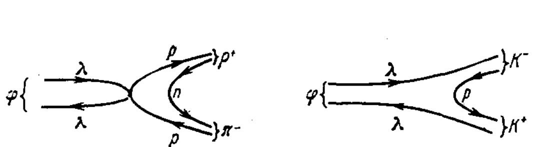
Experimentally, Γ_ψ = 5 keV, Γ_ρ = 6.3 keV, Γ_ω = 0.7

keV and Γ_φ = 1.3 keV, which is close to the prediction. Unfortunately, we cannot say that the theoretical and experimental results are in agreement, since we have not taken into account the mass difference, which is quite substantial in the case of the ψ. If we assume that the width is proportional to the mass, we find that the width of the ψ is smaller than the predicted value by about a factor of three or four. Considering the crudeness of the model, the agreement is not bad, even in this form.

d) **The decay $\varphi \to 3\pi$ and Zweig's rule.** Why are the decays $\psi \to$ hadrons suppressed? Before attempting to answer this question, we point out that a similar suppression has already been well known for a long time; we have in mind the decay $\varphi \to 3\pi$. The width of this decay is only 0.7 MeV, although this decay is not forbidden by SU(3) or isospin selection rules. Comparing this value of 0.7 MeV with the "expected width" ~350 MeV, we see that the decay $\varphi \to 3\pi$ is suppressed by a factor of about 500. (The "expected" width is obtained by starting with an $\omega \to 3\pi$ decay width equal to 10 MeV and estimating how it would be enhanced if the mass m of the decaying meson is increased from 780 MeV to 1020 MeV. This estimate makes use of the fact that $\Gamma \sim m^7$ if the masses of the pions are neglected; in addition, allowance for these masses and for the ρ resonance in the sequence $\omega \to \pi + \rho \to 3\pi$ further enhances the estimated probability. The resulting "expected" width of 350 MeV is apparently somewhat overestimated, since such large widths are more characteristic of heavier mesons like the ρ', whose mass is 1.6 GeV. Thus, it would be more correct to say that the decay $\varphi \to 3\pi$ is suppressed by a factor of not 500, but perhaps 300.)

We note that, despite its small phase space, the decay $\varphi \to K\overline{K}$ has a width ~3 MeV and its matrix element is not suppressed. In the quark model, this means that annihilation of the quarks $\lambda\overline{\lambda}$ which form the φ meson into the pairs of quarks $p\overline{p}$ and $n\overline{n}$ is forbidden, while the decay in which the initial λ and $\overline{\lambda}$ quarks do not annihilate is allowed. An example of a forbidden decay is shown in Fig. 17a, and an example of an allowed decay is shown in Fig. 17b. This regularity, according to which annihilation diagrams such as that of Fig. 17a are suppressed, is known as Zweig's rule.

e) **Zweig's rule and the meson masses.** Since the annihilation $\lambda\overline{\lambda} \to p\overline{p}$, $n\overline{n}$ is weak, the admixture of $p\overline{p}$ and $n\overline{n}$ in the φ meson must be small; the same applies to the admixture of $\lambda\overline{\lambda}$ in the ω meson: $\omega \approx (1/\sqrt{2})\,(p\overline{p} + n\overline{n}) + \epsilon\lambda\overline{\lambda}$, where $\epsilon \sim 5\%$. The small mass splitting between the ρ^0 and ω mesons (m_ρ = 770 ± 10 MeV, m_ω = 782.7 ± 0.6 MeV) indicates that the amplitude for the annihilation quark transition $p\overline{p} \to n\overline{n}$ is also small. (If there were no annihilation, the masses would be degenerate.) The fact that, in spite of the weakness of the annihilation, we are still dealing with the states having definite isospin T = 1, $\rho^0 = (1/\sqrt{2})(p\overline{p} - n\overline{n})$, and T = 0, $\omega = (1/\sqrt{2})(p\overline{p} + n\overline{n})$, and not the states $p\overline{p}$ and $n\overline{n}$, is due to the fact that the isotopic non-invariant mass splitting of the p and n quarks is much smaller than the amplitude for the annihilation transition $p\overline{p} \to n\overline{n}$.



a b

FIG. 17

Zweig's rule is violated much more strongly for the pseudoscalar mesons than for the vector mesons. This is indicated by the strong mass splitting between the η^0 and π^0 mesons.

f) **Zweig's rule and colored gluons.** In the framework of the quark model, one might attempt to provide a qualitative explanation of the nature of Zweig's rule in the following way. The annihilation of a quark-antiquark pair should lead to gluons. Let us consider the annihilation of a pair in a white $^3S_1$ state. Such a pair cannot be converted into a single colored gluon because of the conservation of color. (Such a transition would be allowed in the case of white gluons, which we have rejected.) It cannot be converted into two gluons because of the conservation of charge parity and color. Consequently, the minimum number of gluons into which such a pair can be converted is three. By assuming that the gluon-quark interaction constant is less than unity ($\alpha_S = g^2 \sim 1/2$) and calculating the probability of the decay $\varphi \to 3\gamma_S$, where γ_S denotes a gluon, we can obtain the observed annihilation width of the φ meson ($\Gamma(\varphi \to 3\pi) \sim$ 0.7 MeV). In the spirit of parton calculations, the second step of the process (in this case, the conversion of three gluons into three pions) is not considered here, with the assumption that the probability for the process is determined mainly by the first step, i.e., $\Gamma(\varphi \to 3\gamma_S) \approx \Gamma(\varphi \to 3\pi)$. The calculation is based on the fact that the $\lambda\overline{\lambda}$ system is similar to orthopositronium, and the role of the constant $\alpha = 1/137$ is played by the quantity α_S. If it is assumed that the lowest level of the $\lambda\overline{\lambda}$ pair, corresponding to the φ meson, falls within the narrow "Coulomb" part of the exchange potential, the formulas for three-photon annihilation of positronium can be used (for massless gluons) to calculate the annihilation $\varphi \to 3\gamma_S$. Calculating the width of the decay $\varphi \to \gamma \to e^+e^-$ in a similar way, we find the ratio

$$\frac{\Gamma(\varphi \to e^+e^-)}{\Gamma(\varphi \to 3\pi)} = \frac{\Gamma(\varphi \to \gamma \to e^+e^-)}{\Gamma(\varphi \to 3\gamma_S)} = \frac{18\pi}{5(\pi^2-9)}\,\frac{\alpha^2}{\alpha_S^3} \approx 10\,\frac{\alpha^2}{\alpha_S^3}.$$

For $\alpha_S \approx 1/2$, this gives $\sim 4 \times 10^{-3}$, which is to be compared with the experimental ratio $\sim 2 \times 10^{-3}$.

In the case of pseudoscalar mesons, two-gluon annihilation of the $q\overline{q}$ pair in the $^1S_0$ state is allowed by the conservation of C-parity and color. The annihilation effects are therefore larger in this case. This might provide a partial explanation of the large mass splitting between the η and π^0 mesons and the large admixture of λ quarks in the η meson.

Similar estimates have been made for the decays of the ψ meson. In this case, using the expression for α_S which is obtained in the asymptotically free theory, it is assumed that $\alpha_S(m_\psi) \approx 0.2$. This makes the ratio of the leptonic and hadronic widths of the ψ meson about an order of magnitude larger than the same ratio for the φ meson and brings this ratio closer to the experimental value.

It should be noted, however, that there is no basis for using the value $\alpha_S(m_\psi) \approx 0.2$ to calculate the probability for decay into free gluons with $m_{\lambda_S} = 0$.

g) **The ψ' meson.** In the framework of the quark model, it is natural to assume that the ψ' is a radially excited $^3S_1$ state with the radial quantum number $n_r = 1$.

Adopting the "charmonium" model, the $\psi' \to e^+e^-$ decay width should be proportional to $|\psi(0)|^2$, the probability of finding the quark and the antiquark at the center of the charmonium. For a Coulomb potential, we have $|\psi_{n_r}(0)|^2 = 1/\pi(n_r + 1)^3$ and hence

$$\frac{\Gamma(\psi' \to e^+e^-)}{\Gamma(\psi \to e^+e^-)} \sim \frac{1}{8}.$$

For an oscillator potential,

$$|\psi_{n_r}(0)|^2 = \frac{(2n_r+1)!!}{\pi^{3/2}2^{n_r}n_r!},$$

where the lowest level corresponds to $n_r = 0$ and the next radially excited level has $n_r = 1$. Hence

$$\frac{\Gamma(\psi' \to e^+e^-)}{\Gamma(\psi \to e^+e^-)} \sim \frac{3}{2}.$$

Experimentally, this ratio is of order 1/2 – 1/3, so that, if we are discussing a potential, this potential is apparently intermediate between the Coulomb and oscillator potentials (see Fig. 16c). If all hadronic decays of the ψ' proceeded via $c\bar{c}$ annihilation at small distances, the charmonium model would require that the hadronic width of the ψ' is $\lesssim 30$ keV, which is about an order of magnitude smaller than the experimental value.

The decay $\psi' \to \psi + 2\pi$, which accounts for about half of all hadronic decays of the ψ' meson, can occur even when the c and $\bar{c}$ quarks are far from one another, since the c and $\bar{c}$ quarks are conserved in this decay. This decay is forbidden to a lesser degree than the hadronic decays of the ψ meson; it can proceed via two-gluon exchange (Fig. 18), the gluons in this process having relatively low energies and hence a stronger interaction.

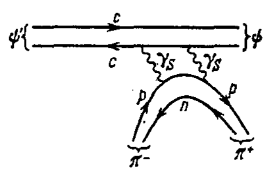


FIG. 18

It is not yet clear what other decay channels the ψ' meson has (apart from the channel $\psi' \to \psi + 2\pi$). As we have already mentioned, the decays of the ψ' into ordinary pions and kaons must be weak in the charmonium model It is therefore mainly the following two possibilities that are discussed: 1) radiative decays, and 2) decays involving the emission of charmed mesons.

h) **Other levels of the charmonium.** The possibility of radiative decays is suggested by an analysis of the levels of the charmonium corresponding to other spin and orbital states of the $c\bar{c}$ pair.

Let us begin by discussing the problem of paracharmonium, i.e., the $^1S_0$ state. Making a comparison of the known vector and pseudoscalar mesons, ρ-π, K^*-K, ω-η and φ-η', the mass of the ground state of paracharmonium (which we denote by η_c) might be expected to be smaller than the mass of the ψ. If the mass difference $m_\psi - m_{\eta_c}$ is small (a value ~100 MeV seems natural), the decay $\psi \to \eta_c\gamma$ will be suppressed by the small value of phase space. The decay $\psi' \to \eta_c\gamma$ can be suppressed only by the orthogonality of the radial ψ-functions for the states with $n_r = 1$ and $n_r = 0$. Calculations show that its expected width is of the order of several keV. Searches for photons of energy of order 700 MeV among the decay products of the ψ', which are now in progress, are of very great interest. The ratio of the decays $\eta_c \to 2\gamma$ and $\eta_c \to$ hadrons should be of order $(\alpha/\alpha_S)^2 \sim 10^{-3}$. However, this estimate may be grossly incorrect.

It is well known that the pseudoscalar η and η' mesons correspond to quark states in which the annihilation mixing $\lambda\bar{\lambda} \longleftrightarrow n\bar{n} \longleftrightarrow p\bar{p} \longleftrightarrow \lambda\bar{\lambda}$ is not small. A large value of the annihilation mixing $n\bar{n} \longleftrightarrow p\bar{p}$ is suggested, in particular, by the large mass difference between the π^0 and η^0 mesons. The large mixing in this case is due to the fact that two-gluon intermediate states are possible for the pseudoscalar mesons, unlike the vector mesons.

In addition, the annihilation mixings $c\bar{c} \longleftrightarrow p\bar{p}$, $c\bar{c} \longleftrightarrow n\bar{n}$ and $c\bar{c} \longleftrightarrow \lambda\bar{\lambda}$ must be small not only for the vector particles, but also for the pseudoscalar particles. This means that the η_c meson must be a practically pure state of the $c\bar{c}$ pair and that the admixture of this pair in the other pseudoscalar mesons η and η' must be negligible, since otherwise strong decays such as $\psi \to \eta\omega^0$ or $\psi \to \eta'\omega^0$ would have to occur with a high probability. This great purity of the η_c state is apparently due to its large mass. However, a thorough understanding of this problem is still lacking.

The other charmonium states correspond to orbital excitations.

Let us recall what the nonrelativistic levels in the Coulomb and oscillator potentials look like:

Oscillator	"Coulomb"
$E_n = \omega(n + 3/2)$,	$E_n = -\frac{1}{2}\frac{m\alpha^2}{n^2}$,
$n = 2n_r + l$	$n = n_r + l + 1$.

Here n is the principal quantum number, n_r is the radial quantum number, and l is the orbital quantum number. The 2S and 2P levels are degenerate in the Coulomb potential. In the oscillator potential, the level with $n_r = 0$ and $l = 1$ is half as high as the level with $n_r = 1$ and $l = 0$. We may therefore expect the P levels of the charmonium to have a mass of order 3.5 GeV.

Electric dipole transitions from the state ψ' (i.e., 2^3S_1) to the P states (and from the latter to the ψ (1^3S_1)) might account for a significant fraction of the decays of the ψ' meson. In this case, photons should be emitted with energy ~200 MeV in the transition $\psi' \to$ P and with energy ~400 MeV in the transition P $\to \psi$.

As to the D state ($l = 2$, $n_r = 0$), this state should occur near the ψ'. It might be produced via the transition $e^+e^- \to \gamma \to {}^3D_1$; however, owing to the centrifugal repulsion, $|\psi(0)|^2$ is much smaller in this state than in the state 2^3S_1, so that the cross section for producing the appropriate meson in colliding e^+e^- beams should also be small.

We emphasize that the probabilities of the above-mentioned radiative transitions between the charmonium levels do not involve $|\psi(0)|^2$ and are not forbidden by Zweig's rule. If monochromatic photons corresponding to radiative transitions between the charmonium levels are not observed, this might be explained by the fact that the charmonium P levels lie above the 2S level. This inversion of the levels with respect to what is predicted by the oscillator model might be due to the existence of a strong short-range attraction which lowers the S levels or a strong spin-orbit interaction which raises the P levels. The transition $\psi' \to \eta_c\gamma$ cannot be eliminated in this way. Searches for this decay are therefore of special interest.

i) Decays of the ψ' into charmed mesons? A detailed discussion of the expected properties of charmed mesons will be given below; for the moment, we note only that the lightest of them (which may be the pseudoscalar mesons $D^0 = c\bar{p}$ and $D^+ = c\bar{n}$) should have masses of order 2 GeV. If these masses are somewhat less than 1.85 GeV = $(1/2)m_{\psi'}$, the decays $\psi' \to D^0\bar{D}^0$ are possible. These decays are not suppressed by Zweig's rule, since they involve conservation of the c and $\bar{c}$ quarks. They are analogous to the decays $\varphi \to K\bar{K}$ and are described by the diagram of Fig. 19. Adopting the quark model, it is easy to obtain the relation between the electromagnetic mass differences in the limit of SU(4) symmetry:

$$m_{D^+} - m_{D^0} = m_{K^0} - m_{K^+} + 2\,(m_{\pi^+} - m_{\pi^0}) \approx 13 \text{ MeV}.$$

This relation takes into account the fact that the charges of the p, n, λ and c quarks are in the ratio 2 : −1 : −1 : 2 and the fact that the mesons have the structure

$$D^+ = c\bar{n},\quad D^0 = c\bar{p},\quad K^0 = n\bar{\lambda},\quad K^+ = p\bar{\lambda},\quad \pi^+ = p\bar{n},\quad \pi^0 = \frac{1}{\sqrt{2}}(p\bar{p} - n\bar{n}).$$

Therefore $m_{D^+} - m_{D^0} = m + 6K$, $m_{K^0} - m_{K^+} = m - 3K$ and $m_{\pi^+} - m_{\pi^0} = (9/2)K$, where m is the mass difference between the p and n quarks, and K is the energy of the Coulomb attraction between the n and $\bar{n}$ quarks. If these relations hold even approximately, the decays $\psi' \to D^0\bar{D}^0$ should be much more probable than the decays $\psi' \to D^+D^-$ because of the large phase space. The D mesons should decay weakly, mainly into several mesons (2–3 pions + 1 kaon, or 3–4 pions) and leptons. Searches for leptons, kaons and peaks in the distributions in the invariant masses of several pions in colliding electron-positron beams should, in the near future, provide an answer to the question of whether the decay $\psi' \to D\bar{D}$ takes place.

j) The 4.15-GeV resonance. These same experiments should also decide whether the 4.15-GeV resonance decays into $D\bar{D}$. The width of this resonance, ~300 MeV, indicates that its decay is not forbidden by any selection rules. Adopting the quark model, it is therefore natural to assume that its mass is much higher than the threshold for producing the pair $c\bar{c}$ (these quarks are, of course, not free, but confined) and higher than the threshold for decay into $D\bar{D}$. It is usually assumed that this resonance corresponds to the $^3S_1$ state with $n_r = 2$. It is then natural to expect that its electron width, which governs its production in e^+e^- collisions and which is proportional to $|\psi(0)|^2$, is approximately the same as for the ψ and ψ' mesons.

3. Quarks and Charmed Particles

The model of four colored quarks predicts the existence of a large number of new particles with non-zero supercharge. However, as SU(4) symmetry is strongly violated, the predicted masses of these particles are somewhat uncertain. From the requirement that the contributions of the p and c quarks cancel in weak processes, we can infer only that these masses should be less than ten GeV. If the ψ mesons are actually $c\bar{c}$ bound states, this fixes the scale of masses of the charmed particles and enables us to predict their masses with far greater accuracy.

a) Vector mesons. In the framework of broken SU(4) symmetry, the ψ meson should belong to a family consisting of 16 vector mesons (Table II). Assuming that the meson masses are equal to the sums of the masses of their constituent quarks, we have

$$m_{\bar{D}^*} = \frac{1}{2}(m_\psi + m_\omega) = 1.94 \text{ GeV},$$
$$m_{F^*} = \frac{1}{2}(m_\psi + m_\varphi) = 2.01 \text{ GeV}.$$

This estimate is obviously highly tentative. In particular, it makes no allowance for the differences between the forces which bind the $c\bar{c}$, $c\bar{p}$ and $p\bar{p}$ systems, for example. Some authors favor quadratic mass formulas. In that case,

$$m^2_{D^*} = \frac{1}{2}(m^2_\psi + m^2_\omega),\quad m_{D^*} = 2.26 \text{ GeV},$$
$$m_{F^*} = \frac{1}{2}(m^2_\psi + m^2_\varphi),\quad m_{F^*} = 2.3 \text{ GeV}.$$

In either case, however, the masses of the charmed vector mesons should be close to 2 GeV.

The ψ' mesons naturally correspond to another family, consisting of 16 heavier vector mesons. At the present time, we have more or less reliable knowledge of only the ρ' mesons, which have $J^P = 1^-$ and mass 1.6 GeV. It would therefore be of great interest to search not only for the supercharged D'^* and F'^* mesons, but also for the superneutral members of this 16-plet: ω', φ', etc.

b) Pseudoscalar mesons. The splitting between the $^1S_0$ and $^3S_1$ levels should evidently lead to masses of the pseudoscalar charmed D and F mesons which, as in the case of the ordinary mesons, are less than the masses of the corresponding vector mesons D^* and F^*, and the mass of the η_c should be less than the mass of the ψ (see Table III for the notation). If there are D mesons among the decay products of the ψ', then $m_D \approx 1.84$ GeV. We can expect the magnitude of the $\eta_c - \psi$ splitting to be smaller than the D-D^* splitting, and the latter to be smaller than the K-K^* splitting (the heavier the quarks, the smaller their gluon "magnetic" moments).

The charmed vector mesons should be heavier than the pseudoscalar mesons in the limit of SU(4) symmetry.

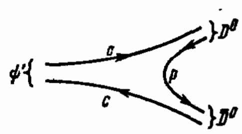


FIG. 19

TABLE II

	$\bar{p}$	$\bar{n}$	$\bar{\lambda}$	$\bar{c}$
p	$\frac{1}{\sqrt{2}}(\omega^0 + \rho^0)$	ρ^+	K^{+*}	$\bar{D}^{0*}$
n	ρ^-	$\frac{1}{\sqrt{2}}(\omega^0 - \rho^0)$	K^{0*}	D^{-*}
λ	K^{-*}	$\bar{K}^{0*}$	φ^0	F^{-*}
c	D^{0*}	D^{+*}	F^{+*}	ψ^0

TABLE III

	$\bar{p}$	$\bar{n}$	$\bar{\lambda}$	$\bar{c}$
p	$\frac{\eta'}{\sqrt{3}}-\frac{\eta}{\sqrt{6}}+\frac{\pi^0}{\sqrt{2}}$	π^+	K^+	$\bar{D}_0$
n	π^-	$\frac{\eta'}{\sqrt{3}}+\frac{\eta}{\sqrt{6}}-\frac{\pi^0}{\sqrt{2}}$	K^0	D^-
λ	K	$\bar{K}^0$	$\frac{\eta'}{\sqrt{3}}-\frac{2\eta}{\sqrt{6}}$	F^-
c	D^0	D^+	F^+	η_c

Considering the violation of this symmetry, however, we cannot exclude the possibility that $m_{F^*} < m_F$, for example. Such an anomalous sign of the "fine-structure" splitting would be an indication of an anomalous sign of the gluon "magnetic" moment of the c quark. We note that $m_\psi > m_{\eta_c}$ must hold even in this case.

c) If $m_{D^*} - m_D > m_\pi$, the strong decay $D^* \to D + \pi$ takes place. If, on the other hand, $m_{D^*} - m_D < m_\pi$, then the main decay of the D^* meson is the radiative decay $D^* \to D + \gamma$. In the case of the F^*, the decay $F^* \to F\pi$ is forbidden by conservation of isospin, so that the decay $F^* \to F\gamma$ should dominate. Similar arguments apply to the fast decays of the pseudoscalar mesons in the case in which they are heavier than the corresponding vector mesons.

d) **Leptonic decays of mesons.** The lightest of the supercharged mesons having given values of charm C, strangeness S and charge Q should be stable, decaying only via weak interactions. The weak supercharge-changing current has the form $(-\bar{n}\sin\theta + \bar{\lambda}\cos\theta)c$. Leptonic decays of the D and F mesons should therefore occur as a result of the interaction

$$\frac{G}{\sqrt{2}}(-\bar{n}c\cos\theta + \bar{\lambda}c\sin\theta)(\bar{\nu}_\mu\mu + \bar{\nu}_e e) + \text{h.c.}$$

and should be similar to the leptonic decays of kaons (K_{l_2}, K_{l_3}, K_{l_4}). The dominant decays should be those whose amplitudes are proportional to $\cos\theta$: $D^+ \to \bar{K}^0\mu^+\nu$, $D^+ \to \bar{K}^0 e^+\nu$, $D^0 \to K^-\mu^+\nu$, $F^+ \to \mu^+\nu\eta$, etc. In addition to two- and three-body decays, there should also be decays involving two or more mesons in the final state, such as $D^+ \to K^-\pi^+e^+\nu$.

Of the weak decays of the vector mesons, two-body decays such as $F^* \to \mu\nu$, which are not suppressed by conservation of helicity, would be of the greatest interest.

By taking into account the violation of SU(4) symmetry only in the phase space, we can make a crude estimate of the widths of decays such as D_{l_3}. For example,

$$\Gamma(D^+ \to \bar{K}^0 e^+\nu) \approx \left(\frac{m_D}{m_K}\right)^5 2\,\mathrm{ctg}^2\theta\,\Gamma(K^+ \to \pi^0 e^+\nu) \approx 2\cdot 10^3\cdot 20\cdot 4\cdot 10^6 \approx 1.6\cdot 10^{11}\ \mathrm{sec}^{-1}.$$

There have so far not been any more precise estimates of the absolute widths of the leptonic decays with allowance for form factors. Estimates of the probabilities for decays of the type D_{l_4} have not yet been made.

However, a knowledge of the weak current enables us to predict certain symmetry relations among the amplitudes for leptonic decays. The lepton amplitudes $\sim\cos\theta$ must satisfy the selection rules

$$\Delta C = \Delta S = -1,\quad \Delta T = 0,\quad \Delta U = 1/2,\quad \Delta V = 1/2;$$

here C is the charm (the value of C is equal to the number of c quarks in a particle), S is the strangeness, and the projections of the T, U and V spins of the quarks are specified in Table IV. The values of the T, U and V spins of the D and F mesons are determined by their quark structure: $D^0 = c\bar{p}$, $D^+ = c\bar{n}$, $F^+ = c\bar{\lambda}$. The selection rule $\Delta T = 0$ implies, for example, that

$$\Gamma(D^0 \to K^-\mu^+\nu) = \Gamma(D^+ \to \bar{K}^0\mu^+\nu).$$

The lepton amplitudes $\sim\sin\theta$ must satisfy the selection rules

$$\Delta C = \pm 1,\quad \Delta S = 0,\quad \Delta T = \frac{1}{2},\quad \Delta U = \frac{1}{2},\quad \Delta V = 0.$$

e) **Non-leptonic decays of mesons.** The non-leptonic decays of the D and F mesons are governed by products of the currents $(\bar{p}n_\theta)$ and $(\bar{\lambda}_\theta c)$. The individual terms of a product with $\Delta C = -1$ must satisfy the selection rules indicated in Table V. The meaning of the numbers which are underlined will be explained later. The selection rules shown here follow directly from the quark structure of the corresponding terms. For example, the selection rule $\Delta U = 1$ for the term $\sim\sin\theta\cos\theta$ follows from the fact that the difference $\lambda\bar{\lambda} - n\bar{n}$ is a component of a U-triplet. The minus sign in this case is due to the fact that the p and c quarks are associated with the orthogonal combinations n_θ and λ_θ in the charged current. We quote several relations which follow from the selection rules $\Delta T = 1$ and $\Delta U = 1$:

$$\Delta T = 1: (K^-\pi^+)_{D^0} + \sqrt{2}(\bar{K}^0\pi^0)_{D^0} = (\bar{K}^0\pi^+)_{D^+},\quad (\pi^+\pi^0)_{F^+} = 0,$$

$$\Delta U = 1: (K^+\bar{K}^0)_{F^+} - \sqrt{\frac{3}{2}}(\pi^+\eta^0)_{F^+} = (\bar{K}^0\pi^+)_{D^+},$$

$$(\bar{K}^0\pi^0)_{D^0} = \sqrt{3}(\bar{K}^0\eta^0)_{D^0};$$

here $(K^-\pi^+)_{D^0}$, for example, denotes the $D^0 \to K^-\pi^+$ decay amplitude.

In order to gain a better understanding of the expected properties of the non-leptonic decays of charmed particles, we shall consider the known properties of the non-leptonic decays of strange particles.

TABLE IV

	p	n	λ	c
T_3	$+\frac{1}{2}$	$-\frac{1}{2}$	0	0
U_3	0	$+\frac{1}{2}$	$-\frac{1}{2}$	0
V_3	$-\frac{1}{2}$	0	$+\frac{1}{2}$	0

TABLE V

	$\cos^2\theta\,(\bar{p}n)(\bar{\lambda}c)$	$\cos\theta\sin\theta\,[(\bar{p}\lambda)(\bar{\lambda}c) - (\bar{p}n)(\bar{n}c)]$	$-\sin^2\theta\,(\bar{p}\lambda)(\bar{n}c)$
ΔS	-1	0	$+1$
ΔT	1	$\underline{\frac{1}{2}}, \frac{3}{2}$	$\underline{0}$, 1
ΔU	1	1	1
ΔV	$\underline{0}$; 1	$\underline{\frac{1}{2}}, \frac{3}{2}$	1

f) **The $\Delta T = 1/2$ rule and octet enhancement.** It is well known that the amplitudes for non-leptonic decays of strange particles are enhanced with respect to the values which they would have if they were products of matrix elements of currents taken from the corresponding leptonic decays. For instance, the $K_S \to \pi^+\pi^-$ decay amplitude is more than an order of magnitude greater than the product of the current factors taken from the $K^0 \to \pi^- e^+ \nu$ and $\pi^+ \to \mu^+ \nu$ decay amplitudes. There is then an enhancement of the amplitudes which satisfy the selection rule $\Delta T = 1/2$ and which are components of an octet. But the amplitudes with $\Delta T = 3/2$ are not weakened. The product of the currents $(n\bar{p})$ and $(p\bar{\lambda})$ gives the non-leptonic Lagrangian $(\bar{n}p)(\bar{p}\lambda)$, which involves transitions with both $\Delta T = 1/2$ and $\Delta T = 3/2$. For certain decays, partial conservation of the axial-vector current leads to a suppression of the transitions with $\Delta T = 3/2$. However, this mechanism is neither univeral nor sufficiently strong. Several ways of achieving a universal dynamical enhancement of the amplitudes with $\Delta T = 1/2$ can be imagined.

The first possibility is the annihilation of p and $\bar{p}$ at small distances, which leads to an effective transition $\lambda \longleftrightarrow n$ and is described by the proton loop shown in Fig. 20a. For $-q^2 \gg m_p^2$, the contribution of this diagram has the form

$$\frac{G}{\sqrt{2}}\,\bar{\lambda}\gamma_\alpha(1+\gamma_5)\lambda_l n\,\frac{g_S q^2}{3(4\pi)^{3/2}}\ln\frac{m_W^2}{-q^2}$$

and represents a sum of monopole and anapole gluon moments; here m_W and m_p are the masses of the W boson and the proton quark, q is the 4-momentum of the gluon, and g_S is the gluon-quark interaction constant, with $g_S^2 = \alpha_S$.

In the four-quark model, the loops involving the p quark (Fig. 20a) and the c quark (Fig. 20b) must cancel with one another in the limit of exact SU(4) symmetry. Their total contribution is proportional to $\ln(m_c^2/q^2)$ for $m_p^2 \ll -q^2 \ll m_c^2$ and tends to zero for $q^2 \gg m_c^2$. If $\alpha_S \lesssim 1$, the effective weak constant of these monoanapole moments is $\sim 0.1G$, and it is not yet clear whether the amplitudes associated with them can be greater than the amplitudes due to a simple four-fermion product of weak quark currents.

Another possibility is an enhancement of the expression

$$(\bar{n}p)(\bar{p}\lambda) - \frac{1}{\sqrt{2}}(\bar{n}n)(\bar{n}\lambda),$$

which has $\Delta T = 1/2$. This expression contains the term $(\bar{n}n)(\bar{n}\lambda)$, which might appear in a more or less natural way from the initial term $(\bar{n}p)(\bar{p}\lambda)$ if there were isovector gluons. In the case of the symmetry $G_S \times G_W$, the gluons are isoscalar, and the term $(\bar{n}n)(\bar{n}\lambda)$ appears only as a result of the annihilation $\bar{p}p \to$ gluons $\to \bar{n}n$ and does not have the required form of a product of V-A currents. In the framework of gauge models of the weak interaction, the introduction of the neutral strangeness-changing current $\bar{n}\lambda$ in the initial Lagrangian leads to difficulties in connection with decays of the type $K_L \to \mu^+\mu$ and the mass difference between the K_L and K_S mesons, which we discussed earlier, and is therefore unsatisfactory.

A third possibility is an enhancement of the states which are antisymmetric with respect to the interchange of $\bar{n}$ and $\bar{p}$ in the expression $(\bar{n}p)(\bar{p}\lambda)$, giving the combination

$$(\bar{n}p)(\bar{p}\lambda) - (\bar{p}p)(\bar{n}\lambda).$$

The fact that this expression gives $\Delta T = 1/2$ can be seen from the fact that the initial state $p\lambda$ has $T = 1/2$, while the final state $np - pn$ has $T = 0$. As we shall see below, the model of colored quarks and colored gluons involves a dynamical enhancement of the antisymmetric amplitudes belonging to the octet representation of SU(3) and a weakening of the symmetric amplitudes belonging to the 27-dimensional representation.

g) **The role of colored quarks in the octet enhancement.** It is well known that the four-fermion V-A amplitude is antisymmetric with respect to interchange of the spins and momenta of the two initial (or final) particles. By virtue of the generalized Pauli principle, this implies that it must be symmetric with respect to the interchange of all the internal degrees of freedom. For white quarks, it is therefore not possible to achieve octet enhancement by antisymmetrization in the local limit. However, this becomes possible for colored quarks if we assume that there is also antisymmetry in the color indices:

$$[(\bar{n}^i p_{i'})(\bar{p}^k \lambda_{k'}) - (\bar{p}^i p_{i'})(\bar{n}^k \lambda_{k'})]\,\varepsilon_{ikl}\varepsilon^{i'k'l'}.$$

This expression is symmetric with respect to the simultaneous interchange of the isotopic ($\bar{n} \longleftrightarrow \bar{p}$) and color ($i \longleftrightarrow k$) variables. It is very significant that this expression is a color invariant and hence gives transitions between white particles. In essence, color plays the same role here as in reconciling the spin and statistics of quarks in constructing baryons from them.

h) **The role of colored gluons in the octet enhancement.** If colored quarks enable us to enhance the antisymmetric amplitudes with $\Delta T = 1/2$, colored gluons provide a dynamical realization of this possibility.

It can be inferred from calculations of diagrams such as that shown in Fig. 21 and more complex diagrams, which allow for the exchange of an arbitrary number of gluons, that the transition amplitudes with $\Delta T = 3/2$ are weakened by a factor ϵ, while those with $\Delta T = 1/2$ are enhanced by a factor ϵ^2, where

$$\varepsilon = \left(\frac{\alpha_S(\mu)}{\alpha_S(m_W)}\right)^{6/25};$$

here m_W is the mass of the W boson, μ is some characteristic hadron mass $\lesssim 1$ GeV, and α_S is the gluon analogue of $\alpha = 1/137$ in electrodynamics. The value of α_S depends on the magnitude of the momentum transfer, since, as we have already mentioned, the interaction becomes weaker at small distances in an asymptotically free theory. In the case of SU(3) symmetry,

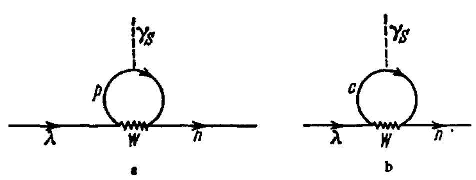


FIG. 20

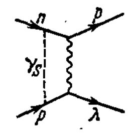


FIG. 21

$$\alpha_S(m_1)=\frac{\alpha_S(m_2)}{1+(25/6\pi)\,\alpha_S(m_2)\ln(m_1/m_2)}.$$

By choosing, for example, $\alpha_S \sim 1$ with $\mu \sim 1$ GeV, we can obtain $\epsilon^3 \sim 5$. Experimentally, the amplitudes with $\Delta T = 3/2$ amount to a few percent of the amplitudes with $\Delta T = 1/2$. Thus, the relative enhancement with $\epsilon^3 \sim 5$ is insufficient. Perhaps the experimentally observed effect has several components: that which we have just discussed, the above-mentioned consequence of partial conservation of the axial-vector current, and transitions $\lambda \longleftrightarrow n$.

i) **The enhancement in SU(4).** In going over from three to four colored quarks, we can provide an SU(4)-symmetric generalization of the octet-enhancement mechanism discussed above.

Consider the product of the two currents $\bar{q}^{\alpha}q_{\beta}$ and $\bar{q}^{\gamma}q_{\delta}$, where $\alpha, \beta, \gamma, \delta = 1, 2, 3, 4$. It is easy to see that the antisymmetric tensor $T^{[\alpha,\gamma]}_{[\beta,\delta]}$ obtained from the product of these currents belongs to a 20-plet. In fact, we have 6 antisymmetric combinations both above and below. But the tensor obtained by antisymmetrization is reducible:

$$6\times\bar{6}=36=1+15+20.$$

All four quarks are different in the product $(\bar{p}n_{\theta})(\lambda_{\theta}c)$. Therefore this product cannot belong to either the 1 or the 15, but can be contained only in the 20. We thus arrive at the conclusion that the 20-plet representation of SU(4) should be enhanced.

Let us now see which SU(3) representations are contained in the 20-plet $T^{[\alpha\gamma]}_{[\beta\delta]}$. The quantity $[\alpha\gamma]$ transforms like $3+\bar{3}$; here $\alpha = 4$ and $\gamma = 1, 2, 3$ corresponds to the representation 3, while $\alpha = 1, 2, 3$ and $\gamma = 1, 2, 3$ corresponds to the representation $\bar{3}$. Hence

$$6\times\bar{6}=(3+\bar{3})(\bar{3}+3)=1+(1+8+3+\bar{3})+(8+6+\bar{6}).$$

Thus, $20 = 8 + 6 + \bar{6}$. Here the octet does not contain the c quark, the sectet contains the $\bar{c}$, and the antisextet contains the c.

We see that enhancement of the octet amplitudes in strange-particle decays should correspond to enhancement of the sextet amplitudes in charmed-particle decays.

j) **T-, U- and V-spin selection rules.** With sextet enhancement, the operators $\bar{p}$ and $\bar{\lambda}_{\theta}$ in the expression $(\bar{p}n_{\theta})(\bar{\lambda}_{\theta}c)$ are antisymmetrized. Now $[\bar{p}\bar{\lambda}]$ has $V = 0$, while $[\bar{p}\bar{n}]$ has $T = 0$. Consequently, the term $\cos^2\theta(\bar{p}n)(\bar{\lambda}c)$ must obey the selection rule $\Delta V = 0$, the term $\sin^2\theta(\bar{p}\lambda)(\bar{n}c)$ must obey $\Delta T = 0$, and the term $\cos\theta\sin\theta[(\bar{p}\lambda)(\bar{\lambda}c) - (\bar{p}n)(\bar{n}c)]$ must obey $\Delta T = 1/2$ and $\Delta V = 1/2$. We note that $\Delta T = 1/2$ always holds for the term $(\bar{p}\lambda)(\bar{\lambda}c)$, but is obtained for the term $(\bar{p}n)(\bar{n}c)$ only after antisymmetrization of $[\bar{p}\bar{n}]$. Similarly, $\Delta V = 1/2$ always holds for the term $(\bar{p}n)(\bar{n}c)$, but is obtained for the term $(\bar{p}\lambda)(\bar{\lambda}c)$ only after antisymetrization of $[\bar{p}\bar{\lambda}]$.

The values of ΔT, ΔU and ΔV which follow from sextet enhancement are underlined in the table of selection rules for the non-leptonic decays of charmed particles which we gave earlier (see Sec. 3e).

The selection rule $\Delta V = 0$ implies that the decay $D^+ \to K^0\pi^+$ is forbidden. In fact, the D^+ is a V-singlet, while the $\bar{K}^0$ and π^+ are components of a V-doublet. The final state with $V = 0$ has the form $(\bar{K}^0\pi^+ - \pi^+\bar{K}^0)$. But this state is antisymmetric with respect to the interchange $\bar{K}^0 \longleftrightarrow \pi^+$ and is therefore forbidden for bosons in a state with zero orbital angular momentum. Putting $\Gamma(D^+ \to \bar{K}^0\pi^+) = 0$ in the relations obtained earlier from the selection rules $\Delta T = 1$ and $\Delta U = 1$, we get

$$\Gamma(D^0\to K^-\pi^+)=2\Gamma(D^0\to\bar{K}^0\pi^0),$$
$$2\Gamma(F^+\to K^+\bar{K}^0)=3\Gamma(F^+\to\pi^+\eta^0).$$

Many analogous relations can be derived for the other decay channels of the D and F mesons and for the decays of charmed baryons.

k) **The transition c $\longleftrightarrow$ p.** If the non-leptonic decays of charmed hadrons were due to the interaction $\bar{c}p + \bar{p}c$, these decays would obey the selection rules

$$\Delta S=0,\quad \Delta T=1/2,\quad \Delta U=0,\quad \overline{\Delta V=1/2}.$$

The corresponding amplitudes would be components of an SU(3) triplet and an SU(4) 15-plet. In this case, the D meson would decay into pions, but not into a $\bar{K}$ meson $\pm$ pions. If it were found that there are few kaons among the decay products of the resonance at 4.15 GeV, the interaction $\bar{c}p + \bar{p}c$ would attract attention. However, this interaction has a number of serious theoretical defects.

In the Weinberg model, the transition $c \longleftrightarrow p$ is due to the sum of the diagrams of Fig. 22. The total contribution of these diagrams vanishes in the limit of exact SU(3) symmetry. Therefore this transition should be weaker than the analogous transition $n \longleftrightarrow \lambda$. This circumstance renders the mechanism involving the transition $c \longleftrightarrow p$ implausible within the framework of the Weinberg model.

l) **The expected decay widths of the D and F mesons.** It is not yet possible to make reliable predictions of the partial widths of the D and F mesons. The total width of the semi-leptonic channels can be estimated by taking the muon width $\Gamma_{\mu} = (1/2)\times 10^{+6}$ sec^{-1} and increasing it by a factor $2(m_c/m_{\mu})^5$, where m_c is the mass of the charmed meson (or quark?). For $m_c \sim 2$ GeV, we obtain $\Gamma_{\text{semilept}} \approx 3\times 10^{12}$ sec^{-1}. It is assumed in this estimate that the c quark decays like a point particle and that the strong interactions redistribute the decay products but do not alter the total decay probability.

It is more difficult to predict the widths of the non-leptonic decays. These decays should be enhanced with respect to the semi-leptonic decays (we recall the sextet enhancement and the large ratio, R = 2.5, of the cross sections for annihilation of e^+e^- into hadrons and into $\mu^+\mu^-$). However, it should be noted that the sextet enhancement should be much weaker than the octet enhancement, since charmed particles are much heavier than strange particles (if allowance is made for the dependence of α_S on the characteristic momentum transfers, the sextet enhancement gives a net factor of order two). Considering all the uncertainties, we can say that

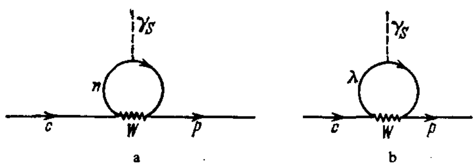


FIG. 22

the fraction of semi-leptonic decays apparently lies somewhere in the range between 1 and 50%.

The non-leptonic decays should be dominated by multi-particle channels with $\bar{n} \sim 4$. Such a multiplicity ($\bar{n}$ = 4, 5) is characteristic of the products of e^+e^- and $\bar{p}p$ annihilation at center-of-mass energies of order 2 GeV.

m) **The transitions $D^0 \longleftrightarrow \bar{D}^0$.** Since the weak interaction does not conserve charm, the transitions $(D^0 = c\bar{p}) \longleftrightarrow (\bar{D}^0 = \bar{c}p)$, which are analogous to the well-known transitions $(K^0 = n\bar{\lambda}) \longleftrightarrow (\bar{K}^0 = \lambda\bar{n})$, should occur in a vacuum. Because of these transitions, there should exist two states having definite masses and lifetimes—a long-lived state D_L^0 and a short-lived state D_S^0. In contrast with what happens in the case of the K_L and K_S mesons, the D_L and D_S mesons are both short-lived: $\tau_L = 1/\Gamma_L \sim \tau_S = 1/\Gamma_S \approx 10^{-13} - 10^{-14}$ sec, and $\Gamma_S - \Gamma_L \ll \Gamma_S + \Gamma_L$. This last statement is based on the fact that only decays of the D^0 and $\bar{D}^0$ into the non-strange hadronic states $2\pi, 3\pi, 4\pi, \ldots$ can interfere with one another. But these decays have small amplitudes $\sim \cos\theta \sin\theta$. The amplitudes $\sim \cos^2\theta$ of the principal decay modes such as $D^0 \to K^-\pi^+\pi^0$ and $\bar{D}^0 \to K^+\pi^-\pi^0$ do not interfere with one another and give identical contributions to the decays of D_L and D_S. (For the moment, we are neglecting possible effects of CP violation and assuming that

$$D_L \approx D_1^0 \equiv \frac{1}{\sqrt{2}}(D^0 - \bar{D}^0), \qquad D_S \approx D_2^0 \equiv \frac{1}{\sqrt{2}}(D^0 + \bar{D}^0),$$

where D_1^0 and D_2^0 have positive and negative CP-parity, respectively.) Consequently, the fact that the decays $D_2 \to 2\pi$ are forbidden has very little effect on the total widths:

$$\frac{\Gamma_S - \Gamma_L}{\Gamma_S + \Gamma_L} \ll \mathrm{tg}^2\theta \sim 5\%.$$

(We recall that the fact that the decay $K_2 \to 2\pi$ is forbidden is decisive in the case of kaons: $\Gamma_L/\Gamma_S \sim 2 \times 10^{-3}$.) The mass difference between the D_L and D_S mesons is especially small. This follows from the fact that the transitions $c\bar{p} \longleftrightarrow \bar{c}p$ are forbidden in the limit of SU(3) symmetry. It is easiest to see that this is so from the sum of the loops shown in Fig. 23. At each of the vertices in the figure, we have indicated the value of the appropriate interaction constant obtained by multiplying the weak currents. The total contribution of the four diagrams clearly vanishes in the limit of SU(3) symmetry ($m_n = m_\lambda$). It is reasonable to expect that $|m_{D_L} - m_{D_S}| \ll \Gamma_{D_S}$.

Owing to the short lifetimes of the D mesons, the experimental observation of oscillation phenomena in their decays would be difficult. However, it would be relatively easy to observe the effects which survive after an integration with respect to time.

As an example, owing to the transitions $D^0 \longleftrightarrow \bar{D}^0$, an initial D^0 meson would decay into the channel $K^+\mu^-\bar{\nu}_\mu$ in some small fraction of the events. It is readily shown that

$$\frac{N(D^0 \to K^+\mu^-\bar{\nu}_\mu)}{N(D^0 \to K^-\mu^+\nu_\mu) + N(D^0 \to K^+\mu^-\bar{\nu}_\mu)} = \frac{1}{2}\,\frac{(\Gamma_S - \Gamma_L)^2 + 4(m_S - m_L)^2}{(\Gamma_S + \Gamma_L)^2 + 4(m_S - m_L)^2}.$$

In the case of non-leptonic decays of D^0 mesons, final states having an anomalous sign of the strangeness appear not only because of the transitions $D^0 \longleftrightarrow \bar{D}^0$, but also because of the small term $\sin^2\theta(\bar{n}c)(\bar{p}\lambda)$ in the non-leptonic Lagrangian. For example,

$$\frac{N(D^0 \to K^+\pi^-)}{N(D^0 \to K^-\pi^+)} \approx \sin^4\theta + \frac{m_S - m_L}{\Gamma_S + \Gamma_L}\sin^2\theta + 2\left(\frac{m_S - m_L}{\Gamma_S + \Gamma_L}\right)^2 + \frac{1}{2}\left(\frac{\Gamma_S - \Gamma_L}{\Gamma_S + \Gamma_L}\right)^2;$$

in this result, we have put $\cos^2\theta \sim 1$ and neglected higher powers of the ratio $(m_S - m_L)/(\Gamma_S + \Gamma_L)$.

n) **D mesons and CP violation.** The D^0-$\bar{D}^0$ system is a beautiful object for exhibiting the effects of CP violation, even if this violation occurs in the super-weak interaction. Allowing for CP violation, we have

$$D_S = D_1 + \varepsilon_D D_2, \quad D_L = D_2 + \varepsilon_D D_1.$$

This leads to a number of interesting effects. For example, in the decay of a resonance such as the ψ' into a pair $D^0 + \bar{D}^0$, there should be a charge asymmetry of the decay products of the D mesons.

Consider the decays into $K^+\mu^-\bar{\nu}$ and $K^-\mu^+\nu$. Denoting the numbers of pairs $(K^+\mu^-\bar{\nu}, K^+\mu^-\bar{\nu})$ and $(K^-\mu^+\nu, K^-\mu^+\nu)$ by N^{++} and N^{--}, the charge asymmetry is given by

$$\delta_D = \frac{N^{++} - N^{--}}{N^{++} + N^{--}} = \frac{4\,\mathrm{Re}\,\varepsilon_D(1 + |\varepsilon_D|^2)}{(1 + |\varepsilon_D|^2)^2 - 4(\mathrm{Re}\,\varepsilon_D)^2} \approx 4\,\mathrm{Re}\,\varepsilon_D.$$

Non-leptonic decays of pairs of D mesons (such as $K^+\pi$ and $K^-\pi$) should also exhibit a charge asymmetry, which is somewhat weaker in this case because of the direct decays of the type $D \to K^+\pi^-$, whose probability is proportional to $\sin^4\theta$.

The discovery of CP-noninvariant effects somewhere other than in the decays of K^0 mesons, where they have so far been observed, might provide the key to finding the mechanism of CP violation.

o) **Baryon SU(4) multiplets.** It is convenient to arrange the four quarks at the vertices of a tetrahedron (Fig. 24).

This emphasizes the symmetry between the quarks. It is also convenient to arrange the baryons built up from the quarks in the form of a tetrahedron (Fig. 25).

The four quarks can be used to construct $4^3 = 64$ three-quark states. These 64 states are divided into four distinct SU(4) multiplets:

$$64 = 4 + 20 + 20 + 20.$$

Here 4 states are completely antisymmetric: $[pn\lambda]$, $[pnc]$, $[pc\lambda]$ and $[cn\lambda]$; 20 states are completely symmetric, and two 20-plets have mixed symmetry.

To determine the number of symmetric states, we observe that they may be of three types: {aaa}, where all the quarks are identical; {aab}, where two quarks are identical; and {abc}, where all the quarks are distinct. There are four combinations {aaa}, twelve (4 × 3)

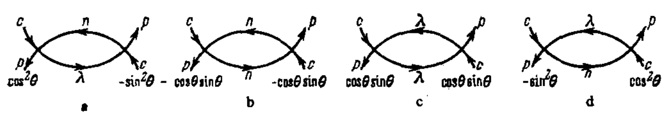


FIG. 23

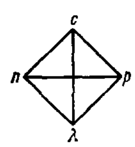


FIG. 24

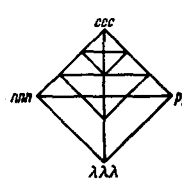


FIG. 25

combinations {aab}, and four combinations {abc}. There are 20 particles in all. The well-known SU(3) decuplet with $J^P = 3/2^+$ should belong to such a 20-plet. The two SU(4) 20-plets of mixed symmetry each contain 12 combinations of the type aab and 8 combinations of the type abc. One of these 20-plets contains the well-known SU(3) octet with $J^P = 1/2^+$.

In order to gain a better understanding of the structure of the baryons, let us enumerate once more all 64 three-quark combinations which lie on the tetrahedron. At the four vertices of the tetrahedron, there are combinations of the type aaa. On the six edges of the tetrahedron, there are 36 combinations of the type aab. Finally, at the four centers of the faces, there are 24 combinations of the type abc. In all, there are 4 + 36 + 24 = 64 particles on the tetrahedron.

p) SU(3) multiplets of charmed baryons. The pyramid of 64 baryon states can be divided in a natural way into SU(3) multiplets having definite values of supercharge (charm). At the base of the pyramid, there are 27 superneutral particles, of which 18 belong to the well-known SU(3) baryon multiplets: the octet with $J^P = 1/2^+$ (Fig. 26a) and the decuplet with $J^P = 3/2^+$ (Fig. 26b), corresponding to the lowest orbital state of the quarks. (The other nine particles—a singlet and an octet—correspond to $l \neq 0$.) At the next level of the pyramid, there are also 27 particles. Each of them contains a single c quark. These particles include a triplet and a sextet with $J^P = 1/2^+$ and a sextet with $J^P = 3/2^+$, corresponding to the lowest orbital state of the quarks. (The other twelve particles—a sextet and two triplets—correspond to $l \neq 0$).

Still higher, there are 9 particles, each containing two c quarks: a triplet with $J^P = 1/2^+$, a triplet with $J^P = 3/2^+$, and a triplet with $l \neq 0$.

Finally, at the top of the pyramid, there is a single particle ccc with $J^P = 3/2^+$.

q) Notation for the baryons. There is as yet no generally accepted notation for the charmed baryons. We shall follow the nomenclature used in the review of Gaillard, Lee and Rosner. The triplet ($\bar{3}$) and sextet (6) of baryons containing a single c quark are shown in Figs. 27a and 27b. The upper indices indicate the charges of the baryons, and the lower indices (in the case of the C baryons) indicate their isospins. We also indicate the quark structure of each baryon, where {...} denotes symmetrization and [...] denotes antisymmetrization of the appropriate quarks. All the baryons shown in Figs. 27a and 27b can have the values $J^P = 1/2^+$. Because of the symmetry of the quark ψ-function, only the baryons shown in Fig. 27b can have the values $J^P = 3/2^+$. We use the same letter to denote baryons which have the same supercharge, isospin and charge, but different values of the spin and mass. For example, we have the Σ^+ (M = 1189, $J^P = 1/2^+$) and the Σ^+ (M = 1385, $J^P = 3/2^+$).

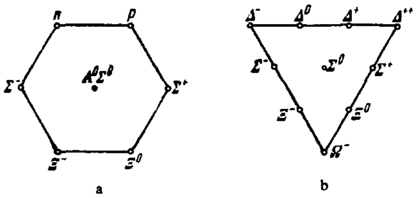


FIG. 26

The baryons which contain two c quarks are shown in Fig. 28. (Gaillard, Lee and Rosner designate them X_u^{++}, X_d^+ and X_s^+, respectively, since they employ a different notation for the quarks: p → u, n → d, λ → s.) The X baryons with $l = 0$ can occur with both $J^P = 1/2^+$ and $J^P = 3/2^+$. In addition, there is a triplet of X baryons with $l \neq 0$. We denote the baryon containing three c quarks by O^{++}. It has $J^P(O^{++}) = 3/2^+$.

r) The baryon masses. The predicted masses of the charmed baryons depend strongly on the particular mass formulas used by various authors. According to some estimates, $M(O^{++}) \sim 10$ GeV. However, it is more natural to expect that $M(O^{++}) \sim 4.5$–5 GeV. This estimate is based on a comparison of the O^{++} baryon consisting of three c quarks with the ψ meson, which consists of the quarks c and $\bar{c}$. The masses of these particles should be in the ratio 3 : 2 if the binding energy per quark is the same for the white baryons and mesons. We are, of course, adopting linear mass formulas here for both the baryons and the mesons. We note that the ratio of the masses of the Ω hyperon ($\lambda\lambda\lambda$) and the φ meson ($\lambda\lambda$) is actually close to 3/2. The same is true of the ratio M_Δ/M_ω. Taking $M(O^{++}) \sim 4.7$ GeV and $M(\Delta) = 1.24$ GeV and assuming an equal-spacing rule for the masses, we can easily estimate the expected masses of the other baryons:

$$M(C) \sim 2.4 \text{ GeV}, \ M(S) \sim M(A) \sim 2.5 \text{ GeV},$$
$$M(T) \sim 2.7 \text{ GeV}, \ M(X_p) = M(X_n) \sim 3.5 \text{ GeV},$$
$$M(X_\lambda) \sim 3.7 \text{ GeV}.$$

Apparently, the lightest of the charmed baryons should be the C_0^+ baryon with $J^P = 1/2^+$—a charmed analogue of the Λ^0 hyperon. According to some estimates, $M(C_0^+) \approx 2.1$ GeV.

V. OTHER MODELS OF THE NEW PARTICLES

In this chapter, we shall consider theoretical descriptions of the ψ mesons other than the charmed-particle hypothesis and the color hypothesis.

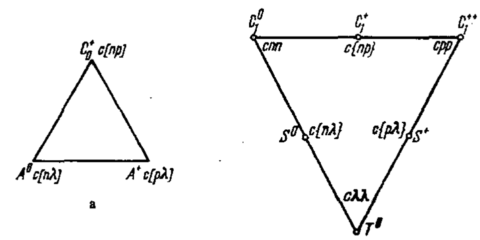


FIG. 27

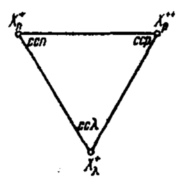


FIG. 28

1. The ψ and ψ' are Intermediate Bosons

The hypothesis that the ψ and ψ' belong to a family of intermediate bosons, i.e., that they are neutral partners of the charged intermediate bosons of the weak interactions, is supported by the data on the widths of the decays of the ψ and ψ' into leptons and hadrons.

If the ψ is an intermediate boson, i.e., if there exists a direct interaction of the ψ with muons and electrons

$$\mathscr{H}_{\psi l}=\sqrt{4\pi}\,g_l\psi_\lambda\,(\bar{u}_e\gamma_\lambda u_e+\bar{u}_\mu\gamma_\lambda u_\mu) \tag{5.1}$$

with an interaction constant $g_e = g_\mu \equiv g_l$ (we are assuming that parity is conserved in this interaction (see Chap. II) and, for simplicity, adopting the hypothesis of μ-e universality), then the $\psi \to$ ee decay width is

$$\Gamma_{ee}=\Gamma_{\mu\mu}=\frac{1}{3}g_l^2M. \tag{5.2}$$

The constant of the four-fermion interaction governed by the exchange of the ψ boson is by definition

$$G=\sqrt{2}\,\frac{4\pi g_l^2}{M^2}=\sqrt{2}\,\frac{12\pi\Gamma_{ee}}{M^3}. \tag{5.3}$$

Substituting in Eq. (5.3) the experimental values of the widths from Table I, we find for the ψ and ψ' bosons

$$G_\psi^l=\frac{0.77\cdot 10^{-5}}{m^2}=0.77G_F,$$
$$G_{\psi'}^l=\frac{0.21\cdot 10^{-5}}{m^2}=0.21G_F, \tag{5.4}$$

where G_F is the universal Fermi weak-interaction constant $G_F = 10^{-5}/m^2$, and m is the nucleon mass. The effective hadron-lepton interaction constants which can be obtained from (5.3) by making the substitution $\Gamma_{ee} \to \sqrt{\Gamma_{ee}\Gamma_h}$ are also rather close to the Fermi constant: $G_{\psi}^{lh} \approx 3G_F$ and $G_{\psi'}^{lh} \approx 1.5\,G_F$.

The fact that these constants are close to G_F suggests that the ψ and ψ' might be intermediate bosons. In that case, the 4.15-GeV resonance, if it exists, would of course have nothing to do with them.

In Sec. 3 of Chap. II we pointed out that, according to the experimental data, the ratios $R_{2\pi^+2\pi^-}$ and $R_{3\pi^+3\pi^-}$ of the cross sections for the processes $e^+e^- \to 2\pi^+2\pi^-$, $3\pi^+3\pi^-$ to the cross section for the process $e^+e^- \to \mu^+\mu^-$ are the same in the region of the ψ meson and in the background in its vicinity. In describing the ψ as a hadron, this fact was explained in a natural way by observing that the isotopic spin of the ψ is equal to zero and that the decay $\psi \to 2\pi^+2\pi^-$ proceeds via a virtual photon. If the intermediate-boson hypothesis is used to account for this fact, it must be assumed that the Hamiltonian describing the interaction of the ψ with hadrons has the form

$$\mathscr{H}_{\psi h}=\sqrt{4\pi}\,(g_l j_{\mu,\,\mathrm{el}}^{h,\,T=1}+g' j_\mu^{h,\,T=0})\,\psi_\mu, \tag{5.5}$$

where the electromagnetic interaction of hadrons is written in the form

$$\mathscr{H}_{\mathrm{el}}=\sqrt{4\pi}\,e\,(j_{\mu,\,\mathrm{el}}^{h,\,T=1}+j_{\mu,\,\mathrm{el}}^{h,\,T=0})\,A_\mu; \tag{5.6}$$

here g_l is the same constant as in (5.1), $j_{\mu,\mathrm{el}}^{h,T=1}$ is the component of the hadronic electromagnetic current corresponding to isospin T = 1, and $j_\mu^{h,T=0}$ is the hadronic current with T = 0 (not necessarily the same as $j_{\mu,\mathrm{el}}^{h,T=0}$, the component of the hadronic electromagnetic current with T = 0).

It is an obvious consequence of (5.1), (5.5) and (5.6) that $R_{H,T=1}$ is the same in the region of the ψ resonance and in the background region for all hadronic states with T = 1. If $R_{H,T=0}$ is to be much greater in the region of the ψ meson than in the background, we require that $|g'| \gg |g|$ and/or $\Sigma_n|\langle 0|j_\mu^{h,T=0}|n\rangle|^2 \gg \Sigma_n|\langle 0|j_{\mu,\mathrm{el}}^{h,T=0}|n\rangle|^2$.

It is known from the decay $\psi' \to \psi\pi\pi$ that the $\psi'\psi\pi\pi$ interaction is characterized by a constant of order unity, i.e., it is strong (or medium-strong; see Chap II). A description of the ψ and ψ' as intermediate bosons is therefore possible only if it is assumed that, in addition to their semi-weak interaction with leptons and hadrons, which is linear in ψ and ψ', the ψ and ψ' also have a strong interaction with hadrons, which is quadratic in ψ and ψ' is ruled out, or at least restricted, by the absence of the decays $\psi' \to \psi l\bar{l}$).

The existence of intermediate bosons having a strong interaction with hadrons which is quadratic in the fields of these bosons was proposed at one time as a description of the weak interactions[1]. Their strong quadratic interaction with hadrons was supposed to provide a cutoff of the growth of the virtual weak interactions at small distances and, in particular, account for the absence of the decays $K_L \to \mu^+\mu^-$ and the small mass difference between the K_L and K_S mesons[1b] (see also also[2]). Adopting the terminology of Appelquist and Bjorken[1c], we shall call this intermediate boson a stenon.

In most of its interactions with hadrons and leptons, the stenon is no different from a hadron-like meson which, because of some selection rules, goes weakly into ordinary hadrons. The processes of single and pair production of ψ mesons in hadronic collisions may have similar cross sections in the two cases and, in view of the presence of a strong quadratic stenon-hadron interaction, the cross section for stenon-nucleon scattering may be of the same order as ordinary hadronic cross sections, i.e., tens of millibarns. As a crude estimate of the cross section for stenon production in nucleon-nucleon collisions at high energies, we may take

$$\sigma\,(N+N\to\psi+\text{anything})\sim g_h^2\sigma_h f\,(M^2), \tag{5.7}$$

where $\sigma_h \sim 2\times 10^{-26}\,\mathrm{cm}^2$ is a quantity of the order of the hadronic total cross sections, $g_h^2 = 3\Gamma_h/M$, and $f(M^2)$ is a factor which takes into account the decrease of the cross section for producing a particle as its mass increases. Taking $f(M^2) \approx m_0^2/M^2$ with $m_0 \sim 1$ GeV in analogy with electroproduction, we obtain the estimate $\sigma(N+N\to\psi+\text{anything}) \sim 10^{-31}\,\mathrm{cm}^2$, which, bearing in mind its crudeness, is not inconsistent with the experimental data[3].

We note that a comparison of the estimate (5.7) with the experimental data leads to the conclusion that the form factor $f(M^2)$ cannot be very small; for example, it cannot decrease exponentially like $f(M^2) \sim \exp(-M/M_0)$ with $M_0 \lesssim 300$ MeV, as is assumed in the statistical theory. This statement applies not only to the stenon model, but generally, since it is to be expected that, whatever approach is adopted, the small parameter which determines the cross section for single ψ-meson production is of the same order of magnitude as the small parameter which determines the decay of the ψ into hadrons, i.e., of order g_h^2 (see Chap. VI). With the stenon hypothesis, the cross section for photoproduction of the ψ on a nucleon, $\gamma + N \to \psi$ + anything, at sufficiently high energies is of order

$$\sigma\,(\gamma+N\to\psi+\text{anything})\sim\alpha g_h^2\sigma_h f\sim 1\ \text{nb}.$$

This value is about an order of magnitude smaller than the experimental value[4]. Thus, the experimental data on ψ-meson photoproduction tends to oppose the stenon hypothesis. However, considering the crudeness of the theoretical estimates, we cannot regard this argument as a completely convincing refutation of this hypothesis.

The stenon differs from the hadronic ψ meson only by the fact that it has a direct interaction with leptons, whereas the interaction of the hadronic ψ meson with muons and electrons proceeds via a virtual photon. This circumstance makes it possible to ascertain whether the ψ meson is a stenon or a hadron by studying the process of electroproduction of the ψ on nucleons[5].

The electroproduction of a hadronic ψ meson on a nucleon is described by the diagram of Fig. 29, and the cross section for the process $e + N \to e + \psi$ + hadrons has the form

$$\frac{d\sigma}{dQ^2\,d\nu}=\frac{\alpha}{\pi}\frac{1}{Q^2}\frac{1}{E}\left(\frac{E-\nu}{\nu}+\frac{\nu}{2E}\right)\sigma_\gamma(Q^2,\ \nu); \qquad (5.8)$$

here $\nu = E - E'$, $Q^2 = 4EE'\sin^2(\theta/2)$, E and E′ are the energies of the initial and final electron, θ is the scattering angle, and $\sigma_\gamma(Q^2, \nu)$ is the cross section for photoproduction by a virtual photon. (We are assuming that $\nu^2 \gg Q^2$ and neglecting the contribution of longitudinally polarized virtual photons.)

For $Q^2 \gg \alpha(g_h/g_l)M^2$, stenon electroproduction is determined by the diagram of Fig. 30 and, with this hypothesis,

$$\frac{d\sigma}{dQ^2\,d\nu}=\frac{g_l^2}{\pi}\frac{1}{(M^2+Q^2)^2}\frac{Q^2}{E}\left(\frac{E-\nu}{\nu}+\frac{\nu}{2E}\right)\sigma_{sc}\ (Q^2,\ \nu), \qquad (5.9)$$

where $\sigma_{sc}(Q^2, \nu)$ is the cross section for scattering of a virtual stenon by a nucleon. Comparing (5.8) and (5.9) we see that $d\sigma/dQ^2$ has very different dependences on Q^2 in the region $\alpha M^2 \ll Q^2 \lesssim M^2$ for these two hypotheses.

Let us estimate the contribution to the total cross section for stenon electroproduction from the diagram of Fig. 30, i.e., let us calculate the integral of (5.9) with respect to Q^2 and ν. We take $\sigma_{sc}(Q^2, \nu) = \sigma_0$ for $Q^2 < m_0^2$ and $\sigma_{sc}(Q^2, \nu) = \sigma_0(m_0^2/Q^2)$ for $Q^2 > m_0^2$. Then (5.9) yields

$$\sigma\,(eN \to e\psi + \text{anything}) \approx \frac{g_l^2}{\pi}\frac{m_0^2}{M^2}\left(1+\frac{1}{2}\frac{m_0^2}{M^2}\right)\sigma_0\left(\ln\frac{E}{\nu_m}-\frac{3}{4}+\frac{\nu_m}{E}-\frac{\nu_m^2}{4E^2}\right), \qquad (5.10)$$

where $\nu_m = (M + 2m)M/2m$, and m is the nucleon mass. For $E = 20$ GeV, $m_0 = 1$ GeV and $\sigma_0 = 10$ mb, we have $\sigma(eN \to e\psi + \ldots) = 0.5$ nb. It is useful to compare Eq. (5.10) with the total cross section for electroproduction calculated under the assumption that the ψ is a hadron. The main contribution to the integral here comes from the region of small Q^2, and the cross section for electroproduction can be expressed in terms of the cross section for photoproduction in the form

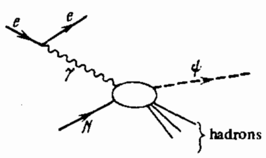


FIG. 29

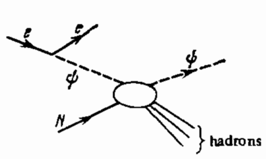


FIG. 30

$$\sigma(eN\to e\psi+\text{anything})=\frac{\alpha}{\pi}\,\sigma(\gamma N\to\psi+\text{anything})\ln\left(\frac{E}{m_e^2}\frac{m_0^2}{\nu_m}\right)\ln\frac{E}{\nu_m}, \qquad (5.11)$$

where m_e is the electron mass, and it is assumed that $E \gg \nu_m$. Substituting numerical values in (5.11) with E of the order of tens of GeV and $\sigma(\gamma N \to \psi + \text{anything}) \approx 20$ nb, we have $\sigma(eN \to e\psi + \text{anything}) \sim 1$ nb, i.e., the contributions of (5.8) and (5.9) are of the same order of magnitude. In the case of stenon electroproduction, it should be noted that the region of small $Q^2 \lesssim \alpha M^2$ is also described by the diagram of Fig. 29, its contribution to the total cross section for electroproduction being given approximately by (5.11).

Experimental data on muon pair production in a neutrino experiment, i.e., measurements of the cross section for the reaction $\nu_\mu + N \to \mu^+\mu^-$ + anything, enable us, within the framework of the stenon hypothesis, to exclude the possibility of a direct interaction of the ψ with neutrinos with a constant $g_\nu \sim g_l$[5]. In fact, if such an interaction occurred, the cross section for stenon production in the neutrino experiment would be of the same order of magnitude as the cross section for electroproduction of the ψ. In spite of possible uncertainties in the values of m_0 and σ_0 in (5.10), it can be estimated that at $E_\nu \sim 100$ GeV we have $\sigma(\nu N \to \nu\psi + \text{anything}) \sim 10^{-34}\ \text{cm}^2$, while the cross section for muon pair production is $\sigma(\nu N \to \mu^+\mu^-\nu + \text{anything}) \sim 7 \times 10^{-36}\ \text{cm}^2$. This figure is 3–4 orders of magnitude above the experimental upper limit for this quantity, $\sigma_{exp} \sim 10^{-38}$–$10^{-39}\ \text{cm}^2$[6].

Several papers have been devoted to attempts to describe the ψ mesons on the basis of various specific schemes for the theory of intermediate bosons[7]. The ψ and ψ' were identified with the neutral components of a triplet of intermediate bosons W^+, $W^{0'}$, W^0 having a strong quadratic interaction with hadrons, which had been proposed earlier by Marshak et al.[1b]. This model[7b] encounters the following difficulties: 1) the charged W boson should have a relatively low mass ~4 GeV, i.e., the linear growth of the cross section for νN scattering as a function of energy should come to an end when E_ν is of order 10–20 GeV, and this is not observed; 2) the decay $\psi' \to \psi$ + hadrons should take place mainly via the channel $\psi' \to \psi + K^0$, which is apparently not in accord with the experimental data; and 3) the decay $\psi \to$ hadrons $+\gamma$ should be much more probable than the decay $\psi \to e^+e^-$. A model has been studied[7c] in which the ψ, ψ' and ψ (4.15) are gauge fields interacting with conserved electron, muon and baryon currents, respectively, and mixing between the ψ, ψ' and ψ (4.15) occurs as a result of an interaction with Higgs mesons. In this case, an appreciable violation of μ-e universality was found in the decays of the ψ', namely $\Gamma^{\psi'}_{\mu\mu} \approx 8\Gamma^{\psi'}_{ee}$, which is not in accordance with the experimental data.

A model has also been proposed[7a] in which an intermediate boson can interact with leptons both directly and via a virtual photon. None of these models appear to be very attractive, and it seems to us that discussions of them should be deferred until experiments provide unambiguous evidence in favor of some particular model.

2. Particles with New Additive Quantum Numbers

Immediately after the discovery of the ψ meson, the hypothesis was put forward[8] that the ψ mesons possess a new quantum number (we shall designate it S′) which is conserved in the strong interaction but which is violated by some new interaction responsible for the decays of the ψ.

A characteristic prediction of such models is the statement that the ψ mesons are linear combinations of fields D which are not actually neutral and which differ from the antiparticles in the value of the quantum number S': $D \neq \bar{D}$, $\psi \sim D + \bar{D}$.

It was assumed that the same situation occurs in the case of ψ mesons as for strange particles—there is the possibility of strong associated production of ψ mesons with conservation of S' and relatively weak decays with violation of S'. Consequently, a strong energy dependence of the cross section for ψ-meson production was predicted. It is now known that the cross section, if it does rise with energy, does not increase very rapidly and, in any case, constitutes a small fraction $\sim 10^{-4}$ of the total cross section at energies up to $\sim 10^3$ GeV. However, it is hardly possible to rule out such a model on the basis of this fact alone, since the suppression of the cross section may be due to the large mass of the ψ mesons.

In this model, the decays into electrons and muons are usually explained by means of a direct interaction of the ψ with leptons, and in this respect the model is similar to the stenon hypothesis (see Sec. 1 of Chap. V). Experiments to test the stenon model would also be useful as tests of the hypotheses regarding the existence of a new quantum number for the ψ meson.

3. The Schwinger Model

Schwinger[9a] put forward the hypothesis that the ψ mesons belong to a family of vector and axial-vector fields which interact relatively weakly with the ordinary hadrons. The possible existence of such mesons had been discussed earlier by Schwinger[9b] in connection with the problems of weak-interaction theory, and the total number of such mesons was predicted to be 18 (nine vector and nine axial-vector fields).

Instead of the usual Cabibbo theory, according to which the W boson interacts with quarks, it is assumed in Schwinger's scheme that there is an interaction with mesons (which in principle cannot be constructed from quarks) having the form

$$W^{+}(v_{12} + \sin\theta_c v_{13} + \cos\theta_c v'_{13}) + \text{c.c.}, \qquad (5.12)$$

where we have written only that part of the Lagrangian which refers to the vector interactions of the charged intermediate boson; the unprimed fields denote the ordinary vector mesons $v_{12} = \rho^+$ and $v_{13} = K^{*+}$, and the primed fields denote new hypothetical fields which are weakly coupled to the ordinary hadrons. The consequences of (5.12) coincide with those of Cabibbo's theory for the charged currents. The Lagrangian can also be chosen so that it involves no $n\lambda$ transitions.

The model of[9] has not been developed in sufficient detail to permit a direct comparison with experiment. From the purely phenomenological point of view, the proposed scheme is similar to the model in which the ψ mesons are the quanta of a new semi-strong interaction.

4. Are the ψ Quasi-Nuclear Mesons?

In interpreting experiments on the annihilation process $e^+e^- \to$ hadrons, allowance must be made for the possible production of so-called quasi-nuclear mesons—relatively weakly bound states of a nucleon and an antinucleon, two nucleons and two antinucleons, a hyperon and an antihyperon, etc. Such mesons have been discussed previously in a number of theoretical papers (see the review[10]) and have been the objective of experimental searches[11]. The hypothesis that the ψ mesons are such quasi-nuclear mesons has been discussed in several recent papers[12].

Since the decay into hadrons is not suppressed in this case (it can occur without a pure gluon intermediate state), the fraction of the e^+e^- channel should be small, $\lesssim \alpha^2 \sim 10^{-4}$. The interpretation of the ψ and ψ' mesons as quasi-nuclei therefore runs into serious difficulties. The small number of strange particles among the decay products of the ψ meson rules out the interpretation of this meson as a bound $\Omega\Omega$ system.

The situation may be more favorable for the quasi-nuclear hypothesis in the case of broad resonances (such as the 4.15-GeV resonance). However, the appropriate analysis has not yet appeared in the literature.

5. Are the ψ Solitons?

In recent years, the idea of spontaneous symmetry breaking has aroused interest in solitons—stable classical solutions of field equations[13-20]. In particular, studies have been made of solutions having the form of very thin and dense surfaces, so-called "walls"[14], as well as bubbles formed by these surfaces, in analogy with domains in solids (the walls are boundaries between regions having different values of the meson condensate). Studies have been made of the meson condensate in nuclei[15], strings analogous to vortices in superconductors[16], and Dirac-monopole solutions, so-called "porcupines"[17].

There have been lively discussions in the literature of whether hadrons are such classical objects. In particular, a model of a hadron as a classical bubble filled with quarks has been widely discussed[18]. Very recently, it has been discovered that there are classical quasi-stable states formed by the "walls" when they approach each other at small distances[19].

It has been proposed[20] to interpret the narrow ψ-meson resonances as peculiar quasi-stable bubbles of the hadronic condensate. However, a direct realization of this idea encounters serious difficulties. In particular, it is not clear why a quasi-stable soliton decays so readily into an e^+e^- pair.

A more detailed analysis of this possibility will undoubtedly become unavoidable if it turns out that quark explanations of the ψ mesons are ruled out experimentally.

The idea of describing the ψ as a vibrational state of hadronic matter similar to the giant resonance in nuclei was put forward in[21].

VI. PRODUCTION MECHANISM OF THE NEW PARTICLES

1. Production of the New Particles in e^+e^- Annihilation

Measurements of e^+e^- annihilation into hadrons at high energies ($E_{c.m.}$ = 3–5 GeV) carried out in 1973–74 revealed a number of unusual properties of this process which did not correspond to the theoretical expectations (see, e.g., the review[1,2]). The strongest discrepancies with the theory were as follows: 1) the total cross section $\sigma_{tot}(e^+e^- \to \text{hadrons})$ remained practically constant within the studied range of energies or of $s = E_{c.m.}^2$ and the quantity

$$R = \frac{\sigma_{tot}(e^-e^- \to \text{hadrons})}{\sigma(e^-e^- \to \mu^-\mu^-)}$$

rose linearly with s, whereas the theoretically expected behavior was $\sigma(e^+e^- \to \text{hadrons}) \sim 1/s$ and R = const; 2) the inclusive spectra of the observed hadron with respect to the energy E, which were expected theoretically to have a scale-invariant behavior $(1/\sigma_{tot})d\sigma/dx = f(x)$, where $x = 2E/\sqrt{s}$ and f(x) is independent of s, turned out to be strongly dependent on s. After the discovery of the ψ and ψ' mesons, the question naturally arose as to whether these mesons are responsible for the unusual properties of the process of e^+e^- annihilation into hadrons and, in interpreting them as particles with hidden charm, whether the behavior of the cross section $\sigma(e^+e^- \to \text{hadrons})$ is a consequence of the production of charmed particles. More accurate measurements of $\sigma_{tot}(e^+e^- \to \text{hadrons})$ carried out after the discovery of the ψ and ψ' mesons have led to the observation of the ψ (4.15) resonance, which has made it even more plausible that the behavior of $\sigma_{tot}(s)$ can be explained by assuming that charmed particles are produced. We shall consider below the theoretical predictions for the process $e^+e^- \to$ hadrons on the basis of this hypothesis. (It is, of course, possible that the behavior of $\sigma_{tot}(s)$ and the production of ψ mesons are related to each other, but have nothing to do with charmed particles as we now understand them.)

In the parton model, the quantity R introduced above is given asymptotically as $s \to \infty$ by the expression

$$R = \frac{\sigma_{tot}(e^+e^- \to \text{hadrons})}{\sigma(e^-e^- \to \mu^-\mu^-)} = \sum_i Q_i^2. \tag{6.1}$$

where Q_i are the charges of the partons (we are assuming that all the charged partons are fermions with spin 1/2). The same expression arises in theories with asymptotic freedom, in which case it is also possible to calculate a correction term to (6.1) proportional to $1/\ln(s/\mu_0^2)$:

$$R = \sum Q_i^2 \left[1 + \frac{a}{\ln(s/\mu_0^2)}\right] \tag{6.2}$$

(where μ_0 is some mass). The coefficient a can be evaluated in specific theories.

In the theory involving three colored quartets of quarks c, p, n and λ with fractional charges, we have $R_\infty = \Sigma_i Q_i^2 = 10/3$ and $a = 12/25$.[3] Although the experimental data in the available energy range are much higher than the asymptotic value of R for this theory, we cannot say at the present time that these data exclude this theoretical possibility. One reason for this is that the asymptotic value of R is approached from above, as is implied by (6.2). Secondly, in the region in which the production of charmed particles sets in, there may be some growth of the cross section for the annihilation $e^+e^- \to$ hadrons as a result of threshold effects and the final-state interaction of the charmed particles, particularly if this interaction is due to long-range forces, as is assumed in the charmonium description of the ψ. If the production of charmed particles plays a major role in the total cross section for the annihilation $e^+e^- \to$ hadrons and if these particles have relatively large masses $\gtrsim 2$ GeV, we can also account for the violation of scale invariance in the inclusive spectra.

It should be stressed that the value $R_\infty = 10/3$ arises in the case of three colored quartets of quarks; for a single quartet, we would have $R_\infty = 10/9$, which is grossly inconsistent with the experimental data. To account for the experimental data, we must therefore assume either that 2/3 of the hadrons produced in e^+e^- annihilation at the available energies are colored (which seems improbable) or that there exists some confinement mechanism for colored states which inhibits their real production but which does not upset the validity of the calculation of R_∞ according to the parton model.

If the growth of R(s) for $\sqrt{s} > 3$ GeV is due to the production of pairs of charmed particles and if the principal decay modes of the charmed particles are those involving the emission of strange particles, kaons must be present in about 40% of the events in the annihilation process $e^+e^- \to$ hadrons at sufficiently high energies ($\sqrt{s} \gtrsim 4$ GeV). If the ψ (4.15) resonance is a bound state of the quarks $c\bar{c}$, it should have a large width for decay into $\bar{D}D$ and possibly into pions (the D mesons are charmed particles). In that case, under the same assumption regarding the decays of charmed particles, the peak in the annihilation cross section at $\sqrt{s} = 4.15$ GeV should be much more pronounced than that in the total cross section if a selection is made of those events that involve kaons in the final state.

The assumption that the ψ (4.15) resonance is due to the production of pairs of charmed particles (and/or the assumption that $\sigma(e^+e^- \to \text{hadrons})$ for $\sqrt{s} \gtrsim 4$ GeV has an appreciable contribution—of order 40%—from the production of charmed particles), together with the usual assumption (see the discussion in Chap. IV) that the decays of charmed particles are dominated by non-leptonic decays, lead to results which are difficult to reconcile with the existing experimental data.

In fact, let us consider the multiplicity of the hadrons produced in e^+e^- annihilation into charmed particles and pions, $e^+e^- \to \bar{D}D + n\pi$. The D meson cannot be much lighter than 1.8 GeV, otherwise the ψ' would decay with a large probability into a $\bar{D}D$ pair and it would not have such a small width. On the other hand, $2m_D < 4.15$ GeV. Consequently, the main processes involving D-meson production in the region of the ψ (4.15) resonance should be $e^+e^- \to \bar{D}D$ and $e^+e^- \to \bar{D}D + \pi$, and the process $e^+e^- \to \bar{D}D\pi$ should be expected to dominate for $E_{c.m.} \approx 4.5$–5 GeV.

With a D-meson mass $m_D \approx 1.8$–2.0 GeV, the average multiplicity of charged particles in non-leptonic decays of the D should be (see also Chap. IV) $\langle n_{ch}\rangle_D \approx 2.5$–3.0. (The multiplicity in this mass region in e^+e^- or $\bar{p}p$ annihilation has similar or even larger values; see, e.g., the review[1].) Taking into account the pions produced in the annihilation process, this implies that the average multiplicity of charged particles for $\sqrt{s} = 4.15$–5 GeV is

$$\langle n_{ch}\rangle \approx 6.$$

Since $\langle n_{ch}\rangle \leq 4$ for $\sqrt{s} < 4$ GeV, we should expect, according to the foregoing arguments, an appreciable enhancement of the charged-particle multiplicity at and above $\sqrt{s} = 4.15$ GeV, by an amount $\Delta\langle n_{ch}\rangle \approx 1$. This is not observed experimentally. This contradiction might be resolved by assuming that a significant fraction (of the order of 50% or more) of D-meson decays are semi-leptonic decays of the type

$$D \to l + \nu + K + m\pi. \qquad (6.3)$$

(Kaon emission is not required for our arguments here; it is assumed in order to obtain agreement with the usual theoretical schemes.) We should expect an average pion multiplicity $\langle m \rangle \sim 1$ in the decays (6.3), so that $\langle n_{ch} \rangle \approx 2$ (including leptons); if the principal decay modes of the D are semi-leptonic decays, the multiplicity of charged particles at $\sqrt{s}$ = 4.15–5 GeV is found to be

$$\langle n_{ch} \rangle \approx 4.6.$$

which is perhaps consistent with the experimental data. In this case, of course, leptons should be seen in a significant percentage (~20–30%) of the events.

The hypothesis that semi-leptonic and perhaps leptonic decays play a dominant role (if no kaon emission occurs) becomes plausible if we assume that the lowest boson states with $|C| = 1$ have spin 1, i.e., that the charmed vector mesons are lighter than the pseudoscalar mesons. With a V-A interaction, there is then no suppression of the decays $D^* \to l + \nu$ for these mesons (provided, of course, that these decays are allowed by the changes in the quantum numbers C and S due to the weak interaction) and no suppression of the decays $D^* \to l + \nu$ + mesons involving the emission of pions and kaons of low energies. This circumstance may enhance the relative probability of semi-leptonic decays.

The assumption that semi-leptonic meson decays play a major role also solves the problem of the "energy crisis" in e^+e^- annihilation (i.e., the fact that the fraction of the energy carried away by neutral particles increases as a function of s), since an appreciable fraction of the energy is carried away by neutrinos in this case. It is possible that this assumption may also provide the means of solving the problem of the violation of scaling in e^+e^- annihilation into hadrons.

In explaining the growth of R(s) with s and the ψ (4.15) resonance as consequences of the production of charmed particles, we are led to definite predictions for the ratio of kaons to pions. Experimentally (see the reviews[1,2]), the ratio K^-/π^- is about 0.2 at $\sqrt{s}$ = 3.0 GeV and changes comparatively little in going to $\sqrt{s}$ = 4.15 and $\sqrt{s}$ = 4.8 GeV. (Such a value is also obtained by integrating the experimental inclusive spectrum $Ed^3N/d^3p \sim e^{-E/T}$ with the same value $T \approx 170$ MeV for both pions and kaons. However, it must be borne in mind that the measurements of the spectrum and of the ratio K^-/π^- at each energy have a relatively poor accuracy ~30–50%.) If, as is usually assumed, the lightest of the charmed particles are the particles with $|C| = 1$ and $|S| = 0$, then, as we have already mentioned, we should expect theoretically that 40% of the events of e^+e^- annihilation into hadrons with $\sqrt{s}$ = 4.15–5 GeV involve two kaons. Taking into account the fact that the remaining 60% of the events at these energies have K^-/π^- = 0.2 as before and that the average π^- multiplicity is $\langle n_{\pi^-} \rangle \approx 2$, we find for the ratio K^-/π^- at $\sqrt{s}$ = 4.15–5 GeV the value

$$\frac{K^-}{\pi^-} = 0.2 \cdot 0.6 + 0.5 \cdot 1/2 \cdot 0.4 = 0.22,$$

i.e., the ratio K^-/π^- should grow by only 10% (or by 25% if $K^-/\pi^- \approx 0.15$ for $\sqrt{s}$ = 3–4 GeV). Such a growth is not excluded by the current experimental data. If the lightest of the charmed particles are the mesons with $|C| = 1$ and $|S| = 1$, then, according to Zweig's rule, the overwhelming majority of their decays should involve the production of η or φ mesons, giving a very specific character to the final non-leptonic states in e^+e^- annihilation (a large number of photons and/or kaons). It is to be hoped that experiments will, in the near future, provide unambiguous answers to all the questions which arise here.

2. The Production of New Particles in Hadron Collisions

a) The signal-to-background ratio in ψ-meson production. We can get a certain general idea of the production mechanism of ψ mesons in hadron collisions from an analysis[4] of the signal-to-background ratio, i.e., the ratio of the number of events in the peak to the number of events outside the peak. The background production of e^+e^- (or $\mu^+\mu^-$) pairs is a purely electromagnetic process and, by comparing the cross section for ψ-meson production with the cross section for producing a virtual photon of similar mass, we can make crude deductions about the force of interaction responsible for ψ-meson production. Moreover, this comparison requires no knowledge of the absolute value of the cross section.

In order to compare the intensities of ψ-meson production in electron-positron and in nucleon-nucleon collisions, let us form the ratio

$$r = \frac{\sigma(NN \to l^+l^-X)_{\text{res}}}{\sigma(NN \to l^+l^-X)_{\text{nonres}}} \frac{\sigma(e^+e^- \to X)_{\text{nonres}}}{\sigma(e^+e^- \to X)_{\text{res}}} \frac{\Delta_1}{\Delta_2}, \qquad (6.4)$$

where Δ_1 and Δ_2 are the experimental resolutions in the mass of the pair in the experiments on NN and e^+e^- collisions, and $l = e, \mu$.

Graphically, we are comparing the ratios of the diagrams a and b of Fig. 31 and the diagrams a and b of Fig. 32. If the ψ mesons were, for example, intermediate bosons which interact semi-weakly with quarks and leptons (see Chap. V), it would be natural to expect $r \sim 1$, since the effect of the structure of the strong interactions cancels to a great extent in calculating r.

Let us form the ratio r for the two experiments in which the ψ particles were discovered. In the experiment using the intersecting e^+e^- rings, the effect-to-background ratio is about 200 for a resolution in the mass of ~2.5 MeV, after introducing radiative corrections. In the experiment of Ting's group, the effect-to-

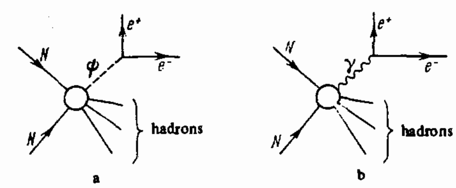


FIG. 31. ψ-meson production in nucleon-nucleon collisions (a) and nonresonant e^+e^- pair production in nucleon-nucleon collisions (b).

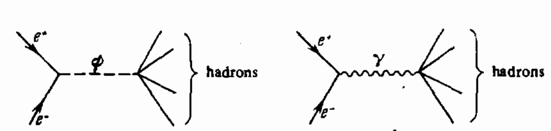


FIG. 32. ψ-meson production in e^+e^- annihilation (a) and nonresonant hadron production in e^+e^- annihilation (b).

background ratio depends on whether we take the background events to the left or to the right of the resonance (i.e., with $m_{ee} < m_\psi$ or $m_{ee} > m_\psi$) (see Fig. 13). From physical considerations, it seems reasonable to compare the numbers of events with $m_{ee} = m_\psi$ and $m_{ee} > m_\psi$, since the radiative decay $\psi \to e^+e^-\gamma$ leads to the production of pairs with mass m_{ee} somewhat less than m_ψ, although it has nothing to do with the strictly nonresonant production of e^+e^- pairs. We have given the formulas for the probability of radiative decay in Sec. 1 of Chap. II, and it can be seen that allowance for such decay is actually important. Moreover, as m_{ee} increases, the background of random coincidences decreases. As a result, we find that the effect-to-background ratio in the experiment of Ting's group is ~ 80 for a resolution of 25 MeV. Thus, the value of r is much greater than 1 in this case, namely $r \sim 5$.

The ratio r can also be estimated by using the data obtained by Lee's group (see[3] in Chap. III) for initial nucleons of high energy. Taking the estimate of the signal-to-background ratio ~ 2 for a resolution of 200 MeV, as given in their work, we have $r \sim 1$. If the value of r actually decreased with energy, we would have here a very interesting effect. However, the small value of r in this case is evidently due to the background of random coincidences. It would be of great interest to determine r more accurately.

b) **The production of ψ mesons in central collisions.** Within the framework of the quark model, the ratio r (see (6.4)) cannot be much less than unity. In fact, if ψ-meson decay into hadrons is possible, then it is also possible to have ψ production in quark annihilation followed by decay of the ψ into e^+e^-. The corresponding diagram of Fig. 33a is the time-reversal of the graph for ψ production in e^+e^- collisions. Similarly, the electromagnetic annihilation of a $q\bar{q}$ pair (Fig. 33b) describes the nonresonant production of lepton pairs in nucleon-nucleon collisions, and the time-reversed process coincides with the nonresonant production of hadrons in e^+e^- annihilation. Therefore $r \sim 1$ if only the quark annihilation graphs are taken into account (certain deviations of r from 1 may be due, for example, to the different electromagnetic charges of the n and p quarks).

Clearly, the annihilation of a quark-antiquark pair can describe only central collisions, i.e., processes involving a large energy release $\sim m_\psi$ in a short period of time $\tau \sim 1/m_\psi$. The fact that r is close to unity indicates that central collisions play an important role. The effect of a decrease in r with energy which we discussed in the preceding section would imply that the role of $q\bar{q}$ annihilation in the process of ψ production becomes even larger.

We note that the value of r may be large even for ψ production in central collisions. The point is that charmed particles may be produced not only in quark annihilation, but also in gluon collisions. It is usually

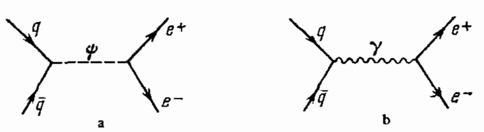


FIG. 33. ψ-meson production in quark annihilation (a) and nonresonant e^+e^- pair production in quark annihilation (b).

assumed that allowance must be made in the wave function of a nucleon at high energy not only for the quarks, but also for the gluons which transfer the interaction between them. It can be inferred from the analysis of the experimental data that the gluons carry about half of the momentum of the nucleons. If the gluon-quark coupling constant is relatively small, the cross section for producing charmed quarks is larger in gluon-gluon collisions than in quark annihilation. Such estimates of the cross section for this process can be found in[5].

It should be borne in mind, however, that the c quarks in gluon-gluon collisions are produced in a charge-even state and can be transformed into ψ mesons only in conjunction with the emission of pions or other particles. The probability of such a transition, which is forbidden by Zweig's rule, is difficult to estimate theoretically. There is evidently a much larger cross section for producing pseudoscalar mesons η_c, consisting of c quarks, in gluon-gluon collisions.

Another possibility for enhancing the cross section for ψ-meson production was pointed out in[6]. At high energy, there may be an admixture of charmed quarks in the nucleon, and the cross section for producing charmed particles or ψ mesons in $c\bar{c}$ annihilation may be dominant. It is usually assumed that the c quarks in the nucleon carry a small fraction of the momentum of the nucleon, and it is predicted in[6] that there is therefore a growth of the cross section for ψ-meson production in the region of small values of the scaling variable x (where $x = p_\psi/p_N$).

To conclude this section, we mention that the cross section for meson production is sometimes estimated[7] by means of the thermodynamic model[8], according to which the cross section is small because of the large mass m_ψ:

$$\sigma(NN \to \psi X) \sim \exp\left(-\frac{\sqrt{m_\psi^2 - p_\perp^2}}{T}\right). \tag{6.5}$$

where $p_\perp$ is the transverse momentum, and $T \sim 160$ MeV is a characteristic temperature.

Since we know that ψ mesons decay into ordinary hadrons with a small width, there is evidently no reason to believe that a thermodynamic equilibrium between the charmed and ordinary quarks is established during the collision time and that the formula (6.5) is applicable. This reservation may not apply to pair production of the ψ.

c) **Multiperipheral models of ψ-meson production.** As we have already mentioned, peripheral production of ψ mesons and charmed particles is also possible in strong interactions. Peripheral production of ψ mesons is often associated[9] with the quark diagrams that are allowed by Zweig's rule (see Chap. IV). An example of such a diagram is shown in Fig. 34a. It is significant that the ψ meson is produced here in conjunction with a pair of charmed mesons $D\bar{D}$. Single production of the ψ (see the diagram of Fig. 34b) is forbidden by this rule. A test of the prediction regarding the associated production of ψ mesons and charmed particles would be of great importance in elucidating the nature of both ψ mesons and Zweig's rule.

Even if associated production of ψ mesons and charmed particles is enhanced dynamically, the production of several heavy particles would require a suffic-

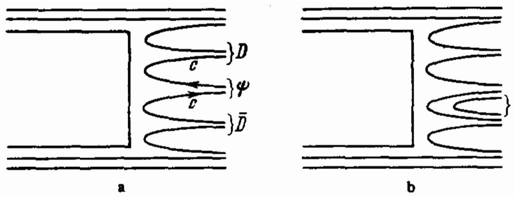


FIG. 34. ψ-meson production in nucleon-nucleon collisions in the quark model. An example of a diagram allowed by Zweig's rule (a) and a diagram forbidden by Zweig's rule (b).

iently large initial energy. We may therefore expect a strong dependence of the cross section on the initial energy.

According to some models, there is no suppression of the cross section at asymptotically large energy as a result of the large ψ-meson mass. Thus, according to[10], the cross section for inclusive production of charmed particles is determined by the cross section for the scattering of charmed particles by nucleons:

$$\frac{\sigma(D)}{\sigma(K)}=\frac{\sigma_{\text{tot}}(\psi N)}{\sigma_{\text{tot}}(\varphi N)}, \qquad (6.6)$$

where $\sigma(D)$ and $\sigma(K)$ are the cross sections for inclusive production of D and K mesons, and $\sigma_{\text{tot}}(\psi N)$ and $\sigma_{\text{tot}}(\varphi N)$ are the cross sections for the interactions of ψ and φ mesons with nucleons (we are assuming the validity of the additive quark model here, so that the cross section $\sigma_{\text{tot}}(\psi N)$ or $\sigma_{\text{tot}}(\varphi N)$ is a sum of the interaction cross sections of the constituent quarks of the ψ or φ meson). The prediction (6.6) refers to the so-called central region of the inclusive spectrum, i.e., to the production of D and K mesons with momenta equal to a relatively small fraction of the momentum of the incident particle.

A priori, there is no reason to expect a suppression of the cross section for scattering of the ψ by the nucleon, and the ratio (6.6) may have a value of order 1. Below, in connection with a discussion of the vector-dominance model, we shall give estimates of the cross sections $\sigma_{\text{tot}}(\psi N)$ and $\sigma_{\text{tot}}(\varphi N)$ according to which the ratio (6.6) is of order 1/5. In that case, according to (6.6), many charmed particles should be produced.

However, the currently available energies are evidently not at all asymptotic, and the cross section for the production of charmed particles may be strongly suppressed for kinematic reasons. The main kinematic effect is connected with the necessity of large momentum transfers in producing heavy particles. If two jets of particles with masses $\sqrt{s_1}$ and $\sqrt{s_2}$ are produced in a collision of two light hadrons, the minimum value of the square of the momentum transfer, $t_{\min}$, is determined by the condition

$$s\mid t_{\min}\mid=4s_1s_2,$$

where s is the square of the total energy of the colliding particles. On the other hand, it is known that the vertices have a pronounced t-dependence.

The value of $t_{\min}$ can become arbitrarily small at very high energies and, if the suppression of the cross section for ψ production is due entirely to the value of $t_{\min}$, the cross section for ψ production can become large at high energies, as predicted, for example, by the relation (6.6).

However, there seems to be a better basis for another point of view, according to which there exists an analogy between the production of heavy particles and the production of pions with large momentum transfers. Actually, we have already encountered this analogy in discussing the thermodynamic model (see (6.5)). In the multiperipheral model, there is no reason to expect the cross section to fall off exponentially with increasing $m_\perp=\sqrt{m^2+p_\perp^2}$, but, as before, there can be a significant dependence only on $m_\perp$. This assumption can be justified by analyzing multiperipheral diagrams of the "ladder" type. In this model, the small value of the cross section for producing particles with large transverse momentum is due to a large deviation from the mass shell in certain rungs of the diagram. It can be seen that the deviation from the mass shell in the propagators is characterized by $m_\perp$.

The existing data on the production of pions, kaons, protons and antiprotons are consistent with a dependence on $m_\perp$ which is universal for all particles, within an accuracy of a factor 2. By generalizing this observation to the case of the production of ψ mesons, we can estimate their production cross section on the basis of the data on the production of pions with large transverse momenta without introducing any unknown parameters. According to estimates obtained in this way by L. Kanetski, the cross section for ψ production is consistent with the existing data. It is significant that the coupling of ψ mesons with ordinary hadrons is not assumed to be small in the calculation. On the other hand, we know that the width of the ψ is small. The agreement between the calculated and experimental results evidently means that ψ mesons are produced either in pairs or together with charmed particles. We have already mentioned this possibility in discussing Zweig's rule.

We conclude this section by quoting some data on the production of strange particles and antiprotons, since we might expect the production of these particles to be analogous to ψ-meson production. In fact, the cross section for nucleon pair production may be small because of their masses, while the cross section for strange-particle production may be small because of selection rules which follow from Zweig's rule.

Experimentally, the cross sections for producing K mesons and Λ hyperons rise within the studied energy range: at 10 GeV the cross section for K^0-meson production is about 1 mb, while at 10^3 GeV it is of order 10 mb; the cross section for Λ-hyperon production is about the same at this first energy, while at 100 GeV it has reached 4 mb.

The inclusive cross section for antiproton production in pp collisions also rises up to the energies of the colliding beams at CERN: it has the value ~0.4 mb at an initial proton energy 50 GeV, ~1 mb at 200 GeV, and 2–4 mb at 1700 GeV.

Finally, the cross section for producing the pair $\Sigma^0\bar{\Sigma}^0$, which is suppressed at relatively low energy because of both the mass of the Σ^0 hyperon and the weakness of the transition of ordinary quarks into strange quarks, attains an appreciable value 0.4 mb at 10^3 GeV. Thus, we see that strangeness does not inhibit strange-particle production at high energy

3. The Production of New Particles in Electromagnetic Interactions

We are interested in ψ-meson production processes at the present time mainly from the standpoint of elucidating the properties of these particles. However, in case of ψ-meson production reactions in hadron-hadron collisions, we have already seen that estimates are strongly dependent on models of the strong interactions.

An analogous situation holds for ψ-meson photoproduction: the interpretation of the experimental data depends on the model. The vector-dominance model is most frequently employed. We shall also require the parton model. We shall first give a brief description of this model. Of course, we make no pretence of completeness in this exposition, whose aim is merely to facilitate an understanding of the material which follows.

a) **The vector-dominance model.** In this model (see, e.g., the review[11a]), it is assumed that the photon is first transformed into a vector meson, which is then scattered by the target. The cross section for photoproduction is thus expressed in terms of cross sections for strong interactions of vector mesons. The coupling constant for a vector meson (V) and a virtual photon can be determined from the electron or muon decay width, and it is usually assumed that this constant remains the same when the photon mass is changed from m_V (in the decay $V \to e^+e^-$, $\mu^+\mu^-$) to zero (in the reaction of photoproduction of a vector meson). The coupling constant for a photon and a vector meson is denoted by $(e/g_V)m_V^2$, and the $V \to e^+e^-$ decay width can be expressed in terms of this coupling constant as follows:

$$\Gamma(V \to e^+e^-) = \frac{\alpha}{g_V^2 4\pi}\frac{m_V}{3}, \tag{6.7}$$

where $\alpha = e^2/4\pi = 1/137$, and m_V is the mass of the vector meson.

The quantity g_V in the vector-dominance model has the interpretation of a coupling constant for the vector meson and hadrons. For example, g_ρ is the coupling constant for a ρ meson and pions. In this case, the ρ-meson exchange diagram reproduces the charge vertex of the photon-pion interaction at $q^2 = 0$. Similarly, the ω meson is associated with the isoscalar current of the n and p quarks, and the φ meson is associated with strange quarks. If there are charmed quarks, it is natural to generalize the vector-dominance model and to assume that g_ψ describes the strong interaction of ψ mesons with c quarks.

As there are several vector mesons—the ψ (3.1), ψ (3.7) and ψ (4.15)—the model becomes complicated in general and makes no pretence of great accuracy. Moreover, as the mass of the vector meson increases, the extrapolation of the photon-meson coupling constant from the point $q^2 = m_V^2$ to the point $q^2 = 0$ becomes more and more dubious.

As we have already mentioned, the photoproduction cross section in the vector-dominance model is expressed in terms of the cross section for the interaction of the vector meson with the target:

$$\frac{d\sigma(\gamma N \to VN)}{dt} = \frac{e^2}{g_V^2}\frac{d\sigma(VN \to VN)}{dt}. \tag{6.8}$$

If the interaction takes place at high energy, it is natural to take the forward scattering amplitude to be purely imaginary. We can then make use of the optical theorem

$$\operatorname{Im} A(VN \to VN) = s\sigma_{tot}(VN) \tag{6.9}$$

and express the differential cross section for forward photoproduction of a vector meson in terms of the total VN interaction cross section:

$$\left.\frac{d\sigma(\gamma N \to VN)}{dt}\right|_{t=0} = \frac{1}{16\pi}\frac{e^2}{g_V^2}\sigma_{tot}^2(VN). \tag{6.10}$$

If the forward scattering amplitude is not purely imaginary, then the right-hand side of (6.10) provides a lower limit for the differential cross section, and an upper bound on the value of $\sigma_{tot}(VN)$ can be obtained experimentally.

The vector-dominance model provides a satisfactory description of the photoproduction of ρ, ω, φ and ρ' mesons (although it may be necessary to allow for an additional dependence on the mass in the case of φ-meson photoproduction). The following values of the coupling constants and cross sections are obtained:

$$\left.\begin{aligned} &\frac{g_\rho^2}{4\pi} \approx 2{,}5, && \sigma_{tot}(\rho, N) \approx 25\ \text{mb}, \quad \sigma_{el} \approx 4.5\ \text{mb},\\ &\frac{g_\omega^2}{4\pi} \approx 18, && \sigma_{tot}(\omega, N) \approx 25\ \text{mb}, \quad \sigma_{el} \approx 5\ \text{mb},\\ &\frac{g_\varphi^2}{4\pi} \approx 10, && \sigma_{tot}(\varphi, N) \approx 10\ \text{mb}, \quad \sigma_{el} \approx 1\ \text{mb},\\ &\frac{g_{\rho'}^2}{4\pi} \approx 25\frac{\Gamma_{2\pi^+ + 2\pi^-}}{\Gamma_{tot}}, && \sigma_{tot}(\rho', N) \approx 30\ \text{mb}, \quad \sigma_{el} \approx 6\ \text{mb}. \end{aligned}\right\} \tag{6.11}$$

b) **The generalized vector-dominance model.** The vector-dominance model can be generalized to the case of a continuous spectrum of hadron masses[11b]:

$$F_\gamma(Q^2, \nu) = \int \frac{\sigma_{e^+e^-}(\varkappa^2)\varkappa^4}{(Q^2 + \varkappa^2)^2} F_h(\varkappa^2, \nu)\, d\varkappa^2, \tag{6.12}$$

where $\sigma_{e^+e^-}$ is the cross section for e^+e^- annihilation into hadrons with mass κ, $F_\gamma(Q^2, \nu)$ is the forward scattering amplitude for a photon with energy ν and 4-momentum squared $-Q^2$, and $F_h(\kappa^2, \nu)$ is the forward scattering amplitude for a hadron system with mass κ and energy ν. The normalization of the amplitudes is such that $\operatorname{Im} F_{\gamma,h} = s\sigma_{\gamma,h}^{tot}$, where s is the square of the total energy in the c.m.s., and $\sigma_{\gamma,h}^{tot}$ is the total cross section for the interaction of a photon or hadron system with the target. In the case of diffraction scattering or of a purely imaginary forward scattering amplitude, F_γ and F_h in Eq. (6.12) can be interpreted as the total cross section for the interaction of a photon or of the hadron system with the target.

The usual variant involving dominance of the ρ, ω and φ mesons corresponds to the replacement of the e^+e^- annihilation cross section in Eq. (6.12) by a sum of several terms proportional to the δ-function:

$$\frac{F_h(\varkappa^2, \nu)}{i2m_N\nu}\sigma_{e^+e^-} = \frac{e^4}{g_\rho^2}\sigma_{tot}(\rho N)\,\delta(\varkappa^2 - m_\rho^2) + \frac{e^4}{g_\omega^2}\sigma_{tot}(\omega N)\,\delta(\varkappa^2 - m_\omega^2) + \frac{e^4}{g_\varphi^2}\sigma_{tot}(\varphi N)\,\delta(\varkappa^2 - m_\varphi^2), \tag{6.13}$$

where $e^2/4\pi = 1/137$, the quantities g_ρ, g_ω and g_φ have been introduced above (see (6.11)), $\sigma_{tot}(\rho N)$, $\sigma_{tot}(\omega N)$ and $\sigma_{tot}(\varphi N)$ are the total cross sections for the interaction of ρ, ω and φ mesons with the nucleon, and we have assumed for simplicity that the amplitude for scattering of a vector meson by a nucleon is purely imaginary.

It is obvious that allowance can also be made for heavier mesons, such as the ρ' meson with mass $m_{\rho'} = 1.6$ GeV, in the same way.

c) **Scaling in eN scattering** and e^+e^- **annihilation.** It is well known that the cross section for electron (or muon) scattering by the nucleon has been found to have a scaling behavior at large momentum transfers Q and energies ν, i.e., in the so-called deep inelastic region. It has been found that there is a significant dependence only on the ratio

$$x=\frac{Q^2}{2m_N\nu}, \tag{6.14}$$

but not on each of the variables Q^2 and ν individually. Knowing the cross section at certain values of Q^2 and ν, the cross section can be determined at other values of Q^2 and ν, provided that the scales of both variables are simultaneously changed in such a way that x remains unchanged.

It has been found that the total cross section for eN scattering has no form factor in Q^2, as though there were point particles inside the nucleon. In other words, the cross section for the interaction of a virtual photon falls off like $1/Q^2$ (it is also possible to understand this behavior of the cross section without the analogy with point particles, starting instead from more general considerations involving an analysis of the space-time description of the scattering process[1,2]).

It would be natural to expect theoretically that the cross section for e^+e^- annihilation into hadrons at high energies also behaves like the cross section for point particles, i.e.,

$$\sigma(\varkappa^2)\sim\frac{c}{\varkappa^2}, \tag{6.15}$$

where the constant c is equal to the sum of the squares of the charges of the quarks. The behavior (6.15) can be justified on the basis of rather general assumptions (see, e.g., the review[2]).

The scaling behavior of the cross section has important consequences for the generalized vector-dominance model (6.12). It can be seen that the amplitude $F_h(\kappa^2, \nu)$ and hence the cross section for the interaction of a hadron system of mass κ with the target must fall off with κ at least like κ^{-2}. Otherwise, the integral with respect to κ^2 on the right-hand side of (6.12) would converge too slowly and would not give $F_\gamma \sim Q^{-2}$. (In principle, at large Q^2 but with a very small ratio $Q^2/2m_N\nu$, i.e., in the diffraction region, we cannot at the present time exclude the possibility that there is no scaling in deep inelastic scattering and that the cross section, for example, does not fall off with Q^2. In that case, the conclusion that $F_h \sim \kappa^{-2}$ does not hold.)

However, the conclusion that the cross section for the interaction of vector mesons falls off as a function of their mass is inconsistent with the data on ρ'-meson photoproduction (see (6.11)). To resolve this inconsistency, we must abandon the assumption that there are only diagonal transitions and consider also non-diagonal transitions such as $\rho N \to \rho' N$. It is then not possible to use the data on the photoproduction of vector mesons for a direct determination of the total cross sections for their interactions with the nucleon, and ρ'-meson production, for example, may be due to diffraction excitation of the ρ meson. By introducing a whole spectrum of vector mesons and allowing non-diagonal transitions between them, it is possible[13] to reconcile scaling in eN scattering and e^+e^- annihilation with Eq. (6.12), but the predictive power of the model is then weakened.

d) **The parton model.** An alternative description of the scattering of a virtual photon (or a weak intermediate W boson) by the nucleon at large energies and momentum transfers is provided by the parton model[14].

It is assumed that if the collision time is small, i.e., if the momentum transfer is large in comparison with the characteristic hadron masses $Q^2 \gg m^2_{\text{char}}$, the photon "sees" a nucleon as consisting of point particles—partons. Scattering by the partons takes place elastically, like scattering by point particles. The partons are characterized by internal quantum numbers (spin, charge, strangeness, etc.) and momenta which at high energies are measured relative to the momentum of the nucleon in the center-of-mass system or in the infinite-momentum system:

$$x=\frac{p_{\text{parton}}}{p_N}. \tag{6.16}$$

It is easy to show from the kinematics of elastic scattering that the variable defined in this way coincides with the scaling variable $x = Q^2/2m_N\nu$.

The quantum numbers of the partons are usually assumed to be the same as those of the quarks, and the distribution in the variable x is determined experimentally.

A distinction is made between the so-called valence partons and the "sea" of parton-antiparton pairs, i.e., it is assumed that the proton consists of three valence quarks ppn, which carry the proton quantum numbers, and $p\bar{p}$, $n\bar{n}$ and $\lambda\bar{\lambda}$ pairs, which have the vacuum quantum numbers.

e) **Partons and quarks.** To the reader who is familiar with only the quark model, it may seem surprising that we discuss the "sea" of quark-antiquark pairs in the nucleon, since the nucleon consists of only three quarks according to the quark model. However, the model of a nucleon as a bound system of three quarks may be useful only for nonrelativistic considerations.

The nucleon makes virtual transitions into states with a large number of quarks. To be sure, the time of these fluctuations is small:

$$\tau_{\text{fluc}}\sim\frac{1}{\Delta m}, \tag{6.17}$$

where Δm is the mass difference between the nucleon and a multi-quark (multi-particle) state; in considering static quantities (the magnetic moment of the nucleon, the β-decay axial constant of the nucleon, etc.), fluctuations which take place within a small time interval may not show up. In deep inelastic scattering, the collision time is small and the fluctuations of the nucleon must be taken into account.

In order to formulate the idea of a collision time more precisely, let us go over to the Breit coordinate system, in which the 4-momentum of the photon (or W boson) has only a third component: $q = (0, 0, q_z, 0)$. Clearly, $q_z = \sqrt{Q^2}$. The dimensions of the spatial region within which the photon is localized in this coordinate system can be estimated from the uncertainty principle:

$$\Delta z\sim\frac{1}{q_z}=\frac{1}{\sqrt{Q^2}}. \tag{6.18}$$

The collision time, during which the nucleon is located within the field of the photon, is also of order

$$\tau_{\text{coll}}\sim\frac{1}{\sqrt{Q^2}} \tag{6.19}$$

and falls off with increasing Q^2.

Moreover, the fluctuation time (6.17) in the laboratory coordinate system cannot be compared directly with (6.19), and allowance must also be made for the relativistic contraction of time of all processes in the moving coordinate system:

$$\tau'_{\rm fluc} \sim \frac{1}{\Delta m}\frac{p_z}{m_N}, \qquad (6.20)$$

where p_z is the nucleon momentum in the Breit coordinate system. The value of p_z is readily determined by making use of the invariance of the product of the 4-momenta of the photon and the nucleon ($p_N q$):

$$p_z = \frac{m\nu}{q_z} = \frac{1}{2}\sqrt{Q^2}\,x^{-1}. \qquad (6.21)$$

It follows from (6.20) and (6.21) that the smaller x at fixed Q^2, the larger the masses of the virtual states of the nucleon which show up in deep inelastic scattering, so that it is natural to expect a large contribution from the sea of parton-antiparton pairs at small x.

Antipartons show up only at small x, and for a different reason. As we have already mentioned, the quantity x is equal to the fraction of the nucleon momentum associated with a parton with which a collision has taken place. If, during a collision, a nucleon consists of a large number of $q\bar{q}$ pairs and three valence partons, the fraction of the momentum associated with each parton is small.

We note that the Regge model of deep inelastic scattering is often used for small values of x. This is connected with the fact that small x corresponds to the case in which the energy of the virtual photon is much greater than its mass and the scattering of the photon may be similar to the scattering of ordinary hadrons at high energy. Equality of the cross sections for the interactions of particles and antiparticles in diffraction scattering of ordinary particles (the Pomeranchuk theorem) corresponds to equality of the distribution functions for partons and antipartons at small x.

It is known from the experimental data that the antiparton contribution may be neglected in a first approximation in the analysis of the total cross sections:

$$\frac{P(\bar{p})+P(\bar{n})}{P(p)+P(n)} \leqslant 0.1,$$

where $P(p, n, \bar{p}, \bar{n})$ are the fractions of the nucleon momentum associated with the partons having the quantum numbers of the corresponding quarks.

It is also known that the partons have approximately half of the nucleon momentum. It is usually assumed that the rest of the momentum is carried by neutral gluons, which play no role in either the weak or the electromagnetic interaction.

f) The space-time description of the parton model. One can also pose the problem of how to reconcile Eq. (6.12) with the scaling behavior of the cross section in the parton model. It turns out that a suppression of the cross section for scattering of hadrons of large mass (see the foregoing discussion of the generalized vector-dominance model) arises here in a natural way[15,16]. In fact, the photon is first transformed into two energetic partons, whose interaction cross section with the target is small, since partons are point particles, and the scattering in a single partial wave at high energy is limited by the unitarity condition, $\sigma(s) \lesssim 1/s$.

In order to interact with the target, an energetic parton must undergo dissociation into slow particles, which is described in the multiperipheral model by a ladder-type diagram (Fig. 35). Consequently, the development of a shower of slow particles requires a time of order

$$\tau_{\rm shower} \sim \frac{1}{m_{\rm char}}\frac{\nu}{\varkappa}, \qquad (6.22)$$

where $m_{\rm char}$ is some characteristic strong-interaction mass, and $\nu/\varkappa$ is the Lorentz factor for transforming from the rest system of the hadron to the laboratory system.

On the other hand, the fluctuation time during which the photon "exists" in the form of hadrons of mass $\varkappa$ is limited by the uncertainty principle and is of order

$$\tau_{\rm fluc} \leqslant \frac{1}{\varkappa}\frac{\nu}{\varkappa}. \qquad (6.23)$$

If $\varkappa > m_{\rm char}$, the total fluctuation time is not long enough for the parton to dissociate into less energetic particles, and the cross section is small.

The degree of suppression of the cross section can be estimated as follows. Suppose that a photon with energy ν is transformed into one parton having almost the same energy and another parton with energy of order $m\omega$, where $\omega = 1/x$. By comparing (6.23) with (6.22) with $Q^2 \sim \varkappa^2$, we see that the less energetic parton in this case is able to produce a shower of slow particles, and scattering takes place, It is easy to see that the phase space for such configurations of partons, where one of them is much more energetic than the other, is of order $p_\perp^2/\varkappa^2$ of the total phase space, where $p_\perp$ is some characteristic transverse momentum.

Thus, the cross section for the interaction of a system of partons of mass $\varkappa$ with the target is suppressed by a factor $\varkappa^{-2}$ in relation to the ordinary hadronic cross sections. As we have already mentioned, such a decrease in the cross section is required in order to reconcile scaling in eN scattering and e^+e^- annihilation with Eq. (6.12).

Since it is known that scaling in deep inelastic scattering sets in for $Q^2 \lesssim 1\ \mathrm{GeV}^2$, the characteristic strong-interaction mass should not be large: $m_{\rm char} \approx m_\rho$ or less. In the language of the space-time description in the laboratory coordinate system, the dominance of the valence partons which we discussed in the preceding section means that no pomeron shower develops at the available energies, and the cross section is determined by the low-energy amplitudes for the interaction of partons. The diffraction regime should set in at higher energy.

g) Photoproduction of ψ mesons and the vector-dominance model. After our brief discussion of the vector-dominance model and the parton model, let us now turn to the application of these models to ψ-meson production processes. In particular, the vector-dominance model enables us to estimate the cross section for the interaction of ψ mesons with nucleons. From the existing experimental data, it is easy to obtain the

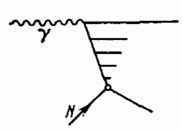


FIG. 35. A diagram of the "ladder" type for deep inelastic scattering of a photon by a nucleon.

following estimates of the total cross section σ_{tot} and the elastic cross section σ_{el} for ψN scattering:

$$\sigma_{tot}(\psi N) \sim 1 \text{ mb}, \quad E_\psi \sim 150 \text{ GeV}, \tag{6.24}$$

$$\sigma_{el}(\psi N) \approx \frac{\sigma_{tot}^2}{16\pi b} \sim 12 \ \mu\text{b}, \tag{6.25}$$

where b is the slope of the diffraction peak, for which we took the value $b = 4 \text{ GeV}^{-2}$ in making our estimate.

It can be seen that the cross section for the ψN interaction determined in this way is much smaller than the cross sections for scattering of other vector particles (see (6.11)). There are various possible explanations of this suppression: the inapplicability of the vector-dominance model, a non-hadronic character of the ψ mesons, or a suppression of the cross section for the interaction of charmed quarks with ordinary quarks. We shall discuss these possibilities below.

h) **Photoproduction of ψ mesons and Regge schemes.** The suppression of the cross section for the interaction of ψ mesons with nucleons, which is a consequence of the application of the vector-dominance model to the experimental data on ψ-meson photoproduction, is also predicted in certain Regge schemes[17-19]. This refers to models which relate the residue of the Pomeranchuk pole to the mass of a particle.

In the model of[20,21], the residue of the Pomeranchuk pole is expressed in terms of the masses of spin-2 particles which lie on appropriate Regge trajectories. For example, in the case of scattering of ρ mesons, which consist of n and p quarks, an important role is played by the exchange of the A_2 and f mesons ($m_{A_2} = 1.31$ GeV, $m_f = 1.26$ GeV), which are constructed from the same quarks. This statement is based on the application of dual models. Similarly, the residue for the φ meson is determined by f′ exchange ($m_{f'} = 1.51$ GeV). As the f′ meson is heavier than the f meson, the cross section for φN scattering is smaller than the cross section for the ρN interaction. This prediction is in agreement with experiment (see (6.11)). There are also natural explanations of certain other facts referring, for example, to the interactions of kaons[22].

Quantitatively, the residues of the Pomeranchuk pole for ψ and φ mesons are predicted to be in the ratio

$$r_\psi : r_\varphi = [1 - \alpha_{f'}(t)]\,[1 - \alpha_{f_c}(t)]^{-1}, \tag{6.26}$$

where $\alpha_{f'}$ and α_{f_c} are the trajectories of the Regge poles with the quantum numbers of the f′ and f_c mesons, which are constructed from charmed quarks, and t is the square of the momentum transfer.

The trajectory $\alpha_{f'}$ is known experimentally:

$$\alpha_{f'} \approx 0.15 + 0.8t. \tag{6.27}$$

The mass of the f_c meson, which consists of c quarks, has not yet been determined, but we can estimate the position of the trajectory α_{f_c} by assuming that α_{f_c} has a linear dependence on t and that the observed resonances ψ (3.7) and ψ (4.15) lie on daughter Regge trajectories. In that case,

$$\alpha_{f_c} = -3.8 + 0.5t. \tag{6.28}$$

We shall also require the trajectory α_f in what follows:

$$\alpha_f = 0.5 + 0.9\,t. \tag{6.29}$$

The ratio of residues (6.26) determines the ratio of the total cross sections for the φN and ψN interactions. The ratio of the cross sections for the photoproduction of φ and ψ mesons also depends on the meson-photon coupling constants g_ψ and g_φ and on the slopes of the diffraction peaks, b_ψ and b_φ. Taking the values of g_ψ and g_φ obtained from the decays $\psi \to e^+e^-$ and $\varphi \to e^+e^-$ and assuming that the slopes b_ψ and b_φ are equal, we find[17-19]

$$\sigma(\gamma N \to \psi N) = \sigma(\gamma N \to \varphi N)\,\frac{r_\psi^2}{r_\varphi^2}\,\frac{g_\varphi^2}{g_\psi^2}\,\frac{b_\varphi}{b_\psi} \sim 15 \text{ nb}, \tag{6.30}$$

which is in excellent agreement with the experimental value of the cross section, $\sigma = 13 \pm 5$ nb.

Similar conclusions are obtained from the model of[23], according to which the amplitude for scattering of a vector meson falls off as a function of its mass.

Let us consider two consequences of models involving a small value of the residue of the Pomeranchuk pole for the ψ meson[24]. First of all, the total cross section for photoproduction of the charmed particles $D\bar{D}$ should be much greater than the cross section for ψ-meson photoproduction:

$$\frac{\sigma_{el}}{\sigma_{tot}^c} \ll 1, \tag{6.31}$$

where $\sigma_{el} = \sigma(\gamma H \to \psi N)$, $\sigma_{tot} = \sigma(\gamma N \to \psi X) + \sigma(\gamma N \to D\bar{D}X)$, and $D\bar{D}$ is a pair of charmed particles. Allowance has been made here for the fact that the transitions of charmed quarks into ordinary quarks are weak, so that only particles with c quarks may be produced in the scattering of ψ mesons.

For $\sigma_{tot}^c = 1$ mb and $b_\psi = 4 \text{ GeV}^2$, we have

$$\frac{\sigma_{el}}{\sigma_{tot}} = \frac{\sigma_{tot}}{16\pi b_\psi} \approx \frac{1}{80} \tag{6.32}$$

and we obtain the following value for the total cross section for the photoproduction of charmed particles in the experiment in which photoproduction of the ψ was studied ([5] in Chap. III):

$$\sigma(\gamma N \to D\bar{D}X) \approx 1 \ \mu\text{b}, \tag{6.33}$$

i.e., a pair of charmed particles should be produced in about one percent of all the events.

This prediction is not valid if the experimental cross section for the ψ-meson photoproduction process $\gamma N \to \psi X$ exceeds the cross section for the elastic photoproduction process $\gamma N \to \psi N$ by a large factor, or if the forward scattering amplitude at the available energies is determined not by the imaginary part but by the real part. In this case, however, the agreement between the theoretical value (6.30) and the experimental data also disappears.

The second consequence of the small residue of the Pomeranchuk pole for the ψ meson which we would like to discuss concerns the interaction with nuclei. If the cross section for the ψN interaction really has the value 1 mb, nuclei should be transparent for ψ mesons. The cross sections for coherent production and for noncoherent production are proportional to $A^{4/3}$ and A, respectively, where A is the atomic number, and the coherent and noncoherent contributions are dominated by inelastic processes.

Strictly speaking, these statements apply to the region of ψ-meson energies E_ψ for which the ψ mesons are produced inside the nucleus, i.e., energies $E_\psi < R_{nuc} m_\psi^2$, where R_{nuc} is the radius of the nucleus.

Screening may set in at higher energies, even if the cross section for the interaction of ψ mesons with nucleons is small. Without giving any detailed arguments, we merely point out that the situation may be the same as in the case of the scattering of photons: although the cross section for the interaction of photons is small, scattering by nuclei at high energy should exhibit screening of the nucleons.

i) **Inclusive photoproduction and electroproduction of charmed particles in the parton model.** In the parton model, it is easier to estimate the inclusive production of particles containing charmed quarks than the cross sections for the individual channels. We shall make an estimate by using Eq. (6.12), assuming for simplicity that the cross section for e^+e^- annihilation into hadrons is the sum of the cross sections for annihilation into the quarks $p\bar{p}$, $n\bar{n}$ and $\lambda\bar{\lambda}$ and the charmed quarks $c\bar{c}$. We shall identify the production threshold for ordinary particles with the ρ-meson mass m_ρ and the production threshold for particles containing c quarks with the ψ-meson mass m_ψ. As a crude estimate of the cross sections for photon scattering, we then obtain

$$\sigma_\gamma(Q^2, \nu) \sim \left(\frac{e_p^2+e_n^2+e_\lambda^2}{m_\rho^2+Q^2}+\frac{e_c^2\sigma_c/\sigma_q}{m_\psi^2+Q^2}\right)\frac{m_\rho^2\sigma(Q^2=0, \nu)}{2/3}, \qquad (6.34)$$

where $e_{p,n,\lambda,c}$ are the charges of the quarks, and σ_c and σ_q are the cross sections for the interactions of ordinary (q) and charmed (c) quarks with the nucleon. The expression (6.34) is normalized to the photoproduction cross section $\sigma(Q^2 = 0)$.

If Eq. (6.34) can be used to estimate the cross section for producing charmed particles, this is so only at very high energies, In fact, we have replaced the quark-nucleon forward scattering amplitude in (6.12) by its asymptotic value is σ_{tot} in deriving (6.34). As we have already mentioned, the so-called dominance of the valence quarks holds at currently available energies; this means that the regime of diffraction scattering has not yet been reached, even for ordinary quarks. There is every reason to believe that this regime will become established at even higher energy for charmed particles.

In estimating the cross section within the accessible energy range, the last term in (6.34) must therefore be multiplied by some function of the photon energy, $f(x')$, with $x' = (Q^2 + m_\psi^2)/m\nu$, this function being less than unity.

To estimate the quantity $f(x')$, we note that Eq. (6.34) actually predicts a violation of scaling in deep inelastic scattering. On the other hand, it is known experimentally that scaling holds with 10% accuracy for $Q^2 \sim m_\psi^2 \sim 10$ GeV^2 up to energies $\nu \sim 50$ GeV. Under these conditions, we might expect that the contribution of charmed-particle production is also not more than ~10%, i.e., $\sigma_c/\sigma_q \lesssim 0.1$.

There are also no anomalies in the energy dependence of the photoproduction cross section. The data are well described by the contribution of several Regge poles, starting with relatively low energies ~5 GeV. The accuracy is better here than in the case of electroproduction, but the contribution of the second term in (6.34) falls off with decreasing Q^2 and it then becomes more difficult to obtain a bound on its magnitude.

j) **Photoproduction and electroproduction of ψ mesons in the parton model.** What has been measured experimentally is not the inclusive production of particles containing charmed quarks, but the cross section for ψ-meson photoproduction. It is more difficult to estimate the cross section for this process in the parton model, since allowance must also be made for the probability that $c\bar{c}$ quarks produced by the scattering are "assembled" into a ψ meson instead of any other state. We shall normalize the cross section for ψ-meson photoproduction to ρ-meson production and assume for the purpose of our estimates that the probability of transitions of ordinary quarks $q\bar{q}$ into a ρ meson and of the quarks $c\bar{c}$ into a ψ meson are characterized by the $\rho \to e^+e^-$ and $\psi \to e^+e^-$ decay widths (such a factor occurs in the generalized vector-dominance model (see (6.12)).

Thus, the photoproduction cross sections for ψ mesons and ρ mesons are found to be in the ratio

$$\frac{\sigma(\gamma N\to\psi X)}{\sigma(\gamma N\to\rho X)} \sim \frac{m_\rho^2}{m_\psi^2}\frac{\sigma_c}{\sigma_q}\frac{\Gamma(\psi\to e^+e^-)m_\rho}{\Gamma(\varphi\to e^+e^-)m_\psi}, \qquad (6.35)$$

where σ_c and σ_q are the cross sections for scattering of charmed and ordinary quarks by the nucleon.

Taking a value ~0.1 for the ratio σ_c/σ_q at an energy $E_\gamma \sim 100$ GeV, as implied by the estimates of the preceding section, we find that the cross section for ψ-meson photoproduction is given by[24]

$$\frac{\sigma(\gamma N\to\psi X)}{\sigma(\gamma N\to\rho X)} \sim \frac{1}{16}\frac{1}{10}\frac{1}{5} \sim \frac{1}{800}, \qquad (6.36)$$

where the individual factors correspond to the various factors on the right-hand side of (6.35), or

$$\sigma(\gamma N\to\psi X) \sim 20 \text{ nb}. \qquad (6.37)$$

This result is close to the experimental value of the cross section at the highest available energy $E_\gamma \sim 150$ GeV (see Chap. III). Considering the crudeness of the estimate, the agreement seems more readily fortuitous. We note also that, strictly speaking, we have obtained an upper estimate of the contribution of the second term in (6.34) from the inclusive cross sections, since neither a violation of scaling nor anomalies in the total photoproduction cross section have been observed. Therefore the estimate (6.37) is actually an upper limit.

The study of ψ-meson electroproduction would be of great importance in testing the theoretical scheme. It follows from Eq. (6.34) that the relative contribution of charmed-particle production to the cross section becomes larger with increasing Q^2 (at least for a fixed value of $x' = (Q^2 + m_\psi^2)/2m_\nu$). The study of ψ-meson electroproduction would also provide a test of the stenon hypothesis for the nature of the ψ meson (see Chap. V).

We note that, since there is no dominance of vector mesons in the parton model, the estimates of the total cross section for the photoproduction of charmed particles which we have given (see Sec. 3i above) do not apply. The dependence on the atomic number may also be different from that discussed above. Unfortunately, it is not clear whether it will be possible to distinguish experimentally between the results for the interaction of ψ mesons with nuclei which follow from the vector-dominance model and the parton model. The point is that some suppression of the cross section for the interaction of charmed quarks is assumed in the parton model, at least at currently accessible energies.

4. The Production of Charmed Particles in Weak Interactions

a) **The production of charmed particles on the valence partons and the "sea" of parton pairs.** To estimate the cross sections for producing charmed particles in neutrino experiments, we can make use of the vector-dominance model or the parton model, as in considering ψ-meson photoproduction. We shall first consider certain qualitative arguments which are common to both models.

The only possible reaction involving the valence partons in which charmed particles are produced in a neutrino beam is

$$\nu_\mu n \to \mu^- c, \tag{6.38}$$

and its corresponding amplitude is proportional to $\sin\theta_C$, where θ_C is the Cabibbo angle, with $\sin^2\theta_C \sim 1/20$.

With an admixture of $\bar{n}$ partons, it is possible to have the reaction

$$\bar{\nu}_\mu \bar{n} \to \mu^+ \bar{c}. \tag{6.39}$$

The possible processes involving λ quarks are

$$\nu_\mu \lambda \to \mu^- c, \quad \bar{\nu}_\mu \bar{\lambda} \to \mu^+ \bar{c} \tag{6.40}$$

and the corresponding amplitude is proportional to $\cos\theta_C$.

Finally, an admixture of c and $\bar{c}$ partons shows up in the transitions

$$\begin{aligned} &\nu_\mu \bar{c} \to \mu^- \bar{\lambda}, \quad \bar{\nu}_\mu c \to \mu^+ \lambda, \\ &\nu_\mu \bar{c} \to \mu^- \bar{n}, \quad \bar{\nu}_\mu c \to \mu^+ n, \end{aligned} \tag{6.41}$$

whose amplitudes are related to those of (6.38), (6.39) and (6.40) by the crossing relations.

The c quarks can also take part in elastic neutrino scattering induced by the neutral currents:

$$\nu_\mu c \to \nu_\mu c. \tag{6.42}$$

In the Weinberg-Salam model, the corresponding amplitude has the form

$$\frac{G}{2\sqrt{2}} \bar{\nu}_\mu \gamma_\alpha (1+\gamma_5) \nu_\mu [\bar{c}\gamma_\alpha c (1 - 4\sin^2\theta_W) + \bar{c}\gamma_\alpha\gamma_5 c], \tag{6.43}$$

where θ_W is the Weinberg angle, whose experimental value is apparently given by $\sin^2\theta_W \sim 0.3$, and γ_α and γ_5 are the usual Dirac matrices.

Roughly speaking, the production of charmed particles on the valence quarks corresponds to fragmentation of the target, i.e., to production of charmed hadrons with small momentum in the laboratory coordinate system, while the interaction with partons from the "sea" of quark-antiquark pairs corresponds to fragmentation of the current or of the intermediate boson, i.e., to production of mesons with large momentum in the laboratory coordinate system.

We have enumerated the various elementary processes involving the production of charmed particles in the language of the parton model. Of course, a similar analysis is possible in the vector-dominance model[19,25]. According to this model, the intermediate boson is first transformed into a vector meson containing a charmed quark, $F^* \sim \lambda\bar{c}$ or $D^* \sim n\bar{c}$. The amplitude for a transition into D^* contains a factor $\sin\theta_C$. Scattering of a λ ($\bar{\lambda}$) or n ($\bar{n}$) quark by the target corresponds to the processes (6.38) or (6.40), while scattering of charmed quarks corresponds to the processes (6.41).

It is usually assumed (see, for example, the discussion of ψ-meson photoproduction in the vector-dominance model) that the interaction cross section of charmed quarks is smaller than that of p, n and λ quarks. In that case, the amplitudes for the processes (6.41) can be neglected. Similarly, it is usually assumed in the parton model that the admixture of c partons in the nucleon is smaller than that of λ partons, and (6.41) is again neglected.

If the cross sections for scattering of the F^* and $\bar{F}^*$ by the target are equal, diffraction production of charmed particles is the same in a neutrino beam and in an antineutrino beam. Similarly, it is usually assumed in the parton model that the momentum distributions of the λ and $\bar{\lambda}$ partons are the same.

b) **Estimates of the cross sections for charmed-particle production in the vector-dominance model.** To calculate the cross section for charmed-particle production in the vector-dominance model[19,25], it is necessary to know the constant for the transition of a W boson into an F^* meson and the residue of the Pomeranchuk pole for the F^* meson (in what follows, we shall neglect the transition $W \to D^*$, whose amplitude is proportional to $\sin\theta_C$). The value of the residue r_{F^*} can be determined in the model of tensor-meson dominance (see Chap. IV):

$$\frac{r_{F^*}}{r_\rho} = \frac{1}{2}(r_1 + r_2), \quad r_1 = \frac{1-\alpha_f(t)}{1-\alpha_{f'}(t)}, \quad r_2 = \frac{1-\alpha_f(t)}{1-\alpha_{f_c}(t)}, \tag{6.44}$$

where the notation is the same as in the section concerning ψ-meson photoproduction. In deriving the relations (6.44), use is also made of the additive quark model: it is assumed that the total cross section for the interaction of the F^* with the target is the sum of the cross sections for the interactions of the λ and c quarks which form the F^* meson.

The constant g_{F^*} for the transition of a W boson into an F^* meson cannot be determined experimentally at the present time. In the limit of SU(4) symmetry, g_{F^*} is equal to the constant g_ρ for the transition of a W boson into a ρ meson. It might be expected that this consequence of SU(4) symmetry is not violated very strongly. In fact, the constants for the transitions of a photon into ψ and ρ^0 mesons are related by the equation $g_\psi = (2\sqrt{2}/3)g_\rho$, or $g_\psi \approx g_\rho$, in the limit of SU(4) symmetry. Experimentally, $2g_\psi \sim g_\rho$, i.e., the increased mass of the ψ meson has led to a change in its constant by a factor of two. The F^* meson should be lighter than the ψ meson, and the violation of SU(4) symmetry might be expected to be smaller in this case. To obtain a crude estimate, we shall take $1.5g_{F^*} = g_\rho$.

Assuming that the trajectories α_f, $\alpha_{f'}$ and α_{f_c} are described by Eqs. (6.29), (6.27) and (6.28), respectively, we obtain the following estimate of the cross section for diffraction production of charmed particles:

$$\sigma^c_{\text{diff}} \sim \frac{1}{4}\sigma^{\text{tot}}_{\text{diff}}, \tag{6.45}$$

where $\sigma^{\text{tot}}_{\text{diff}}$ is the total cross section for diffraction production. We stress that $\sigma^{\text{tot}}_{\text{diff}}$ is not the same as the total cross section for the interaction of the neutrino (or antineutrino) with the nucleon. If the cross section were dominated by ρ-meson exchange, the total cross sections $\sigma_{\text{tot}}(\nu N)$ and $\sigma_{\text{tot}}(\bar{\nu}N)$ would be approximately equal. Experimentally, they differ by about a factor of three, i.e., by the maximally allowed amount in the scaling region. This difference between neutrino and antineutrino interactions corresponds to the dominance of valence partons (see Sec. 3e) and relatively weak diffraction transitions.

The value of $\sigma_{\rm diff}^{\rm tot}$ can be estimated from the relation

$$\frac{(1/3)\,\sigma_{\rm valence}-\sigma_{\rm diff}^{\rm tot}}{\sigma_{\rm valence}+\sigma_{\rm diff}^{\rm tot}}\approx\frac{\sigma_{\rm tot}(\bar{\nu}N)}{\sigma_{\rm tot}(\nu N)}. \qquad (6.46)$$

Taking $\sigma_{\rm tot}(\bar{\nu}N) = 0.40\sigma_{\rm tot}(\nu N)$ (the experimental values are 0.37 ± 0.02 at the CERN energies and 0.36 ± 0.05 at the Batavia energy), we obtain

$$\sigma_{\rm diff}^{\rm tot}(\nu N)\approx 0.1\sigma_{\rm tot}(\nu N), \qquad (6.47)$$

$$\sigma_{\rm diff}^{\rm tot}(\bar{\nu}N)\approx 0.3\sigma_{\rm tot}(\bar{\nu}N). \qquad (6.48)$$

Accordingly,

$$\sigma_{\rm diff}^{c}(\nu N)\approx 0.025\sigma_{\rm tot}(\nu N) \qquad (6.49)$$

and

$$\sigma_{\rm diff}^{c}(\bar{\nu}N)\approx 0.075\sigma_{\rm tot}(\bar{\nu}N). \qquad (6.50)$$

Detailed estimates of the cross sections for charmed-particle production can be found in[25b], according to which $\sigma_{\rm diff}^{c}(\bar{\nu}N) = 0.1\sigma_{\rm tot}(\bar{\nu}N)$.

c) **The Q^2-dependence of the cross section for diffraction production.** Strictly speaking, the estimate (6.49) and (6.50) refers to the value $Q^2 = 0$. As Q^2 increases, the absolute value of the cross section for charmed-particle production becomes smaller, but its relative contribution to the total cross section rises:

$$R_{\rm diff}=\frac{\sigma_{\rm diff}^{c}}{\sigma_{\rm diff}^{\rm tot}}\sim\frac{(Q^2+m_\rho^2)^2}{(Q^2+m_c^2)^2}, \qquad (6.51)$$

where the Q^2-dependence is due to the propagators of the vector particles. At $Q^2 = 1\ {\rm GeV}^2$, the ratio $R_{\rm diff}$, as given by (6.51), is larger than its value at $Q^2 = 0$ by a factor of about 7.

However, we cannot rely upon the estimate (6.51) literally, since the vector-dominance model does not provide a satisfactory description of the Q^2-dependence of the total cross section: according to this model, the cross section falls off like Q^{-4}, but experimentally it falls off like Q^{-2}.

For this reason, it is important to emphasize that the growth with Q^2 of the relative contribution to the total cross section from charmed-particle production is not associated with the specific model used above, but is of a more general character. A similar conclusion can be reached, for example, by considering Eq. (6.12). As we have already explained (see Sec. 3c), the scaling behavior of the electroproduction cross section implies that the integral on the right-hand side of (6.12) has good convergence at high energies and is determined by the integrand near the lower limit. Thus, the cross section for the interaction of a virtual photon (or W boson) is proportional to Q^{-2} for ordinary particles and proportional to $(Q^2 + m_c^2)^{-1}$ for charmed particles, since the integration with respect to the mass κ begins at m_c in the second case. Then

$$R_{\rm diff}\sim\frac{Q^2+m_\rho^2}{Q^2+m_c^2}. \qquad (6.52)$$

A measurement of the Q^2-dependence of the ratio $R_{\rm diff}$ would evidently serve as a test of whether the mechanism of diffraction production is actually operative.

In connection with our discussion of the generalized vector-dominance model, we note that the estimates of the absolute value of the cross section at $Q^2 = 0$ are smaller in this case by a factor m_ρ^2/m_c^2 than in the ordinary vector-dominance model. The reason for this is that, according to this model, the cross section for scattering of a vector meson falls off as a function of its mass. There is also a factor m_ρ^2/m_c^2 according to the parton model, but this factor is not strictly due to the mass of the vector particle, but rather to the large deviation from the mass shell.

In the generalized vector-dominance model, we can give the following crude estimate of the cross section for charmed-particle production:

$$\frac{\sigma_{\rm diff}^{c}}{\sigma_{\rm diff}^{\rm tot}}\sim\frac{1}{4}\,\frac{\langle Q^2\rangle+m_\rho^2}{\langle Q^2\rangle+m_c^2}, \qquad (6.53)$$

where $\langle Q^2\rangle$ is the average value of the square of the momentum transfer in the diffraction processes, and we are retaining our previous estimates of the cross sections for quark interactions. It is usually assumed that diffraction scattering is characterized by values of the scaling variable in the range $x \lesssim 0.1$. Using the fact that $x = Q^2/2m\nu$ and that the average value $\langle\nu\rangle$ for neutrino interactions is approximately $(1/2)E_\nu$, we obtain the estimate $\langle Q^2\rangle \approx 0.025\cdot 2m_N E_\nu$.

d) **The cross section for charmed-particle production in the parton model.** Let us consider the production of charmed particles within the framework of the parton model in greater detail. If we neglect the masses of charmed particles, the cross section can be written in analogy with the ordinary case:

$$\frac{d^2\sigma^c(\nu N)}{dx\,dy}=\frac{G^2 s}{2\pi}\{\sin^2\theta_C\,[n(x)+p(x)]+\cos^2\theta_C\cdot 2[\lambda(x)+(1-y)^2\bar{c}(x)]\}, \qquad (6.54)$$

where $n(x)$, $p(x)$, $\lambda(x)$ and $\bar{c}(x)$ are the distribution functions of the n, p, λ and $\bar{c}$ partons in the nucleon, $y = \nu/E_\nu$, $x = Q^2/2m_N\nu$, $s = 2m_N E_\nu$, and θ_C is the Cabibbo angle.

Integrating the cross section (6.54) with respect to x and y, we find

$$\frac{\sigma^c(\nu N)}{\sigma_{\rm tot}(\nu N)}=\frac{P(\lambda)+(1/3)P(\bar{c})+(1/2)\,{\rm tg}^2\theta_C\,[P(n)+P(p)]}{(1/2)[P(p)+P(n)]+(1/3)P(\bar{p})+(1/3)P(\bar{n})}, \qquad (6.55)$$

where $P(q)$ is the total momentum of the partons of the appropriate type.

As we have already mentioned, almost nothing is known experimentally about the distribution function $\lambda(x)$ $(c(x))$. The production of charmed particles would be the first experiment that is sensitive to an admixture of λ partons in the nucleon. We shall consider briefly the various theoretical estimates of $\lambda(x)$.

1) It can be assumed that the fraction of the momentum belonging to the λ partons is proportional to the fraction of the momentum belonging to the kaons in inclusive production in the strong interactions. The nucleon is viewed here as dissociating in a vacuum into a leading nucleon, together with pions and kaons. In hadron-hadron collisions involving small momentum transfer, the vacuum fluctuation is transformed into ordinary hadrons; in deep inelastic processes, the photon (or W boson) interacts with the constituent quarks of

the hadrons. The momentum of the λ quarks is then proportional to the momentum of the kaons:

$$\frac{P(\lambda)}{P(\bar{n})} \approx \frac{P(\lambda)}{P(\bar{p})} \approx \frac{P(K)}{P(\pi)} \sim \frac{1}{5} - \frac{1}{10}. \tag{6.56}$$

However, we cannot exclude the possibility that the final hadronic state in strong interactions is depleted of kaons in the decay of resonances or heavy clusters. These decays occur for a relatively long time, longer than the time of a deep inelastic collision. Other estimates are therefore frequently employed.

2) The admixture of λ quarks can be estimated by making use of relations of the type (6.6):

$$\frac{\lambda(x)}{\bar{n}(x)} \approx \frac{\lambda(x)}{\bar{p}(x)} \approx \frac{\sigma(\lambda N)}{\sigma(\bar{p}N)} \approx \frac{\sigma(\psi N)}{\sigma(\rho N)} \sim \frac{1}{2.5}. \tag{6.57}$$

In fact, in the Breit coordinate system, the cross section for a deep inelastic process can be expressed in terms of the parton densities $\bar{p}(x)$, $\bar{n}(x)$ or $\lambda(x)$. In the laboratory system, the same cross section can be expressed in terms of the quark scattering amplitude. It follows from the equivalence of these two approaches that the content of partons of each type in the nucleon is proportional to the cross section for the interaction of these partons with nucleons, $\sigma(\bar{p}N)$, $\sigma(\bar{n}N)$ or $\sigma(\lambda N)$. Using also the additive quark model, the cross sections for the interactions of quarks can be expressed in terms of cross sections for the interactions of hadrons, such as $\sigma(\varphi N)$ or $\sigma(\rho N)$.

It is well known, however, that such relations may be violated rather strongly at the available energies. For example, similar arguments lead to the conclusion that all hadronic total cross sections are equal[15].

3) Finally, the simplest possible way of estimating the admixture of λ quarks consists in the assumption that the "sea" of parton pairs is SU(3)-symmetric, i.e., that

$$\bar{n}(x) = \bar{p}(x) = \lambda(x) = \bar{\lambda}(x). \tag{6.58}$$

Detailed model calculations of the parton distribution functions have been made on the basis of this assumption. These calculations made use of further assumptions, in particular, duality. We refer to[26] as an example and quote the simplest analytic expressions proposed by Farrar (cited in[9b]):

$$p(x) = \frac{0.2(1-x)^7}{x} + \frac{1.89(1-x)^7}{x^{1/2}} + 5(1-x)^3,$$
$$n(x) = \frac{0.2(1-x)^7}{x} + \frac{1.03(1-x)^7}{x^{1/2}} + 5(1-x)^3,$$
$$\bar{p}(x) = \bar{n}(x) = \lambda(x) = \bar{\lambda}(x) = \frac{0.2}{x}(1-x)^7.$$

It follows from our previous discussion that the contribution of the λ quarks is evidently overestimated in relation to the contributions of the $\bar{n}$ and p quarks.

e) **The production of charmed particles by neutral currents.** So far, we have discussed only the production of charmed particles by charged currents. Processes involving the production of new particles in neutrino reactions as a result of neutral currents are also possible. We have already written the appropriate term in the Hamiltonian within the framework of the Weinberg-Salam model (see (6.43)).

The vector part of the weak current is proportional to the electromagnetic current of the charmed particles, and the corresponding contribution to the cross section can be expressed in terms of the cross section for electroproduction of charmed particles. We shall assume that the electroproduction cross section has a Q^2-dependence of the form $\sigma_\gamma^c(Q^2) = [m_\psi^2/(Q^2 + m_\psi^2)]\sigma_\gamma^c(0)$ (see the discussion in Sec. 3i). We assume for simplicity that the matrix element of the axial-vector current is equal to the matrix element of the vector current. This hypothesis is justified in the parton model for large momentum transfers and energies.

Thus, it can be shown that

$$\frac{1}{2}(\sigma_{\nu N}^c + \sigma_{\bar{\nu}N}^c) = \frac{3G^2}{32\pi^3\alpha}[(-\cos^2\theta_W + \frac{5}{3}\sin^2\theta_W)^2 + 1]\, m_\psi^2 m_N E \sigma_\gamma^c(Q^2=0). \tag{6.59}$$

Unfortunately, the total cross section for photoproduction of charmed particles is also unknown experimentally (various theoretical estimates were given in Sec. 3 of this chapter). Taking the most optimistic estimate (6.33),

$$\sigma_\gamma^c(Q^2=0) = 10^{-2}\sigma_\gamma^{tot} \qquad (Q^2=0), \tag{6.60}$$

we have

$$\sigma_\nu^c \sim 0.1\sigma_\nu^{tot} \tag{6.61}$$

and the cross section for the production of charmed particles by neutral currents is rather large.

However, the derivation of (6.61) is not completely consistent. In fact, Eq. (6.60) was obtained in the framework of the vector-dominance model. The parton model evidently leads to a cross section which is somewhat smaller, i.e., $\sigma_\nu^c \sim 10^{-2}\sigma_\nu^{tot}$. On the other hand, it is inconsistent to retain the contribution of the axial-vector current in the framework of the vector-dominance model[25]. If $\sin^2\theta_W \sim 0.3$, this reduces the estimate (6.61) by a factor of about 6.

Thus, it seems more probable that the cross section for producing charmed particles is no more than 1–2% of the total cross section:

$$\sigma_\nu^c \leqslant 10^{-2}\sigma_\nu^{tot}. \tag{6.62}$$

We quote an estimate of the cross section for ψ-meson production (see also[25]):

$$\frac{1}{2}[\sigma(\nu N \to \nu\psi X) + \sigma(\bar{\nu}N \to \bar{\nu}\psi X)] \approx 6\cdot 10^{-42}E \text{ cm}^2, \tag{6.63}$$

where the energy of the neutrino (antineutrino) is measured in GeV. We see that the expected cross section for ψ-meson production in neutral currents is small. Experimentally, ψ production has not been observed in neutrino reactions.

f) **Allowance for the masses of charmed particles.** To compare the numerical estimates with experiment at the available energies, it would be important to take threshold effects into account. The parton mass is usually neglected. As a first approximation, let us now retain the mass of the charmed quark, assuming as before that the ordinary partons are massless.

If a neutrino collides with a parton which carries a fraction x′ of the nucleon momentum, the law of conservation of momentum takes the form

$$q + p_N x' = p_c, \tag{6.64}$$

where q is the momentum transferred to the hadrons, and p_c is the momentum of the charmed quark which is produced.

Squaring (6.64), we obtain

$$x' = \frac{Q^2 + m_c^2}{2m_N\nu}, \qquad (6.65)$$

where we have replaced the mass of the charmed quark by the mass of the lightest charmed particle, m_C (we take $m_C = 2$ GeV in what follows). The combination $x' = (Q^2 + m_c^2)/2m_N\nu$ also arises naturally in considering the lifetime of the virtual hadron state into which the W boson is transformed, so that the expression (6.65) is perhaps meaningful.

Instead of (6.54), we then obtain

$$\sigma_{\nu N}^c = \frac{G^2 s}{2\pi} \int\limits_{m_c^2/s}^{1} x' \left(1 - \frac{m_c^2}{sx'}\right)^2 \{2\lambda(x') \cos^2\theta_C + [n(x') + p(x')] \sin^2\theta_C\} dx',$$
$$m_c^2/sx' \leqslant y \leqslant 1, \qquad (6.66)$$

where the factor depending on m_C comes from the expression for the cross section for the process $\nu\lambda \to \mu c$, retaining the mass of the c quark but neglecting the other masses (we have also neglected the admixture of c quarks in the nucleon).

The differential cross section $d\sigma^c/dQ^2$ is frequently measured experimentally. The following expression for $d\sigma^c/dQ^2$ can be derived:

$$\frac{d\sigma^c}{dQ^2} = \frac{G^2}{\pi} \cos^2\theta_C \int\limits_{(Q^2+m_c^2)/s}^{1} \left(1 - \frac{m_c^2}{sx'}\right) \lambda(x') dx'$$
$$+ \frac{G^2}{2\pi} \sin^2\theta_C \int\limits_{(Q^2+m_c^2)/s}^{1} \left(1 - \frac{m_c^2}{sx'}\right) [n(x') + p(x')] dx'. \qquad (6.67)$$

The second term here corresponds to the interaction with the valence partons and contributes to the neutrino cross section, but not to the antineutrino cross section. The first term describes scattering by $\lambda\bar{\lambda}$ quark pairs (the distribution function of the λ partons appears for the neutrino, while that of the $\bar{\lambda}$ partons appears for the antineutrino).

To conclude this section, we quote numerical estimates of the cross section for charmed-particle production in the neutrino experiment at Batavia involving searches for specific violations of scaling associated with the production of heavy particles. Choosing $s = 80$ GeV^2 and $Q^2 = 1.4$ GeV^2 as the characteristic values of the kinematic variables, we have

$$\frac{d\sigma^c}{dQ^2} \Big/ \frac{d\sigma^{\text{tot}}}{dQ^2} = \begin{cases} 0.09, & s = 80\,\text{GeV}^2, \quad Q^2 = 1\,\text{GeV}^2, \\ 0.1, & s = 80\,\text{GeV}^2, \quad Q^2 = 4\,\text{GeV}^2 \end{cases} \qquad (6.68)$$

in the neutrino beam, and

$$\frac{d\sigma^c}{dQ^2} \Big/ \frac{d\sigma^{\text{tot}}}{dQ^2} = \begin{cases} 0.08, & s = 80\,\text{GeV}^2, \quad Q^2 = 1\,\text{GeV}^2, \\ 0.1, & s = 80\,\text{GeV}^2, \quad Q^2 = 4\,\text{GeV}^2 \end{cases} \qquad (6.69)$$

in the antineutrino beam. In obtaining this estimate, we have made use of the model calculations of the parton density carried out in[26b] under the assumption that $\lambda(x) = \bar{\lambda}(x) = \bar{n}(x) = \bar{p}(x)$, which evidently overestimates the λ and $\bar{\lambda}$ densities.

g) **Comparison of the theoretical predictions with the experimental data.** As we have already mentioned in Chap. III, certain new phenomena have been observed in the neutrino experiments as Batavia, namely muon pair production and a violation of scaling. Although these effects may of course have nothing to do with charmed-particle production, it seems natural to compare the theoretical expectations with the first experimental observations.

If we assume that muon pair production is actually due to the production and subsequent decay of charmed particles, then the observed magnitude of the effect corresponds to a rather large probability for leptonic decay of charmed particles. If the cross section for producing charmed particles is about 5% of the total cross section, then the probability for leptonic decay is $\sim 20\%$.

Theoretically, it is difficult to understand the observed violation of scaling. The effective-mass distribution of the hadrons produced in $\bar{\nu}$N collisions was studied in[27], and it was found that there is an abundance of events with hadron masses $W > 5$ GeV ($W^2 \approx 2m\nu$). If this enhancement is attributed to charmed-particle production, then

$$\frac{d\sigma^c}{dQ^2} \Big/ \frac{d\sigma_{\text{tot}}}{dQ^2} \Big|_{Q^2 \leqslant 1\,\text{GeV}^2} \sim 100\%, \quad \frac{d\sigma^2}{dQ^2} \Big/ \frac{d\sigma_{\text{tot}}}{dQ^2} \Big|_{Q^2 \sim 4\,\text{GeV}^2} \sim 50\%. \qquad (6.70)$$

We see that the theoretical estimates (6.69) deviate strongly from the data.

First, the relative contribution of the cross section for producing heavy particles exhibits a decrease with Q^2 instead of an increase. This can be understood theoretically only as a threshold effect, and in that case the effect may become stronger with increasing energy. Second, the magnitude of the effect is much greater than the estimate.

h) **Models involving "extra" quarks.** The most natural interpretation of the results of[27] is the production of new particles on the valence partons with a large coupling constant, i.e., one not containing the factor $\sin\theta_C$. It is then possible to understand the magnitude of the effect, but it becomes necessary to abandon the model involving the four quarks p, n, λ and c.

If we accept the formulas for the production cross sections of heavy particles in the parton model which were obtained above, the assumption that new particles are produced on the valence partons can be tested according to the form of $d\sigma/dQ^2$ or of the invariant-mass spectrum of the hadrons produced in $\bar{\nu}$N collisions. In order to make a comparison with experiment, relations of the type (6.67) must be integrated with respect to the experimental energy spectrum of the antineutrino. Considering the preliminary character of the data, we shall not make a detailed comparison of the theory with experiment, but we shall consider certain general properties of models involving charmed quarks.

Clearly, more than four quarks can be introduced in the theory. For example, the doublets with respect to the weak-interaction group $SU(2)_W$ in the Weinberg-Salam model might have the form

$$(p,\, n_1),\ (p',\, n_2),\ (p'',\, n_3) \qquad (6.71)$$

where p, p′, and p″ are the proton-type quark and two new quarks with charge 2/3, and $n_{1,2,3}$ are field combinations with charge $-1/3$. Thus, in addition to the usual p, n and λ quarks, three new quarks p′, p″ and n′ are introduced.

The known quarks might also belong not to doublets, but to triplets of the group $SU(2)_W$:

$$(p,\, n_\theta,\, \Delta),\ (p',\, \lambda_\theta,\, \Delta'), \qquad (6.72)$$

where $n_\theta = n\cos\theta_C + \lambda\sin\theta_C$, $\lambda_\theta = \lambda\cos\theta_C - n\sin\theta_C$, and Δ and Δ' are new quarks with charge $-4/3$.

By choosing the representations of the group $SU(2)_W$ in the form (6.71) or (6.72), one ensures that there are no neutral $\bar{\lambda}n$ currents to first order in the weak interaction and that these currents are suppressed in second order, as in the four-quark model. The prescription for constructing models which involve no λn transitions is simple: all linearly independent combinations of quarks having the same charge must belong to the same multiplet of the weak-interaction group.

As a rule, new leptons must be introduced in addition to the new quarks. This requirement is due to the experimentally observed universality of the beta-decay constants of the neutron and the muon. Since unique values of the coupling constants are calculated in gauge theories, the hadron and lepton multiplets must have the same structure in order to account for the universality of the weak constant. The introduction of new quarks therefore requires the introduction of new leptons.

An exception to this rule was pointed out by Barnett[28], who proposed that the experiment of[27] reveals an interaction with right-handed quarks. The usual two doublets (p, n_θ) and (p', λ_θ) are supplemented by the doublets

$$(p_R,\ n_R' \cos\varphi + \lambda_R' \sin\varphi)\,(p_R',\ n_R' \sin\varphi - \lambda_R' \cos\varphi), \qquad (6.73)$$

where n' and λ' are new quarks, and the index R indicates that the interaction has not the usual V − A form, but a V + A form. It is not necessary to introduce any new leptons in this case, and universality of the weak-interaction constant is saved here by the fact that right-handed leptons and quarks play a significantly different role.

The spectroscopy of ψ mesons and charmed particles in the six-quark model was studied by Harari[29]. Of the general properties of models involving "extra" quarks, we note here only that a larger asymptotic value than in the four-quark model is predicted for the ratio

$$R = \frac{\sigma(e^+e^- \to \text{hadrons})}{\sigma(e^+e^- \to \mu^+\mu^-)},$$

and that strange particles may not be produced here as a result of the decay of charmed mesons and baryons.

The future attitude to models involving extra quarks will evidently depend on the development of the experimental data.

VII. CONCLUDING REMARKS

At the present time (March 1975), experiment has not yet had the final say about the theoretical models which we have discussed. What will happen if experiment confirms the ideas which underlie the quark interpretation of the ψ mesons?

The discovery of α, β and γ rays at the end of the last century marked the beginning of the study of the strong, weak and electromagnetic interactions of elementary particles at high energies. Since then, investigations in these three areas have been interrelated, both experimentally and theoretically. As an example, it suffices to recall the history of how Dirac formulated quantum electrodynamics and Fermi formulated the theory of the weak interaction.

However, it is only in the past decade that there has been hope of creating a unified theory of all the interactions of elementary particles. If experiments actually confirm the existence of charmed particles, this will mean that there is a serious basis for these hopes.

The discoveries of the domain of strange particles a quarter of a century ago and the domain of resonances 15 years ago involved an extraordinary extension of our ideas about the world of hadrons. However, the theoretical significance of these discoveries was obscure at the time. In contrast, the discovery of charmed particles, if these particles are to be discovered, may mark the beginning of a phase of elementary-particle physics which is analogous to the phase in which a unified quantum theory of atoms and molecules was formulated 50 years ago.

Even if the outlines of a future unified theory of elementary particles which we can perceive at the moment are entirely illusory, even in this most unpretentious variant, the discovery of supercharged particles would greatly advance our understanding of many phenomena. These include the weak neutral currents, the $\Delta T = 1/2$ rule in non-leptonic decays, the violation of CP invariance, Zweig's selection rule, and the nature of the weak interactions at small distances.

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Translated by N. M. Queen