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SYMMETRIES AND RESONANCES (THEORETICAL)

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INTRODUCTION

The standard review of resonances, symmetries, or the quark model which was a regular feature of international conferences over the past several years is now out of date. Many excellent reviews already available in the literature summarize the predictions of symmetry schemes and the quark model and their comparison with experiment.[1] Instead of rehashing this old material, this discussion summarizes the state of the art in a few words and concentrates on new applications and on questions of current interest.

Internal symmetries are no longer fashionable in theoretical physics and theorists tend to forget them in jumping on such new bandwagons as the parton model and inclusive reactions. However, symmetries are relevant to these new areas as well, and paradoxes can arise if they are not taken into account properly. Some examples will be presented. Sections I and II consider the parton model and inclusive reactions, respectively, Sec. III reviews the quark model, and Secs. IV and V discuss spectroscopy and the new exotic particles being looked for at NAL.

I. SYMMETRIES, QUARKS, AND PARTONS

The parton model[2] has been extensively used in the analysis of deep-inelastic-scattering experiments. Quarks are candidates for partons, and analysis of these experiments has been proposed to determine whether partons have integral or fractional electric charge.[3,4] However, difficulties arise in application of the parton model to internal symmetries.[5] The discussion will begin with a simple "pedestrian" example of how the quark-parton model is used to make experimental predictions. We then discuss internal symmetries and difficulties.

The difficulties discussed here are intimately related to the "asymptopia paradox." Does scaling occur because asymptopia has

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already been reached and no new physics can occur at higher energies? If quarks are produced at higher energies, this is new physics and the threshold energy provides a scale. Then is the scaling observed at present energies purely accidental? Many physicists worry about this question, but there are no convincing answers. The present discussion adds a few more worries and gives no answers.

A. Inequalities Based on the Quark-Parton Model

The suggestion that quarks are partons has been used to derive inequalities between various inelastic form factors. For example, Nachtmann[3] has shown that the ratio of the deep-inelastic electromagnetic form factors of the proton and neutron is limited by the values 4 and 1/4 if only isospin symmetry is assumed and the values 4 and 1/3 if SU(3) is assumed. A simple example shows how such inequalities result from the assumption that hadrons are made up of quarks and antiquarks. The quarks will be denoted by p, n, and λ. The electric charges are $Q_p = +2/3$, $Q_n = Q_\lambda = -1/3$.

Nachtmann's inequalities are easily understood by constructing explicit quark models to give the limiting cases. The neutron and proton are members of the same isospin doublet. The limits in the quark model are obtained by finding the isospin doublets that have the maximum and minimum ratios of the square of the charge. These are just the quark doublet and the antiquark doublet, respectively, that give the ratios 4 and 1/4. That is,

$$\frac{Q_q^2(I_z=+\frac{1}{2})}{Q_q^2(I_z=-\frac{1}{2})} = \frac{Q_p^2}{Q_n^2} = 4, \qquad (1a)$$

$$\frac{Q_{\bar{q}}^2(I_z=+\frac{1}{2})}{Q_{\bar{q}}^2(I_z=-\frac{1}{2})} = \frac{Q_{\bar{n}}^2}{Q_{\bar{p}}^2} = \frac{1}{4}. \qquad (1b)$$

No isodoublet can be constructed from quarks and antiquarks whose value of $Q^2(I_z=+\frac{1}{2})/Q^2(I_z=-\frac{1}{2})$ is outside these limits.

Quark models for the nucleon can be constructed to give either the ratio (1a) or the ratio (1b) for the deep-inelastic form factor. The ratio (1a) is given by a model in which a single "active" quark absorbs the photon and the nucleon also contains two spectator quarks coupled to isospin zero to give the correct isospin and baryon number, i.e., ratio (1a) results when

$$|N_a\rangle = q_a(q_s q_s)_{I=0}, \quad (2a)$$

where the subscripts a and s on the quark labels denote active and spectator. The ratio (1b) is obtained from a model in which an active antiquark absorbs the photon and the nucleon contains four spectator quarks coupled to isospin zero so that

$$|N_b\rangle = \bar{q}_a(q_s q_s q_s q_s)_{I=0}. \quad (2b)$$

The wave functions (2a) and (2b) have the maximum and minimum values (1a) and (1b) for ratios of squared charge that can be obtained from any isodoublet made from quarks and antiquarks. Thus, Nachtmann's inequality at the SU(2) level is clear.

When SU(3) is considered, the wave functions (2) must also satisfy requirements of SU(3) symmetry; i.e., they must transform under SU(3) like a member of an octet. The wave function (2a) is pure octet, since the two-quark state with zero isospin is classified in the $\bar{3}$ representation of SU(3) while the active quark is in the 3 representation. The only isodoublet contained in the product $\bar{3} \times 3$ is in an SU(3) octet. The wave function (2b), however, is not an SU(3) octet state. The four-quark state with I = 0, Y = 4/3 is in the $\bar{6}$ representation of SU(3) while the antiquark is in the $\bar{3}$ representation. The product $\bar{6} \times \bar{3}$ contains two isodoublet states, one in an octet and one in the $\bar{10}$ representation. The wave function (2b) contains components belonging to both representations. The octet state which is easily constructed with the aid of U-spin Clebsch-Gordan coefficients, is

$$|N_8\rangle = \sqrt{\frac{2}{3}}\,\bar{q}_a(q_s q_s q_s q_s)_{\{I=0,\ Y=\frac{4}{3}\}} + \sqrt{\frac{1}{3}}\,\bar{\lambda}_a(q_s q_s q_s \lambda_s)_{\{I=1/2,\ Y=-\frac{1}{3}\}}, \quad (3a)$$

where q denotes a nonstrange quark. The ratio of the squared charge values for the active antiquark in the octet state (3a) is then

$$\frac{Q_8^{\,2}(I_z=+\frac{1}{2})}{Q_8^{\,2}(I_z=-\frac{1}{2})} = \frac{\frac{2}{3}Q_{\bar{n}}^{\,2} + \frac{1}{3}Q_{\bar{\lambda}}^{\,2}}{\frac{2}{3}Q_{\bar{p}}^{\,2} + \frac{1}{3}Q_{\bar{\lambda}}^{\,2}} = \frac{1}{3}. \quad (3b)$$

This is the Nachtmann SU(3) limit. Note that it depends on the presence of strange quark-antiquark pairs required by SU(3) in the wave function.

B. G Parity and Isospin in the Parton Model

Is the parton model compatible with conservation laws from such internal symmetries as G parity and isospin? To consider this question we first examine a process in which the initial state is an eigenstate of G—e.g., the deep inelastic scattering on a pion by the isovector component of the vector current. This initial state has odd G parity. Therefore G conservation allows only multipion final states having an odd number of pions, as shown in Fig. 1. Final states having an even number of pions are forbidden.

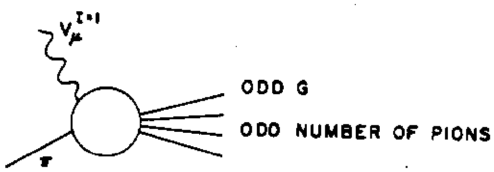

Fig. 1. Deep inelastic scattering from G eigenstate.

In the quark-parton model, the parton is not an eigenstate of G. If a single parton absorbs the incident momentum and the final multipion state is produced by interactions between the struck parton and other partons, both even and odd numbers of pions can be produced. Thus a naive application of the parton model would seem to be incompatible with G conservation.

Incorporating G parity in the parton model is straightforward elementary quantum mechanics. The initial and final states in the process are eigenstates of the conserved G parity, but the transition is described in a model with intermediate states that are not eigenstates of G. Figure 2 indicates schematically how G conservation is included in the model. For each parton-model intermediate state that is not an eigenstate of G parity, there exists a G-conjugate intermediate state produced by operating on the first intermediate state with the G operator, i.e.,

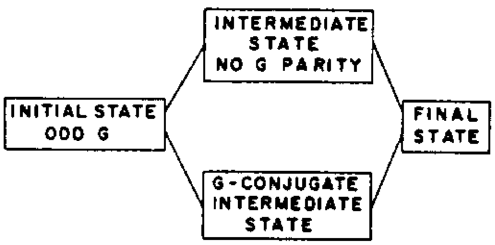

Fig. 2. G conservation with intermediate states that are not G eigenstates.

$$\psi_{int}^{conj} = G\psi_{int}^{part} . \quad (4)$$

This conjugate state has the same space-time structure as ψ_{int}^{part} but each parton is replaced by its G-conjugate antiparton. This interference between the contributions from such pairs of G-conjugate intermediate states is crucial for incorporating G conservation in the parton model. The amplitudes from both intermediate states must be added coherently to give the transition amplitude for the final state. The two contributions give amplitudes having the same magnitude for each final state; however, the two amplitudes add constructively for the allowed final states having odd numbers of pions, while there is destructive interference and cancellation for the forbidden even numbers of pions.

Suppose the interference between these two amplitudes is neglected as in the most naive application of the parton model. Nonvanishing cross sections would be obtained for the forbidden even-G states, while the cross sections for the allowed odd-G states would be in error by a factor of two because of the neglect of the constructive interference term. At very high energies at which large numbers of pions are produced, the probability for producing n pions should be a smooth function of n as indicated in Fig. 3, which shows both the true cross section σ_n for the production of n pions and the cross section calculated with the neglect of interference terms. In this case the neglect of the interference term does not introduce a serious error for average properties or gross structure of the cross section—e.g., for calculation of the total cross section, the average multiplicity, or the cross section within a bin containing many values of n. The error is observed only in the detailed fine structure, i.e., in the difference between the cross sections for producing 2n pions and 2n + 1 pions.

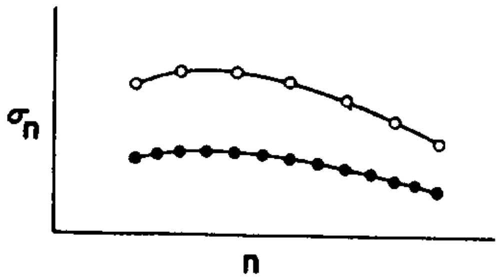

Fig. 3. Cross section for production of n-pion final state. Open circles denote true cross section. Solid circles denote cross sections that would be calculated with neglect of interference.

Consider now the description of the total cross section as the imaginary part of a forward-scattering amplitude as indicated in Fig. 4. The parton model assumes an impulse approximation in which the struck parton propagates freely in the intermediate state and then emits the final photon or lepton pair. In this simple picture the intermediate state is not an eigenstate of G. The total cross section calculated in this way apparently includes exclusive cross sections to forbidden states, as if the interference between G-conjugate pairs of intermediate states were neglected. If the fine structure of the total cross section is described by the smooth variation of Fig. 3, only a negligible error in the total cross section would result from neglecting interference terms.

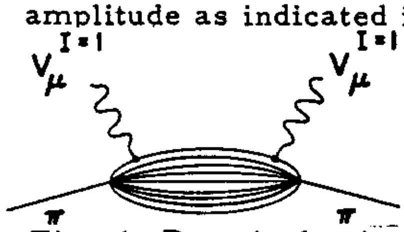

Fig. 4. Deep inelastic scattering from a pion as a forward amplitude.

If there is a systematic difference between the calculated even and odd cross sections so that the averaging procedure of Fig. 3 does not give the proper total cross section, it becomes necessary to take into account the interference terms indicated in Fig. 2. This could be done by suitably symmetrizing the intermediate states shown in Fig. 4—i.e., by choosing the G-parity eigenfunctions as a basis for intermediate states

$$\psi_\pm = \tfrac{1}{2}[1 \pm G]\psi_{\text{int}}^{\text{part}}. \tag{5}$$

If this basis is used for intermediate states, the total cross section includes contributions in which the incident lepton pair or photon is absorbed by one parton and the subsequent emission comes from its G-conjugate parton in the G-conjugate wave function. Whether this symmetrization violates the spirit of the impulse approximation is a matter of taste. However, this is a perfectly consistent procedure and may be relevant to cases in which one wishes to define a modified impulse approximation which has "manifest internal symmetry conservation" (analogous to manifest covariance) at all stages in the calculation. This point is discussed further in Sec. I C.

A similar problem arises with isospin. Consider the deep inelastic scattering of an isoscalar current on a simple deuteron model which consists of a neutron and a proton. Since the deuteron has isospin zero, only states of isospin zero can be produced. However, a final state in which the proton has absorbed the momentum and the neutron has not is not an eigenstate of isospin. Isospin conservation is preserved by adding the contributions indicated in Figs. 5a and 5b coherently to give transitions only to properly symmetrized final states.

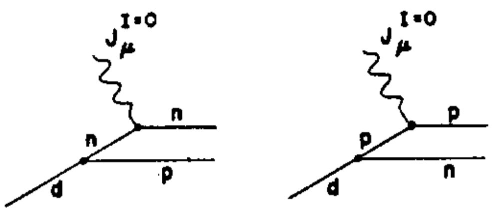

Fig. 5. Deep inelastic scattering of isoscalar current from deuteron.

In the description of the total cross section as a forward amplitude, as indicated in Fig. 6, manifest isospin invariance is consistent with the impulse approximation only if symmetrized wave functions are used for the intermediate states. In this symmetrized description, the photon is absorbed by a nucleon and emitted by a nucleon, but the active nucleon can be either a neutron or a proton and can change from one to another between the absorption and the emission. This charge change is necessary to preserve manifest isospin invariance. Whether it is consistent with the spirit of the impulse approximation is not clear and is discussed below in Sec. I C.

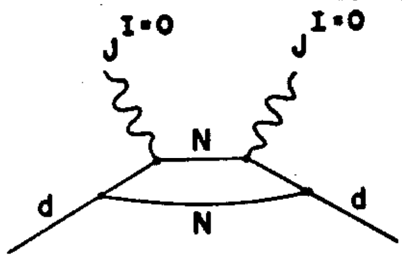

Fig. 6. Forward scattering of isoscalar current by deuteron.

We thus conclude that there are two correct approaches to internal symmetries in the parton model: either forget about them, or treat them properly. But the two approaches should not be mixed by forgetting symmetries when doing the calculation and remembering them when interpreting the answer. In the G-parity example, forgetting symmetry is all right for total cross sections and gross features of the pion multiplicity, but fine structure effects such as the even-odd pion difference are lost. If such precise fine structure associated with internal symmetries is desired, then the internal symmetry must be treated properly. This involves coherent addition or symmetrization of amplitudes before they are squared to get cross sections.

Note that the electric charge of a parton is an internal symmetry quantum number and therefore the determination of the parton charge from analyses that do not properly take account of internal symmetries can be risky.[5]

C. Hidden Internal Symmetries

Consider deep inelastic electron scattering on a deuteron in a world where all nuclei are made of deuterons and all low-lying nuclear states have isospin zero. Experiments below the threshold for producing states of nonzero isospin would not detect the presence of isospin, which would be a "hidden internal symmetry." Only the isoscalar component of the electromagnetic current could contribute to transitions between states of zero isospin. The isovector component would be undetectable in this energy range. Since an isoscalar probe cannot distinguish between a proton and a neutron, low-energy experiments could not distinguish between the two components of the deuteron. Inelastic scattering would show that a deuteron with electric charge +1 was composed of two identical spin-$\frac{1}{2}$ objects which might be called nucleons, and that only the <u>symmetric</u> states of the two-nucleon system were observed. <u>The nucleon would thus appear to be a spin-$\frac{1}{2}$ particle with the fractional electric charge of 1/2 and peculiar statistics.</u>

Only after the excitation of the first I = 1 states of the two-nucleon system would it become clear that the fractional charge and peculiar statistics simply masked the existence of an additional degree of freedom and of two kinds of nucleons, each having integral charge. These are described as states of a single particle—the nucleon—by the formal introduction of a new internal degree of freedom and a new internal SU(2) symmetry, the isospin. The peculiar statistics reduce to ordinary Fermi statistics when the new internal degree of freedom is included, since states that are antisymmetric in isospin are required to be symmetric in the other degrees of freedom.

Two important characteristics of this model underly its peculiar properties: (1) the <u>binding energy</u> of the two nucleons in the deuteron and (2) the <u>symmetry energy</u> which separates the I = 0 and I = 1 states of the two nucleon system. These two energies are normally independent of one another. The peculiar properties of this model result from very high values of <u>both</u> the binding energy <u>and</u> the symmetry energy, for which condition there is a large range of experimental excitation energies below the thresholds of both deuteron breakup and excitation of I = 1 states.

This deuteron model reminds one of the quark model[6] for baryons, which has fractional charges and peculiar statistics. The Han-Nambu three-triplet model[7] can give the same results with integral charges and ordinary Fermi statistics by introducing a new hidden degree of freedom and a hidden symmetry, in this case an SU(3) symmetry rather than SU(2). Again the low-lying states are all required to be singlets in the new internal symmetry; i.e., the symmetry energy must be high so that only singlet states are observed. As long as only singlet states in the internal symmetry are observed, the different members of a given internal symmetry multiplet appear to be identical and have a fractional charge equal to the average charge of the members of the multiplet, just as in the deuteron model. States that are not singlets in the hidden SU(3) degree of freedom are sometimes called "charmed" states. Since these have not been observed, the symmetry energy is assumed to be above present accelerator energies to avoid disagreement with experiment.

Thus we cannot distinguish between the Han-Nambu model and the quark model with fractional charges and peculiar statistics until the energy exceeds the binding energy so one can see the constituents directly or exceeds the symmetry energy so charmed states are excited. Even if the Gell-Mann-Zweig quark model is correct, the Han-Nambu model cannot be disproved until quarks or other fractionally-charged objects are actually observed. Conversely, if the Han-Nambu model is correct, the Gell-Mann-Zweig quark model cannot be disproved until either the symmetry energy or the binding energy is reached and the charmed states or the integrally-charged triplets are observed.[8]

We now return to the validity of the impulse approximation in the calculation of the total cross section by the diagram of Fig. 6. If the symmetry energy is zero or negligibly small, the I = 0 and I = 1 components of the intermediate states of Fig. 6 are degenerate. Then no error is introduced by using the states of the naive impulse approximation, in which the active nucleon remains a neutron or remains a proton and does not change its charge in the intermediate state. If the symmetry energy is high, however, only the I = 0 intermediate state should be used, and the naive impulse approximation introduces an error. But a high symmetry energy defines, through the uncertainty principle, a <u>short time</u> which characterizes the exchange of charge between the two nucleons.[5] For photon energies below the symmetry energy, this "charge-exchange time" is short compared to the lifetime of the intermediate state, and one can say that the naive impulse approximation is not valid because the collision is not fast compared to characteristic times defined by the bound state. However, the exact meaning of a modified impulse approximation using symmetrized intermediate states of definite isospin is not clear. Simple models for two-nucleon systems with large symmetry energy require exchange forces that interchange proton and neutron. These are not represented clearly in the naive parton picture.

D. Quarks and Partons

Additional difficulties arise if quarks are assumed to be physical objects which have not yet been observed simply because their mass is too high; e.g., they might have a mass of 50 GeV and be waiting for NAL. If these very heavy quarks exist and are very strongly bound in hadrons, how can scaling occur below the threshold for quark production and be described by a model in which quarks move quasi-freely even though they are very strongly bound? The various arguments and theorems connected with this point are beyond the scope of this review. I simply include some provocative remarks about the commonly used expression

$$\langle p\,|\overline{\psi}(x)\psi(0)\,|p\rangle, \tag{6}$$

where p is a state of a proton at rest and $\overline{\psi}(x)\psi(0)$ is a bilocal operator which appears formally to be connected with quark fields.[9] However, if this operator creates physical quarks, one can prove that it cannot give any contribution to cross sections below the threshold for the production of fractionally-charged states.

The catastrophe results from the following two assumptions which are natural if $\psi(x)$ is a quark field that creates fractionally charged states. (1) The matrix element (6) can be evaluated by inserting a complete set of intermediate states between the two ψ's. That is,

$$\langle p\,|\overline{\psi}(x)\psi(0)\,|p\rangle = \sum_i \langle p\,|\overline{\psi}(x)\,|i\rangle\langle i\,|\psi(0)\,|p\rangle. \tag{7}$$

(2) The operators $\psi(x)$ and $\overline{\psi}(x)$ satisfy commutation relations with the electric charge operator Q, namely

$$[Q, \overline{\psi}(x)] = q\overline{\psi}(x), \tag{8a}$$

$$[Q, \psi(x)] = -q\psi(x), \tag{8b}$$

where q is a fraction, such as 1/3 or 2/3.

<u>Theorem</u>: If assumptions (1) and (2) above are valid, the time dependence of the matrix element (6) has no fourier components below the frequency corresponding to the threshold energy

for the production of fractionally charged particles.

Proof: The time dependence of the matrix element (6) is exhibited explicitly by writing

$$\langle p|\bar{\psi}(\vec{x}, t)\psi(0, 0)|p\rangle = \langle p|e^{iHt/\hbar}\psi(\vec{x}, 0)e^{-iHt/\hbar}\psi(0, 0)|p\rangle \tag{9a}$$

Using Eq. (7), we obtain the Fourier transform of the time dependence.

$$\langle p|\bar{\psi}(\vec{x}, t)\psi(0, 0)|p\rangle = \sum_i \langle p|\bar{\psi}(\vec{x}, 0)|i\rangle\langle i|\psi(0, 0)|p\rangle e^{-i(E_i - E_p)t/\hbar} . \tag{9b}$$

From the commutation relation (8b), it follows that

$$\langle i|[Q, \psi(x)]|p\rangle = [Q_i - Q_p]\langle i|\psi(x)|p\rangle = -q\langle i|\psi(x)|p\rangle. \tag{10}$$

Thus the matrix element (10) differs from zero only for intermediate states i having the electric charge

$$Q_i = Q_p - q. \tag{11}$$

Since q is fractional, only intermediate states of fractional charge can contribute to the sum on the right-hand side of (9b). The smallest Fourier component in this sum then comes from the least energetic fractional-charge state that satisfies Eq. (11) and has an energy above the threshold for producing fractionally charged particles.

This theorem is clearly unaffected by interposing γ matrices between the two spinors in expression (6) or including SU(3) indices and λ matrices.

We therefore conclude that if the Fourier transform of expressions of the form (6) appear in theoretical relations for total cross sections for inclusive processes, and if assumptions (1) and (2) are valid, these expressions cannot contribute to the cross section below the threshold for production of fractionally charged particles. Assumptions (1) and (2) are very reasonable if the spinor field $\psi(x)$ is interpreted as being that of a physical particle, such as a quark, which has a third-integral charge. In this case, all derivations that take bilocal operators of the form (6) seriously are unable to explain scaling below the threshold for production of fractionally charged objects.

This difficulty is avoided by saying that the relevant physical quantity is a bilocal operator J(x, 0) and that there is no physical meaning in the factorized expression[10]

$$J(x, 0) = \bar{\psi}(x)\psi(0). \tag{12}$$

The spinor fields $\psi(x)$ are not interpreted as physical operators but merely as mathematical crutches that aid us in defining commutators. In that case, Eq. (7) does not hold and the theorem is not valid. So far there is nothing obviously wrong with this point of view.

However, if real quarks exist and will be found some day, our theorem shows that expressions of the form (6) have nothing to do with the scaling that is presently observed in electroproduction experiments below the quark-production threshold.

II. INTERNAL SYMMETRIES AND INCLUSIVE REACTIONS

The big band wagon of 1971 was inclusive reactions and multiparticle processes.[11] Here are two areas in which internal symmetries are relevant.

A. Isospin and U-spin Relations in Inclusive Reactions

Consider the inclusive reaction

$$A + B \to C_{IM} + X, \tag{13}$$

where C_{IM} denotes a set of states within the same isospin multiplet with isospin I, and M is the eigenvalue of I_z. Peshkin[12] has shown that isospin invariance constrains the variation of the cross section σ_{IM} for the inclusive process (13) as a function of the index M. Peshkin's theorem is the analog in isospin space of the statement in angular momentum space that an initial state containing only s and p waves cannot produce final-state angular distributions containing higher powers of $\cos\theta$ than $\cos^2\theta$.

The theorem states that if $I_{AB}^{(max)}$ is the maximum isospin present in the initial state, the inclusive cross section σ_{IM} can be written as a polynomial in M of degree $2I_{AB}^{(max)}$. That is,

$$\sigma_{IM} = \sum_{n=0}^{2I_{AB}^{(max)}} a_{In} M^n, \tag{14}$$

where the coefficients a_{In} are undetermined parameters. Since the number of independent cross sections σ_{IM} is $2I + 1$, and the number of free parameters a_{In} in the polynomial (14) is $2I_{AB}^{(max)} + 1$, the condition for obtaining nontrivial relations between cross sections from Peshkin's theorem is

$$I_C > I_{AB}^{(max)}. \tag{15}$$

For initial states having maximum isospin $\frac{1}{2}$, nontrivial relations can be obtained for inclusive production of isovector particles. For initial states having maximum isospin 1, relations are obtainable for inclusive production of particles having $I \geq \frac{3}{2}$. With available beams $I_{AB} = \frac{1}{2}$ is obtained only by the use of isoscalar targets—i.e., deuterons. With nucleon targets, only relations for production of particles having $I \geq \frac{3}{2}$ are obtained.

Since isospin conservation is accepted in strong interactions, these relations test consistency of experimental data analysis; e.g. the treatment of deuteron data or of resonances and background. They can also test diffractive excitation models which assume exchange only of isospin zero between beam and target. In these models Eq. (14) holds separately in the fragmentation regions of A and B with $I_{AB}^{(max)}$ replaced by I_A and I_B respectively. Relations for photoproduction processes also test the assumption that the photon has only isoscalar and isovector components.

Relations obtained by a straightforward application of Peshkin's theorem (14) are

$$\sigma(K^{\pm}d \to \pi^{+}X) + \sigma(K^{\pm}d \to \pi^{-}X) = 2\sigma(K^{\pm}d \to \pi^{0}X), \tag{16a}$$

$$\sigma(K^{\pm}d \to \Sigma^{+}X) + \sigma(K^{\pm}d \to \Sigma^{-}X) = 2\sigma(K^{\pm}d \to \Sigma^{0}X), \tag{16b}$$

$$\sigma(K^{\pm}d \to \Delta_{M}X) = a_0 + a_1 M, \tag{16c}$$

where a_0 and a_1 are unknown parameters, and

$$\sigma(K^{\pm}N \to \Delta^{+}X) + \frac{1}{3}\sigma(K^{\pm}N \to \Delta^{-}X) = \sigma(K^{\pm}N \to \Delta^{0}X) + \frac{1}{3}\sigma(K^{\pm}N \to \Delta^{++}X) \tag{16d}$$

Relations identical to (16a—d) are obtained for analogous reactions with beams of protons or antiprotons instead of kaons and also for fragmentation of kaons or nucleons in diffractive excitation models[13]. A relation for photoproduction or electroproduction of Δ's is

$$\sigma(\gamma d \to \Delta^{+}X) + \frac{1}{3}\sigma(\gamma d \to \Delta^{-}X) = \sigma(\gamma d \to \Delta^{0}X) + \frac{1}{3}\sigma(\gamma d \to \Delta^{++}X), \tag{16e}$$

where the photon can also be virtual. All photoproduction relations in this paper also apply to electroproduction.

The identical argument can be applied to U spin or V spin. We replace C_{IM} by C_{UM} or C_{VM}, which denote a U spin or V spin multiplet. Here M is the eigenvalue of U_z or V_z. Predictions then follow from U-spin or V-spin invariance—i.e., from SU(3) symmetry. The degree to which they are satisfied experimentally may give some insight into SU(3) symmetry breaking.

Because the photon is a U-spin scalar, the U-spin analog of the condition (15) for nontrivial relations is most easily satisfied in photoproduction on protons, for which the initial state has $U = \frac{1}{2}$. However, relations involving octet particles in the final state are not easily obtainable because the Λ, Σ^0, π^0, and η are not eigenstates of U spin and V spin. Thus, relations involving decuplet baryons are the ones most easily obtained.[14] These are

$$\sigma(\gamma p \to \Delta^0 X) + \sigma(\gamma p \to \Xi^{*0} X) = 2\sigma(\gamma p \to Y^{*0} X), \tag{17a}$$

$$\sigma(\gamma p \to \{\Delta^-, Y^{*-}, \Xi^{*-} \Omega^-\}X) = a_0 + a_1 M, \tag{17b}$$

$$\sigma(\gamma n \to Y^{*-} X) + \frac{1}{3}\sigma(\gamma n \to \Omega^- X) = \sigma(\gamma n \to \Xi^{*-} X) + \frac{1}{3}(\gamma n \to \Delta^- X), \tag{17c}$$

$$\sigma(K^\pm p \to Y^{*-} X) + \frac{1}{3}\sigma(K^\pm p \to \Omega^- X) = \sigma(K^\pm p \to \Xi^{*-} X) + \frac{1}{3}\sigma(K^\pm p \to \Delta^- X), \tag{17d}$$

$$\sigma(\pi^\pm p \to Y^{*-} X) + \frac{1}{3}\sigma(\pi^\pm p \to \Omega^- X) = \sigma(\pi^\pm p \to \Xi^{*-} X) + \frac{1}{3}\sigma(\pi^\pm p \to \Delta^- X). \tag{17e}$$

The same approach can be used for multiparticle inclusive processes. Equalities are not easily obtained because multiparticle states are not generally isospin eigenstates. However, useful inequalities can be derived by noting that all observable cross sections are positive. Let σ_M denote the sum of the cross sections for all states of the multiparticle system having a given value of M. Then

$$\sigma_M \equiv \sum_I \sigma_{IM} \geq \text{any one } \sigma_{IM}, \tag{18}$$

and combining Eqs. (14) and (18) leads to inequalities setting lower bounds on a given σ_M. For example, in pion-nucleon inclusive production from initial states with $I_{AB}^{(max)} = \frac{1}{2}$, Eqs. (14) and (18) give[14]

$$\sigma(K^\pm d \to \pi^+ nX) + \sigma(K^\pm d \to \pi^0 pX) \geq \frac{2}{3}\sigma(K^\pm d \to \pi^+ pX) + \frac{1}{3}\sigma(K^\pm d \to \pi^- nX), \tag{19a}$$

$$\sigma(K^\pm d \to \pi^0 nX) + \sigma(K^\pm d \to \pi^- pX) \geq \frac{2}{3}\sigma(K^\pm d \to \pi^- nX) + \frac{1}{3}\sigma(K^\pm d \to \pi^+ pX). \tag{19b}$$

The analogs of these equations apply also to nucleon and kaon fragmentation in diffractive excitation models with zero isospin exchange.

For initial states with $I_{AB}^{(max)} = 1$, e.g., in such reactions as

$$K^{\pm} + p \to \pi + N + X, \tag{20a}$$

$$p + p \to \pi + N + X, \tag{20b}$$

$$\bar{p} + p \to \pi + N + X, \tag{20c}$$

$$\gamma + d \to \pi + N + X, \tag{20d}$$

the inequalities obtained are

$$\sigma(\pi^{+} nX) + \sigma(\pi^{0} pX) \geq \frac{1}{3}\left[\sigma(\pi^{+} pX) - \sigma(\pi^{-} nX)\right], \tag{21a}$$

$$\sigma(\pi^{0} nX) + \sigma(\pi^{-} pX) \geq \frac{1}{3}\left[\sigma(\pi^{-} nX) - \sigma(\pi^{+} pX)\right]. \tag{21b}$$

Similarly, different charge states for any multiparticle inclusive process from an initial state with $I_{AB}^{(max)} = 1$ are related by the inequality

$$\sigma_M(I_{AB}^{(max)} = 1) \geq \left(\frac{M + m}{2\mathcal{M}}\right)\left[\sigma_{\mathcal{M}}\left(\frac{\mathcal{M} + M}{\mathcal{M} + m}\right) - \sigma_{-\mathcal{M}}\left(\frac{\mathcal{M} - M}{\mathcal{M} - m}\right)\right], \tag{22}$$

where $\mathcal{M}$ is a particular value of M, m is any allowed value of M, and $|M| < |\mathcal{M}|$ and $|m| < |\mathcal{M}|$. The inequality (22) can be applied to multipion inclusive processes, such as the two-pion reactions

$$K^{\pm} + p \to \pi(k_1) + \pi(k_2) + X, \tag{23a}$$

$$p + p \to \pi(k_1) + \pi(k_2) + X, \tag{23b}$$

$$\bar{p} + p \to \pi(k_1) + \pi(k_2) + X, \tag{23c}$$

$$\gamma + d \to \pi(k_1) + \pi(k_2) + X, \tag{23d}$$

where the momenta k_1 and k_2 distinguish between the two pions. From these one obtains[14] the inequalities

$$\sigma_1(I_{AB}^{(max)} = 1) \geq \tfrac{1}{2}[\sigma_2 - \sigma_{-2}], \qquad (24a)$$

$$\sigma_1(I_{AB}^{(max)} = 1) \geq \tfrac{1}{4}[\tfrac{3}{2}\sigma_2 - \tfrac{1}{2}\sigma_{-2}], \qquad (24b)$$

$$\sigma_0(I_{AB}^{(max)} = 1) \geq \tfrac{1}{4}[\frac{2}{3}\sigma_2 - 2\sigma_{-2}], \qquad (24c)$$

$$\sigma_{-1}(I_{AB}^{(max)} = 1) \geq \tfrac{1}{2}[\sigma_{-2} - \sigma_2], \qquad (25a)$$

$$\sigma_{-1}(I_{AB}^{(max)} = 1) \geq \tfrac{1}{4}[\tfrac{3}{2}\sigma_{-2} - \tfrac{1}{2}\sigma_2], \qquad (25b)$$

$$\sigma_0(I_{AB}^{(max)} = 1) \geq \tfrac{1}{4}[\frac{2}{3}\sigma_{-2} - 2\sigma_2], \qquad (25c)$$

where for the case of the two-pion reactions (24),

$$\sigma_{\pm 2} \equiv \sigma(AB \to \pi^{\pm}\pi^{\pm}X), \qquad (26a)$$

$$\sigma_{\pm 1} \equiv \sigma(AB \to \pi^{\pm}\pi^{0}X) + \sigma(AB \to \pi^{0}\pi^{\pm}X), \qquad (26b)$$

$$\sigma_0 \equiv \sigma(AB \to \pi^{+}\pi^{-}X) + \sigma(AB \to \pi^{-}\pi^{+}X) + \sigma(AB \to \pi^{0}\pi^{0}X). \qquad (26c)$$

The inequalities (22), (24) and (25) also hold for the multipion inclusive processes

$$A + B \to n\pi + X; \qquad I_{AB}^{(max)} = 1, \qquad (27)$$

where σ_M includes all the n-pion states of charge M. These inequalities hold for any set of fixed values of the momenta of the outgoing particles. Thus in a given experiment they can be tested at each point in the energy spectrum and angular distribution.

Note that the inequalities (20d) and (23d) could test for the presence of an isotensor component in the electromagnetic current. Equalities, such as Eq. (16e), that involve resonance production are useless for such tests because ambiguities in separating resonances from background are always greater than the effect tested. Extensive tests of isospin properties of currents by inequalities in exclusive reactions have been recently described by Pais.[15]

These model-independent relations follow from isospin and U-spin invariance, respectively, and will hold in any model, such as the Mueller-Regge model,[16] if the model does not violate isospin or SU(3) symmetry. Additional model-dependent symmetry relations have been obtained from particular models.[17] These usually follow from assumptions that limit the quantum numbers in a particular channel to those of allowed (nonexotic) Regge trajectories, or to be those of the Pomeron in the case of a diffractive process.

B. What is Exotic in Inclusive Reactions?

Internal symmetry also plays a role in the discussion of what is exotic in an inclusive reaction. Extending the Harari-Freund conjecture[18] for duality in two-body reactions to inclusive cross sections suggests that processes "having exotic quantum numbers should be constant or show limiting behavior; i.e., the contributions of secondary Regge trajectories should cancel one another. But which combination of the quantum numbers of a three-particle system should be used to define exotic? Various combinations have been suggested. Chan et al.[19] first suggested that an inclusive process should be considered exotic whenever the quantum numbers of the complex ($AB\bar{C}$) are exotic. This has been followed by the group at Stony Brook and can be called the Stony Brook party line. Suggestions that the Stony Brook party line has internal inconsistencies[20–22] have led to alternative criteria which could be called the CERN,[20] Berkeley,[23] and MIT[24] party lines. Without considering the merits of each individual criterion, I wish to point out the close relation between internal symmetries and the alleged inconsistencies in the original Chan criterion. A clear contradiction cannot be found without applying the criterion to hypothetical results of "gedanken experiments" involving η, η', or hyperon beams or meson or hyperon targets and relating these to observable experiments by SU(3) or factorization.[21] Such arguments would break down if the results of these experiments violate SU(3) or factorization in a violent and unexpected way. Although this does not seem very reasonable, it is still an open loophole.

This point has been illustrated in the example of the reactions[25]

$$K^{\pm} + N \to \pi^{\mp} + X, \tag{28a}$$

for which experimental data are available.[26] Consider the region in which the pion is a fragment of the kaon. In the standard Mueller formalism,[16] the inclusive cross section is given by the amplitude for the six point function

$$\{K^{\pm} + \pi^{\mp}\} + N \to \{K^{\pm} + \pi^{\mp}\} + N, \tag{28b}$$

where the braces indicate that in the fragmentation region of the kaon, Regge-pole exchanges are considered between the nucleon and the $K\pi$ complex, as shown in Fig. 7. The Regge trajectories

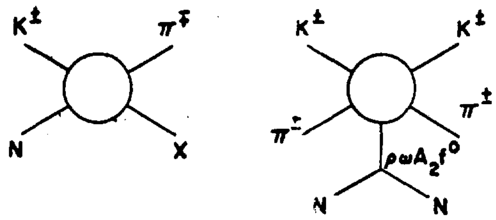

Fig. 7. The inclusive reaction $K^{\pm}N \to \pi^{\mp}X$.

considered in addition to the Pomeron are the usual isoscalar and isovector trajectories with even and odd signature, characterized by the quantum numbers of ρ, ω, A2, and f^0. In the four different charge states of reaction (28), the three-particle system in Eq. (28b) is exotic by Chan's criterion in three cases, namely,

$K^+\pi^+p$	exotic	(29a)
$K^+\pi^+n$	exotic	(29b)
$K^-\pi^-n$	exotic.	(29c)

The fourth case, which is not exotic, is

$K^-\pi^-p$	not exotic.	(29d)

Chan et al.[19] then require the couplings of the four trajectories to the Kπ system and the nucleon to be adjusted so that their contributions cancel in the three exotic cases (29a), (29b) and (29c) but add constructively in the nonexotic case (29d). This is achieved by making all four contributions equal in magnitude and adjusting phases to add in the nonexotic case (29d). Since the three exotic cases are obtained from (29b) by reversing either signature or isospin or both, each of the three cases has two positive and two negative equal contributions and they all vanish.

So far there are no problems. Difficulties arise when we attempt to apply the same criterion to include the entire baryon octet and not just the nucleon. The $K^-\pi^-\Sigma^-$, $K^-\pi^-\Sigma^0$, $K^-\pi^-\Xi^-$, and $K^-\pi^-\Xi^0$ states are all exotic, but the nucleon case is related to these other baryons by the SU(3) sum rule[27]

$$A(p) - A(n) - A(\Xi^-) + A(\Xi^0) = A(\Sigma^+) - A(\Sigma^-) = \tfrac{1}{2}[A(\Sigma^0) - A(\Sigma^-)], \quad (30)$$

where A(B) is the amplitude for the $K^-\pi^-B$ case. Thus A(p) must vanish if $A(\Sigma^-)$, $A(\Sigma^0)$, $A(\Xi^-)$ and $A(\Xi^0)$ vanish, and SU(3) kills the good result for the nucleon case.

The good nucleon result can be kept while making Regge contributions cancel for all exotic kaon-pion-hyperon combinations only at the price of introducing couplings that strongly violate SU(3) between the Regge trajectories and the baryon octet. Although this seems unreasonable, there is still no experimental evidence for the coupling of these trajectories to strange baryons and the possibility exists that they may violate SU(3) in this violent fashion.

Difficulties also arise[25] if the same criterion is used with meson targets. The $K^-\pi^-\pi^0$ state is exotic, and the ρ, ω, and A2 do not couple to the π^0. Only the f^0 contribution remains, and it must decouple from the Kπ system. But the ρ, ω, and A2 contributions are required to be equal to the f^0 contribution in the nonexotic nucleon case (29d). Thus if $A(\pi^0) = 0$, $A(p) = 0$ and the good nucleon result is killed by factorization.

This case is an example of an "exotic fragmentation vertex," i.e., the $K^\pm\pi^\mp$ system has exotic quantum numbers. Similar difficulties have been pointed out[21] for the case in which the exotic fragmentation vertex is a meson-baryon system rather than two mesons. Again the Chan criterion leads to difficulties, but only if SU(3) is assumed for processes that have not yet been observed in the laboratory.

One attempt to preserve the Chan criterion has been the suggestion that all secondary trajectories are decoupled whenever the fragmentation vertex is exotic—e.g. that none of the reactions (29) have contributions from the secondary Regge trajectories.[22] In that case the cross sections for the two cases (29a) and (29d) should be equal because they are all represented by Pomeron exchange, which has even signature and is isoscaler. However, experimental data[26] indicate that they are not equal. The cross section for case (29d) is twice that for (29a). Hence the suggestion that secondary Regge trajectories cancel whenever the fragmentation vertex is exotic, regardless of whether the three-particle system is exotic, seems to be excluded by experiment.[25] However, the choice between the Stony Brook, CERN, Berkeley, and M.I.T. party lines is still open, and whether any one of them is right or all are wrong can be settled only by more experiments.

III. THE QUARK MODEL

The quark model continues to give a good and puzzling description of many regularities of hadron dynamics.[1,28] Nobody understands why the quark model works as well as it does, but nobody can convincingly dismiss all of the successes of the quark model as being completely accidental.

A. Successful Predictions of the Quark Model

The successes of the quark model are conveniently classified under the headings of (1) symmetry, (2) spectroscopy, (3) exoticity, and (4) universality.

1. Symmetry. The quark-triplet representation of SU(3) furnishes a convenient basis for SU(3) symmetry. The assumption

that quarks are universal building blocks leads naturally to SU(3), not only as a strong-interaction symmetry for the classification of hadrons but also as the algebra of the weak-interaction charges, and also gives the observed simple transformation properties of the electromagnetic current under SU(3).[28] People who dislike quarks say that it is possible to have SU(3) without them. But no other theory or model explains why the <u>same</u> mysterious SU(3) group appears everywhere in strong, electromagnetic, and weak interactions. There are also the higher symmetries such as SU(6), which are very difficult to explain without some kind of quark structure.

2. <u>Spectroscopy.</u> The predictions of the quantum numbers of the lowest-lying states for both baryons and mesons are an important success of the quark model. The spins and parities 0^-, 1^-, $\frac{1}{2}^+$, and $\frac{3}{2}^+$ of the lowest states are usually taken for granted. However, the odd parity of the low-lying mesons and their spins of 0 and 1 are difficult to obtain otherwise but arise naturally from the assumption that they are composed of fermions and antifermions in an s state. The $\frac{1}{2}^+$ and $\frac{3}{2}^+$ baryon states are also difficult to obtain otherwise[29] but arise naturally from three spin-$\frac{1}{2}$ objects.

3. <u>Exoticity.</u> A very large number of the successes of the quark model can be summed up by a general requirement of "absence of exotics." The quark model gives a few very simple rules defining what is exotic. Baryons are postulated to be made of three quarks. Mesons are postulated to made of quark-antiquark pairs. Weak, electromagnetic, and strong couplings or transitions are postulated to be sums of single-quark couplings or transition amplitudes. All other states, couplings, and transitions that do not fit these postulates are called exotic. The principle that "everything exotic is forbidden" leads to a remarkable number of successful experimental predictions. The postulates cannot be justified from any detailed model. There are no satisfactory answers to such questions as why four quarks are not bound more strongly than three or why something that looks like the impulse approximation should be valid for couplings and transitions. Yet from these postulates follow a large number of successful predictions and an explanation of the systematics of the experimental data, and these can hardly be accidental. The existence of the magic number 3 in baryon spectroscopy, limiting the maximum isospin to $\frac{3}{2}$ and the maximum strangeness to -3 is one particular striking example.[30] Thus there seems to be some validity in the quark-model exoticity rules, even if we do not understand where they come from.

Many other apparently disconnected predictions of the quark model can also be considered as absence of exotics. For example, the $\Delta S = \Delta Q$ rule in semi-leptonic weak interactions follows from the assumption that there are no exotic weak currents—i. e., no currents involving transitions of more than a single quark in a baryon. Similarly, the absence of an isotensor electromagnetic current is simply the vanishing of another exotic coupling. The requirement that Δ photoproduction be pure M1 is again a statement of the absence of an exotic amplitude. The E2 or L2 amplitudes normally allowed for a transition between a $\frac{1}{2}^+$ and a $\frac{3}{2}^+$ state are forbidden for transitions between two $\frac{1}{2}^+$ states. They are therefore exotic because they require transitions by more than a single spin-$\frac{1}{2}$ quark.

4. Universality Rules. The quark model goes beyond the symmetry schemes in giving relations between properties of mesons and properties of baryons. The most striking of these relations have been for the ratios of meson-baryon to baryon-baryon total cross sections.[31—34] Although there is no direct relation between the quark model and Regge theory, quark-model predictions are conveniently stated with the use of a Regge picture to label amplitudes having definite t-channel quantum numbers.[35, 36] The quark-model predictions state that each trajectory is universally coupled to all hadrons and that the coupling constants are given by a quark-counting recipe that includes a weighting factor depending on the trajectory and on the internal quantum numbers of the quark. The ρ trajectory is coupled to the isospin, the ω trajectory to the baryon number of nonstrange quarks, etc.

B. Universality Predictions and Serpukhov Data

It is interesting to review universality predictions in the light of present data, including the new total cross sections from Serpukhov.[37] The ω and ρ universality relations[27, 33, 38] hold up extremely well up through Serpukhov energies,[34, 39] as shown in Fig. 8. Because Serpukhov data for negative particles on liquid deuterium[40] are not yet available, we choose the linear combination of ρ and ω universality relations which require data only from proton targets.[34] The result is

$$\sigma(\overline{p}p) - \sigma(pp) = 3[\sigma(K^-p) - \sigma(K^+p)] - [\sigma(\pi^-p) - \sigma(\pi^+p)]. \quad (31)$$

The data and the lines through them have been taken directly from the Serpukhov paper and have been used to predict the total-cross-section differences $\sigma(\overline{p}p) - \sigma(pp)$ and $\sigma(K^-n) - \sigma(K^+n)$. The first is seen from Fig. 8 to be in remarkable agreement with experiment. The second is a prediction which will be tested by new data for kaons on deuterium.[40]

It is particularly interesting that the Serpukhov data show a very different energy decrease for the pion-nucleon and kaon-nucleon cross-section differences, a pattern suggesting that the ω

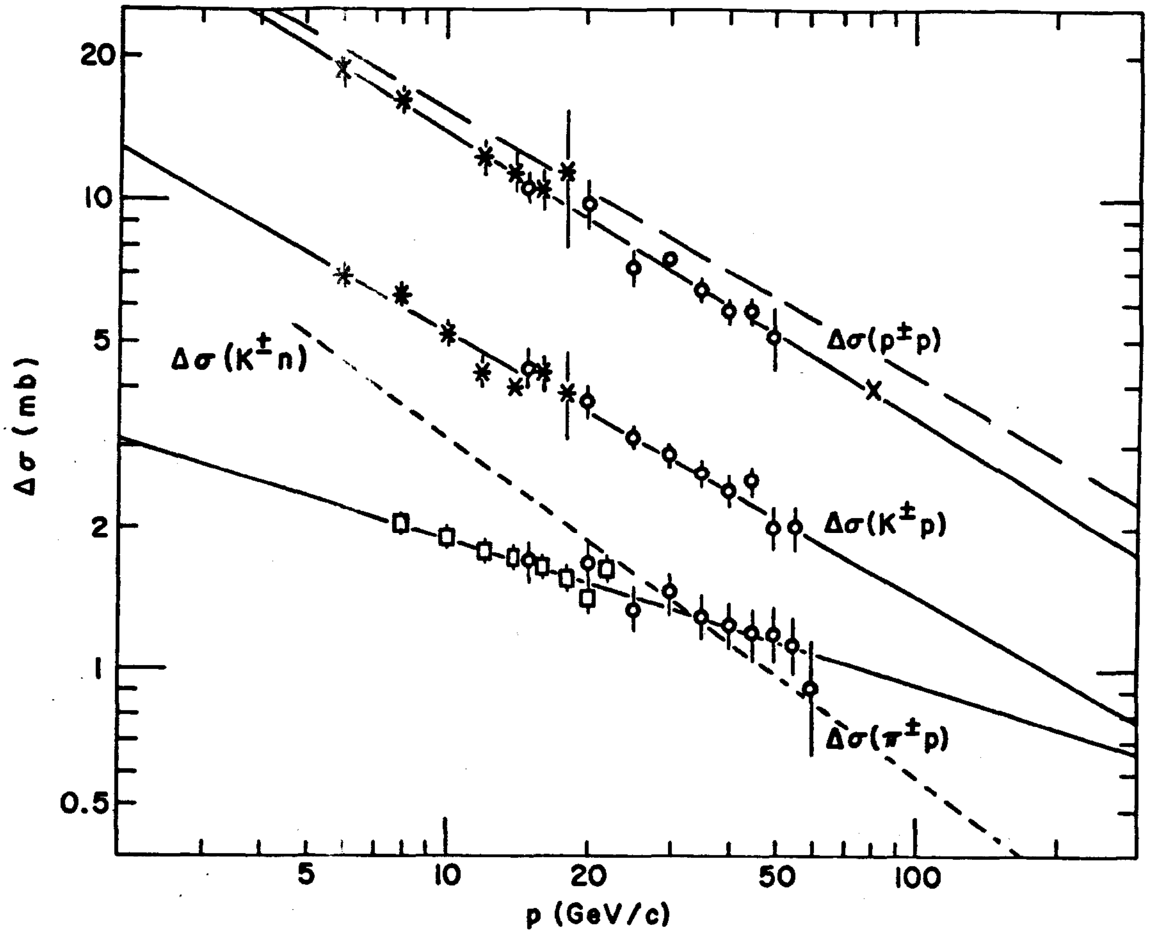

Fig. 8. Experimental tests and predictions from ρ-ω universality. The data and the straight-line fits to the data are from Ref. 37. The points and solid lines are for $\Delta\sigma(\pi^+p) = \sigma_{tot}(\pi^-p) - \sigma_{tot}(\pi^+p)$, $\Delta\sigma(p^+p) = \sigma_{tot}(\bar{p}p) - \sigma_{tot}(pp)$, and $\Delta\sigma(K^{\pm}p) = \sigma_{tot}(K^-p) - \sigma_{tot}(K^+p)$. The long-dashed line represents $3[\sigma_{tot}(K^-p) - \sigma_{tot}(K^+p)]$. Crosses (x) give $\Delta\sigma(p^{\pm}p)$ as predicted from Eq. (31). The short-dashed line gives the prediction for $\Delta\sigma(K^{\pm}n) = \sigma_{tot}(K^-n) - \sigma_{tot}(K^+n)$.

and ρ trajectories have very different intercepts. Despite these differences, the fit to the universality relation (31) indicates that the quark-model universal couplings are still valid for each trajectory. The validity of universality despite the difference in energy behavior will have a dramatic test in the kaon-neutron cross sections since the prediction obtained from the pion-proton and kaon-proton cross sections by use of ρ universality leads to the peculiar curve shown in Fig. 8, in which the kaon-neutron difference should actually cross the pion-nucleon difference somewhere in the Serpukhov energy range.

It has been suggested that the ρ and ω poles are still exchange degenerate and that the decrease with energy of the pion-nucleon cross section difference is due to an additional contribution from a ρ-Pomeron cut.[41] Such an explanation would still be consistent with the predictions of Fig. 8 within experimental errors because the ratio of the couplings between a ρ-Pomeron cut and pions and kaons would be approximately equal to the ratio of the couplings of the ρ pole [even including SU(3) breaking in the Pomeron coupling to pions and kaons gives an effect that is buried in the experimental errors] and the ρ contribution for the nucleon-nucleon case is too small to be sensitive to any deviation due to the cut.

It is interesting that these relations that involve only amplitudes with odd-signature exchange in the t channel are fitted by straight lines going through the lower-energy and Serpukhov-energy regions. They do not show the change in energy behavior seen, for example, in the $K^{\pm}p$ and pp total cross sections. This change is shown most dramatically (Fig. 9) in a plot of $\sigma(K^+p)$ vs $\sigma(pp)$ with the points taken at the same value of p_{lab}. This plot was originally published[42] to show that below 20 GeV the cross section $\sigma(K^+p)$ is constant while $\sigma(pp)$ is definitely decreasing, in contrast to the hand-waving duality statements that both cross sections should be constant because both channels are free of s-channel resonances.[43] Figure 9 shows this plot with the Serpukhov points added, and the difference is dramatic. Above 20 GeV, one sees that $\sigma(pp)$ suddenly becomes constant while $\sigma(K^+p)$ suddenly begins to increase, as reflected by the knee in the curve. This difference in their energy dependence suggests that very different dynamics may underly the even- and odd-signature amplitudes, whether or not the quark model is correct. This would be the case if the peculiar energy dependence can all be blamed on the Pomeron contributions.

The even-signature universality, which predicts the Levin-Frankfurt[31] ratio of $\frac{3}{2}$ for nucleon-nucleon to pion-nucleon cross sections, is not in as good agreement with experiment as the odd-signature predictions. There is a consistent discrepancy of about 20% over the entire energy range, the baryon cross section being larger than predicted. So far there is no satisfactory explanation of this deviation, although one might expect odd-signature relations to be better satisfied than even-signature relations because the odd-signature amplitudes are coupled to conserved quantities such as baryon number, isospin, and hypercharge.[28,29,35] Thus in models that contain seas of quark-antiquark pairs in addition to the three valence quarks or single quark-antiquark pair, the added pairs affect the quark counting for even-signature amplitudes but not for odd-signature ones.

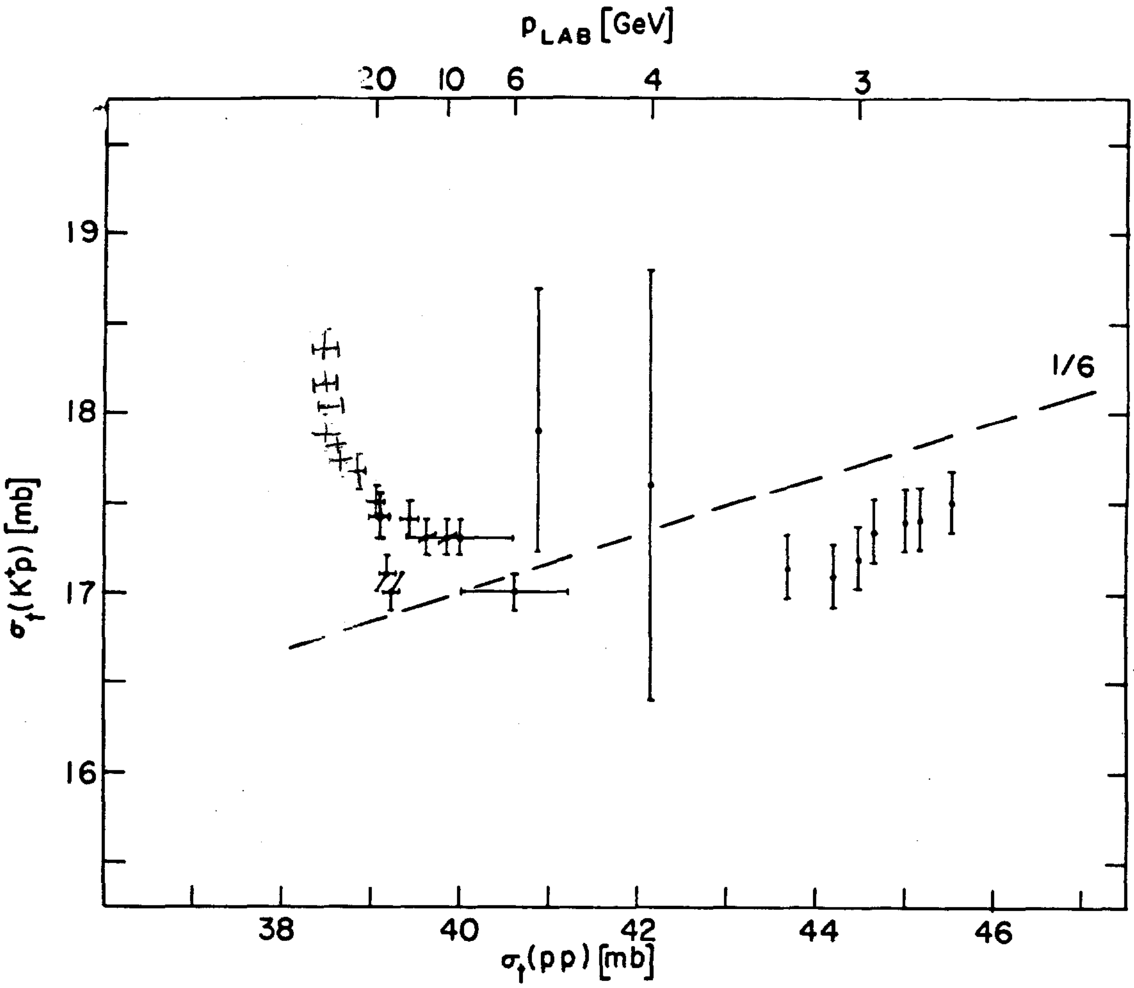

Fig. 9. Plot of $\sigma_{tot}(Kp)$ vs $\sigma_{tot}(pp)$. Serpukhov points from Ref. 37 are added to the plot given in Ref. 42.

The Levin-Frankfurt ratio comes from the even-signature isoscalar exchange amplitudes and has contributions from both the Pomeron and f^0 trajectories. It is therefore tempting to examine the energy dependence of this ratio to see whether the discrepancy can be blamed on one of these two trajectories while the other follows the quark-model rule. Such attempts do not work and indicate that a deviation of 20% in favor of the baryon couplings is present in both the Pomeron and f^0 couplings. This would be consistent with a picture in which both the f^0 and Pomeron couplings are determined by the same quark-counting recipe and the discrepancy from the simple quark-model result is due to the presence of additional quark-antiquark pairs. It would also arise in models of f meson dominance of the pomeron.[44]

C. Symmetry Relations for Production of Neutral Hyperons and Vector Mesons

Among the quark-model relations for scattering amplitudes are a series of predictions for neutral-hyperon and vector-meson production with strangeness exchange.[45—47] These have the peculiar property that some of the quark-model predictions are in very good agreement with experiment while others are in striking disagreement.[48, 44] The ratio of ρ to ω production agrees with theory while the production ratios Λ/Σ and $\Sigma/Y^*(1385)$ are in strong disagreement with experiment. However, predictions relating the vector meson polarization density matrices in reactions where they are produced with Λ, Σ, and Y^* agree with experiment.[49] The same predictions are obtained without quarks in a model involving octet-meson exchanges and the use of SU(3) and SU(6) symmetries in relating different couplings.[47] One possible explanation of the discrepancy has been the suggestion of configuration mixing in the baryon octet[50] —namely that there is a certain small admixture of decuplet analogous to the d-wave admixture in the deuteron.[51,52]

Using the conventional mixing-angle formalism, we can write the physical Σ as a linear combination of singlet and octet components,[50] namely

$$|\Sigma\rangle = |\Sigma_8\rangle \cos\theta + |\Sigma_{10}\rangle \sin\theta. \qquad (32a)$$

Since the mixing is expected to be small, we can try to treat it by first-order perturbation theory. This gives

$$\sin\theta = \frac{V}{\Delta E} \approx \frac{150-200}{1000} \approx 0.15-0.20, \qquad (32b)$$

where V is the SU(3)-breaking interaction, estimated as 150—200 MeV from the mass splittings in the baryon octet, and ΔE is the mass difference between the $|\Sigma_8\rangle$ and the $|\Sigma_{10}\rangle$. The $|\Sigma_{10}\rangle$ is assumed to be in the lowest $J = \frac{1}{2}^+$ decuplet, i.e., the decuplet containing the $\Delta(1910)$, so we take $\Delta E = 1000$ MeV.

The mixing (32) of 15—20% in amplitude or 2—4% in probability is sufficiently small to have been unnoticed until now. This mixing will produce a deviation δ_Σ from the Gell-Mann-Okubo mass formula, and calculation by second-order perturbation theory gives

$$\delta_\Sigma = V^2/\Delta E = 22.5-40 \text{ MeV}. \qquad (33a)$$

From SU(3), the mixing angle and the mass shift for the Σ are the same as for the Ξ, while the Λ and nucleon are not mixed because there are no states in a decuplet to mix with them. This compares very well with the experimental value

$$\delta_{exp} = 3M_\Lambda - M_{\Sigma^0} - 2M_n - 2M_{\Xi^0} = 30.7 \text{ MeV}. \qquad (33b)$$

Although this rough calculation cannot be taken too seriously, it indicates that mixtures of 15—20% in amplitude are quite reasonable, and suggests looking for experiments that can reveal the presence of mixing. The obvious places to look are phenomena for which the dominant octet contribution is suppressed or forbidden, so that the smaller decuplet contribution can be seen. This is the case in Σ production by strangeness exchange, and may account for the observed discrepancy. Quantitative tests of this explanation are difficult, since there are too many unknown quantities—e.g., the matrix element of the transition to the decuplet component.

D. General Theoretical Remarks

As the new accelerator at NAL begins to go into operation and there are many thoughts about searching for quarks, it is well to recall that all of the good predictions of the quark model are also obtained from the three-triplet model of Han and Nambu.[7] This model avoids three difficulties of the standard quark model at the price of postulating the existence of states in addition to quarks that have not been observed. The three advantages of the model are (1) no need for fractionally-charged particles, (2) no need for peculiar statistics to obtain a baryon wave function that is symmetric both in space and also in SU(6), and (3) a natural mechanism for the saturation of the baryons at three quarks by postulating strong binding only in the totally symmetric state. Whether these three advantages outweigh the disadvantage of postulating many additional unobserved states is a matter of taste. However, it is certainly not obvious that the three-triplet model is more unreasonable than the quark model.

With all the successes of the quark model, there has been no progress in the theoretical attempt to justify the quark model from first principles. All the conventional hand-waving arguments against the naive model with slow nonrelativistic motion despite strong binding can be countered by equally convincing hand-waving arguments.[28—30] However, it is not sensible to try to construct a quark theory from first principles because so many properties of the quarks are not known a priori. Consider, for example, the following questions.

1. What is the quark mass?
2. How many quark triplets are there?
3. What is the spin of the quark?

4. What are the electric charges of the quarks?
5. What are the statistics of the quark?
6. What kind of equation should be used to describe quark dynamics, e.g., Bethe-Salpeter, Dirac with self-consistent potential, or what?
7. Which approximation methods should be used in solving the dynamical equations?
8. What is the average potential determining the hadron wave function?
9. What is the basic quark-quark interaction (e.g., four-vector or world-scalar)?
10. Are hadron wave functions nearly pure quark-antiquark or three-quark or can there be a large sea of additional quark-antiquark pairs (e.g., partons)?
11. What is the quark-quark scattering amplitude?
12. What is the relation between hadron-hadron scattering and the quark-quark scattering amplitude?
13. What is the electromagnetic form factor of the quark?
14. What is the relation between the quark form factor and the electromagnetic coupling to a hadron?
15. What are the weak form factors of the quarks?
16. What is the relation between the quark form factors and weak hadron form factors?
17. What is the relation between strong three-point vertex functions and quark structure?

These seventeen questions and many others must be answered in at least one of two possible ways before one can embark on a well defined program for quark-model calculations from first principles. This means choosing one out of more than seventeen different models, each of which has equal a priori probability of being correct. The particular model chosen thus has a probability less than $<2^{-17}$ of being the correct one. This is not an attractive research program.

The fruitful approach is to avoid asking these questions and to use the quark model only for general hadron properties which do not depend on the detailed answers to such questions.

IV. RESONANCE SPECTROSCOPY

The most significant development in resonance spectroscopy over the past year has been that theorists now no longer feel the need to dream up new and even crazier explanations for the A2 splitting.[53] However, many interesting questions of resonance spectroscopy remain open and require attention from both theorists and experimentalists.

A. The Vector Mesons

Many unresolved questions remain in the spectrum of the lowest lying even-parity bosons. There is the question of whether all the states predicted by the quark model with p-wave orbital excitation are really where, and there are the usual controversies concerning the 1^+ mesons involving the existence of the A1 and the existence and structure of the Q bump. There are also the questions of the decay modes of the 1^+ states, for which both s- and d-wave decays are possible and for which both seem to be present but not in the mixture predicted by the most naive applications of the quark model.[54]

B. Quark Model vs Daughter Trajectories

There is also the question of whether there exist meson resonances[29] that are not predicted by the quark model but are predicted by Regge-pole models having daughter trajectories,[55] e.g., the Veneziano model.[56] The 0^+ state degenerate with the ρ and ω fits the daughter-trajectory models and is not inconsistent with the quark model although the degeneracy can only be viewed as a peculiar accident.[29] The real test comes with the $1^-(\rho')$ and 0^+ states which daughter-trajectory models predict to be degenerate with the f^0. These would not be easily explained in the quark model. So far there is no evidence for the ρ' but there seems to be evidence for a 0^+ s-wave $\pi\pi$ resonance[57] degenerate with the f^0.

The recent observation of the s-wave resonance[57] under the f^0 has an additional peculiar feature. The angular distribution of the decay pions shows strong forward and backward peaks but a very flat angular distribution between them. The secondary peak in the equatorial region, which would be present in a d-wave decay, seems to be exactly canceled by the s-wave component. If this effect is systematic and not a result of the peculiar kinematics of a particular experiment, it raises not only the question of the existence of the s-wave resonance degenerate with the f^0 but also the additional question of why the two resonances should be produced with just the exact mixture, including phase, so that the s wave cancels the intermediate peak in the d wave.

A simple quark-model description of $\pi\pi$ scattering[58] gives just this angular distribution in a straightforward way with no free parameters. In the center-of-mass system, with the incident momenta in the z direction, the resonant state is assumed to be a quark-antiquark pair in a p wave and in the triplet spin state. If both L_z and S_z are conserved in the resonance production process, as is commonly assumed in the quark model, the resonant state has $L = 1$, $S = 1$, $L_z = S_z = 0$. This is a coherent mixture of $J = 0$ and

$J = 2$, which decays with a $\cos^4\theta$ angular distribution, according to the standard quark-model prescription.[54] However, this model assumes that the 0^+ resonance is degenerate with the f^0 and therefore leaves no explanation for the 0^+ state degenerate with the ρ. Further experimental investigation of this regularity would be of interest, as would a search for similar effects in the angular distribution of the two pions emitted in the decay of the g meson.

C. Nucleon-Antinucleon Annihilation

Additional experimental evidence for possible new interesting structures in boson spectroscopy have come from studies of nucleon-antinucleon annihilation. The natural-parity states studied in two-meson channels, and particularly the peculiarities of the two-kaon final states,[59] indicate coherent production of degenerate isoscalar and isovector states. The absence of the K_1K_1 final state while the K_1K_2 continues to be observed up to energies at which many partial waves are present indicate that only odd-parity states contribute to the $K^0\bar{K}^0$ final state. However, the pronounced forward-backward asymmetry observed in the K^+K^- final state indicates that both even- and odd-parity contributions are present in the charged mode. The absence of even-parity contributions in the neutral final state can come only from cancellation of isoscalar and isovector contributions in the particular even-parity partial waves that contribute to the charged final state.

It has also been suggested that nucleon-antinucleon annihilation is dominated by the Regge recurrences of the pion.[60] This would show up in multiparticle final states and suggest that states with odd numbers of pions should dominate over states with even numbers of pions, which have the wrong G parity to be produced by the Regge recurrences of the pion.

Further evidence for peculiar coherence effects between the isoscalar and isovector amplitudes in proton-antiproton annihilation come from the experimental observations of strong ρ-ω interference in various final states containing only pions.[61] Such interference involves states of opposite G parity and therefore can come only from the isoscalar and isovector components of the initial state.

Interesting structure in the nucleon-antinucleon system is also indicated by phenomena observed in annihilation at rest and at low energies in deuterium. These have been interpreted as a possible resonance below the nucleon-antinucleon annihilation threshold.[62]

D. Baryon Spectroscopy

Baryon spectroscopy continues to be consistent with various

SU(3) predictions and with the quark model, with no conclusive evidence in those areas in which different versions of the quark model would give different predictions.[63]

V. WHY SEARCH FOR EXOTIC PARTICLES?

This last section summarizes remarks at a recent NAL seminar given to provide theoretical background for a discussion of the various particle-search proposals presented to NAL. The best answer to the question: "Why search for new particles?" is "Why not?" The relevant question is not why but how and what to search for. As one goes through the suggestions to look for quarks, monopoles, vector bosons, heavy leptons, tachyons,[64] and partons, one might wonder whether there is not a much more exciting new kind of particle waiting to be discovered at NAL but not in any proposal because the theorists have not yet thought of it. How to find such a particle is left as an exercise for the reader—with no constructive suggestions other than to keep an open mind and avoid being brainwashed by theoretical arguments, including this one.

The objects of searches proposed today are widely different and have different theoretical motivations. Their one common denominator is that all theoretical motivations are weak and the probability that any one of the particles will be found is rather small. As guide lines for evaluation of particle searches let us consider two examples.

1. The search for the Ω^-. At the time the existence of the Ω^- was proposed, it was considered a wild speculation by the theoretical establishment. However, it was a well-defined prediction. There was a well-defined symmetry which had had certain successes and which predicted that this particle should exist, giving its mass as well as its quantum numbers, and these predictions suggested reactions in which it might be produced and gave rough estimates of cross sections (which might be off by one or two orders of magnitude). The finding of the Ω^- was certainly a great milestone in the establishment of SU(3) symmetry. However, failure to find the Ω^- after a considerable period of search would have been a fairly decisive blow to SU(3) symmetry. Thus the Ω^- was an ideal subject for a search. A positive result was reasonably exciting but a negative result would also have important significance.

2. A search for a neutral boson with a mass near the pion mass, coupled only to strange baryons and decaying only into five photons. This is a particle I invented for this discussion; it has no deeper theoretical background. However, there is no good evidence against the existence of such a particle. The purpose of this example is to point out that the existence of many kinds of exotic particles can be proposed without contradicting present experimental

evidence but also without any theoretical motivation whatsoever. The most exciting possibilities awaiting discovery at NAL may be just of this type—namely, those for whose existence no theorist has yet discovered a reason. The problem is not whether to look for such particles but how. The present example shows that it is fairly easy to dream up a million different particles, each requiring a different expensive experimental technique for a search and each equally plausible on theoretical grounds. We cannot search for all of them; yet one of them may be there. How do we decide what to do?

All of the examples under serious consideration at NAL fall somewhere between the two cases discussed above. They are not as well defined as the Ω^-, but they also have somewhat more theoretical justification than the ad hoc boson. However, none of them is sufficiently well defined to make a negative result in any experiment conclusive proof against their existence. A negative result can only say that a particular experiment is not the right way to look for the particle and cannot indicate whether or not they may be found in some other experiment.

As an example of the difficulty that plagues all these searches, consider the mass of the quark. Suppose a theorist who believes in quarks were asked for an estimate of the quark mass with error bars. The only honest answer would be

$$M_Q = \infty \pm \infty. \tag{34}$$

The simple reason for this ridiculous result is not widely appreciated. The quark mass has a lower bound because quarks have not yet been seen, and in any case the mass must be positive. There is no upper bound on the quark mass.[65] Furthermore, there is no theoretical reason why one allowed value of the quark mass is more probable than any other. The inevitable conclusion that all values of the quark mass from the lower bound up to infinity are equally probable leads to Eq. (34).

It is very exciting if the quark is "just around the corner," i.e., if its mass is just too high to be discovered by the present generation of accelerators but is waiting for the next one. But there is no theoretical justification for this optimism. Even if quarks are there, there is no reason why their mass should not be at least two orders of magnitude higher than anything that can be produced at NAL. There are those who argue that a negative result of a particle search is significant because it can give new limits on the properties of the particle and may eventually lead to contradictions with theory. In the case of the quark, however, Eq. (34) gives the present limits on the quark mass and will not be appreciably affected by any negative result at NAL.

There will certainly be quark searches at NAL because it is very easy experimentally to look for particles with fractional charge. The excitement if they should be found certainly justifies the relatively small expenditure of money, manpower, and machine time, even if the probability of finding them is quite small and a negative result proves nothing.

With this general background in mind, it seems reasonable to classify the proposed searches according to the following four criteria.

1. Who needs it?

2. If not found, so what? Unfortunately most cases can be dismissed with the comment: "Try harder." That is, all that has been shown is that a particular experiment that has not found them is the wrong way to look for them.

3. Craziness index. This indicates the degree of personal prejudice on the part of the establishment against certain aspects of these particles—e.g., against fractional charge, imaginary mass, etc. Fractional charge is crazy because leptons have integral charge and behave like elementary Dirac particles to a precision (g - 2) greater than anything that can be said about strong interactions. This remarkable result is lost with a composite lepton model,[28] and the suggestion that leptons have three times the fundamental charge seems crazy. Particles with imaginary mass have space-like momenta. A tachyon produced by a beam hitting a target will appear in a different Lorentz frame to have been emitted from elsewhere before the experiment and to arrive at the target just in time to be absorbed at the exact instant at which the beam hits it. This seems crazy enough to me. (A word of caution: Parity nonconservation and CP violation once had very high craziness indices.)

4. Signature. If it is easy to verify that a particular particle is there (i.e., if it has a clear signature), it is much easier to look for it.

Table I summarizes the various particles whose search has been suggested and lists their properties.[66] All these searches are "long shots" and involve a certain element of gambling. The low probability of finding anything must be considered in the context of the size of the reward, the investment, and the possible spin-off. Broad exploratory searches, which also have the possibility of finding unexpected objects, should be encouraged. Remember that Columbus looked for India and found America and that Fermi looked for transuranium elements and found fission.

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TABLE I. Guide to inconclusive experiments and hypothetical particles.

Particle	Who needs it?	If not found, so what?	Craziness index	Signature
		I. PEDAGOGICAL EXAMPLES		
Ω^-	MGM & YN	Kills SU(3)	No	Good
$M^0(150) \rightarrow 5\gamma$	Nobody, but why not?	Nobody cares	Not particularly	Missing mass
		II. PROPOSED SEARCHES AT NAL		
Tachyons	Nobody, but why not?	Nobody cares. Try harder	Very	Good
Quarks	Dalitz	Try harder	Fair	Good (fractional charge)
Monopoles	Dirac-Schwinger	Try harder	Moderately	Good
Intermediate bosons	Yukawa	Try harder, but credibility falls	No	Good
Heavy leptons	Nobody, but why not?	Look elsewhere (spectroscopy)	No	Good
Partons	Bjorken-Paschos	Ask Bjorken-Paschos	No	Good
Han-Nambu triplets	Dalitz might settle for these	Try harder	Less than quarks	Missing mass best
Superheavy nuclei	Nuclear physicists	Try harder	No	Chemistry. Not clean
		III. THE REALLY EXCITING SEARCH		
?	Nobody has thought of it	It will be found; it's there	Who knows, the theorists have not thought of it yet	

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