Diquarks

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It is becoming increasingly clear that the concept of a diquark (a two-quark system) is important for understanding hadron structure and high-energy particle reactions. According to our present knowledge of quantum chromodynamics (QCD), diquark correlations arise in part from spin-dependent interactions between two quarks, from quark radial or orbital excitations, and from quark mass differences. Diquark substructures affect the static properties of baryons and the mechanisms of baryon decay. Diquarks also play a role in hadron production in hadron-initiated reactions, deep-inelastic lepton scattering by hadrons, and in e^+e^- reactions. Diquarks are important in the formation and properties of baryonium and mesonlike semistable states. Many spin effects observed in high-energy exclusive reactions pose severe problems for the pure quark picture of baryons and might be explained by the introduction of diquarks as hadronic constituents. There is considerable controversy, not about the existence of diquarks in hadrons, but about their properties and their effects. In this work a broad selection of the main ideas about diquarks is reviewed.

CONTENTS

I. INTRODUCTION

A. What are diquarks?

There exists today a leading candidate theory of the strong interaction, namely, quantum chromodynamics (QCD). However, in common with other "realistic" continuum field theories, QCD can at present be evaluated systematically only in a perturbation expansion. Unfortunately, such an expansion is inadequate for treating many interesting problems involving hadrons, because the running coupling constant varies from weak to strong. Thus there is no small expansion parameter for these problems.

Because of the inadequacy of QCD perturbation theory, physicists have resorted to inventing models to describe the strong interactions. In fact, the use of quark

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and diquark models to calculate properties of hadrons dates from shortly after the original papers on quarks by Gell-Mann (1964) and Zweig (1964), whereas QCD dates from several years later (Nambu, 1966; Fritzsch and Gell-Mann, 1972; Fritzsch *et al.*, 1973; Gross and Wilczek, 1973; Weinberg, 1973).

Creutz (1980) pioneered evaluating QCD nonperturbatively by numerical methods, using a lattice version of the theory (Wilson, 1974; Kogut and Susskind, 1975). Aside from questions of principle, lattice calculations suffer because an enormous amount of computer time is necessary to achieve very modest results. Thus, at present, calculations with lattice gauge theory are not a satisfactory substitute for calculations with phenomenological models.

Among the useful phenomenological ideas is the notion of a diquark. Gell-Mann (1964) first mentioned the possibility of diquarks in his original paper on quarks. Later, Ida and Kobayashi (1966) and Lichtenberg and Tassie (1967) introduced diquarks in order to describe a baryon as a composite state of two particles, a quark and diquark. Around the same time, states having some or all of the quantum numbers of diquarks were introduced in certain group-theoretical schemes by Bose (1966), Bose and Sudarshan (1967), and Miyazawa (1966, 1968).

It is too simple to regard a diquark as a point particle with the quantum numbers of two quarks, although this oversimplified picture often leads to predictions in qualitative agreement with experiment. In this review, *any two-quark system is a diquark.*

Included in this definition of a diquark are more specialized definitions, including the following two important ones: (1) A diquark is any system of two quarks considered collectively. (2) A diquark is a two-quark correlation in a hadron containing more than two quarks. There is some overlap between these definitions.

As an example of definition (1), if, in a high-energy collision, a quark is knocked out of a baryon, a diquark remains. (Of course, subsequently, both the ejected quark and the spectator diquark will convert to hadrons.) As an example of definition (2), diquark correlations occur in the nucleon because of spin-dependent forces arising in QCD.

A diquark has the quantum numbers of a two-quark system. In its ground state a diquark has positive parity and may be an axial vector (spin 1) or a scalar (spin 0). An axial-vector diquark is often called a vector for short.

Quantum chromodynamics leads us to view a baryon as being made of three valence quarks plus a sea of gluons and quark-antiquark pairs. However, it turns out that it is a good approximation to calculate the static properties of a baryon by treating it as a bound state of just three constituent valence quarks or of a constituent quark and a diquark. A constituent quark consists of a current quark plus the gluons and quark-antiquark pairs which are dragged along as the current quark moves. Although, according to QCD, a current quark is pointlike, a constituent quark has a size greater than zero, as we shall discuss in Sec. II. Whether it is useful to use a model with constituent or current quarks depends on the momentum transfer at which the quark is probed.

Likewise, a diquark has a size which is greater than zero. Pire (1993) distinguishes between a large diquark, which he calls a "fat" diquark, and a small diquark, which he calls a "thin" diquark. A large diquark, sometimes called a constituent diquark, is composed of two constituent quarks. A small diquark is not a point like a current quark, but is composed of two nearby current quarks and perhaps also some glue and pairs. Nevertheless, a small diquark is sometimes called a current diquark in the literature, and we occasionally use this term. A semantic difficulty is that the word "constituent" is used in two ways in the literature and in this review: first, as in "constituent quark" or "constituent diquark," and second, as in "constituent of a hadron." Thus, for example, in the parton model, a current quark is an elementary constituent of a hadron.

The size of even large (constituent) diquarks is often neglected, just as the size of constituent quarks is neglected in many calculations. For example, in a quark-diquark potential model for describing baryons, one can take the potential to depend on the distance between the quark and diquark, which actually means between the center of mass of the quark and the center of mass of the diquark. Often, good results are obtained in the point approximation to quarks and diquarks, but this should not lead the reader to assume that we believe that either constituent quarks or any kind of diquarks, even so-called current diquarks, are really pointlike.

The mass of a constituent quark is larger than the mass of a current quark, perhaps by around 300 MeV. A complication is that the mass of a constituent quark is probably not a constant, but depends on the hadron in which the quark is bound. Often, however, authors take the constituent mass of a quark of given flavor to be a fixed constant, so as to avoid introducing too many free parameters into their models. Of course, not all authors use the same quark constituent masses.

Likewise, the mass of a large diquark should be larger than the sum of the masses of the current quarks it contains and should also depend on the environment in which the diquark finds itself. A small (or current) diquark may have a somewhat smaller mass, but in most models the mass is still taken to be larger than the sum of the current masses of the two quarks it contains. In addition, somewhat different mass estimates are attached to scalar and to vector diquarks. Again, to reduce the number of parameters, most authors take the mass of a diquark of given color, flavor, and spin to be a constant, but different authors use different masses.

No one is surprised that a constituent quark has different properties from a current quark, and neither should one be surprised that a large diquark is different from a small one. The constituent-quark and large-diquark pictures are good for processes at small momentum transfer. At large momentum transfer, a constituent quark is resolved into a current quark plus gluons and quark-antiquark pairs of the sea. At intermediate momentum transfer a large diquark is resolved into a small diquark plus pairs and gluons. At sufficiently large momentum transfer, a small diquark is presumably resolved into two current quarks plus additional pairs and gluons of the sea. However, the small-diquark picture turns out to be occasionally useful at even quite large momentum transfer. In this review we do not always specify whether we mean a large or small diquark, and hope that our meaning will be reasonably clear from context.

An examination of the effect of spin-dependent forces between quarks in a nucleon indicates that there are only modest diquark correlations and that a large diquark is, on the average, not very much smaller than the nucleon itself. On the other hand, in some models of deep-inelastic scattering and other high-energy processes, the best agreement with experiment is obtained if a diquark is quite small compared to a nucleon. In these cases, what may be happening is that a small diquark is selected by the particular collision process being considered (Pire, 1993). For those cases in which two current quarks are relatively far apart, the process may not occur.

A two-quark correlation in a baryon can be treated in a three-quark formalism with explicit correlations in the wave function or in a formalism with only quark and diquark degrees of freedom. A diquark *model,* in contrast to a diquark *correlation,* is a model in which two quarks are approximated as a single particle, which may have a size larger than zero and have a form factor. The possibility of excited diquark states is usually neglected. Although in this review we sometimes treat diquark correlations in three-quark formalisms, our emphasis is on diquark models.

In almost any dynamical scheme, the heavier any two quarks are, the smaller will be their mean separation. For this reason, a baryon composed of two heavy quarks and a light quark has strong diquark correlations. If one quark is radially or orbitally excited, then again diquark correlations are present. In the case of equal quark masses, the two unexcited quarks are closer to each other on the average than they are to the excited quark. Even in the ground state of some baryons, diquark correlations occur because, as we have already remarked, of the existence of spin-dependent forces. In a nucleon, for example, these forces cause the mean distance between a u and a d quark to be smaller than the mean distance between two like quarks.

Spin-dependent forces do not require that diquark correlations occur in the Δ or Ω baryons, which have spin $\frac{3}{2}$. Diquark correlations in these baryons exist in a few *ad hoc* models, for example, in a model of Mitra (1968) with extremely short-range quark-quark forces, and in a model of Halse (1991), in which quarks in a baryon are arranged (on the average) in geometric patterns. However, it is not clear whether either of these models has anything to do with QCD. Whether any diquark correlations exist in the Δ and the Ω is an open question.

As we have remarked, when viewed in different contexts diquarks have apparently different properties. We have already mentioned the possibility that a large diquark is different from a small diquark. Furthermore, all diquark models are oversimplified. Nevertheless, diquark models are useful as a class because they take into account some of the nonperturbative aspects of QCD and explain certain phenomena that otherwise are much more difficult to understand.

At this point we remark that we have chosen to collaborate on this review because we bring together a wide spectrum of experience and viewpoints on diquarks. For this reason we believe we can give a fair discussion of many different aspects of the subject. The other side of the coin is that sometimes we do not agree among ourselves about the nature of diquarks. In those cases where we discuss different approaches which may be inconsistent with each other without stating which approach we favor and why, the reader can assume that we do not all unconditionally agree on a unified statement. Diquark research is being pursued actively at present, and we hope that as a result of future work, the properties of diquarks will become clearer.

Although we have chosen in this introductory section to illustrate the idea of diquarks mainly with baryons, diquarks have been considered in a much wider context, including the description of exotic hadrons and calculations of hadron production by hadrons and leptons. These topics and others related to diquarks are treated in this review.

Definitions of diquarks that differ from ours have appeared in the literature. For example, in some grand unified theories such as SU(5) (Georgi and Glashow, 1974), gauge bosons, called diquarks, of mass perhaps 10^{15} GeV exist and are emitted when a quark transforms into an antiquark. As another example, topological diquarks exist in a model of Chew and Poénaru (1984). As a third example, there are diquarks in a three-triplet model with unconfined color (Nambu and Han, 1976). We shall not consider in detail any diquarks except two-quark systems and single-particle approximations to such systems.

Well over 500 papers have been written on diquarks in the last 25 years, and many additional papers have appeared in which diquarks have not been mentioned explicitly, but in which some effects of two-quark correlations have been taken into account. We do not have the space to summarize all the work described in the papers that treat diquarks explicitly or implicitly, or even to quote a large fraction of them. Szczekowski (1989) has written a review on diquarks in which additional references may be found. Skytt and Fredriksson (1991) have compiled a large list of references to papers on diquarks, with a letter key to their subject matter. This list is updated periodically. References can also be found in the proceedings of three conferences emphasizing diquarks

(Anselmino and Predazzi, 1989, 1993; Goeke *et al.*, 1993). This activity attests to the vitality of the subject. In order for this review to be reasonably self-contained, we have included updated versions of some topics previously reviewed (Szczekowski, 1989).

B. Plan of the review

In the following sections, we discuss effects of diquarks on various aspects of physics involving hadrons. In Sec. II we discuss static properties of baryons both in explicit diquark models and, to a lesser extent, in three-quark models in which diquark correlations are taken into account either explicitly or implicitly. Among the properties of baryons we discuss are mass spectra, decay rates, magnetic moments, the ratio of the axial-vector coupling to vector coupling G_A/G_V, and the charge radius of the neutron.

In Sec. III we discuss models of exotic mesons, emphasizing those models in which an exotic meson is a bound state of a diquark and antidiquark. We treat models of baryonium states and models of exotic states that contain heavy quarks. We also include models in which diquark intermediate states contribute to the decay of ordinary mesons.

In Sec. IV we start discussing the possible dynamical role of diquarks in many processes involving hadrons. We begin by considering lepton-nucleon deep-inelastic scattering and show that diquarks as nucleon constituents induce scaling violations in the nucleon structure functions which may be compatible with experimental information. We compare quark-diquark models of the nucleon with deep-inelastic-scattering data to obtain information about some of the diquark physical properties, like size, mass, and relative scalar-vector abundance. We also discuss diquark momentum distributions, form factors, and fragmentation.

We treat deep-inelastic scattering on nuclei in Sec. V, together with the role of diquarks and more complex quark clusters in dibaryons and heavier nuclei and in quark-gluon plasmas. We consider possible effects of diquarks on the structure functions of bound nucleons. In the dibaryon case we consider six-quark systems, which, unlike the situation in the deuteron, are not clustered into two color singlets of three quarks each, but rather are composed of three diquarks or larger colored clusters. In the case of more complex nuclei, we discuss a model in which diquarks contribute to pairing forces in nuclei. We also discuss the question of whether diquarks can exist in a quark-gluon plasma, and, if so, what their signatures might be.

Diquarks also may play an important role in hadron production in lepton-lepton and hadron-hadron collisions, where the abundance of final baryons produced can be explained by the fragmentation of diquarks rather than of quarks. We treat hadron production in e^+e^- annihilation in Sec. VI and in hadronic inclusive processes in Sec. VII. Both the soft (at small transverse momentum) and hard (at large transverse momentum) production of mesons and baryons in hadron-hadron interactions can be understood in the framework of quark-diquark models of the baryons. We briefly discuss many of these cases.

Most modern Monte Carlo programs dedicated to describing multiple production use the notion of diquarks as a handy tool to avoid a lot of formal and substantial complications. These are not diquark models in a strict sense. In the following we discuss briefly one such model (the Lund model), but others exist (such as JETSET), which we do not discuss for brevity.

We can describe hadronic exclusive reactions at high energies and large angles in terms of interactions among the hadron constituent quarks and gluons, including diquarks. Although we are still far from being able to calculate these processes in a satisfactory way, in some simple cases we can approximately compute cross sections and other observables. The introduction of diquarks may simplify the treatment of such processes in that it reduces the number of active constituents (and, therefore, the number of degrees of freedom) and helps in solving some spin problems, which appear to be a serious problem for the pure quark scheme. We discuss the role of diquarks in exclusive processes in Sec. VIII.

Quantum chromodynamics is a theory of color. Because a diquark and an antiquark have the same color, a diquark in a baryon should behave in similar fashion to an antiquark in a meson. This situation gives rise to a broken supersymmetry between mesons and baryons. We discuss meson-baryon supersymmetry in Sec. IX.

Finally, in Sec. X we briefly present a summary and our conclusions.

II. PROPERTIES OF BARYONS

A. Models of hadrons

Because QCD perturbation theory is inadequate for treating many problems in hadron physics and lattice calculations are so time consuming, phenomenological quark models have been used extensively. Among them are potential models (Greenberg, 1964; Morpurgo, 1965; Dalitz, 1965), bag models (Chodos *et al.*, 1974; Johnson and Thorn, 1976), and string (Eguchi, 1975) or flux-tube models (Isgur and Paton, 1983).

In the potential-model approach, one regards constituent quarks as being bound in a confining potential, which may be "motivated" by both perturbative and lattice QCD. Two papers (Godfrey and Isgur, 1985; Capstick and Isgur, 1986) contain extensive references to the literature. The nonrelativistic three-quark treatment of baryons, including the quark-diquark approximation, has recently been reviewed by Richard (1992).

According to QCD perturbation theory, the potential between two quarks or a quark-antiquark pair is approximately Coulomb-like at short distances. At large distances, one can use lattice QCD in the valence approximation (Weingarten, 1982) to calculate the static quark-antiquark meson potential (see, e.g., Stack, 1983) and the three-quark baryon potential (Thacker *et al.*, 1988). Although three-body forces exist in baryons at large distances, the potential is approximately a sum of two-body linear potentials (Carlson *et al.*, 1983), as proposed by Isgur and Karl (1977). In the one-gluon approximation, the quark-quark potential in a baryon is equal to half the quark-antiquark potential in a meson, and lattice work shows that this relation is rather good even nonperturbatively. Phenomenologically, the approximation also seems to be good.

A string or flux-tube picture also emerges from lattice gauge calculations. The binding force in a meson arises from a string or tube of color-electric flux between quark and antiquark, and the forces among quarks in a baryon arise from three strings meeting at a vertex. If the string breaks, a quark appears at one broken end and an antiquark at the other. The string picture thus includes a natural mechanism for a meson to decay into two mesons or for a baryon to decay into another baryon plus a meson, but it is not a good approximation at short distances.

Hadrons have also been treated in bag models. An advantage of bag models is that they are fully relativistic. They have several disadvantages, however. First, the bag has a sharp boundary, which is quite different from the linearly rising potential expected from lattice calculations. Second, it is difficult to take center-of-mass motion properly into account. Third, it is not clear how to take into account the degrees of freedom of the bag.

B. Diquark correlations in baryons

Potential, bag, and string models all give rise to diquark correlations in some baryons. We first consider diquark correlations in a baryon containing two heavy quarks QQ ($Q=c$ or b) and one light quark q ($q=u$, d, or s). This is a clean case for diquarks because the dynamics requires the mean separation of the QQ pair to be significantly smaller than the mean distance of the q from the center of mass of the QQ pair. Furthermore, there is no need to antisymmetrize the wave function under the interchange of the coordinates of a heavy and a light quark. It follows that the quark-diquark approximation is a good one for a qQQ baryon. Diquark correlations were demonstrated in this case in a three-quark potential-model calculation by Fleck *et al.* (1988). Savage and Wise (1990) and White and Savage (1991) were fully justified in treating qQQ baryons in the diquark approximation, when using the heavy-quark symmetry of Voloshin and Shifman (1987) and Isgur and Wise (1989). Thus far, however, no baryon containing two heavy quarks has been observed experimentally.

Diquark correlations should also exist in a qqQ baryon with only one heavy quark. Fleck *et al.* (1988) have shown explicitly in their three-quark model that a qqQ baryon is mostly made of a qQ diquark and a light quark q. The fact that the baryon wave function must be antisymmetric under the interchange of the two light quarks (if they have the same flavor) does not destroy the diquark effect.

Next we consider a baryon containing three light quarks, with one quark radially or orbitally excited with respect to the other two (the diquark). Then the mean size of the diquark is small compared to the mean distance between the excited quark and the center of mass (c.m.) of the diquark, as shown by several authors, including Eguchi (1975) using a string model, Johnson and Thorn (1976) using a bag model, and Fleck *et al.* (1988) using a potential model. Martin (1986, 1989) proved semiclassically that, at large angular momentum, a quark-diquark configuration of a baryon has minimum energy. Furthermore, if the potential is linear at large distances, an approximately linear Regge trajectory results. (If several hadrons, each composed of two particles—a quark and a diquark in our case—lie on a linear trajectory, the squares of their masses are proportional to the orbital angular momentum L of their constituents.)

If we use the QCD result that, in lowest-order approximation, a diquark in a baryon interacts like an antiquark in a meson because both belong to the $\bar{3}$ of color SU(3), we obtain that the slopes of baryon Regge trajectories are the *same as the corresponding slopes of meson trajectories.* This striking prediction, which is well confirmed by experiment, is *simply explained by diquark correlations in orbitally excited baryons.* (See Sec. IX for further discussion of this point.)

We now turn to the case of a ground-state spin-$\frac{1}{2}$ baryon composed of three light quarks. Then the spin-dependent color-hyperfine (colormagnetic) interaction of QCD (De Rújula *et al.*, 1975), when treated beyond lowest-order perturbation theory, leads to diquark correlations. The color-hyperfine interaction is attractive in a spin-0 state of two quarks and repulsive in a spin-1 state. In the nucleon, the two like quarks have spin 1, whereas the two unlike quarks (u and d) are in a linear combination of spin-1 and spin-0 states. The attractive spin-dependent force in the spin-0 state makes the mean distance between u and d smaller than the mean distance between two u quarks (proton) or two d quarks (neutron). On the other hand, in the Δ, each pair of quarks is in a spin-1 state, and so we do not expect diquark correlations.

According to the three-quark potential model of Fleck *et al.* (1988), diquark correlations in the proton are small. The work of Fleck *et al.* indicates that a scalar diquark in a nucleon is a rather large object, almost as large as the nucleon itself. Ram and Kriss (1987) and Gromes (1988), using somewhat different potential models, also conclude that constituent diquarks are comparable in size to hadrons. On the other hand, Narodetskii *et al.* (1992), using a potential model that emphasizes the color-hyperfine interaction, obtain a smaller constituent diquark.

A problem for constituent models of the nucleon with a small scalar diquark is that a scalar (ud) diquark ought to be bigger than the pion, given that the attractive spin force is weaker in a qq system than in a $q\bar{q}$ system. The solution to this problem could be the one suggested by, e.g., Betman and Laperashvili (1985), namely, that the measured (big) pion radius reflects the size of the constituent $q\bar{q}$ pair including its cloud of gluons and virtual $q\bar{q}$ pairs. On the other hand, the scalar (ud) radius appearing in some models of the nucleon is half the root-mean-square distance between the centers of the u and d quark clouds, and might be considerably smaller than the measured pion radius. Such a picture is in accordance with our discussion (in the Introduction and in Sec. II.C) of the sizes of current and constituent quarks.

C. Baryon static properties

We argued in the last section that, according to our present understanding of hadron physics, diquark correlations must be present in many baryons. Whether this fact justifies treating baryons with only quark and diquark degrees of freedom, rather than with the degrees of freedom of three quarks (or whether one should use a combination of the two pictures), is another question, and one which is more difficult to answer. However, we believe that an important test of a phenomenological model is how well its predictions agree with experiment. Quark-diquark models of baryons frequently pass this test.

One weakness in some calculations of baryon properties with constituent diquark models is the neglect of the Pauli principle. By this we mean that some authors neglect to antisymmetrize the wave function of a multiquark hadron with respect to two identical quarks, one of which is in the diquark while the other is outside it. Of course, nobody grossly neglects the Pauli principle by giving a diquark quantum numbers that are forbidden to a two-quark state.

In the constituent quark-diquark model, if one ignores the Pauli principle between quarks inside and outside the diquark, one predicts the existence of low-lying baryon states, such as a flavor octet of spin 3/2, which are not observed experimentally. Clearly, such an approximation is a bad one. (On the other hand, in parton models with current quarks, with or without small diquarks, neglecting the Pauli principle seems to be a good approximation.)

Because antisymmetrization complicates the diquark approximation, various authors have used other schemes that include the main effects of antisymmetrization. It was pointed out in early work (Lichtenberg, 1969) that one can take into account effects of the Pauli principle by including an exchange term in the quark-diquark potential. Ono (1972, 1973) introduces forces depending on both flavor and spin to break the unwanted degeneracy. Goldstein and Maharana (1980) calculate the baryon spectrum using a QCD-motivated quark-diquark potential model. They emphasize that the effect of the Pauli principle is to require a quark-exchange term in the quark-diquark potential. Goldstein (1989) gives a good discussion of baryon states in the model.

Buck *et al.* (1992) also use quark-exchange terms in treating baryons as bound states of quarks and diquarks in the model of Nambu–Jona-Lasinio (1961a, 1961b). We do not have the space to discuss work in the framework of the Nambu–Jona-Lasinio model and refer interested readers to the paper of Buck *et al.* (1992) for details and additional references.

If one calculates baryon properties with diquark correlations in a three-quark formalism, it is easier to include the effects of the Pauli principle exactly. Many papers have appeared that include diquark correlations while fully respecting the Pauli principle. A paper by Greenberg and Resnikoff (1967) is an early example.

In both the three-quark and quark-diquark models, baryons can be classified as multiplets of broken $SU(6)\times O(3)$. A multiplet is denoted by (N,L^P), where N is the SU(6) multiplicity, L is the total orbital angular momentum, and P is the parity. A light quark belongs to a six-dimensional multiplet of flavor-spin SU(6), while a light diquark belongs to a 21-dimensional multiplet. Because the diquark belongs to a **21** in its ground state, a baryon can belong only to a **56** or **70**. If exchange forces are included, the lowest-energy baryons containing only light quarks belong to the multiplets $(56,0^+)$, $(70,1^-)$, and $(56,2^+)$. Remarkably, these are the multiplets for which there is best experimental evidence. Three-quark models lead to a richer spectrum of states, containing all the multiplets of quark-diquark models, and, in addition, such multiplets as $(70,1^+)$ and $(20,L^P)$, which are forbidden in quark-diquark models if the diquark is unexcited. We believe that the principal lesson to be learned is that the lowest-lying baryon excited states have configurations that correspond to the excitation of a single quark, and the remaining two quarks can be well approximated by a diquark.

In the SU(6) scheme with no diquark correlations, the ground-state baryons belong to the $(56,0^+)$ multiplet. However, the color-hyperfine interaction contributes to the breaking of SU(6) and leads to a mixing of other configurations, including the $(70,0^+)$. Anyone who includes $(70,0^+)$ mixing in the baryon ground state is in effect including diquark correlations. Several authors, for example, Isgur (1977), Carlitz *et al.* (1977), and Le Yaouanc *et al.* (1977), have emphasized the effects of such mixing on baryon properties without mentioning diquarks explicitly. Earlier, Dalitz (1965) and Faiman and Hendry (1968), although treating baryon ground state as belonging to a pure $(56,0^+)$, included mixing in excited states. The above authors present strong evidence that including mixing, i.e., diquark effects, is necessary to get agreement with experiment.

We now briefly consider string and bag models. In the string model of Eguchi (1975), baryons with $L > 0$ have a quark-diquark structure and lie on linear Regge trajectories. According to Eguchi, exchange terms in the potential energy arise "naturally" in his model. A string model is also considered by Cutkosky and Hendrick (1977). In their model, a linear configuration of the constituents is favored, "with two quarks relatively close together." The diquark radius is a parameter, and the amount of exchange forces depends on that parameter. Only for rather large exchange forces does the spectrum of Cutkosky and Hendrick resemble the "minimum" quark-diquark spectrum. See also Burden and Tassie (1982) for another string model calculation.

Johnson and Thorn (1976) apply the MIT bag model to hadrons (see also Mulders *et al.*, 1979). For baryon states with large L, an elongated cigar-shaped bag is plausible; that is, such configurations have lower energies than states with the same quantum numbers in a spherical bag. The condition that the outward pressure of the color flux balance the inward pressure of the bag determines the length of the bag as a function of L. A state with two quarks at one end of the bag and one quark at the other has the lowest energy for a given L. Furthermore, because the color flux connecting opposite ends of the bag is the same for mesons and baryons, these hadrons lie on Regge trajectories that are asymptotically (for large L) linear and have the same slope (see Sec. IX). The calculated asymptotic slope agrees quite well with experiment.

Misra (1980) and Franklin (1980) use harmonic-oscillator potentials to describe interquark forces in baryons, but use a distorted harmonic-oscillator basis; i.e., the harmonic-oscillator wave function is allowed to have a different size for different pairs of quarks. The best agreement with experiment occurs when the wave function contains diquark correlations. One feature of these papers is that the same constituent-quark wave functions that are used to calculate baryon masses also describe deep-inelastic lepton-nucleon scattering. This idea was carried further by Dziembowski and Franklin (1990; see also Dziembowski, 1993). The calculations show that the mean distance between a u and a d quark in a nucleon is somewhat smaller than the mean distance between two like quarks. Gunion and Soper (1978) reached a similar conclusion about the transverse distance by looking at structure functions in deep-inelastic electron-proton scattering.

Capstick and Isgur (1986) calculate baryon masses in a potential model in which spin-dependent interactions are included in the potential (with a cutoff) rather than evaluated as a perturbation. In such a calculation, diquark correlations exist (except in spin-$\frac{3}{2}$ states with equal-mass quarks) because the mean distance between two quarks is affected by their spin-dependent interactions. The improved agreement with the data listed by the Particle Data Group (Hikasa *et al.*, 1992) relative to a perturbative treatment of the spin-dependent terms is evidence for diquark correlations in the wave functions.

The properties of baryons can be calculated in a parameter-free way in a quark-diquark model (Lichtenberg *et al.*, 1982, 1983). One takes the static quark-quark potential to be equal to half the quark-antiquark potential, and the quark-diquark potential to be equal to the quark-antiquark potential. Then one calculates the properties of the diquark in terms of parameters determined from fits to meson data. Next one calculates the properties of baryons as bound states of a quark and diquark, again with no free parameters. In that way, one obtains the properties of baryons by solving two two-body problems rather than one three-body problem.

One interesting aspect of the above calculation, which seems generic to calculations in both three-quark and quark-diquark models, is that the nucleon radius turns out to be too small (about 0.5 fm rather than the experimental value of 0.8 fm). (If the interaction is changed so that the proton size is right, the spacing between excited energy levels becomes too small.) We can understand this result if the constituent quarks themselves have a size of order 0.5 fm. Of course, one can invoke a cloud of pions or quark-antiquark pairs to enlarge the nucleon, but such an explanation is outside the constituent-quark model.

In the above treatment, the wave function of a quark and diquark in a baryon is similar to the wave function of a quark and antiquark in a meson. Gottlieb (1985) has compared meson and baryon wave functions calculated on a lattice in the valence approximation. In the baryon case, he makes a diquark approximation by restricting two quarks to be on the same lattice site. His calculated meson and baryon wave functions are quite similar, a result which lends support to potential models and also to the idea of an approximate supersymmetry between mesons and baryons (Sec. IX).

Quite a different approach to diquark correlations in baryons has been taken by Cahill and collaborators (Burden *et al.*, 1989; Cahill, 1989; Cahill *et al.*, 1989; Praschifka *et al.*, 1989). These authors use a functional-integral approach to transform QCD from a form involving quark and gluon fields into another form involving bilocal meson and diquark fields. The mesons and diquarks have form factors and so are not point particles. Burden *et al.* (1989) have applied this formalism to calculate the structure and mass of the nucleon as an effective bound state of a quark and diquark. Somewhat related work has been done in two-dimensional QCD by Ebert and Kaschluhn (1991, 1993). A paper by Cahill (1991) has many references to the relevant literature. See also Cahill's (1992) recent review of progress in obtaining an effective action for hadron physics with the functional-integral method. Because of Cahill's review, we do not need to discuss the method in more detail here.

It is difficult to say how calculations in the functional-integral approach are related to the phenomenological ones we have previously discussed, because the two formalisms are so different. As far as we can see, the two methods (in four dimensions) give results that are consistent with each other.

Let us now consider a model of baryon ground-state wave functions that uses a three-quark picture with diquark correlations. We use the proton to illustrate the model. We build the proton wave function from isospin states, denoted by ϕ_{12}, and spin states, denoted by χ_{12}, which are either symmetric (s) or antisymmetric (a) under the interchange of the first two quarks:

$$\phi_{12}^{s}=(2uud-udu-duu)/\sqrt{6}\ ,\quad \phi_{12}^{a}=(udu-duu)/\sqrt{2}\ , \tag{2.1}$$

$$\chi_{12}^{s}=(2\alpha\alpha\beta-\alpha\beta\alpha-\beta\alpha\alpha)/\sqrt{6}\ ,\quad \chi_{12}^{a}=(\alpha\beta\alpha-\beta\alpha\alpha)/\sqrt{2}\ , \tag{2.2}$$

where α and β denote spin up and spin down, respectively. Using these states, we build two different proton wave functions p and p' which contain diquark correlations and yet satisfy the Pauli principle. These are

$$p=N(f_{12}\phi_{12}^{s}\chi_{12}^{s}+f_{13}\phi_{13}^{s}\chi_{13}^{s}+f_{23}\phi_{23}^{s}\chi_{23}^{s})\ , \tag{2.3}$$

$$p'=N'(g_{12}\phi_{12}^{a}\chi_{12}^{a}+g_{13}\phi_{13}^{a}\chi_{13}^{a}+g_{23}\phi_{23}^{a}\chi_{23}^{a})\ , \tag{2.4}$$

where N and N' are normalization constants, and f_{12} and g_{12} are spatial wave functions which are symmetric under the interchange of the first two quarks. These wave functions, which are related to wave functions written down in the SU(6) formalism by Isgur and Karl (1980), can be used to calculate baryon properties.

The inclusion of diquarks usually has a small influence on baryon masses but can strongly affect other static properties of baryons, such as magnetic moments (Isgur, 1977; Isgur and Karl, 1980), the ratio of the axial vector to vector coupling constants in β decay (Isgur, 1977), and the neutron charge radius (Carlitz *et al.*, 1977; Isgur, 1977; Dziembowski *et al.*, 1981; Isgur *et al.*, 1981; Dziembowski, 1993).

The simple quark model gives good qualitative results for the magnetic moments of baryons, but there are quantitative discrepancies with the data, of up to $0.2\mu_N$ (μ_N denotes nuclear magnetons). See recent reviews by Franklin (1989) and Brekke and Rosner (1988). Some time ago Franklin (1968; see also Franklin *et al.*, 1981) calculated baryon magnetic moments using three-quark wave functions with implicit diquark correlations. With Franklin's wave functions, the ordering of the quarks is significant. If two quarks in a baryon have the same flavor, they are the first two; but if all three quarks have different flavors, the two lightest are the first two. The magnetic moments of spin-$\frac{1}{2}$ baryons are calculated to be

$$\mu_B=\tfrac{1}{3}(2\mu_1+2\mu_2-\mu_3),\quad \mu_B=\mu_3\ . \tag{2.5}$$

The first of these equations applies if the first two quarks have total spin 1, and the second if the first two quarks have total spin 0. These same expressions describe the magnetic moments of spin-$\frac{1}{2}$ ground-state baryons in both the simple symmetric quark model and in Franklin's model with diquark correlations.

However, if we use the wave function of either Eq. (2.3) or (2.4), then our results do depend on the amount of diquark correlation. Let us evaluate the ratio R of the proton-to-neutron magnetic moment ($R=\mu_p/\mu_n$) and the ratio λ of the axial-vector-to-vector coupling constant ($\lambda=|G_A/G_V|$). We assume that quarks have Dirac moments and we neglect the d-u mass difference. The expressions are formally the same whether we use Eq. (2.3) or (2.4) for the proton wave function, provided we use an equation for the neutron obtained by substituting d for u and vice versa, as required by isospin symmetry. We get

$$R=-(4+5\mathcal{J})/(2+4\mathcal{J}),\quad \lambda=(2+3\mathcal{J})/(2+\mathcal{J})\ , \tag{2.6}$$

where $\mathcal{J}$ is an overlap integral satisfying $|\mathcal{J}|\leq 1$. The value $\mathcal{J}=1$ corresponds to no diquark correlations, while $\mathcal{J}=0$ corresponds to a pointlike diquark. Diquark correlations strongly affect the values of R and λ.

We see from Eq. (2.6) that if $\mathcal{J}=1$, we get the SU(6) results

$$R=-\tfrac{3}{2},\quad \lambda=\tfrac{5}{3}\ , \tag{2.7}$$

in comparison with the experimental values

$$R=-1.46,\quad \lambda=1.25\ . \tag{2.8}$$

We therefore have the situation that the SU(6) result is in rather good agreement with experiment for R, but in serious disagreement for λ. Changing $\mathcal{J}$ improves the value of λ, but worsens the value of R. For example, if we set $\mathcal{J}=0.286$, as required to fit the experimental value of λ, then we get $R=-1.73$, in serious disagreement with experiment.

Similar expressions for R and λ occur in different formalisms. Some time ago Lipkin (1969) noted that configuration mixing increases the discrepancy between the calculated value of R and experiment. In an SU(6)$\times$O(3) formalism with mixing, Isgur and Karl (1980) obtain

$$R=-\tfrac{3}{2}-\tfrac{1}{2}\tan^2\phi\ , \tag{2.9}$$

where ϕ measures the amount of $(70,0^+)$ mixed into the $(56,0^+)$ ground state. It is seen from this expression that mixing worsens the value of R.

The fact that diquark correlations *by themselves* increase the disagreement between the quark model value of R and the experimental value is not a good argument against diquarks in the nucleon. We should not expect these magnetic-moment calculations to be in good agreement with experiment, because relativistic effects, which are omitted, are important for magnetic moments. For example, certain relativistic effects (correlated quark-antiquark pairs, or mesons, in the wave function) go in the opposite direction from diquark correlations, and so can restore the good agreement with experiment (Lichtenberg and Namgung, 1989). Lipkin (1990) has obtained good agreement with the experimental value of R by adding to the nucleon wave function a component that contains a gluon plus three valence quarks in a color octet combined to make an overall color singlet.

Diquark correlations also lead to a negative neutron charge radius ρ, as can be seen from the following argument (Carlitz *et al.*, 1977; Isgur, 1977; Isgur *et al.*, 1981): Because of the color-hyperfine interaction, the two d quarks, having spin 1, repel each other, and so at least one gets pushed to the periphery of the neutron, while the u quark remains closer to the center. Thus the neutron charge distribution is positive at small distances and negative at large distances, and this leads to a negative root-mean-square charge radius ρ. The value of ρ cannot be given simply in terms of I, but must be calculated with a definite radial wave function. This has been done by Isgur *et al.* (1981), who calculated the neutron charge distribution and found it to be in rather good agreement with experiment. For a similar analysis see Dziembowski *et al.* (1981). Note that other possible mechanisms exist to give the neutron a negative charge radius. For example, the neutron wave function contains an amplitude for a virtual proton near the origin and a π^- further from the origin, and this configuration also leads to a negative charge radius.

We see that diquark correlations in the proton and neutron wave functions have a substantial effect on their static properties. However, relativistic and other effects, which are difficult to take into account or require additional parameters, prevent us from being able to use the measured static properties of the nucleon to make precise statements about the magnitude of diquark correlations.

D. Baryon decays

There have been many more calculations of baryon decays with three-quark models than with quark-diquark models. In the framework of the former, Isgur *et al.* (1978) point out that in a three-quark model certain SU(6)-violating baryon decays occur because the color-hyperfine interaction mixes some $(70,0^+)$ configuration into the $(56,0^+)$ nucleon ground state. The fact that these SU(6)-forbidden decays have been observed experimentally is additional evidence that diquark correlations occur in the nucleon. Although diquark correlations are undoubtedly present, they do not seem to be large. For example, Forsyth and Cutkosky (1981) calculate that substantial diquark clustering will make some baryon decay widths too large, but a small amount of diquark correlation is allowed.

In the remainder of this section, we confine ourselves to quark-diquark models.

Ono (1972, 1973) calculates some baryon decay rates in a quark-diquark model which contains an SU(3) flavor sextet of (axial) vector diquarks but not a flavor triplet of scalar diquarks. Ono obtains good agreement with experiment in a number of instances, but, unfortunately, does not discuss why scalar diquarks are absent in his model. On the other hand, Hayashi *et al.* (1978) uses a quark-diquark scheme with only a scalar diquark, but not a vector, to obtain relations among hyperon nonleptonic decays. Their results for both S-wave and P-wave decays are qualitatively in agreement with experiment. We think it is a better approximation to include scalar diquarks only than vectors only, but it is still better to include both, with a larger scalar than vector amplitude. Shito (1980) improves the model of Hayashi *et al.* by including both scalar and vector diquarks. Goldstein (1989) uses the diquark model to give a brief qualitative discussion of baryon decays with either γ or π emission, but does not give quantitative calculations.

Bediaga *et al.* (1985), using a quark-diquark model, estimate the lifetime of the charmed baryon Λ_c as a function of the lifetime of the pseudoscalar charmed meson D^+. As a preliminary step, the authors (Bediaga *et al.*, 1984, 1985) evaluate the decay constants of pseudoscalar and vector mesons and of scalar and vector diquarks in terms of the pion-decay constant f_π. Using the experimental value of the D^+ as input, the authors obtain a value of the Λ_c lifetime which is in good agreement with the data.

Weak nonleptonic decays of baryons have been calculated with a quark-diquark model by Dosch *et al.* (1989). These authors assume that the flavor-octet baryons consist of a quark and a scalar diquark and get a natural explanation for an enhanced $|\Delta I|=\frac{1}{2}$ amplitude, in agreement with experiment. Here ΔI is the difference between the isospin of the initial and final states. Stech (1987, 1989) had already pointed out the relevance of virtual diquarks in weak decays of mesons (see Sec. III.C). The $|\Delta I|=\frac{1}{2}$ enhancement arises because the Pauli principle requires a scalar diquark containing a u and a d quark to have $I=0$. Dosch *et al.* (1989) calculate the dominant terms (pole terms) of P-wave hyperon decays and obtain good qualitative agreement with experiment (see Table 1 of their paper). As far as we know, the work of Dosch *et al.* provides the only qualitatively simple explanation to date for the $|\Delta I|=\frac{1}{2}$ rule in baryon decay. However, more work is needed to justify the absence (or small effects) of vector diquarks.

Efimov *et al.* (1990) have used a quark-diquark model of baryons to calculate both static and decay properties of baryons. The calculation of static properties, such as baryon magnetic moments, is not particularly good, but the authors achieve remarkably good agreement with experiment for the decay rate of a member of the spin-$\frac{3}{2}$ baryon decuplet into a pion plus a member of the spin-$\frac{1}{2}$ baryon octet. We should not fault Efimov *et al.* for their poor magnetic-moment calculations, because, as we have already remarked, the explanation for the observed baryon magnetic moments lies elsewhere.

E. Exotic baryons

So far, we have considered normal baryons (three-quark systems). There also exists the possibility of exotic baryons, containing more than three quarks. The simplest of such baryons are states containing four quarks and an antiquark (five-quark or pentaquark states). There is no good experimental evidence for the existence of five-quark baryon states, and because of lack of space we do not discuss such states. For the interested reader, we simply refer to a few theoretical papers (Gignoux *et al.*, 1987; Lipkin, 1987; Fleck *et al.*, 1989; Leandri and Silvestre-Brac, 1989) through which additional references can be traced.

III. PROPERTIES OF MESONS

A. Baryonium and other exotic mesons

In the constituent-quark model, an ordinary meson is a bound state of a quark-antiquark pair. However, it is plausible from QCD that exotic mesons should exist which cannot be so described. A manifestly exotic meson is one whose quantum numbers are such that it cannot be a quark-antiquark pair. Otherwise, the exotic nature of the meson is hidden.

Among exotic mesons we include a meson composed of two quarks and two antiquarks (for short, a four-quark state), a meson composed only of gluons (glueball), and a meson composed of a quark and antiquark plus gluons (hybrid). To date, we have some experimental evidence for the existence of only a few exotics, and all or nearly all of these are not manifestly exotic. If a meson is exotic but not manifestly exotic, theoretical interpretation in addition to experimental information is necessary to identify it. The interpretation is difficult because idealized four-quark states can mix with glueballs, hybrids, and ordinary mesons.

Of exotic mesons, only the four-quark states are candidates for being composed of a diquark and antidiquark. (Some authors use the term "diquonium" for a state composed of a diquark-antidiquark pair.) Another possibility is that a four-quark state is a "molecule" composed of a lightly bound or resonant state of two mesons. There have been many discussions of this possibility, of which we mention a recent paper by Törnqvist (1991) and references therein. Close (1993) has given a good discussion of some of the issues involved in whether four-quark states are primarily composed of two diquarks or of two-meson molecules. Still a third possibility is that a four-quark state is primarily composed of two quark-antiquark pairs, each of which is in a color-octet state. Here, we focus primarily on diquark-antidiquark systems.

The possibility of diquark-antidiquark states was mentioned quite early (Lichtenberg and Tassie, 1967), but was not treated in detail because of the absence of any evidence for exotic mesons. Although the experimental situation is not much better today, a large number of authors have considered exotic mesons as diquark-antidiquark bound states. For brevity, we often call a color-antitriplet diquark a color triplet, and we call a color-antisextet antidiquark a color sextet. Gromes (1988) has emphasized that color-triplet and color-sextet diquarks are in general not gauge-invariant concepts without a specification of how they are coupled to form color singlets. Depending upon the dynamics, there could be mixing between color-triplet and color-sextet configurations. We do not consider questions of gauge invariance further.

Among the candidates for diquark-antidiquark mesons are the so-called baryonium states, which are defined in this context as mesons with an appreciable probability to decay into a baryon-antibaryon pair. A baryonium state defined in this way may or may not be a diquonium state: it could be a bound or resonant state of a baryon-antibaryon pair. Stimulated by preliminary experimental evidence for a large number of baryonium states, many theorists considered them candidates for four-quark states. Even though much of the experimental evidence for baryonium resonances has not been confirmed, we review here some of the theoretical arguments for the existence of four-quark states. It is probable that such states exist, but that they are hard to produce, hard to recognize, or both. For example, if they have large decay widths or mix appreciably with two-quark states, identification becomes difficult. We cannot quote all the papers that have appeared on this subject. For a review, see Montanet *et al.* (1980).

Freund and Rosner (1992) have given an argument based on a string theory approximation to QCD that the number of mesons and baryons per unit energy interval should become equal at high energy. However, if baryons consist only of three-quark states and mesons only of two-quark states, the number of baryons per energy interval will eventually exceed the number of mesons in the same interval. The only plausible way known to prevent this is to include four-quark (i.e., two quarks, two antiquarks) meson states, five-quark (four quarks, one antiquark) baryon states, etc.

Much of the analysis of four-quark states has been done with an effective simple Hamiltonian,

$$H=-\sum_{i<j} a\mathbf{F}_i\cdot\mathbf{F}_j\mathbf{S}_i\cdot\mathbf{S}_j/(m_im_j)\ , \tag{3.1}$$

where a is a positive constant, $\mathbf{F}_i$ is the color operator, $\mathbf{S}_i$ is the spin, and m_i is the mass of the ith quark. In Eq. (3.1) an additive constant term which depends on the flavors of the quarks is omitted.

The above Hamiltonian is suggested by the form of the color-hyperfine interaction of QCD. If all quarks are considered as free particles inside a bag, then it might be enough to consider just the Hamiltonian in Eq. (3.1). However, it cannot be a good approximation to neglect all other terms suggested by QCD, so that the principal reason for using this Hamiltonian is that it allows one to calculate the energy of any n-quark state with a given wave function in a straightforward way. However, even the color-hyperfine term is oversimplified, because the effect of the spatial wave function is neglected. Diquark correlations in this wave function require (at least) that the constant a in the Hamiltonian be replaced by a_{ij}. Sometimes, when hadrons containing only u, d, and s quarks are treated, the Hamiltonian (3.1) is further simplified by the neglect of the flavor dependence arising from the quark masses in the denominator.

Silvestre-Brac (1992) has given a systematic treatment of four-quark mesons using the Hamiltonian of Eq. (3.1). In his work, which we consider in more detail in Sec. III.B, he includes mixing of sextet-antisextet and triplet-antitriplet diquarks. This same Hamiltonian has been applied to other exotic hadrons, such as five-quark baryons and six-quark dibaryons (see Sec. V.A).

We now return to earlier work. Johnson and Thorn (1976), using the MIT bag model (Chodos *et al.*, 1974), consider four-quark mesons. In their picture, for those states which lie on the leading Regge trajectory, the bag is deformed into a cigar shape, with a color source at each end. The color sources may consist of a diquark and antidiquark, each with color multiplicity of three or six, connected by a color-electric flux. Alternatively, each source may consist of a "color-octet meson," which is defined to be a quark and an antiquark in a color-octet state. Jaffe (1978; see also Aerts *et al.*, 1980) gives an extensive treatment of four-quark mesons in the same picture, remarking that the idea of diquarks simplifies his discussion "immensely."

Johnson and Thorn and Jaffe calculate the asymptotic slope (at large orbital angular momentum L) of the Regge trajectory of a four-quark state. They find that the Regge slope is inversely proportional to the square root of the quadratic Casimir operator of the color source at each end of the elongated bag. A diquark in a color-$\overline{\mathbf{3}}$ state has a Casimir operator $F^2=\frac{4}{3}$, while in a color-$\mathbf{6}$ state it has $F^2=\frac{10}{3}$. Alternatively, a color-octet meson may be at each end of the bag, in which case $F^2=3$. The authors conclude that the Regge trajectory of states composed of a diquark-antidiquark pair, each in a color-triplet state, has the same slope as ordinary meson and baryon trajectories. On the other hand, the trajectory of exotics made of color-sextet diquarks should have a slope only $\sqrt{4/10}=0.63$ as large, and the trajectory of exotics made of color-octet mesons should have a slope $\frac{2}{3}$ as large as the ordinary slope. The slopes 0.63 and 2/3 are almost the same, and there is little hope that they can be distinguished experimentally. There is an additional complication that at small values of L the trajectories are expected to have some curvature.

Chan and Høgaasen (1977), Hendry and Hinchliffe (1978), and Barbour and Ponting (1979) discuss ideas similar to those of Johnson and Thorn (1976) and Jaffe (1978). Hendry and Hinchliffe adopt a string picture, with a color-$\overline{\mathbf{3}}$ diquark at one end and a color-$\mathbf{3}$ antidiquark at the other. Hendry and Hinchliffe use an SU(6) (of flavor and spin) classification scheme, in which color-triplet diquarks belong to a **21** multiplet. They briefly discuss color-$\mathbf{6}$ diquarks, but suggest that four-quark states composed of them lie higher in energy.

This point of Hendry and Hinchliffe is supported by the work of Gromes (1988). In a calculation of properties of four-quark states, Gromes gives good arguments that a sextet diquark has a considerably larger mass than a triplet diquark. The neglect of this mass difference in treatments of states with the Hamiltonian of Eq. (3.1) is likely to lead to some results that are qualitatively incorrect.

Lipkin (1978) uses a phenomenological potential behaving approximately as the logarithm of the distance between quarks. With such a potential, Lipkin is able to relate the properties of mesons, baryons, and baryonium states (the last being considered as composed of a diquark-antidiquark pair).

Chan and Høgaasen do not use an explicit model but rather argue that QCD requires four-quark states to exist. They do assume that a confining potential exists and is independent of spin and flavor. They also assume the existence of the usual color-hyperfine interaction. Although many of the four-quark states should have broad decay widths and be indistinguishable from the continuum background, the authors state that certain quark configurations should have narrower widths and be more amenable to observation. These configurations consist of a diquark and antidiquark in a state with $L\neq 0$. The diquark and antidiquark, however, are in their lowest states, with no internal orbital angular momentum. As in the models of other authors, the states lie on Regge trajectories whose (common) slope is determined by the quadratic Casimir operator. Therefore, in this picture, too, mesons made of $\mathbf{3},\overline{\mathbf{3}}$ diquarks lie on trajectories with common slope equal to the universal slope of ordinary mesons and baryons, while the slope of the trajectories of mesons made of $\mathbf{6},\overline{\mathbf{6}}$ diquarks is only 0.63 times as large.

The argument of various authors that Regge trajectories of exotics made of color-sextet, color-triplet, and color-octet clusters should have different slopes may be in error. Lattice gauge calculations (Griffiths *et al.*, 1985; Hasenfratz, 1987) indicate that gluons screen all colored clusters to triplets and that, therefore, the string tension and the Regge slope are universal, at least for multiquark states containing only light quarks.

It is remarkable that several groups of authors, working independently at around the same time with different models (bag, string, qualitative QCD, potential), obtain the same picture of four-quark states, although details of the spectrum vary. In all of the models, the spectrum of diquark-antidiquark states is very rich. This entire picture is criticized by Isgur (1989) as based on too simple assumptions, too crude approximations, or both (see also Kalman and Misra, 1982). Much of Isgur's criticism is justified. Nevertheless, diquark-antidiquark states still ought to exist, but, as Pennington (1991) points out, such states are likely to have typical hadronic widths and overlap in energy. According to Pennington, an amplitude analysis of $p\bar{p}$ scattering will be required to discover these states. It is clear that additional theoretical and experimental work will be necessary before we have a good understanding of the nature of four-quark states.

B. Exotic mesons containing heavy quarks

Just as there might exist diquonium states composed of light quarks, there also might exist similar states composed of any combination of light and heavy quarks. Rosenzweig (1976) and others (e.g., Kenny *et al.*, 1976; Chao, 1981) propose that some of the higher-lying levels observed in the charmonium ($c\bar{c}$) spectrum might not be charmonium, but four-quark states composed of a charmed diquark and a charmed antidiquark. We now believe that most, if not all, of the four-quark candidates are ordinary $c\bar{c}$ states. Nevertheless, states composed of heavy diquark-antidiquark pairs might exist. As with light four-quark mesons, if they exist, they are either hard to produce, hard to identify, or both.

An interesting case of exotic mesons (e.g., Chao, 1980) is $Q\,Q\,\bar{q}\,\bar{q}$, where Q is a heavy quark and q is a light one. A number of authors (Ader *et al.*, 1982; Ballot and Richard, 1983; Lipkin, 1986; Zouzou *et al.*, 1986; Heller and Tjon, 1987; Carlson *et al.*, 1988) agree that if the ratio of the mass of the heavy quark to that of the light one is sufficiently high, then the state is stable against decay into the two mesons $Q\bar{q}+Q\bar{q}$. Moreover, the configuration is that of a color-$\bar{\mathbf{3}}$ diquark QQ bound to a color-$\mathbf{3}$ antidiquark $\bar{q}\,\bar{q}$.

We give here a simple argument why this is the case. At small distances between two quarks (or a quark and an antiquark), the potential between them is an attractive Coulomb potential (we are neglecting spin-dependent forces). Sufficiently heavy quarks are close enough together to feel this potential more than the linear confining potential. But the binding energy of two heavy quarks in a Coulomb potential is proportional to the mass m_Q of the heavy quark. On the contrary, the binding energy of a bound state of $Q\bar{q}$ in the same approximation is proportional only to the reduced mass of the system. The effect of the confining potential is only to lessen the binding energy, and its effect is stronger in the $Q\bar{q}$ system than in the QQ system. The fact that the effective coupling strength is a factor 2 greater for $Q\bar{q}$ than for QQ cannot overcome the effect of m_Q, provided the latter is sufficiently large. Therefore, the system consists of a small, tightly bound QQ diquark bound to two light antiquarks. The actual situation is complicated by the mixing of color $\mathbf{6}$ and $\bar{\mathbf{6}}$ states and by the color-hyperfine interaction, but the essential argument is unaffected.

Carlson *et al.* (1988) have performed model calculations for $c\,c\,\bar{q}\,\bar{q}$, $b\,c\,\bar{q}\,\bar{q}$, and $b\,b\,\bar{q}\,\bar{q}$. They conclude that $b\,b\,\bar{q}\,\bar{q}$ is stable against strong decay, $c\,c\,\bar{q}\,\bar{q}$ is unstable, and $b\,c\,\bar{q}\,\bar{q}$ is uncertain. Of the states containing bb, the state $b\,b\,\bar{u}\,\bar{d}$ should have the lowest mass, with the bb diquark having spin 1 and the $\bar{u}\bar{d}$ antidiquark having spin 0. Thus the quantum numbers of the state are $J^P=1^+$. The state is manifestly exotic (since it contains two b quarks), and its discovery would be a triumph for our present understanding of diquarks.

Ma *et al.* (1992) suggest that there might be a bq-$\bar{b}\bar{q}$ state near the $B\bar{B}$ threshold. The authors claim that such a state might explain the observation by Alexander *et al.* (1990) of an anomalously large number of high-momentum J/ψ mesons produced near the $\Upsilon(4S)$ state. In order to obtain a large three-gluon decay width, the authors choose their new state to consist of color-sextet diquarks. However, subsequent data (Bortoletto *et al.*, 1992) do not confirm high-momentum J/ψ production in excess of continuum production. We conclude that there is at present no evidence for a four-quark state near the $B\bar{B}$ threshold.

Silvestre-Brac (1992) treats diquonium states composed of all possible flavor (excluding the t quark), spin, and isospin combinations of two quarks and two antiquarks. He uses a model Hamiltonian of the form of Eq. (3.1). With this simple Hamiltonian, Silvestre-Brac is able to calculate which diquonium states are bound, i.e., which states have masses below the sum of the masses of the two mesons which contain the same quarks. Of the 584 diquonium states computed, Sylvestre-Brac finds that 110 are stable, 15 of which are bound by at least 100 MeV. An important problem with the calculation, as the author himself recognizes, is that the model Hamiltonian leads to degeneracy between the π and η mesons, whereas, experimentally, the pion mass is 140 MeV, while the η mass is 547 MeV. If the author excludes those diquonium states with quantum numbers that permit a pion as a decay product (assuming conservation of isospin), he finds that the most interesting diquonia in his model are $u\,d\,\bar{b}\,\bar{b}$, $u\,d\,\bar{c}\,\bar{c}$, $u\,d\,\bar{c}\,\bar{b}$, $u\,d\,\bar{s}\,\bar{b}$, $u\,d\,\bar{s}\,\bar{s}$, $u\,d\,\bar{s}\,\bar{c}$ with isospin $I=0$, spin $S=1$, $u\,s\,\bar{u}\,\bar{s}$, $u\,d\,\bar{c}\,\bar{b}$, $u\,d\,\bar{s}\,\bar{b}$, $u\,d\,\bar{s}\,\bar{c}$ with $I=0$, $S=0$; and $u\,s\,\bar{b}\,\bar{b}$ with $I=\frac{1}{2}$, $S=1$. The model predicts that all of these diquonium states are bound by at least 100 MeV against strong decay into two mesons; but, as the author admits, the model probably overestimates the binding energies. We think that some of these states may not even be bound.

Some of the calculations we have discussed in this subsection are quite sound, for example, the calculations that say that $q\,q\,\bar{Q}\,\bar{Q}$ states are bound for a sufficiently heavy Q. The best candidate is $u\,d\,\bar{b}\,\bar{b}$ with $I=0$, $S=1$. Unfortunately, most of the calculations are so oversimplified that they can at best provide only a guide to where further work is needed. However, we expect that with further experimental and theoretical effort, the situation will become clearer.

C. Weak decays of mesons composed of quark and antiquark

Stech and collaborators (Stech, 1987, 1989; Neubert and Stech, 1989; Stech and Xu, 1991) have given arguments that weak nonleptonic decays of ordinary (quark-antiquark) mesons are mediated by color-triplet diquark virtual states. In this picture, a meson undergoes a transition to a virtual scalar diquark and antidiquark, and these latter particles in turn convert with unit probability into final states consisting of mesons. A recent paper by Neubert and Stech (1991) gives a good analysis of the weak decays of strange mesons and contains additional references.

With their model, the authors are able to account for the $|\Delta I|=\frac{1}{2}$ rule in the decay of K mesons. In the decay, the $|\Delta I|=\frac{1}{2}$ amplitude is enhanced because a scalar diquark containing u and d quarks must have isospin zero. We saw in Sec. II that Dosch *et al.* (1989) were able to explain the $|\Delta I|=\frac{1}{2}$ rule in baryons on the basis of diquarks. Thus scalar diquarks are the common dynamical mechanism that gives a simple explanation in both mesons and baryons for the $|\Delta I|=\frac{1}{2}$ rule, which is otherwise hard to understand. Unfortunately, as we already noted in Sec. II, we must understand the absence of effects of vector diquarks before we can say we really understand the $|\Delta I|=\frac{1}{2}$ rule.

Stech and collaborators also invoke diquarks to account for the fact that the D^0 meson has a shorter lifetime than the D^+. The model leads to a suppression of the D^+ $(c\bar{d})$ decay rate relative to that of the D^0 $(c\bar{u})$, because in the former case two amplitudes destructively interfere. Because the diquark-antidiquark intermediate states mainly lead to multimeson final states, the usual treatment of exclusive two-body D $D_s(c\bar{s})$ decays does not have to be modified very much except in channels with small kinetic energy in the final state. In B-meson decays, diquark-antidiquark intermediate states should usually lead to final states with a baryon-antibaryon pair. A measurement of such states should give an indication of how important the mechanism is in B decay.

D. Exotic meson decays

Chan and Høgaasen (1977) and Jaffe (1978) use the quark-line diagrams of Harari (1969) and Rosner (1969) to argue that diquark-antidiquark states with large L are likely to decay into baryon-antibaryon pairs. This occurs naturally only if the diquarks are in color-$\bar{3}$ states, and these are the states which should be classified as baryonium states. For the case in which the diquark belongs to a color sextet, a four-quark state should have only a small probability to decay into a baryon-antibaryon pair.

Hendry and Hinchliffe (1978) have given the argument, based on a string picture, that a meson decaying into a baryon and antibaryon is likely to be a diquark-antidiquark exotic state. If a diquark is attached to an antidiquark by a string, and the string breaks, the simplest decay will be into a baryon-antibaryon pair. In such a configuration, the model of Hendry and Hinchliffe is similar to that of Johnson and Thorne and Jaffe.

A difficulty with such schemes is that, if the quantum numbers permit, four-quark exotics can mix with hybrids, ordinary excited mesons, and glueballs. Such mixing is likely to change the decay characteristics of four-quark states and therefore make them difficult to recognize.

IV. DEEP-INELASTIC LEPTON-NUCLEON SCATTERING

A. General features

We consider here and in the next sections several phenomenological implications of the idea that two quarks act coherently as a single constituent in nucleons and in dynamical processes involving hadrons in general and baryons in particular. We start with the processes in which the internal structure of nucleons was first explored, namely, deep-inelastic scattering (DIS).

The experiments on deep-inelastic lepton-nucleon scattering and the discovery of scaling in the nucleon structure functions led to the parton model, i.e., to the picture of nucleons as composite objects made of pointlike partons. Further data have established the nature of charged partons as spin-$\frac{1}{2}$ constituents, or quarks. Nevertheless, two-quark correlations exist in baryons, and one can usefully introduce diquarks into this picture. In order to accommodate spin-0 (or scalar) diquarks and spin-1 [or (axial) vector] diquarks as nucleon constituents, we briefly recall the formalism of the parton model of deep-inelastic lepton-nucleon scattering.

The cross section for unpolarized deep-inelastic lepton scattering by a nucleon of four-momentum P via a virtual-photon of four-momentum q can be written in terms of two structure functions, $F_1(\nu,Q^2)$ and $F_2(\nu,Q^2)$ (see, e.g., Leader and Predazzi, 1983), where $q^2=-Q^2$. We also let $P^2=m_N^2$, $x=Q^2/(2P\cdot q)$, and $P\cdot q=m_N\nu$. The interaction of the virtual photon with the nucleon is described as a sum of the interactions of the photon with all the parton constituents of the nucleon taken as free particles. If x is the fraction of the nucleon momentum carried by a parton, we can express F_1 and F_2 in terms of the number density of partons and explain scaling in the variable x.

Now consider diquarks as (extended) elementary constituents. We denote by $S(x)$ $[V(x)]$ the number density of scalar [vector] diquarks in a nucleon. Then the diquark contributions to the structure functions are (Esaibegyan and Matinyan, 1974; Pavkovic, 1976b; Fredriksson *et al.*, 1983a; Linkevich, Savrin, Sanadze, and Skachkov, 1983; Linkevich, Savrin, and Skachkov, 1983; Leader and Anselmino, 1988; Anselmino, Caruso, Leader, and Soares, 1990; Meyer and Mulders, 1991)

$$
\begin{aligned}
F_1^{(S)} &= 0 , \\
F_2^{(S)} &= e_S^2 S(x) x D_S^2 , \\
F_1^{(V)} &= \frac{1}{3} e_V^2 V(x) \left[1+\frac{\nu}{2m_N x}\right] D_2^2 , \\
F_2^{(V)} &= \frac{1}{3} e_V^2 V(x) x \left\{ \left[\left[1+\frac{\nu}{m_N x}\right] D_1 - \frac{\nu}{m_N x} D_2 \right.\right. \\
&\qquad \left. + 2 m_N \nu x \left[1+\frac{\nu}{2m_N x}\right] D_3 \right]^2 \\
&\qquad \left. + 2 \left[D_1^2 + \frac{\nu}{2m_N x} D_2^2 \right] \right\} ,
\end{aligned}
\tag{4.1}
$$

where $D_S(Q^2)$ and $D_i(Q^2)$ $(i=1,2,3)$ are form factors in the electromagnetic couplings of the virtual photon with a scalar diquark and a vector diquark, respectively. There are also contributions from transitions between scalar-vector ($S-V$) and vector-scalar ($V-S$) diquarks:

$$\begin{aligned} F_1^{(S-V)} &= \frac{1}{2} e_S^2 S(x) x^2 m_N^2 \left[1 + \frac{\nu}{2m_N x} \right] D_T^2 \\ F_2^{(S-V)} &= \frac{1}{2} e_S^2 S(x) x^2 m_N \nu D_T^2 \\ F_1^{(V-S)} &= \frac{1}{6} e_S^2 V(x) x^2 m_N^2 \left[1 + \frac{\nu}{2m_N x} \right] D_T^2 \\ F_2^{(V-S)} &= \frac{1}{6} e_S^2 V(x) x^2 m_N \nu D_T^2 \ , \end{aligned} \tag{4.2}$$

where $D_T(Q^2)$ is a transition form factor. For more details, including the diquark contribution to the polarized structure functions, see Anselmino, Caruso, Leader, and Soares (1990). In the limit of pointlike diquarks the form factors are given by $D_S(0)=1$, $D_1(0)=1$, $D_2(0)=1+\kappa$, $D_3(0)=0$, $D_T(0)=0$, where κ is the vector-diquark anomalous magnetic moment.

Scaling in x is not violated by pointlike scalar diquarks. Pointlike vector diquarks, however, lead to strong scaling violations, which are incompatible with the data. As diquarks are extended objects, a realistic comparison with experiment should take into account their form factors D_S, D_i, and D_T. These are expected to decrease as inverse powers of Q^2. As seen from Eqs. (4.1) and (4.2), scaling violations are lessened by the Q^2 dependence of the diquark form factors. The existence of powerlike corrections to the logarithmic scaling violations induced by QCD is indeed confirmed by the data (Aubert *et al.*, 1985; Milsztajn *et al.*, 1991; Botje, 1992; Virchaux and Milsztajn, 1992). We also know that, within the parton model, aside from logarithmic QCD corrections, the Callan-Gross relation, $F_2^{(q)}=2xF_1^{(q)}$, requires the following large Q^2 behavior of the so-called R ratio,

$$R = \frac{F_2}{2xF_1} \left[1 + \frac{2m_N x}{\nu} \right] - 1 \sim \frac{1}{Q^2} \ . \tag{4.3}$$

With diquarks it is still possible to fulfill this relation while keeping scaling violations proportional to $1/Q^2$ or smaller (Anselmino, Caruso, Leader, and Soares, 1990) with the assumption, suggested by a perturbative QCD analysis of the diquark form factors (Vainshtein and Zakharov, 1978; Kroll and Schweiger, 1989), $D_S \sim Q^{-2}$, $D_1 \sim D_2 \sim Q^{-4}$, $D_3 \sim Q^{-6}$. In addition, the choice $D_S \sim Q^{-2}$, $D_1 = D_2 \sim Q^{-2}$, $D_3=0$ (Fredriksson *et al.*, 1983a; Leader and Anselmino, 1988) satisfies Eq. (4.3) in the same manner.

The choice $D_1=D_2$, $D_3=0$ implies that vector diquarks have no anomalous magnetic moment and that $F_2^{(V)} \simeq 2xF_1^{(V)}$. Hence, vector-diquark contributions to F_1 and F_2 satisfy the Callan-Gross relation to leading order and show that the relation $F_2=2xF_1$ does not necessarily imply that partons have spin $\frac{1}{2}$, as is often stated.

Recently, Dziembowski (1993) discussed DIS from a different point of view. First, the nucleon wave function is considered in the rest system with a distorted harmonic-oscillator basis (Franklin, 1968), which allows for diquark correlations. The deformation parameter is adjusted so as to obtain agreement with the proton and neutron charge radii. The result is that a nucleon has a positive core consisting of a spin-0 constituent (*ud* scalar diquark) and an outer layer containing the (polarized) third quark. The wave function is then transformed to momentum space, light-cone variables are chosen, and a QCD evolution is performed. Rather good agreement is found with the experimental ratio of the neutron-to-proton structure functions and the proton spin asymmetry. This work presents a unified picture of low-energy and high-energy properties of the nucleon in terms of a quark-diquark picture. Dziembowski's paper is related to the earlier work of Franklin (1980), Close and Thomas (1988), and Dziembowski and Franklin (1990).

B. Extracting diquark momentum distributions and form factors

We now consider more detailed phenomenological implications of diquarks in DIS. The aim is to review the published efforts to get best fits for the admixture of diquark states in nucleons, diquark form factors, and x distributions. As will be made clear, there are many interesting approaches to this task, but so far a consensus has been reached only for qualitative findings such as the dominance of scalar diquarks over vectors in nucleons.

The main ambiguity is related to how one should include perturbative QCD contributions. At one extreme, diquarks can be made responsible for only those features of the data that remain after an overall best fit of perturbative contributions. At the other extreme, one might investigate how much of the data diquarks can reproduce without any perturbative gluon correction to the usual quark-parton model. A reasonable but more elaborate approach is to try a simultaneous best fit of both diquark and perturbative QCD parameters. In general, diquarks in DIS take phenomenologically into account all interactions between the two quarks in the diquark, both perturbative and nonperturbative ones. At large values of Q^2 only the leading, $1/Q^2$, corrections are significant. The diquark is then resolved into two quarks, and these leading corrections can be related to the $1/Q^2$ higher-twist corrections one obtains in the usual operator product expansion analysis of DIS.

In most circumstances, scalar diquarks are more important than vector diquarks in nucleons because of the QCD spin forces (for early models of DIS with vector diquarks only, see Esaibegyan and Matinyan, 1974; Anisovich, 1975; Anisovich *et al.*, 1976; Kawabe, 1983). The scalar-vector transition terms are usually neglected (except for neutrino scattering; see below), and the three vector form factors are often assumed for simplicity to obey $D_1 = D_2$ and $D_3 = 0$. However, the influence of vector diquarks is not generally negligible in the proton, because the vector (uu) diquark in the proton is enhanced by a factor of 16 compared to the scalar (ud), due to its charge. In certain kinematical regions, this could balance the differences in form factors and momentum distributions between vectors and scalars.

Most of the early diquark workers attempted to make plausible the idea that diquarks can explain the *full* Q^2 dependences in F_2 and F_1. A fair agreement with the early SLAC data of, e.g., Bodek *et al.* (1979) was claimed, both with SU(6)-symmetric diquarks (Schmidt and Blankenbecler, 1977; Fernandez-Pacheco *et al.*, 1978; Donnachie and Landshoff, 1980) and with dominating scalars (Pavkovic, 1976a, 1976b; Abbott *et al.*, 1979, 1980; Carlitz and Creamer, 1979).

This original optimism was questioned by Close and Roberts (1981). They argued that, because a neutrino cannot interact directly with a scalar diquark, the neutrino induced data should be quite different from data taken with charged-lepton beams. The SU(6)-symmetric model of Donnachie and Landshoff (1980) was shown to be in serious disagreement with the data (Bodek *et al.*, 1979; de Groot *et al.*, 1979).

Dismissing the assumption that neutrinos do not interact at all with diquarks (Kawabe, 1982), plausible alternatives are to question either the validity of the SU(6) symmetry or the assumption by Close and Roberts about how diquarks interact with neutrinos. Fredriksson *et al.* (1982, 1983a, 1983b) assume that SU(6) is severely broken in favor of the scalar (ud) diquark and that this diquark interacts with a neutrino by turning into a vector (uu) diquark. In this model, diquarks do not break up after the interaction, as assumed by Donnachie and Landshoff, and it is postulated that the nucleon is always in a $q(ud)$ state with a scalar (ud) diquark. Vector diquarks occur in the model only in "accidental" interactions involving quark pairs other than the scalar (ud). The scalar (ud) is more pointlike than the vector (ud) and (uu), and the scalar x distribution is broader and more like that of a single quark. Using

$$D(Q^2) = \left[1 + \frac{Q^2}{Q_0^2} \right]^{-1} \tag{4.4}$$

for the form factors, one finds, from the data available a decade ago, that $Q_0^2 \simeq 10$ (GeV)^2 for the scalar and only $\simeq 2$ (GeV)^2 for the vector. Hence in this model a scalar diquark is much smaller than a nucleon, and the vector-diquark radius is more than twice that of the scalar. The vector x distribution peaks at $x \simeq \frac{2}{3}$, suggesting that vector diquarks consist of two almost uncorrelated quarks.

In principle, one can extract the quark and diquark x distributions and form factors from the data on F_1 and F_2 by computing various sums and differences (Close and Roberts, 1981; Fredriksson *et al.*, 1982), if one attributes all the observed Q^2 dependence to diquarks. The quality of those data does not allow more model-independent conclusions than the ones quoted above. A precise model with quarks, "genuine" scalar diquarks, and "accidental" vector diquarks can become quite complex, especially if one wants to avoid inconsistencies like double counting of quarks and diquarks.

Recently, a detailed analysis of DIS data, in the framework of the quark–scalar-diquark nucleon model of Fredriksson *et al.* (1982), was performed by Dugne and Tavernier (1992). They find excellent agreement with the data, both with a value $Q_0^2 = 10$ (GeV)^2 and $Q_0^2 = 3$ (GeV)^2 [see Eq. (4.4)], provided they include perturbative QCD corrections. They also present interesting results on the Drell-Yan process. A similar analysis, with almost equally good results, has been performed within a model with both scalar and vector diquarks (Tavernier, 1992). Such a model, however, does not take into account the inelastic contributions of the diquarks. This shows how difficult it might be to extract unambiguous results from the data. Some more detailed analyses are, nevertheless, in preparation (Caruso, 1992).

When $Q^2 \to \infty$, all diquarks should be resolved, and one should obtain the usual quark model. High-Q^2 data can be used to deduce the x distributions for both u and d quarks in the proton. They differ in that a d quark carries, on the average, substantially less momentum than a u quark, as if the d quark were staying mostly inside a (ud) diquark, which has about the same momentum as the single u quark. One can estimate the x distribution of the single u quark by subtracting that of the d quark from the total u distribution.

If diquarks are responsible for the full Q^2 dependence, the structure functions should approach the (scale-invariant) quark values as Q^2 increases. Recent high-Q^2 data (Benvenuti *et al.*, 1989, 1990) seem to push the scale-invariant regime far beyond 10 (GeV)^2. The data can be interpreted as evidence either that perturbative QCD effects dominate only at $Q^2 \gg 10$ (GeV)^2, or that the scalar (ud) diquark is effectively smaller than anticipated (or both).

C. Other comparisons with the data

It is useful to take appropriate combinations of the structure functions F_1 and F_2 from proton and neutron (deuterium) targets and to minimize ambiguities associated with extended vector diquarks by using only data with, say, $Q^2 \gg 2$ (GeV)^2.

(a) A much quoted combination of structure functions is $3(F_2^{lp} - F_2^{ln})$, where l stands for a charged lepton. In the quark model, with an isospin-symmetric sea, this combination gives the difference between the x distribution of a u and a d quark in the proton, while in a diquark model (without vectors) it gives the distribution of the single (i.e., nondiquark) u quark. One can thus study the x and Q^2 dependence in F_2 due to nondiquark contributions (Ekelin and Fredriksson, 1985a).

Several comments are in order. First, in the high-x region no relevant Q^2 dependence results from comparing the medium-Q^2 SLAC data (Bodek *et al.*, 1979) with the high-Q^2 EMC (European Muon Collaboration) data (Aubert *et al.*, 1983a, 1983b). In addition, neutrino data (Allasia *et al.*, 1985) show a weaker Q^2 dependence in the difference $F_2^{\nu p}-F_2^{\nu n}$ than in the structure functions themselves. This is not expected from perturbative QCD and could be a signature of scalar diquarks. Second, the high-x dependence is not the $(1-x)^3$ expected from dimensional counting rules (Brodsky and Farrar, 1973; Matveev *et al.*, 1973). Rather, the falloff with x is linear at $x>0.4$, as if the single u quark had only one spectator—the scalar diquark. In addition, the kinematic limit for this linear falloff does not seem to be at $x=1$, as expected with massless current quarks, but rather at $x\simeq 0.87$, as if the single u quark could not carry the full proton momentum because of a (diquark) spectator rest mass, of the order of 300 MeV (not incompatible with other estimates of the scalar diquark mass).

Turning to more recent results, the BCDMS (Bologna-CERN-Dubna-Munich-Saclay) collaboration data (Benvenuti *et al.*, 1990) have a Q^2 dependence in $F_2^{\mu p}-F_2^{\mu n}$ at high x, compatible with expectations from perturbative QCD. However, the recent EMC data from deuterium (Aubert *et al.*, 1987), although less precise than the BCDMS data, do not show the same Q^2 dependence. The issue needs clarification, since the BCDMS data obviously disagree with the notion of a very small scalar diquark in nucleons, while the EMC data do not. The very recent NMC (New Muon Collaboration) data (Amaudruz *et al.*, 1992b) seem to fit well the perturbative QCD expectation for μp and μD scattering separately, but an even better fit is achieved when a substantial higher-twist effect is added (Botje, 1992). Since this Q^2 contribution differs somewhat from contributions expected from diquarks, a reanalysis of the NMC data in terms of diquarks could clarify this important issue.

The data discussed above are known not to obey the Gottfried sum rule (Gottfried, 1967). Rather $3\int(F_2^{\mu p}-F_2^{\mu n})dx/x\simeq 0.7$ (Amaudruz *et al.*, 1991) instead of unity (as expected from the number of contributing quarks). This can be accounted for by (vector) diquarks persisting at very high Q^2 (Anselmino *et al.*, 1992, and references therein). This explanation, however, seems unlikely, as it does not correspond to our idea of a diquark being a correlated pair of quarks. See, for example, the very recent paper by Caruso and Leader (1992). An isospin-asymmetric sea can also contribute to a deviation of the integral from unity, and still other explanations are possible (see Forte, 1993).

(b) A second combination of structure functions is the ratio F_2^n/F_2^p. As $x\to 1$ and $Q^2\to\infty$, the value 0.25 is expected from the dominance of the single (nondiquark) quark (Carlitz and Creamer, 1979). The same value is expected in the quark model (Close, 1973). Diquark models predict a strong Q^2 dependence, with the ratio decreasing with increasing Q^2 for all x, whereas in perturbative QCD, the dominant Q^2 dependence factorizes in the structure functions (see, e.g., Duke and Owens, 1984) and hence cancels in the ratio. A comparison between the SLAC and EMC data reveals a clear Q^2 dependence, in favor of diquark models. This trend persists in the recent NMC data (Allasia *et al.*, 1990; Amaudruz *et al.*, 1992a). For analyses of this type using neutrino data, see Allen *et al.* (1981). An analysis of neutrino data (Jones *et al.*, 1989) claims a disagreement between the "data" on the (valence) quark ratio $d(x)/u(x)$ and the expectations from two different diquark models. It is not clear, however, how the diquark calculations for charged leptons were compared with neutrino results, given the delicate role played by the transition of a scalar (ud) into a vector (uu), which was mentioned earlier.

(c) A third combination is the ratio $R=\sigma_L/\sigma_T$ of the cross sections for longitudinally and transversely polarized virtual photons, given in Eq. (4.3). Equations (4.1)–(4.3) show that R is a measure of the ratio of diquarks to quarks, with charges and form factors included, if vector diquarks can be neglected. This was pointed out by Abbott *et al.* (1979), who also analyzed the early SLAC data and showed that the measured R values were not inconsistent with diquarks. See also the earlier, but less conclusive, study by Ono (1974a, 1974b, 1975) and the one by Peng and Zou (1980).

Four more recent diquark model analyses have been presented. One of them (Ekelin and Fredriksson, 1985a) claims agreement with the EMC and SLAC data, while another (Whitlow *et al.*, 1990; see also Dasu *et al.*, 1988) claims disagreement with a reanalysis of all the existing SLAC data. Lipniacka (1991) supports the view that (scalar) diquarks can explain these EMC and SLAC data, as well as the high-p_T proton production data from the CERN SFM (split-field-magnet) experiment (see Sec. VII.B for details). However, Bosted *et al.* (1992) argue that their new SLAC data from an aluminum target leads to R values that are much smaller than those expected in the model of Ekelin and Fredriksson. This issue obviously needs experimental clarification.

We conclude from these examples that the data existing a decade ago made a good case for models with diquarks as the only source of scale breaking in the structure functions. More recent data, however, give evidence for the perturbative QCD contributions being at least as important in this respect as diquarks or higher twists.

D. Diquark fragmentation

An additional way to learn about diquarks in DIS is to study the fragmentation of a two-quark system when the third quark has been knocked out. This process is best analyzed in neutrino scattering, as a neutrino most likely interacts with a d quark (and not as strongly with a diquark), while an antineutrino interacts with a u quark. The inclusive spectra of hadrons in the backward (c.m.)

region then mirror the fragmentation of a two-quark system. One assumes that quark and diquark fragmentations are universal, i.e., do not depend on the process that knocks out the leading constituent—an assumption that is supported by the data.

Here we discuss those studies of two-quark fragmentation in DIS that illuminate the discussion of diquarks as dynamical objects. One approach is to estimate how often the two quarks end up in the same hadron by comparing the fragmentation functions with those of single quarks (taken from, for example, e^+e^-). Any difference between the two-quark fragmentation function and twice that of a single quark is an indication of a collective diquark effect. When analyzed in several fragmentation models of quarks and diquarks (Sukhatme *et al.*, 1982; Bartl *et al.*, 1982, 1983; Chang *et al.*, 1983; Ito *et al.*, 1985; Noda and Tashiro, 1985; Noda *et al.*, 1985), the data indicate that a uu system, left over from a neutrino-proton collision, stays together in about half of the events. A few authors (Hanna *et al.*, 1981, 1982; Moriyasu and Wolin, 1983) argue in favor of a lower, and even negligible, uu breakup probability. It is more difficult to study a ud system in antineutrino-proton collisions, since the ud can "disappear" into a neutron (see, for example, Chang *et al.*, 1983). An estimate (Chang, 1981) suggests that a ud system stays together more often than a uu system.

A model-independent way of probing differences between a ud and a uu system is to extract the ratio of the two fragmentation functions for $uu \rightarrow$any pion and $ud \rightarrow$any pion (Fredriksson and Larsson, 1983). In the data of Allen *et al.* (1983) this ratio is significantly higher than unity for all values of the Feynman-x scaling variable of the pions, suggesting that a uu system behaves more often like two independent quarks than does a ud system. This, in turn, suggests a dynamical scalar-diquark effect.

One can also compare two-quark with single-quark fragmentation functions in specific cases where they should be almost equal. One example is the fragmentation process diquark$\rightarrow$proton, which should, in principle, resemble quark$\rightarrow$meson. Any difference, once the diquark stays together and ends up in a final baryon, should be due to kinematic mass effects and final-state resonance decays. An accurate study requires a detailed event generator, like the Lund Monte Carlo program (Andersson, Gustafson, Ingelman, and Sjöstrand, 1982; Andersson, Gustafson, and Sjöstrand, 1982; see also Andersson *et al.*, 1977, and the review by Andersson, 1991). Both within this model, which has no particular preexisting diquarks inside nucleons, and in other schemes (see, e.g., Kinoshita *et al.*, 1982a, 1982b), one finds that the fragmentation functions for diquarks$\rightarrow$baryons are different from those for quarks$\rightarrow$mesons. The leading baryon in the first process is, on the average, more energetic than the leading meson in the second. As a consequence, fewer hadrons are produced in diquark jets than in quark jets (Bardadin-Otwinowska *et al.*, 1982). In this respect, diquarks fragment like heavy quarks. Thus the heavy-quark fragmentation formalism of, for example, Peterson *et al.* (1983) also seems appropriate for describing diquarks fragmenting into hadrons.

In general, diquark fragmentation schemes give a consistent description of how a two-quark system fragments into baryons, without necessarily implying a dynamical diquark effect in nucleons. One exception is deep-inelastic muon-proton scattering, which, up to Q^2 values around 100 (GeV)^2, produces significantly more protons than antiprotons in the forward fragmentation region (Aubert *et al.*, 1981). This is hard to understand unless preexisting diquarks are knocked out directly by the muon. The data have been reproduced by Toyoda and Tsai (1988) within a model with a pointlike scalar (ud) diquark. However, Arneodo *et al.* (1987) claim that their data on the correlation between protons produced in the forward and backward regions are consistent with models in which no significant diquark production occurs in the forward direction.

The bulk of data from charged-lepton beams can nevertheless be used for studying diquark fragmentation in the backward region without specifying whether strong diquark correlations exist in nucleons. Here one has to sum over the quarks in the target (for an early model, see Fontannaz *et al.*, 1978). The information on diquark fragmentation so far seems to agree well with the general picture from neutrino data, reflecting the universality of the fragmentation mechanism. Hadron-hadron collisions can also be used in this respect, although fragmentation functions cannot easily be extracted from the data in a model-independent way (see Sec. VII).

Thus the data apparently give evidence for the collective behavior of the two quarks in the target fragmentation region, once a single quark is knocked out of a nucleon. The effect is stronger for a ud than for a uu pair and suggests that there are substantial spin effects in a two-quark system. There are, however, no strong indications in the fragmentation data that this effect is due to preexisting diquarks in nucleons.

Concluding this section on deep-inelastic scattering, we note that the evidence presented so far in favor of diquark effects implies that a ud diquark is substantially smaller than the whole nucleon and its uu or dd pairs. Because the experiments considered in this section probe quarks and diquarks in the form of partons, the data tell us about the properties of small (current) diquarks. Unfortunately, it is not a straightforward task to extract accurately from presently available data the relative importance of diquarks and gluons in deep-inelastic scattering. A lot more theoretical and phenomenological work is obviously needed in this field. We hope that the HERA accelerator will soon produce data at Q^2 values that are high enough to show gluon effects without the interference of diquarks, enabling us to disentangle these two contributions to deep-inelastic scattering. A number of issues related to quark hunting at HERA have been discussed recently (Anselmino *et al.*, 1993).

V. MULTIQUARK SYSTEMS AND NUCLEAR MATTER

Here we review a handful of imaginative ideas about the possible influence of diquarks in various exotic forms of hadronic matter—from dibaryons to supernovas. At present, there is a lack of conclusive experimental results in these areas, but we hope this situation will change in some of them. In addition, none of the models described here contains essentially new ideas about diquarks themselves. Rather, the models take for granted the importance of diquarks, especially scalar diquarks, in baryonic matter.

A. Dibaryons

Only one bound-dibaryon state is known in nature, the deuteron. However, when we discuss dibaryons in a diquark picture, we do not refer to two color-singlet nucleons bound, like the deuteron, by residual strong interactions, but rather to more complicated six-quark configurations in color space, like three diquarks, a diquark and quadriquark, or a quark and a pentaquark. At present, there is no strong experimental evidence for the existence of any such state.

Dibaryons were considered by many authors, with early papers appearing in 1977 and 1978 (Chan and Høgaasen, 1977; Jaffe, 1977; Chan *et al.*, 1978; Lichtenberg *et al.*, 1978; Mulders *et al.*, 1978, 1980). In a number of these models, dibaryons are colored quark clusters, each cluster being a color $\mathbf{3}$, $\bar{\mathbf{3}}$, $\mathbf{6}$, or $\bar{\mathbf{6}}$. When only $\mathbf{3}$ and $\bar{\mathbf{3}}$ clusters are considered, the interaction between clusters can be related to that of quarks in mesons and baryons. In most models, the dibaryons built out of colored clusters have masses appreciably larger than that of the deuteron, so that if they exist, they should appear only as resonances. Such resonances have been analyzed in detail by Konno and collaborators (Konno and Nakamura, 1982; Konno *et al.*, 1987) within a diquark-cluster model. A similar analysis has been presented by Kondratyuk *et al.* (1987).

Using the color-hyperfine Hamiltonian given in Eq. (3.1), Silvestre-Brac and Leandri (1992) have considered the systematics of six-quark states. As we have pointed out, that Hamiltonian is considerably oversimplified, but its use permits the authors to calculate the mass of a six-quark state relative to the sum of the masses of the lightest two baryons with the same total quantum numbers. Such calculations provide little more than an indication of where improved calculations might be interesting.

A six-quark state that has received much theoretical attention is the so-called H dibaryon ($uuddss$), sometimes called a dilambda, proposed by Jaffe (1977). In the H, the quarks can arrange themselves in a configuration with strong attractive color-hyperfine interaction, and Jaffe conjectures that this state might have a mass less than that of two Λ baryons and so be stable against strong decay. Because of the color-hyperfine interaction, diquark correlations should exist in the H.

Jaffe's analysis suffers from the use of an oversimplified Hamiltonian, according to which quarks behave as free particles inside a bag, except for the color-hyperfine interaction. Since Jaffe's original paper, many authors have calculated the mass of the H in various models, some concluding that it is bound (with respect to two Λ's) and others that it is not. The results are evidently quite model dependent. So far, the H has not been observed.

Among the many papers on the H, we mention the following as having interesting features: Fleck *et al.* (1989), using both a nonrelativistic constituent-quark model and a bag model, calculate that the H is bound in the SU(3) flavor limit, but unbound when SU(3) flavor breaking is included. This paper has a good discussion of the problem and references to the earlier literature. Golowich and Sotirelis (1992) have made a particularly careful calculation of the H mass in the bag model, including diquark correlations and terms of order α_s^2. Their best estimate is that the H is bound. We also call attention to two lattice gauge calculations, one by Mackenzie and Thacker (1985), which finds that the H is unbound, and a second by Iwasaki *et al.* (1988), which says that the H is bound. The calculation of the properties of the H is particularly difficult, and we have to conclude that none of the calculations to date is sufficiently reliable to settle the issue of whether it is stable against decay into two Λ's.

A decade ago, Fredriksson and Jändel (1982) proposed the existence of a so-called demon deuteron, a dibaryon made of three scalar (ud) diquarks. This dibaryon was assumed to have a large nuclear collision cross section in order to explain the anomalously large cross sections of some nuclear fragments apparently seen in cosmic rays and high-energy heavy-ion collisions. Such high cross sections of nuclear fragments have not received further experimental confirmation.

Nevertheless, we briefly consider the model, according to which the dibaryon contains three scalar diquarks in P states and has the spin-parity assignment $J^P=0^-$. A three-body system with $J^P=0^-$ cannot have all three diquark constituents in mutual P waves. However, the QCD Lagrangian allows for the gluon field to carry momentum, and, in phenomenological models like the MIT bag model, gluon momentum can be simulated by having the bag itself oscillate relative to one or more of the constituents. This idea was criticized by Lipkin (1982) and Close (1983), who argued that all observed baryon resonances could be explained in the usual quark model, without the help of bag or gluon-field oscillations. A configuration with bag oscillations would therefore be too massive and short-lived to be of experimental interest. Some arguments against this view were presented (Fredriksson and Jändel, 1983). None of the suggestions for detecting the demon deuteron experimentally (Seth *et al.*, 1983; Pal and Dasgupta, 1985) has so far produced positive results. Thus a light exotic state of three (ud) diquarks (and glue) with quantum numbers $J^P=0^-$, although not definitely excluded, is presently lacking experimental support.

B. Diquarks in normal atomic nuclei: Low energy

Any diquark model of nucleons implies the existence of diquarks in normal nuclear matter. There are claims that diquarks help us to understand the detailed behavior of nuclear matter, as probed in both low- and high-energy processes. By low energy we mean the normal nuclear-physics regime, while high energy refers to the DIS regime (see Sec. IV). Because the two approaches depend on different aspects of diquarks, we treat them separately.

In the QCD-based nuclear-structure model of Bleuler and collaborators (Bleuler *et al.*, 1983; Petry *et al.*, 1985; Hofestädt *et al.*, 1987), the nucleus is a generalized MIT bag of quarks. According to the model, quark-quark interactions produce scalar (ud) correlations (diquarks) inside nuclei. The dynamics and the Pauli principle see to it that the number of diquarks matches the number of remaining single quarks. The color forces make the quarks and diquarks join into nucleons.

The generalized bag has a size that "regulates" itself to minimize the total mass. This requirement leads to the correct empirical formula for the radius of a nucleus with A nucleons. The nuclear size corresponds to a vacuum pressure that ensures that there are as many diquarks as single quarks. At higher nuclear densities, one expects more diquarks (see Sec. V.D). The diquarks are essential for the phenomenon of pairing, i.e., for the strong correlations in the nuclear ground state and the characteristic band gap to the lowest excitations. This phenomenon can be described by a theory analogous to the BCS theory of superconductivity. In the approach used by Bleuler *et al.*, the diquarks are equivalent to the Cooper pairs in a superconductor, and several parallels can be drawn between nuclear pairing and superconductivity. For reviews of this pairing model, see Bleuler and Werner (1989) and Petry and Scholl (1989). The quark-quark force needed in this model is argued (Petry *et al.*, 1985) to originate from instantons coupling to quarks, along the lines suggested by 't Hooft (1976). According to 't Hooft, such a force arises naturally in QCD and favors a scalar two-quark configuration. Petry *et al.* find a best-fit value of about 1 fm for the range of this force, i.e., a loosely bound scalar diquark. However, Betman and Laperashvili (1985), using particle-physics data in an analysis of 't Hooft instantons, find a considerably smaller diquark radius of about 0.3 fm.

C. Diquarks in normal atomic nuclei: High energy

The structure functions of partons, in free protons as well as in nucleons in heavier nuclei, are probed in deep-inelastic scattering (DIS), as we have discussed in Sec. IV. Until 1982, it was generally taken for granted that the structure functions measured on a nucleus (mass number A, charge Z) should be related to the individual nucleon structure functions simply as

$$F_2^A(x)=\frac{1}{A}[ZF_2^p(x)+(A-Z)F_2^n(x)] \tag{5.1}$$

and that there would be no fundamental difference between, for example, deuterium and iron targets, apart from nuclear shadowing and Fermi motion effects.

However, the experimental study of the ratio $F_2^{\mu\mathrm{Fe}}/F_2^{\mu\mathrm{D}}$ showed that the structure functions extracted from iron and deuterium targets are considerably different over the entire x range (Aubert *et al.*, 1983c). This so-called EMC effect was confirmed and refined by other experiments, supporting the view that a free nucleon and a nucleon bound in a nucleus are different. Among other explanations (see Barone and Predazzi, 1987), it was suggested (Fredriksson, 1984) that the EMC effect arises from "swelling diquarks." In this context, the entire Q^2 dependence of F_2 at $Q^2 \gg 2$ (GeV)^2 comes from the $(ud)_0$ form factor, which scales in r^2Q^2, where r^2 is the mean-square diquark radius. Hence

$$F_2^{lA}(x,Q^2)=F_2^{lB}(x,kQ^2)\ , \tag{5.2}$$

where $k=r_A^2/r_B^2$, and indices A and B refer to the two nuclei. A diquark in a nucleus is surrounded by the color fields of the constituents of neighboring nucleons, and these disturbances make it less tightly bound than in a free nucleon. It is reasonable to parametrize the data as $F_2^{\mu\mathrm{Fe}}/F_2^{\mu\mathrm{D}}=k^{-\alpha(x)}$, where $k=r_{\mathrm{Fe}}^2/r_{\mathrm{D}}^2$. The function $\alpha(x)$ is obtained from the iron data. Comparison with the EMC data leads to a value $k\simeq 1.2-2$ and the diquark radius in iron being 10–45% larger than in deuterium.

Cleymans and Thews (1985) compare this model to a model in which Q^2 rescales due to a perturbative QCD evolution of structure-function moments in swelling nucleons. After performing a slightly different fit of $\alpha(x)$ to the data, Cleymans and Thews conclude that the QCD-based model is preferred by the data, but that the diquark model is not ruled out.

D. Diquarks in a quark-gluon plasma

The possible relevance of diquarks in a quark-gluon plasma was first suggested by Ekelin and Fredriksson (1985b). Because it is energetically favorable for two quarks in a spin-0 configuration to form a diquark, one should consider the possibility that the deconfined quarks "pair up" in the plasma. If there are fewer single quarks than diquarks in the plasma when it decays, hadronization cannot go entirely through baryon and meson formation, and one expects multiquark states such as dibaryons as an experimental signature. If the diquark radius exceeds the screening distance for quark forces in the plasma, the diquark will decay into two quarks. Thus, at very high temperature, one does not expect diquarks to play a dynamical role. However, there may be plasma conditions when the screening radius is large enough to allow for a diquark component.

Chiral symmetry is broken in normal hadrons. It is not known whether this symmetry is restored in a plasma and, if so, whether the transition occurs at the same point as deconfinement, or whether there exists an intermediate phase of hadronic matter, in the form of a plasma with massive (constituent) quarks and diquarks. Such a plasma structure would be in between normal hadronic matter and a chirally symmetric plasma phase with light (current) quarks and diquarks. If constituent masses, say, $m_q = 250-350$ MeV and $m_D = 500-700$ MeV, are relevant, then diquarks are kinematically favored if they are bound two-quark states, as the scalar ones might be. On the other hand, in the chirally symmetric plasma, such a binding does not occur, and the diquark component is kinematically disfavored ($m_q = 0-10$ MeV and $m_D = 200-300$ MeV). It is nevertheless favored by Bose statistics. Thus diquarks might exist in either type of plasma.

A thermodynamical approach to the quark-diquark-gluon plasma at finite temperature and density has been presented for plasmas with constituent ("high") masses (Ekelin, 1986a) and with current ("low") masses (Ekelin, 1986b). The plasma is considered an ideal relativistic gas of quarks, antiquarks, diquarks, antidiquarks, and gluons in thermal and chemical equilibrium. Then one can calculate the number densities and contributions to the pressure, energy density, and entropy density of the various plasma components as functions of the temperature and the quark chemical potential. By supplementing an evolution relation for the plasma, one obtains these properties as functions of temperature alone. In this approach, interactions in the plasma are represented by a bag pressure and by the processes $(ud) \leftrightarrow u + d$, $g \leftrightarrow q + \bar{q}$, $q \leftrightarrow q + g$, $g \leftrightarrow g + g$, etc. An interesting point for the case of current masses is the possibility of Bose condensation of diquarks in the high-density plasma.

Another approach to the possible importance of diquarks in the quark-gluon plasma is due to Donoghue and Sateesh (1988), who compare the energy of a diquark gas to that of a quark gas at zero temperature for various densities. The interaction between diquarks is modeled by a $\lambda\phi^4$ theory, with the coupling constant λ related to the Δ-nucleon mass difference. Results are given for a range of λ values. Whereas the quark gas is treated as a free Fermi gas, the diquark gas is described by a Gaussian momentum distribution with a width chosen to minimize the total energy. The results indicate that at low densities the diquark gas is energetically favored as compared to the quark gas, whereas at high densities diquark repulsion outweighs the effects of kinematics (constituent-quark binding) and Fermi statistics. Sateesh (1992) has recently presented some phenomenological implications of this model.

On the other hand, Geist (1988) suggests that diquarks might not exist in a plasma because screening effects would break up any close quark pair. This absence of bound diquarks would lead to an experimental signature of deconfinement in the form of a suppression of high-p_T nucleons in heavy-ion collisions. In most diquark models, in fact, the primary source of high-p_T nucleons in hadronic collisions is diquark scattering (see Sec. VII).

It has been speculated that diquarks could be important components in the core of a would-be supernova. Fredriksson (1989) discusses their role for the explosive phase and argues that condensation of diquarks in a quark-diquark plasma might be responsible for the ultimate collapse of the infalling stellar matter. Kastor and Traschen (1991) calculate the influence of diquarks for various stages of the evolution within the approach of Donoghue and Sateesh. Several features that have previously been achieved with nucleons only are reproduced with a stellar core of quarks and diquarks. A similar approach has recently been advocated by Horvath (1992).

We conclude this section by noting that further theoretical work is obviously needed in order to make plausible the idea that diquarks play a role in exotic matter, given that they are important in nucleons. It may be quite some time before relevant experimental data become available, except, possibly, about exotic dibaryons.

VI. ELECTRON-POSITRON ANNIHILATION

A. General features

When an electron and a positron annihilate at high energy (but much below the Z^0 mass) to create a hadronic system, they form a virtual photon, which then decays into hadrons. The hadronic part of the reaction is supposed to take place in two steps. First, a quark-antiquark pair is created electromagnetically from the photon; second, both quarks "fragment" more or less independently into hadrons. This fragmentation consists of repeated steps in which $q\bar{q}$ pairs are formed from the strong QCD field and break apart. When a quark and an antiquark fly apart, the intermediate color field grows in strength until it breaks with the creation of another $q\bar{q}$ pair, etc. At the end, it is energetically favorable for quarks to form stable hadrons or resonances. This chain of events is illustrated by Fig. 1(a).

Analogously, diquarks, or, more generally, any multiquark configuration, can be created either directly in pairs from the virtual photon or from the QCD vacuum during the fragmentation of a quark. The smaller the diquarks, the more probable is their direct creation [Fig. 1(b)]. The rate of pair production of pointlike diquarks depends only on their charge, mass, and spin, while for extended diquarks it is suppressed by the form factors. The fragmentation of directly produced diquarks can be either estimated from other processes or related in some model-dependent way to the fragmentation of quarks (see Sec. IV.D).

In order to analyze in detail the second way of creating diquarks [Fig. 1(c)], one needs a model for quark fragmentation, giving, for instance, the inclusive cross sections of various hadrons. It is often assumed, for simplicity, that the diquark form factors do not directly

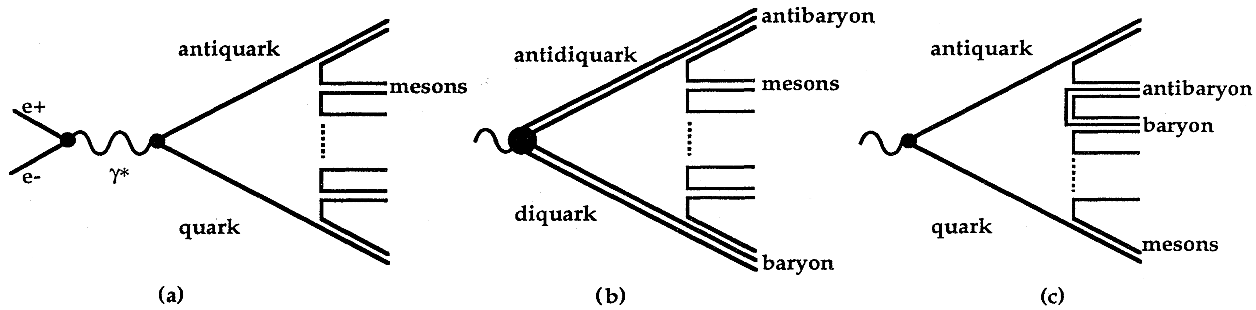

FIG. 1. Dominating parton processes in the one-photon approximation of e^+e^- annihilation into hadrons, according to elementary diquark models: (a) A quark-antiquark pair is created and fragments into mesons. (b) A diquark-antidiquark pair is created and fragments into mesons and baryons. (c) A quark-antiquark pair fragments into mesons and baryons. The baryons result from diquark-antidiquark creation in the quark jets.

suppress the creation of a diquark pair in a quark jet, but that diquarks are rarer than light quarks because of their higher masses. One can take diquark spin into account, and also diquark breakup after formation.

A sum over all the ways a photon can create hadrons should, in principle, be carried out. The corresponding cross sections should then be added to get the inclusive hadron yields. We assume from now on that we need to take into account only quark and diquark pairs (which then fragment into hadrons).

For the various diquark models and their relevance to e^+e^- reactions, we refer the reader to reviews by Mättig (1988), Saxon (1989), and Hofmann (1989).

B. Diquarks from the virtual photon

The first treatment of the process of Fig. 1(b) is due to Pavkovic (1976c, 1976d), although he quotes Bjorken as the first to suggest diquark effects at the parton level. For pointlike diquarks, the lowest-order (one-photon exchange) contribution to the total cross section of $e^+e^- \to$hadrons (Pavkovic, 1976c) are

(i) for scalar (spin-0) diquarks,

$$\sigma(e^+e^-\to S\bar{S})=\frac{\pi e_S^2\alpha^2}{3W^2}\left[1-\frac{4m_S^2}{W^2}\right]^{3/2}; \tag{6.1}$$

(ii) for vector (spin-1) diquarks with zero anomalous magnetic moment,

$$\sigma(e^+e^-\to V\bar{V})=\frac{\pi e_V^2\alpha^2}{3m_V^2}\left[1-\frac{4m_V^2}{W^2}\right]^{1/2}\times\left[1-\frac{m_V^2}{W^2}-\frac{12m_V^4}{W^4}\right]; \tag{6.2}$$

(iii) for a mixed spin-0/spin-1 pair,

$$\sigma(e^+e^-\to S\bar{V},\bar{S}V)=\frac{\pi e_S e_V\alpha^2}{6m_V^2}\left[1-2\frac{m_V^2+m_S^2}{W^2}+\frac{(m_V^2-m_S^2)^2}{W^4}\right]^{3/2}C\,. \tag{6.3}$$

Here α is the fine-structure constant and W is the total e^+e^- c.m. energy. The diquarks have masses m_S, m_V and charges e_S, e_V, and C is an *a priori* unknown constant that reflects the $S\leftrightarrow V$ spin-flip transition. In comparison, for quarks we have

$$\sigma(e^+e^-\to q\bar{q})=\frac{4\pi e_q^2\alpha^2}{3W^2}\left[1-\frac{4m_q^2}{W^2}\right]^{1/2}\left[1+\frac{2m_q^2}{W^2}\right] \tag{6.4}$$

[and the same for $\sigma(e^+e^-\to\mu^+\mu^-)$, provided m_q is replaced by m_μ and e_q^2 by 1].

The total hadronic cross section can be obtained as a sum over quarks and diquarks, including color as well as flavor. It is customary to introduce the ratio

$$R=\frac{\sigma(e^+e^-\to\text{hadrons})}{\sigma(e^+e^-\to\mu^+\mu^-)}\,. \tag{6.5}$$

Because diquarks are not point particles, the above formulas should be multiplied by weight factors equal to the squared diquark electromagnetic form factors, as pointed out by Pavkovic (1976c). This makes comparison with the data complicated, because, as discussed in Sec. IV, vector diquarks have, in general, three different form factors. Neglecting vector diquarks, Pavkovic claims that the low-energy data on R are consistent with the presence of light and almost pointlike scalar diquarks.

The first phenomenological search for direct-diquark contributions in e^+e^- annihilation is due to Meyer (1982). He includes the process of Fig. 1(b) to predict inclusive yields of protons, antiprotons, and Λ's, assuming that diquarks are pointlike and that their production is described by only one parameter. The momentum spectra of baryons at $W=30-34$ GeV are consistent with a diquark yield not exceeding 7.5% of the quark production.

Following Pavkovic, Ekelin *et al.* (1984) assume that only scalar diquarks contribute to e^+e^- into hadrons. The conventional (spacelike) form factor $F_{(ud)_0}(Q^2) = (1+Q^2/Q_0^2)^{-1}$ (see Sec. IV) cannot be continued into the timelike region ($W^2=-Q^2>0$) in a unique way. The authors assume that the timelike diquark form factors remain constant ($=1$) from the production threshold up to some $W^2=Q_0^2$, after which they drop like Q_0^2/W^2. In this scheme, only the (uc) diquark is relevant for the data on R because of its high value of the squared charge (16/9). It contributes a bump in R at $W\simeq 6$ GeV, which the authors claim is visible in the data. Another effect of scalar diquarks in this scheme is the deviation of the angular distribution of the hadronic system from the form $(1+\cos^2\theta)$ typical of spin-$\frac{1}{2}$ constituents. Scalar partons have instead $f_{2\,\text{jet}}(\theta)\propto(1-\cos^2\theta)$. One finds that the (uc) diquark dominates the deviations from the pure quark value. The authors sum all quark and diquark contributions to get the energy dependence of the often-used parameter α in the empirical formula

$$f_{2\,\text{jet}}(\theta)\propto 1+\alpha(W)\cos^2\theta \ . \tag{6.6}$$

A substantial deviation from the naive quark model prediction of $\alpha(W)=1$ is found, especially at $W<10$ GeV, which is claimed to be in agreement with the data. Other effects, such as kinematic mass effects, gluonic processes, effects related to hadronization, and difficulties in identifying jets, could result in $\alpha(W)<1$, but a systematic analysis is still lacking in the normal quark-parton model.

Directly produced diquarks should manifest themselves in the production of fast baryons. However, fast baryons are difficult to distinguish experimentally from mesons. One needs to specify how diquarks fragment into baryons and to estimate the contribution due to diquarks produced in quark jets. A few qualitative predictions (Fredriksson *et al.*, 1983b; Ekelin *et al.*, 1984) for baryon production can be made, particularly in the region 5–10 GeV, where charmed diquarks might appear. Events with fast baryons should have a jet distribution closer to $(1-\cos^2\theta)$ than to that of average events. In addition, a fast baryon should be correlated to a fast back-to-back antibaryon. Any trend in the jet distribution and in the back-to-back correlation, which vanishes at higher energies or lower baryon momenta, would be a signal of directly produced diquarks.

Other evidence of direct diquarks could come from three-jet events in which a scalar and a vector diquark are created [see Eq. (6.3)], after which the vector breaks into two quarks. This process leads to a three-jet event with one diquark jet and two quark jets (Fredriksson *et al.*, 1983b). A fast baryon is produced in the diquark jet and a slower one near the direction of the two quarks. An enhanced baryon yield has, in fact, been observed in three-jet events, but this can also be accounted for in perturbative QCD.

We conclude that small diquarks should reveal themselves as leading partons in e^+e^- annihilation, with weights depending on their charges and form factors. A good candidate seems to be the charmed, scalar (uc) diquark. Production of other diquarks may be suppressed by low charges, high masses, or small effective form factors.

C. Diquarks in quark jets

The occurrence of diquark pairs during the fragmentation of a quark jet was first suggested by Ilgenfritz *et al.* (1978) and Bartl *et al.* (1980), and this process has been widely used by the Lund group (Andersson, Gustafson, and Sjöstrand, 1982) for computing inclusive hadron yields. The Lund group has made successful predictions of the bulk of e^+e^- data, as well as of data from other inclusive high-energy hadronic processes. We shall review only those results that are of interest for testing diquark effects. Several other fragmentation schemes exist, but most of them make similar use of diquarks as the source of baryons (see, e.g., Bartl *et al.*, 1980; Migneron *et al.*, 1982a, 1982b; Itō *et al.*, 1985; Noda and Tashiro, 1985).

The formalism of the Lund model and similar quark fragmentation schemes is an adaption of Schwinger's (1951, 1954) treatment of pair creation in an electromagnetic field and was first applied to strong interactions by Brezin and Itzykson (1970). The relative probabilities W_F, W_B for a parton-antiparton pair to be created in the field are given by

$$W_F \propto \sum_{n=1}^{\infty}\frac{1}{n^2}\exp\left[-\frac{n\pi m_q^2}{\kappa}\right] \tag{6.7}$$

for quark pairs and

$$W_B \propto \frac{1}{2}\sum_{n=1}^{\infty}\frac{1}{n^2}(-1)^{n-1}\exp\left[-\frac{n\pi m_S^2}{\kappa}\right] \tag{6.8}$$

for scalar-diquark pairs. Here κ is the strength of the color field, and $m_q(m_S)$ is the quark (scalar-diquark) current mass. From the quark mass and the field energy ($\kappa\simeq 0.2$ GeV^2), one can fit the diquark masses to the relative baryon yields. This has been done by the Lund group (Andersson, Gustafson, and Sjöstrand, 1982) and by Ekelin *et al.* (1983) in two slightly different diquark approaches.

The Lund group applies Eq. (6.7) to quarks and to scalar and vector diquarks, while Ekelin *et al.* use Eq. (6.8) for scalar diquarks and neglect vector diquarks. Both groups fit the relative inclusive yields of baryons (about 8% summing over all baryons), in order to find best-fit values for the diquark masses. All diquarks are treated as pointlike, and spatial extensions are assumed to reflect themselves in the empirical diquark masses. Diquarks heavier than the vector (ss) have not been considered, because their contributions in Eqs. (6.7) and (6.8) are suppressed by the mass. Charmed baryons, for instance, can be produced only from a leading charmed quark picking up a light diquark or from a directly produced charmed diquark picking up a light quark. The two approaches predict a strong suppression of spin-$\frac{3}{2}$ baryons, in accord with the data (see, e.g., Abachi *et al.*, 1987; Klein *et al.*, 1987a). In the first model this comes about because of the high vector-diquark masses, while in the other it results from the (postulated) absence of vectors.

An application of the quark recombination model of Casher *et al.* (1979), in which two quarks from a diquark end up in different hadrons, gives rise to so-called popcorn events (Andersson, Gustafson, Ingelman, and Sjöstrand, 1982; Andersson, Gustafson, and Sjöstrand, 1985). Such events are illustrated by Fig. 2. Aihara *et al.* (1985) observe that $p\bar{p}$ pairs created in 29-GeV collisions tend to appear on the same side of the quark jet, which in the Lund model suggests that more than 80% of all baryons come from popcorn events. Ekelin and Fredriksson (1986) remark that this could also be explained by a lack of repulsive spin forces between a created scalar diquark and its antidiquark. It should be noted, though (Hofmann, 1989), that the Lund model reproduces the transverse momentum spectra of mesons and baryons, thanks, in part, to the inclusion of frequent popcorn events.

Quark recombination does not account for the data in which a baryon is accompanied by its own antibaryon, for example, $\Lambda\bar{\Lambda}$. It rather favors mixed events, such as $\Lambda\bar{p}$. Accordingly, the relatively low $\Lambda\bar{\Lambda}$ correlation at $W=10.5$ GeV has been interpreted (Albrecht *et al.*, 1989), within the Lund model, as evidence of a popcorn probability of about 70%. On the other hand, the Λ_c inclusive yield at the same energy suggests a conventional diquark production mechanism (Bowcock *et al.*, 1985; Bortoletto *et al.*, 1988). The issue is not yet settled.

It has been argued that the observed Ω^- yield (Klein *et al.*, 1987b) is too large compared with that of a hypothetical "spin-$\frac{1}{2}$ Ω" to be reproduced by any model with suppressed vector diquarks. A counterargument (Fredriksson, 1989) is that, in a model with only scalar

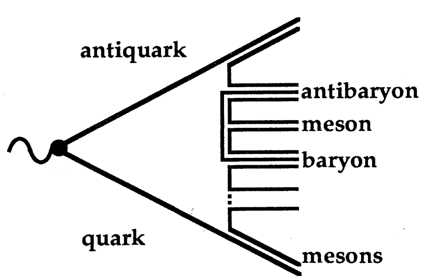

FIG. 2. Popcorn event resulting from diquark breakup after the creation of a diquark-antidiquark pair in a quark jet. The vertical double lines (not labeled) are diquarks.

diquarks, an Ω^- can be created from a decaying $s(sc)$ octet baryon and that a leading scalar (sc) diquark has a chance of surviving at rather high energies, due to the smallness of a system of two heavy quarks. Again, further work is necessary to clear up this disagreement.

In this section, we have compared several diquark models with data from e^+e^- reactions. Although the inclusion of diquarks seems to improve agreement with experiment, no really consistent picture has emerged. Among the open questions are the importance of vector diquarks, of directly produced diquarks, and of so-called popcorn events. The best case for future experimental study seems to be the predictions for diquarks produced directly from the virtual photon. An accurate study of fast baryons at $W<10$ GeV could probe the importance of possible charmed (scalar) diquarks and diquark-related three-jet events.

VII. HADRON-HADRON COLLISIONS

A. General features

The diquark concept plays an important role in the description of hadron-hadron interactions. In collisions involving baryons (mainly protons), two of the valence quarks may interact collectively and contribute to the scattering process as a single entity. Many applications of this idea have been presented, and we summarize here the phenomenological models and the relevant experimental information, for both hard and soft hadronic processes.

In a hard hadronic interaction, $A+B\rightarrow C+X$, the inclusive production of a final particle C with a large transverse momentum p_T is described, according to the parton model, in terms of the elementary scatterings among the initial hadron constituents. Consequently, two of the constituents of hadrons A and B undergo a hard scattering, $a+b\rightarrow c+d$, and then one of the final partons fragments into the observed hadron C. In the first applications of such a scheme (Field and Feynman, 1977), all the elementary constituent interactions are assumed to be of the same empirical form. In QCD, the factorization theorem allows one to justify the above picture, and the interactions of the elementary constituents, quarks and gluons, can be computed in the framework of perturbative QCD (see, e.g., Leader and Predazzi, 1983). The number density of partons inside hadrons and the fragmentation functions of partons into the final hadrons can be derived from other processes—for example, deep-inelastic scattering and e^+e^- annihilation—and their Q^2 evolution is given by perturbative QCD. Diquarks can fit into this scheme as independent partonic constituents of baryons. Indeed, as we shall see, such an idea has been used to explain the abundant large-p_T proton production in pp collisions.

In low-p_T soft hadronic production we cannot use perturbative QCD to compute the overall cross sections in terms of the elementary ones, as we are not in the regime where the running strong-coupling constant is small. However, many QCD-motivated models of soft hadronic production can be found in the literature. The general idea is that during the collision between two hadrons, their constituents, valence quarks and diquarks, become separated and subsequently hadronize, i.e., fragment into a jet of hadrons, more or less independently. The "separation" of the constituents is described by phenomenological distribution functions of quarks and diquarks in the nucleons, and the overall normalization is often a free parameter. Many examples exist in the literature of quark-diquark fragmentation chains, quark-diquark cascade models, quark-diquark fusion, or recombination schemes. The consensus on the hadronization properties of quarks and diquarks is that they obey universality, i.e., are independent of the jet production mechanism and are the same in different processes, both in soft and hard interactions. These include hadron-hadron collisions at small and large transverse momentum, e^+e^- annihilation, and deep-inelastic lepton-nucleon scattering (Fontannaz *et al.*, 1978; Capella *et al.*, 1979, 1980; Göttgens *et al.*, 1981; Bakken *et al.*, 1982; Hanna *et al.*, 1982; Grishin *et al.*, 1984; Baldin *et al.*, 1987).

B. Hard hadronic processes

Data on the production of protons at large transverse momentum p_T and different c.m. polar angles θ in pp and p-nucleus interactions at Fermilab and the CERN ISR (Intersecting Storage Rings) have been presented in a series of papers (Antreasyan *et al.*, 1979; Breakstone *et al.*, 1984b, 1985, 1987a, 1987b; Straub *et al.*, 1992a, 1992b). The yields of protons and antiprotons relative to mesons have been analyzed from the p_T dependence of the ratios $\sigma(p)/\sigma(K^+)$, $\sigma(\bar{p})/\sigma(K^-)$, $\sigma(p)/\sigma(\pi^+)$, $\sigma(\bar{p})/\sigma(\pi^-)$. Moreover, the θ dependence of the proton and antiproton production fractions $R(p) \equiv \sigma(p)/\sigma(\text{all pos.})$ and $R(\bar{p}) \equiv \sigma(\bar{p})/\sigma(\text{all neg.})$, at fixed values of p_T, has been examined.

The strong p_T and θ dependences of $R(p)$ and the large difference from $R(\bar{p})$ are difficult to explain in the usual parton model. If protons and positive mesons are produced via the same elementary scattering mechanism, say, the perturbative QCD production of a large-p_T quark, any difference between proton and meson production must result from the final hadronization process. But one does not expect a strong p_T and θ dependence in the fragmentation chain (Breakstone *et al.*, 1984b, 1985). Therefore, if the production mechanisms for protons and positive mesons were the same, one would expect that $R(p)$ would be approximately constant in all kinematical variables, as observed for $R(\bar{p})$. Furthermore, one would expect that the ratio $\sigma(K^+)/\sigma(\pi^+)$ (Breakstone *et al.*, 1984a) would be comparable in magnitude to $R(\bar{p})$. But these expectations are in contradiction with experiment. A similar situation appears in the large-p_T production of protons in π^-p collisions (Frisch *et al.*, 1983).

Such a drastic difference between proton and positive-meson production at large p_T, both in pp and π^-p collisions, suggests that the underlying elementary processes in the two cases are dynamically different. It turns out (Minakata and Shimizu, 1980; Laperashvili, 1982; Ekelin and Fredriksson, 1984; Larsson, 1984; Breakstone *et al.*, 1985; Kim, 1988) that the effect of diquarks as proton constituents and their active contribution, as single particles, to the scattering process can explain the experimental data. While meson production at large p_T is dominated by the usual hard scattering of quarks and gluons, proton production is due in large part to the scattering of a diquark inside an initial proton by a constituent (quark, gluon, or diquark) in the other hadron.

According to the model of Breakstone *et al.* (1985), the proton has a sizable component of scalar (ud) diquarks. Their interactions with quarks, gluons, and other diquarks is computed in lowest-order perturbative QCD, and their composite, extended nature is taken into account by a form factor. Ekelin and Fredriksson (1984) reach the same results by assuming the proton to be predominantly made of a u quark and a scalar (ud) diquark. Following Field and Feynman (1977), one assumes the quark and diquark elastic interactions with the other constituents to be of the same empirical form. Hence the observed p_T and θ dependences come through the diquark from factor. The same model explains the large-p_T production of protons in π^-p collisions (Larsson, 1984). Kim (1988) makes a similar calculation for pp interactions and includes data from Serpukhov (Abramov *et al.*, 1980, 1985). From an analysis of single and double high-p_T proton production, both in pp and pn collisions, Sulyaev (1989) concludes that, besides scalar diquarks, vector diquarks should also exist inside nucleons.

Further experimental information on pp collisions at the CERN ISR has supported the idea of diquarks as partonic constituents. The observed (but scarce) production of Δ^{++} at large transverse momentum requires the presence of an additional (uu) scattering component, so that both scalar and vector diquarks seem to be required inside a nucleon (Breakstone *et al.*, 1987b). The correlation between the production of large-p_T π^+, K^+ and p at 90° and the production of forward p and Λ^0 favors an effective diquark scattering mechanism (Smith *et al.*, 1987). A collective behavior of two quarks in the large-p_T production of protons at the CERN ISR was already noticed by Drijard *et al.* (1979), who pointed out that the spectator jet originates from single-quark fragmentation only.

Diquarks also play a role in explaining the abundant large-p_T deuteron production observed in 70-GeV pp interactions at Serpukhov (Abramov *et al.*, 1987). Efremov and Kim (1987) show that a good description of the data can be achieved within the quark–scalar-diquark model of the proton from a double quark-diquark scattering at the elementary level. In such a case, two (ud) diquarks are emitted with close momenta, which then form the observed deuteron.

C. Soft hadronic processes

The abundant large-p_T proton production in pp interactions and its explanation in terms of hard scattering of diquarks has prompted further experimental investigations also in soft hadronic processes.

Inclusive hadron distributions in the small-p_T reactions $pp \to$hadron+anything at 360 GeV have been analyzed in the framework of a quark-diquark fragmentation model (Bailly *et al.*, 1986, 1987a, 1987b). According to the model, a quark or a diquark is knocked out of the target and then hadronized into the observed particles. A knowledge of the quark and diquark distributions in the nucleon and their hadronization mechanisms is required for the analysis. The former are parametrized with the help of dual Regge models, and the latter are taken from either the Field-Feynman or Lund fragmentation models, with the parameters obtained from $e^+e^- \to$hadrons. Some primordial motion of the constituents in a proton, the intrinsic transverse momentum, is also taken into account. Both the p_T distributions and the Feynman-x p_T correlations are studied. It turns out that the quark-diquark model gives a better description of the data than a model with three independent quarks. In particular, the model without diquarks fails to reproduce the data on $pp \to \Lambda^0 X$.

Evidence in favor of diquarks is also found in combined analyses of soft and hard processes. Forward proton production in pp collisions at the CERN ISR, accompanied by a large-p_T pion in the same hemisphere, can be interpreted by a mechanism in which a single quark is knocked out of the proton and fragments into the pion, provided the remaining two quarks act as a diquark that eventually ends up in a fast forward nucleon (Beavis and Desai, 1981; Hanna *et al.*, 1981).

Other experiments indicate collective behavior of pairs of quarks. Charmed-meson production in 400-GeV pp interactions at the CERN SPS (Superconducting Proton Synchrotron) shows an abundance of leading D^+ as compared with $\bar{D}^0$. One can explain this by assuming the proton to be composed of a quark and a diquark. The final recombination mechanism then favors the creation of D^+ (Aguillar-Benitez *et al.*, 1988). However, a recent Fermilab experiment (Kodama *et al.*, 1991) finds that hadronization plays a minor role in shaping the kinematical distributions of charmed mesons and that the data are in agreement with perturbative QCD predictions. One reaches a conclusion in favor of quark-diquark configurations by looking at single- and two-particle charged-pion production, in the target fragmentation region, in 70-GeV K^+p interactions (Barth *et al.*, 1981). The comparative study of Σ^+, Σ^-, Δ^{++}, and Λ^0 inclusive productions in pp interactions at Fermilab (Okusawa *et al.*, 1988) is in agreement with the valence quark-diquark picture of the proton. For example, the Feynman-x distributions of Σ^+ and Δ^{++}, whose wave functions have a (uu) pair in common, are remarkably similar. Evidence for diquarks comes also from the investigation of the properties of proton diffractive dissociation in proton-proton interactions at low transverse momentum and high energies, mainly $pp \to pp\pi^+\pi^- m\pi^0$ ($m=0,1,2$) (Asai *et al.*, 1990).

Phenomenological models to describe soft multiparticle production via a quark-diquark fragmentation chain in proton-proton and proton-nucleus interactions have been considered by many authors (Capella *et al.*, 1979, 1980; Capella and Tran Thanh Van, 1980, 1981, 1982; Ranft and Ritter, 1983). Diquarks as single entities appear also in the framework of quark-diquark cascade models to explain small-p_T hadronic production in high-energy interactions (Kinoshita *et al.*, 1980, 1982a, 1982b; Misra *et al.*, 1982; Tashiro *et al.*, 1987) and in a preasymptotic mechanism for nucleon-antinucleon annihilation (Zakharov and Kopeliovich, 1989). Diquarks, in quark-diquark fragmentation chains or quark-diquark fusion or recombination schemes, are often introduced in the description of the hadronic production of baryons, in particular hyperons (see Donnachie, 1980; Shimizu, 1980; Fisjak and Kistenev, 1981; Cooke, 1984; and, essentially, all Monte Carlo–based schemes).

We conclude this section by remarking that the evidence supporting the notion of a collective behavior of pairs of quarks inside a proton is indeed impressive. However, the precise nature of a diquark still remains debatable. Whereas in some hard scattering processes, the diquark seems to behave much like an elementary object, in other instances it can be considered simply as an economical way of dealing with the complicated multiparticle dynamics of hadron interactions. We believe that diquark effects are most important at intermediate momentum transfer and tend to disappear with increasing momentum transfer.

VIII. EXCLUSIVE PROCESSES

A. QCD and exclusive processes

We turn now to yet another application of the idea of a diquark as a single baryonic constituent, with a discussion of exclusive hadronic processes. We consider interactions at momentum transfer Q^2 of the order of a few $(\mathrm{GeV})^2$. For such intermediate Q^2 values, we explore the consequences of the idea that the proton is not (yet) seen as three quarks, but as a quark-diquark state, with the diquark acting as an (almost) elementary constituent. We show that this idea often leads to better agreement with experiment than the usual three-quark description of baryons. We also give applications to small-angle exclusive reactions.

We begin by briefly recalling the usual description of exclusive hadronic interactions in terms of the interactions among constituents in the pure quark perturbative QCD scheme of Brodsky and Farrar (1975), Lepage and Brodsky (1980), Mueller (1981), Chernyak and A. R. Zhitnitsky (1984), and Botts and Sterman (1989). According to this scheme the $A+B \to C+D$ high-energy and large-angle c.m. helicity scattering amplitudes are given by

$$H_{\lambda_C\lambda_D;\lambda_A\lambda_B}(\theta)=\sum_{a,b,c,d;\lambda_a,\lambda_b,\lambda_c,\lambda_d}\int\prod_i[dx_i]\Psi_C^*(x_c)\Psi_D^*(x_d)\hat{H}_{\lambda_c\lambda_d;\lambda_a\lambda_b}(x_a,x_b,x_c,x_d;\theta)\Psi_A(x_a)\Psi_B(x_b)\,, \tag{8.1}$$

where i and λ_i $(i=a,b,c,d)$ denote, respectively, the whole set of constituents of hadron $I(=A,B,C,D)$ and their helicities. The quantity $\hat{H}_{\lambda_c\lambda_d;\lambda_a\lambda_b}(\theta)$ is the helicity amplitude describing the elementary constituent interaction, $a+b\to c+d$. The $\Psi_I(x_i)$ are the spin, momentum, and color hadronic wave functions in terms of the constituents. A hadron I is a collection of n_I partons, each carrying a fraction of its four-momentum, so that, for each i, $[dx]=\prod_{j=1}^{n_I}dx_j\delta(1-\sum_{j=1}^{n_I}x_j)$.

Equation (8.1), with quarks and gluons as constituents, is supposed to hold in the large momentum-transfer limit, $Q^2\to\infty$, where the strong-coupling constant is small and all masses can be neglected. Then the elementary process $a+b\to c+d$ can be computed in lowest-order perturbative QCD. The leading hadronic configurations are those with a minimum number of collinear constituents—in most cases, the valence quarks only. We have not explicitly written here the Q^2 dependence of the hadronic wave functions coming from QCD evolution. In order to evaluate Eq. (8.1), one should know all the elementary helicity amplitudes, including their relative phases, because of possible interference effects. If one does not know these amplitudes and phases, one must either calculate them in an approximate way or make an assumption about their form.

The actual computation of the scattering amplitudes for a physical process $A+B\to C+D$ according to Eq. (8.1) has been carried out only in some simple cases like exclusive $\gamma\gamma$ reactions (Brodsky and Lepage, 1981a; Farrar *et al.*, 1985; Maina, 1990), nucleon Compton scattering (Maina and Farrar, 1988; Farrar and Zhang, 1990), electromagnetic form factors (Lepage and Brodsky, 1979; Chernyak and I. R. Zhitnitsky, 1984), and some charmonium decays (Farrar *et al.*, 1985; Duncan and Mueller, 1980; Chernyak and Zhitnitsky, 1982; Damgaard *et al.*, 1985; Chernyak *et al.*, 1989). For some other processes, like pp elastic scattering, such a task might be prohibitive, due to the large number of elementary diagrams that have to be summed (Farrar and Neri, 1983; Farrar *et al.*, 1985).

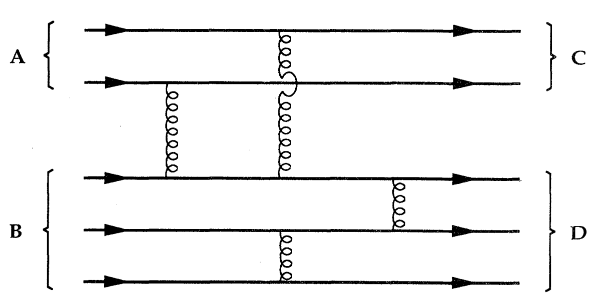

FIG. 3. One of the many Feynman diagrams contributing to the elementary processes involved in meson-baryon large-angle elastic scattering.

Let us consider a typical diagram describing the elementary interaction contributing to, say, meson-baryon elastic scattering, as shown in Fig. 3.

Using dimensional arguments, one can see that the fixed-angle (θ), large-energy ($\sqrt{s}$) c.m. elastic cross section behaves like

$$\frac{d\sigma}{dt}\sim f(\theta)s^{2-n}\,, \tag{8.2}$$

where n is the total number of elementary constituents taking part in the elementary interactions, $n=n_A+n_B+n_C+n_D$. Equation (8.2) reproduces the so-called dimensional counting rules (Brodsky and Farrar, 1973; Matveev *et al.*, 1973). We have neglected the extra logarithmic s dependence coming from the strong-coupling constant and the QCD evolution of the hadronic wave function (Lepage and Brodsky, 1980). Equation (8.2) also agrees with our previous statement that only configurations with a minimum number of constituents contribute; i.e., each extra quark or gluon adds a factor s^{-2} to the right-hand side of Eq. (8.2).

Another remarkable property of Eq. (8.1) is its helicity structure. Because of the gluon-quark coupling, helicity is conserved along each fermion line in the elementary Feynman diagrams (see, e.g., Fig. 3). One also assumes that the hadron helicity equals the sum of the constituent helicities (as is natural for constituents all moving parallel to the parent hadron). Then, for each exclusive reaction $A+B\to C+D$, the sum of the initial helicities equals the sum of the final ones (Brodsky and Lepage, 1981b):

$$\lambda_A+\lambda_B=\lambda_C+\lambda_D\,. \tag{8.3}$$

The helicity-conservation rule, Eq. (8.3), might be broken by terms proportional to m_q/E_q (which allow helicity flips) and by terms proportional to the intrinsic transverse momentum k_T of the quarks inside the hadrons, when the sum of the constituent helicities does not equal the hadron helicity.

Equation (8.2) and the dimensional counting rules are in good agreement with the data (Anderson *et al.*, 1973; Stone, Chanowski, Gray *et al.*, 1977; Jenkins *et al.*, 1978; Stone, Chanowski, Gustafson *et al.*, 1978; Arnold *et al.*, 1986). However, the helicity-conservation rule (8.3) is a source of trouble when its consequences are compared with the existing spin data in exclusive hadronic reactions.

B. Spin problems and the diquark solution

We start by considering pp elastic scattering. Equation (8.3) implies that some helicity amplitudes involving helicity flips, like $H_{++;+-}$ and $H_{++;--}$, must vanish. This has consequences for physical quantities (for the relationship between the scattering amplitudes and the observables, see e.g., Bourrely *et al.*, 1980). For example, it turns out that the proton polarization P is zero:

$$P \equiv \frac{d\sigma^{\uparrow} - d\sigma^{\downarrow}}{d\sigma^{\uparrow} + d\sigma^{\downarrow}} = 0 \; . \tag{8.4}$$

Here $d\sigma^{\uparrow(\downarrow)}$ is the elastic cross section of unpolarized protons resulting in a final proton with spin up ($\uparrow$) or down ($\downarrow$) with respect to the scattering plane (the other final proton spin is not observed).

Analogously, from Eq. (8.3) it follows that the double spin asymmetries A_{NN} and A_{LL} are related, at the c.m. scattering angle $\theta = \pi/2$, by

$$2A_{NN}(\pi/2) - A_{LL}(\pi/2) = 1 \tag{8.5}$$

with

$$A_{NN} \equiv \frac{d\sigma^{\uparrow\uparrow} - d\sigma^{\uparrow\downarrow}}{d\sigma^{\uparrow\uparrow} + d\sigma^{\uparrow\downarrow}} \; , \tag{8.6}$$

where $\uparrow$ and $\downarrow$ are the spins of the initial protons (again, up and down with respect to the scattering plane). The final spins are not observed. The quantity A_{LL} is defined as in Eq. (8.6), with the spin quantized along the direction of the incoming proton. The large-angle pp elastic data at Q^2 of a few $(\mathrm{GeV})^2$ contradict Eqs. (8.4) and (8.5) (Crabb *et al.*, 1978; Crosbie *et al.*, 1981; Auer *et al.*, 1984; Cameron *et al.*, 1985). Particularly surprising is the increase of the proton polarization P with Q^2, while in the free-quark scheme, Eq. (8.1), the expectation is $P = 0$.

Similar considerations hold for the helicity density-matrix elements of the final ρ^- vector meson in the process $\pi^- p \to \rho^- p$:

$$\rho_{\lambda\lambda'}(\rho^-) \sim \sum_{\lambda_p \lambda_{p'}} H_{\lambda\lambda_{p'};\lambda_p} H_{\lambda'\lambda_{p'};\lambda_p} \; . \tag{8.7}$$

Again, Eq. (8.3) predicts all nondiagonal terms of $\rho(\rho^-)$ to be zero:

$$\rho_{\lambda\lambda'} = 0, \quad \lambda \neq \lambda' \; , \tag{8.8}$$

whereas the reported experimental values for the production of ρ^- at $\theta = \pi/2$ in the c.m. system indicate $\rho_{1-1} = 0.32 \pm 0.10$ (Heppelman *et al.*, 1985).

Another example of the problems raised by the rule of helicity conservation comes from η_c and χ_0 decays into $p\bar{p}$, which have been measured (Baglin *et al.*, 1986; Baltrusaitis *et al.*, 1986). Consider the η_c decay. Its quantum numbers, $J^{PC} = 0^{-+}$, together with parity and angular momentum conservation, require $L = S = 0$ for the final $p\bar{p}$ system. On the other hand, if the decay proceeds via quark and hard gluon interactions, the final proton and antiproton must have opposite helicities, which is forbidden in an $S = 0$ state. Thus $\eta_c \not\to p\bar{p}$ in the perturbative QCD scheme of Eq. (8.1), contrary to experimental observation. A similar argument holds for $\chi_0 \to p\bar{p}$.

The previous examples show that the asymptotic perturbative QCD scheme does not adequately account for observed spin effects at intermediate Q^2 values in exclusive reactions. Higher-order and nonperturbative corrections are still important. This has prompted a series of papers attempting to overcome the above difficulties by modifying Eq. (8.1) to include diquarks as baryonic constituents.

The introduction of diquarks as extended elementary partons has two consequences. First, it modifies the dimensional power rules (8.2) by effectively decreasing the number of constituents. For example, for pp elastic scattering, the behavior of the cross section is

$$\frac{d\sigma}{dt} \sim s^{-6} F^4(s) \tag{8.9}$$

instead of s^{-10}, as from Eq. (8.2). At very large values of s, the diquark form factor $F(s)$ is expected to behave as s^{-1}, so that the pure quark result is asymptotically recovered. At intermediate energy values, when $F(s) \simeq 1$, Eq. (8.9) clearly deviates from the quark counting result. The data support this transition from the s^{-6} to the s^{-10} behavior (Anselmino *et al.*, 1987).

A second consequence of diquarks as constituents is the violation of the helicity-conservation rule, Eq. (8.3). Such a violation can only come from couplings between gluons and those partons that allow helicity flips, such as vector diquarks. Again, at very large Q^2 values, if a diquark resolves into two quarks, the helicity-conservation rule is recovered, while at intermediate Q^2 values, where the diquarks act as elementary objects, helicity conservation can be strongly violated and solve the quark model spin problems.

The quark-diquark model for baryons has been applied to several kinds of physical processes. Photon-nucleon processes include proton and neutron electromagnetic form factors, $p - \Delta$ and $p - S_{11}$ transition form factors, Compton scattering ($\gamma p \to \gamma p$), and annihilation ($\gamma\gamma \to p\bar{p}$) (Anselmino *et al.*, 1989; Kroll *et al.*, 1991a, 1991b, 1992). The simultaneous treatment of these processes fixes the parameters of the model, typically contained in the quark-diquark nucleon wave function and in the diquark form factors. Good agreement can be obtained with the data on Compton scattering, the proton magnetic form factor G_M^p, and the $p - \Delta, p - S_{11}$ transition form factors.

In the pure quark model, the predictions are equally good for G_M^p (Chernyak and I. R. Zhitnitsky, 1984) but lie much below the data for the Compton scattering (Farrar and Zhang, 1990). Actually, the three-quark proton wave functions needed to obtain good agreement with the data for G_M^p show a strong asymmetry in the sharing of the proton momentum by the quarks. Such asymmetries follow from QCD sum rules (Chernyak and I. R. Zhitnitsky, 1984), but are also given by diquark clustering (Dziembowski and Franklin, 1990). In addition, the proton electric form factor G_E^p and the neutron form factors G_M^n and G_E^n have been predicted, but the present data do not yet lead to unambiguous conclusions (see, however, Kroll *et al.*, 1991a). The values found for the annihilation $\gamma\gamma \to p\bar{p}$ are smaller than the data, which, however, are taken at a rather low Q^2 value. A sizable proton transverse polarization is predicted in Compton scattering. Such a polarization is zero in the pure quark model, and its experimental measurement would be of great interest.

Other processes considered in the quark-diquark model for the proton are charmonium decays into $p\bar{p}$, such as η_c, $\chi_{0,1,2}\to p\bar{p}$ (Anselmino, Caruso, and Forte, 1991). Vector diquarks as proton constituents allow nonzero decays of $\eta_c, \chi_0 \to p\bar{p}$, contrary to the pure quark scheme in which they are zero. The data give us a large nonzero value for the η_c and, at the moment, only a very large upper bound for the χ_0.

The simultaneous treatment of several processes in the same energy range and in similar physical situations leads to a determination of the parameters of the model, like the relative abundance of scalar and vector diquarks. The quark-diquark model gives good numerical results for $\chi_{1,2}\to p\bar{p}$ decays and similar values for $\chi_0\to p\bar{p}$. The pure quark model also succeeds in describing $\chi_{1,2}\to p\bar{p}$, by making use of the asymmetric three-quark proton wave function mentioned above, but gives zero for $\chi_0\to p\bar{p}$ (Damgaard *et al.*, 1985; Chernyak *et al.*, 1989). However, in the quark-diquark model, the actual value of $\Gamma(\eta_c\to p\bar{p})$, although different from zero, turns out to be still much smaller than the experimental result. The $\eta_c\to p\bar{p}$ decay rate, based admittedly on very few events, seems to be anomalously large.

Several other decay channels of the η_c exhibit unusual, not understood features. This suggests that different decay mechanisms (glueballs?) might be at work there (Anselmino, Caruso, and Murgia, 1990; Anselmino, Genovese, and Predazzi, 1991). The definite observation of the $\chi_0\to p\bar{p}$ decay would be a much stronger indication in favor of diquarks. A simplified diquark model has recently been applied, with success, to the description of $J/\psi\to$baryon-antibaryon decays (Kada and Parisi, 1993). Moreover, the $p\bar{p}$ invariant-mass distribution, in the $J/\psi\to\gamma p\bar{p}$ process, has been computed, both in the framework of pure quark and in quark-diquark models, by Carimalo and Ong (1991). In both cases the results look bad, but better data are needed.

Another application of diquarks in exclusive reactions is $p\bar{p}$ annihilation into hyperons and heavy-flavor baryon and meson pairs, $p\bar{p}\to Y\bar{Y}$, $\Lambda_c\bar{\Lambda}_c$, D^+D^-, etc. (Kroll and Schweiger, 1987, 1989). Results are in good agreement with the data. Particularly relevant are processes requiring the annihilation of two quark-antiquark pairs, like $p\bar{p}\to\Sigma^-\bar{\Sigma}^-$ or $p\bar{p}\to D^+D^-$. In such cases diquarks play a crucial role, via the elementary annihilation of a diquark-antidiquark pair into diquark-antidiquark or $q\bar{q}$ pairs. A similar mechanism is invoked by Klempt (1988) to explain $p\bar{p}$ annihilation at rest into two strange mesons.

We now turn to pp elastic scattering, where spin effects provided early motivation for the introduction of diquarks in exclusive processes. Although the quark-diquark picture is simpler than the pure quark scheme because the former contains a smaller number of constituents, still the actual computation of the helicity amplitudes for $pp\to pp$, in the version of Eq. (8.1) modified by diquarks, is extremely complicated. It has not been carried out except in the end-point model, in which the whole proton momentum is carried by just one constituent, quark or diquark. Such a computation shows how a good description of pp elastic scattering at intermediate Q^2 values might be achieved in a diquark scheme, and how it is possible to obtain nonzero results for the nondiagonal elements of the helicity density matrix of ρ mesons produced in $\pi p\to\rho p$ processes (Anselmino *et al.*, 1987).

Zakharov (1989) discusses spin effects in nucleon-nucleon elastic scattering at small angles. He models the nucleons as being made of a quark and a scalar (ud) diquark and considers Pomeron and Odderon contributions to the scattering amplitudes in the lowest QCD approximation. (Pomeron and Odderon contributions correspond to two-gluon and three-gluon exchanges, respectively, in charge-conjugation even and odd configurations.) There is a phase difference between these two contributions as well as a nonzero transverse momentum of the quarks in the nucleon, so that the sum of the parton helicities does not equal the nucleon helicity. Therefore Zakharov obtains a sizable value for the polarization in pp and $p\bar{p}$ scatterings, as supported by the data and contrary to a similar computation in the pure quark model.

A diquark spectator model, in the framework of quark additive models, is used by Zrałek *et al.* (1979) to describe the forward meson-baryon strangeness-exchange reactions $K^-p\to M^0+\Lambda^0,\Sigma^0,\Sigma^{*0}$, where M^0 is a neutral nonstrange meson. Although the quark-diquark structure of the baryons is crucial for obtaining agreement with the data, the diquarks are spectators that do not play an active dynamical role.

Diquarks as extended elementary constituents seem to be a useful phenomenological way of modeling higher-order and nonperturbative effects in order to achieve a better description of many hadronic exclusive reactions. The picture that emerges from all the above applications is one in which, at Q^2 values of a few $(\mathrm{GeV})^2$, the proton is essentially a quark-diquark state. The scalar diquarks are more abundant, less extended, and lighter than the vector ones. A small component of spin-1 diquarks, however, is required to explain many observed spin effects, unless one resorts to effects of parton mass and transverse momentum. We conclude this section by pointing out that the treatment of exclusive processes in the framework of constituent models and perturbative QCD is still far from being understood in a unique and well-defined computational scheme.

IX. SUPERSYMMETRY OF MESONS AND BARYONS

A. Diquarks as supersymmetric partners of antiquarks

The supersymmetry that we treat in this review has nothing to do with supersymmetric extensions of the

Standard Model. We discuss here an approximate supersymmetry between certain mesons and baryons which arises because of an underlying approximate dynamical supersymmetry of diquarks and antiquarks in QCD. Not only do the supersymmetric partners of this approximate symmetry differ in spin, but they also differ in baryon number (Hwa and Lam, 1975, 1976).

An antiquark in a meson and a diquark in a baryon both belong to a $\overline{3}$ multiplet of color SU(3), and each interacts with a quark to form an overall color-singlet configuration. It follows from perturbative and lattice QCD that quark-antiquark and quark-diquark static color-electric interactions are the same. In the approximation that the static color interaction is dominant, the transformation replacing an antiquark by a diquark in a hadron does not appreciably affect the hadron's properties. As a consequence, there is an approximate dynamical supersymmetry between a baryon, which is a fermion, and a meson, which is a boson. This supersymmetry follows from the underlying approximate supersymmetry between antiquark and diquark.

An approximate supersymmetry between mesons and baryons was first proposed by Miyazawa (1966) before the formulation of QCD, and he used neither the word "supersymmetry" nor the word "diquark." As far as we know, this is the first application of supersymmetry to particle physics. Robson (1976) uses symmetry considerations to obtain relations between the properties of mesons and baryons, but his symmetry scheme is different from Miyazawa's.

Calculations in lattice QCD show that the static quark-antiquark potential is approximately equal to the static quark-diquark potential (Thacker *et al.*, 1988) and that the wave function of a quark-antiquark pair in a meson is similar to the wave function of a quark-diquark pair in a baryon (Gottlieb, 1985). However, the supersymmetry is broken for at least three reasons.

(1) A diquark and an antiquark have different masses, and therefore, even with only static forces, there are kinematical differences between a bound antiquark-quark system and a bound diquark-quark system.

(2) The interaction between colored particles not only contains a static interaction but also nonstatic spin-dependent and velocity-dependent terms. Because the spin and the mass of a diquark are different from those of a quark, the spin-dependent and velocity-dependent terms are different in the two cases.

(3) According to QCD, a quark is pointlike, and, up to now, only an upper limit to its size has been measured. On the other hand, a diquark is not a point particle and may be nearly of hadronic size. The finite size of the diquark must affect its interaction with a quark, especially when the distance between quark and diquark is not large compared to the size of the diquark.

This third approximation may not be as bad as it seems. Although a current quark is pointlike, a constituent quark, clothed with its sea of pairs and gluons, should, like a diquark, have a size larger than zero, as we have already discussed in Sec. II.

The fundamental multiplet of SU(3/3) supersymmetry contains two flavor triplets, one with spin $\frac{1}{2}$ and the other with spin 0. The spin-$\frac{1}{2}$ particles are quarks, and the spin-0 particles are scalar antidiquarks. Likewise, there is a supersymmetry between a flavor $\overline{3}$ multiplet of antiquarks $(\bar{s},\bar{d},\bar{u})$ and a flavor $\overline{3}$ multiplet of scalar diquarks $(ud-du, us-su, ds-sd)$.

In a second paper, Miyazawa (1968) classifies the vector and pseudoscalar mesons with the octet and decuplet of baryons in a single supermultiplet of the supersymmetry algebra SU(6/21). The fundamental constituents of SU(6/21) are a quark belonging to a **6** of SU(6) and a diquark, belonging to a **21**.

Golowich and Haqq (1981) and Gao and Ho (1982, 1983), apparently unaware of earlier work, also devised quark-diquark supersymmetry schemes. Golowich and Haqq propose that a quark has a scalar supersymmetric partner with the quantum numbers of a scalar diquark. This scalar may be either elementary or composite. In the model, a baryon wave function has a three-quark amplitude and a quark-scalar amplitude. A drawback of this scheme is that the quark-scalar amplitude is an additional parameter which is adjusted to help fit the data. In the work of Gao and Ho, diquarks are introduced explicitly. In their first paper, they concentrate on predicting properties of states composed of diquark-antidiquark pairs. In their second paper, they apply their scheme to multiparticle production.

Catto and Gürsey (1985) show that Miyazawa's scheme can be realized within the framework of QCD if the particle belonging to the **21** is a diquark. They use supersymmetry to explain the fact that Regge trajectories of mesons and baryons are approximately parallel (see Sec. II). They also emphasize that if the supersymmetry is good, then exotic mesons containing two quarks and two antiquarks should also exist, as these exotics are in the same multiplet with mesons and baryons. In a second paper, Catto and Gürsey (1988) reduce the SU(6/21) scheme to a smaller algebra, which does not include the exotic mesons.

Some of the papers we discussed in the previous sections, although not making explicit use of supersymmetry to relate mesons and baryons, use diquark-antiquark analogies to obtain properties of baryons from the properties of mesons. In effect, the authors of these papers are making use of broken supersymmetry between diquark and antiquark.

B. Supersymmetry of hadrons containing heavy and light quarks

The supersymmetry of mesons and baryons containing only light quarks is quite badly broken. For example, the pion and nucleon are in the same supermultiplet of SU(6/21) and so, if the symmetry were unbroken, would have the same mass. However, the masses and other properties of the pion and nucleon are so different that it does not seem useful to classify them in the same supermultiplet.

In hadrons containing one heavy quark, supersymmetry ought to hold to a better approximation than in hadrons containing only light quarks (Lichtenberg, 1990), because in the former case the main contribution to the hadron mass comes from the heavy quark. In addition, if we regard a baryon with a heavy quark as being composed of a light diquark and a heavy quark, we do not need to symmetrize the wave function under the interchange of either of the quarks in the diquark with the external heavy quark. Furthermore, in a meson containing a light antiquark and a heavy quark, strong and electromagnetic quark-antiquark annihilations do not occur, and this makes the decays of such mesons more similar to the corresponding decays in baryons with one heavy quark.

The spectator model of heavy-quark decays also leads to an approximate supersymmetry of mesons and baryons containing one heavy quark. According to the spectator model, the weak decay of a heavy quark is independent of the light quarks around it. Hence a baryon and a meson containing a heavy b quark will have the same lifetime for weak decay. The c quark is not heavy enough for the spectator model to be a good approximation, nor is the spectator model good for strong hadron decays.

A different realization of supersymmetry arises in the heavy-quark formalism of Voloshin and Shifman (1987) and Isgur and Wise (1989). As pointed out by Georgi and Wise (1990), this is really a heavy-color symmetry. Therefore it includes a supersymmetry between QQq baryons and $\bar{Q}q$ mesons. As we have already noted in Sec. II, Savage and Wise (1990) and White and Savage (1991) have exploited this symmetry, although they did not mention supersymmetry explicitly.

X. SUMMARY AND CONCLUSIONS

Our most important goal in this review was to present a summary as complete and succinct as possible of what has been done on the subject of diquarks and, especially, of the state of this subject today, nearly 30 years after this exploration began. A problem, however, is that completeness and succinctness do not go together easily, and the risk is that we have cut too much in order to make the subject concise. One consequence may be that we have not given appropriate credit to all those who deserve it. A second problem is that completeness and succinctness go even less easily together with clarity. The consequence here is that we have covered a number of topics only in very qualitative terms, sometimes only quoting results rather than giving the analysis that led to them.

In this review we have tried as much as possible to avoid the use of heavy mathematical formalism in the hope that the reader will not shy away from the subject. Neither have we produced in figure form any of the published fits of diquark models to experimental data. Occasionally, we might have provided a somewhat oversimplified view of some aspect of the problem, but the alternative would have only been to get involved in too complex calculations.

With the previous cautionary remarks, we hope to have been able to convey in a sufficiently clear way our main message, namely, that diquarks, in their third decade, are alive and well and promise to be with us much longer. Nevertheless, some of the properties of diquarks remain somewhat obscure and others controversial, and their role has not yet been entirely clarified. But one thing is undeniable—that they are very useful phenomenological tools in the low- and intermediate-Q^2 regions where perturbative QCD is not fully operational. Much less obvious is their role in the very-large-Q^2 domain. From the contradictory comments made from time to time concerning the role of diquarks in reproducing large-Q^2 data, we infer either that we have not yet reached a point where the experimental precision is sufficiently good, or that something is not well understood on the theoretical side (or both).

Many questions arise. Among these, we just mention the most intriguing one in the present context: Will perturbative QCD become valid as $Q^2 \to \infty$ or, alternatively stated, will diquark effects really die off as $Q^2 \to \infty$? It may be a long time before this question can be resolved, or it may never be resolved, like the old question "Where is asymptopia?" (which may in fact be related). But we expect that HERA and the new hadronic accelerators, LHC and SSC (if and when built), will open entirely new fields of exploration where, we believe, diquarks will still play an important role in phenomenological analyses.

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