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High Energy Spectroscopy*

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I. INTRODUCTION

How can spectroscopy be studied with a new high energy accelerator? A new machine should discover new things and it is not easy to suggest how to look for them, particularly since the most exciting new discoveries have unexpected and surprising properties. Suggestions from theorists are of dubious value. Even when they are right their advice is usually useless and following it exactly usually leads to missing something crucial. But something equally crucial can be missed by ignoring their advice. After each discovery it usually turns out that some theorist predicted it. But dozens of equally plausible suggestions also made at the same time led nowhere and it was by no means obvious which approach would be fruitful. This makes life difficult for experimentalists and program committees trying to decide what experiments to do. But if their tasks were easier and the outcome of experimental investigations could be predicted in advance, research would be much less exciting.

The recently discovered new charmonium spectroscopy presents an instructive example of these difficulties. At the 1975 Palermo Conference I was given credit[1] for predicting the discovery of these particles on the basis of the analysis[2] shown in Table I.1 of the new particle search proposals in 1972 at Fermilab. The conclusions were that the searches for quarks, monopoles, tachyons, etc. were not apt to lead anywhere and that the really exciting search would discover a

TABLE I.1 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
$M^0(150) \to 5\gamma$	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 <u>there</u>	Who knows, the theorists have not thought of it yet	

particle not listed in these proposals and which the theorists had not thought of. This prediction is not strictly correct if the new particles[3] discovered since November 1974 are indeed bound states of charmed quarks and antiquarks as they seem to be today. Such states were proposed by theorists[4] a long time ago and their properties were investigated in detail. However, in 1972 there were no charm search proposals at Fermilab. Even in the summer of 1974 when charm searches suddenly became fashionable and theorists suggested ways of looking for charm,[5] there was no suggestion that charmonium or hidden charm would be found long before charm itself or that the most fruitful search would be for very narrow states produced in electron-positron annihilation. The reason why these suggestions were not made is instructive. Two crucial missing links in our understanding of hadron properties prevented the appropriate suggestions from being made and taken seriously. These were the existence of neutral weak currents[6] and the mysterious selection rule attributed to Zweig, Okubo, Iizuka and others.[7, 8, 9, 10]

In 1971 hadron spectroscopy was well described by the conventional quark triplet with three quarks and no fourth quark was needed to describe the observed states. The motivation for charm came entirely from weak interactions where a number of attractive looking theories encountered difficulties in predicting the existence of neutral weak currents[11] in flagrant contradiction with experiment. The introduction of a fourth charmed quark with the GIM mechanism[12] cancelled out all the strangeness changing neutral currents and removed the disagreement with experiment. But the strangeness conserving neutral currents were not cancelled and there was no experimental evidence for such weak neutral currents. There was also no convincing evidence against them, but most particle physicists assumed that this was simply a problem of experimental techniques. Sensitive experiments testing strangeness-changing neutral currents were much easier than tests of strangeness-conserving neutral currents, and there was no obvious reason why one should be absent while the other was present. Thus a model which looked attractive to theorists did not seem attractive to experimentalists because it predicted all kinds of unobserved experimental results and then had to introduce various ad hoc cancellations to get rid of them. Furthermore the same theorists of the Harvard group who proposed the charm model to get rid of strangeness changing neutral currents had more complicated models[13] with additional heavy leptons that could get rid of all neutral currents. There was a general proliferation of models each introducing either new quarks, new leptons or new ad hoc couplings of electromagnetic and weak currents. They were all equally believable and each suggested different experiments to test its validity. It was hard for an unprejudiced experimentalist to know which model should be taken seriously or whether the whole picture of gauge theories was worth considering seriously at all.[14]

Everything changed with the discovery of the weak neutral currents.[6] It was now clear that nature had placed the strangeness conserving and strangeness violating neutral currents on a completely different basis and the most natural explanation for this difference came from the GIM mechanism[12] which required the existance of charm. So the charm model suddenly jumped from being one of many dubious theoretical models with ad hoc assumptions not justified by experiment to the simplest and most reasonable model available which would explain a very striking and important new experimental result.[15] Attention immediately turned to charm searches.

The charmonium states, bound states of a charmed quark-antiquark pair were also predicted, and it was also realized that the decay of these states would be inhabited by the same OZI selection rule which prevents a strange quark-antiquark pair from disappearing in the ϕ meson decay to produce final states without strange quarks. However, estimates of the suppression factor were off by a large factor because the width of the $\phi \rightarrow \rho\pi$ decay was the only experimental evidence available for the strength of transitions violating the selection rule. Why the charmonium states are so much narrower is still not understood.

It is now 2 1/2 years since the J particle was produced at Brookhaven by Sam Ting and collaborators. But even though we recognize the importance of Ting's discovery and great effort has gone into subsequent investigations we still know very little about the production mechanism for the J in those experiments.

The OZI rule allows this J production only with an accompanying pair of charmed particles. But there is no evidence for this charmed pair, and the J production seems to go via some mechanism[10] which violates the OZI rule.

Except for this absence of charmed pairs we know very little about the final state in the reaction which includes the J. Thus, it is very difficult to estimate production cross sections for other new objects in hadronic experiments and any extrapolation of Ting's results for such estimates contain so many unknown factors that they are extremely unreliable. Since the narrow width of the J is not understood all estimates of the strength of couplings of new objects to ordinary hadron channels are unreliable. Isabelle experiments might provide new insight into these fundamental uncertainties.

All properties of the charmonium states were predicted well except for the most striking property, the very narrow width which was crucial in their discovery. Similar theoretical considerations and difficulties can be expected to arise in predicting the properties of states to be discovered with new high energy accelerators. So theoretical guidelines should not be dismissed but should be considered with the view that they may be even 90% correct, but a crucial 10% may be missing.

This talk consists of two parts. The first part, Sections II and III gives some general guidelines for high energy spectroscopy pointing out crucial differences between the search for new objects at high masses and conventional low energy spectroscopy. The second part, Sections IV, V and VI reviews some puzzles and open questions in conventional spectroscopy which might find solutions in experiments at higher energies. Some specific predictions or suggestions are made as examples, but they are presented primarily to stimulate thinking of experimentalists along new lines rather than to provide specific instructions which should be followed literally.

II. SIGNAL AND NOISE IN HIGH MASS SPECTROSCOPY

Resonances with masses in the several GeV range have very many open decay channels. Their branching ratios into any one exclusive channel are of the order of 0.1%. Since the signature for the detection of such a resonance generally picks a particular decay mode, the signal is proportional to the branching ratio and is very small. The crucial factor in discovering and confirming such high mass resonances is the signal to noise ratio.

It is useful to define a figure of merit F(P,T) for the production of particle P, by observing a characteristic T of the final state which may either be used as a trigger or as a signature for picking out events. The trigger T may be either the full final state like the electron pair in the decay of the J, or one of the particles produced inclusively in the decay such as a single muon. The figure of merit is defined by the relation

$$F(P,T) = \sigma(P+X) \cdot BR(T)/\sigma(T+X) \qquad (2.1)$$

where $\sigma(P+X)$ and $\sigma(T+X)$ denote the cross sections for inclusive production of the particle P and the trigger T in the reaction under consideration and BR(T) denotes the branching ratio for the appearance for the trigger T in the decay of the particle P.

Examination of Eq. (2.1) shows that the optimization of the figure of merit may best be achieved by finding a trigger T with low inclusive production. The characteristics of the signal appearing in the numerator will not be changed very much by choosing a different trigger or a different production mechanism. However, the denominator may be reduced by a large factor by choosing a trigger for which the background is low. Possibilities for improving F(P,T) by reducing the noise seem to be more favorable than by enhancing the signal. We examine three possible approaches to noise reduction.

1) Production of a low noise signal. The signal can be produced by a mechanism which naturally has a low background, as in the production of the Ψ as a very narrow resonance in e^+e^- annihilation.

2) A low noise signal signature. An exclusive decay channel can be found which has a low production background as in the detection of the J particle by its leptonic decay mode. The particular case of ϕ signatures is of interest.

3) Use of background signature. Since many partial waves in the background can appear at the high mass available and only a few in the signal, the background may have a characteristic structure which enables cuts in selected kinematic regions of the multiparticle phase space to reduce the noise by a large factor.

2.1 Production of Low Noise Signal

The production of a new particle with a very low background is possible for a narrow s-channel resonance whose cross section is very much enhanced over the background in a narrow energy region. This approach can be used only for the production of resonances having the quantum numbers available in the initial state. It is particularily suitable for the production of vector meson resonances in electron-positron annihilation.

In a proton-proton colliding beam machine, it would be useful only for nonstrange resonances with baryon number two. If such exotic objects should indeed exist, at high mass, they might be found most easily in this way. For example if an exotic "molecule" containing six nonstrange quarks and a charmed quark-antiquark pair (i.e. a deuteron-J/ψ bound state) should exist with a mass below the mass of two nucleons and a J/ψ, it might be observed in this way. Whether it is worth searching for such objects depends on how much trouble is involved in looking for them and on the general state of dibaryon and molecular exotics at the time.

For states which do not have the quantum numbers of the photon or of the meson-baryon, nucleon-nucleon or nucleon-antinucleon system, some possibilities exist for production via the decays of states which do have these quantum numbers; e.g. in the production of the positive parity charmonium states by radiative decay of the ψ' and the production of charmed particle pairs by the decays of higher vector resonances.

For states not easily produced in this way and available only in inclusive production there is no simple mechanism for reducing the multiparticle background by choice of a particular production mechanism. This applies to most cases of hadronic resonance production, as in J production where no one production mechanism seems to be superior by any large factor.

2.2 Low Noise Triggers and φ signature spectroscopy

The triggers which have low inclusive production cross section in normal hadronic processes include photons and leptons produced by electromagnetic interactions. These are suppressed by powers of α relative to hadron production. Some examples are the lepton pairs used as the signature for the discovery of the J particle, the photons used as a signature to discover even parity charmonium states produced by the decay of the ψ' and the two-photon and multiphoton channels used for the possible detection of the pseudoscalar mesons.

In addition to these electromagnetic triggers which have already been used successfully, particles like the φ and f' which are suppressed by the OZI rule in nonstrange hadron reactions might be used successfully. These appear as signatures for states whose branching ratios into decay channels involving φ and f' are not suppressed by significant factors over other decays. φ signature spectroscopy looks attractive for states decaying into a φ because inclusive φ production without kaons is forbidden for nucleon-nucleon and pion-nucleon reactions and the background should be small. Typical suppression factors observed experimentally for φ production are a factor of 500 below _ production in pion-nucleon reactions[16] at 6 GeV/c or a factor of 100 below pion production at Fermilab energies.[17] The ϕ is easily detected in the $K^+ K^-$ decay mode at high energies because the Q of the decay is so low that both kaons will pass together in the same arm of a spectrometer and will not trigger a Cerekov detector set for pions.[17] An even smaller background would be expected in $\phi\phi$ spectroscopy for states expected to decay into two ϕ's. Examples of such states are isoscalar bosons even under charge conjugation which have the structure of a quark-antiquark pair, either strange, charmed, or some new heavy quark.

"Strangeonium" states of a strange quark-antiquark pair are allowed by the OZI rule to decay into $\phi\phi$ and should have a comparatively strong branching ratio. Such strangeonium states are of general interest since no such states above the ϕ or f' are well known. Our present knowledge of charmonium spectroscopy is at present much better than strangeonium because the low noise electromagnetic signature of lepton pairs and photons enables charmonium to be seen much more easily. Even if $\phi\phi$ spectroscopy does not lead to the discovery of any new charmonium or "x-onium" states made from heavy quarks of type x, the development of strangeonium spectroscopy would add to our understanding of hadron dynamics.

The ϕ decay of charmonium or x-onium is singly forbidden by OZI or other quark line rules and is therefore on the same footing as all other hadronic decays which are also at least singly forbidden. Estimates of the $\phi\phi$ branching ratios for these particles are of the order of 0.1%, which is probably only a small factor below the $\rho\rho$ branching ratio. The $\phi\phi$ background should be very much lower than the $\rho\rho$ background and therefore can provide a fruitful trigger for such states.* The most interesting of such states at present are the pseudoscalar states of charmonium or of the new heavier quarks if they are there.

Single ϕ spectroscopy would be useful also in observing decays of higher strange resonances such as K^*, Λ^*; Σ^*, Ξ^*; * and ϕ^* which could decay into lower resonances with the same quantum numbers by ϕ emission above the threshold. Nonstrange baryon resonances at high masses have been observed by the technique of pion-nucleon phase shift analysis. ϕ spectroscopy may enable the discovery of corresponding resonances with different quantum numbers not accessible to phase shift analysis.

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*One estimate for $(x\bar{x})_{C=+} \to \phi\phi$, is based on the analogy with $\psi \to \phi\eta$ which also involves annihilation of a heavy quark pair and creation of two strange quark pairs. Another is based on the analogy $\psi \to \rho\pi$ and $(x\bar{x}) \to \rho\rho$ and used SU(3) to relate $\phi\phi$ to $\rho\rho$.

States like the $F^{\pm}$ meson containing both charm and strangeness might be observed by the decay into a ϕ and a pion or lepton pair. The $\phi\pi$ decay mode might also be useful in the search for the exotic[18,19] four-quark states discussed in section V.

The $\phi\pi$ decay mode is particularly interesting in searches for new objects, because $\phi\pi$ decay is forbidden by the OZI rule for <u>any</u> boson constructed from a quark-antiquark pair. Thus resonances in the $\phi\pi$ system indicate either a new object like a four-quark system, an OZI-violating strong decay of a conventional boson, or a weak decay into a system containing a strange quark-antiquark pair.

A partial list of states which might be detected by ϕ signature spectroscopy are

<u>Single ϕ spectroscopy</u>:

$$K^* \to K + \phi \tag{2.2a}$$

$$\Lambda^* \to \Lambda + \phi \tag{2.2b}$$

$$\Sigma^* \to \Sigma + \phi \tag{2.2c}$$

$$\Xi^* \to \Xi + \phi \tag{2.2d}$$

$$\Omega^* \to \Omega^- + \phi \tag{2.2e}$$

$$\phi^* \to K + \bar{K} + \phi \tag{2.2f}$$

$$F^{\pm} \to \pi^{\pm} + \phi \tag{2.2g}$$

$$F^{\pm} \to \text{leptons} + \phi \tag{2.2h}$$

$$\tilde{F}^{o}_{I} \to \pi^{o} + \phi \tag{2.2i}$$

$$A^{S}_{\pi} \to \pi^{o} + \phi \tag{2.2j}$$

<u>ϕ-ϕ spectroscopy</u>

$$\eta_c \to \phi + \phi \tag{2.3a}$$

$$\text{charmonium } (c\bar{c})_{C=+} \to \phi + \phi \tag{2.3b}$$

$$\text{strangeonium } (s\bar{s})_{C=+} \to \phi + \phi \tag{2.3c}$$

$$\text{x-onium } (x\bar{x}, \text{ where x is a new heavy quark})_{C=+} \to \phi + \phi \tag{2.3d}$$

Above 3 GeV the possibility of observing 3ϕ decay arises. Vector meson states like the ψ' and other higher members of the ψ family can decay into three vector mesons. The dominant 3V final state would be $\omega\rho\rho$ but 3ϕ would be of the same order of magnitude in the SU(3) symmetry limit. The 3ϕ state would have a unique signature and a very low background.

The use of ϕ triggers can thus lead to various kinds of interesting physics. The first step is the understanding of ϕ production itself, by examining the other particles produced along with the ϕ and looking for ϕx resonances. Understanding the mechanisms for ϕ production can provide insight into models for particle production, even if no new phenomena or resonances are found. But chances are that some part of the production will be due to decays of higher resonances, and at this stage any resonance with a ϕ-decay mode is interesting.

Other signatures which had been suggested for low noise background[20] are K_1K_1 and $\Lambda\bar{\Lambda}$. Although these states are not as forbidden as the ϕ, their inclusive cross sections are much lower than those for pions and the background can be further reduced by looking in appropriate kinematic regions. For example, one can look in a forward direction with incident beams that will not produce one of these particles by the strongest peripheral nonexotic exchanges; e.g. looking at forward K^+ with an incident K^- beam. However, there would be a background from decay products of diffractively produced resonances. The observation that hyperon beams can be produced with incident protons with an intensity of Σ^- equal to that of π^- in the forward direction[21] is a warning against depending too much on the absence of such production.

2.3 Background signatures

The signal to noise ratio can be improved by the alternative approach of characterizing peculiar signatures for the background in order to enable its removal from the signal. This approach is based on the fundamental difference between the spectroscopies of the high mass resonances and old low-lying resonances. The conventional low-lying resonances show up as peaks in cross sections with particular decay angular distributions against a comparatively smooth and structureless background. At high mass the background may have a more striking and easily identified structure than the signal.

High mass resonances are states of low angular momentum decaying primarily into multi-particle channels. Their decays reflect the low angular momentum by containing very few parital waves all having relatively low angular momentum. The background on the other hand can have very large angular momenta and a sharp structure in momentum and angular distributions are present in the signal. A small portion of the multi-particle phase space could include a very large portion of background events. In this case the signal to noise ratio would be improved by a cut excluding this small volume of phase space. The exact kind of cut to be effective depends on the individual case and could be most easily decided by examining the background and looking for its most striking features.

Consider for example the search for a new particle in a particular four-particle decay channel by looking for peaks in the mass spectrum, e.g. looking for a charmed baryon decaying into $\Lambda 3\pi$. The problem is how to use the angular distributions of these four particlesin the center-of-mass system of the four particle cluster (hopefully the rest system of the new particle) as a means of distinguishing between signal and background. Three axes are relevant for examining the angular distributions, (1) the direction of the incident beam momentum, (2) the direction of the momentum of the four-particle clusters, and (3) the normal to the production plane. Signatures which characterize the new particle appear most clearly in angular distributions with respect to the direction of the momentum of the four-particle cluster or with respect to the production plane. But signatures for the noise will show up in angular distributions with respect to the incident beam direction.

Background from uncorrelated particles whose mass happen accidently to fall in the desired range should have angular distributions with respect to the incident beam direction similar to those for single-particle inclusive productions. They should be peaked in the forward and backward directions with a rapidly falling cutoff in transverse momentum. Background events could show forward-backward asymmetry or a tendency to be concentrated in cones forward and backward relative to the direction of the incident beam. The signal from decay of a D meson of spin zero should show a completely isotropic angular distribution with respect to any axis. Particles of non-zero spin might have some anisotropy in their angular distributions if they are polarized in production. But these will involve only low order spherical harmonics and will not concentrate large numbers of events in a small region of phase space. Thus a cut eliminating events in which one or more particles appear within a narrow cone forward and/or backward with respect to the incident beam direction could reduce the background considerably with a neglible effect upon any signal coming from the decay of a low angular momentum state.

As an example consider a four particle decay into a baryon and three pions of a state produced by a high energy accelerator beam hitting a fixed target. This state appears as a four particle cluster with a low mass in the several GeV region but with total laboratory momentum in the 100 GeV range. In the center-of-mass system of the cluster the momenta of the baryon and of the pions are all small and of the same order of magnitude. In the laboratory the baryon has a much larger momentum than the pions because of the effect of the mass on the Lorentz transformation. If the baryon is not a proton and cannot be a leading particle the inclusive momentum distribution for the baryon and the pions can be expected to be very different in the relevant ranges. In particular the momentum distribution for high momentum hperon or anti-hyperons could be falling rapidly in this region while the momentum distribution for relatively low momentum pions could be rising. This would appear in the center-of-mass system for the multi-particle cluster as baryons being preferentially emitted backward and pions preferentially emitted forward. Cutting out events in which all pions are in the forward hemisphere would thus appreciably reduce the background, but would only remove one eigth of the signal. Using a cone instead of a hemisphere would interferere even less with the signal and still substantially reduce the background.

III. QUARKONIUM SPECTROSCOPY

Among the new exciting states hopefully waiting to be discovered are sets of positronium-like mesons made of a quark-antiquark pair with the same flavor. These include "strangeonium" states like the ϕ and f' of a strange quark-antiquark pair, charmonium states like the J/ψ family, and states made from quarks of new flavors as yet undiscovered.

3.1 Flavor dependence of the spectrum

Strangeonium ($s\bar{s}$) spectroscopy is still in its infancy, and is not yet as well developed as charmonium spectroscopy, even though strangeness was known over two decades before charm. The reason for the comparatively slow development of strangeonium spectroscopy is the absence of a good signature having a high figure of merit like the electromagnetic signatures used to detect charmonium states. The dominant decay modes of the strangeonium states are $K\bar{K}X$ which are allowed by the OZI rule and which also appear in the background. As a result the higher strangeonium states are expected to be broad, have comparatively low branching ratios to electromagnetic channels, and no striking signature different from background below the $\phi\phi$ threshold.

Charmonium ($c\bar{c}$) has given rich experimental results because the dominant OZI allowed decay channel, $D\bar{D}$, is closed for a large set of low-lying states including the radially excited s-wave (the ψ') as well as the lowest p states. Thus these states are all narrow and have appreciable branching ratios and couplings to electromagnetic channels like e^+e^-, $\mu^+\mu^-$ $\gamma\gamma$ and γX. The vector mesons states are therefore easily produced in e^+e^- annihilation and photoproduction experiments, and can also be detected by leptonic decay modes if produced by other means. Other states can be produced by cascade decays of the higher vector mesons and recognized by the presence of photons from the decay which produced them or from their own decays.

Higher x-onium states from heavier quarks with new flavors are expected in many theoretical models, but none had been seen at the time of the conference. At the time of this writing evidence for such a state has been reported.[23] Eichten and Gottfried[24] have pointed out that such states should show an even richer spectrum than charmonium, because of theoretical arguments showing that more states lie below the OZI-allowed threshold for increasing quark mass. This threshold for the decay of an $(x\bar{x})$ meson is at twice the mass of the lowest $(x\bar{u})$ state; e.g. $2M_K$ for strangeonium and $2M_D$ for charmonium. Eichten and Gottfried argue that the lowest vector state, analogous to the ϕ for strangeonium and the ψ for charmonium, is farther below the threshold as the quark mass increases, continuing the trend seen in the ϕ and the ψ. Thus the range of excitation energy available for narrow OZI-forbidden resonances increases with quark mass.

3.2 Quarkonium production mechanisms

Quarkonium production for states with flavors absent in the initial state is forbidden in strong interactions by the OZI rule. Electromagnetic $(x\bar{x})$ pair creation is not suppressed and is comparable to other $(q\bar{q})$ production if the x-quark has an electric charge. However, the production of $(x\bar{x})$ from a single photon occurs only for states with the same quantum numbers as the photon, namely odd-C vector mesons.

Processes involving the Pomeron might not be suppressed by OZI. In the SU(3) limit the Pomeron couples equally to strange and nonstrange quarks, and a factorizable Pomeron carries no information on strangeness from one vertex to another. This is borne out by the total cross section for ϕN scattering, which has no OZI-suppression factor, and is only lower than $\sigma(KN)$ by the same amount that $\sigma(KN)$ is below $\sigma(\pi N)$.

This small effect is naturally understood as SU(3) breaking in the couplings of the Pomeron to strange and nonstrange quarks, and is not related to the connected and disconnected quark diagrams of the OZI rule. Thus in a multiperipheral process, the f' is emitted by a Pomeron about as easily as any other tensor meson. In the particular case of double Pomeron exchange,[25] one should expect to see f' production comparable to f production. In a Mueller diagram for the central region,[5] one should also expect comparable ϕ and ω production and comparable f and f' production if the Pomeron is approximately an SU(3) singlet as commonly believed.

There is no contradiction in the violation of OZI rule by the Pomeron, since the connected quark diagrams used to describe Reggeon exchanges do not apply to the Pomeron. However,in models where the Pomeron is "built" from other trajectories,[26] there may be some "memory" of quantum numbers propagated a small distance down the multiperipheral chain and a consequent respect for OZI at moderate energies and low multiplicities. This question is still open. It could be tested by looking for the f' in processes where the f is produced by a mechanism which seems to be double Pomeron exchange, or by looking at the ϕ/ω ratio in the central plateau.

Experimental data on ψ photoproduction seem to indicate that the coupling of the ψ to the Pomeron is considerably less than that of ordinary strange and nonstrange mesons. This must be taken into account in estimating production cross sections for new particle production by Pomeron exchange. But this flavor dependence in Pomeron couplings should not be confused with the OZI rule which is determined by the topological character of quark diagrams.

Hadronic production of quarkonium states may have a very different dependence on the spin and parity quantum numbers than electromagnetic production, which favors vector mesons. There are suggestions that the OZI rule holds much better for vector mesons than for pseudoscalars. In QCD, where the rule is broken by annihilation of a quarkonium pair into gluons, three gluons are required to annihilate a vector state, while a pseudoscalar can go into two gluons. There are also experimental arguments which show that OZI violating processes are stronger in the pseudoscalar state than in the vector state. The absence of ideal mixing in the lowest pseudoscalar nonet is evidence for OZI violation, since the interaction which mixes strangeonium and nonstrangeonium effectively[9] violates OZI. More recently there is experimental evidence from radiative decays that the OZI-violating transition between charmonium states and light quark states is stronger in the pseudoscalar state than in the vector state.[27]

In radiative decays of charmonium to a photon and light quarks, there are two possible transitions (a) The photon is emitted by the charmonium system before the transition into light quarks. In this case the photon cannot carry away isospin and the final light quark state must have isospin zero; (b) The photon is emitted by the light quark system after the OZI-violating transition of the charmonium into light quarks.

$$(c\bar{c}; I=0, J^P = 1^-) \rightarrow (c\bar{c}; I=0, J^P = J_f^{P_f}) + \gamma \rightarrow (q\bar{q}; I=0, J^P = J_f^{P_f}) + \gamma \tag{3.1a}$$

$$(c\bar{c}; I=0, J^P = 1^-) \rightarrow (q\bar{q}; I=0, J^P = 1^-) \rightarrow (q\bar{q}; I=I_f, J^P = J_f^{P_f}) + \gamma. \tag{3.1b}$$

In case (a) the photon carries away its angular momentum and parity before the OZI violation, and the violation occurs in a system having the space-spin quantum numbers of the final state. In case (b) the OZI violation occurs in a system having the space-spin quantum numbers of the initial state before the photon carries away angular momentum and parity. The photon can now carry away isospin zero or one, and the final state can be both isoscalar and isovector. Thus the isospin properties of the final state contain information on the space-spin state in which the OZI violation occured.

In the particular case of $\psi \rightarrow P\gamma$ decays, the $\pi^o\gamma$ state can only be produced by the transition (4b) with emission of an isovector photon after the OZI violation has occured in the initial vector state. The $\eta\gamma$ and $\eta'\gamma$ states can be produced by either transition (3.1a) or (3.1b) with isoscalar photons emitted either before or after OZI violation. Experimentally the $\eta\gamma$ and $\eta'\gamma$ decays are much stronger than the $\pi^o\gamma$ decay,[27] by a factor of about 30. So OZI violation in the pseudoscalar state seems to be much stronger than in the vector state.

We can use this information to estimate the production of the pseudoscalar charmonium state η_c in pp collisions. Assuming that the difference between η_c production and J production is only in the OZI violating charmed pair creation, and that the difference between the strength of the violation in vector and pseudoscalar states is given by the argument of radiative decays above, we obtain

$$\frac{\sigma(pp \to \eta_c X)}{\sigma(pp \to JX)} \sim \left| \frac{A(q\bar{q} \to c\bar{c};\, J^P = 0^-)}{A(q\bar{q} \to c\bar{c};\, J^P = 1^-)} \right|^2 \sim \frac{BR(\psi \to \eta'\gamma)}{BR(\psi \to \pi^0\gamma)} \sim 30\,. \tag{3.2}$$

3.3 How to look for new quarkonium states

The charmonium experience shows that e^+e^- colliding beams provide a very effective means for discovering and studying the properties of vector mesons which are directly produced as s-channel resonances, and of other states produced by electromagnetic decays of these vector mesons. Hadronic beams can produce these vector states, but very little information about their properties are obtained in a simple way because of the enormous background. If the SPEAR and DESY results were not available to complement the information obtained from the Brookhaven experiment, we would know very little about the nature of the J particle, and there would be very little evidence that it is indeed a charmonium state.

Hadronic beams might provide additional information on the properties of other states not easily seen with e^+e^-, such as the pseudoscalars. So far the η_c has been seen only in one experiment at DESY and only in the $\gamma\gamma$ decay mode. There is interest in seeing the hadronic decay modes, and any ingenious method for seeing such decay modes with hadronic production would constitute a real breakthrough in x-onium spectroscopy. If the estimate (3.2) of the hadronic production cross section is reasonable, there may be some hope for detecting the η_c via the $\phi\phi$ decay mode after production in pp collisions. The figure of merit for this process can be estimated by comparison with the detection of the J in the e^+e^- decay mode.

$$\frac{F(\eta_c, \phi\phi)}{F(J, ee)} = \frac{\sigma(pp \to \eta_c X)}{\sigma(pp \to JX)} \cdot \frac{BR(\eta_c \to \phi\phi)}{BR(J \to ee)} \cdot \frac{\sigma(pp \to eeX)}{\sigma(pp \to \phi\phi X)} . \tag{3.3}$$

Since the decay $\eta_c \to \phi\phi$ is similar in nature to the decay $J/\psi \to \phi\eta$, we can assume

$$BR(\eta_c \to \phi\phi) \sim 2\, BR(J/\psi \to \eta\phi) \sim (1/35)\, BR(J \to ee) , \tag{3.4}$$

where we have introduced a factor 2 because only about 50% of the η wave function, the ss piece, contributes to the $\eta\phi$ decay mode of the J/ψ, and we have substituted the experimental values for the branching ratios. Combining Eqs. (3.2), (3.3) and (3.4) then gives

$$\frac{\sigma \cdot BR(pp \to \eta_c X \to \phi\phi X)}{\sigma \cdot BR(pp \to JX \to eeX)} \sim (30/35) \sim 1 , \tag{3.5a}$$

$$\frac{F(\eta_c, \phi\phi)}{F(J, ee)} \sim \frac{\sigma(pp \to eeX)}{\sigma(pp \to \phi\phi X)} \quad . \tag{3.5b}$$

Thus if the $\phi\phi$ background is no more than the lepton pair background, it should be just as easy to see $\eta_c \to \phi\phi$ as it is to see $J/\psi \to$ lepton pairs.

Results from the double arm spectrometer experiment at Fermilab[12] showed no $\phi\phi$ events, while the same run observed about 100 events of $J/\psi \to \mu^+\mu^-$. This is still consistent with the result (3.5) of equal signal/noise and comparable signals for the two processes, because the spectrometer had a much lower acceptance for ϕ's than for muons. The absence of any $\phi\phi$ signal confirms that the background is low, and that any further experiments with increased sensitivity might see a small signal without appreciable background. Note that even 3 events for $\phi\phi$ at 2.8 GeV with no background would constitute serious supporting evidence for the existence of the η_c, whereas several hundred events in another decay mode against a background of thousands of events would be ambiguous.

Similar arguments would apply to the detection of higher x-onium pseudoscalars via the $\phi\phi$ decay mode. Note that x-onium pseudoscalars above 6 GeV would also have a $\psi\ \psi$ decay mode which might be detectable in a four lepton final state.

The $\pi A2$ decay mode of the η_c has also been suggested as a possible useful signature.[8] A detailed analysis of the hadronic decays of the η_c has been given by Quigg and Rosner.[28]

IV. STRANGENESS AND SPIN MASS SPLITTINGS

An open problem in hadron spectroscopy is how to describe the regularities in mass splittings occurring in hadron multiplets and supermultiplets. The basic difficulty in deriving any formula for masses from symmetry breaking is the absence of an underlying theory. Simple formulas are obtained by postulating simple transformation properties of the symmetry breaking, e.g. octet splitting for SU(3) which gives the Gell-Mann-Okubo mass formula. But there is no theory to tell whether the formula applies to linear masses, quadratic masses, some exotic power of the mass, the S-matrix, or to "reduced" matrix elements with certain kinematic factors removed. The original folklore suggested linear mass formulas for baryons and quadratic formulas for mesons. These gave good agreement with experiment for SU(3) and SU(6) mass formulas. But the quark model gave results which related baryon mass splittings to meson mass splittings, in particular, the naive assumption that the difference between strange and nonstrange quarks relates meson and baryon splittings as well as mesons and baryons among themselves. Within the meson and baryon supermultiplets these quark model relations are equivalent to SU(6) relations. But between mesons and baryons they give something new, which agrees with experiment when linear masses are used. The situation was summarized at the 1966 Berkeley conference[29] by the "crazy mass formula"

$$K - \pi \overset{Q}{=} K^{*}-\rho \overset{L}{=} \Sigma^{*} - \Delta \overset{L,Q}{=} \Xi - \Sigma , \qquad (4.1)$$

where the L above the equality implies that linear masses should be used and the Q above the equality implies that quadratic masses should be used.

While there are many ways to derive some of these equalities, no credible model includes both the linear and quadratic relations involving the same vector meson mass splitting. But the experimental agreement with the crazy formula is sufficiently impressive to suggest that it cannot be wholly accidental.

The discovery of charm allows a similar formula to be written for the charmed states by simply replacing all strange quarks in (4.1) by charmed quarks. The result is

$$D - \pi \overset{Q}{=} D^{*} -\rho \overset{L}{=} C_1^{*} - \Delta = \qquad , \qquad (4.2)$$

where the last equality is left open since the doubly charmed baryon analogous to the Ξ has not yet been found. This formula also agrees with experiment, as shown in Table 4.1. Thus changing a nonstrange quark in the ρ to a strange or to a charmed quark produces a linear mass shift which is equal to that produced by the corresponding change of a quark in the Δ, while the shift in squared mass is equal to that produced by the corresponding quark change in the pion.

An interesting relation between the spin splittings of the masses of strange and nonstrange baryons was given by Federman, Rubinstein and Talmi[30] in 1966

$$(1/2)(\Sigma + 2\Sigma^{*} - 3\Lambda) = \Delta - N. \qquad (4.3)$$

Experimentally the left and right hand sides of this relation are 307 and 294 MeV, which is rather good agreement. This relation follows from the assumption that the mass differences are due to two-body forces which are spin dependent. The right hand side is just (3/2) the difference between the interaction of two nonstrange quarks in the triplet and singlet spin states when these quarks are bound in a nonstrange baryon. The left hand side is the same difference for a nonstrange quark pair bound in a hyperon (the particular linear combination chosen causes the contribution from the strange quark interaction to cancel out). The experimental agreement indicates that the assumptions of two-body forces and SU(6) spin couplings in the wave functions are good approximations.

Here again the relation can be extended to charm by replacing strange quarks everywhere with charmed quarks.

$$(1/2)(C_1 + 2C_1^{*} - 3C_o) = \Delta - N \qquad (4.4)$$

Since the present experimental information on charmed baryons[22] gives a mass of 2260 for the C_o and a mass of 2500 for a broad peak interpreted to be the unresolved C_1-C_1^* combination, it is convenient to rewrite eq. (4.4) as

$$(C_1 + 2C_1^{*})/3 = C_o + (2/3)(\Delta - N). \qquad (4.5)$$

The left hand side is a weighted average of the C_1 and C_1^* masses, which can be roughly approximated by the value 2500 MeV for the unresolved peak. The left hand side is 2456 MeV, which is in reasonable agreement. So the spin interactions of the ordinary u and d quarks in charmed hadrons are the same as in nucleons and hyperons.

We see that charm really behaves very much like strangeness, and that we don't understand it either!

TABLE 4.1 Experimental Tests of Crazy Mass Formula

a) Strangeness Splittings

	$K-\pi$	$\overset{Q}{=}$	$K^*-\rho$	$\overset{L}{=}$	$\Sigma^*-\Delta$	$\overset{L,Q}{=}$	$\Xi-\Sigma$
ΔM(GeV)	0.35 GeV		0.12		0.15		0.12
ΔM^2 (GeV)^2	0.22		0.20				

b) Charm Splittings

	$D-\pi$	$\overset{Q}{=}$	$D^*-\rho$	$\overset{L}{=}$	$C^*-\Delta$
ΔM(GeV)	1.72		1.23		1.26 (if $M_{C^*}=2.5$)
ΔM^2 (GeV)^2	3.3		3.4		

V. ARE THERE EXOTIC HADRONS?

5.1 Naive Exotics and Saturation

A simple-minded quark model suggests that the quark-antiquark interaction is attractive in all states, because bound states are found as mesons with all the quantum numbers allowed for the $q\bar{q}$ system. Then if two positive pions are brought together, there should be a strong attraction between the quark in one pion and the antiquark in the other to produce a doubly charged bound state with I=2 below 300 MeV. Since no such exotic bound state or resonance has been found the naive model fails and some saturation mechanism is needed to explain the absence of naive exotics around the dipion mass.

The presently accepted colored quark model with forces from exchange of an octet of colored gluons provides a saturation mechanism in which the $q\bar{q}$ and 3q states behave like neutral atoms.[31] Different parts of the bound state wave function attract and repel an external particle and the net force exactly cancels. Thus theory and experiment now agree on the absence of naive exotics. But the possibility exists of higher exotics. Molecular-type exotics in which attraction results from spatial polarization of one hadron by another have been considered, but crude calculations indicate that the force is insufficient to produce binding.[31,32] Rosner[33] has postulated the existence of exotics from the point of view of finite energy sum rules and duality. This approach has been carried further by other theorists and experiments have been suggested in a search for exotics by baryon exchange processes.

So far there is no evidence for exotic mesons with masses below 2 GeV. This has been taken as evidence against the $qq\bar{q}\bar{q}$ configuration for low-lying states. Although $qq\bar{q}\bar{q}$ states without exotic quantum numbers also exist, these were not taken seriously as possible configurations for the known states, because there was no good theoretical reason why such states should be present and their exotic partners should be absent. But now there seems to be evidence that the low-lying 0^{++} nonet is indeed such a $qq\bar{q}\bar{q}$ state,[34] and there are new convincing theoretical reasons why only states with nonexotic quantum numbers are seen.[18]

5.2 Color-Spin (Magnetic) Exotics and the Flavor Antisymmetry Principle

Recently Jaffe[18] has suggested the existence of exotics bound by the "magnetic-type" spin dependent forces arising naturally in the colored-quark-gluon (QCD) models. The prediction rests on much more general grounds than the specific M.I.T. bag model used in Jaffe's original derivation. The essential physical input is that the N-Δ mass difference is much larger than the binding energy of the deuteron:

$$M_\Delta - M_N \gg M_n + M_p - M_d \tag{5.1}$$

where n, p and d denote neutron, proton and deuteron, not quarks, and this equation shows that there are problems of ambiguities in both the pnλ and uds notations.

The physics of eq. (5.1) is that the dominant spin-independent (color charge) forces which bind quarks into hadrons saturate at the $q\bar{q}$ and 3q states and the residual forces between color singlet hadrons is only of the order of 2 MeV like the deuteron binding energy. However, the spin dependent force responsible for the mass difference between the N and Δ is very much larger, of order 300 MeV. Thus if two hadrons are brought very close together so that the quarks in one can feel the interactions of the quarks in the other, there is only a very weak force if the wave functions of the individual hadrons are not changed. However, if the spins of the quarks are recoupled to optimize the spin dependent interactions between the quarks in different hadrons, binding energies of the order of 300 MeV are available and could give rise to bound exotics. In the quark-antiquark system, the ρ-π mass splitting shows that 600 MeV is gained by changing the spins from S=1 to S=0.

Jaffe has simply used the N-Δ and ρ-π mass splittings as input for the strength of the spin dependent interaction and calculated its effect in binding exotic configurations. Only one further ingredient is needed, the color dependence of the interaction. In color singlet $q\bar{q}$ and 3q systems, every $q\bar{q}$ pair is in a color singlet state and every qq pair is in the antisymmetric color triplet state. Exotic configurations, even if they are overall color singlets, can have some $q\bar{q}$ pairs in the color octet state and some qq pairs in the symmetric sextet state. The interactions in these states are not obtainable from observed masses, and are obtained by assuming that the color dependence of the interaction is that obtained from the spin-dependent part of the one-gluon exchange potential in QCD. Evidence supporting this interaction is the agreement with qualitative features of the low-lying hadron spectrum not obtained in any other way, in particular the sign of the N-Δ and Λ-Σ mass splittings.[35] With this form for the interaction, its contribution to the binding of exotic hadron states is easily calculated by the use of algebraic techniques.

One result of the algebraic derivation is simply expressed as the "flavor-antisymmetry principle."[19] <u>The binding force between two quarks of different flavors in the optimum color and spin state is stronger than the binding force between two quarks of the same flavor.</u> Although the forces are assumed to be flavor-independent, their color and spin dependence appears as a flavor dependence because of the generalized Pauli principle. For maximum binding the state should be overall symmetric in color and spin together. Thus if the quarks are in the same orbit, and therefore symmetric in space, they must be flavor antisymmetric. This is seen in the N-Δ example where the $I = 1/2$ state is lower than the $I = 3/2$ state even with isospin independent forces, because the Pauli principle requires the correlation between spin and isospin of (1/2,1/2) and (3/2,3/2) for a color singlet state.

The flavor antisymmetry principle requires the most strongly bound state of a system of quarks and antiquarks to have quarks and antiquarks separately in the most antisymmetric flavor state allowed by the quantum numbers. Thus for example the lowest state of the six quark system has the configuration (uuddss) with no more than two quarks of any one flavor.

The general question of dibaryon bound states and resonances as six quark systems has been considered by Jaffe,[36] with the prediction of a low-lying six quark state as a bound state or resonance of the $\Lambda\Lambda$ system. The exact values of the masses of these states calculated by Jaffe can be questioned because of uncertainties in parameters appearing in the bag model but certain qualitative features are reasonably clear. The spin-dependent force between quarks in the two baryons will be strongest in the $\Lambda\Lambda$ system because of the flavor-antisymmetry principle. The exact values of the masses depend not only on the strength of the spin-dependent interaction, but also on other effects not included in the model calculation and difficult to estimate. However, if these other effects do not depend strongly on flavor, dibaryon bound states or low-lying resonances are most likely to be found in the $\Lambda\Lambda$ system.

It is interesting to note that multiquark binding lies outside the conventional SU(6) classification of hadrons. In the SU(6) symmetry limit the nucleon and the Δ are degenerate and the color-magnetic forces responsible for multiquark binding are absent. The existence of magnetic multiquark exotics requires SU(6) symmetry breaking, and may be related to other SU(6)-breaking effects in addition to the mass differences. One possible effect is the finite neutron charge radius, which vanishes in the SU(6) symmetry limit. Carlitz et al.[37] have suggested that this results from the same spin-dependent interaction which gives rise to the mass splittings and have made a quantitative estimate which agrees with experiment. It is interesting to note that the sign of the neutron charge radius is seen immediately from the flavor antisymmetry principle. In the SU(6) symmetry limit the spatial separation between any quark pair in the neutron is the same as that of any other pair and there is no spatial charge distribution. Breaking SU(6) with the "flavor-antisymmetric" interaction provides a stronger attractive force between quarks of different flavors and distorts the SU(6) wave function to bring the ud pairs in the neutron closer together than the dd pair. Thus the negatively charged d quarks are farther out on the average than the odd u quark which likes to be closer to the differently flavored d quarks, and the charge distribution is negative at large radius and positive at smaller radius.

So far there is no experimental evidence for a strongly bound $\Lambda\Lambda$ state, and there is some evidence against it.[38] Hypernuclei with two Λ's have been observed,[39] and are bound by only about 5 MeV more than the binding of two single Λ's. A $\Lambda\Lambda$ bound state with a much stronger binding energy would be expected to be formed in such hypernuclei. The failure to observe this transition might be explained by selection rules or barrier penetration factors. But any such mechanism presenting formation of a bound state by two Λ's present in the same nucleus for a time equal to the Λ decay lifetime should produce even greater inhibition in any experiment where the two Λ's are produced in a strong interaction collision and are close together for a much shorter time. There may be many-body effects in the hypernucleus which invalidate this argument; e.g. repulsive cores in the Λ-nucleon interaction might prevent the two Λ's from coming too close together in the presence of a finite nucleon density. But except for such effects, the existence of the lightly bound $\Lambda\Lambda$ hypernuclei suggest that strongly bound $\Lambda\Lambda$ states are not easily produced even if they exist.

For the $qq\overline{q}\overline{q}$ system flavor antisymmetry gives two very interesting qualitative predictions.[18,19]

1. The lowest states do not have exotic quantum numbers.
2. The lowest states which have both charm and strangeness include exotics.

These predictions are simply derived by noting that a four body system must have two bodies with the same flavor if there are only three flavors. Since the flavor-antisymmetry principle requires the flavors of the quark pair and of the antiquark pair to be different in the lowest states, the two bodies with the same flavor must be a quark-antiquark pair. The flavor quantum numbers of this pair cancel one another and the quantum numbers of the system are those of the remaining pair and therefore not exotic. Prediction 1 gives a natural explanation for the absence of low-lying states with exotic quantum numbers, while allowing low-lying four-quark states with nonexotic quantum numbers. Jaffe has called such states "crypto-exotic". Prediction 2 follows from the observation that the flavor antisymmetry principle is easily satisfied with exotic quantum numbers when there are four flavors. Thus exotic states with both charm and strangeness may be found in the same mass range as the lowest F and F* mesons with both charm and strangeness.

VI. WHAT ARE STRANGENESS AND BARYON NUMBER?

The new charm degree of freedom[4] provides a new quantum number like electric charge, strangeness and baryon number. But understanding charm is difficult when we still do not understand the old internal degrees of freedom.[40] We have some understanding of the role of electric charge in particle interactions and dynamics even though we do not understand why electric charge is quantized and universal. But our understanding of baryon number and strangeness is much weaker. There is no theory like quantum electrodynamics in which baryon number or strangeness appear as coupling constants defining the strengths of interactions. There is no formula analogous to the Rutherford formula for Coulomb scattering describing the dependence of strong interaction scattering on baryon number and strangeness.

A few phenomenological models and symmetries like the quark model and SU(3) symmetry give rough descriptions of the dependence of total cross sections on baryon number and strangeness. But these descriptions are highly inadequate and the difference between mesons and baryons and between strange and non-strange hadrons are not really understood. Furthermore, many of the models developed work in only one area of hadron physics and are incompatible with models used in other areas. For example, the quark model used in describing hadron strong interactions is not the same as the quark model used in weak interactions.

Consider, for example, the description by conventional models of the difference between pion and kaon wave functions. The quark model says that both are made from a quark-antiquark pair.[14] But weak interaction quarkists explain the ratio of the $\pi \to \mu + \nu$ and $K \to \mu + \nu$ decay requiring the wave functions at the origin to be very different as described by Weisskopf-Van-Royen[41] formula

$$\frac{|\psi_K(0)|^2}{|\psi_\pi(0)|^2} = \frac{M_K}{M_\pi} .$$

Strong interaction quarkists say that the difference between pion and kaon wave functions is measured by the difference between their scattering cross sections on nucleons. These differ by less than 20%. Recent data at high energies show that πp and Kp differential cross sections approach equality with increasing momentum transfer. This suggests equality within 20% of the mean square radii of pion and kaon wave functions and nearly identical short distance behavior, in sharp contrast with the weak quarkist result.

The very precise experimental data[42] now available on pion, kaon and nucleon total cross sections give us some information about the difference between the interactions of strange and nonstrange particles with matter. Careful examination of the data show very clearly that there is a difference between strange and nonstrange particles and that there are puzzles not explained by the quark model. This is strikingly shown in linear combinations of cross sections which have no Regge component and are therefore conventionally assumed to be pure pomeron. The K^+p and pp channels are exotic and have no contribution from the leading Regge exchanges under the common assumption of exchange degeneracy. The following linear combinations of meson-nucleon cross sections are constructed to cancel the contributions of the leading Regge trajectories

$$\sigma(\phi p) = \sigma(K^+p) + \sigma(K^-p) - \sigma(\pi^-p) \tag{6.1a}$$

$$\Delta(\pi K) = \sigma(\pi^-p) - \sigma(K^-p). \tag{6.1b}$$

Figure 6.1 shows these two quantities on the conventional plot of cross section versus P_{lab} on a log scale.

$\sigma(\phi p)$ as defined by Eq. (6.1a) is the quark model expression for $\sigma(\phi p)$; i.e., the cross section for the scattering of a strange quark-antiquark pair on a proton. The very simple energy behavior of this quantity as seen in Fig. 6.1 is striking. It shows a monotonic rise beginning already at 2 GeV/c. That total cross sections rise at high energies was first noticed by Serpukhov data from 20—50 GeV/c, but the older data at lower energies already show this rising behavior in $\sigma(\phi p)$. If anyone suggested something particularly fundamental about this cross section for strange quarks on a nucleon before the Serpukhov data were available and concluded that its rising cross section indicated that all cross sections would eventually rise he would naturally have been disregarded as crazy. But now that the whole picture up to 200 GeV/c is available we may conclude that there is indeed something simpler and more fundamental about the cross sections for strange quarks on a proton target. Understanding this simpler behavior may help us to understand the more complicated energy behavior of the other cross sections.

The quantity $\Delta(\pi K)$ defined by Eq. (6.1b) represents the difference in the scattering of a strange particle and a nonstrange particle on a proton target. In the quark model this is the difference between the scattering of a strange quark and a nonstrange quark on a proton target after the leading Regge contributions have been removed. This difference between strange and nonstrange also has a very simple energy behavior, decreasing constantly and very slowly (less than a factor of 2 over a range P_{lab} of two orders of magnitude). So far there is no good explanation for why strange and nonstrange mesons behave differently in just this way.

Since the two quantities (6.1) have no contribution from the leading Regge trajectories they represent something loosely called the pomeron. However, their energy behaviors are different from one another and also from that of the quantities $\sigma(K^+p)$ and $\sigma(pp)$ which should also be "pure pomeron." However the following linear combinations of $\sigma(K^+p)$ and $\sigma(pp)$ have exactly the same energy behavior as the meson-baryon linear combinations (6.1)

$$\sigma_1(pK) = \frac{3}{2}\,\sigma(K^+p) - \frac{1}{3}\,\sigma(pp) \qquad (6.2a)$$

$$\Delta(MB) = \frac{1}{3}\,\sigma(pp) - \frac{1}{2}\,\sigma(K^+p). \qquad (6.2b)$$

These quantities are also plotted in Fig. 6.1.

The equality of the quantities (6.2) and the corresponding quantities (6.1) suggest that the pomeron, defined as what is left in the total cross sections after the leading Regge contributions are removed by the standard prescription, consists of two components, one rising slowly with energy and the other decreasing slowly. The coefficients in Eq. (6.2) were not picked arbitrarily but were chosen by a particular model. In this model the rising component of the total cross section is assumed to satisfy the standard quark model recipe exactly.

$$\sigma_1(Kp) = \sigma_1(\pi p) = \frac{2}{3}\,\sigma_1(pp) = \frac{2}{3}\,\sigma_1(Yp) = \frac{2}{3}\,\sigma_1(\Xi p), \qquad (6.3a)$$

where Y denotes a Λ or Σ hyperon. The falling component has been assumed to satisfy the following relation

$$\sigma_2(Kp) = \frac{1}{2}\,\sigma_2(\pi p) = \frac{2}{9}\,\sigma_2(pp) = \frac{1}{3}\,\sigma_2(Yp) = \frac{2}{3}\,\sigma_2(\Xi p). \quad (6.3b)$$

This particular behavior is suggested by a model in which the correction to a simple quark-counting recipe comes from a double exchange diagram involving a pomeron and an f coupled to the incident particle.[43]

We thus see unresolved problems in the total cross-section data associated with the questions of what is the difference between strange and nonstrange particles and what is the nature of the pomeron. Note that Eq. (6.1b) defines the difference between the scattering of a nonstrange quark and a strange quark while Eq. (6.2b) can be interpreted as the difference between the scattering of a quark in a baryon and a quark in a meson. The fact that the strange-nonstrange difference and the meson-baryon difference are equal and have the same energy behavior over such a wide range is a puzzle which may be explained by pomeron-f double exchange but may also indicate something deeper.

A very good fit to the experimental total cross section data up to 200 GeV/c has been obtained with the two components (6.2) and (6.3) parametrized by simple power behavior. This gives a formula with five parameters which were adjusted to fit the data,[43]

$$\sigma_{tot}(Hp) = C_1\sigma_1(Hp) + C_2\sigma_2(Hp) + C_R\sigma_R(Hp) \quad (6.4)$$

where $C_1 = 6.5$ mb., $C_2 = 2.2$ mb., $C_R = 1.75$ mb.,

$$\sigma_1(Hp) = N_q^H(P_{lab}/20)^{\epsilon} \quad (6.5a)$$

$$\sigma_2(Hp) = N_q^H N_{ns}^H(P_{lab}/20)^{-\delta} \quad (6.5b)$$

$$\sigma_R(Hp) = (N_{\bar{n}}^H + 2N_{\bar{p}}^H)(P_{lab}/20)^{-\frac{1}{2}} \quad , \quad (6.5c)$$

N_q^H is the total number of quarks and antiquarks in hadron H (N_q^H = 2 for mesons and 3 for baryons), N_{ns}^H is the total number of non-strange quarks and antiquarks in hadron H and $N_{\bar{n}}^H$ and $N_{\bar{p}}^H$ are the total number of $\bar{n}$ and $\bar{p}$ antiquarks in hadron H, $\epsilon = 0.13$ and $\delta = 0.2$.

The dependence of the individual terms in Eqs. (6.5a) and (6.5b) on the quantum numbers of H are determined by the model and discussed in ref. 43. The explicit form for the energy dependence is chosen to minimize the number of free parameters. Thus power behavior is chosen rather than logarithmic for the two components of the Pomeron, because two parameters are sufficient to describe a power and at least three are needed to describe logarithmic behavior. The Regge term was chosen to minimize the number of free parameters by assuming exact duality and exchange degeneracy for the leading trajectories with the conventional intercept of one-half.

The extension of the formula(6.4)to the real part of the amplitude is a straightforward application of analyticity and crossing, which is particularly simple for terms with power behavior[44] and gives the following expression for the ratio of the real to imaginary parts of the Hp amplitude

$$\rho(Hp) = \frac{C_1\sigma_1(Hp)\tan(\pi\epsilon/2) - C_2\sigma_2(Hp)\tan(\pi\delta/2) - C_R\sigma_R(\overline{H}p)}{\sigma_{tot}(Hp)} . \quad (6.6)$$

The total proton-proton cross section and the real part of the forward scattering amplitude have been recently measured[45] at ISR. Table 6.1 show that the new data in the energy range equivalent to P_{lab} = 500 to 2000 GeV/c are in excellent agreement with predictions from the five parameter formula (6.4) - (6.6) with no adjustment of the values of these parameters from already published values fixed by fits to data below 200 GeV/c. Table 6.1 also lists predictions for higher energies and shows remarkable agreement with results from Cosmic Ray experiments[46] up to $P_{lab} = 40{,}000$ GeV/c. Whether these agreements confirm the validity of the oversimplified two-component model is unclear. However, the formula can certainly be used as a simple parameterization of the data and a guide to the physics of further experiments. The ISR group fit their data with a seven parameter formula.[45]

The good fits obtained to very high energy data indicate that these rather crude approximations are nevertheless adequate up to these energies. As long as this reasonable fit continues models containing more detailed assumptions will not be easily tested by the available data. For example, as long as a good fit is obtained with power behavior for the first component the necessity for logarithmic terms will be difficult to demonstrate since a considerably better fit is required to justify the use of additional parameters. The same is true for more detailed or realistic descriptions of the Regge component, since breaking exchange degeneracy or choosing a value different from one-half for the intercept necessarily requires more parameters. However, as soon as data appear which fail to fit this formula, the underlying assumptions are so simple that the physics of the disagreement should be readily apparent. The nature of the disagreement might suggest, for example, that the rise of the cross sections is logarithmic rather than a power, that exchange degeneracy is breaking down, or that the Regge intercept is not one-half. There may also be a breakdown of the two-component pomeron picture if the dependence on the quantum numbers of hadron H no longer satisfies the simple relations of the model. Thus, regardless of the validity of the two component pomeron description, the formula (6.4) should be a valuable guide to the analysis of data on high energy total cross sections and real parts of scattering amplitudes.

TABLE 6.1 Theoretical Predictions and experimental data for $\sigma_{tot}(pp)$ and $\rho(pp)$

P_{lab} (GeV/c)	$\sqrt{s}$ (GeV)	$\sigma_{tot}(p\bar{p})$ Theory (mb)	$\sigma_{tot}(pp)$ Theory (mb)	$\sigma_{tot}(pp)$ Experiment (mb)	$\rho(pp)$ Theory	$\rho(pp)$ Experiment
498	30.6	41.8	40.0	40.1 ± 0.4	.025	.042 ± .011
1064	44.7	42.8	41.6	41.7 ± 0.4	.064	.062 ± .011
1491	52.9	43.5	42.5	42.4 ± 0.4	.079	.078 ± .010
2075	62.4	44.3	43.5	43.1 ± 0.4	.092	.095 ± .011
4600	92.9	46.8	46.2	47.0 ± 0.8	.118	
10000	137.	49.8	49.5	50.6 ± 1.2	.138	
25000	217.	54.3	54.0	53.8 ± 2.2	.156	
40000	274.	56.9	56.7	55.0 ± 3.0	.163	
100000	433.	62.7	62.6		.174	

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FIGURE CAPTION

Fig. 1: Plots of Eqs. (6.1) and (6.2).

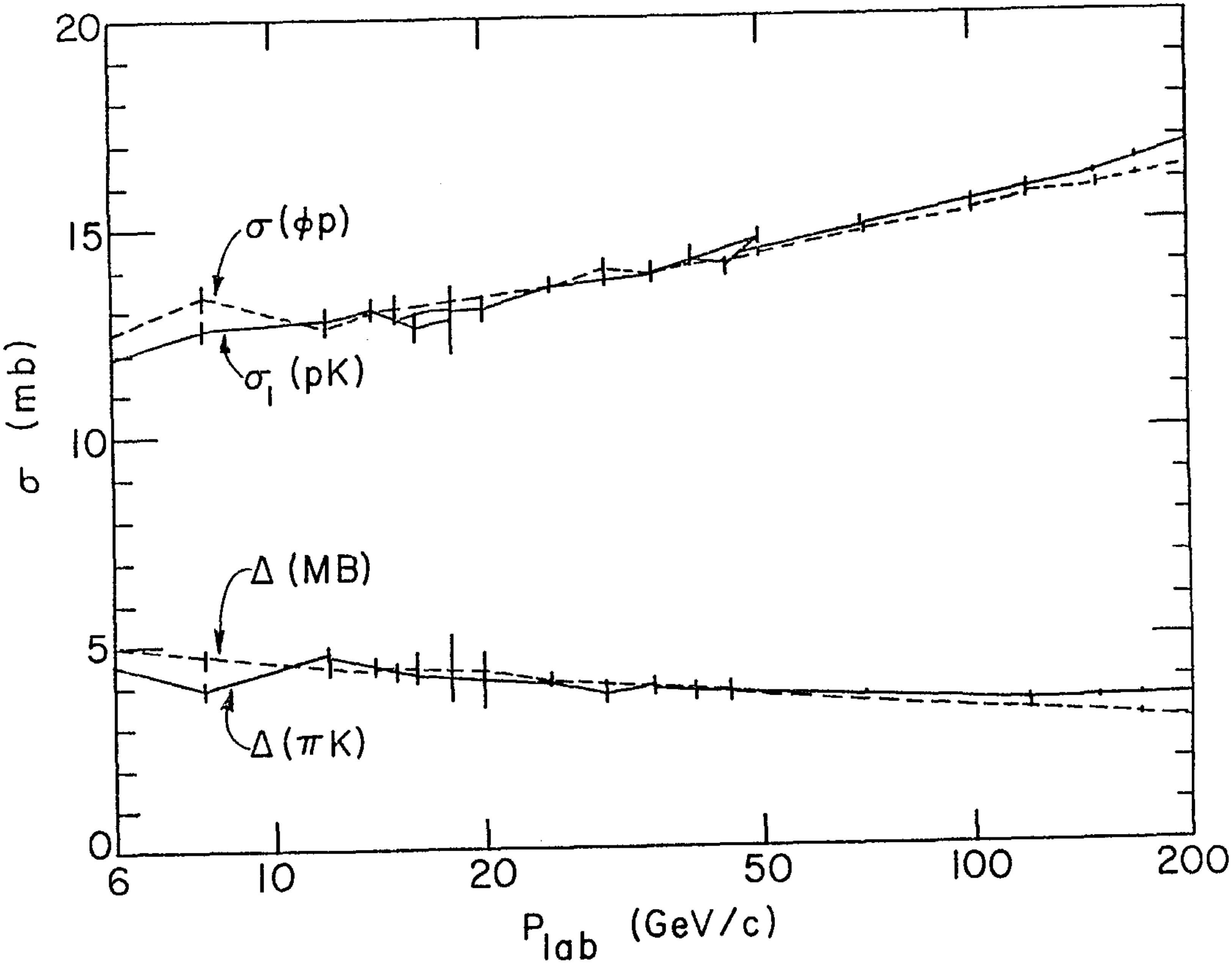

Fig. 1