# SPECTROSCOPY AFTER THE NEW PARTICLES

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## 1. INTRODUCTION

Theorists have abandoned conventional spectroscopy and are busy inventing wrong theories for the new particles, so there have been no recent theoretical contributions on the conventional hadron spectrum. *Rosner's London review*[1] of the status of the conventional hadron spectrum is still up to date at this time. But the discovery of new particles[2] which do not fit into our previous picture of the hadron spectrum tells us something is missing in our description of hadron structure and hadron spectroscopy. We therefore reexamine conventional spectroscopy and look for the puzzles and paradoxes which have arisen in attempting to describe the properties of the known particles. These may offer clues to the missing elements necessary for the description of the new particles.

Predictions of the existence of new degrees of freedom from the properties of conventional particles have been discussed by Cabibbo[3] who compared the present theoretical situation to the irregularities observed in the orbit of the planet Uranus which led to the prediction and the discovery of the new planet Neptune. Our guide to searches for new planets, new particles, new symmetries and new degrees of freedom is based on the two key questions: (1) Who needs it and (2) who cares if it isn't found? Evaluation of a proposed theory must consider two *crucial aspects*: (1) the uniqueness of the theory and (2) the constraints imposed by external experimental information.

As examples of this guide we ask the key questions in the cases of (1) the search for the planet Neptune, (2) the search for the $\Omega^-$ and (3) the search for SU(3) as a higher symmetry for elementary particles.

The planet Neptune was needed by those who believed in Newton's laws of motion and gravitation and the precise astronomical observations of the orbits of other planets. If Neptune were not found either Newton's laws or the astronomical observations would have to be discarded. The prediction was unique; no competing models explained the irregularities in the orbit of Uranus. There were very many experimental constraints provided by the laws of motion and gravitation and by all the data on planetary motion which fit together to predict the existence of Neptune.

The $\Omega^-$ was needed by those who believed in SU(3) symmetry as a classification scheme for hadrons because SU(3) successfully classified 35 observed particles in four SU(3) multiplets with only one vacant slot remaining to be filled. If the $\Omega^-$ did not exist the symmetry scheme which had already fit the other particles would have to be discarded. The scheme was unique. No competing classification existed for the known particles. The existence of 35 particles having the proper eigenvalues of spin, parity, isospin and hypercharge to fit into those SU(3) multiplets provided a sufficient number of experimental constraints.

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The search for SU(3) was completely different from these two other searches. After isospin and hypercharge became established there were indications that a higher symmetry should exist to provide a classification scheme for particles having the same spin and parity. But there were very many possibilities and very few experimental constraints. The SU(3) symmetry group was found only after an eight-year journey via all the wrong symmetries.[4)] Today the transition from the SU(2) x U(1) of isospin and strangeness to SU(3) seems as simple as 2 + 1 = 3. But the theorists did not know they had SU(2). Isospin was considered to be a rotation in a 3-dimensional space and its natural generalization was rotations in spaces of higher dimensions, like R(4), R(5), R(6), R(7) and R(8). None of these higher rotation symmetries worked. The particle theorists needed to learn something new, namely the unitary groups which were not known to them at that time.

Today, particle theorists have learned all about the SU(n) groups and all about quarks, and the experimental situation suggests a higher symmetry beyond SU(3). The theorists add more and more quarks to make larger and larger SU(n) groups. At the present rate they may end up with SU(Rosenfeld) with all the particles in the Rosenfeld table classified in the fundamental representation of SU(n) where n is the number of particles in the table.

The quark model was born when the number of observed hadrons became larger and larger and the idea of making them all out of a small number of basic building blocks seemed attractive. But quarks have not been found. Instead experimentalists have found new phenomena not easily described by a simple model containing only three quarks. The theorists add more quarks: four, five, six, eight, nine, twelve and eighteen quark models have been suggested. Again there seem to be too many so-called elementary objects. Perhaps they are not elementary but are made of even smaller sub-quarks. As experiments do not find quarks, theorists find excuses for adding even more quarks. Perhaps this is not the right approach to go beyond SU(3).

Perhaps the quarks are like the fictitious 3-dimensional isospace which was so useful in the early days for understanding the implications of isospin symmetry. A realization of a given symmetry which is useful but has no real physical basis can impede the search for a higher symmetry. Maybe today we need to stop adding more quarks and to learn something new.

## 2. THE MINIMUM NUMBER OF ELEMENTARY BUILDING BLOCKS

The search for elementary building blocks for hadrons[5)] began with the Fermi-Yang model[6)] which suggested that the pion was a bound state of a nucleon-antinucleon pair. The discovery of strange particles suggested the addition of another basic building block to carry the strangeness quantum number. The Sakata model[7)] added the $\Lambda$ hyperon to the nucleon to make three basic building blocks. Goldhaber[8)] suggested the addition of the K meson as the basic strange building block, with the $\Lambda$ as a KN bound state. No detailed dynamics of such peculiar strongly bound systems was seriously proposed. The problem was to find the minimum number of basic building blocks from which all observed particles could be constructed with proper eigenvalues for all conserved quantum numbers.

The discovery of SU(3) multiplets with the baryons classified in an octet rather than a triplet led to the replacement of the Fermi-Yang-Sakata model by the quark model having as basic building blocks a triplet of spin-1/2 particles with baryon number 1/3 and the same isospin and strangeness quantum numbers as the Sakata triplet. The Goldhaber model was discarded.

The suggestion that a color degree of freedom[9,10)] was desirable to describe many hadron properties led to the expansion of the 3-quark model into the 9-quark model. The suggestion that charm[3,11)] is also needed to describe hadrons leads to even more quarks. These degrees of freedom which add new types of quarks can be described in a general way[12] by defining "generalized color" and "generalized charm." Generalized color denotes models with additional quarks that are all SU(3) triplets labeled by a new quantum number called "color," with an additional $SU(n)_{color}$ symmetry transformation between triplets of different colors. The requirement that all observed states be color singlets leads to a classification with exactly the same multiplet structure as the conventional 3-quark model. The original proposal has three colors but any number can be considered.

Generalized charm denotes models with additional quarks that are all SU(3) singlets and have new additive quantum numbers. The original charm proposal added a fourth quark and one new quantum number called charm and defined an SU(4) symmetry. Generalized charm adds n quarks and n new quantum numbers to make a system of $3 + n$ quarks and a higher symmetry $SU(3 + n)$. These additional quarks are postulated to have a high mass to explain why bound states containing charmed quarks are not observed in the conventional hadron spectrum.

We now return to the problem of finding the minimum number of basic building blocks needed to describe the observed states. Let us examine models like the Goldhaber model which have elementary bosons as well as fermions and transform their basic fermions into quark-like states with color and baryon number 1/3. The minimum number of basic states is achieved by choosing the $\Lambda$ and $K$ as the elementary building blocks, with the nucleon as a $\Lambda\bar{K}$ bound state. The quantum numbers of all the old particles can be constructed from these building blocks. When the spin and antiparticle states are counted, this model has a total of eight basic states in contrast to 12 in the Sakata and quark models.

The "quarkized" version of this $\Lambda K$ model has a strange $\lambda$ quark with baryon number 1/3 and a scalar or pseudoscalar boson with isospin 1/2 denoted by $\kappa$ as basic building blocks. The constituent nonstrange n and p quarks are bound $\lambda\bar{\kappa}$ states. All states of the old hadron spectrum can be constructed from these constituent quarks just as in the conventional quark model. In addition to the conventional states of the hadron spectrum constructed from constituent quarks, states of the $\kappa$ bosons without quarks can also be considered. These additional states can accommodate a spectrum of new particles. The $\kappa\bar{\kappa}$ system has two neutral states with isospin 0 and 1, respectively. These might correspond to the $\psi$ and $\psi'$ states. Note that the $\kappa$ has no SU(3) classification and the new particles formed from $\kappa$'s would not appear in SU(3) multiplets. Note also that the lowest state of this system is $0^+$ rather than $0^-$ as in a quark model.

In this picture the strange quark $\lambda$ is elementary and the nonstrange quarks n and p are composite, while SU(3) symmetry and current algebra do not exist at the fundamental level but arise as a phenomenological symmetry of constituent quarks, analogous to the Wigner supermultiplet symmetry in nuclei. Color can be introduced by coloring the strange quarks. The nine colored constituent quarks of the conventional three triplet model then appear as three elementary strange quarks and six composite nonstrange bound systems of a colored strange quark and a strange boson of isospin 1/2. This model has only five basic building blocks, a colored quark triplet and a boson isospin doublet. The fermions carry the color and the bosons carry the isospin. When antiparticle and spin states are counted this model has only 16 basic states instead of the 36 of the colored quark model.

Thus the minimum number of elementary building blocks necessary to produce the observed spectrum including the new resonances is much less than the number found in models with colored and charmed quarks. Models with elementary bosons as well as elementary fermions allow all states to be constructed with a much smaller number of basic particles. Such models might also be suggested by supersymmetries[13] which place bosons and fermions together in supermultiplets. The purpose of this discussion is not to present a serious argument for this particular model. It is rather to point out the possibility of new directions including those having no simple SU(3) description, to counteract brainwashing by quark models in which everything is made from elementary spin-1/2 objects, and to emphasize the possibility of a unified explanation of the new particles and the old particles.

## 3. CHARM AND COLOR

Let us now examine the charm and color degrees of freedom and ask who needs them.

### 3.1 Who Needs Charm?

People who like gauge theories like charm. But now that two kinds of weak interactions are observed experimentally charm is needed to unify weak interactions and weak interactions. The present quark structure cannot give a unified description of the two interactions produced by charged and neutral currents. The weak part of the Hamiltonian is written as the sum of two terms, one for the charged current and one for the neutral current,

$$H_{weak} = G_{ch} J^{\dagger}_{ch} J_{ch} + G_{neut} J^{\dagger}_{neut} J_{neut} \,. \tag{3.1}$$

But there is no way to relate the strengths of the two interactions. For example the requirement that $G_{ch}$ and $G_{neut}$ should be equal has no meaning because there is no unique normalization for the charged and neutral hadron currents relative to one another and no unique normalization relative to the lepton current. The Cabibbo theory provides such a normalization for the charged currents but cannot work for strangeness conserving neutral currents. This is most easily seen by examining the linear combinations of $n$ and $\lambda$ quarks rotated by the Cabibbo angle

$$n_c \equiv n \cos\theta + \lambda \sin\theta \tag{3.2a}$$

$$\lambda_c \equiv n \sin\theta - \lambda \cos\theta \tag{3.2b}$$

where $\theta$ is the Cabibbo angle. The charged current has components with the form

$$J_{ch} = \cos\theta(\overline{p}n) + \sin\theta(\overline{p}\lambda) = \overline{p}n_c \tag{3.3}$$

while the particular state $\lambda_c$ defined by Eq. (3.2b) is completely decoupled from the charged current weak interaction. The leptonic part of the charged current has exactly the same form and same normalization as the charged hadron current (3.3).

The neutral current which conserves strangeness cannot be expressed entirely in terms of the states $p$ and $n_c$ but must also include a contribution from the $\lambda_c$ in order to conserve strangeness. The most general strangeness-conserving neutral current is most easily written in terms of the original unrotated $(pn\lambda)$ states

$$J_{neut} = \alpha(\overline{p}p) + \beta(\overline{n}n) + \gamma(\overline{\lambda}\lambda) \tag{3.4}$$

where the coefficients $\alpha$, $\beta$ and $\gamma$ are arbitrary. There is no obvious way to determine them or to relate them either to the lepton currents or to the charged current.

The necessity or desirability of unifying the charged and neutral weak interactions suggests that there is something missing in this simple quark description (3.3) and (3.4). The addition of a fourth charmed quark enables a very simple and elegant description[3] in which the charged and neutral currents have a well-defined relative normalization and the unification of the two interactions is straightforward. However, this SU(4) description is by no means unique and does not satisfy the criterion of having many experimental constraints. There are no experimental constraints forcing such an SU(4) description. Many other possibilities exist with no experimental indication that any one is superior to the others. The SU(4) charm scheme is only one of many ways to introduce a new internal degree of freedom suggested by the presence of neutral strangeness-conserving currents together with the Cabibbo charged strangeness-violating currents.

### 3.2 Who Needs Color?

Many reasons have been proposed for introducing color, and not all of them are compatible. Color is needed by

1) People who like ordinary fermi statistics for quarks[9] and do not like baryon models with three spin-1/2 quarks in symmetric rather than in antisymmetric states.

2) People who like integral electric charge.[10]

3) People who believe Adler's argument for color[14], based on the current-algebra-PCAC calculation of the decay $\pi^0 \to \gamma\gamma$. Adler's result is proportional to the sum of the squares of the charges of all elementary fermions in the theory. The numerical experimental value for the width of this decay agrees with predictions from a 3-color model and disagrees with models having no color degree of freedom.

4) People who want to push up the ratio $R \equiv e^+e^- \to \text{hadrons}/e^+e^- \to \mu^+\mu^-$, whose present experimental value exceeds the prediction from the simple quark model.[15] The addition of new internal degrees of freedom pushes this ratio up, just as in $\pi^0 \to \gamma\gamma$.

5) People who worry about the saturation of hadrons at the quark-antiquark and three-quark levels and want a model which explains why states like $qq\bar{q}$ and $4q\bar{q}$ are not found. Colored models provide a natural description of this saturation.[12,16]

6) People who like non-Abelian gauge theories and quark confinement.[17] However these people require the color symmetry to be an exact symmetry of nature not broken by weak or electromagnetic interactions. They are unable to incorporate integrally-charged quarks into this framework and must have fractional charges.

7) People who want to explain the new particles with colored states that are not color singlets.[18] They require the Han-Nambu integrally-charged model in order to explain the production of the particles in $e^+e^-$ annihilation and therefore require breaking of color symmetry by electromagnetism.

There is a definite conflict between those who want to explain the new particles as color-octet states and those who want to use color in a non-Abelian gauge theory as an exact symmetry of nature. There is also a conflict between the explanation of new particles as simple color-octet states and the success of the color model in explaining saturation at the quark-antiquark and three-quark system. This point is discussed in detail in the following section.

## 4. THE COLORED QUARK MODEL FOR SATURATION

How color can lead to saturation in the hadron spectrum is illustrated by an instructive example[12,16] of a fictitious nuclear physics world with all observed nuclei composed of deuterons, all observed states with isospin zero and the neutron and proton not yet discovered. The experimentally observed properties of the deuteron would indicate that it is a two-body system, whose constituents might be called nucleons. In this isoscalar deuteron world, the isospin degree of freedom would not be known, only isoscalar properties of states would be measurable and there would be no way to distinguish between the neutron and the proton. Thus the two nucleons in the deuteron would appear to be identical particles, with the deuteron electric charge divided evenly between them. Each nucleon would thus have a charge of 1/2.

$$Q_N = 1/2 Q_d = +1/2 \ . \tag{4.1}$$

The ground state of the deuteron and all the observed isoscalar excited states have wave functions symmetric in space and spin rather than antisymmetric as expected for two fermions. Thus the nucleon appears to be a particle with fractional charge and peculiar statistics.

At this point a clever theorist might suggest that there are really two kinds of nucleons with different colors. The red nucleon denoted by $p$ has a charge of $+1$ and the blue nucleon denoted by $n$ has charge zero. A hidden SU(2) symmetry transforms between different colored nucleon states and all the low-lying states would be singlets in this $SU(2)_{color}$ degree of freedom. Their wave functions would be antisymmetric in color and Fermi statistics would require them to be symmetric in space and spin. Thus the colored nucleon would have integral charges and satisfy ordinary Fermi statistics.

This colored deuteron world is an exact analog of the colored quark model for hadrons. Consider the $\Omega^-$ which consists of three identical strange quarks in the quark model and has a spin of 3/2 and an electric charge of $-1$. The charge of the $\Omega^-$ is assumed to be divided equally between the three strange quarks. Thus

$$Q_\lambda = 1/3 Q_{\Omega^-} = -1/3 \ . \tag{4.2}$$

The spin-3/2 state is totally symmetric in spin and a totally symmetric spatial wave function is assumed for the lowest state. Thus the strange quark appears to have a fractional charge and peculiar statistics.

Now suppose that there are three kinds of $\lambda$ quarks denoted by $\lambda_1$, $\lambda_2$ and $\lambda_3$ with three different colors. If there is a hidden $SU(3)_{color}$ symmetry and all low-lying states are required to be singlets in $SU(3)_{color}$, they are antisymmetric in the color degree of freedom and required by Fermi statistics to be symmetric in the other degrees of freedom. Thus introducing color eliminates the peculiar statistics. Integral charges are also obtainable by setting

$$Q_{\lambda_1} = Q_{\lambda_2} = 0 \tag{4.3a}$$

$$Q_{\lambda_3} = -1 \ . \tag{4.3b}$$

We have chosen the $\Omega^-$ for this simple example because it contains only strange quarks. The same arguments hold for baryons containing both non-strange and strange quarks. These have quark model wave functions totally symmetric in the SU(6) and spatial degrees of freedom. They satisfy normal statistics if they are required to be antisymmetric in an additional color degree of freedom.

The assumption that only color singlet states are observed leads to exactly the same predictions for the classification of hadron states as the conventional quark model without color. For all observed properties of conventional hadrons, there is no difference between models with integrally-charged and fractionally-charged quarks. The strong interactions of both kinds of charged colored quarks are identical. Differences arise only in the electromagnetic and possibly in weak interactions, and are observable only in states which are not color singlets. Since our discussion of saturation depends only on strong interactions it holds equally for integrally-charged and fractionally-charged quarks.

The colored quark model also answers three puzzles which are unsolved in the conventional nonrelativistic quark model. These are:

1) The triality puzzle. The quark-antiquark interaction is attractive in all possible channels as indicated by the existence of bound quark-antiquark states for all possible values of quantum numbers. An antiquark should therefore be attracted by the three quarks in a baryon to make a state of three quarks and one antiquark. No such bound states with non-zero triality are observed.

2) The exotics puzzle. The above argument also holds for states like $qqq\bar{q}$ or $4q\bar{q}$ which have proper triality and are not excluded by any new principle preventing the observance of fractionally-charged states. The attractive interactions necessary to bind three quarks into a baryon and a quark-antiquark pair into a meson should bind these exotic states even more strongly than normal hadrons and predict unobserved exotic mesons and baryons near or below the pion and nucleon masses.

3) The meson-baryon puzzle. The quark-antiquark and quark-quark interactions must both be attractive to bind both mesons and baryons. However the quark-quark interaction must be considerably weaker than the quark-antiquark interaction to prevent di-quarks from being observed with masses comparable to mesons. Low-lying bound states occur in the three-quark system which has three quark-pair interactions, rather than only one as in a meson. Conventional simple interactions have quark-quark and quark-antiquark interactions which are equal in magnitude and are either opposite in sign or both attractive, depending on the behavior of the interaction under charge conjugation. The interaction necessary to bind mesons and baryons does not have a simple behavior under charge conjugation but rather requires that the even and odd components satisfy a very peculiar relation.

All these puzzles have been shown to be solved by a model with three triplets of colored quarks interacting via the Yukawa interaction produced by exchange of an octet of colored vector gluons. In this model mesons and baryons are both bound with the desired interaction strength. The difference between the quark-quark and quark-antiquark interactions is given naturally by an algebraic coefficient from the color algebra. States which are singlets in $SU(3)_{color}$ behave like neutral atoms with no interaction with any external quarks. The individual components of the bound colored singlet states have both attractive and repulsive interactions with an external particle. These cancel exactly as in the case of a neutral atom. Thus there are no bad triality states and no exotic states in this model, only the desired states appear -- or more correctly -- only the states that were desired before the discovery of the new particles.

These colored quark models bound by an octet of colored gluons cannot simply incorporate the new particles. The color octet states of a quark-antiquark pair which have been suggested as possible configurations for the new particles[18)] are unbound in this particular model because the quark-antiquark interaction is repulsive in all color octet states. This repulsion is needed to prevent an antiquark from being bound to a baryon to make a $3q\bar{q}$ state. If an attractive interaction is postulated in certain color octet channels to provide states for the new particles, there is no longer a cancellation between attraction in color singlet and repulsion in color octet states and the color singlet baryons no longer behave as neutral atoms. Thus there is a conflict between the saturation properties of the colored-gluon-Yukawa interaction and the possibility of describing the new particles as color octet bound states. With simple interactions one can have either saturation or a description of the new particles, but not both. The models which propose color for the new particles disregard the saturation question for ordinary particles. The difficulty can be overcome by more exotic models as discussed below.

## 5. SPIN SPLITTINGS IN THE MESON SPECTRUM

The discovery of new particles of spin one and the absence so far of new spin-zero particles raises the question of why the lowest new particles should be vectors rather than pseudoscalars as in the conventional spectrum. However, this question can be reversed to why the lowest in the conventional spectrum are pseudoscalars. There is no answer to this question. All models describe the spin splittings with an external parameter inserted by hand. Thus there may be some interest in examining spin splitting in the conventional hadron spectrum to look for clues in understanding the new particles. These spin splittings are indeed very peculiar. The $\rho$-$\pi$ splitting which is an s-wave hyperfine splitting in atomic or positronium language is of the same order as the $\rho$-A2 splitting which is an orbital splitting. The splitting between the A2 and the B which is a hyperfine splitting in the p wave is consistent with zero. Thus the s-wave hyperfine splitting is large and of the same order as the orbital splitting while the p-wave hyperfine splitting is zero. These are not explained in the conventional quark model but simply inserted as external spin-dependent interactions each with its own strength parameter.

All these peculiar spin splittings can be explained by adding a repulsive core[19] to the quark-antiquark potential in the triplet spin state. Such short-range hyperfine interactions have been suggested[20], but are normally treated by perturbation theory which neglects the triplet-singlet splitting resulting from differences in the wave functions, and considers only the interaction. A strong repulsive core which cannot be treated by perturbation theory gives an energy shift from the added kinetic energy of the wave function forced to vanish at the origin. This energy is on the scale of the orbital splitting and independent of the strength of the core. This effect occurs only in the s wave where the wave function is otherwise appreciable at the repulsive core and not in the p wave where the wave function already vanishes at small distances.

A rough estimate of the effect of a repulsive core on the meson spectrum is obtained from a harmonic oscillator model with the effect of the core in the $^3S$ state simulated by using a wave function with s-wave angular dependence and p-wave radial dependence to make it vanish at the origin. This leads to the following spectrum

$$E_\pi = \frac{3}{2}\,\hbar\omega \quad \text{(normal s wave)} \tag{5.1a}$$

$$E_\rho = \frac{11}{6}\,\hbar\omega \quad \text{(s wave with p-wave radial dependence)} \tag{5.1b}$$

$$E_\delta = E_{A1} = E_{A2} = E_B = \frac{5}{2}\,\hbar\omega \quad \text{(normal p wave)}\ . \tag{5.1c}$$

The $\rho$ thus appears roughly midway between the $\pi$ and the p-wave states.

We now examine the effect of such a repulsive core on the interaction in color octet states in a colored quark model[16]. If the repulsive core has the same color behavior as the rest of the interaction it has opposite signs in the color octet and color singlet states. The quark-antiquark interaction in the $^3S$ color octet state then has a very strong short-range attraction and a repulsive potential barrier. This leads to a color octet of bound vector mesons having a much smaller size than conventional mesons. There are no bound color octet pseudoscalar states since the repulsive core is not present in the $^1S$ color singlet pseudoscalars and no corresponding attractive interaction appears in the $^1S$ color octet state. If the new particles are such "collapsed" states of much smaller radial size than conventional hadrons the difference in the wave functions could explain the suppression of radiative transitions between collapsed new particles and normal states and the smaller photoproduction cross sections for the new particles.

Similar arguments have been given in the past for the nucleon-nucleon interaction. The original Fermi-Yang model for the pion was supported by the argument that the repulsive core present in the nucleon-nucleon interaction would reverse sign and become attractive in the nucleon-antinucleon interaction and could lead to bound states. The possibility that collapsed states of complex nuclei could be produced by a very short-range attractive interaction within the repulsive core[21] has also been suggested.

Similar qualitative effects are obtained from a strong short-range hyperfine interaction which is attractive in the singlet spin state and repulsive in the triplet spin state, rather than being completely absent in the singlet state. The attractive short-range potential in the spin singlet state becomes repulsive in the color octet state and may bind collapsed vector states but not collapsed pseudoscalars.

This discussion gives another example of how considering old and new particles together and using the unexplained puzzles of the conventional hadron spectrum can lead to new approaches in understanding the new particles.

## 6. THREE KINDS OF QUARKS

There are many puzzles and paradoxes behind the usual assertion that the quark model gives a very good description of the hadron spectrum. Everyone has his own quark model which is different from other quark models. Each gives a good description of some aspects of the hadron spectrum but not of others, and different approaches appear to be incompatible. This paradox is illustrated by the question, "Where are the friends of the $\rho$ meson?". Which are the isovector mesons which belong in the same family or supermultiplet as the $\rho$?

The SU(6) classification puts the $\rho$ and the $\pi$ in the same supermultiplet. Current algebra relates vector and axial vector currents and relates vector particles ($\rho$) with the axial vector (A1). Exchange degeneracy and duality place the $\rho$ and the A2 on degenerate Regge trajectories. Chiral symmetry describes pion emission with the axial charge operator $Q_5$. Thus any state A coupled to the $\rho\pi$ system has a nonvanishing matrix element $\langle A|Q_5|\rho\rangle$.

$$A \to \rho + \pi \to \langle A|Q_5|\rho\rangle \neq 0 \tag{6.1a}$$

$$\rho \to A + \pi \to \langle A|Q_5|\rho\rangle \neq 0 \,. \tag{6.1b}$$

Since $Q_5$ is a generator of the chiral symmetry algebra any resonance A which satisfies Eq. (6.1a) or Eq. (6.1b) must be in the same chiral symmetry multiplet as the $\rho$. Similarly, for any state A' coupled to the $A\pi$ system

$$A' \to A + \pi \to \langle A'|Q_5|A\rangle \neq 0 \,. \tag{6.1c}$$

The $\rho$, $\pi$, $\omega$, A1, B, A2 and $\eta$ all satisfy the condition (6.1a), (6.1b) or (6.1c). They must all be in the same chiral symmetry multiplet together with all higher states which decay to these by successive pion emission. This paradox has been treated by representation mixing in which each particle has many components classified in different chiral symmetry representations.

All these pictures are partially correct. Each one gives a mass formula relating the $\rho$ mass to the masses of other bosons but each picture involves different bosons. How do we put these different pictures together?

The hydrogen atom provides an instructive example of three completely different approaches to composite models with unknown basic building blocks. Suppose the proton and the electron had not yet been discovered but the low-lying states of hydrogen spectrum had been measured by studying the emission and absorption of low-energy photons. One can imagine the following trialog between an atomic physicist, an S-matrix theorist and a field theorist regarding the structure of the hydrogen atom.

Atomic physicist: The hydrogen atom consists of a proton and electron which will eventually be discovered. The hydrogen spectrum is given by solving a Schroedinger equation with a Coulomb interaction. This model also gives correct values for radiative transitions. The relativistic version with the Dirac equation gives the fine structure splittings in agreement with experiment.

S-matrix theorist: Nobody has ever seen an electron or a proton. The observable quantities are photons, the normal stable state of the hydrogen atom and the S matrix for photon-hydrogen scattering. The objects which atomic physicists call excited states are poles in the photon-hydrogen scattering amplitude. Protons and electrons are fictitious objects which will never be found. They are only useful mathematical objects in drawing diagrams to explain the observed quantum numbers of the hydrogen atom and to explain the absence of states having exotic quantum numbers.

Field theorist: A correct description must begin with a Lagrangian for the electron, proton and electromagnetic fields whose quanta are the partons that make up the hydrogen atom. The true wave function for the hydrogen atom contains very many partons including electrons, antielectrons, protons, antiprotons and the photons which are vector gluons whose interactions bind the partons together. Consistent calculations using quantum electrodynamics and perturbation theory give exact descriptions of all properties of the hydrogen atom. The atomic physics approach is a crude approximation which considers only the valence particles, the electron and the proton. It gives reasonable results for some observable quantities because $\alpha$ is a small number. But the full apparatus of the field theory is necessary to calculate effects of the infinite parton sea, like the Lamb shift and vacuum polarization.

In the case of the hydrogen atom the field theorist has the ultimate correct description. But in hadron physics where interactions are strong and unknown the question is still open. Each approach works in some areas of hadron dynamics but not in others. Field theory and its parton model has been useful in treating high energy deep inelastic and hadron-hadron scattering, but not for hadron spectroscopy and intermediate energy scattering where the atomic physics and S-matrix approaches have been very successful.

The difference between the three approaches to hadron structure is seen explicitly in comparing the three approaches to spin.

1. The static approach of the atomic physicist describes spin with Pauli matrices or the Dirac matrices $\sigma$. Static spin combined with SU(3) gives the static SU(6) symmetry which provides a successful classification of the hadron spectrum. However static SU(6) fails as a symmetry for transitions since it forbids the most common strong decays.

$$\left.\begin{array}{l}\rho \not\to 2\pi \\ \Delta \not\to N+\pi\end{array}\right\} \text{ in static SU(6) .} \qquad \begin{array}{r}(6.2a)\\(6.2b)\end{array}$$

2. The chiral approach of the field theorist and parton model considers chirality defined by the Dirac matrix $\gamma_5$ which is equivalent to $\sigma_z$ and helicity in the infinite momentum frame with $p_z = \infty$. This description is

good for the applications of current algebra and PCAC. However it is bad for classification because there is no simple way to include rotational invariance and require states of different helicity of the same particle to have the same mass and it suggests the existence of parity doublets which are not found experimentally.

3. The Regge approach of the S-matrix theorist places particles on trajectories along which the spin changes. The Regge approach gives a different spin spectrum from SU(6) and suggests the existence of particles on daughter trajectories.

The folklore which says that the harmonic oscillator quark model of the atomic physicist and the linear Regge trajectories of the S-matrix theorist agree in their predictions of the hadron spectrum is very misleading. The contradictions are immediately evident when spin is considered. The SU(6) harmonic oscillator quark model has degenerate $\pi$ and $\rho$ masses. They are not on degenerate Regge trajectories since their spin is different; rather they are required to be on different trajectories with a spacing of one unit. Degenerate $\pi$ and $\rho$ trajectories would not give degenerate masses. Neither the masses nor the trajectories should be expected to be exactly degenerate. But the puzzle is that the real world is exactly half way in between these two cases. Neither the masses nor the trajectories are degenerate. The spacing between the $\pi$ and $\rho$ trajectories is neither zero nor one but 1/2. Why?

Static and chiral approaches to spin have been combined by the definition of W spin and $SU(6)_W$. The spin flip operators in the W spin scheme are defined to flip both $\gamma_5$ and $\sigma_z$.

$$W_\pm \equiv \beta\sigma_\pm \tag{6.3}$$

These operators flip helicity in the infinite momentum frame and are equivalent to ordinary spin operators in the rest frame except for a phase factor. They are invariant under Lorentz transformations in the z direction and useful for all processes where a Lorentz frame can be defined with all momenta in the z direction. W-spin conservation allows the strong decays (6.2) forbidden by static spin conservation. The associated $SU(6)_W$ symmetry leads to good results for transitions as long as there is no transverse momentum. However the W-spin formulation breaks down in the presence of transverse momentum. Hadron states with finite orbital excitation always have transverse momentum, since

$$L_z \neq 0 \rightarrow p_x,\ p_y \neq 0\ . \tag{6.4}$$

## 7. WHO NEEDS MELOSH? WHERE IS JACKSON?

The Melosh[1,22] transformation has been suggested as an answer to all the questions regarding spin and transverse momentum, but papers on the Melosh transformation tend to obscure simple physics with complicated formalism. The relevant question is "Where is Jackson?".

Consider the emission of a pion by a quark. The process is particularly simple in the rest frame of the final quark state with the z axis in the direction of the momenta of the incoming quark and the outgoing pion. For virtual pion exchange in a scattering process this frame is commonly called the Jackson frame. Since the pion is spinless and carries no angular momentum in the z direction the z component of the total angular momentum is equal to the z-component of the quark spin in this frame. Thus $\sigma_z$ cannot flip and is conserved. Applying SU(3) and parity conservation shows that $SU(6)_W$ is automatically conserved for pion emission in the Jackson frame of the quark.

But where is Jackson? Experiments see hadrons not quarks. They measure hadron momenta and not the quark momenta inside the hadron. They cannot find the quark Jackson frame and present their results in the hadron Jackson frame where there is no transverse momentum in initial and final states of the hadron which emits the pion. But this frame is not the quark Jackson frame and the quarks inside the hadron can have transverse momentum. The Melosh school say they cannot find Jackson either but that there is an unknown transformation between the hadron Jackson frame and the quark Jackson frame.

Big deal.

If you can't find Jackson you have to work in a frame where helicity flip is allowed and $SU(6)_W$ is broken. Two kinds of terms thus appear in the description of pion emission and give two unrelated reduced matrix elements, one for nonflip transitions which conserve $SU(6)_W$ and one for flip transitions which violate $SU(6)_W$. If we could find Jackson we could relate these two matrix elements. But so far no one knows how to find Jackson. The existence of these two independent terms was already known in the first applications of the quark model to pionic decays. The two were then called direct and recoil terms.

Who needs Melosh?

## 8. BEYOND THE SINGLE-QUARK TRANSITION -- THE ZWEIG RULE MYSTERY

The big puzzle in the quark model description of hadron transitions has not been considered by the Melosh school, namely, the success of the description of all transitions by single quark operators. This Levin-Frankfurt rule[5,23], first applied to high energy scattering to give the ratio of 3/2 for nucleon-nucleon to pion-nucleon scattering, postulates that only one active quark in each hadron is responsible for any transition and all other quarks are spectators which remain in the same state. This single quark rule has been applied with uniform success but without justification for transitions which cover the complete spectrum of strong, electromagnetic and weak transitions. The active quark can emit or absorb a pion, a photon, a lepton pair or a Reggeon. The Zweig rule[24] which says that a quark-antiquark pair cannot disappear in a transition is a special case of the Levin-Frankfurt rule which requires either the quark or the antiquark to be a spectator and remian in the final state. Thus both cannot disappear. But both the Levin-Frankfurt rule and its special case of Zweig's rule can be defined only in some kind of Born approximation. In any theory these rules are violated in higher orders by a succession of allowed transitions. Consider for example the Zweig forbidden decay of the f' into two pions

$$f' \not\to 2\pi \,. \tag{8.1}$$

But the following transitions are all allowed by Zweig's rule.

$$f' \to K\bar{K} \tag{8.2a}$$

$$K\bar{K} \to \pi\pi \,. \tag{8.2b}$$

Thus in higher order we can have

$$f' \to K\bar{K} \to \pi\pi \,. \tag{8.3}$$

There is no way to forbid the transition (8.3) if the transitions (8.2) are allowed. The only way to suppress the transition (8.3) is to claim that it is somehow of higher order and therefore smaller than the Born term. But why should a Born approximation be valid for strong interactions?

To obtain some insight into the mysterious validity of the single quark transition we look for cases where this approximation breaks down and we see the higher order two-quark transition. The best place to look is in total cross sections for high energy scattering where precise data are available at the 1% level and small breaking effects of the Levin-Frankfurt approximation can be found. This approximation, also called quark-model additivity in this context, breaks down at the 15-20% level in total cross sections[5,25]. For example,

$$\sigma(\pi p) < (2/3)\sigma(pp) \qquad (8.4a)$$

rather than being equal as required by the Levin-Frankfurt approximation. However there is also a consistent difference of the same order between $\pi N$ and $KN$ total cross sections

$$\sigma(Kp) < \sigma(\pi p) \ . \qquad (8.4b)$$

Recently an empirical relation between these inequalities has been shown to agree with experiment over a wide energy range[25].

$$\sigma(\pi^- p) - \sigma(K^- p) = 1/3\sigma(pp) - 1/2\sigma(K^+ p) \ . \qquad (8.5)$$

The agreement with experiment[26] is shown in Fig. 8.1. The two inequalities (8.4) thus seem to be empirically related even though one is a breakdown of quark-model additivity and the other is a breakdown of SU(3) symmetry. The relation (8.5) can be derived from double exchange models described by two-quark operators which discriminate against strangeness. Examples of such models are the double exchange of a pomeron and an $f$ as a cut or as two legs of a triple Regge diagram. However the experimentally observed energy dependence is not that of a pomeron-$f$ cut. The difference $\sigma(\pi^- p) - \sigma(K^- p)$ which should be due to this additional contribution decreases very slowly from 5 mb at 2 GeV/c to 3.5 mb at 200 GeV/c. A triple-Regge diagram might give this energy dependence[27].

Thus another intriguing puzzle is added to the question of why the single-quark transition works so well, namely, why does strangeness dependence of the total cross section seem to be related to quark number dependence. This reminds us that we really do not understand strangeness. Attributing the difference between strange and nonstrange particles to constituent strange and nonstrange quarks simply shifts the question to the quark level. Perhaps the answer is in a model like the one proposed in Sect. 2 which gives a different structure to strange and nonstrange particles. In this model with elementary quarks and elementary bosons a double-scattering contribution with one scattering off any quark in the hadron and the second off any boson satisfies the relations (8.4) and (8.5).

## 9. NEW PARTICLES AND OLD SYMMETRIES

"This year's sensation is next year's calibration"...remark by V. L. Telegdi one year after the discovery of CP violation when kaon beam experimentalists were already using the CP-violating $2\pi$ decay mode to calibrate their apparatus.

The rapidly accumulating experimental data on production and decay of the new particles furnish a new laboratory for the study of old symmetries, quark models and empirical rules like Zweig's rule. For such high mass states all quasi-two-body and quasi-three-body decay channels related by symmetries are open and well above threshold. These final states have such large momenta that kinematic breaking effects of thresholds, barrier factors and mass differences can be neglected. This rich source of experimental data with many related decay modes should enable conclusive and significant tests of theoretical predictions.

One important open question in conventional tests of SU(3) symmetry is whether the observed suppression of kaon production relative to pion production can all be explained by kinematic factors resulting from the $K\pi$ mass difference, or whether there is also an inherent SU(3) breaking. A clean test of this point is obtainable by comparing $2\pi$ and $K\bar{K}$ decay modes of high mass states. An SU(3) singlet state must have $C = G = +$ in order to decay into two pions. There are both theoretical and experimental indications that such states should exist as new scalar and tensor mesons. For such a state SU(3) predicts

$$\Gamma(\chi_{singlet} \to \pi^+\pi^-) = \Gamma(\chi_{singlet} \to K^+K^-) \quad (9.1a)$$

Any deviation from this prediction would indicate either an inherent SU(3) breaking in the decay vertex function or the presence of an octet component mixed with the singlet in the initial state.

If the new particle is not an SU(3) singlet but a member of a nonet other members should be found. Equality between pionic and kaonic decay modes is predicted for the sum of the decay rates over all neutral nonstrange members of a nonet or octet.

$$\sum_{\chi} \Gamma(\chi \to \pi^+\pi^-) = \sum_{\chi} \Gamma(\chi \to K^+K^-) \quad (9.1b)$$

where the summation is over all neutral nonstrange states in the octet or nonet. The sum rule (9.1b) holds in the SU(3) symmetry limit independent of nonet mixing angles. It could be broken by mass differences between the different $\chi$ states in the octet or nonet. However such mass breaking effects are easily taken into account without introduction of arbitrary factors.

The relations (9.1) can also be obtained from U-spin invariance without requiring the full SU(3). Thus electromagnetic contributions to the decays should also satisfy these relations.

The new vector particles have $C = G = -1$ and are forbidden to decay into two pseudoscalar mesons if they are SU(3) singlets. The best test of the equality of strange and nonstrange couplings for such states is in the vector-pseudoscalar decay mode. The relations analogous to (9.1) are

$$\Gamma(\psi_{singlet} \to \rho^+\pi^-) = \Gamma(\psi_{singlet} \to K^{*+}K^-) \quad (9.2a)$$

$$\sum_{\psi} \Gamma(\psi \to \rho^+\pi^-) = \sum_{\psi} \Gamma(\psi \to K^{*+}K^-) \quad (9.2b)$$

The baryon-antibaryon decay modes can also be used

$$\Gamma(\psi_{singlet} \to \bar{p}p) = \Gamma(\psi_{singlet} \to \Sigma^+\overline{\Sigma^+}) \quad (9.2c)$$

$$\sum_{\psi} \Gamma(\psi \to \bar{p}p) = \sum_{\psi} \Gamma(\psi \to \Sigma^+\overline{\Sigma^+}) \quad (9.2d)$$

Equations (9.2) also follow only from U-spin invariance and apply to electromagnetic contributions.

If the relations (9.1) and (9.2) are found to be valid experimentally, decay modes involving the $\eta$, $\eta'$, $\omega$ and $\phi$ states can also be studied to determine values for mixing angles and to test the validity of Zweig's rule.

Quasi-three-body decays are also simply related by SU(3) and additional constraints are imposed by Zweig's rule. The $\rho K\bar{K}$, $\omega K\bar{K}$, $K^*\bar{K}\pi$, $\phi K\bar{K}$ and $\phi\pi\pi$ decays are all related by SU(3) to the observed $\omega\pi\pi$ decay and should provide significant tests. Zweig's rule forbids the $\phi\pi\pi$ decay mode for any initial singlet or nonet state.

$$\Gamma(\psi \to \phi\pi\pi) = 0 \quad \text{by Zweig's Rule .} \tag{9.3}$$

Thus the strength of the $\phi\pi\pi$ decay mode immediately gives an indication of the validity of Zweig's rule.

For an initial state which is an SU(3) singlet there are two SU(3) invariant amplitudes if Zweig's rule is not assumed, corresponding to the octet and singlet states of the vector mesons. All decay amplitudes can be expressed in terms of two which are taken as input. The $\omega\pi\pi$ and $\phi\pi\pi$ modes are convenient for input since the former has been observed and the latter should vanish if Zweig's rule holds. The SU(3) predictions are then[28)]

$$A_0(\phi K^+K^-) = -(1/\sqrt{2})A_0(\omega\pi^+\pi^-) \tag{9.4a}$$

$$A_0(\omega K^+K^-) = -(1/2)A_0(\omega\pi^+\pi^-) - (1/\sqrt{2})A_0(\phi\pi^+\pi^-) \tag{9.4b}$$

$$A_0(\rho^0K^+K^-) = -(1/2)A_0(\omega\pi^+\pi^-) + (1/\sqrt{2})A_0(\phi\pi^+\pi^-) \tag{9.4c}$$

$$A_0(K^{*0}K^-\pi^+) = -(1/\sqrt{2})A_0(\omega\pi^+\pi^-) + A_0(\phi\pi^+\pi^-) \tag{9.4d}$$

$$A_0(\rho^0\pi^0\eta_8) = -(1/\sqrt{3})A_0(\omega\pi^+\pi^-) + (\sqrt{2/3})A_0(\phi\pi^+\pi^-) \tag{9.4e}$$

$$A_0(\phi\eta_8\eta_8) = +(2\sqrt{2/3})A_0(\omega\pi^0\pi^0) - (1/3)A_0(\phi\pi^0\pi^0) \tag{9.4f}$$

$$A_0(\omega\eta_8\eta_8) = +(1/3)A_0(\omega\pi^0\pi^0) + (2\sqrt{2/3})A_0(\phi\pi^0\pi^0) \tag{9.4g}$$

where $A_0(VPP)$ denotes the amplitude for the decay of a unitary singlet into any particular VPP state. The final state can be defined either by the momenta of the three mesons and the polarization of the vector meson or by any partial wave amplitude. The relations (9.4) are independent of all kinematic variables and hold in any region of the Dalitz plot.

Other high mass resonances such as the newly reported H meson[29)] also offer the possibility of spectroscopy with many open channels and fewer kinematic ambiguities. The H meson should be observable in the $K^0\bar{K}^0$ mode as $K_sK_s$ and give a very clean signal with little background. If the H has the quantum numbers of the f meson except for spin, as expected if it is the Regge recurrence, it should also have the same mixing angle. For this case the $K_sK_s$ width is related to the observed $2\pi$ width by SU(3)

$$\Gamma(H \to K_sK_s) = \frac{1}{4}\Gamma\,(H \to \pi^0\pi^0)\ . \tag{9.5}$$

The Regge recurrences of the A2 and the f' should also be expected and observable in $K\bar{K}$ modes. The analog of f-A2 interference should also be seen in the Regge recurrences of these states.

## 10. THE f-A2 INTERFERENCE

An interesting new tool which can be useful in hadron spectroscopy is interference between resonances which have a common decay mode. The $\rho$ and $\omega$ are examples of states which are nearly degenerate in mass and for which interesting interference effects have already been observed. These particles have no common strong decay mode and interference is observable only through the small electromagnetic contribution to the $\omega \to 2\pi$ decay. Both the $\rho$ and the $\omega$ are coupled to the $K\bar{K}$ system and would show strong interference effects in the $K\bar{K}$ decay mode if the channel were open. In an SU(3) symmetric world with equal kaon and pion masses these $K\bar{K}$ decays would have the same strength as the $\rho \to 2\pi$ decay.[29)]

For higher resonances such as the tensor mesons a similar nonet structure is observed with nearly degenerate isovector and isoscalar nonstrange bosons. These are now above the $K\bar{K}$ threshold and interference effects can therefore be seen. For higher nonets well above the $K\bar{K}$ threshold the $K\bar{K}$ decay mode should approach the two pion decay mode in strength and allow strong interference effects to be observed.

Because the f and A2 have different isospin the relative phase of the charged and neutral kaon pair decay modes is different in the two cases. Thus if the f and A2 are produced coherently the interference contributions observed in the kaon pair decay channel have opposite signs in the $K^+K^-$ and $K^0\bar{K}^0$ decay modes. This charge asymmetry can be very striking in some cases.

Interference effects have been reported and used in the analysis of the reaction[30)]

$$\pi^- + p \to M^0 + n \qquad (10.1a)$$

where $M^0$ denotes a neutral nonstrange meson which decays in the $K\bar{K}$ mode.

$$M^0 \to K + \bar{K} \, . \qquad (10.1b)$$

Both the f and the A2 can be produced in this reaction and interference effects can be observed. However some coherence is lost in averaging over polarization states and angular distributions if the two states are produced by different reaction mechanisms. With pion beams the f is produced primarily by pion exchange while the A2 cannot be produced by pion exchange because of G parity and must be produced by some other exchange such as $\rho$ exchange. The coherence is therefore reduced and is model dependent. Once the properties of the resonances are well established this model dependence can be used as a test of the models for the reaction mechanism.[31)] However if the parameters of the resonance are not well established the description of such a reaction becomes ambiguous and complicated.

Stronger interference effects with less model dependence are obtainable with kaon and photon beams which do not have a definite G parity and can produce both the f and the A2 by the same mechanism. This occurs in the reaction

$$K^- + n \to M_p^0 + Y \qquad (10.2a)$$

where N and Y denote any nucleon and hyperon states allowed by the conservation laws, and $M_p^0$ is the coherent linear combination of f and A2 which contains only a $p\bar{p}$ quark-antiquark pair and does not contain an $n\bar{n}$ component. This state decays only into charged kaons. The neutral kaon decay mode is forbidden because the f and A2 decay amplitudes exactly cancel one another in the neutral mode.

$$M_p^0 \to K^+ + K^- \qquad (10.2b)$$

$$M_p^0 \not\to K^0 + \bar{K}^0 \, . \qquad (10.2c)$$

This result is simply seen in quark diagrams but follows from the assumptions of SU(3) symmetry and ideal mixing without requiring any specific quark model assumptions.[32)] Effects of SU(3) symmetry breaking can be included by inserting physical masses and widths for the f and A2 instead of assuming the degeneracy of the symmetry limit. Figures 10.1a and 10.1b show predicted cross sections obtained by simply adding Breit-Wigner resonance curves with relative magnitude and phase predicted by SU(3) but with the mass and width given by the experimental data. The contribution from the f' is also included since it is also determined relative to the f and A2 by SU(3) except for a phase which depends on the signature of the exchange. Figure 10.1a has the f' phase corresponding to odd signature exchange such as vector exchange while Fig. 10.1b has the phase corresponding to even signature exchange such as scalar and tensor. The suppression of the $K^0\bar{K}^0$ mode under the f-A2 peak is very marked. There is also a surprising effect of the interference between the overlapping tails of the f and A2 and f' resonances. Exact quantitative features of those curves are not reliable in this region since the simple Breit-Wigner description may not hold in the tail. However the difference between the charged and neutral decays near the peak of the f' may be observable. Figures 10.1 also show the curve which would be obtained if there were no A2 contribution and no isovector-isoscalar interference. The interference effect in the tails of the f and f' are still noticeable.

Similar effects are observable in the production of neutral mesons by the $\gamma\gamma$ reaction[33)]

$$\gamma + \gamma \to M^0 \to K + \overline{K} . \tag{10.3a}$$

This can be studied with lepton colliding beams

$$e^{\pm} + e^- \to e^{\pm} + e^- + M^0 . \tag{10.3b}$$

Again SU(3) determines the relative amplitude and phase of the f, A2 and f' and the predicted cross section is shown in Fig. 10.2. The $K^0\overline{K^0}$ is strongly suppressed in the f-A2 region but not as strongly as in the case of the $K^-$ reactions (10.2). The predicted shift of the $K^0\overline{K^0}$ peak relative to $K^+K^-$ might be observable experimentally. A large $K^+K^-/K^0\overline{K^0}$ ratio is predicted in the f' region and might also be observable. Such interference effects may give new insight into the nature of these particles. Once the technique is established they may serve as a powerful tool for unscrambling reaction mechanisms.

## 11. TESTS OF ZWEIG'S RULE BY $\rho - \omega$ AND f-A2-f' INTERFERENCE

Interference effects which measure small amplitudes directly rather than their squares can provide sensitive tests of Zweig's Rule. The forbidden f' production amplitude in the $\pi N$ reaction (10.1) has been detected by f-f' interference.[34)] Similar interference between the tail of an allowed f amplitude and the peak of the forbidden f' amplitude should appear in the $K^-N$ reaction (10.2a). Since backward f' production is forbidden, f-f' interference in the allowed $K\overline{K}$ decay mode should be similar to that observed in the $\pi N$ reaction (10.1), but tests the validity of Zweig's rule at a baryon vertex rather than at a meson vertex. Forward f' production is allowed, and Zweig's rule can be tested in the forbidden $\pi\pi$ decay mode by f-f' interference.

The f-A2 interference discussed in the reaction (10.2) and the analogous $\rho$-$\omega$ interference also really test Zweig's rule. This is most easily seen from an SU(3) rotation of the familiar selection rule forbidding $\phi$ and f' production without accompanying strange particles. Interchanging the roles of n and $\lambda$ quarks we see that Zweig's rule also forbids production of the state $M_n^0$ which consists of an $n\overline{n}$ quark-antiquark pair without accompanying particles containing n and $\overline{n}$ quarks. It thus forbids the production of states containing only $M_n^0$ and charged kaons. Neutral nonstrange mesons produced with charged kaons must be in the $M_p^0$ state which is an equal and fully coherent mixture of $\rho$ and $\omega$ or of f and A2. Maximum interference should be observed in decays. One example is seen in the relations (9.4) between $\psi$ decays. If Zweig's rule (9.3) holds the amplitudes (9.4b) and (9.4c) become equal and exhibit maximum interference in the $K^+K^-\pi^+\pi^-$ decay mode. A similar effect should occur in the decay of <u>any</u> SU(3) singlet meson in this mode and also in $\overline{p}p$ and $e^+e^-$ annihilation in this mode.

Different experiments which produce the $\rho$ and $\omega$ only via the $M_p^0$ state should all exhibit the same interference and the same $\pi^+\pi^-$ effective mass distributions. Thus results from several experiments can be combined to get better statistics.

## 12. CONCLUSION

We still have much to learn about the old particles!

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