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7. Toward the Bag

We have already mentioned in § 3.4.2 some consequences of quark confinement, in the context of an extremely stylized description of confinement : the boundary condition that the Dirac wavefunctions vanish on a static spherical surface. The static cavity approximation, as it is called, is a principal technical assumption in the formulation of the MIT bag model (Chodos, et al., 1974 ; Chodos, Jaffe, Johnson, and Thorn, 1974 ; De Grand, Jaffe, Johnson, and Kiskis, 1975). Of the bag model itself, which has been extremely influential in hadron spectroscopy, there exist several fine reviews, including those by Johnson (1975, 1977, 1979); De Tar (1980), Hasenfratz and Kuti (1978), Jaffe and Johnson (1977), and Jaffe (1977d, 1979) as well as the summary in Close (1979). A different but related picture of quark confinement, known as the SLAC bag, was put forward by Bardeen, et al. (1975) ; see also Giles (1976). Our interest here is much more restricted : to understand how the mechanism of quark confinement (see Wilson, 1974 ; Nambu, 1976 ; Mandelstam, 1980 ; 't Hooft, 1980 ; Adler, 1981 ; Bander, 1981) thought to operate in QCD may give rise to hadronic bags. In the absence of a compelling argument, I follow the usual practice of giving two incomplete arguments. The heuristic discussions are themselves quite standard, and can be found in similar form in many places, including Kogut and Susskind (1974), Lee (1980) and Gottfried and Weisskopf (1981).

7.1. An Electrostatic Analog

It is typical in field theories that the coupling constant depend upon the distance scale. This dependence can be expressed in terms of a dielectric constant ε . We define

$$\varepsilon(r_0) \equiv 1 \tag{7.1}$$

and write

$$g^2(r) = g^2(r_0)/\varepsilon(r). \tag{7.2}$$

We assert that the implication of asymptotic freedom (Gross and Wilczek, 1973 ; Politzer, 1973 ; see also 't Hooft, 1973ab and Khriplovich, 1969) is that in QCD the effective color charge decreases at short distances and increases at large distances. In other words, the dielectric "constant" will obey

$$\varepsilon(r) > 1, \qquad \text{for} \qquad r < r_0 , \tag{7.3a}$$

$$\varepsilon(r) < 1, \qquad \text{for} \qquad r > r_0 . \tag{7.3b}$$

Indeed, to second order in the strong coupling we may write

$$\varepsilon(r) = \left[1 + \frac{1}{2\pi}\frac{g^2(r_0)}{4\pi}(11 - 2n_f/3)\ln(r/r_0) + \mathcal{O}(g^4)\right]^{-1} \tag{7.4}$$

in QCD, where n_f is the number of active quark flavors.

Let us now consider an idealization based upon electrodynamics. In Quantum Electrodynamics, we choose

$$\varepsilon_{\text{vacuum}} = 1 , \tag{7.5}$$

and can show (e.g. Landau and Lifshitz, 1960) that physical media have $\varepsilon > 1$. The displacement field is

$$\underset{\sim}{D} = \underset{\sim}{E} + 4\pi \underset{\sim}{P} \tag{7.6}$$

and atoms are polarizable with $\underset{\sim}{P}$ parallel to the applied field $\underset{\sim}{E}$, so that $|\underset{\sim}{D}| > |\underset{\sim}{E}|$. Since the dielectric constant is defined through

$$\underset{\sim}{D} = \varepsilon \underset{\sim}{E}$$

in these simple circumstances, we conclude that $\varepsilon > 1$. For a thorough treatment, see Dolgov, Kirzhnits, and Maksimov (1981).

Now let us consider, in contrast to the familiar situation, the possibility of a dielectric medium with

$$\varepsilon_{\text{medium}} = 0 , \tag{7.7}$$

a perfect dia-electric, or at least

$$\varepsilon_{\text{medium}} \ll 1 , \tag{7.8}$$

a very effective dia-electric medium. We can easily show that if a test charge is placed within the medium, a hole will develop around it.

To see this, consider the arrangement depicted in Fig. 34(a), a positive charge distribution ρ_+ placed in the medium. Suppose that a hole is formed. Then because the dielectric constant of the medium is less than unity, the induced charge on the inner surface of the hole will also be positive. The test charge and the induced charge thus repel, and the hole is stable against collapse. In normal QED, the induced charge will be negative, as indicated in Fig. 34(b), and will attract the test charge. The hole is thus unstable against collapse.

The radius R of the hole can be estimated on the basis of energetics. Within the hole the electrical energy W_{in} is finite and independent of the dielectric constant of the medium. The displacement field is radial and hence continuous across the spherical boundary. Thus it is given outside the hole by

$$\underset{\sim}{D}_{out}(r>R) = \hat{\underset{\sim}{r}}\, Q/r^2, \tag{7.9}$$

where Q is the total test charge. The induced charge density on the surface of the hole is

$$\begin{aligned} \sigma_{induced} &= (1-\varepsilon)|\underset{\sim}{D}(R)|/4\pi\varepsilon \\ &= (1-\varepsilon)Q/4\pi\varepsilon R^2, \end{aligned} \tag{7.10}$$

which has the same sign as Q, as earlier asserted. Outside the hole, the electric field is determined by the total interior charge

$$Q + (1-\varepsilon)Q/\varepsilon = Q/\varepsilon, \tag{7.11}$$

so that

$$\underset{\sim}{E}_{out}(r>R) = \hat{\underset{\sim}{r}}Q/\varepsilon r^2 \tag{7.12}$$

The energy stored in electric fields outside the hole is then

$$\begin{aligned} W_{out} &= \frac{1}{8\pi}\int d^3\underset{\sim}{r}\;\; \underset{\sim}{D}_{out}(r)\cdot\underset{\sim}{E}_{out}(r) \\ &= \frac{1}{2}\int_R^\infty r^2 dr\, Q^2/\varepsilon r^4 = Q^2/2\varepsilon R. \end{aligned} \tag{7.13}$$

As the dielectric constant of the medium approaches zero, W_{out} becomes large compared to W_{in}, so that the total electrical energy

$$W_{el} \equiv W_{in} + W_{out} \rightarrow W_{out}, \quad \text{as} \quad \varepsilon\rightarrow 0. \tag{7.14}$$

One must consider as well the energy required to hew such a hole out of the medium. For a hole of macroscopic size, it is reasonable to suppose that

$$W_{hole} = \frac{4\pi R^3}{3}\,v + 4\pi R^2 s + \cdots, \tag{7.15}$$

where v and s are non-negative constants. The total energy of the system,

$$W = W_{el} + W_{hole} \tag{7.16}$$

can now be minimized with respect to R . In the regime where the volume term dominates W_{hole} , the minimum occurs at

$$R = \left(\frac{Q^2}{2\varepsilon} \cdot \frac{1}{4\pi v} \right)^{1/4} \neq 0, \tag{7.17}$$

for which

$$W_{el} \approx \left(\frac{Q^2}{2\varepsilon} \right)^{3/4} (4\pi v) \tag{7.18}$$

and

$$W_{hole} \approx \frac{1}{3} \left(\frac{Q^2}{2\varepsilon} \right)^{3/4} (4\pi v)^{1/4}, \tag{7.19}$$

so that

$$W \approx \frac{4}{3} \left(\frac{Q^2}{2\varepsilon} \right)^{3/4} (4\pi v)^{1/4}. \tag{7.20}$$

Thus, in a very effective dia-electric medium, a test charge will induce a bubble or hole of finite radius. Notice, however, that in the limit of a perfect dia-electric medium

$$W \to \infty \qquad \text{as} \qquad \varepsilon \to 0. \tag{7.21}$$

An isolated charge in a perfect dia-electric thus has infinite energy. This is the promised analog of the argument used in § 5.2.1. to wish away isolated colored objects.

If instead of an isolated charge we place a test dipole within the putative hole in the medium, we can again show that the minimum energy configuration occurs for a hole of finite radius about the test dipole. In this case, however, the field lines need not extend to infinity, so the hole radius remains finite as $\varepsilon \to 0$, and so does the total energy of the system. The analogy between the exclusion of chromoelectric flux from the QCD vacuum and the exclusion of magnetic flux from a superconductor is now obvious. To separate the dipole charges to $\pm\infty$ requires an infinite amount of work, as shown in the previous example. This is the would-be analog of quark confinement. For a recent attempt to deduce an effective dia-electric theory from QCD, see Nielsen and Patkós (1981).

Two issues arise in this line of reasoning. One is the question of quark (or as we have phrased it here, charge) confinement. The other is what form does the sourceless QCD vacuum take if it is analogous to a perfect, or very effective, dia-electric medium ? Is the QCD vacuum unstable against the formation of domains containing dipole pairs in the electrostatic model, corresponding to gluons in color-singlet spin-singlet configurations ?

7.2. A String Analog

Suppose, as discussed in § 5.2.2 , that color-electric flux lines are squeezed into a flux tube. This effect can be parametrized by the statement that a region of space of volume

$$V = \sigma \cdot r \tag{7.22}$$

containing color-electric flux contributes a term

$$W_{bag} = B \cdot V = B\sigma r \tag{7.23}$$

to the total energy of the world, where B is a positive constant. The effect of the "bag pressure" B will be to compress the flux lines as much as possible.

The region of color-electric field emanating from a source of charge Q contains an energy density $E^2/8\pi$, where the electric field strength is

$$E = 4\pi Q/\sigma \tag{7.24}$$

if the flux lines are confined within an area σ . The energy stored in the field is thus

$$W_{field} = \frac{E^2}{8\pi} V = \left(\frac{4\pi Q}{\sigma}\right)^2 \cdot \frac{\sigma r}{8\pi}$$

$$= 2\pi Q^2 r/\sigma. \tag{7.25}$$

The total bag plus field energy is

$$W = W_{bag} + W_{field} = \left(B\sigma + \frac{2\pi Q^2}{\sigma}\right) r, \tag{7.26}$$

which can be minimized with respect to the area σ , whereupon

$$\sigma_0 = Q(2\pi/B)^{1/2}. \tag{7.27}$$

At this minimum, the energy density per unit length is

$$k = B\sigma_0 + 2\pi Q^2/\sigma_0 = 4\pi Q^2/\sigma_0 \; . \tag{7.28}$$

For a quark-antiquark pair, the replacement

$$Q^2 \rightarrow 4\alpha_s/3 \tag{7.29}$$

(compare (5.31a)) leads to

$$k = 16\pi\alpha_s/3\sigma_0 \tag{7.30}$$

or

$$\sigma_0 = 16\pi\alpha_s/3k \equiv \pi d^2 . \tag{7.31}$$

Recalling from (5.21) that

$$k^{-1/2} \approx \tfrac{1}{2} \text{ fm} \tag{7.32}$$

we find that

$$d \approx 2 \text{ fm } (\alpha_s/3)^{1/2} . \tag{7.33}$$

For a strong coupling constant $\alpha_s \approx 1$, the radius of the flux tube is

$$d \approx 1 \text{ fm}, \tag{7.34}$$

a reasonable hadronic dimension. We have therefore contrived a situation in which a flux tube of finite radius is a stable configuration. It remains to show that this situation actually obtains in QCD.

Both in the present discussion and in § 5.3.2. we have neglected quark masses. Their inclusion is interesting as a matter of principle and is of some practical importance for particles composed of heavy quarks. Within the framework of an extended bag, the problem has been addressed by Johnson and Thorn (1976) and by Johnson and Nohl (1976) ; see also Chodos and Thorn (1974). Their work suggests that the Regge trajectories of particles composed of massive quarks should be shallow at low spins, but should approach a universal slope as $J \to \infty$. Some evidence for the first half of this statement was noted in § 4.1 , in connection with Fig. 19.

If chromoelectric confinement is indeed the origin of the string picture, we also gain an understanding of the equality of the Regge slopes of the mesons and baryons, which is apparent from Figs. 19 and 20. In an elongated bag, both mesons :

$$q \text{ ———————— } \bar{q} \tag{7.35}$$

and baryons :

$$q \text{ ———————— } (qq) \tag{7.36}$$

are $[\underset{\sim}{3}] - [\underset{\sim}{3}^*]$ color configurations. They must therefore have the same chromoelectric flux density, hence the same amount of stored energy per unit area, hence the same Regge slope. It will thus be of considerable interest to learn the Regge trajectories of baryons containing several strange quarks or a heavy quark.

7.3. Quark Nonconfinement ?

If we assume in view of the heuristic arguments reviewed above that unbroken QCD is indeed a confining theory, how might we accommodate the observation of free quarks ? At first sight it seems straightforward to consider a spontaneously broken color symmetry which endows gluons with small masses and permits quark liberation. This has been explored by De Rujula, Giles, and Jaffe (1978), for example. Georgi (1980) has countered that a small mass term in the Lagrangian need not, in the face of strong quantum corrections, lead to a spontaneous symmetry breakdown. This possibility is open to discussion (De Rujula, Giles, and Jaffe, 1980). Okun and Shifman (1981) have argued that this style of partial confinement is incompatible with the known evidence for asymptotic freedom and with the absence of fractionally-charged hadrons. A different pattern of spontaneous symmetry breaking has been advocated by Slansky, Goldman, and Shaw (1981). Evidently the experimental search for fractionally-charged matter and the theoretical search for proofs or evasions of confinement are research topics of no little importance.

8. Glueballs and Related Topics

The possibility of quarkless states, composed entirely of gluons, would seem to be unique to a non-Abelian field theory such as QCD—as opposed to the elementary quark model. In this short introduction to glueballs I shall try to explore the four important questions :

i/ Should glueballs exist in QCD ?

ii/ What are their properties ?

iii/ How can they be found ?

iv/ Are they found in nature ?

Since most of what we believe to be the solution to QCD is abstracted from the elementary quark model, and because the quark model provides no guidance for quarkless states, the answers given to all of these questions will be partial and frustratingly vague. In the course of explaining these partial answers, one naturally encounters some other issues of significance : violations of the Zweig rule, deviations from ideal mixing, and the continuing problem of the pseudoscalar masses. A common thread will be seen to run through all these topics, and to tie them to the properties of glueballs.

The search for quarkless states has become intense, and several candidate states have appeared. I am not prepared to endorse any of these claims, at least not yet, but I shall have a little bit to say about the experimental situation. This will include some general and specific suggestions for experimental studies.

The subject of glueballs is a newly active one, which remains to be distinctly defined by experimental observations and by theoretical predictions of greater clarity. The modest aim of this Section is merely to underline the importance of the topic, and to introduce some of the issues involved. As for multiquark states, understanding the role of quarkless states in hadron spectroscopy remains in the future.

8.1. The Idea of Glueballs

If color is confined, color singlet states composed entirely of glue may exist as isolated hadron resonances. This is in essence the argument for the existence of quarkless states, as emphasized quite early by Fritzsch and Gell-Mann (1972). If one assumes the existence of gluons, the gauge interaction among gluons, and the confinement of color, this conclusion cannot easily be challenged. After the recognition of asymptotic freedom and the increasingly explicit formulation of QCD (among them Fritzsch, Gell-Mann, and Leutwyler, 1973 ; Gross and Wilczek, 1973b ; Weinberg, 1973), many authors have analyzed, in one or another framework, the possibilities for glueballs. A partial bibliography includes the papers by Freund and Nambu (1975), Fritzsch and Minkowski (1975), Bolzan, Palmer, and Pinsky (1976), Jaffe and Johnson (1976), Willemsen (1976), Kogut, Sinclaïr, and Susskind (1976), Veneziano (1976), Robson (1977), Roy and Walsh (1978), Koller and Walsh (1978), Ishikawa (1979ab), Bjorken (1979, 1980), Novikov, et al. (1979, 1980abcd, 1981), Zakharov (1980ab), Suura (1980), Donoghue (1980,1981), Roy (1979,1980), Soni (1980), Berg (1980), Coyne et al. (1980), Carlson, et al. (1980, 1981), Bhanot and Rebbi (1981), Bhanot (1981), Shifman (1981), Barnes (1981).

Bjorken (1979, 1980) has emphasized the apparent inevitability of color singlet, flavor singlet multigluon states within QCD. In pure (sourceless) QCD, with no fermions, the existence of glueballs follows at once from our assumptions stated above. This may be argued in any of the pictures we have discussed before. A "most attractive channel" analysis is implicit in the work of Barnes (1981). Bag arguments, of the sort given in § 7, lead to the conclusion that the color-singlet configuration is energetically favored, whereas colored states require infinite energy. The string picture of § 5.2.2 is also easily transplanted, with gluon sources replacing quark sources and glueballs replacing $(q\bar{q})$ mesons. The larger flux density between octet sources (cf. Table 12) than between triplets implies flatter Regge trajectories and hence a smaller level density for glueballs than for $(q\bar{q})$ states. In pure QCD there will be among the glueballs a lightest glueball state which, it is reasonable to expect, must be stable.

The introduction of massive quarks (stage II of Bjorken, 1980) does little but provide new sources of glueball production. Quarkonium states may now decay according to

$$^1S_0\,(Q\bar{Q}) \to gg, \qquad (8.1)$$

a colorless, $J^{PC} = 0^{-+}$ final state, and

$$^3S_1(Q\bar{Q}) \rightarrow \begin{cases} ggg, & (8.2a) \\ gg\gamma & (8.2b) \end{cases}$$

colorless hadronic states with $J^{PC} = 1^{--}$ for the three-gluon semifinal state and $J^{PC} = 0^{++}, 0^{-+}, 2^{++}$, etc. for the $gg\gamma$ semifinal state. Again the lightest gluon will be stable, because all $(Q\bar{Q})$ states are— by assumption— extremely massive.

Extending QCD to the light-quark sector raises two questions that go directly to the heart of the matter : what is the mass scale for glueballs, and how prominently will they appear in the spectrum of hadrons ? Given the small mass of the pion it is essentially a certainty that the lightest glueball will be unstable. We must then ask whether the quarkless states will become so broad as to be lost in a general continuum, whether they will mix so strongly with $(q\bar{q})$ and $(q\bar{q}g)$ states as to lose their identity, or whether they will remain relatively pure glue states of modest width. Until definite theoretical predictions can be given, we may conclude only that the observation of glueballs would support the notion that gluons exist and interact among themselves. Not finding glueballs, at the present level of understanding of QCD, has a less obvious significance.

8.2. The Properties of Glueballs

Some characteristics of quarkless states such as their flavor properties are unambiguous, but many others including masses and decay widths are predicted rather indecisively. It is reasonable to attempt to enumerate few-gluon states by analogy with Landau's (1948) classification of two-photon states, which incorporates the restrictions of Yang's (1950) theorem. This has been done by Fritzsch and Minkowski (1975), Barnes (1981), and within the bag model by Donoghue (1980). A pair of massless vector particles can be combined to yield states with

$$J^{PC} = \begin{cases} (\text{even} \geq 0)^{++} , & (8.3a) \\ (\text{even} \geq 0)^{-+} , & (8.3b) \\ (\text{even} \geq 2)^{++} , & (8.3c) \\ (\text{odd} \geq 3)^{++} . & (8.3d) \end{cases}$$

Many papers (e.g. Robson, 1977 ; Coyne, et al., 1980) treat the gluons as massive vectors and arrive at longer lists of two-gluon states. Similarly, extra states may arise in the bag model unless spurious modes associated with the empty bag are eliminated (Donoghue, Johnson, and Li, 1981). The lowest-lying two-gluon configurations should therefore include $J^{PC} = 0^{++}, 2^{++}, 0^{-+}$, and 2^{-+} states.

A variety of estimates of varying degrees of sophistication have been made for the masses of these states. Keeping in mind that the scalar ground state has precisely the quantum numbers of the vacuum and may therefore be appreciably mixed or even subsumed into the vacuum, let us list some representative predictions. The bag model (Donoghue, 1981) suggests that

$$M(0^{++}) \approx M(2^{++}) \approx 1 \text{ GeV}/c^2, \tag{8.4}$$

neglecting hyperfine effects, and that

$$M(0^{-+}) \approx M(2^{-+}) \approx 1.3 \text{ GeV}/c^2, \tag{8.5}$$

again neglecting hyperfine effects. The QCD sum rules of the ITEP Group (Novikov, et al., 1979, 1980abcd, 1981 ; Zakharov, 1980ab ; Shifman, 1981) lead to slightly larger values :

$$M(0^{++}) \approx M(2^{++}) \approx 1.2\text{–}1.4 \text{ GeV}/c^2, \tag{8.6}$$

and

$$M(0^{-+}) \approx 2\text{–}2.5 \text{ GeV}/c^2, \tag{8.7}$$

but with an important gluon component in η' (958). The effective potential calculations of Suura (1980) and Barnes (1981) lead to degenerate pseudoscalar and scalar states, with masses supposed to be on the order of 1–2 GeV/c^2. Barnes (1981) concludes that

$$M(2^{++})/M(0^{++}) \approx 1.8, \tag{8.8}$$

with his description of hyperfine forces.

Three-gluon bound states are more complicated to analyze, especially in terms of the dynamics. It will suffice to give one estimate (Donoghue, 1981) of the masses, obtained in the bag model upon neglecting spin-spin forces :

$$M(0^{++}) \approx M(1^{+-}) \approx 1.45 \text{ GeV}/c^2, \qquad (8.9)$$

and

$$M(0^{-+}) \approx M(1^{-+}) \approx M(1^{--}) \approx M(2^{-+}) \approx M(2^{--}) \approx 1.8 \text{ GeV}/c^2. \qquad (8.10)$$

The general conclusion is that a host of states are to be expected, and that it is plausible that many exist in the region between 1 and 2 GeV/c^2. However all calculations have at least some degree of arbitrariness in the overall mass scale.

A simple lattice argument has also been presented (Kogut, Sinclair, and Susskind, 1976) for glueball masses in the 1-2 GeV/c^2 region. Figure 35(a) shows the minimal lattice configuration for a meson : a single link. On the other hand, on a rectangular lattice the minimal quarkless state consists of a closed loop made up of four links, as shown in Fig. 35(b). Consequently one may suppose that the mass of a typical ground-state glueball is approximately four times the mass of a typical ground-state meson and thus on the order of 1-2 GeV/c^2.

With respect to quantum numbers let us note that apart from the 1^{-+} level, which cannot occur as a $(q\bar{q})$ state, all of the glue states resemble ordinary mesons. Their distinctive property is that pure glue states must be flavor singlets. With this restriction, the allowed decay modes follow from standard selection rules, although branching ratios may be strongly influenced by phase space effects and by the preëminence of quasi-two body final states.

It is quite possible, as we shall now discuss, that glueballs may be narrower structures than $(q\bar{q})$ mesons of comparable mass. This suspicion is tied up with the validity of the so-called Zweig rule and the mixing of glueballs with ordinary mesons, to which we now turn.

8.3. Gluons and the Zweig Rule

In § 3.3.2 we concluded on the basis of simple mass formulas, that φ (1019) is essentially a pure $(s\bar{s})$ state. This conclusion is sustained by an examination of the decay modes of φ, which are collected in Table 21. The total width is

$$\Gamma(\varphi) = 4.1 \pm 0.2 \quad \text{MeV}. \qquad (8.11)$$

Decays into $K\bar{K}$ are inhibited by the limited phase space available, and the relative rates for the charged and neutral final states are understood in terms of p-wave kinematics. For the suppression of the 3π $(\pi\rho)$ mode, however, a dynamical explanation must be sought.

The suppression of nonstrange decays can be accounted for, if not explained from first principles, by the rule (Okubo, 1963 ; Zweig, 1964ab ; Iizuka, Okado, and Shito, 1966 ; Iizuka, 1966) that decays which correspond to connected quark-line diagrams are allowed, but those which correspond to disconnected diagrams are not. This is made concrete for the case of φ decay in Fig. 36. The dissociation and subsequent dressing of the $(s\bar{s})$ pair is allowed (a), but the quarkless semifinal state reached by $(s\bar{s})$ annihilation is not. One may attribute the small observed rate for $\varphi \to 3\pi$ either to light-quark impurities in the φ wavefunction or to violations of the Zweig rule.

Additional evidence in favor of the rule comes from the remarkable metastability of $\psi(3097)$, for which (Particle Data Group, 1980)

$$\Gamma(\psi \to \text{hadrons}) = 45 \text{ keV}, \tag{8.12}$$

and of Υ (9433), for which (Schamberger, 1981)

$$\Gamma(\Upsilon \to \text{hadrons}) = 28 \text{ keV}. \tag{8.13}$$

The Zweig rule thus provides a notable mnemonic for forbidden decays. It is of interest to ask whether there is a dynamical basis for the rule within QCD, and whether there may be other manifestations of violations of the rule.

To this end, recall the outstanding failure of our description of meson masses : the problem of the η and η' masses. In the language of singlet and octet mixing we found it possible in § 3.3.2 to parametrize $M(\eta)$ and $M(\eta')$ in terms of two free parameters : a flavor-singlet mass M_1 and a mixing angle θ . The resulting wave-functions imply relations between decay and reaction rates that are imperfectly respected by the data, as noted in § 5.5.1. In the other-wise successful quark language we were not able to understand the $\pi^0-\eta-\eta'$ splitting or the high mass of the η' . If we interpret the failure as pertaining only to isoscalar states, it is sensible to consider the possibility that virtual annihilations into glue states

$$q\bar{q} \rightarrow \text{glue} \rightarrow q'\bar{q}' \tag{8.14}$$

may influence the masses of $(q\bar{q})$ states. (See among others De Rujula, Georgi, and Glashow, 1975, Isgur, 1976). Such transitions of course cannot affect flavored states.

In the $u\bar{u}$, $d\bar{d}$, $s\bar{s}$ basis the mass matrix of the pseudoscalar mesons can be written as

$$\mathcal{M} = \begin{pmatrix} 2m_u - 3\delta M_{c.m.} + A & A & A \\ A & 2m_u - 3\delta M_{c.m.} + A & A \\ A & A & 2m_s - 3(m_u/m_s)^2 \delta M_{c.m.} + A \end{pmatrix} \tag{8.15}$$

in the notation of eqns. (5.52, 53, 57, 58), where A represents the flavor-independent amplitude for the process (8.14). Recognizing that virtual annihilations cannot affect the π^0 mass, we recast (8.15) in a basis of $(u\bar{u} \mp d\bar{d})/\sqrt{2}$, $s\bar{s}$ as

$$\mathcal{M}=\begin{pmatrix} 2m_u-3\delta M_{c.m.} & 0 & 0 \\ 0 & 2m_u-3\delta M_{c.m.}\ 2A & \sqrt{2}\,A \\ 0 & \sqrt{2}\,A & 2m_s-3(m_u/m_s)^2\delta M_{c.m.}+A \end{pmatrix}, \tag{8.16}$$

which retains the expected result $M(\pi^{\pm})=M(\pi^{0})$.

The remaining two-by-two isoscalar mass matrix suggests a common origin for the $\pi^0-\eta-\eta'$ splitting and the deviations from ideal mixing. With the parameters of § 5.3., the sum of η and η' masses is reproduced with the choice

$$A = 198\ \text{MeV}/c^2, \tag{8.17}$$

for which

$$M(\eta) = 408\ \text{MeV}/c^2 \tag{8.18a}$$

and

$$M(\eta') = 1109\ \text{MeV}/c^2. \tag{8.18b}$$

The wavefunctions implied are

$$|\eta'\rangle = 0.44\left(\frac{u\bar{u}+d\bar{d}}{\sqrt{2}}\right) + 0.81\ s\bar{s}, \tag{8.19a}$$

and

$$|\eta\rangle = 0.81\left(\frac{u\bar{u}+d\bar{d}}{\sqrt{2}}\right) - 0.44\ s\bar{s}, \tag{8.19b}$$

which are similar to those (3.112) given by the linear mass formula in the singlet-octet picture. The masses (8.18) are considerably improved over those produced in § 5.3 , but they are still not perfect. At any rate, we have succeeded in raising the η-η' center of gravity by invoking virtual annihilations, and have thus been able to begin to reconcile the constituent picture with the symmetry approach. Note also that if physical glue states do exist, the mass matrix must be enlarged and the mixing pattern may be considerably more complicated.

The success of our earlier description of the vector meson masses argues that no appreciable annihilation is required there. For heavy mesons such as ψ and Υ , the analogy with ortho- and para-positronium seems apt. A coupling constant argument then suggests that in the asymptotically free regime

$$\Gamma(^3S_1 \to \text{glue})/\Gamma(^1S_0 \to \text{glue}) = \alpha_s \times \text{numerical factors} \qquad (8.20)$$

A power of small coupling constant may inhibit mixing of vector states with gluons in quarkonium, but this is a tenuous argument for the light mesons. In the following § 9 we shall review an argument in favor of the Zweig rule that does not depend upon powers of the coupling constant.

Among the orbitally-excited mesons, there is also room for virtual annihilations. One should in general be alert for the possibility whenever a breakdown of ideal mixing is signalled by the nondegeneracy of the isovector and would-be $(u\bar{u}+d\bar{d})/\sqrt{2}$ states. For a recent look at the 1^{++} and 2^{++} nonets, see Schnitzer (1981ab).

If virtual transitions such as (8.14) occur, they may account for violations of the Zweig rule. This mechanism for Zweig-rule violations naturally suggests the pattern

$$\Gamma(Q\bar{Q} \to \text{hadrons}) < \Gamma(\text{glue} \to \text{hadrons}) < \Gamma(q\bar{q} \to \text{hadrons}), \quad (8.21)$$

where the Zweig-inhibited quarkonium decay rate is of order A^2, the decay rate of a glueball into light quarks is of order A^1, and the rate for Zweig-allowed dissociation of a light quark pair is of order A^0. The possibility therefore exists that a pure glue state will be relatively more stable than a light-quark meson of comparable mass.

8.4. Searching for Glueballs

As strongly-interacting particles, glueballs should be produced routinely in hadronic collisions, where they may be sought out using the techniques of traditional meson spectroscopy. Special kinematic selections may enhance the glueball signal over ordinary mesonic background. An obvious choice with the CERN $p\bar{p}$ collider at hand is an investigation of "Double-Pomeron events", which yield hadronic states in the central region of rapidity with vacuum quantum numbers. If, as Freund and Nambu (1975) and others have suggested, there is a deep connection between the Pomeron and quarkless states, such a selection may be of more than merely kinematical benefit.

Another favorable situation may be in the decays of heavy quarkonium according to (8.2b), leading to transitions of the form

$$\psi \to \gamma + G \qquad (8.21)$$

where G denotes a glueball. This is not only a case in which the general arguments of § 8.1 lead us to expect that glueballs may be produced, but also one that permits inclusive as well as exclusive searches and lends itself to comparison with

$$\psi \to (\omega, \varphi) \text{ + anything} \qquad (8.22)$$

in which the anything is presumably composed of quarks.

Interest in quarkonium decay has been increased by the recent observation of suggestive structures in

$$\psi \to \gamma \text{ + hadrons}. \qquad (8.23)$$

According to Scharre (1981), there is evidence in the Crystal Ball Experiment at SPEAR for two new states. The first, named iota (1440) is seen in the cascade decay

$$\begin{aligned}\psi \to \gamma + \iota&(1440)\\ &\hookrightarrow K\bar{K}\pi \neq KK^{*}\end{aligned} \qquad (8.24)$$

with a mass $M_i = 1440^{+10}_{-15}$ MeV/c^2 and a width of $\Gamma_i = 50^{+30}_{-20}$ MeV. The state has $J^{PC} = 0^{-+}$ and the combined branching ratio for the cascade is

$$B(\psi \to \gamma i)\, B(i \to K\bar{K}\pi) \approx 4 \times 10^{-3}. \tag{8.25}$$

The second is seen in

$$\begin{array}{l} \psi \to \gamma + \theta(1640) \\ \qquad\qquad \hookrightarrow \eta\eta , \end{array} \tag{8.26}$$

with a mass M_θ = 1650 ± 50 MeV/c^2 and a width of $\Gamma_\theta = 220^{+100}_{-70}$ MeV. The decay angular distribution favors $J^{PC} = 2^{++}$, and the product of branching ratios is

$$B(\psi \to \gamma\theta)\, B(\theta \to \eta\eta) \approx 5 \times 10^{-4}. \tag{8.27}$$

An upper limit exists for the decay of θ into $\pi^0 \pi^0$:

$$B(\theta \to \pi\pi) \lesssim B(\theta \to \eta\eta) . \tag{8.28}$$

We have seen above that scalar, pseudoscalar, and tensor glueballs are to be expected in this mass range. In addition, an analysis (Billoire, et al., 1979) of the spin-parity content of the gluon pair in (8.2b) suggests that 2^{++} formation is favored with equal but smaller probabilities for 0^{++} and 0^{-+} configurations. At the same time, radial excitations of the low-lying mesons are to be expected in precisely this region (Cohen and Lipkin, 1979). Thus it is easily possible that any new states be traditional $(q\bar{q})$ states, or mixtures of $(q\bar{q})$ with glueballs, or other exotic possibilities (Close, 1981), as well as states of pure glue. How can these possibilities be distinguished ?

Without going into details, let us note that Chanowitz (1981), Ishikawa (1981), Donoghue, Johnson, and Li (1981), Lipkin (1981b), and Cho, Cortes, and Pham (1981) have examined the case that $\iota(1440)$ is a glueball. Opinion is divided. Chanowitz (1981) has shown that a large number of seemingly contradictory experiments may be reconciled if, in addition to the 1^{++} E(1420) there is a nearby pseudoscalar state for which $\iota(1440)$ is the obvious candidate. He further argues that $\iota(1440)$ has the characteristics of a glueball, but does not concern himself with the $\eta'(1400)$ - see Table 18. If we accept the spin-parity assignments, then there are at least two isoscalar states around 1400 MeV/c^2. The conclusion that $\iota(1440)$ is pseudoscalar and not axial (as E(1420)) removes a potential embarrassment for the two-gluon glueball interpretation. Lipkin (1981), on the other hand, argues that the absence of an appreciable $\iota \to \eta\pi\pi$ signal is inconsistent with a flavor singlet assignment. Obviously there are many experimental questions to settle, among them the spin-parity assignments and the relationship between η' and ι.

Another obvious test may be available in two-photon reactions

$$e^+e^- \to e^+e^- + \text{hadrons}, \qquad (8.29)$$

at least for the pseudoscalar state (s). A $(q\bar{q})$ state should decay into two photons, whereas a pure glue state should not, in lowest order. This inspires a search for the reactions

$$e^+e^- \rightarrow e^+e^- \, \eta'(1400) \quad \hookrightarrow \eta\pi\pi \tag{8.30a}$$

$$e^+e^- \rightarrow e^+e^- \, \iota(1440) \quad \hookrightarrow K\bar{K}\pi \tag{8.30b}$$

Given an estimate for the two-photon decay rate, standard techniques (described in Quigg, 1980) lead to the two-photon production cross section. The rate for production of a $J^{PC} = 0^{-+}$ $\iota(1440)$ is shown in Fig. 37 under the assumption that

$$\Gamma(\iota \rightarrow \gamma\gamma) = 1 \text{ keV}. \tag{8.31}$$

At the energies accessible at PEP and PETRA, an ample cross section is to be expected.

If a prominent signal is observed, one may conclude that the hadronic state is not an axial vector meson and that it is not dominantly a glueball. But if no signal is found, what then ? I see four possibilities :

i/ the hadron is an axial state,

ii/ the hadron is a glueball,

iii/ the hadron is a $(q\bar{q})$ state with a small width for two-photon decay,

iv/ the hadron is a mixed $(q\bar{q})$-glueball state.

The first and second points are self-evident. The third is more problematical. I believe a two-photon width of 0.1-1 keV is reasonable for a radially-excited pseudoscalar, but I cannot convince myself that this represents the full range of "reasonable" possibilities. If glueballs exist, the fourth possibility seems to me the most reasonable one. It has been studied in some detail by Donoghue, Johnson, and Li (1981), by Rosner (1981b), and by Cho, Cortes, and Pham (1981).

In general we may expect some degree of mixing between nearby (or overlapping) hadrons. The simplest case of one glueball and one $(q\bar{q})$ state can be parametrized as

$$|h_1\rangle = |q\bar{q}\rangle \cos\theta + |G\rangle \sin\theta \,, \tag{8.32a}$$

$$|h_2\rangle = -|q\bar{q}\rangle \sin\theta + |G\rangle \cos\theta \,, \tag{8.32b}$$

in an obvious notation. The decay $\psi \to \gamma$ + glue would then lead to a line shape characteristic of

$$|h_1\rangle \sin\theta + |h_2\rangle \cos\theta \,, \tag{8.33}$$

whereas two-photon collisions would excite

$$|h_1\rangle \cos\theta - |h_2\rangle \sin\theta \,. \tag{8.34}$$

Thus we are led to ask whether, for example, the $f^0(1270)$ seen in the decay $\psi \rightarrow \gamma + f^0$ is identical with that observed in two-photon collisions. To answer such questions requires careful measurements of line shapes and branching ratios in both kinds of reactions, as well as in peripheral and central hadron collisions. There is much to be learned here, but the experimental work called for is demanding and meticulous.

Before leaving the subject of glue, let us note that there may be other manifestations of degrees of freedom beyond those of quark and antiquark. The specific possibility of "vibrational modes" has been raised by Giles and Tye (1977, 1978) and by Buchmüller and Tye (1980). States with constituent gluons ($q\bar{q}g$) states have been examined by Horn and Mandula (1978) and by de Viron and Weyers (1981) ; see also Close (1981). Glue-Bearing baryons (qqq g) have been considered by Bowler, et al. (1980).

9. The Idea of the 1/N Expansion in QCD

The search for small parameters which can play the part of expansion parameters is a central element of the process of approximation and model making that is theoretical physics. In many physical situations, extremes of energy or distance suggest highly accurate and readily improved approximation schemes. In classical electrodynamics the indispensable far-field approximation is applicable when the size of a radiator is negligible compared to the distance between the radiator and receiver. The Born approximation for the scattering of charged-particle beams from atomic electrons is trustworthy for beam energies greatly in excess of the atomic binding energy. In Quantum Chromodynamics, a perturbative treatment (which is to say an expansion in powers of the strong coupling parameter $\alpha_s(Q^2)$) is expected to be reliable when the invariant momentum transfer Q^2 is large compared to a characteristic mass scale denoted by Λ^2.

For the problem of hadron structure, no similar expansion is applicable. All of the relevant energies of the problem are on the order of the naturally occurring scale. In a typical hadron, the separation of the quarks is simply the hadronic size of approximately 1 fm — hardly a regime in which perturbative QCD is likely to make any sense. We may, of course, simply await the day when a very heavy quarkonium family is found, and then happily apply conventional perturbative measures. That insouciant course however leaves untouched the problem of the structure of all the hadrons now known, so other actions are called for.

The strategy of the 1/N expansion is a familiar one. When confronted with a problem we cannot solve, we invent a related problem that we can solve . If this is done adroitly, the new problem will not only be simpler but will also capture the physical essence of the original one. More specifically, the 1/N expansion represents an attempt to introduce a parameter that permits a simplification of the calculation at hand. Problem 24 introduces an elementary example.

For QCD, this simplification is achieved ('t Hooft, 1974ab) by generalizing the color gauge group from $SU(3)_c$ to $SU(N)_c$ and considering the limit in which N becomes very large. Although SU(N) is in general more complicated than SU(3), the hadron structure problem is simplified by two observations :

i/ At any order in the strong coupling constant, some classes diagrams are found to be combinatorially negligible.

ii/ The remaining diagrams have common consequences, in large-N perturbation theory.

This technique does not entirely free us from the constraints of perturbative analysis. Since we shall find, by inspection, that entire classes of combinatorially-favored diagrams have common features to all orders in the coupling constant, we shall have to assume that the content of the theory is accurately represented by the set of all diagrams. For QCD, the reliability of the 1/N expansion is inferred from the fact that $SU(N)_c$ QCD seems to resemble the world we observe. Clear introductions to the method, with allusions to other physical situations, are given by Coleman (1980), and by Witten (1979b, 1980 ab).

The combinatorial analysis of $SU(N)_c$ QCD is most transparent in terms of the double line notation introduced for this purpose by 't Hooft (1974a), which is illustrated in Fig. 38. Several examples will suffice to make the main points.

Consider first the lowest-order vacuum polarization contributions to the gluon propagator, the quark loop illustrated in Fig. 39(a) and the gluon loop pictured in Fig. 39(b), in conventional notation. These are redrawn in the double line notation in Fig. 39 (c,d). For an initial gluon of type $i\bar{j}$, only a single color configuration is possible for the quark loop intermediate state : a quark of color i and an antiquark of color $\bar{j}$. For the gluon loop, however, the index k is free to take on any value 1, 2, ..., N. Thus the gluon loop diagram has a combinatoric factor N associated with it. This illustrates the general rule that gluon loops dominate over quark loops by a factor of N, as $N \rightarrow \infty$.

The presence of the factor N would seem to imply that the gluon loop diagram diverges as $N \rightarrow \infty$. This can be cured by choosing the coupling constant to be $g/\sqrt{N}$, with g fixed as $N \rightarrow \infty$. Then for any value of N, the contribution of the gluon loop goes as

$$\left(\frac{g}{\sqrt{N}}\right)^2 \times N \rightarrow g^2, \tag{9.1}$$

a smooth limit.

That this device solves the divergence problem in general is indicated by an analysis of diagrams with more than one loop. The two-loop diagram depicted in Fig. 40 in (a) standard and (b) two-loop notation is immediately seen to be proportional to

$$\left(\frac{g}{\sqrt{N}}\right)^4 \times N^2 \rightarrow g^4. \tag{9.2}$$

Similarly, the three-loop diagram of Fig. 41 obviously goes as

$$\left(\frac{g}{\sqrt{N}}\right)^6 \times N^3 \longrightarrow g^6. \tag{9.3}$$

The situation is different for nonplanar graphs, however. The simplest such graph is shown in Fig. 42. The double-line notation makes it apparent that this graph contains but a single, tangled color loop, and therefore goes as

$$\left(\frac{g}{\sqrt{N}}\right)^6 \times N \longrightarrow g^6/N^2, \tag{9.4}$$

and is therefore suppressed by $1/N^2$ compared to its planar counterpart at the same order in g^2. It is generally the case that nonplanar graphs are reduced by $1/N^2$, as $N \rightarrow \infty$.

These combinatorial arguments select planar graphs as an important subclass. To evaluate and sum all the graphs thus selected is no trivial task. Instead, we may identify their common features and speculate that these survive confinement. It is possible in this way to establish the following results in the large-N limit :

i/ Mesons are free, stable, and noninteracting. For each allowed combination of J^{PC} and flavor quantum numbers, there are an infinite number of resonances.

ii/ Zweig's rule is exact. Singlet-octet mixing (through virtual annihilations) and meson-glue mixing are suppressed. Mesons are pure $(q\bar{q})$ states, with no quark-antiquark sea.

iii/ Meson-meson bound states, which would include particles with exotic quantum numbers, are absent.

iv/ Meson decay amplitudes are proportional to $1/\sqrt{N}$, so mesons are narrow structures.

v/ The meson-meson elastic scattering amplitude is proportional to $1/N$ and is given, as in Regge theory, by an infinite number of one-meson exchange diagrams.

vi/ Multibody decays of unstable mesons are dominated by resonant, quasi-two body channels whenever they are open. The partial width of an intrinsically k-body final state goes as $1/N^{k-1}$.

vii/ For each allowed J^{PC} there are infinitely many glueball states, with widths of order $1/N^2$. They are thus more stable than $(q\bar{q})$ mesons, interact feebly with $(q\bar{q})$ mesons, and mix only weakly with $(q\bar{q})$ states.

Until QCD is actually solved, we will not know how closely the $N \to \infty$ limit of $SU(N)_c$ resembles the case of interest, which is color-SU(3). The preceding list of large-N results does bear, however, a quite striking resemblance to the world described earlier in these lectures. To the extent that the 1/N expansion faithfully represents the consequences of QCD, much of the foregoing phenomenology is explained, and many of the model approximations are justified.

To see how conclusions (i)-(vii) may be reached, let us consider the 1/N derivation of the Zweig rule. A possible mechanism for the Zweig-forbidden decay of $(q\bar{q})$ state is shown in Fig. 43, the process

$$(q\bar{q}) \longrightarrow gg \longrightarrow q'\bar{q}' \longrightarrow \text{mesons} . \tag{9.5}$$

This is shown in standard notation in Fig. 43(a), and in double-line notation in Fig. 43(b). In the latter case I have tied together the ends of the quark and antiquark lines in mesons to emphasize that the mesons are color singlets. The Zweig-forbidden decay amplitude contains a single color loop. It therefore goes as

$$\left(\frac{g}{\sqrt{N}}\right)^4 \times N \sim g^4/N. \tag{9.6}$$

At the same order in the strong coupling constant, the allowed decay is illustrated in Fig. 44. In the double-line representation, it is seen to contain two color loops. The allowed amplitude is therefore proportional to

$$\left(\frac{g^2}{\sqrt{N}}\right)^4 \times N^2 \sim g^4. \tag{9.7}$$

Thus at each order in perturbation theory, the Zweig-forbidden decay is down by a power of 1/N in amplitude compared with the Zweig-allowed decay. Since this reasoning does not rely upon the smallness of the strong coupling constant, which may well be appropriate for the ψ and Υ families, it is an appealing argument for the inhibition of $\psi \to \rho\pi$. The 1/N expansion has also been applied to the problem of baryon structure by Witten (1979b).

To close this brief section on the 1/N expansion, let us briefly return to the difficulty of understanding the η' mass. A clear statement of the puzzle of the flavor-singlet pseudoscalar meson, which is known as the U(1) problem, was given by Gell-Mann, Oakes, and Renner (1968) and by Weinberg (1975). What seems a promising phenomenological explanation is the influence of virtual states composed of glue alone, as described in § 8.3. A formal solution to the U(1) problem was given by 't Hooft (1976), who argued that the U(1) current has an anomaly which leads to a physical nonconservation of the U(1) charge. This removes the raison d'être for a ninth light pseudoscalar. The relationship between the intuitive and formal approaches was exhibited in the context of the 1/N expansion by Witten (1979a, 1980c), Di Vecchia (1979), and Veneziano (1979b).

10. Regrets

One cannot reach the end of a course such as this without contemplating what might have been, or what should have been. There are a number of subjects that I have been forced by the pressure of time to omit. Here I attempt to make amends by providing a brief bibliography for some of the topics I had hoped to discuss.

10.1. The Masses of Quarks

At various points in the analysis of hadron masses we have had occasion to refer to the effective masses of confined quarks. Several important issues have thus been swept under the rug, or at best talked around. One is how QCD behaves in the limit of vanishing quark masses, for which the Lagrangian will have an exact SU(n) $\otimes$ SU(n) chiral symmetry operating independently on the left-and right-handed parts of the quark fields for the n massless flavors. That this is approximately so in Nature is evidenced by the success of soft-pion theorems (see Adler and Dashen, 1968 ; Renner, 1968 ; Lee, 1972). The Lagrangian will also have the chiral U(1) symmetry which leads to the puzzle of the η' mass dealt with in § 9.

In the limit of zero up-, down-, and strange-quark masses, QCD possesses an octet of exactly conserved axial currents. It is believed that the corresponding chiral symmetry must be spontaneously broken along the lines described by Nambu and

Jona-Lasinio (1961ab). Accordingly, in the world of three massless quark flavors there should be eight massless Goldstone (1961) bosons which we identify with the pseudoscalar octet. See also Nambu (1960). The pattern of the spontaneous breaking of the chiral symmetry of QCD has been discussed from the point of view of the 1/N expansion by Coleman and Witten (1980).

Nonzero values of the quark masses which appear in the Lagrangian are thought to arise from the spontaneous breakdown of the $SU(2) \otimes U(1)$ gauge symmetry of the electroweak interactions by means of the Higgs (1964ab) mechanism (for an elementary discussion, see Quigg, 1981), or through dynamical symmetry breaking (Weinberg, 1976, 1977 ; Susskind, 1979 ; Farhi and Susskind, 1981). It then follows that the π , K , and η are only approximately massless, although they are presumed to retain some memory of their chiral origin. The Lagrangian ("current quark") masses have been studied by Leutwyler (1974ab), Pagels (1975), and Langacker and Pagels (1979), among others.

If the masses of the up and down quarks are not identical — a possibility we have entertained in connection with electromagnetic mass differences of hadrons — there may be a number of observable violations of isospin symmetry. The effect upon $\rho\omega$ mixing was mentioned in passing in § 3.1.3 , and many other applications are discussed by Gross, Treiman, and Wilczek (1979), Isgur, Rubinstein, Schwimmer, and Lipkin (1979), Langacker (1979, 1980), and Shifman, Vainshtein, and Zakharov (1979d).

For recent attempts to understand the chiral nature of the pion within the framework of QCD and confinement, consult Pagels (1979), Pagels and Stokar (1979), Donoghue and Johnson (1980), Goldman and Haymaker (1981), and Haymaker and Goldman (1981).

10.2. Decays and Interactions of Hadrons

Important support for flavor-SU(3) symmetry and for specific multiplet assignments derives from the systematic study of hadron decay rates and hadron-hadron reaction rates. The quark model, with or without specific dynamical assumptions, makes many predictions that are sharper than those of SU(3) alone. Entry to the extensive literature on these subjects may be gained via the lecture notes by Rosner (1981a) and the book by Close (1979).

10.3. QCD Sum Rules

A very different and extremely provocative approach to hadron spectroscopy has been pioneered by a group from the Institute for Theoretical and Experimental Physics in Moscow. I regard my omission of their method of analysis as particularly unfortunate. For the students at Les Houches, although not for posterity, this void was filled by informative seminars by John Bell and Eduardo de Rafael (but see in part Bourrely, Machet, and de Rafael, 1981). A short course is provided by the following articles : Shifman, Vainshtein, and

Zakharov (1979abcd) ; Voloshin (1979) ; Leutwyler (1981) ; Reinders, Rubinstein, and Yazaki (1981) ; Bell and Bertlmann (1981) ; and Ioffe (1981).

10.4. Relation to Other Pictures of Hadrons

Finally, and still more telegraphically, I wish to note a few articles which pertain to other approaches to hadron structure and their connections with the schemes I have discussed. Renormalization group techniques for quarks and strings are reviewed by Kadanoff (1977). The theory of dual models and strings is summarized by Scherk (1975). Parallels between QCD, especially in the $1/N_{color}$ expansion, dual theories, and the Reggeon calculus are drawn by Veneziano (1976, 1979a).

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PROBLEMS

1. Consider bound states composed of fundamental scalar particles (denoted σ). The quantum numbers of σ are $J^{PC} = 0^{++}$. For $(\sigma\sigma)$ composites,

   a) Show that a bound state with angular momentum L (i.e. an orbital excitation) must have quantum numbers

   $$C = (-1)^L ; \quad P = (-1)^L .$$

   b) Allowing for both orbital and radial excitations, construct a schematic mass spectrum of $(\sigma\sigma)$ bound states. Label each state with its quantum numbers J^{PC}.

   c) Now suppose that the fundamental scalars have isospin I. Compute C, P, and G for $(\sigma\sigma)$ bound states, and redo part b).

2. Consider bound states composed of fundamental spin-1/2 particles (denoted f), with isospin = 1/2. For $(f\bar{f})$ composites,

   a) Show that a bound state with angular momentum L must have quantum numbers

   $$C = (-1)^{L+s} ; \quad P = (-1)^{L+1} ; \quad G = (-1)^{L+s+I} ,$$

where s is the spin of the composite system, and I is its isospin.

b) Allowing for both orbital and radial excitations, construct a schematic mass spectrum of ($f\bar{f}$) bound states. Label each state with its quantum numbers J^{PC}.

3. The η -meson (550 MeV/c^2) has quantum numbers $J^{PC} = 0^{-+}$ and isospin zero. Its principal decay modes, and branching fractions, are

$$\gamma\gamma \qquad 38\ \%$$

$$\pi^0\pi^0\pi^0 \qquad 30\ \%$$

$$\pi^+\pi^-\pi^0 \qquad 24\ \% \quad .$$

We wish to understand the surprising competition of photonic and hadronic decay modes. Show that the hadronic decays are isospin-violating. Analyze the $3\pi^0$ and the $\pi^+\pi^-\pi^0$ decays separately. What qualitative explanation can be offered for the relative decay rates ?

4. The permutation group on three objects admits three representations : symmetric (S), antisymmetric (A), and mixed (M). For the first two, the group elements are

Representation \ Element	I	(12)	(13)	(23)	(123)	(132)
S	1	1	1	1	1	1
A	1	−1	−1	−1	1	1

When baryon wavefunctions are constructed from three isospin quarks, $|I=1/2,\ I_z=\pm 1/2\rangle$, the antisymmetric representation cannot be formed. Consider the M representation of I = 1/2 final states, which may be built by first coupling quarks 1 and 2 to isospin 0 or 1, and then coupling the third quark. Use as a basis the two states $|1/2, 1/2\rangle_1$ (symmetric in (12)) and $|1/2, 1/2\rangle_0$ (antisymmetric in (12)). Denote these states as $|1\rangle$ and $|0\rangle$, and use $\begin{pmatrix} |1\rangle \\ |0\rangle \end{pmatrix}$ as a basis vector for M. Find the 2 × 2 matrices representing the action of the permutations listed in the table for the M representation.

5. The flavor symmetry $SU(2)_{\text{isospin}}$ and the rotational symmetry $SU(2)_{\text{spin}}$ may be combined systematically in the group SU(4). In nuclear physics, this symmetry group provides the basis for classification into "Wigner supermultiplets". The fundamental representation of SU(4) is

$$\underset{\sim}{4} \equiv \begin{pmatrix} u_\uparrow \\ d_\uparrow \\ u_\downarrow \\ d_\downarrow \end{pmatrix} .$$

Using the notation $(\underset{\sim}{2I+1},\ \underset{\sim}{2s+1})$ for the isospin × spin decomposition of SU(4) representations, we may write

$$\underset{\sim}{4} = (\underset{\sim}{2}, \underset{\sim}{2}),$$

which shows that the $\underset{\sim}{4}$ of SU(4) transforms as a doublet under isospin rotations and as a doublet under spin rotations.

a) Using the techniques for SU(N) computations developed for example in Chapter 3 of Close (1979) or in Bacry (1967), work out the SU(4) content of the product

$$\underset{\sim}{4} \otimes \underset{\sim}{4} \otimes \underset{\sim}{4} .$$

Characterize each SU(4) representation in the product by its Young tableau, symmetry properties, and dimension.

b) Give the $(\underset{\sim}{2I+1}, \underset{\sim}{2s+1})$ content of each of the SU(4) representations in your expansion of $\underset{\sim}{4} \otimes \underset{\sim}{4} \otimes \underset{\sim}{4}$.

Reference : Lipkin (1966).

6. Compute the magnetic moments of $\Delta^{++}, \Delta^{+}, \Delta^{0}, \Delta^{-}$. Assume that the magnetic moment of a quark is given by

$$\mu_i = e_i \hbar / 2 m_i c$$

where $i = u,d$ and e_i is the quark charge in units of $|e|$. Further assume that $m_u = m_d$.

References : Close (1979), Chapter 4 ; Kokkedee (1969), Chapter 11.

7. Work out the explicit SU(6) wavefunctions for the strange members of the baryon octet :

The expressions will be the analogs of

$$|p\uparrow\rangle = (1/\sqrt{18})(2u_\uparrow d_\downarrow u_\uparrow - u_\downarrow d_\uparrow u_\uparrow - u_\uparrow d_\uparrow u_\downarrow$$
$$- d_\uparrow u_\downarrow u_\uparrow + 2d_\downarrow u_\uparrow u_\uparrow - d_\uparrow u_\uparrow u_\downarrow$$
$$- u_\uparrow u_\downarrow d_\uparrow - u_\downarrow u_\uparrow d_\uparrow + 2u_\uparrow u_\uparrow d_\downarrow)$$

and

$$|n\uparrow\rangle = (-1/\sqrt{18})(2d_\uparrow u_\downarrow d_\uparrow - d_\downarrow u_\uparrow d_\uparrow - d_\uparrow u_\uparrow d_\downarrow$$
$$- u_\uparrow d_\downarrow d_\uparrow + 2u_\downarrow d_\uparrow d_\uparrow - u_\uparrow d_\uparrow d_\downarrow$$
$$- d_\uparrow d_\downarrow u_\uparrow - d_\downarrow d_\uparrow u_\uparrow + 2d_\uparrow d_\uparrow u_\downarrow)$$

8. Using your explicit wavefunctions, express the magnetic moments of the strange baryons in terms of μ_u, μ_d, μ_s.

9. In the SU(3) symmetry limit, the quark magnetic moments are proportional to quark charge :

$$\mu_d = \mu_s = -\tfrac{1}{2}\mu_u .$$

Using the proton moment,

$$\mu_p = 2.793 \text{ n.m.}$$

as input, predict the numerical values of the magnetic moments of the octet baryons. Compare with the measured values given by the Particle Data Group (1980).

Reference for problems 7-9 : Thirring (1966).

10. Consider the weak decay of a Λ -hyperon (with four-momentum p_Λ) into a proton (with four-momentum p) and a π^- (with four-momentum q). In general, the Feynman amplitude for the decay will have vector and axial vector terms. We write the general form for the amplitude as

$$\mathcal{M} = \bar{u}_p(p)(A+B\gamma_5)\gamma_\mu u_\Lambda(p_\Lambda) \cdot q^\mu .$$

Work in the rest-frame of the Λ $(\underset{\sim}{p}_\Lambda = 0)$ and let the proton momentum lie along the $\hat{z}$-direction. Compute the decay angular distributic for a Λ with net polarization P_Λ along an arbitrary direction $\underset{\sim}{\hat{n}}$. Shov that it takes the form

$$\frac{d\sigma}{d\Omega} = \text{constant} \times (1 + \alpha P_\Lambda \, \underset{\sim}{\hat{q}} \cdot \underset{\sim}{\hat{n}}),$$

so a measurement of the decay angular distribution determines αP_Λ. Express the asymmetry parameter α in terms of A, B, M_p, and M_Λ .

11. Now consider the decay of an unpolarized Λ . Show that a measurement of the proton's helicity leads to a determination of the asymmetry parameter α .

References for Problems 10 and 11 : Gasiorowicz (1966), c.33 ; Cronin and Overseth (1963) ; Okun (1965).

12. The magnetic dipole transitions among charmed mesons may be relatively immune from recoil effects, because of the large masses and small mass differences.

a) Neglecting the small phase space difference, and approximating $\mu_c = 0$, calculate the ratio $\Gamma(D^{*0} \to D^0 \gamma)/\Gamma(D^{*+} \to D^+ \gamma)$.

b) Redo your calculation assuming $\mu_c = -\frac{2}{3}\mu_s = 0.41$. Continue to use $\mu_u = -2\mu_d = \frac{2}{3}\mu_p$.

c) Now using the masses, branching ratios, and momenta given in the Particle Data Group (1980) meson table, compare your predictions with experiment. You will need to use isospin invariance for the strong decay amplitudes, and to correct the strong decay rates for phase space differences.

d) Assume the masses of the charmed-strange mesons are F^+ : 2030 MeV/c^2 and F^{*+} = 2140 MeV/c^2. Using μ_c and μ_s as in part b), estimate the absolute width for the decay $F^* \to F\gamma$. What branching ratio do you expect ?

13. Derive the connection between $|\Psi(0)|^2$ and the leptonic decay rate of a $(q\bar{q})$ vector meson . It is convenient to proceed by the following steps :

a) Compute the spin-averaged cross section for the reaction $q\bar{q} \to e^+e^-$. Show that it is

$$\sigma = \frac{\pi \alpha^2 e_q^2}{12E^2} \cdot \frac{\beta_\ell}{\beta_q} (3-\beta_\ell^2)(3-\beta_q^2),$$

where E is the c.m. energy of a quark and β is the speed of a particle.

b) The annihilation rate in a $^3S_1$ vector meson is the density $\times$ relative velocity $\times \frac{4}{3}$ (to undo the spin average) $\times\ \sigma$, or

$$\Gamma = |\Psi(0)|^2 \times 2\beta_q \times \frac{4}{3} \times \sigma.$$

c) How is the result modified if the vector meson wavefunction is

$$|V^0\rangle = \sum_i c_i |q_i \bar{q}_i\rangle \ ?$$

d) Now neglect the lepton mass and the quark binding energy and assume the quarks move nonrelativistically. Show that

$$\Gamma(V^0 \to e^+e^-) = \frac{16\pi\alpha^2}{3M_V^2} |\Psi(0)|^2 (\sum c_i e_i)^2.$$

e) How is the result modified if quarks come in N_c colors and hadrons are color singlets ?

References : Van Royen and Weisskopf (1967) ; Pietschmann and Thirring (1966), Jackson (1976).

14. The Han-Nambu (1965) model is an integer-charge alternative to the fractional-charge quark model, with quark charges assigned as

color \ flavor	u	d	s
R	0	-1	-1
G	1	0	0
B	1	0	0

a) Show that below the threshold for color liberation, the ratio

$$R \equiv \sigma(e^+e^- \to \text{hadrons})/\sigma(e^+e^- \to \mu^+\mu^-)$$

is R = 2, as in the fractional-charge model, and that R = 4 if color can be liberated.

b) Consider the reaction

$$\gamma\gamma \to \text{hadrons} \ ,$$

viewed as $\gamma\gamma \to q\bar{q}$. Show that with fractionally charged quarks

$$\sigma(\gamma\gamma \to \text{hadrons}) \propto \sum_i e_i^4 = \frac{2}{3} \ ,$$

and that in the Han-Nambu model

$$\sigma(\gamma\gamma \to \text{hadrons}) \propto \begin{cases} 2 & \text{below color threshold} \\ 4 & \text{above color threshold} \end{cases}$$

References : Close (1979), c.8 ; Chanowitz (1975) ; Lipkin (1979a) ; see also Okun, Voloshin, and Zakharov (1979).

15. Consider the electromagnetic interaction of two classical charged particles, with charges of q_1 and q_2, masses m_1 and m_2, and positions $\underset{\sim}{r}_1$ and $\underset{\sim}{r}_2$. In the static limit the interaction Lagrangian is the familiar Coulomb Lagrangian,

$$\mathcal{L}_{int;NR} = -q_1 q_2 / r$$

where $\underset{\sim}{r} \equiv \underset{\sim}{r}_1 - \underset{\sim}{r}_2$ is the relative coordinate. Derive the interaction Lagrangian through order $(v/c)^2$, and show that it may be written in the form obtained by Darwin in 1920 :

$$\mathcal{L}_{int} = -\frac{q_1 q_2}{r}\left\{1 - \frac{1}{2c^2}\left[\underset{\sim}{v}_1\cdot\underset{\sim}{v}_2 + (\underset{\sim}{v}_1\cdot\hat{\underset{\sim}{r}})(\underset{\sim}{v}_2\cdot\hat{\underset{\sim}{r}})\right]\right\}.$$

The derivation is most gracefully carried out in the Coulomb gauge.

Reference : Jackson (1975), c. 12.

16. a) Show that the magnetic field due to a classical particle with magnetic dipole moment $\underset{\sim}{\mu}$ at the origin of coordinates is

$$\underset{\sim}{B}(\underset{\sim}{r}) = \frac{8\pi}{3}\underset{\sim}{\mu}\,\delta^3(\underset{\sim}{r}) + \frac{3\hat{\underset{\sim}{r}}(\hat{\underset{\sim}{r}}\cdot\underset{\sim}{\mu}) - \underset{\sim}{\mu}}{r^3}$$

b) Now consider the (classical) interaction of a static nucleus with magnetic moment $\underset{\sim}{\mu}_N$, fixed at the origin, with an electron (with magnetic moment $\underset{\sim}{\mu}_e$ and electric charge e) orbiting about it with angular momentum $\underset{\sim}{L}$. Show that the interaction energy is given by the hyperfine Hamiltonian

$$\mathcal{H}_{HFS} = -\frac{8\pi}{3}\, \boldsymbol{\mu}_e \cdot \boldsymbol{\mu}_N\, \delta^3(\mathbf{r})$$

$$+ \frac{1}{r^3}\left[\boldsymbol{\mu}_e \cdot \boldsymbol{\mu}_N - 3(\hat{\mathbf{r}}\cdot\boldsymbol{\mu}_e)(\hat{\mathbf{r}}\cdot\boldsymbol{\mu}_N) - \frac{e}{mc}\,\mathbf{L}\cdot\boldsymbol{\mu}_N\right].$$

Discuss the origin of each term.

References : Fermi (1930), Jackson (1975), c. 5.

17. The Darwin Lagrangian for two charged particles is given by the interaction Lagrangian $\mathcal{L}_{int}$ of Problem 15 plus the free-particle Lagrangian expanded to order $1/c^2$,

$$\mathcal{L}_{free} = \frac{1}{2}(m_1 v_1^2 + m_2 v_2^2) + \frac{1}{8c^2}(m_1 v_1^4 + m_2 v_2^4).$$

a) Introduce relative coordinates $\mathbf{r} = \mathbf{r}_1 - \mathbf{r}_2$ and $\mathbf{v} = \mathbf{v}_1 - \mathbf{v}_2$ and c.m. coordinates. Write out the Lagrangian $\mathcal{L}_{Darwin} = \mathcal{L}_{free} + \mathcal{L}_{int}$ in the reference frame in which the velocity of the center of mass vanishes and evaluate the canonical momentum components $p_x = \partial\mathcal{L}/\partial v_x$, etc.

b) Compute the Hamiltonian to first order in $1/c^2$ and show that it is

$$\mathcal{H} = \frac{p^2}{2}\left(\frac{1}{m_1} + \frac{1}{m_2}\right) + \frac{q_1 q_2}{r} - \frac{p^4}{8c^2}\left(\frac{1}{m_1^3} + \frac{1}{m_2^3}\right)$$

$$+ \frac{q_1 q_2}{2 m_1 m_2 c^2}\left(\frac{p^2 + (\mathbf{p}\cdot\hat{\mathbf{r}})^2}{r}\right).$$

Compare with the various terms in eqn. (42.1) on p. 193 of Bethe and Salpeter (1957). Discuss the agreements and disagreements.

References : Jackson (1975), problem 12.12 ; Berestetskii, Lifshitz, and Pitaevski (1971), pp. 280-284 ; Breit (1930) ; Heisenberg (1926).

18. By coupling together first the quarks and antiquarks separately, show that the colorspin for a collection of n constituents is given by

$$\sum_{i<j} \langle \underset{\sim}{\lambda}_i \cdot \underset{\sim}{\lambda}_j \, \underset{\sim}{\sigma}_i \cdot \underset{\sim}{\sigma}_j \rangle = -4 \underset{\sim}{G}^{(6)\,2}_{tot} + \frac{4}{3} s_{tot}(s_{tot}+1)$$

$$+ 8 \left[\underset{\sim}{G}^{(6)\,2}_{quarks} + \underset{\sim}{G}^{(6)\,2}_{antiq} \right] - 4 \left[\underset{\sim}{G}^{(3)\,2}_{quarks} + \underset{\sim}{G}^{(3)\,2}_{antiq} \right]$$

$$- \frac{8}{3} \left[s_{quarks}(s_{quarks}+1) + s_{antiq}(s_{antiq}+1) \right] - 8n \, ,$$

where the labels (quarks, antiquarks) refer to the collective representations of quarks and antiquarks. Verify that for a state composed only of quarks you recover (5.81).

19. Consider quark-antiquark states. Using SU(6) techniques, identify the colorspin representations containing color singlets, and compute the expectation value of the colorspin operator. Compare with the results in Table 15 for $0^- - 1^-$ splitting.

20. Enumerate the $SU(6)_{\text{colorspin}}$ representations that can be formed out of two quarks and two antiquarks. Give the $SU(3) \otimes SU(2)$ decomposition of each. Compute $\underset{\sim}{G}^{(6)2}$ for each representation.

Reference for Problems 18-20 : Jaffe (1977ab).

21. Show that quarkonium level spacings independent of the constituent mass occur in a logarithmic potential, $V(r) = \lambda \log(r/r_0)$.

Reference : Quigg and Rosner (1977).

22. Using the Schrödinger equation (6.3), prove the identity

$$\frac{|u'(0)|^2}{4\pi} = |\Psi(0)|^2 = \frac{\mu}{2\pi\hbar^2}\left\langle\frac{dV}{dr}\right\rangle$$

for a system with reduced mass μ .

Reference : This result is apparently due to Fermi and to Schwinger, in unpublished work. A general derivation appears in § 2.2 of Quigg and Rosner (1979).

23. By evaluating the identity just derived in semiclassical approximation, show that for a general nonsingular potential

$$|\Psi_n(0)|^2 \simeq \left(\frac{2\mu}{\hbar^2}\right)^{3/2} \frac{E_n^{1/2}}{4\pi^2}\frac{\partial E_n}{\partial n} .$$

References : Krammer and Léal Ferreira (1976) ; Quigg and Rosner (1978c) ; Bell and Pasupathy (1979).

24. Consider the Schrödinger equation for s-wave bound states of a $1/r$ potential in N space dimensions:

$$[\nabla^2 + 2\mu(E + \alpha/r)]\Psi(r) = 0 \quad . \tag{1}$$

(a) Show that the radial equation is

$$\left[\frac{d^2}{dr^2} + \frac{(N-1)}{r}\frac{d}{dr} + 2\mu(E + \alpha/r)\right]\Psi(r) = 0 \quad . \tag{2}$$

(b) Now take the limit of large N, so that $(N-1) \rightarrow N$. Introduce a reduced radial wavefunction

$$u = r^{N/2}\psi \tag{3}$$

and a scaled radial coordinate

$$R = r/N^2 \quad . \tag{4}$$

Show that the Schrödinger equation becomes

$$\left[\frac{1}{N^2}\frac{d^2u}{dR^2} - \frac{u}{4R^2} + 2\mu(N^2E + \alpha/R)u\right] = 0 \quad . \tag{5}$$

(c) Apart from the factor N^2 which sets the scale of E, this equation describes a particle with effective mass μN^2 moving in an effective potential

$$V_{eff} = \frac{1}{8\mu R^2} - \frac{\alpha}{R} \quad . \tag{6}$$

Find the energy of the ground state in the limit as $N \rightarrow \infty$, for which the kinetic energy vanishes. Show that it is given by the absolute minimum of V_{eff}, so that

$$E_{N\to\infty} = -2\mu\alpha^2/N^2 \quad . \tag{7}$$

Corrections to (7) may be computed by expanding V_{eff} about the minimum and treating the additional terms as perturbations.

(d) The exact solution to the exact eigenvalue problem (2) is easily verified to be

$$E_{exact} = -2\mu\alpha^2/(N-1)^2 \quad . \tag{8}$$

Show that the exact eigenvalue can be recast in the form of an expansion in powers of 1/N as

$$E_{exact} = -\frac{2\mu\alpha^2}{N^2}\sum_{j=1}^{\infty} jN^{1-j}$$

$$= E_{N\to\infty}\left\{1 + \sum_{j=2}^{\infty} jN^{1-j}\right\} ,$$

so that the $N\to\infty$ result may form the basis for a systematic approximation scheme. How many terms must be retained to obtain a 1% approximation for N = 3?

Reference: Mlodinow and Papanicolaou (1980).

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Table 1

Contributions to Electromagnetic Mass Differences in Nonstrange Baryons

Particle	N_d	$\sum_{i<j} e_i e_j$	$\left\langle \sum_{i<j} e_i e_j \underset{\sim}{\sigma}_i \cdot \underset{\sim}{\sigma}_j \right\rangle$
p	1	0	4/3
n	2	-1/3	1
Δ^{++}	0	4/3	4/3
Δ^{+}	1	0	0
Δ^{o}	2	-1/3	-1/3
Δ^{-}	3	1/3	1/3

Table 2

Contributions to Electromagnetic Mass Differences in Nonstrange Mesons

Particle	N_d	$\langle e_q e_{\bar{q}} \rangle$	$\langle e_q e_{\bar{q}} \underset{\sim}{\sigma}_q \cdot \underset{\sim}{\sigma}_{\bar{q}} \rangle$
"η"	1	-5/18	5/6
π^+	1	2/9	-2/3
π^0	1	-5/18	5/6
π^-	1	2/9	-2/3
ω	1	-5/18	-5/18
ρ^+	1	2/9	2/9
ρ^0	1	-5/18	-5/18
ρ^-	1	2/9	2/9
$\langle \rho^0 \| M \| \omega \rangle$	1	1/6	1/6

Table 3

Properties of the Quarks

Quark	I	I_3	S	B	Y=B+S	$Q=I_3+Y/2$
u	1/2	1/2	0	1/3	1/3	2/3
d	1/2	-1/2	0	1/3	1/3	-1/3
s	0	0	-1	1/3	-2/3	-1/3

Table 4

Contributions to Meson Masses for particles containing strange quarks

Particle	N_s	N_d	$\langle e_q e_{\bar{q}} \rangle$	$\langle e_q e_{\bar{q}} \sigma_q \cdot \sigma_{\bar{q}} \rangle$
K^+	1	0	2/9	-2/3
K^o	1	1	-1/9	1/3
$\bar{K}^o$	1	1	-1/9	1/3
K^-	1	0	2/9	-2/3
K^{*+}	1	0	2/9	2/9
K^{*o}	1	1	-1/9	-1/9
$\bar{K}^{*o}$	1	1	-1/9	-1/9
K^{*-}	1	0	2/9	2/9
φ	2	0	-1/9	-1/9

Table 5

Baryon masses (MeV/c^2)

Particle	Model		Experiment	
p	938.28		938.28	
n	939.58		939.57	
Σ^+	1125.95	-7.98	1189.36	-7.98 ± 0.08
Σ^-	1133.93	4.88	1197.34	4.88 ± 0.06
Σ^0	1129.05		1192.46	
Λ	1127.84		1115.60	
Ξ^0	1314.90		1314.9	
Ξ^-	1321.60		1321.62	

Table 6. Magnetic Moments of the Baryon Octet. Numerical values are given in nuclear magnetons. Underlined quantities are inputs.

Baryon	Quark Model	Exact SU(3) $\mu_u = -2\mu_d = -2\mu_s = 2\mu_p/3$	$\mu_u = -2\mu_d$ $\mu_s \equiv \mu_\Lambda$	$\mu_u \neq -2\mu_d$ $\mu_s \equiv \mu_\Lambda$	Measured values
p	$(4\mu_u - \mu_d)/3$	$\mu_p = \underline{2.793}$	$\underline{2.793}$	$\underline{2.793}$	2.793
n	$(4\mu_d - \mu_u)/3$	$-\frac{2}{3}\mu_p = -1.862$	-1.862	$\underline{-1.913}$	-1.913
Λ	μ_s	$-\mu_p/3 = -0.931$	$\underline{-0.614}$	$\underline{-0.6138}$	-0.6138 ± 0.0047 [a] -0.6129 ± 0.0045 [b]
Λ-Σ^0	$(\mu_d - \mu_u)/\sqrt{3}$	$-\mu_p/\sqrt{3} = -1.612$	-1.612	-1.633	$-1.82^{+0.18}_{-0.25}$ [c]
Σ^+	$(4\mu_u - \mu_s)/3$	$\mu_p = 2.793$	2.687	2.673	2.33 ± 0.13 [d]
Σ^0	$(2\mu_u + 2\mu_d - \mu_s)/3$	$\mu_p/3 = 0.931$	0.825	0.791	0.46 ± 0.28 [e]
Σ^-	$(4\mu_d - \mu_s)/3$	$-\mu_p/3 = -0.931$	-1.037	-1.091	-1.41 ± 0.25 [f]
Ξ^0	$(4\mu_s - \mu_u)/3$	$-2\mu_p/3 = -1.862$	-1.439	-1.436	-1.250 ± 0.014
Ξ^-	$(4\mu_s - \mu_d)/3$	$-\mu_p/3 = -0.931$	-0.508	-0.494	-0.75 ± 0.06

a) Schachinger, et al. (1978)

b) Cox, et al. (1981)

c) Dydak, et al. (1977)

d) Settles, et al. (1979) and Particle Data Group (1980)

e) Defined by eq. (3.59)

f) Roberts, et al. (1975) and Particle Data Group (1

g) Handler, et al. (1980)

Table 7

Axial Charges g_A/g_V in Semileptonic Decays of Baryons

Decay	Cabibbo/SU(3)	Quark Model	Experiment	
$n \to pe\nu$	D+F	5/3	1.254 ± 0.007	a)
$\Lambda \to pe\nu$	F+D/3	1	0.62 ± 0.05	a)
			$\pm(0.734 \pm 0.03)$	b)
$\Sigma^- \to \Sigma^0 e\nu$	F	2/3		
$\Sigma^\pm \to \Lambda e\nu$	pure axial	pure axial	$g_V^-/g_A^- = -0.10 \pm 0.22$	a)
$\Sigma^- \to ne\nu$	F-D	-1/3	$\pm(0.385 \pm 0.070)$	a)
$\Xi^- \to \Xi^0 e\nu$	F-D	-1/3		
$\Xi^- \to \Lambda e\nu$	F-D/3	1/3		

a) Particle Data Group (1980)

b) Jensen, et al. (1980)

Table 8. M1 Transitions in Mesons

Process	Photon Energy ω (GeV)	Decay Rate Γ (keV)	$3\pi\Gamma/\omega^3\mu_N^2$	$\lvert\mathcal{M}_{flavor-spin}\rvert^2$ Quark model	SU(3) $3\pi\Gamma/\omega^3\mu_N^2\theta$	θ	$\mu_s \neq \mu_d$ $3\pi\Gamma/\omega^3\mu_N^2$	θ
$\omega\to\pi^0\gamma$	0.380	889 ± 50 a)	5.87 ± 0.35	$(\mu_u-\mu_d)^2$	7.80	0.75 ± 0.04		
		789 ± 92	5.21 ± 0.61			0.67 ± 0.08		
$\rho\to\pi\gamma$	0.375	67 ± 7 b)	0.461 ± 0.038	$(\mu_u+\mu_d)^2$	0.87	0.53 ± 0.04		
$K^{*+}\to K^+\gamma$	0.309	62 ± 14 c)	0.753 ± 0.170	$(\mu_u+\mu_s)^2$	0.87	0.87 ± 0.20	1.56	0.48 ± 0.11
$K^{*0}\to K^0\gamma$	0.309	75 ± 35 a)	0.921 ± 0.422	$(\mu_d+\mu_s)^2$	3.47	0.27 ± 0.12	2.39	0.39 ± 0.18
$\varphi\to\pi^0\gamma$	0.501	5.7 ± 2 a)	0.017 ± 0.006	0	0	-		
$\omega\to\eta\gamma$	0.199	$3^{+2.5}_{-1.8}$ a)	$0.138^{+0.115}_{-0.084}$	Q: $(\mu_u+\mu_d)^2/2$	0.43	$0.32^{+.28}_{-.20}$		
				L: $3(\mu_u+\mu_d)^2/4$	0.65	$0.21^{+.18}_{-.13}$		
$\rho^0\to\eta\gamma$	0.194	50 ± 13 a)	2.48 ± 0.64	Q: $(\mu_u-\mu_d)^2/2$	3.90	0.63 ± 0.17		
				L: $3(\mu_u-\mu_d)^2/4$	5.85	0.42 ± 0.11		
$\varphi\to\eta\gamma$	0.362	62 ± 9 a)	$0.47+0.07$	Q: $2\mu_s^2$	1.73	0.27 ± 0.04	0.75	0.63 ± 0.09
				L: μ_s^2	0.87	0.54 ± 0.08	0.38	1.24 ± 0.18
$\varphi\to\eta'\gamma$	0.060			Q: $2\mu_s^2$	1.73		0.75	
				L: $3\mu_s^2$	2.60		1.13	
$\eta'\to\rho^0\gamma$	0.164	83 ± 30 a,d)	6.81 ± 2.46	Q: $3(\mu_u-\mu_d)^2/2$	11.70	0.58 ± 0.21		
				L: $3(\mu_u-\mu_d)^2/4$	5.85	0.29 ± 0.11		
$\eta'\to\omega\gamma$	0.159	7.6 ± 3 a,d)	0.68 ± 0.27	Q: $3(\mu_u+\mu_d)^2/2$	1.30	0.53 ± 0.21		
				L: $3(\mu_u+\mu_d)^2/4$	0.65	1.05 ± 0.41		

a) Particle Data Group (1980) b) Berg, et al. (1980a) c) Berg, et al. (1980b,1981)

d) Adjusted by Rosner (1980) using total η' width measured by Binnie, et al. (1979) and by Abrams, et al. (1979)

Table 9. Electromagnetic Charge Radii of Mesons

Particle	Beam Momentum (GeV/c)	$\langle r_{EM}^2 \rangle$, fm^2	Reference
π^-	100	0.31 ± 0.04	Dally, et al. (1977)
	250	0.43 ± 0.03	Dally, et al. (1980a)
	Combined fit	0.39 ± 0.04	Dally, et al. (1980a)
K^-	250	0.28 ± 0.05	Tsyganov (1979) ; Dally, et al. (1980c)
	π/K Comparison	0.25 ± 0.05	Dally, et al. (1980b)
K°	30 - 100	-0.054 ± 0.026	Molzon, et al. (1978)

Table 10. Light Mesons as Quark-Antiquark Bound States

State	Mixing ?	J^{PC}	I = 1	I = 0		I = 1/2
^1S_o		0^{-+}	$\pi(140)$	$\eta(549)$	$\eta'(958)$	K(496)
$^3S_1$	$^3D_1$	1^{--}	$\rho(776)$	$\omega(784)$	$\varphi(1019)$	$K^*(892)$
$^1P_1$		1^{+-}	B(1231)	H(1190)[a]		$Q_B(1355)$[b]
^3P_o		0^{++}	$\delta(981)$	$\varepsilon(1300)$?[h]	$S^*(980)$	$\chi(1500)$?
$^3P_1$		1^{++}	$A_1(1240)$[a]	D(1285)	E(1418)	$Q_A(1340)$[b]
$^3P_2$	$^3F_2$	2^{++}	$A_2(1317)$	$f^o(1273)$	$f^*(1516)$	$K^{**}(1430)$
$^1D_2$		2^{-+}	$A_3(1660)$			L(1765)?
$^3D_1$	$^3S_1$	1^{--}	$\rho'(1600)$		$\varphi'(1634)$	$K^*(1650)$?
$^3D_2$		2^{--}				
$^3D_3$	$^3G_3$	3^{--}	g(1700)	$\omega(1670)$	$\varphi_3(1870)$[c]	$K^*(1753)$[d]
$^1F_3$		3^{+-}				
$^3F_2$	$^3P_2$	2^{++}	$f\pi(1700)$[e]		$\theta(1640)$[g]?	
$^3F_3$		3^{++}				
$^3F_4$		4^{++}	$K^+K_s(2060)$[f]	h(2040)		$K^*(2070)$[d]

a) Dankowych, et al. (1981)
b) Leith (1977)
c) Armstrong, et al. (1981)
d) Aston, et al. (1981) ; Cleland et al.(1980b);Dorsaz(1981)
e) Cashmore (1980) : see also Montanet (1980)
f) Cleland, et al. (1980a) ; see also Montanet (1980)
g) Seen in $\psi \to \gamma\eta\eta$; Scharre (1981)
h) According to Wicklund,et al.(1980)the pole lies at 1425MeV,

Table 11

SU(6) Classification of the Baryon Resonances

N	$SU(6)_L^P$	$^{(2J+1)}SU(3)$	Members
0	$\underline{56}_0^+$	$^2[\underline{8}]$	N(939), Λ(1115), Σ(1193), Ξ(1318)
		$^4[\underline{10}]$	Δ(1232), Σ(1385), Ξ(1533), Ω(1672)
1	$\underline{70}_1^-$	$^2[\underline{1}]$	Λ(1405)
		$^4[\underline{1}]$	Λ(1520)
		$^2[\underline{8}]$	N(1535), Λ(1670), Σ(1750), Ξ(1684) ?
		$^2[\underline{8}]$	N(1700), Λ(1870)
		$^4[\underline{8}]$	N(1520), Λ(1690), Σ(1670), Ξ(1820) ?
		$^4[\underline{8}]$	N(1700) Σ(1940) ?
		$^6[\underline{8}]$	N(1670), Λ(1830), Σ(1765)
		$^2[\underline{10}]$	Δ(1650)
		$^4[\underline{10}]$	Δ(1670)
2	$\underline{56}_2^+$	$^4[\underline{8}]$	N(1810), Λ(1860)
		$^6[\underline{8}]$	N(1688), Λ(1815), Σ(1915), Ξ(2030) ?
		$^2[\underline{10}]$	Δ(1910)
		$^4[\underline{10}]$	
		$^6[\underline{10}]$	Δ(1890)
		$^8[\underline{10}]$	Δ(1950), Σ(2030)
	$\underline{56}_0^+$	$^2[\underline{8}]$	N(1470) Σ(1660)
		$^4[\underline{10}]$	Δ(1690)

Table 12

Value of the Color Casimir Operator in Small Representations of SU(3)

Representation	$\langle T^2 \rangle$
[1]	0
[3] or [3*]	4/3
[6] or [6*]	10/3
[8]	3
[10] or [10*]	6
[27]	8

Table 13

"Interaction energies" for few-quark systems

Configuration	$\left\langle \sum_{i<j} \underset{\sim}{T}^{(i)} \cdot \underset{\sim}{T}^{(j)} \right\rangle$
$(q\bar{q})$ $[\underset{\sim}{1}]$	$-4/3$
$(q\bar{q})$ $[\underset{\sim}{8}]$	$+1/6$
(qq) $[\underset{\sim}{3}^{*}]$	$-2/3$
(qq) $[\underset{\sim}{6}]$	$1/3$
(qqq) $[\underset{\sim}{1}]$	-2
(qqq) $[\underset{\sim}{8}]$	$-1/2$
(qqq) $[\underset{\sim}{10}]$	$+1$
$(qqqq)$ $[\underset{\sim}{3}]$	-2

Table 14

Baryon masses including the color hyperfine interaction. For definitions see (5.44) and (5.45).

Baryon	$\Delta E_{HFS} / \delta M_{c.m.}$	N_s	Fitted mass (MeV/c^2)
N(939)	-3	0	939
Λ(1116)	-3	1	1123
Σ(1193)	$1 - 4\, m_u/m_s$	1	1189
Ξ(1318)	$-4\, m_u/m_s + m_u^2/m_s^2$	2	1345
Δ(1232)	$+3$	0	1232
Σ^*(1384)	$1 + 2\, m_u/m_s$	1	1383
Ξ^*(1533)	$2\, m_u/m_s + m_u^2/m_s^2$	2	1539
Ω(1672)	$3\, m_u^2/m_s^2$	3	1701

Table 15

Meson masses including the color hyperfine interaction.
For definitions see (5.52) and (5.53).

Meson	$\Delta E_{HFS} / \delta M_{c.m.}$	N_s	Fitted mass (MeV/c^2)
$\pi(138)$	-3	0	138
$K(496)$	$-3\, m_u/m_s$	1	489
$\eta(549)$	$-\frac{9}{4} - \frac{3}{4}\left(\frac{m_u}{m_s}\right)^2$	1/2	297
$\eta'(958)$	$-\frac{3}{4} - \frac{9}{4}\left(\frac{m_u}{m_s}\right)^2$	3/2	616
$\rho(776)$	1	0	776
$\omega(784)$	1	0	776
$K^*(892)$	m_u/m_s	1	894
$\varphi(1020)$	m_u^2/m_s^2	2	1034

Table 16

Some properties of the Heavy Quarks

Quark	I	Q	Charm	Beauty	Truth
c	0	2/3	1	0	0
b	0	-1/3	0	1	0
t	0	2/3	0	0	1

Table 17

Symmetry properties of flavor and color-spin wavefunctions for three quarks

Flavor SU(3)			Color-spin SU(6)			$G^{(6)^2}$	$SU(3)_{color} \otimes SU(2)_{spin}$
[10]	□□□	S	20	□ □ □	A	21/4	[1], (4) [8], (2)
[8]	□□ □	M	70	□□ □	M	33/4	[1], (2) [8], (4) ⊕ [8], (2) [10], (2)
[1]	□ □ □	A	56	□□□	S	45/4	[8], (2) [10], (4)

Table 18

Candidates for Radially-Excited Pseudoscalars

State	I	Seen In	Remarks
$\pi'(1342)$	1	$\varepsilon\pi$	Bonesini, et al. (1981)
$\eta(1275)$	0	$\eta\pi\pi$	Stanton, et al. (1979)
$\eta'(1400)$	0	$\eta\varepsilon$	Stanton, et al. (1979) via Close (1981)
$\iota(1440)$	0	$\psi\to\gamma+(K\bar{K}\pi)$	Scharre (1981)
K'(1400)	1/2	$K\pi\pi$ $(K\varepsilon)$	Brandenburg, et al. (1976) Aston, et al. (1981)

Table 19

Some properties of the $^3S_1$ ψ states (from Particle Data Group, 1980)

Level	$\Gamma(\psi \to e^+e^-)$, keV	Γ_{tot}, keV
$\psi(3097)$	4.60 ± 0.42	63 ± 9
$\psi'(3685)$	2.05 ± 0.23	215 ± 40
$\psi(4029)$	0.75 ± 0.15	52 ± 10 MeV
$\psi(4159)$	0.77 ± 0.23	78 ± 20 MeV
$\psi(4415)$	0.49 ± 0.13	42 ± 10 MeV

Table 20

Some properties of the $^3S_1$ Υ states (from the review by Schamberger, 1981)

Level	$\Gamma(\Upsilon \to e^+e^-)$, keV	Γ_{tot} , keV
Υ (9433)	1.17 $\pm$ 0.05	35.5^{+8}_{-6}
Υ' (9993)	0.54 $\pm$ 0.03	$\approx \Gamma(\Upsilon)$
Υ'' (10323)	0.37 $\pm$ 0.03	
Υ''' (10546)	0.27 $\pm$ 0.02	$\sim$15 MeV

Table 21

Decay modes of the $(s\bar{s})$ state, $\varphi(1020)$ (Particle Data Group, 1980)

Channel	Branching Fraction (%)	q_{max} (MeV/c)
K^+K^-	48.6 ± 1.2	127
$K_L K_S$	35.2 ± 1.2	111
$\pi^+\pi^-\pi^0$	14.7 ± 0.7	462
$\eta\gamma$	1.5 ± 0.2	362
$\pi^0\gamma$	0.14 ± 0.05	501
e^+e^-	0.031 ± 0.001	510
$\mu^+\mu^-$	0.025 ± 0.003	499

CAPTIONS

Fig. 1 : Energy levels for A = 7 nuclei (from Ajzenberg-Selove, 1979). The diagrams for individual isobars have been shifted vertically to eliminate the neutron-proton mass difference and the Coulomb energy, taken as $E_C = (0.6\ \mathrm{MeV})\, Z(Z-1)/A^{1/3}$. Energies in square brackets represent the approximate nuclear binding energy $E_N = M(Z,A) - ZM_p - (A-Z)M_n - E_C$, minus the corresponding quantity for ^7Li. Note the one-to-one correspondence between levels of the mirror nuclei ^7Li and ^7Be.

Fig. 2 : Energy levels for A = 11 nuclei (from Ajzenberg-Selove, 1975). Notation as in Fig. 1, with binding energies referred to ^{11}B.

Fig. 3 : Energy levels for A = 14 nuclei (from Ajzenberg-Selove, 1976). Notation as in Fig. 1, with binding energies referred to ^{14}N.

Fig. 4 : The isospin quarks.

Fig. 5 : Isospin assignments of the nucleons and nucleon resonances.

Fig. 6 : The weight diagram for the fundamental $[\underset{\sim}{3}]$ representation of SU(3).

Fig. 7 : Weight diagram for the vector meson nonet = $[\underset{\sim}{1}] \oplus [\underset{\sim}{8}]$.

Fig. 8 : Decompositions of the fundamental quark triplets with respect to the SU(2) subgroups U-spin and V-spin.

Fig. 9 : The $J^P = 3/2^+$ baryon decimet.

Fig. 10 : The $J^P = 1/2^+$ baryon octet.

Fig. 11 : Action of the I-spin, U-spin, and V-spin raising and lowering operators on the fundamental triplets of quarks $[\underset{\sim}{3}]$ and antiquarks $[\underset{\sim}{3}^*]$.

Fig. 12 : Properties of the lowest mode of a fermion confined within a rigid sphere. (a) Fermion momentum as a function of its mass m and the sphere radius R. (b) Ratio of the fermion mass to the energy of its lowest confined mode.

Fig. 13 : Lowest-mode energy of a massless fermion confined to a rigid, static sphere of radius R (see eqn. (3.145)).

Fig. 14 : Lowest-mode energy of a fermion of mass m confined within a rigid, static sphere of radius 4/3 fm.

Fig. 15 : Magnetic moment of the confined fermion in units of the Dirac moment for a free fermion with mass equal to the energy ω of the confined fermion.

Fig. 16 : Axial charge of a nucleon composed of equal-mass quarks confined within a rigid spherical cavity, as a function of the dimensionless parameter mR.

Fig. 17 : Overlap factor ϑ defined in eqn. (3.155) as measured in various M1 decays of mesons. The dashed entry for $\omega \to \pi^0 \gamma$ is for the reanalysis by Ohshima (1980).

Fig. 18 : Method of pole extrapolation for studying the interactions of unstable target particles.

(a) Pion scattering from a virtual pion.

(b) Measurement of the electromagnetic form factor of a virtual pion. Additional diagrams are required to contribute by gauge invariance.

Fig. 19 : Regge trajectories of the natural-parity mesons. Uncertain states are indicated by open circles.

Fig. 20 : Expected SU(6) multiplets of baryons.

Fig. 21 : Regge trajectories of the nucleon, Δ , and Λ resonances.

Fig. 22 : The ratio R = $\sigma(e^+e^- \to \text{hadrons})/\sigma(e^+e^- \to \mu^+\mu^-)$ compared with the predictions of the quark-parton model.

(a) W $<$ 4 GeV (after Spinetti, 1979) ; (b) 8 GeV $<$ W $<$ 40 GeV (after Schamberger, 1981).

Fig. 23 : The quark-quark-gluon interaction in QCD.

Fig. 24 : A baryon configuration which is not considered in the sum over two-body forces.

Fig. 25 : Attempting to separate a quark and antiquark results in the creation of a quark-antiquark pair from the vacuum, so that color is always neutralized locally.

Fig. 26 : A massless quark and antiquark connected by a linear string.

Fig. 27 : Meson states in flavor SU(6), decomposed into $SU(4)_{udsc} \otimes U(1)_b \otimes U(1)_t$. The additive quantum numbers are denoted by B(beauty) and T(truth).

Fig. 28 : $J^P = 1/2^+$ baryon states in flavor SU(6). The circled states occur twice, as do those that lie in both $[\underset{\sim}{6}]$ and $[\underset{\sim}{3}^*]$ of $SU(3)_{uds}$. There are 70 states in all.

Fig. 29 : $J^P = 3/2^+$ baryon states in flavor SU(6). There are 56 states in all.

Fig. 30 : The spectrum of charmonium ($c\bar{c}$). Branching fractions (in percent) are shown for the important classes of decays (Particle Data Group, 1980 ; Himel, et al., 1980a ; Oreglia, et al., 1980 ; Schamberger, 1981 ; Scharre, 1981). Charm threshold is indicated at twice the D meson mass.

Fig. 31 : The spectrum of upsilon ($b\bar{b}$) states. Branching fractions (in percent) are shown for the important classes of identified decays (Schamberger, 1981). Beauty threshold is indicated schematically.

Fig. 32 : Lower bounds for leptonic decays of Υ and Υ' (after Rosner, et al., 1978) together with the data cited in Table 20. The bounds are computed from eqn. (6.45) using ψ leptonic widths 1σ below the central values and assuming $m_b / m_c \gtrsim 3$.

Fig. 33 : A possible spectrum of strangeonium ($s\bar{s}$) levels. Identification of E(1418) and $\varphi(1634)$ as pure $s\bar{s}$ states may be disputed. The dotted 0^{-+} entry is impressionistic, having been invented from the φ mass and the $\pi - \rho$ splitting, appropriately rescaled.

Fig. 34 : Charge induced by a positive test charge placed at the center of a hole in a dielectric medium. (a) Dia-electric case $\varepsilon_{medium} < 1$ hoped to resemble QCD ; (b) Dielectric case $\varepsilon_{medium} > 1$ of normal electrodynamics.

Fig. 35 : (a) A single link between two quarks in lattice gauge theory. (b) The smallest closed loop, corresponding to a quarkless excitation.

Fig. 36 : The Okubo-Zweig-Iizuka rule applied to φ -decay. The connected diagram (a) is allowed ; the disconnected diagram (b) is forbidden.

Fig. 37 : Cross sections for the two-photon reactions $e^+e^- \longrightarrow e^+e^- +$ hadrons. The cross section for excitation of $\iota(1440)$ is computed under the assumption that $\Gamma(\iota \to \gamma\gamma) = 1$ keV, and so should be multiplied by $\Gamma(\iota \to \gamma\gamma)$ /(1 keV). The cross section for the reaction $e^+e^- \to \mu^+\mu^-$ ("one unit of R") is shown for reference.

Fig. 38 : Double-line notation for quarks, gluons, and their interactions useful for $1/N_c$ analyses.

Fig. 39 : Lowest order vacuum polarization contributions to the gluon propagator. (a) quark loop ; (b) gluon loop ; (c) quark loop in the double-line notation ; (d) gluon loop in the double-line notation.

Fig. 40 : A two-loop diagram in (a) conventional and (b) double-line notation.

Fig. 41 : A three-loop diagram in (a) conventional and (b) double-line notation.

Fig. 42 : A nonplanar graph in (a) conventional and (b) double-line notation.

Fig. 43 : A mechanism for OZI-forbidden decay, at order g^4, in (a) conventional and (b) double-line notation.

Fig. 44 : OZI-allowed decay of a meson, at order g^4, in (a) conventional and (b) double-line notation.

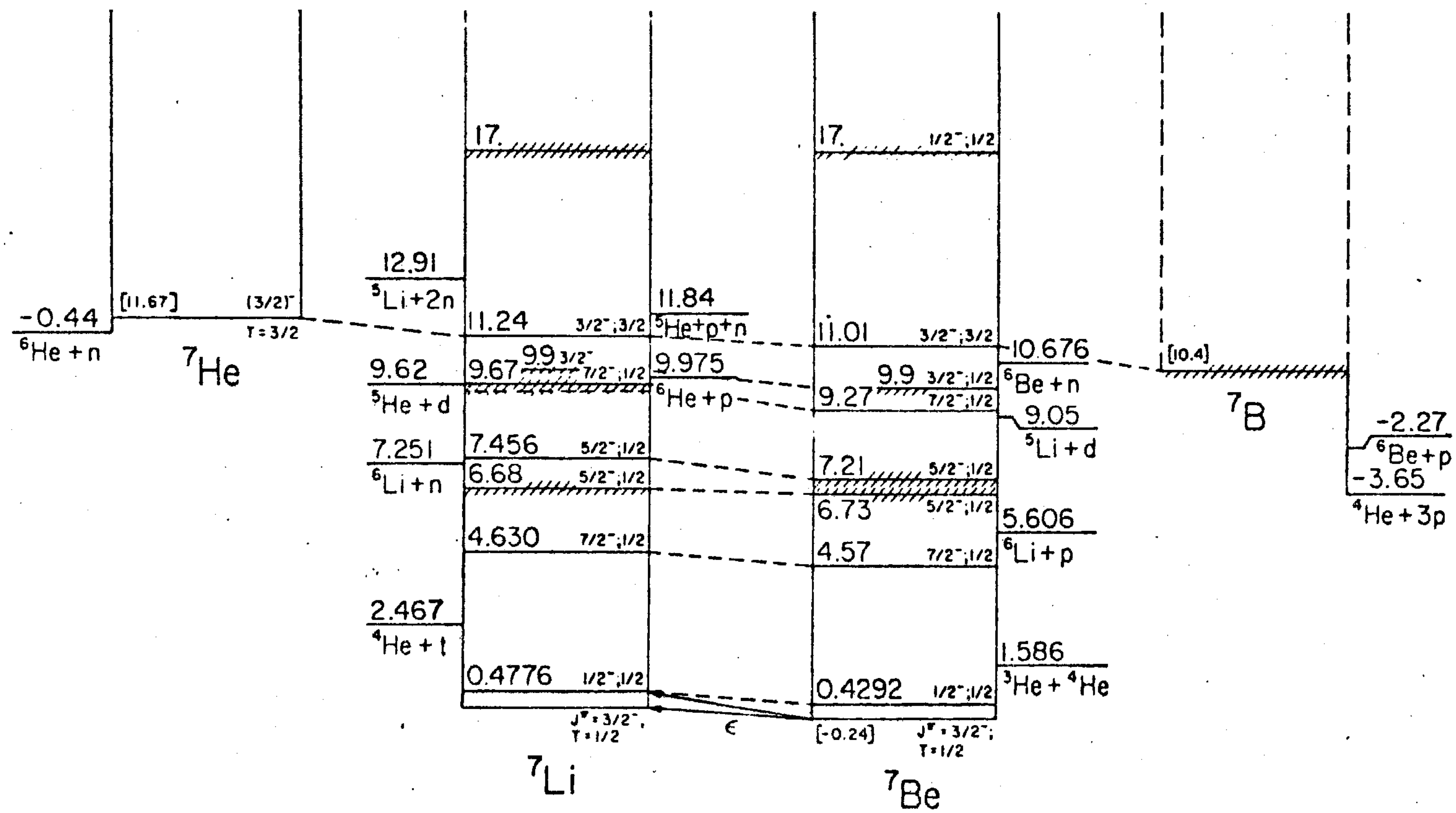
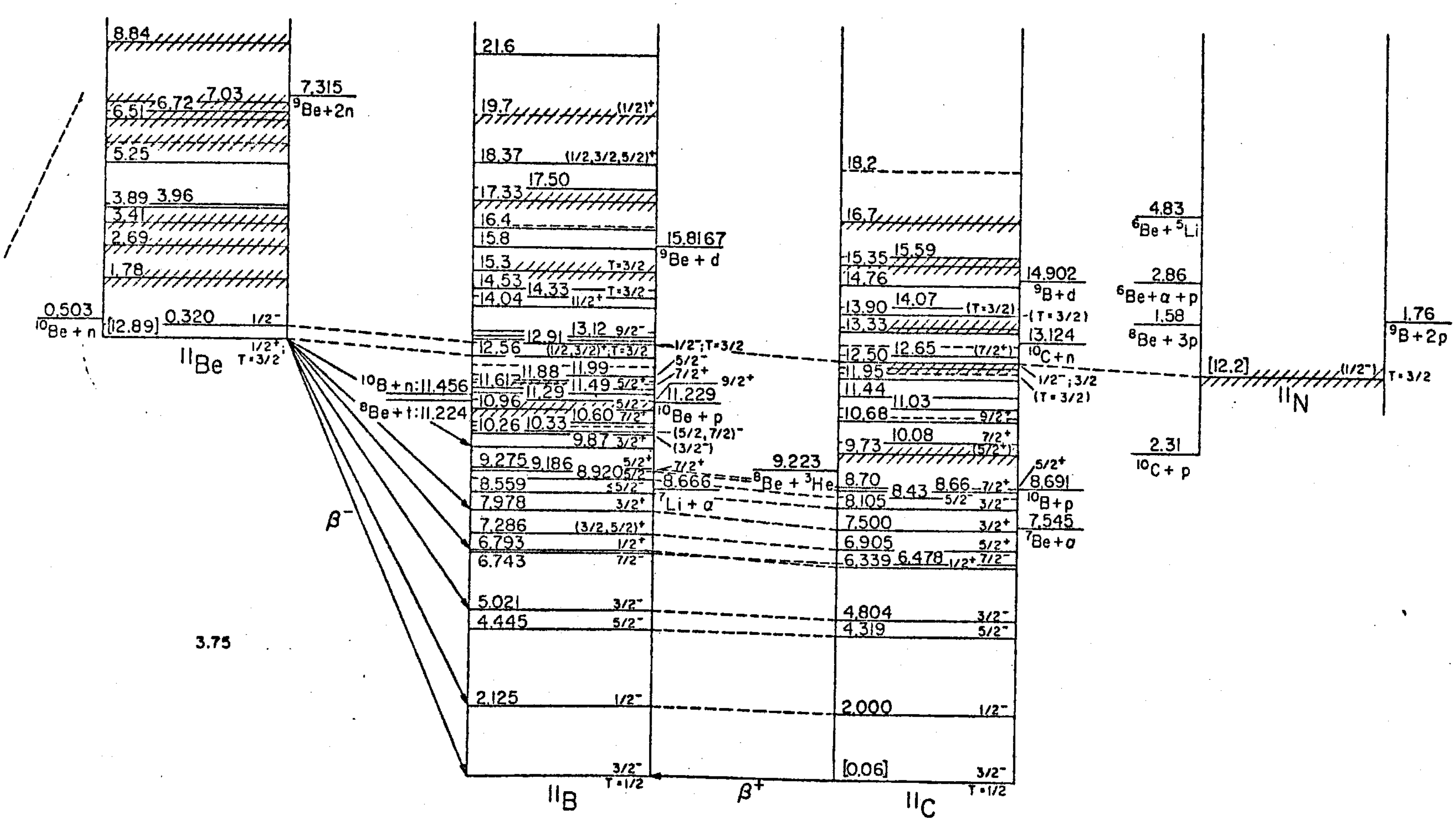
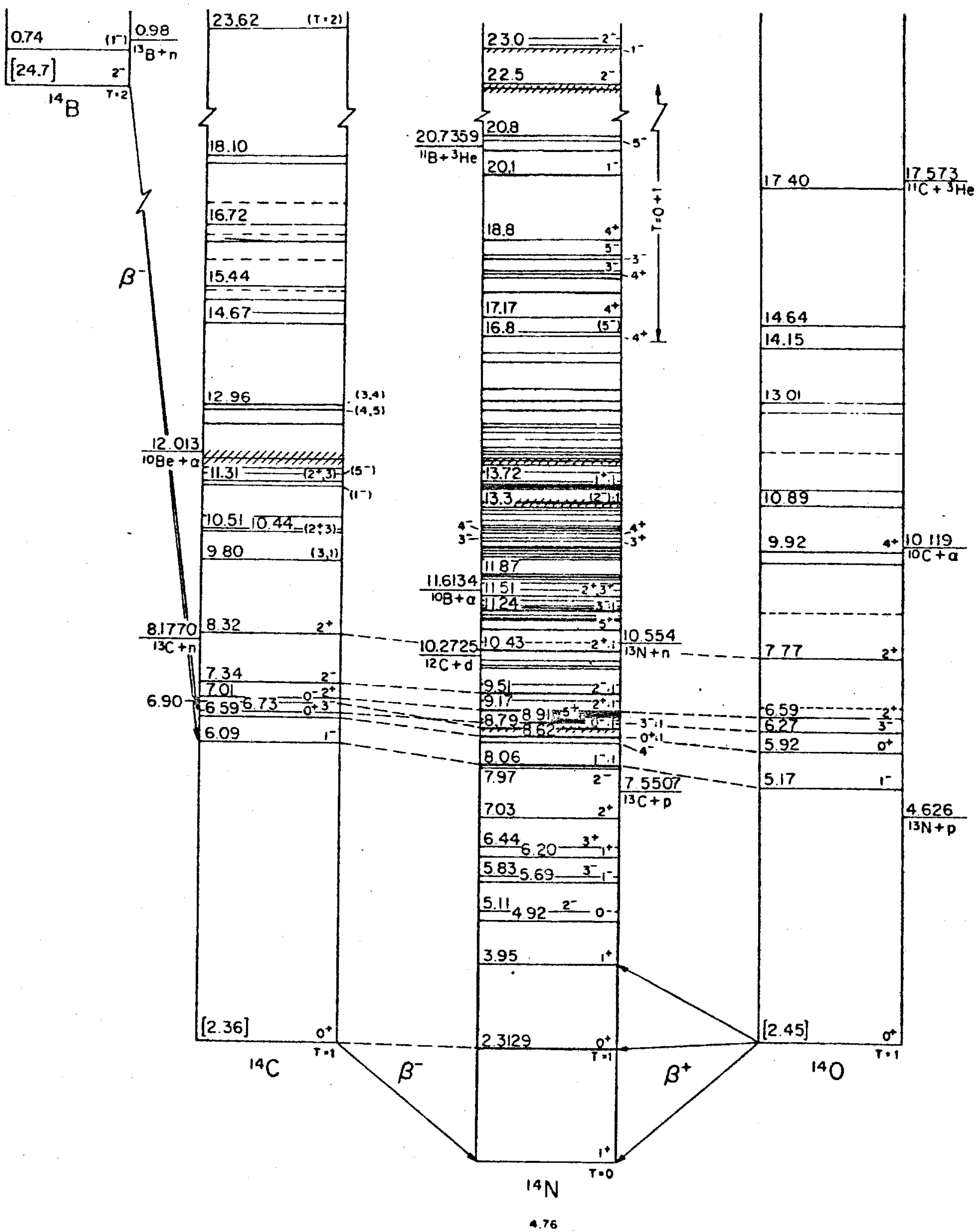

-0.44
^{6}He + n
[11.67]
(3/2)$^-$
T = 3/2
^{7}He
12.91
^{5}Li+2n
9.62
^{5}He + d
7.251
^{6}Li + n
2.467
^{4}He + t
17.
11.24
3/2$^-$;3/2
9.67
9.9
3/2$^-$
7/2$^-$;1/2
7.456
5/2$^-$;1/2
6.68
5/2$^-$;1/2
4.630
7/2$^-$;1/2
0.4776
1/2$^-$;1/2
J$^\pi$ = 3/2$^-$;
T = 1/2
^{7}Li
11.84
^{5}He+p+n
9.975
^{6}He + p
ε
17.
1/2$^-$;1/2
11.01
3/2$^-$;3/2
9.9
3/2$^-$;1/2
9.27
7/2$^-$;1/2
7.21
5/2$^-$;1/2
6.73
5/2$^-$;1/2
4.57
7/2$^-$;1/2
0.4292
1/2$^-$;1/2
[-0.24]
J$^\pi$ = 3/2$^-$;
T = 1/2
^{7}Be
10.676
^{6}Be + n
9.05
^{5}Li + d
5.606
^{6}Li + p
1.586
^{3}He + ^{4}He
[10.4]
^{7}B
-2.27
^{6}Be + p
-3.65
^{4}He + 3p

Fig. 1



11Be
8.84
7.03
6.72
6.51
7.315
9Be+2n
5.25
3.89 3.96
3.41
2.69
1.78
0.503
10Be + n
[12.89]
0.320
1/2-
1/2+; T=3/2
10B + n:11.456
8Be+t:11.224
β-
11B
21.6
19.7
(1/2)+
18.37
(1/2,3/2,5/2)+
17.33
17.50
16.4
15.8
15.8167
9Be + d
15.3
T=3/2
14.53
14.33
14.04
11/2+
12.91
13.12
9/2-
12.56
11.612
11.88
11.99
11.29
11.49
11.229
10.96
10.26
10.33
10.60
9.87
10Be + p
(5/2, 7/2)-
(3/2-)
9.275
9.186
8.920
8.666
8.559
7.978
7Li + α
7.286
(3/2, 5/2)+
6.793
6.743
5.021
4.445
2.125
3/2-
T=1/2
9.223
8Be + 3He
β+
11C
18.2
16.7
15.35
15.59
14.76
13.90
14.07
(T=3/2)
13.33
12.50
12.65
11.95
11.44
11.03
10.68
9.73
10.08
8.70
8.43
8.66
8.105
7.500
6.905
6.339
6.478
4.804
4.319
2.000
[0.06]
14.902
9B+d
13.124
10C+n
1/2-; 3/2
(T=3/2)
8.691
10B+p
7.545
7Be+α
4.83
6Be + 5Li
2.86
6Be+α+p
1.58
8Be + 3p
2.31
10C + p
[12.2]
11N
(1/2-)
T=3/2
1.76
9B + 2p
3.75

Fig. 2



0.74 (1⁻)
[24.7] 2⁻
T=2
14B
0.98
13B+n
β⁻
23.62 (T=2)
18.10
16.72
15.44
14.67
12.96 (3,4) (4,5)
12.013
10Be+α
11.31 (2⁺,3) (5⁻) (1⁻)
10.51 10.44 (2⁺,3)
9.80 (3,1)
8.1770
13C+n
8.32 2⁺
7.34 2⁻
7.01 0⁻ 2⁺
6.90
6.59 6.73 0⁺ 3⁻
6.09 1⁻
[2.36] 0⁺
T=1
14C
β⁻
23.0 2⁻ 1⁻
22.5 2⁻
20.7359
11B+3He
20.8 5⁻
20.1 1⁻
18.8 4⁺
5⁻ 3⁻
3⁻ 4⁺
17.17 4⁺
16.8 (5⁻) 4⁺
T=0+1
13.72 1⁺;1
13.3 (2⁻);1
4⁻ 3⁻ 4⁺ 3⁺
11.87
11.6134
10B+α
11.51 2⁺,3⁺
11.24 3⁻;1
5⁺
10.2725
12C+d
10.43 2⁺;1
10.554
13N+n
9.51 2⁻;1
9.17 2⁺;1
8.79 8.91 5⁺
8.62 0⁻ 3⁻;1 0⁺;1 4⁻
8.06 1⁻;1
7.97 2⁻
7.5507
13C+p
7.03 2⁺
6.44 6.20 3⁺ 1⁺
5.83 5.69 3⁻ 1⁻
5.11 4.92 2⁻ 0⁻
3.95 1⁺
2.3129 0⁺
T=1
1⁺
T=0
14N
4.76
β⁺
17.40
17.573
11C+3He
14.64
14.15
13.01
10.89
9.92 4⁺
10.119
10C+α
7.77 2⁺
6.59 2⁺
6.27 3⁻
5.92 0⁺
5.17 1⁻
4.626
13N+p
[2.45] 0⁺
T=1
14O

Fig. 3

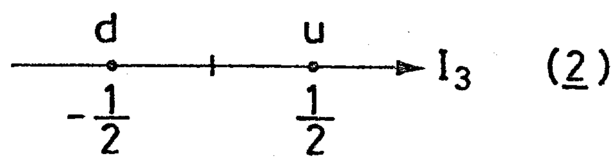

d
u
I_3
(2)
$-\frac{1}{2}$
$\frac{1}{2}$

Fig. 4

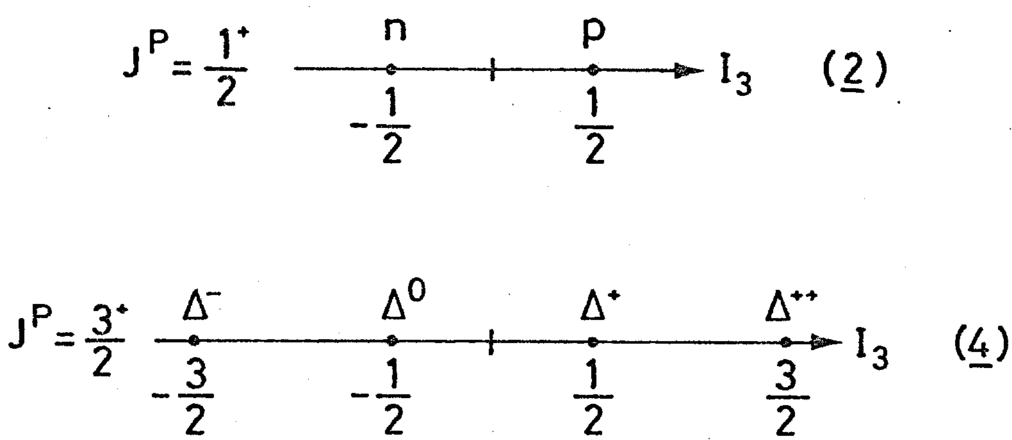

$J^P = \frac{1^+}{2}$
n
p
I_3
(2)
$-\frac{1}{2}$
$\frac{1}{2}$
$J^P = \frac{3^+}{2}$
Δ^-
Δ^0
Δ^+
Δ^{++}
I_3
(4)
$-\frac{3}{2}$
$-\frac{1}{2}$
$\frac{1}{2}$
$\frac{3}{2}$

Fig. 5

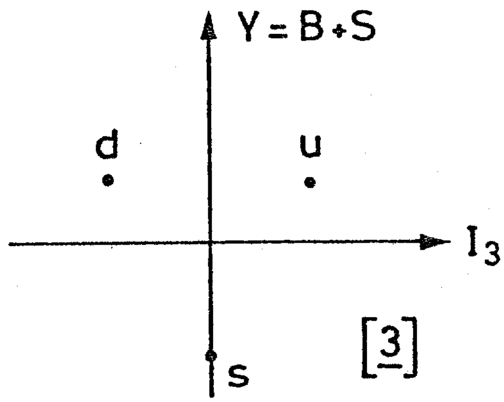


Fig. 6

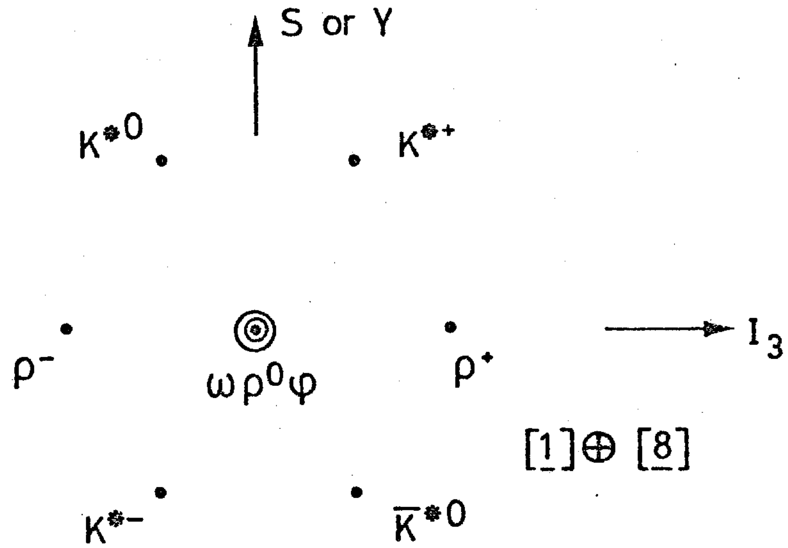

S or Y
K^{*0}
K^{*+}
ρ^-
$\omega\rho^0\varphi$
ρ^+
I_3
[1]⊕ [8]
K^{*-}
$\overline{K}^{*0}$

Fig. 7

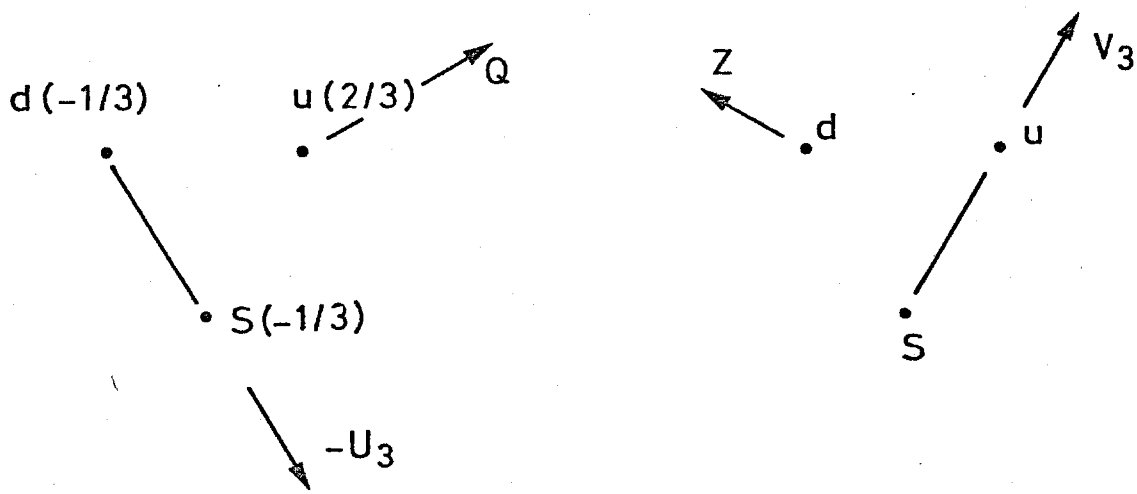

d (−1/3)
u (2/3)
Q
Z
d
V_3
u
S(−1/3)
S
$-U_3$

Fig. 8

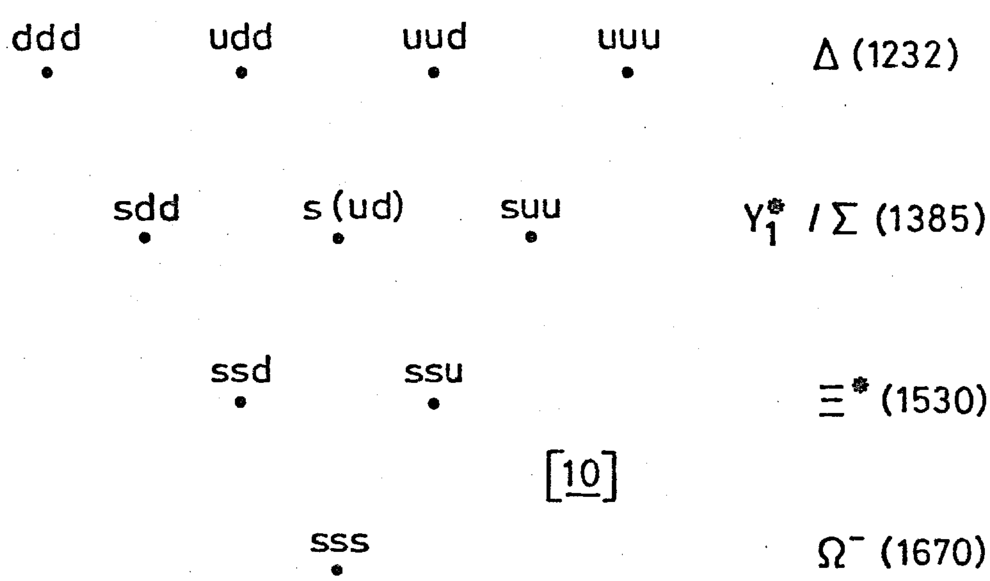

ddd
udd
uud
uuu
Δ (1232)
sdd
s (ud)
suu
Y_1^* / Σ (1385)
ssd
ssu
Ξ* (1530)
[10]
sss
Ω⁻ (1670)

Fig. 9

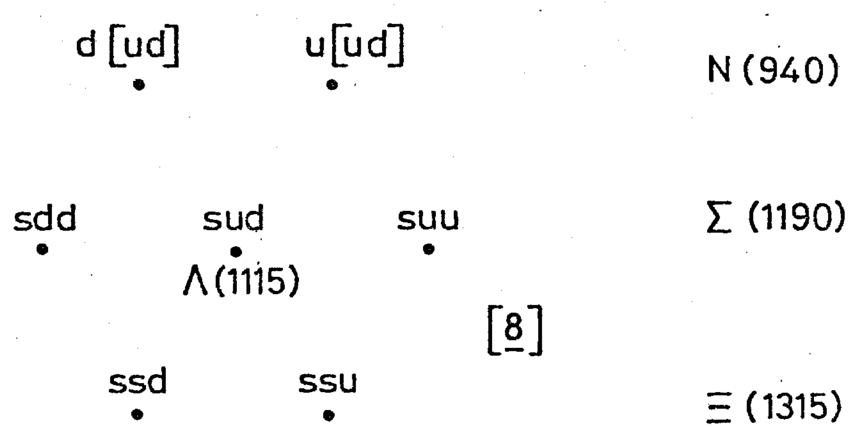

d [ud]
u[ud]
N (940)
sdd
sud
suu
Σ (1190)
Λ(1115)
[8]
ssd
ssu
Ξ (1315)

Fig. 10

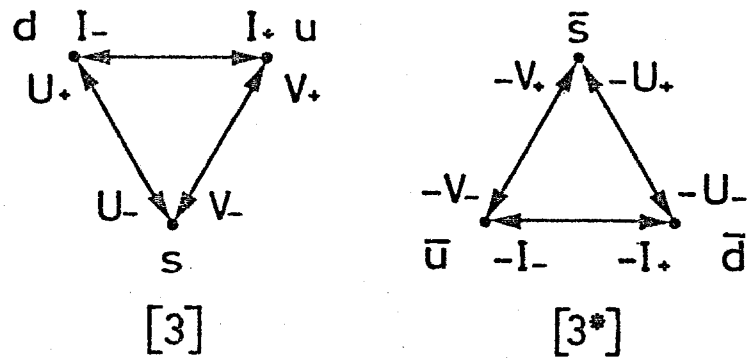


Fig. 11

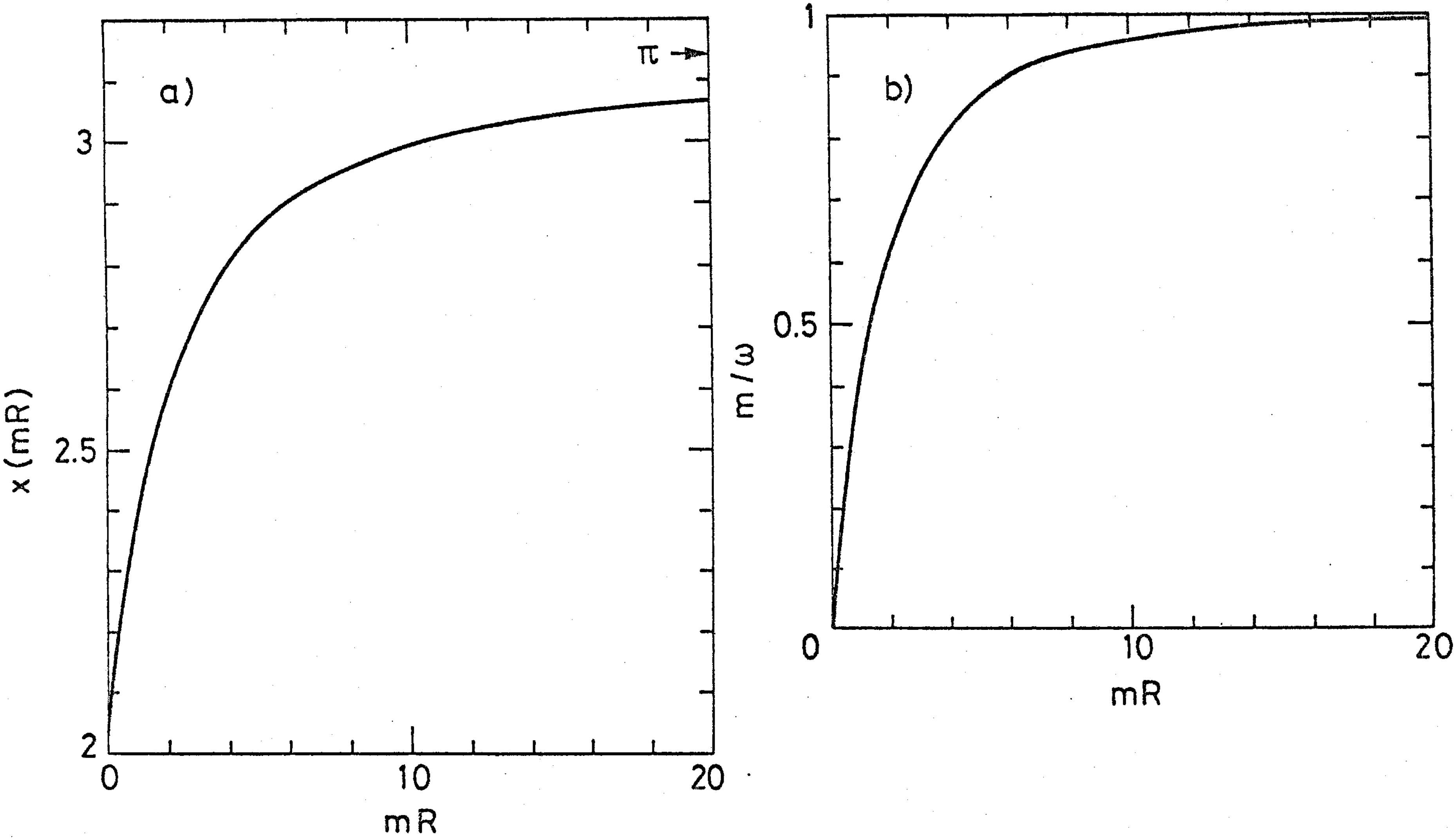


Fig. 12

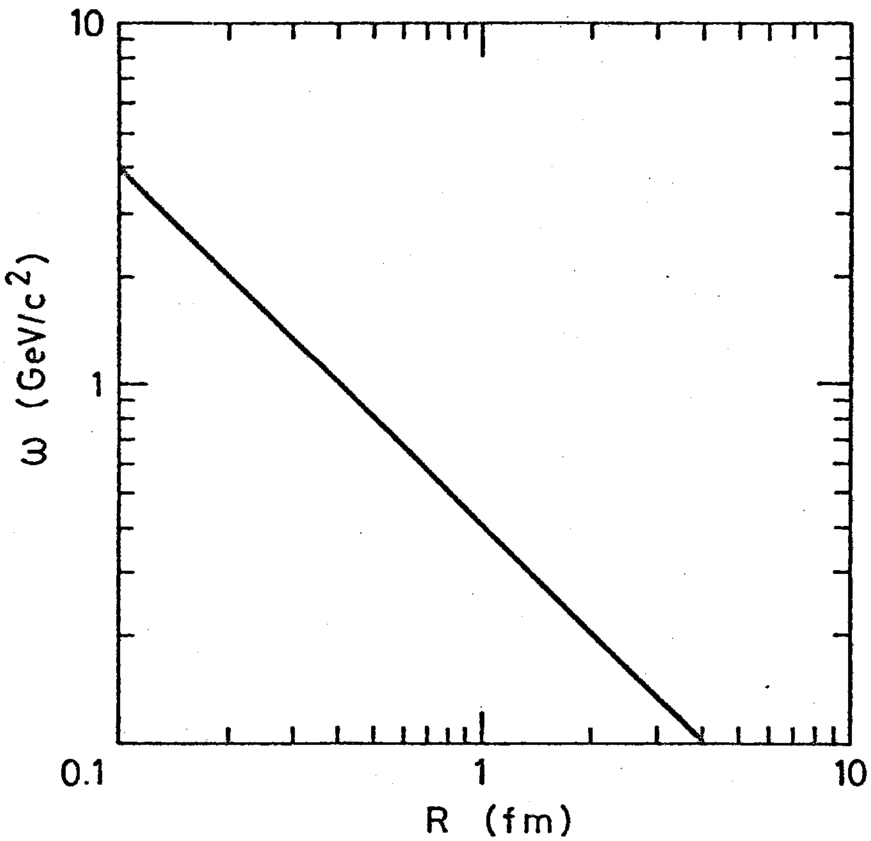


Fig. 13

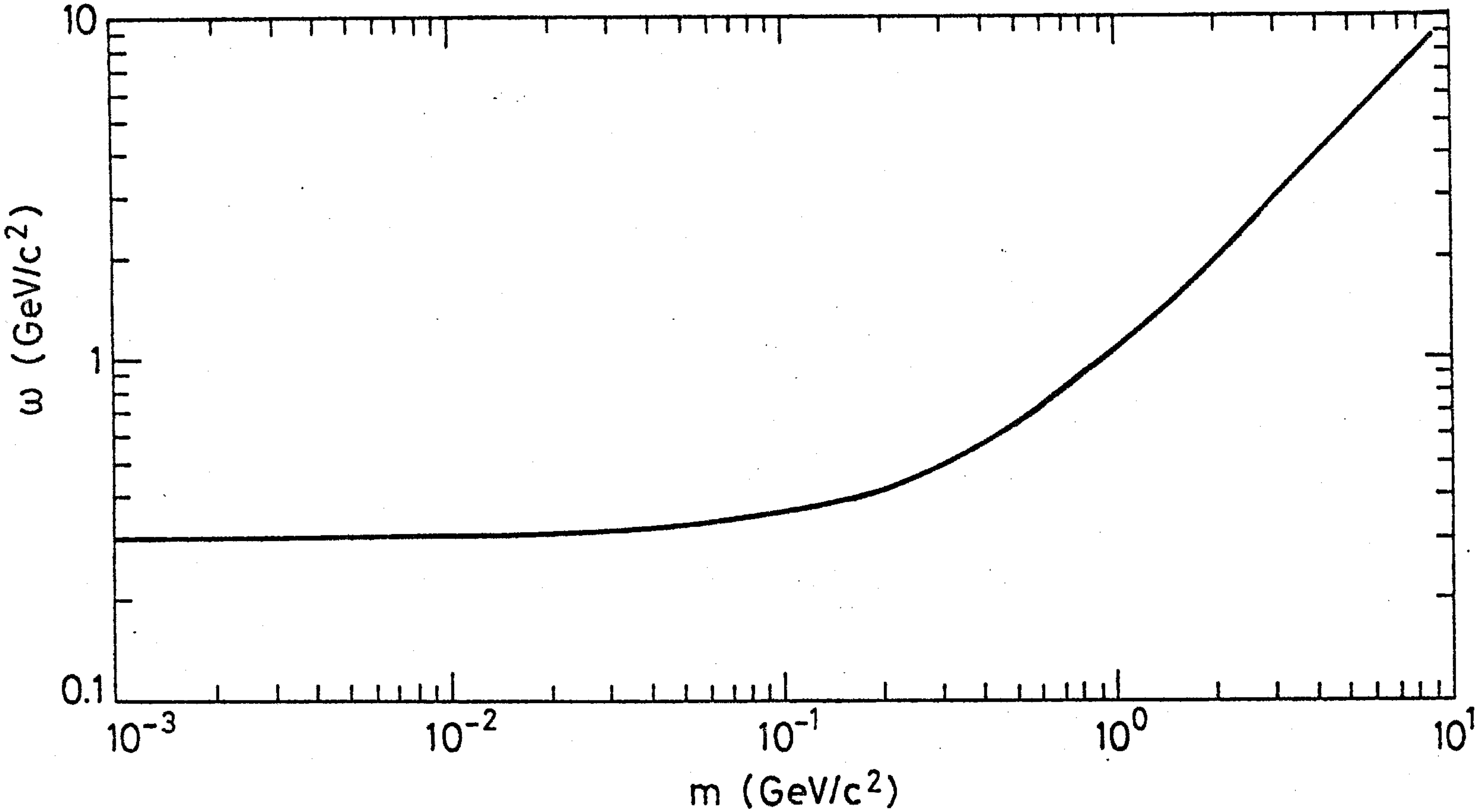


Fig. 14

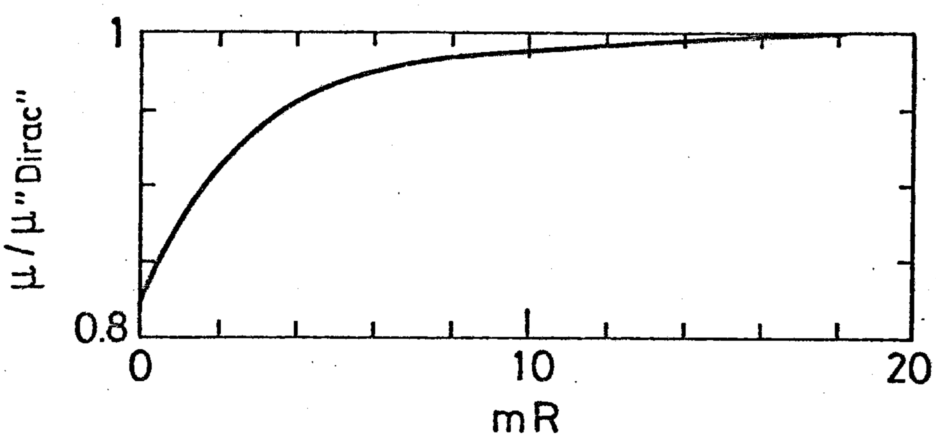


Fig. 15

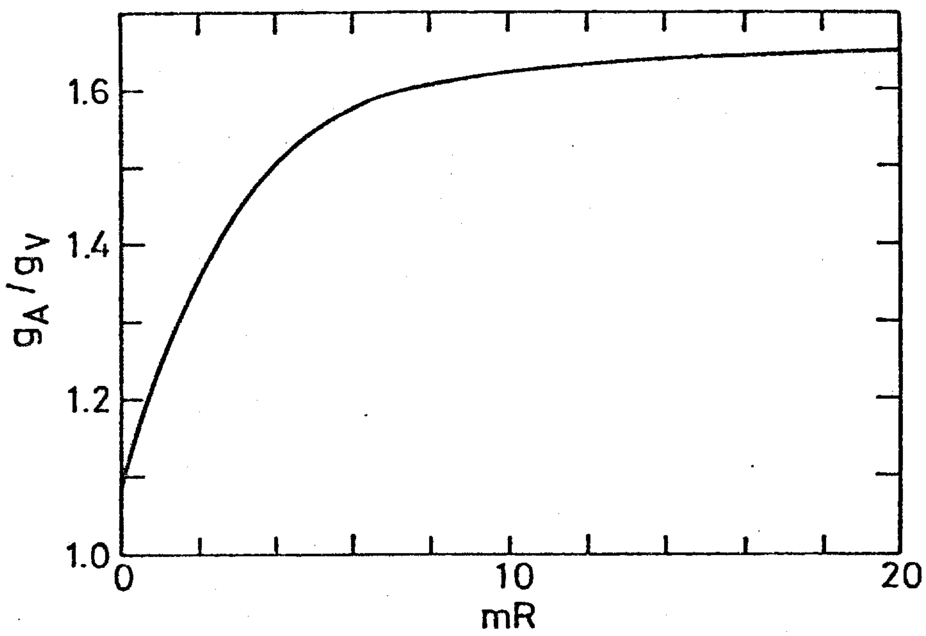


Fig. 16

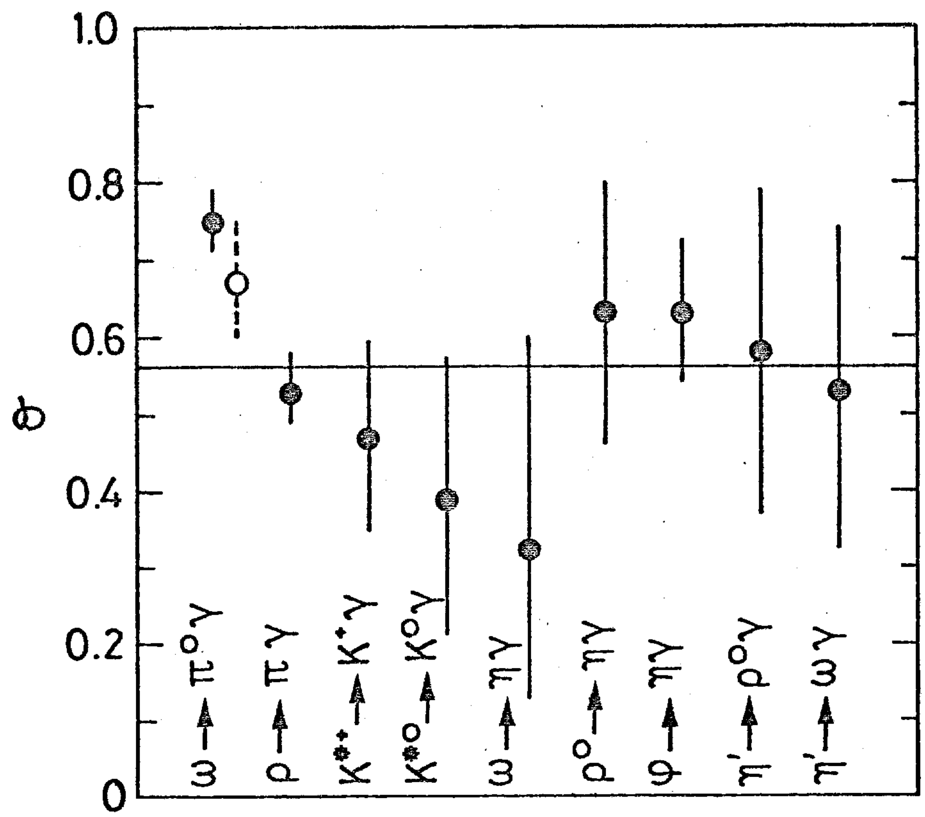


Fig. 17

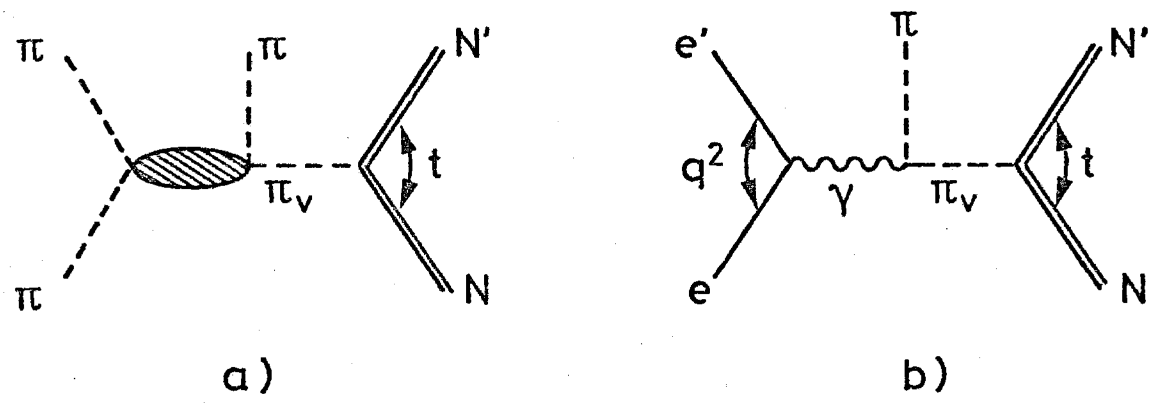


Fig. 18

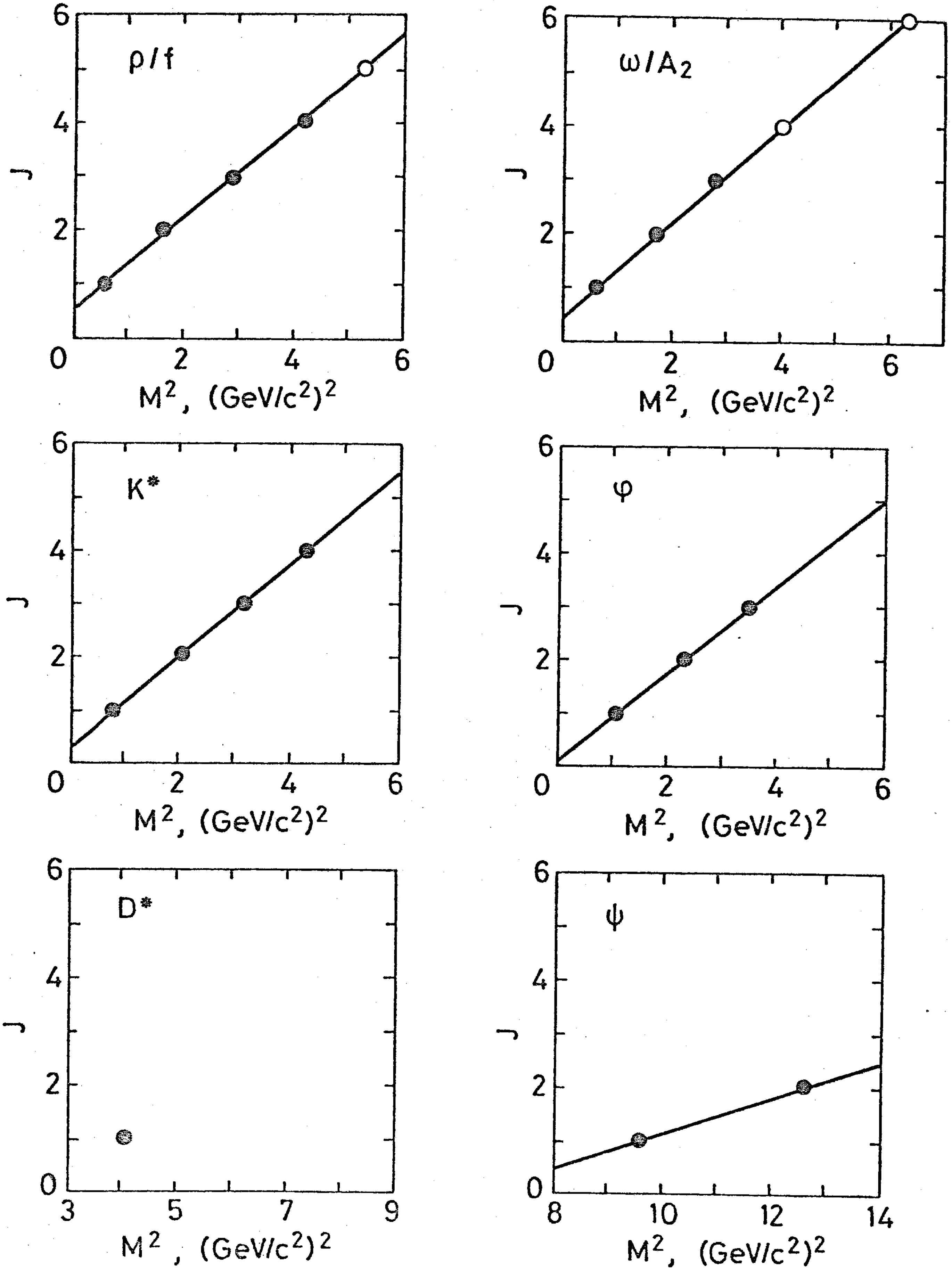
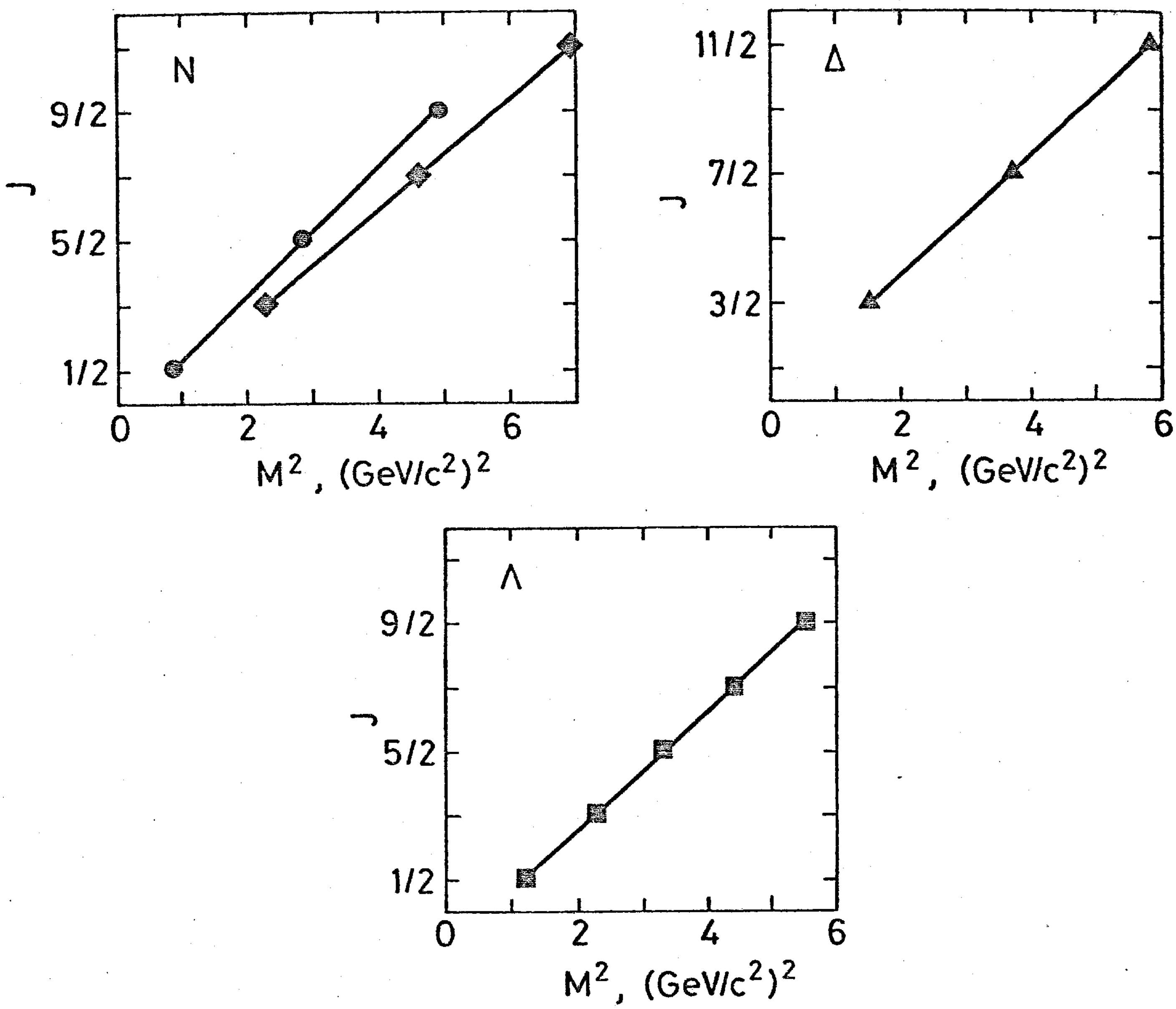


Fig. 19



N
J
9/2
5/2
1/2
0
2
4
6
M², (GeV/c²)²
Δ
J
11/2
7/2
3/2
0
2
4
6
M², (GeV/c²)²
Λ
J
9/2
5/2
1/2
0
2
4
6
M², (GeV/c²)²

Fig. 21

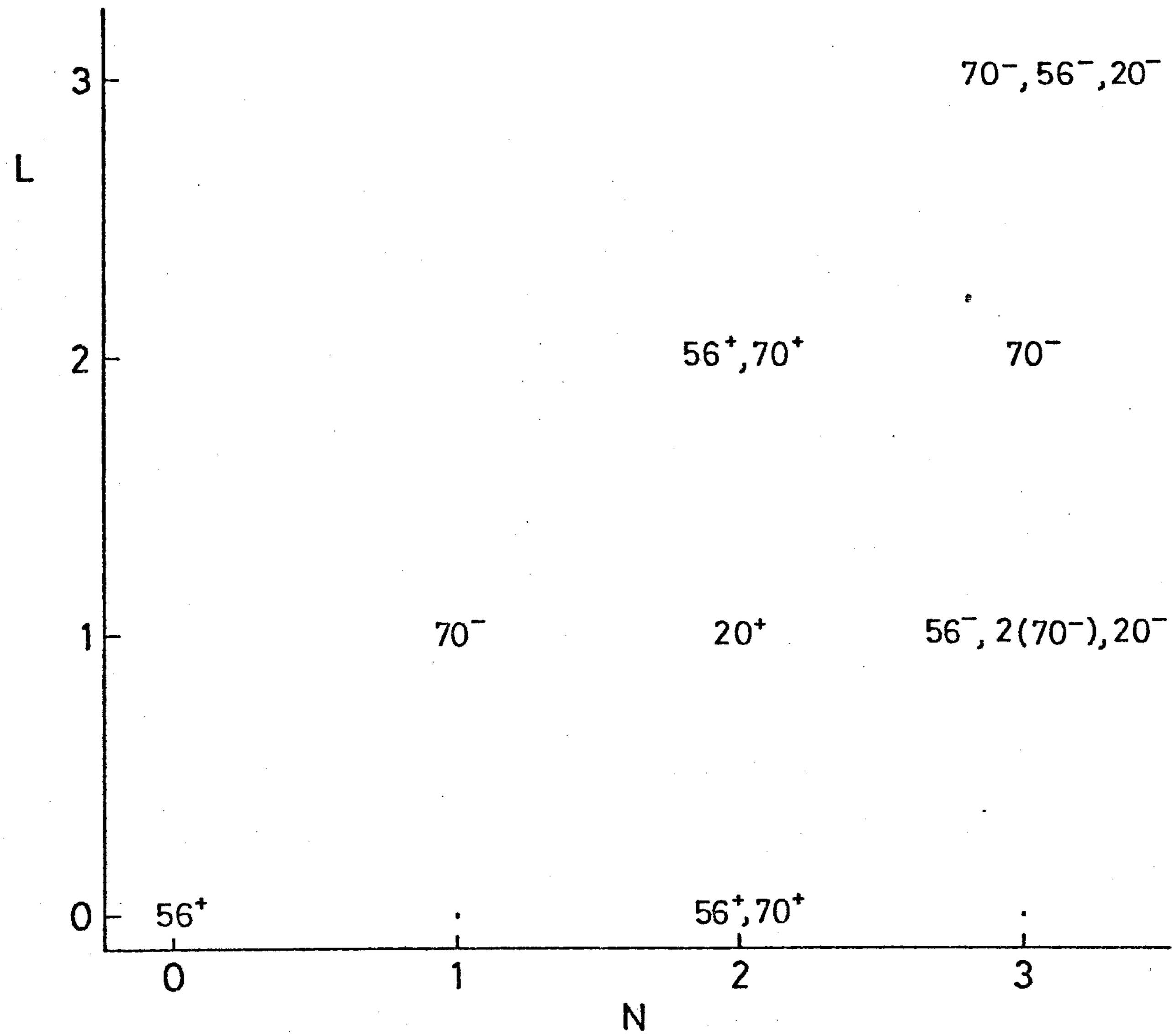
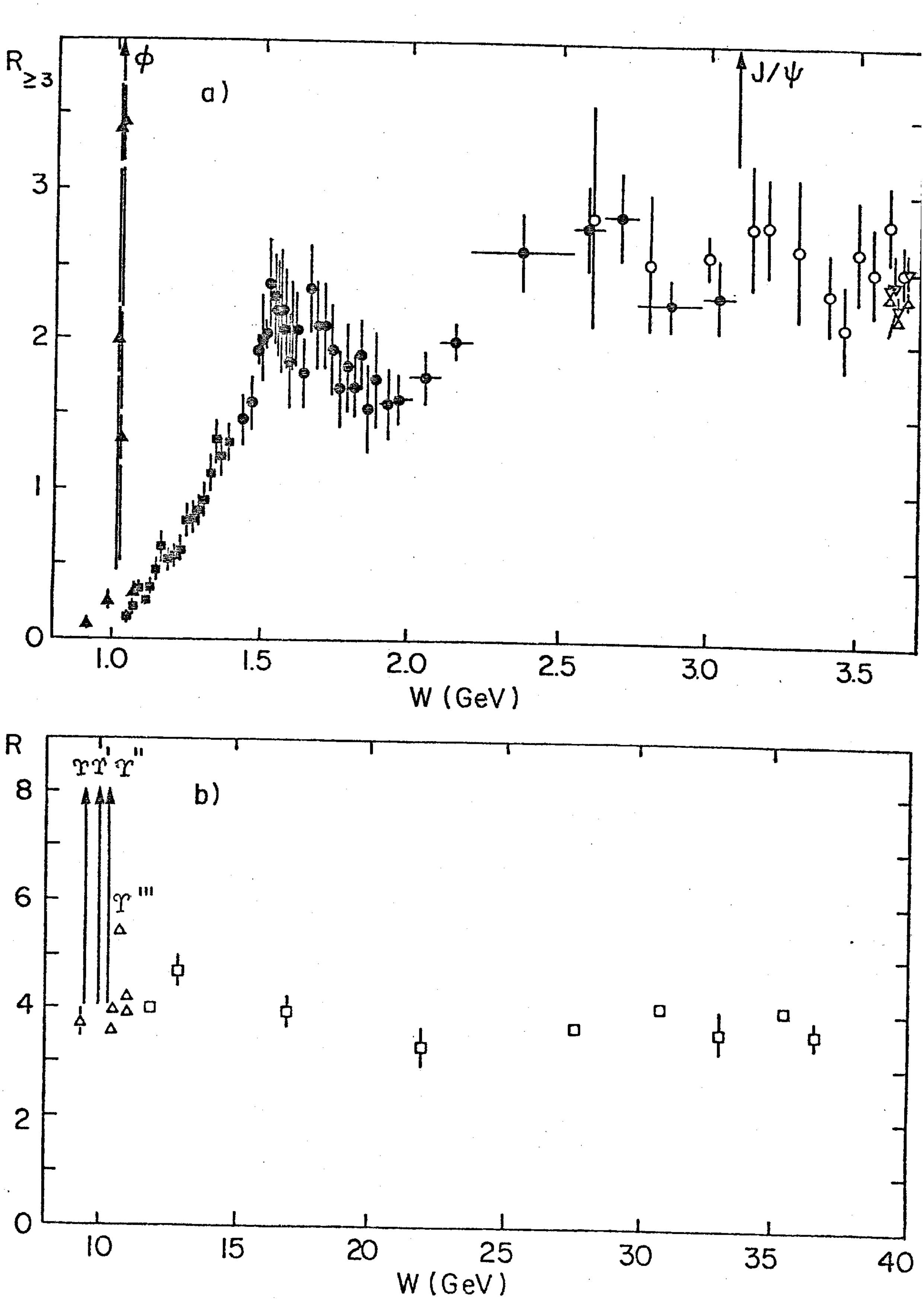


Fig. 20



$R_{\geq 3}$
ϕ
a)
J/ψ
0
1
2
3
1.0
1.5
2.0
2.5
3.0
3.5
W (GeV)
R
$\Upsilon \Upsilon' \Upsilon''$
b)
Υ'''
0
2
4
6
8
10
15
20
25
30
35
40
W (GeV)

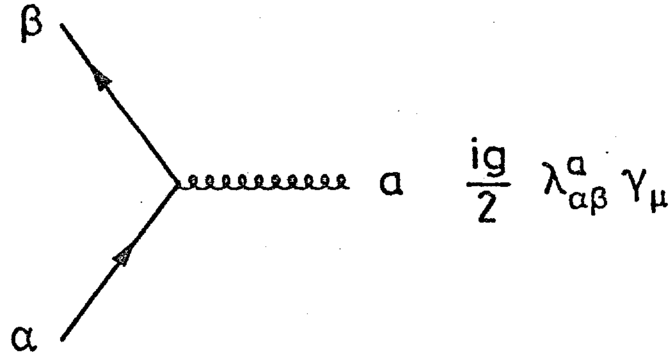


Fig. 23

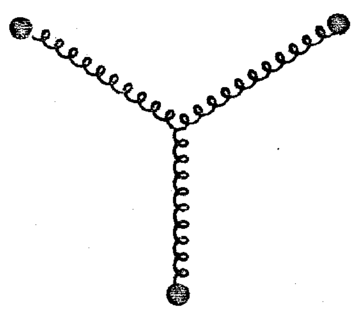


Fig. 24

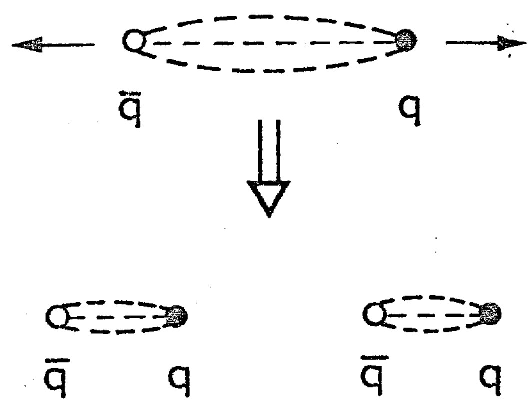


Fig. 25

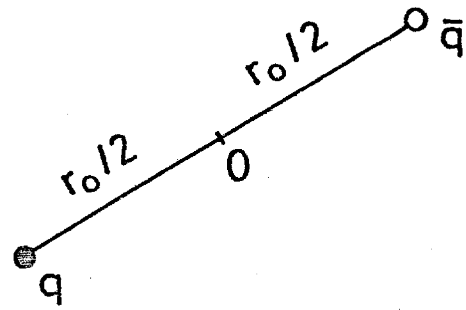
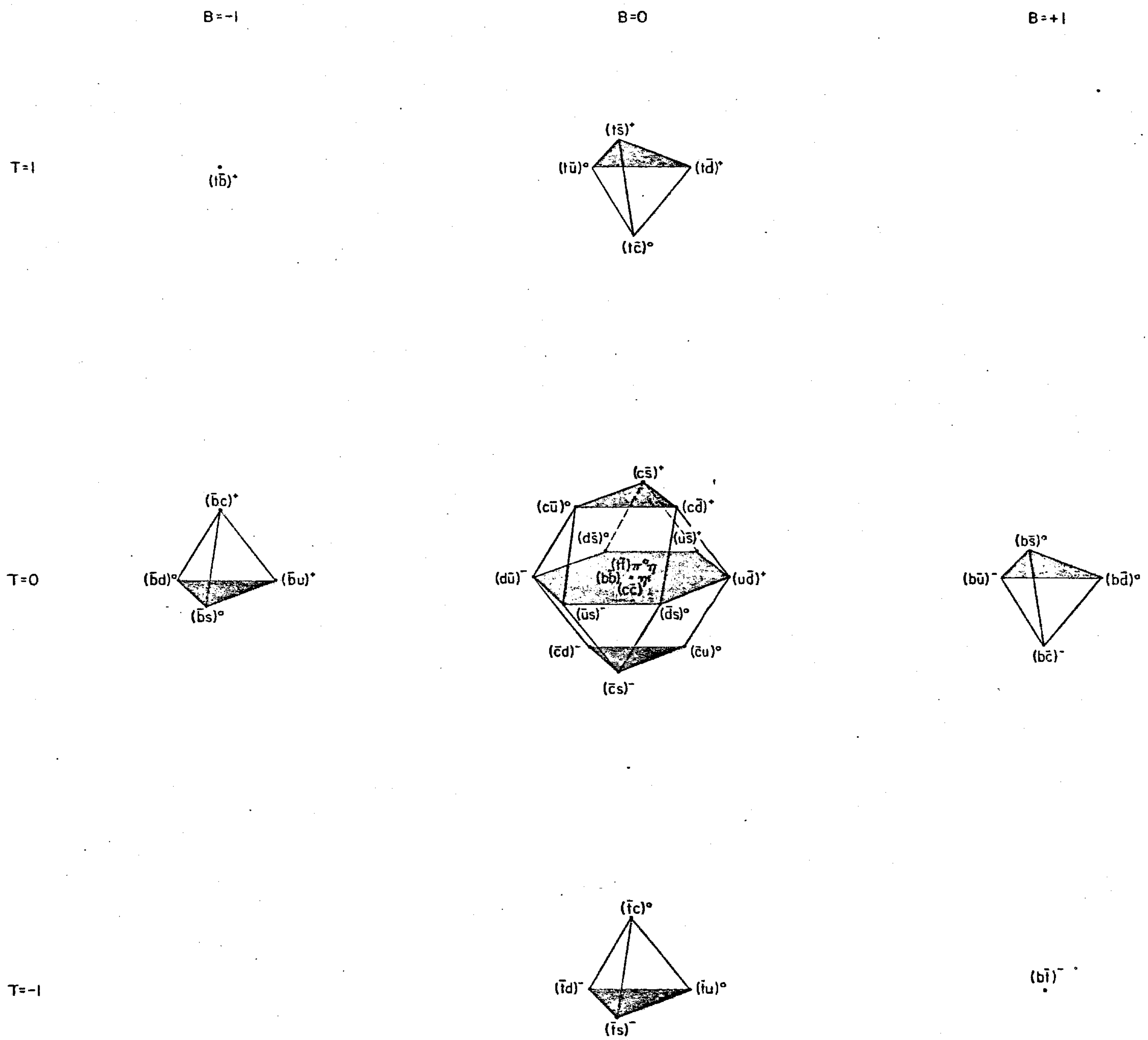


Fig. 26



B=-1
B=0
B=+1
T=1
T=0
T=-1
$(t\bar{b})^+$
$(t\bar{s})^+$
$(t\bar{u})^0$
$(t\bar{d})^+$
$(t\bar{c})^0$
$(\bar{b}c)^+$
$(\bar{b}d)^0$
$(\bar{b}u)^+$
$(\bar{b}s)^0$
$(c\bar{s})^+$
$(c\bar{u})^0$
$(c\bar{d})^+$
$(d\bar{s})^0$
$(\bar{u}s)^+$
$(d\bar{u})^-$
$(u\bar{d})^+$
$(t\bar{t})\,\pi^0\,\eta$
$(b\bar{b})\;\eta'$
$(c\bar{c})$
$(\bar{u}s)^-$
$(\bar{d}s)^0$
$(\bar{c}d)^-$
$(\bar{c}u)^0$
$(\bar{c}s)^-$
$(b\bar{s})^0$
$(b\bar{u})^-$
$(b\bar{d})^0$
$(b\bar{c})^-$
$(\bar{t}c)^0$
$(\bar{t}d)^-$
$(\bar{t}u)^0$
$(\bar{t}s)^-$
$(b\bar{t})^-$

Fig. 27

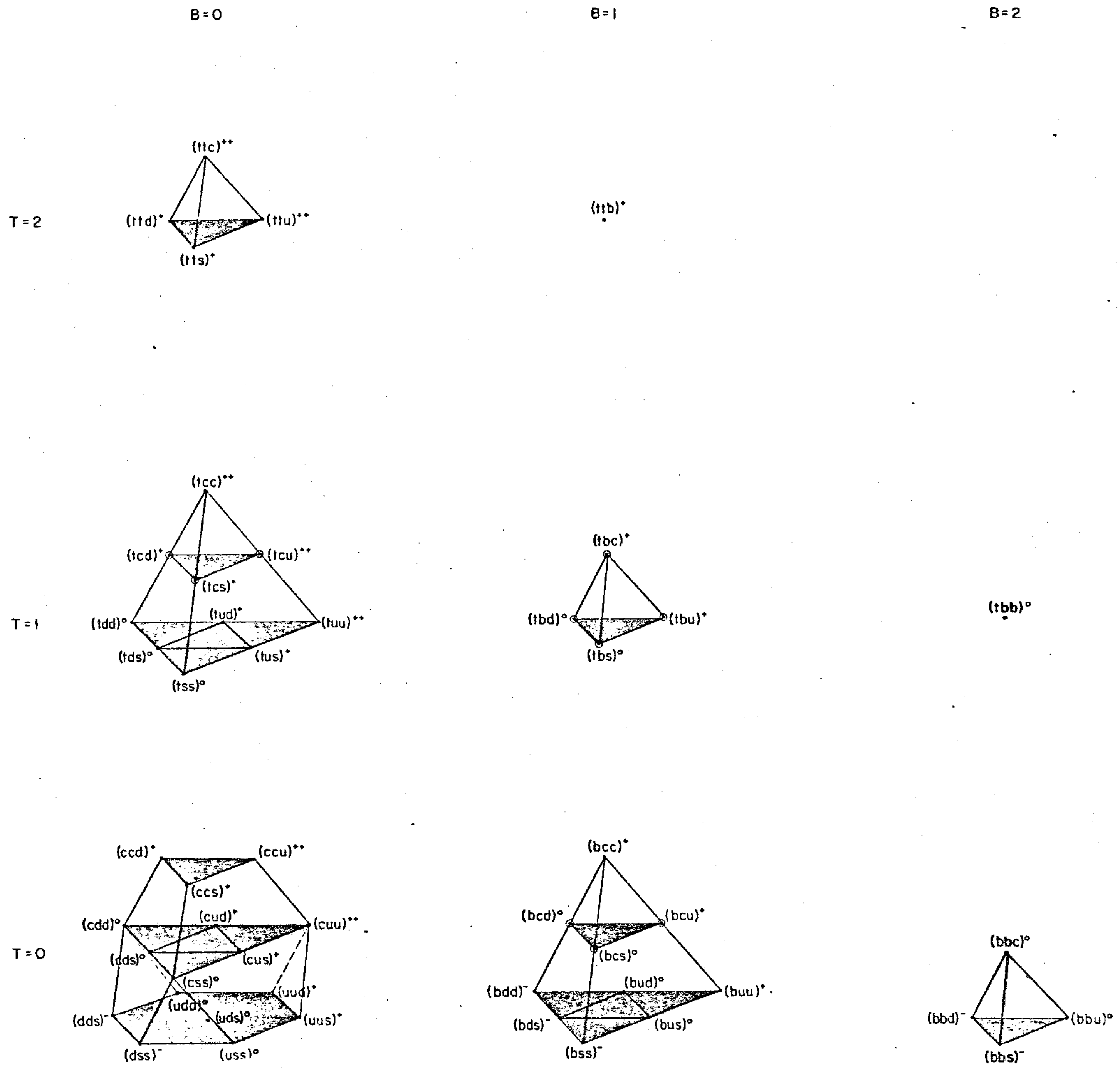
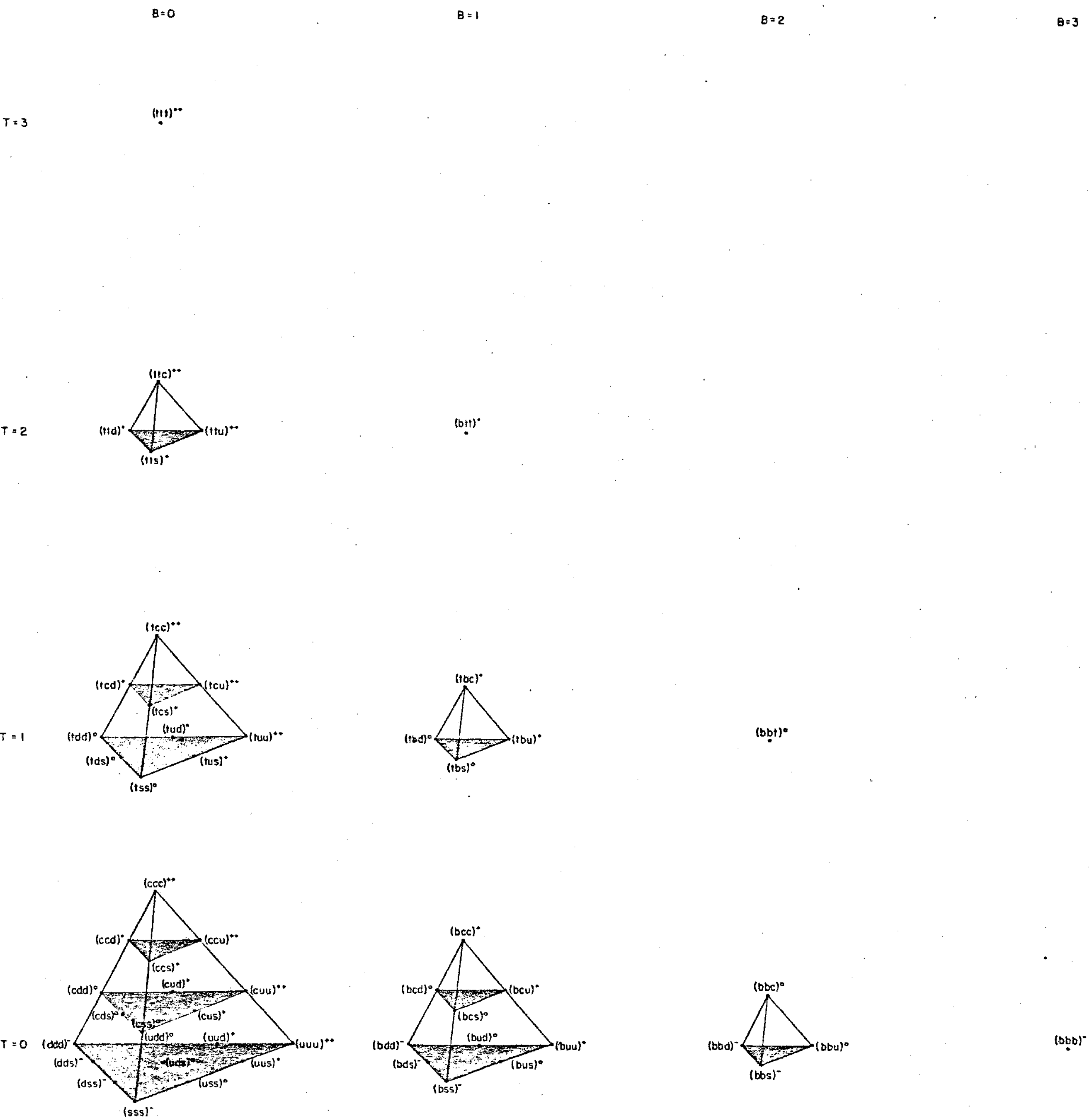

B=0
B=1
B=2
T=2
(ttc)⁺⁺
(ttd)⁺
(ttu)⁺⁺
(tts)⁺
(ttb)⁺
T=1
(tcc)⁺⁺
(tcd)⁺
(tcu)⁺⁺
(tcs)⁺
(tdd)°
(tud)⁺
(tuu)⁺⁺
(tds)°
(tus)⁺
(tss)°
(tbc)⁺
(tbd)°
(tbu)⁺
(tbs)°
(tbb)°
T=0
(ccd)⁺
(ccu)⁺⁺
(ccs)⁺
(cdd)°
(cud)⁺
(cuu)⁺⁺
(cds)°
(cus)⁺
(css)°
(uud)⁺
(udd)°
(uds)°
(dds)⁻
(uus)⁺
(dss)⁻
(uss)°
(bcc)⁺
(bcd)°
(bcu)⁺
(bcs)°
(bdd)⁻
(bud)°
(buu)⁺
(bds)⁻
(bus)°
(bss)⁻
(bbc)°
(bbd)⁻
(bbu)°
(bbs)⁻

Fig. 28



B=0
B=1
B=2
B=3
T=3
(ttt)++
T=2
(ttc)++
(ttd)+
(ttu)++
(tts)+
(btt)+
T=1
(tcc)++
(tcd)+
(tcu)++
(tcs)+
(tud)+
(tdd)°
(tuu)++
(tds)°
(tus)+
(tss)°
(tbc)+
(tbd)°
(tbu)+
(tbs)°
(bbt)°
T=0
(ccc)++
(ccd)+
(ccu)++
(ccs)+
(cud)+
(cdd)°
(cuu)++
(cds)°
(cus)+
(css)°
(udd)°
(uud)+
(ddd)-
(uuu)++
(dds)-
(uds)°
(uus)+
(dss)-
(uss)°
(sss)-
(bcc)+
(bcd)°
(bcu)+
(bcs)°
(bud)°
(bdd)-
(buu)+
(bds)-
(bus)°
(bss)-
(bbc)°
(bbd)-
(bbu)°
(bbs)-
(bbb)-

Fig. 29

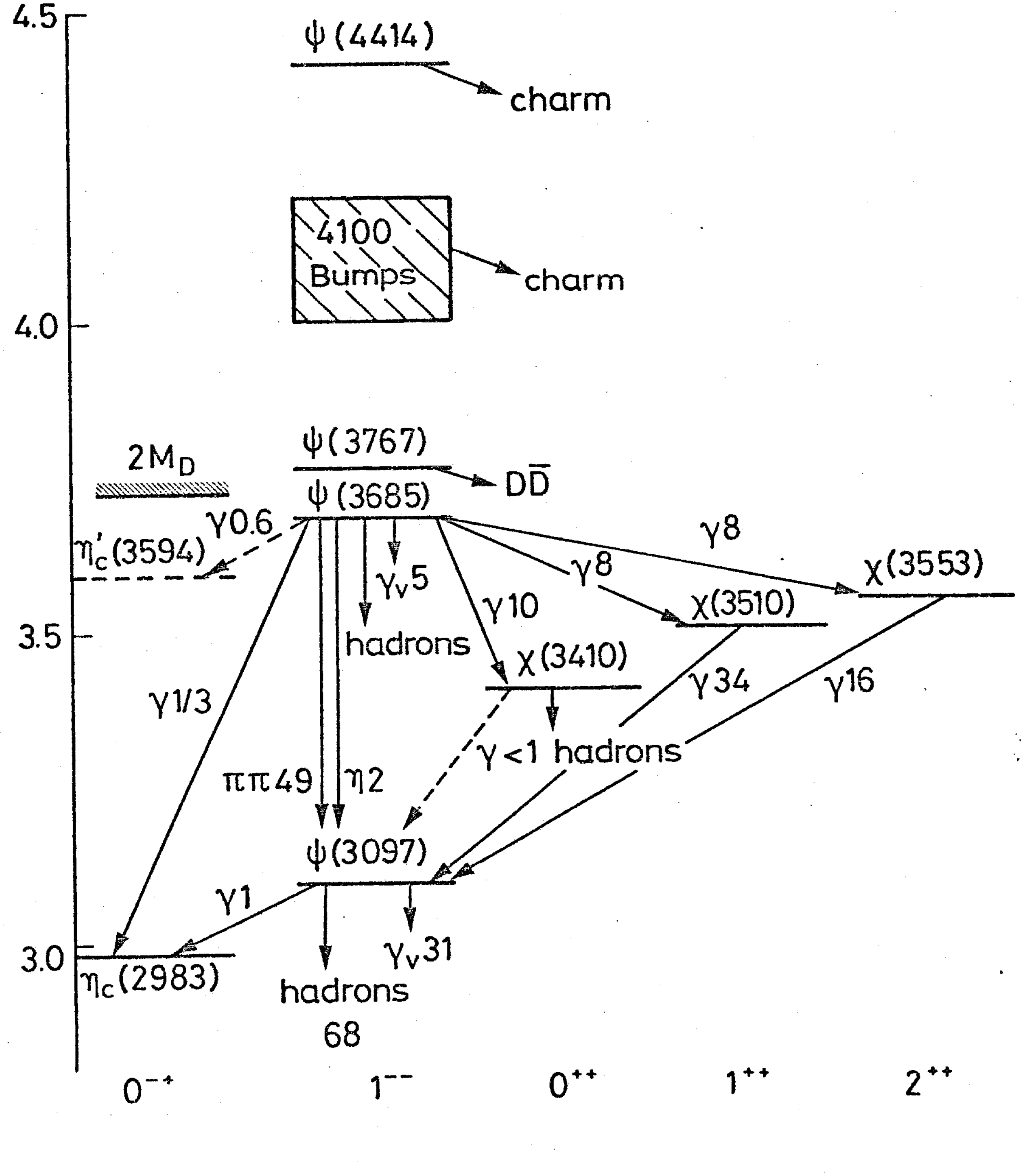
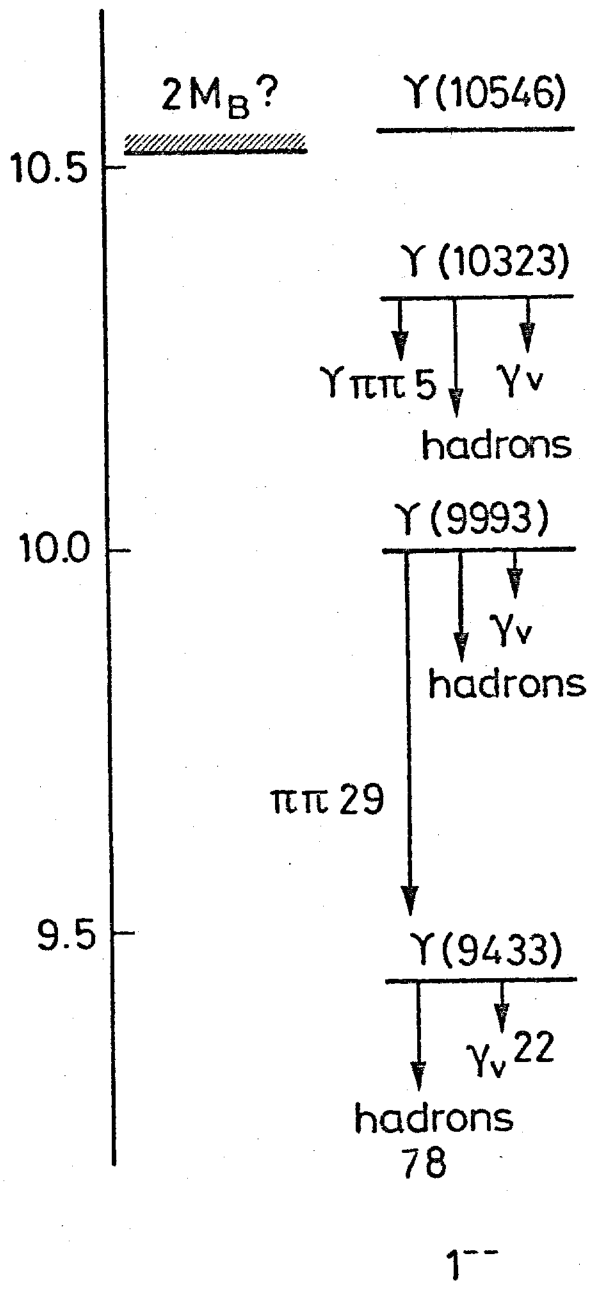
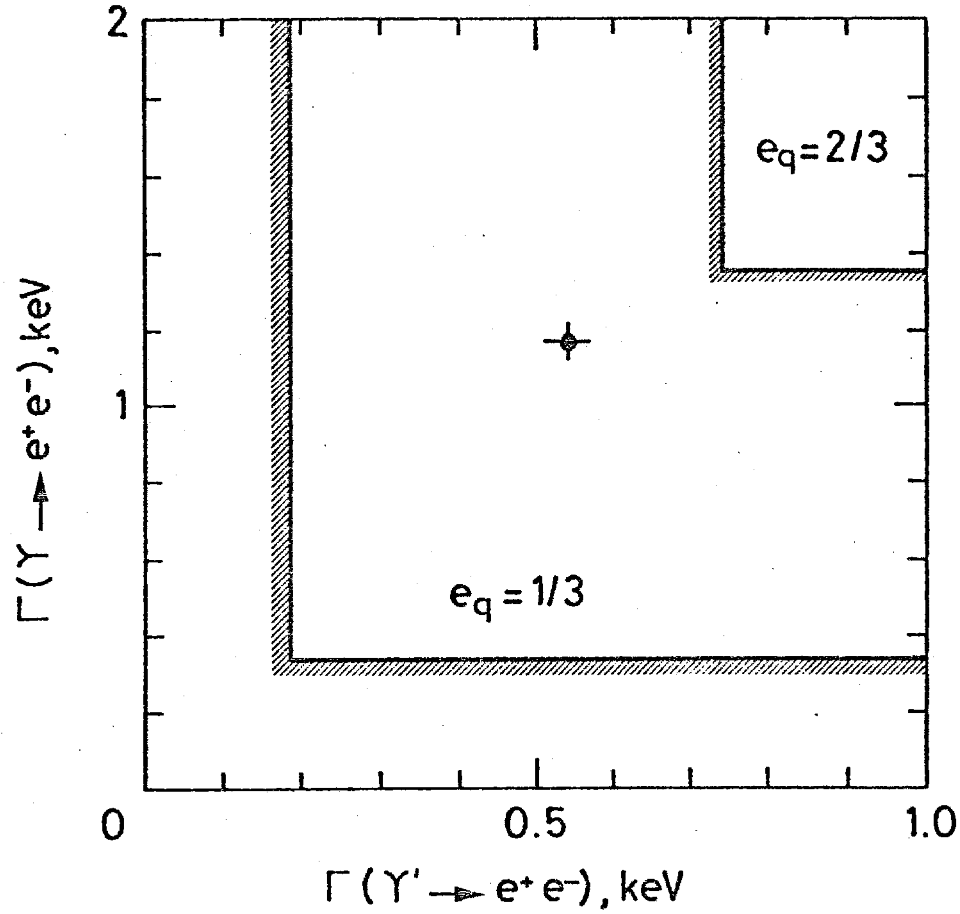

4.5
4.0
3.5
3.0
ψ(4414)
charm
4100
Bumps
charm
ψ(3767)
2M_D
DD̄
ψ(3685)
γ0.6
η′c(3594)
γ_v5
hadrons
γ10
γ8
γ8
χ(3553)
χ(3510)
χ(3410)
γ1/3
γ34
γ16
γ<1 hadrons
ππ49
η2
ψ(3097)
γ1
ηc(2983)
hadrons
68
γ_v31
0⁻⁺
1⁻⁻
0⁺⁺
1⁺⁺
2⁺⁺

Fig. 30



2MB?
Υ(10546)
10.5
Υ (10323)
Υππ 5
γν
hadrons
Υ (9993)
10.0
γν
hadrons
ππ 29
9.5
Υ(9433)
γν 22
hadrons
78
1--

Fig. 31



2
1
0
0.5
1.0
$e_q = 2/3$
$e_q = 1/3$
$\Gamma(\Upsilon \to e^+e^-)$, keV
$\Gamma(\Upsilon' \to e^+e^-)$, keV

Fig. 32

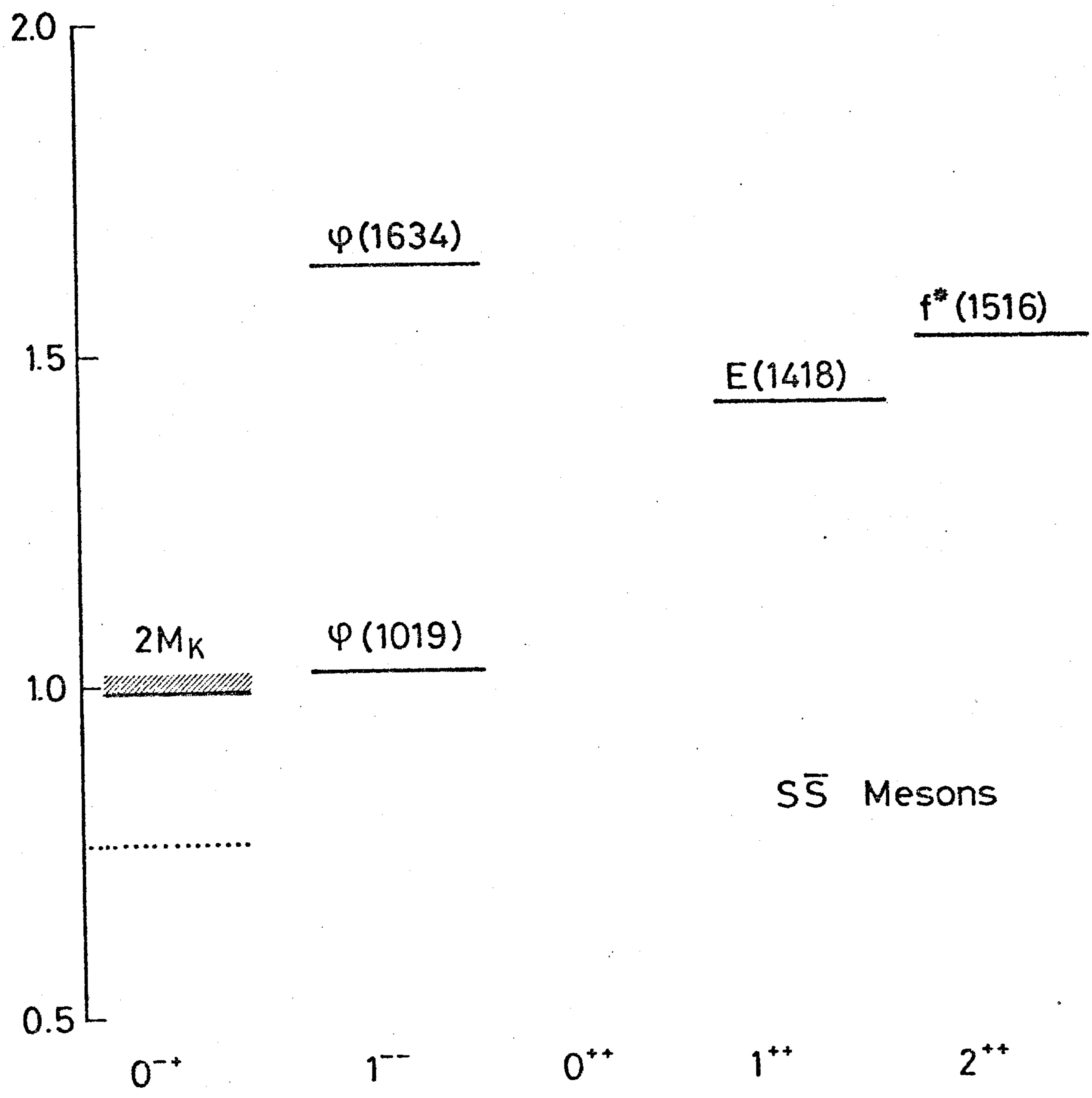


Fig. 33

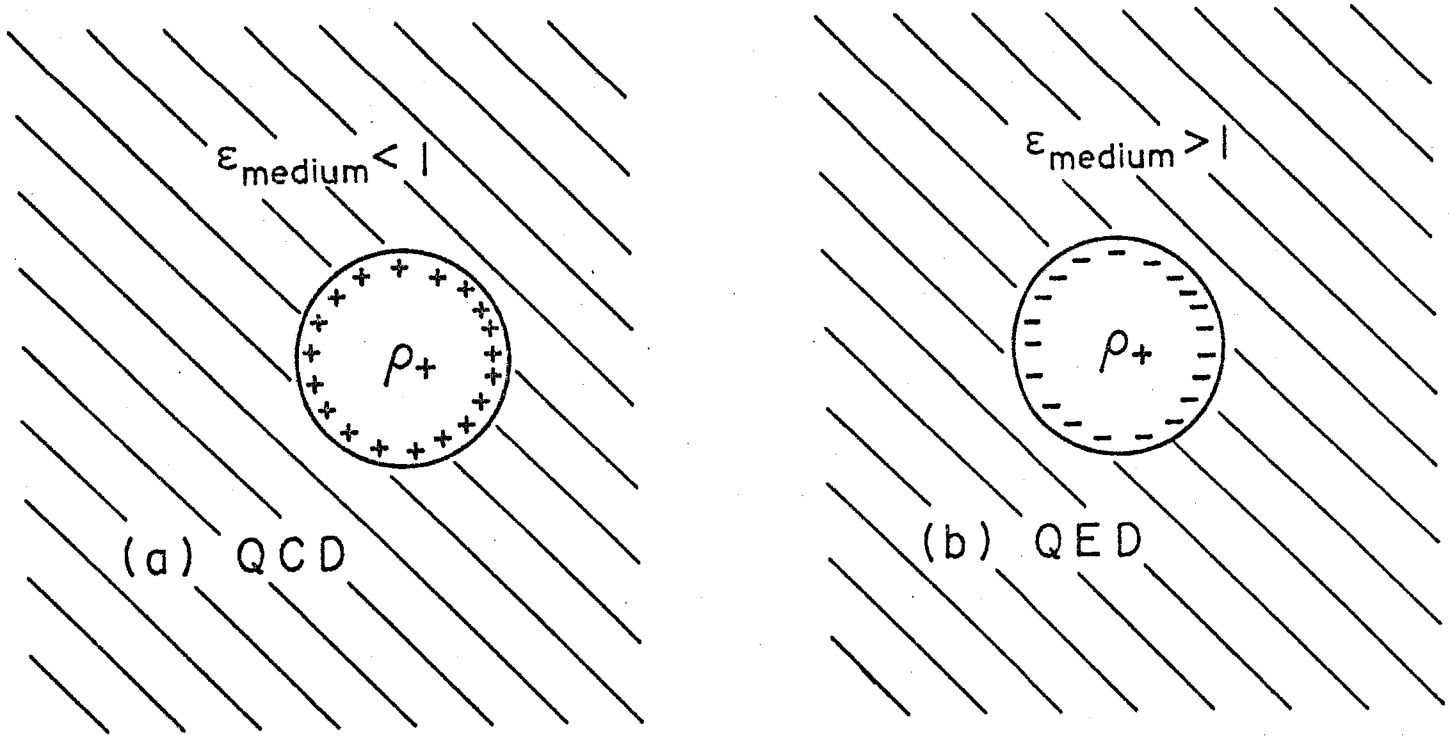


Fig. 34

Fig. 35

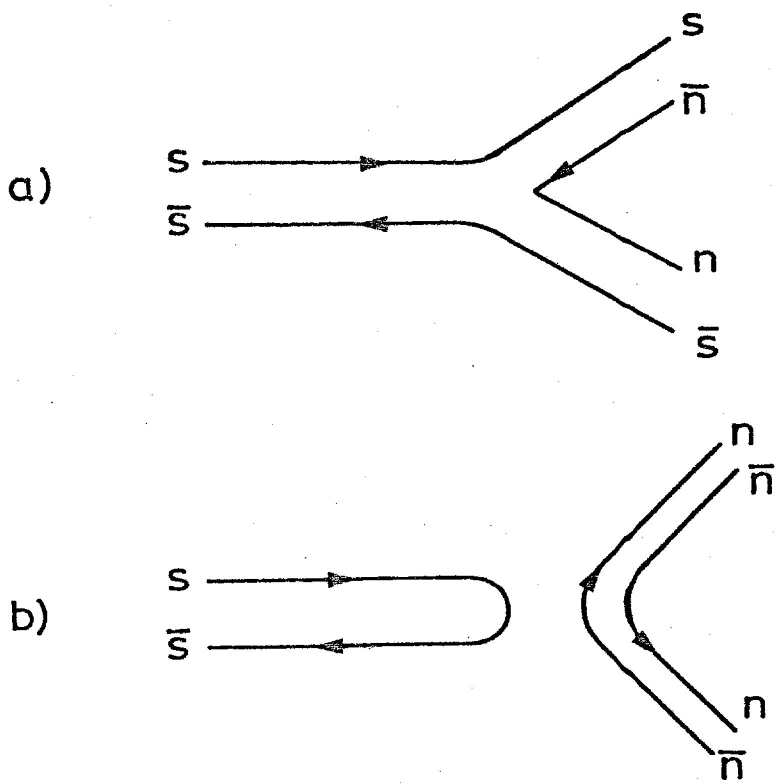
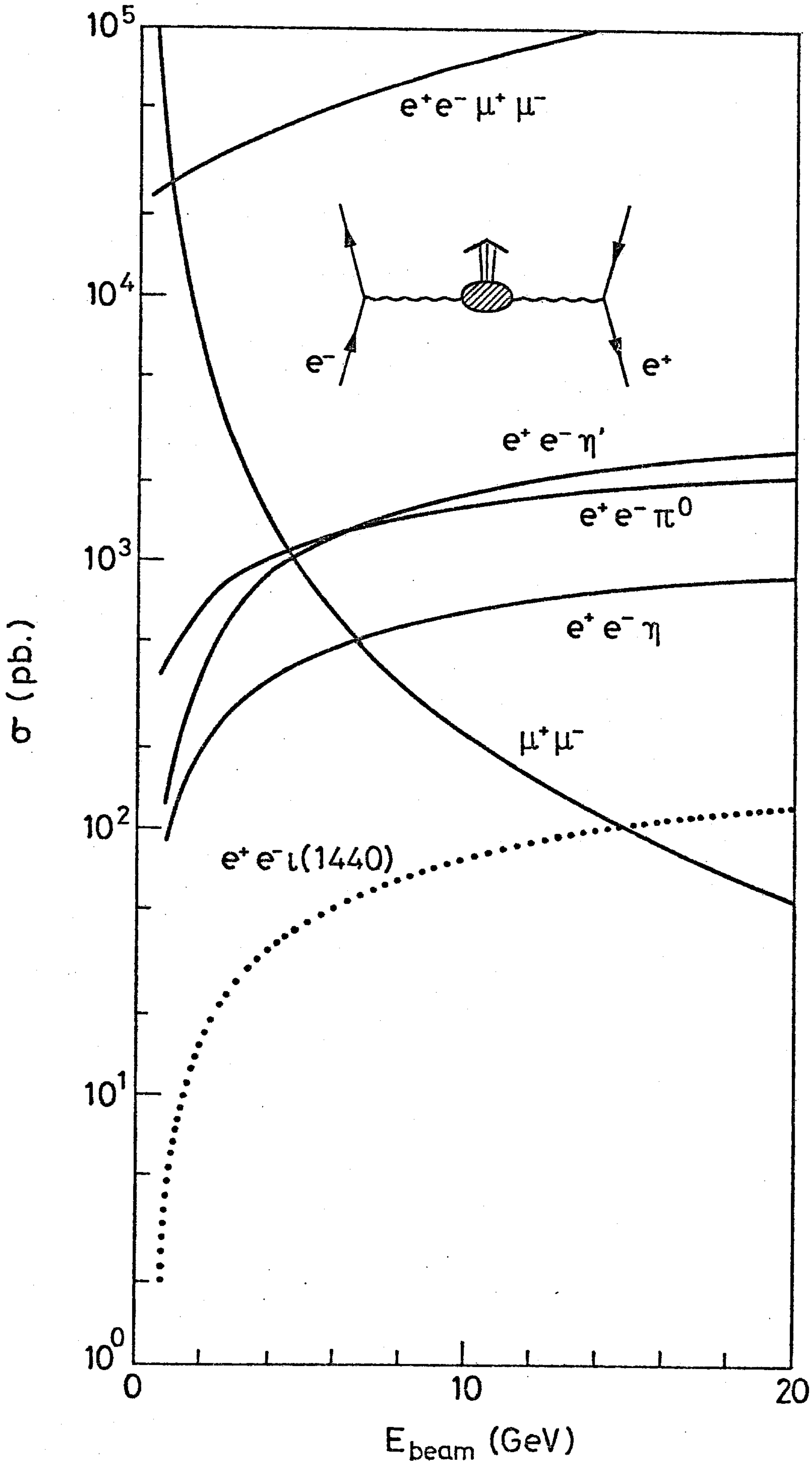


Fig. 36



10^5
10^4
10^3
10^2
10^1
10^0
$e^+e^-\mu^+\mu^-$
e^-
e^+
$e^+e^-\eta'$
$e^+e^-\pi^0$
$e^+e^-\eta$
$\mu^+\mu^-$
$e^+e^-\iota(1440)$
σ (pb.)
0
10
20
E_{beam} (GeV)

Fig. 37

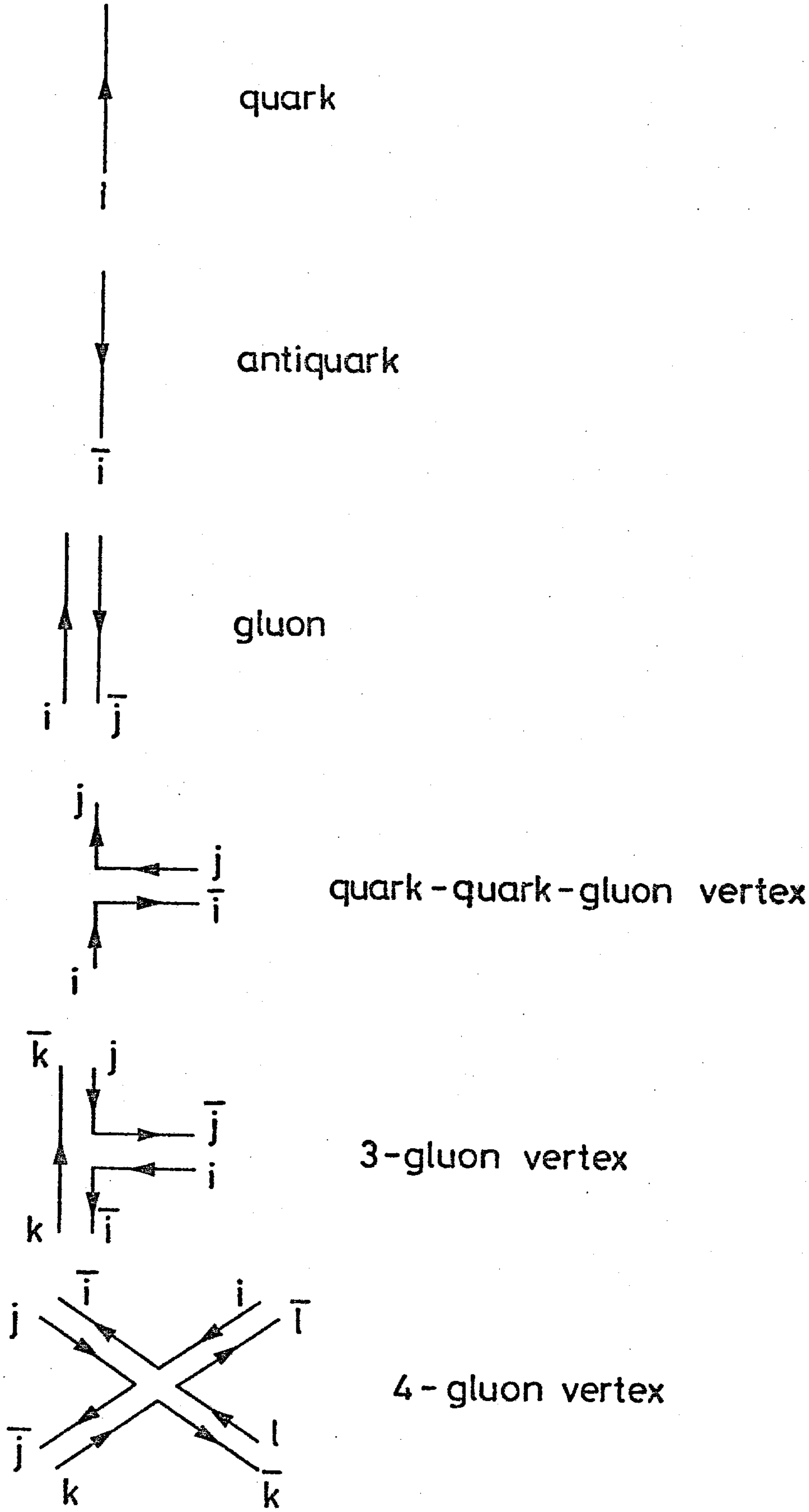

quark
i
antiquark
$\bar{i}$
gluon
i
$\bar{j}$
quark-quark-gluon vertex
j
j
$\bar{i}$
i
3-gluon vertex
$\bar{k}$
j
$\bar{j}$
i
k
$\bar{i}$
4-gluon vertex
$\bar{i}$
i
j
$\bar{l}$
$\bar{j}$
l
k
$\bar{k}$

Fig. 38

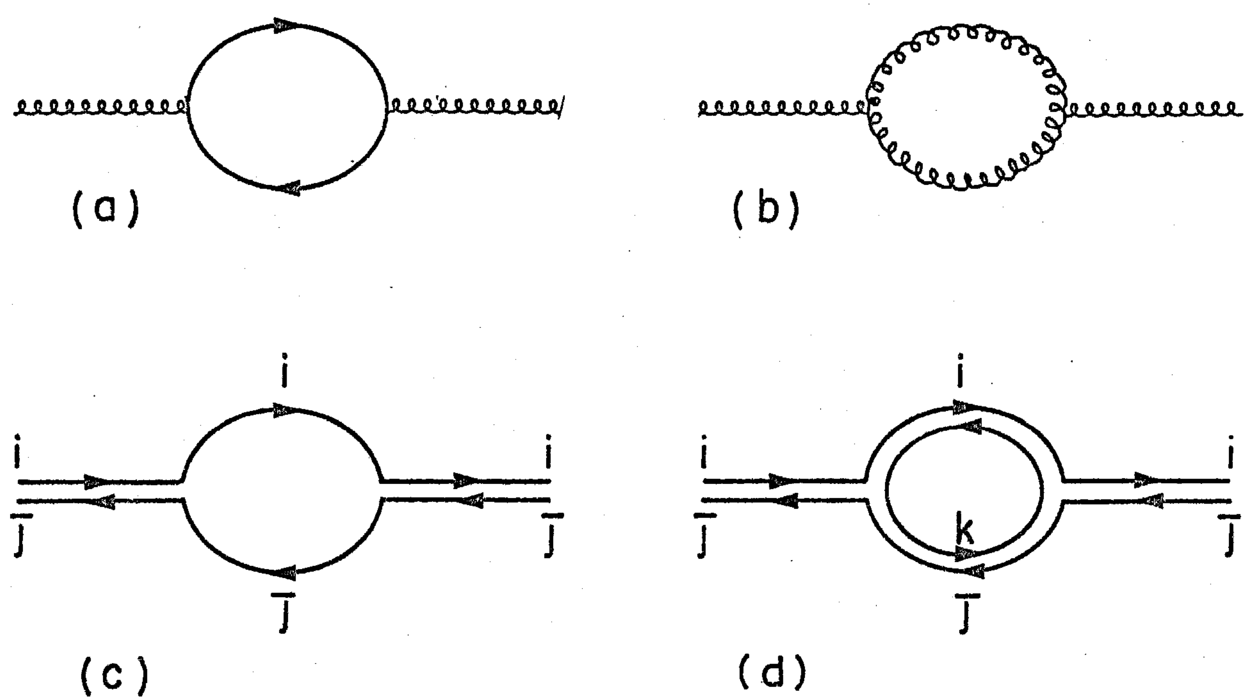
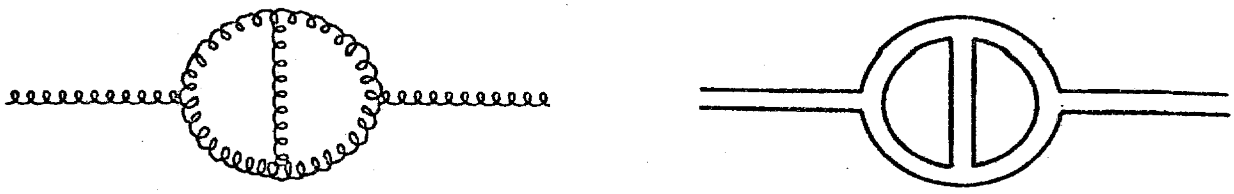
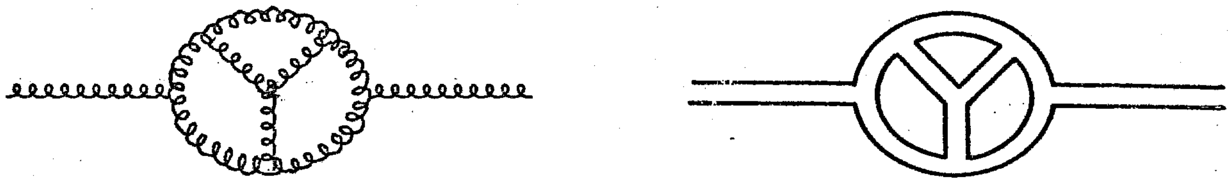


Fig. 39



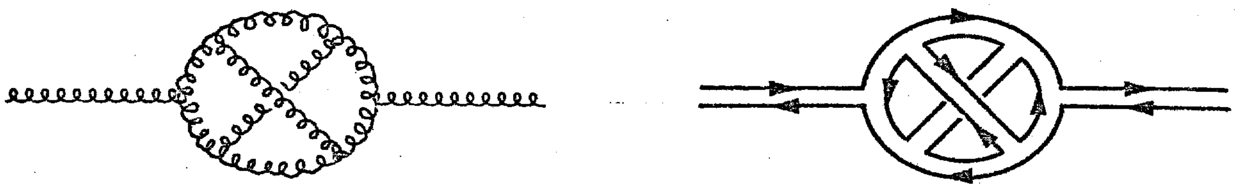
(a) (b)

Fig. 40



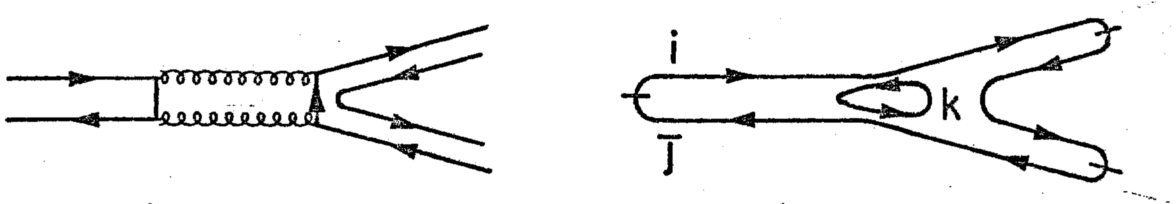
(a) (b)

Fig. 41



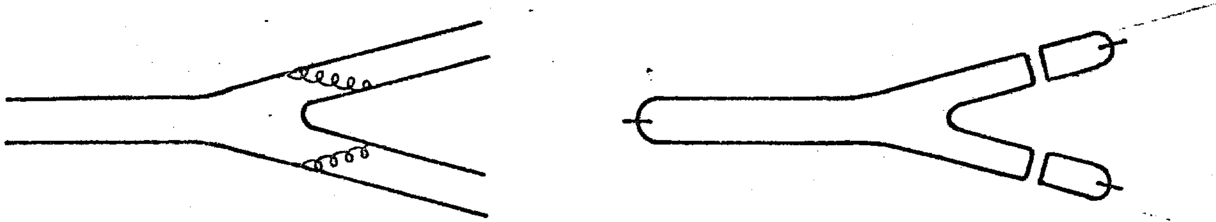
(a) (b)

Fig. 42



(a) (b)

Fig. 43



(a) (b)

Fig. 44