DEUTSCHES ELEKTRONEN-SYNCHROTRON DESY

DESY 78/52
September 1978

$Q^2\bar{Q}^2$ STATES WITH RELATIVELY NARROW WIDTHS

by

Seiji Ono

NOTKESTRASSE 85 · 2 HAMBURG 52

To be sure that your preprints are promptly included in the
HIGH ENERGY PHYSICS INDEX ,
send them to the following address (if possible by air mail) :

DESY
Bibliothek
Notkestrasse 85
2 Hamburg 52
Germany

$Q^2\bar{Q}^2$ STATES WITH RELATIVELY NARROW WIDTHS

by

Seiji Ono$^{+)}$

Deutsches Elektronen-Synchrotron DESY, Hamburg

Abstract:

Using the mass formulas which correctly predict the mass of mesons and baryons the mass of diquark states is computed. From this mass spectrum the existance of the observed narrow baryonia and wide baryonia can be naturally understood. Other relatively narrow $Q^2\bar{Q}^2$ states are predicted to exist.

+) Address from Sept. 19, 1978;
Inst. für Theoret. Physik, E,
TH-Aachen, 51 Aachen,
Germany.

I. Introduction

Over the last few years there have been discovered in experiments a number of new states in the nucleon antinucleon spectrum[1-5] mainly around the $N\bar{N}$ threshold. Some of states have wide width but others are unexpectedly narrow. Since they strongly couple to the antinucleon nucleon system but do not decay into ordinary mesons very fast, they are regarded as candidates of $Q^2\bar{Q}^2$ states.

The reluctance of these states for decay into mesons can be explained[6-8] by the distance between the diquark and the antidiquark caused by the angular momentum barrier. However, this separation is not sufficient to explain the narrow widths of these mesons with large mass, because the rotationally excited states can easily move down to the lower states by emitting π or other mesons. The typical width of this cascade decay is expected to be of the order of 100 MeV from the experience for ordinary meson and baryon resonances. It is very hard to introduce a new selection rule to stop these processes.

The narrow resonance was found[1] as low as 1395 MeV. If this is the S-wave diquonium state, a new selection rule to stop the meson decay must be introduced. However, such selection rule is likely to make too many unobserved baryonia narrow. In this paper we show that due to kinematical reason limited number of diquonium states become narrow.

In §2 the formula to obtain the diquonium state is explained. The S-wave diquonium states and the orbitally excited states are studied in §3 and in §4, respectively. In §5 the concluding remarks are given.

II. One color gluon exchange interaction.

It is known that the mass spectrum of ordinary mesons and baryons can be explained using the nonrelativistic potential although the motion of quarks inside the hadron is relativistic in many cases.

The electromagnetic mass splitting was successfully explained[9,10] by using nonrelativistic wave functions for S-wave baryons. The splitting is caused by the Coulomb and magnetic hyperfine interactions.

The mass splitting in the SU(3) multiplet (or SU(4)) of the S-wave hadrons is caused by one colored gluon exchange. The contributions of the color electric energy to the mass for all mesons are the same and those for all baryons are also the same, because they are all color singlets. Thus, the remaining term is the color magnetic interaction which has a form $\xi\, \vec{\sigma}_1 \cdot \vec{\sigma}_2 / m_1 m_2$. m_1 and m_2 are masses of the quarks. Such a simple interaction is sufficient to predict the correct mass of almost all S-wave hadrons[11].

In order to calculate the mass of the colored diquark we have to write the color dependence explicitely. Let us assume that the mass of the any bound state composed of quarks and antiquarks is described by

$$M = \sum_i m_i + A \sum_{i>j} \lambda_i^a \lambda_j^a - \xi \sum_{i>j} \lambda_i^a \lambda_j^a \frac{\vec{S}_i \cdot \vec{S}_j}{m_i\, m_j} . \tag{1}$$

The first term is the sum of the quark mass. The second term corresponds to the color electric interaction and the third term corresponds to the color magnetic interaction. The color factor $\lambda_i^a \lambda_j^a$ has the following value for the two body system.

$$\lambda_1^a \lambda_2^a = \frac{4}{3} \text{ for } \{6\},\ -\frac{8}{3} \text{ for } \{3\},\ -\frac{2}{3} \text{ for } \{8\},\ \frac{8}{3} \text{ for } \{1\}. \tag{2}$$

For the antiquark the sign of λ_i must be changed. Using the set of parameters[11]

$$m_u = m_d = 336 \text{ MeV}, \quad m_s = 510 \text{ MeV}, \quad m_c = 1680 \text{ MeV}$$

$$\left.\begin{aligned} A_{q-q} &= -\frac{M_b}{8} = -9.6 \text{ MeV}, \\ \frac{\xi_{q-q}}{m_u^2} &= \frac{3}{8} \frac{\xi_b}{m_u^2} = 18.3 \text{ MeV}, \end{aligned}\right\} \text{ for baryon}$$

$$\left.\begin{aligned} A_{q-\bar{q}} &= -\frac{3}{16} M_m = 10.7 \text{ MeV}, \\ \frac{\xi_{q-\bar{q}}}{m_u^2} &= -\frac{3}{64} \frac{\xi_m}{m_u^2} = 29.7 \text{ MeV}, \end{aligned}\right\} \text{ for meson} \tag{3}$$

an excellent agreement with the experimental data for S-wave mesons and baryons including charmed hadron is obtained. The difference of the color electric and color magnetic energy between the baryon and meson can be understood as the difference between the quark-quark wave function and the quark-antiquark wave function.

Chan et al[6] used the same parameters for baryons and mesons. However, if same parameters are used the agreement with the experimental data becomes much worse. It seems to be quite natural to assume that the wave function of the baryon is different from that of the meson. Therefore, we use the different values (3) of parameters A and ξ for the interaction between a quark and a quark from that between a quark and an antiquark.

In the ref.6 the very large value of A (102 MeV) which was estimated by using the less reliable values of diquonium masses was used. In order to assume the large color electric interaction, the large quark mass becomes necessary. We do not want to change the quark mass $m_u = m_d = 336$ MeV which is determined by the magnetic moment of baryons. This mass gives the correct photoproduction and electroproduction amplitudes[10] and the correct mass spectrum for S-wave hadrons[11].

The quasi-static Breit-Fermi-Coulomb potential for the S-state is[12]

$$\alpha_s \sum_{i>j} \lambda_i^a \lambda_j^a \left\{ \frac{1}{r} - \frac{\pi}{2}\left(\frac{1}{m_i^2} + \frac{1}{m_j^2}\right) \delta(\vec{r}_i - \vec{r}_j) + \cdots \right\} . \tag{4}$$

If we use the harmonic-oscillator wave function with the spring constant determined by the electromagnetic mass splitting of the baryon[10] to obtain the expectation value of this interaction, the second term becomes as large as the first in magnitude but with the opposite sign. This means that this expansion has no meaning and we cannot determine even the sign theoretically. We simply take the phenomenological value (3). In our model the color magnetic energy is much more important than the color electric energy. From eq.(1) one gets the following mass of the S-wave diquark.

$$\left.\begin{array}{l}\alpha(I=1,\ S=0,\ \text{color } 6)\ ;\ 2m_u+\frac{4}{3}A_{8\text{-}8}+4\ \frac{\xi_{8\text{-}8}}{m_u^2}=732\ \mathrm{MeV},\\ \alpha^*(I=1,\ S=1,\ \text{color } \bar{3})\ ;\ 2m_u-\frac{8}{3}A_{8\text{-}8}+\frac{8}{3}\ \frac{\xi_{8\text{-}8}}{m_u^2}=747\ \mathrm{MeV},\\ \beta(I=0,\ S=0,\ \text{color } \bar{3})\ ;\ 2m_u-\frac{8}{3}A_{8\text{-}8}-8\ \frac{\xi_{8\text{-}8}}{m_u^2}=552\ \mathrm{MeV},\\ \beta^*(I=0,\ S=1,\ \text{color } 6)\ ;\ 2m_u+\frac{4}{3}A_{8\text{-}8}-\frac{4}{3}\ \frac{\xi_{8\text{-}8}}{m_u^2}=634\ \mathrm{MeV}\end{array}\right\} \quad (5)$$

III. The S-wave diquonium state

Since we do not introduce any kind of the new selection rule, the S-wave diquonium states must decay into meson states fast and it is probable that they are not resonances at all. For the S-wave diquonium decay into meson states no quark pair is needed to be created and this decay mode is called superallowed.

If a diquark and an antidiquark get close each other the color electric and color magnetic interaction between the diquark and the antidiquark becomes important. We estimate this energy using eq.(1). We assume that inside the S-wave $Q^2\bar{Q}^2$ system the rms distance between two quarks equals that between a quark and an antidiquark. Roughly speaking two quarks and two antiquarks make a tetrahedron (Fig.1) if we neglect the relative motion in the S-state.

Let us assume that i = 1 and 2 expresses quarks and i = 3 and 4 expresses antiquarks. One gets the following exppectation value.

$$\lambda_1^a \lambda_3^a = \begin{cases} \frac{10}{3} & \text{for } \{6 - \bar{6}\} \\ \frac{4}{3} & \text{for } \{\bar{3} - 3\} \end{cases} \quad (6)$$

It is easy to calculate the mass of the diquonium by using eq.(1) and parameters (2). The mass spectrum is shown in Fig.2.

In the S-wave diquonium state the difference between the T-diquonium (color 3 x $\bar{3}$) and the M-diquonium is not very clear. They are usually mixed. This four body bound state is a system of two quark-antiquark pairs ($q\bar{q}$ - $q\bar{q}$) and at the same time this is a diquark-antidiquark system

$(qq - \overline{qq})$. They are mixed and there exists no special pair in this four body system as shown in Fig.1. For simplicity we classify these states as the diquark-antidiquark systems.

As is seen from Fig.2 there exists a very strong spin-spin force. If spins of a β^* diquark and $\bar{\beta}^*$ diquark are parallel, the mass of $\beta^*\bar{\beta}^*$ is approximately equal to 2 x mass of ρ meson $\sim$ 1500 MeV and if they are antiparallel, the mass is roughly equal to 2 x mass of π meson $\sim$ 300 MeV. If the β^* and $\bar{\beta}^*$ diquarks are brought close each other from distance with the spin parallel, the interaction between these two particles are very little besides confining potential, because the mass of $[\beta^*\bar{\beta}^*]_{J=2}$ is only a little bit larger than the sum of masses of β^* and $\bar{\beta}^*$. If these two are brought close each other with the spin antiparallel, they attract very strongly and due to the huge binding energy the mass of the diquark almost cancelled. Similar situation exists for $\alpha^*\bar{\alpha}^*$ diquonium states but the interaction between them is smaller than that for the $\beta^*\bar{\beta}^*$.

Super allowed decays of S-wave diquonium states are shown in Fig.2. These decays will take place very fast. According to the optical model as the imaginary part gets large, the resonance and bound state pole move away from the physical region and these are destroyed[13]. (However, Kerbikov et al[14] obtained the different result.) It is quite possible that some of the S-wave diquonia are not resonances at all.

However, the situation is different for the diquonium states with J = 2. Let us consider the $\beta^*\bar{\beta}^*$, $I^G(J^P) = 0^+(2^+)$ state. This particle is not easy to decay into $\pi\pi$ because the quark spin flip is necessary twice and the produced $\pi\pi$ must be in hte D-state. Decay into 3π, $\eta\pi$ and $\rho\pi$ are forbidden by the G parity conservation and the decay into $\omega\pi$ is forbidden by the isospin consevation. Decays into $\rho\rho$ and $\omega\omega$ are forbidden kinematically. However, due to the large width of ρ, the decay into $\rho\rho$ will not be completely forbidden. Probably this resonance mainly decay into $\rho 2\pi$. This is not superallowed but an allowed decay which needs one quark pair creation between a quark and an antiquark. The width can be as large as 50-100 MeV.

A part of this decay process can take place through a cascade decay in the following way.

$$[\beta^*\bar{\beta}^*]_{J=2} \rightarrow [\alpha\bar{\beta}^* + \bar{\alpha}\beta^*]_{J=1} + \pi \qquad (7)$$
$$\hookrightarrow \rho\pi$$

This process is allowed kinematically (See Fig.2). In short the decay $\beta^* \to \alpha + \pi$ takes place inside the diquonium. In order to get a crude estimate of this decay rate, let us assume that π is produced by one of quarks inside the diquark. Assuming that the initial and final diquarks are free and using the static approximation we obtain

$$\Gamma_{\alpha^* \to \beta\pi} \approx \frac{k^3}{2\pi}\left(\frac{f_q}{m_\pi}\right)^2 \frac{4}{3}, \quad f_q = \frac{3}{5} f_{\pi N}, \quad \frac{f_{\pi N}^{\;2}}{4\pi} = 0.082 . \tag{8}$$

i.e.,

$$\Gamma_{\alpha^* \to \beta\pi} \approx (4.3 \ \mathrm{GeV}^{-2}) \times k^3 . \tag{9}$$

where k is the momentum of the emitted π . As for the explanation of this method see refs.15-17.

If the free α^* diquark (mass = 747 MeV) decays into the free β diquark (mass = 552 MeV), one gets k = 117 MeV and $\Gamma_{\alpha^* \to \beta\pi}$ = 6.9 MeV. If this decay takes place in the cascade decay (7), the emitted π has momentum k = 236 MeV. Assuming that eq. (9) still works one gets $\Gamma_{\beta^*\bar{\beta}^* \to \alpha\bar{\beta}^* + \pi}$ $\sim$ 56 MeV. Probably this will be overestimated because the decay process occurs inside the diquonium. Other cascade decay is not allowed except $[\beta^*\bar{\beta}^*]_{J=2} \to [\beta^*\bar{\beta}^*]_{J=1} + \omega$ which has a very small phase space.

The resonance $\beta^*\bar{\beta}^*$ (M-diquonium) is very likely to mix with $\alpha^*\bar{\alpha}^*$ (T-diquonium) with the same quantum number. The latter decays almost in the same way exept that this moves down to the $[\alpha^*\bar{\beta} \pm \bar{\alpha}^*\beta]_{J=1}$ state by emitting π .

The physical states are given by

$$\left.\begin{aligned} G_1 &= \cos\theta \,|\beta^*\bar{\beta}^*\rangle - \sin\theta \,|[\alpha^*\bar{\alpha}^*]_{I=0}\rangle \\ G_2 &= \sin\theta \,|\beta^*\bar{\beta}^*\rangle + \cos\theta \,|[\alpha^*\bar{\alpha}^*]_{I=0}\rangle \end{aligned}\right\} \tag{10}$$

For $\theta = -60^\circ$ G_1 cannot go to $\rho\rho$ and G_2 cannot go to $\omega\omega$ in a superallowed way. In this case only $G_1 \to \omega\omega$ or $\omega 3\pi$ are allowed but the decay rate is small and G_1 can be narrow.

The $\alpha^* \bar{\alpha}^*$,$2^+(2^+)$ has similar properties and mainly decay into $\rho\rho$ or $\rho 2\pi$. It can also go to $[\alpha^* \bar{\alpha}^*]_{J=1}$ state by emitting π . The $\alpha^* \bar{\alpha}^*$,$1^-(2^+)$ will decay into $\rho\omega$ or $\omega\pi\pi$. The decay $\rho\pi$ be somewhat suppressed, because the final π must have angular momentum 2.

Our conclusion is that in the S-wave diquonium states the one with spin 2 will have smaller decay width and these will be seen around the threshold of $\rho\rho$ or $\rho\omega$.

IV. The orbitally excited state.

For the quark-antiquark system the mass splitting of states within the same non zero orbital angular momentum is caused only by the spin-orbit force (e.g., A_2, A_1, B and δ), because the spin-spin force is proportional to $\delta(\vec{r}_1-\vec{r}_2)$ and does not contribute to the ℓ-excited state. For the diquark-antidiquark system the spin-spin force does not become zero because the diquark is not a point particle.

Let us assume that the wave function of the diquark-antidiquark in the ℓ-wave is given by

$$\Psi_\ell^m(\vec{r},\vec{r}_d,\vec{r}_a) = Y_\ell^m(\Omega) R_\ell(r) Y_0^0(\Omega_d) R_0(r_d) Y_0^0(\Omega_a) R_0(r_a) , \tag{11}$$

where $\vec{r}$ is the distance between the center of mass of the diquark and that of the antidiquark. $\vec{r}_d(\vec{r}_a)$ is the distance between the center of mass of the diquark (antidiquark) and one of the quark(antiquark) inside the diquark (antidiquark). Let us use the harmonic-oscillator wave function for $R_\ell(r)$, $R_0(r_d)$ and $R_0(r_a)$ because the calculation is simple and analogous results will be expected for other potentials.

$$\left.\begin{aligned} R_\ell(r) &= \sqrt{\frac{2^{\ell+2}}{\sqrt{\pi}\,(2\ell+1)!!\,R^{2\ell+3}}}\; r^\ell\, e^{-\frac{r^2}{2R^2}} \\ R_\ell(r_{d,a}) &= \sqrt{\frac{2^{\ell+2}}{\sqrt{\pi}\,(2\ell+1)!!\,r_0^{2\ell+3}}}\; (r_{d,a})^\ell\, e^{-\frac{(r_{d,a})^2}{2r_0^2}} \end{aligned}\right\}$$

The expectation value of $\delta(\vec{r}_1 - \vec{r}_3) = \delta(\vec{r} + \vec{r}_a - \vec{r}_d)$ is

$$I_\ell \equiv \int |Y_\ell^m(\Omega)\, R_\ell(r)|^2 \times |Y_0^0 R_0(r_d)|^2 \times |Y_0^0 R_0(r_a)|^2 \delta(\vec{r}+\vec{r}_a-\vec{r}_d)\, d\vec{r}\, d\vec{r}_d\, d\vec{r}_a$$

$$= \frac{1}{\sqrt{(2\pi)^3 \left(1 + \frac{R^2}{2r_0^2}\right)^{2\ell+3} r_0^3}} \qquad (13)$$

Therefore,

$$K_\ell \equiv \frac{I_\ell}{I_0} = \frac{1}{\left(1 + \frac{1}{2}\frac{R^2}{r_0^2}\right)^\ell} \qquad (14)$$

is the damping factor of the color magnetic force. The ratio R/r_0 is not known and must be determined phenomenologically. From the geometrical consideration ane gets $R = \sqrt{2}\, r_0$ by using Fig.1. From the condition that the rms radius between two quarks is equal to that between a quark and an antiquark one obtains $R = \sqrt{2/3}\, r_0$ for the later calculation because we need large distance between the diquark and the antidiquark for the ℓ-excited state to stop the meson decay. For damping factor of the color electric force is much more complicated. For simplicity we use the same damping factor eq.(14). The color electric force is not very important in our model anyway.

The spin-orbit force can be written as

$$H_{LS} = \eta_\ell\, \lambda_1^a \lambda_2^a\, \vec{\ell} \cdot (\vec{S}_1 + \vec{S}_2 + \vec{S}_3 + \vec{S}_4) \qquad (15)$$

The coefficient η_ℓ sen itively depends on the form of the potential which we do not know. The mass splitting among A_2, A_1, B, and δ are caused by the spin-orbit force between the quark and the antiquark.

$$H'_{LS} = \eta'_\ell\, \lambda_1^a \lambda_2^a\, \vec{\ell} \cdot (\vec{S}_1 + \vec{S}_2)\,,\ \lambda_1^a \lambda_2^a = \frac{16}{3} \quad \text{for} \quad \{1\} \qquad (16)$$

Compared to the experimental data one gets η'_1 = 21 MeV and η'_2 = 3.3 MeV. In order to calculate the diquonium mass splitting we simply put $\eta_\ell = \eta'_\ell$, because this will give a good explanation of the baryonium spectrum. We find that the spin-spin force is much more important than the spin-orbit force even for the P- and D-states.

Finally we have to determine the level spacing between the S-state and the P-state and that between the S-state and the D-state. We can roughly estimate the level spacing by making comparison with the quark-antiquark mass levels, since the potential is caused by the gluon exchange. The gluon couples with color. As long as color is the same, the potential is not much different. From the mass spectra of nonstrange mesons, strange mesons, charmonia, and bottonia we know that the level splitting depends on the mass of the quark weakly. We assume $P - S = 550$ MeV and $D - S = 844$ MeV for the T-diquonium.

For the M-diquonium potential we do not have any imformation. The narrow width of the diquonium state can only be explained by the small spacing. Therefore, we assume $P - S = 240$ MeV and $D - S = 374$ MeV for M-diquonium.

In this way we obtain the mass spectrum of the diquonium states. Results are shown in Fig.3 and Fig.4 for the P-state and for the D-state, respectively. We do not show the mass spectrum of the 2S-state because no narrow resonance is likely to exist.

(a) <u>The P-state</u>

The superallowed decay rate is not so large as the S-state because of the centrifugal barrier. However, most states have large widths because they can drop down to lower states by emitting π or other mesons. Since π or other mesons are a color singlet, the transition between M- and T-diquonium not allowed. The mixing of T and M diquonium decreases as the angular momentum increases, because the color changing ($\bar{3} \to 6$ + gluon, $6 \to \bar{3}$ + gluon) one gluon exchanging is of short range.

The $[\alpha^* \bar{\alpha}^*]_{\ell=1}$ diquonium states have the widest widths (more than 100 MeV) because they are the heaviest in the P-state and the constraint which comes from the isospin is lenient. Their masses are large enough to decay into $N\bar{N}$ states creating a quark pair. Similary $[\alpha^*\bar{\beta} \pm \bar{\alpha}^*\beta]_{\ell=1}$ and $[\alpha\bar{\beta}^* \pm \bar{\alpha}\beta^*]_{\ell=1}$ states have wide widths.

The $[\beta^*\bar{\beta}^*]_{\ell=1}$ diquonium states have narrower widths than others. A few of them with large mass can move down to $[\alpha\bar{\beta}^* \pm \bar{\alpha}\beta^*]_{\ell=0,1}$ by emitting π and the width can be relatively large ($\lesssim$ 50 MeV). The

$[\beta^* \bar{\beta}^*]_{\ell=1}$ with mass less than 1400 MeV can only drop down to lower $[\beta^* \bar{\beta}^*]_{\ell=0}$ emitting ω or η mesons.

The $[\beta^* \bar{\beta}^*]_{\ell=1}$, $0^+(1^-)$, 1174 MeV, state is especially interesting because this quantum number does not exist in quark-antiquark states. Any cascade decay mode are forbidden by G parity conservation etc. This has the small decay rate ($\sim$ several MeV) to $\rho 2\pi$ and 2η. The decay process

$$\beta^* \bar{\beta}^* (1174) \rightarrow 2 \times \beta^* \bar{\beta}^* (313)$$
$$\hookrightarrow 2\pi$$

will be also possible to occur but the decay rate will be small ($\sim$ a few MeV) because it needs the quark pair creation twice.

The $[\beta^* \bar{\beta}^*]_{\ell=0}$, $0^-(1^-)$, 1046 MeV, state also cannot make any cascade decay. The decay rate to $\rho\pi$ will be small ($\sim$ a few MeV) because of the separation between two diquarks due to the angular momentum. Other low lying $\beta^* \bar{\beta}^*$ resonances $0^+(0^-)$ and $0^+(2^-)$ can drop down to $[\beta^* \bar{\beta}^*]_{\ell=0}$, $0^+(0^+)$ state by emitting η. Therefore, these two are somewhat broader than $0^+(1^-)$ and $0^-(1^-)$ states but they will be still narrow ($\lesssim$ 10 MeV).

(b) The D-state

The mass spectrum is analogous to the P-states. The splitting due to a color gluon exchange between the diquark and antidiquark becomes small.

The $\alpha^* \bar{\alpha}^*$, $\alpha^* \bar{\beta} \pm \bar{\alpha}^* \beta$ and $\alpha \bar{\beta}^* \pm \bar{\alpha} \bar{\beta}^*$ resonances have large decay widths because of cascade decays and the decay to $N\bar{N}$ by creating one quark pair.

The $\beta^* \bar{\beta}^*$ and $\beta \bar{\beta}$ states have narrower widths than others, because of the kinematical restriction and the constraint which comes from the isospin. The $\beta \bar{\beta}$ state has a small decay rate to $N\bar{N}$ and the $\beta^* \bar{\beta}^*$ state has a small decay rate to $2 \times [\beta^* \bar{\beta}^*]_{\ell=0}$ creating two quark pairs. Pavlopoulos et al[1] found narrow resonances of mass 1684 MeV, 1646 MeV and 1395 MeV. Assuming that atomic S-states as initial states and that they are electric dipole transitions these resonances must be $J^{PC} = 2^{++}$, 1^{++}, 0^{++} or 1^{+-}. Let us assign these states to the $\ell = 2$ $\beta^* \bar{\beta}^*$ states. Five states

($0^+(2^+)$, $0^+(1^+)$, $0^+(0^+)$, $0^+(2^+)$ and $0^-(1^+)$) out of seven can correspond to them.

The 1684 MeV resonance corresponds to the $0^+(2^+)$ state. The $0^+(1^+)$ and $0^+(0^+)$ resonances locates very close each other and they will be observed as one resonance in their experiment because the spectrometer resolution is 19 MeV. These two resonances (1642 MeV and 1653 MeV) corresponds to the 1650 MeV peak. The $0^+(2^+)$, 1411 MeV and $0^-(1^+)$, 1378 MeV resonances will be observed as one peak because the experimental resolution is 34 MeV at this energy. These corresponds to 1395 MeV peak.

The $0^+(4^+)$, 1750 MeV and $0^+(3^+)$, 1710 MeV resonances must be also observed in their γ-ray energy spectrum. However, these resonances must be produced by E3 and E2 transitions and the signal will be much smaller. Acutually they found a peak at 1739 MeV which was explained as the $\pi^- P$ radiative capture. The $0^+(4^+)$ and $0^+(3^+)$ resonances must be hidden around this peak.

We do not know the precise production mechanism of the diquonium from the atomic $P\bar{P}$ state. If the form of the interaction allows only to couple to the T-diquonium, the production of $\beta^*\bar{\beta}^*$ diquonium is understood by the small mixing with the $\alpha^*\bar{\alpha}^*$ diquonium with the same quantum number. In this energy region there is no other narrow resonance.

The cascade decay modes of the $\beta^*\bar{\beta}^*$ resonances are shown n Fig.4. The reason why these have much smaller widths than $\alpha\bar{\alpha}$, $\alpha^*\bar{\alpha}^*$, $\alpha^*\bar{\beta} \pm \bar{\alpha}^*\beta$ and $\alpha\bar{\beta}^* \pm \bar{\alpha}\beta^*$ are the following. (i)The $\Delta\ell = 0$ transition is forbidden kinematically. (ii)For the most resonances the $\Delta\ell = 1$ transition is also forbidden kinematically. Some resonances can drop down to the P-stste with the small phase space. The decay rate will not very large. (iii) The $\Delta\ell = 2$ transitions are allowed. The decay rate depends sensitively on the form of the potential. If the diquark-antidiquark potential is of the harmonic-oscillator type, the decay rates become 10 - 100 MeV. If the potential is of the 1/r type, the decay rate, which is proportional to the square of the overlapping integral, becomes about one twentieth for the $\Delta\ell = 1$ transition and one hundredth for the $\Delta\ell = 2$ transition (e. g. see Figs.3,4 in Ref. 18) compared to those for the harmonic-oscillator potential. Therefore, they can be narrow if the confining potential is not very sharply rising in the outer resion like the harmonic-oscillator type. These resonances have the following small decay rate into two diquania.

$$\left.\begin{array}{l}\beta^* \bar{\beta}^* (1378) \to \beta^* \bar{\beta}^* (739) + \beta^* \bar{\beta}^* (313),\ 1400 \to 739 + 313, \\ 1411 \to 313 + 313,\ 1413 \to 313 + 739,\ 1642 \to 313 + 313,\ 1653 \to 739 + 739 \\ 1675 \to 313 + 313,\ 1708 \to 739 + 739,\ 1752 \to 313 + 313.\end{array}\right\} \quad (17)$$

The $\beta\bar{\beta}$, $0^+(2^{++})$, 1935 MeV resonance is also narrow by the similar reason. (i) The $\Delta\ell = 0$ transition is not allowed. (ii) The only allowed $\Delta\ell = 1$ transition $\beta\bar{\beta} \to [\alpha^*\bar{\beta} \pm \bar{\alpha}^*\beta]_{\ell=0} + \pi$ has a small phase space. (iii) This can decay into the $N\bar{N}$ state. The phase space is also small. This resonance corresponds to S(1935). The quantum number is not established but Carter et al[2] gave argument for $J^{PC} = 2^{++}$ which agrees with our prediction.

V. Concluding remarks

For the S- state the highest spin state ($J = 2^+$) has the narrowest width in the S-state diquonium as is explained before. For the P-state and D-state similar result will be deduced because the large angular momentum must go somewhere if it decays into meson states. Since ρ and ω are easier to be produced than π , it loses the phase space easily. They are not the narrowest but narrower than those which have approximately the same masses and lower spins. Thus they are easier to be detected. (Based on the somewhat different reason Dalitz[19] argued that baryon resonances on the leading trajectory stands out in $\sigma(\pi^+P)$ as a consequence of their high spins.)These resonances can decay into $N\bar{N}$ state more easily than into the meson states by one quark pair creation.

Experimentally[20] two such resonances are known. The T(2190), $1^+(3^-)$ resonance corresponds to the highest angular momentum states in the P-wave $\alpha^*\bar{\alpha}^*$ states. The predicted mass 2153 MeV roughly corresponds to the experimantal value 2192 $\pm$ 10 MeV. The U(2350), $0^+(4^+)$ resonance corresponds to that in the D-wave. The predicted mass 2385 MeV should be compared to the experimental value 2350 $\pm$ 25 MeV.

The $[\beta^*\bar{\beta}^*]$, 3^- and 4^+ for the P-state and D-state, respectively are other highest angular momentum states. Since they are M-diquania, they have very small coupling to $N\bar{N}$.

It will be worth while to seek P-wave narrow diquanium states (1314, 1174, 1104 and 1046 MeV in Fig. 3). These resonances can be produced by the decay of the atomic $P\bar{P}$ state. Since they are in P-states the mixing between M- and T-diquanium becomes rather large and they can couple to the $P\bar{P}$ states. Especially the search of the narrow resonance $0^{+}(1^{-})$, 1174 MeV is important to check the diquonium picture because this quantum number does not exist in $Q\bar{Q}$ states.

The S(1935) resonance $[\beta\bar{\beta}]_{\ell=2}$ is the lowest narrow T-diquonium in our model. Because of the large coupling of the T-diquonium to the $N\bar{N}$ system this resonance is easiest to be detected. Our model also predict another T-diquonium, $[\beta\bar{\beta}]_{\ell=1}$, $0^{-}(1^{-})$, 1628 MeV with the width of order 50 - 100 MeV. This resonance decays into $\omega\pi$ or $\rho\pi$.

Acknowledgement

The author would like to thank Professors J. Daboul, H. Joos and F. Gutbrot for discussions and hospitalities at DESY.
He wishes to thank the Alexander von Humbolt Foundation for the financial support.

References

1. P. Pavlopoulos, et al, Phys. Letters 72B, 415 (1978).
2. A. A. Carter, et al, Phys. Letters 67B, 117 (1977).
3. P. Benkheiri, et al, Phys. Letters 68B, 483 (1977).
4. W. Brückner, et al, Phys. Letters 67B, 222 (1977).
5. T. A. Armstrong, Preprint A5E (1978).
6. Chan Hong-Mo, et al, Phys. Letters 76B,634 (1978); 72B, 121, 400 (1978).
7. C. Rosenzweig, Phys. Rev. Letters 36, 697 (1976).
8. R. L. Jaffe, Phys. Rev. D17, 1444(1978).
9. T. Minamikawa, et al, Progr. Theoret. Phys. 37 & 38, 56 (1966).
10. S. Ono, Nucl. Phys. B107, 522 (1976).
11. S. Ono, Phys. Rev. D17, 888 (1978).
12. A. De Rújula, H. Georgi, and S. L. Glashow, Phys. Rev. D12, 147 (1975), J. D. Jackson, Proceeding of the Summer Institute on Particle Physics, SLAC, Aug. 1976.
13. F. Myhrer and A. W. Thomas, Phys. Letters 64B, 59 (1976).
14. B. O. Kerbikov et al, ITEP-61 (1978).
15. J. J. J. Kokkedee "The Quark Model", W. A. Benjamin (1969).
16. G. Morpurgo, in "Theory and Phenomenology in Particle Physics", Part A (1969) P. 84, Academic Press New York and London.
17. J. Uretsky, In High Energy Theoretical Physics (et H; Aly. Univ. of Beirut 1967) P.285.
18. S. Ono, Phys. Rev. D9, 2670 (1974).
19. R. H. Dalitz, An invited paper presented at the Summer Symposium, "New Directions in Hadron Spectroscopy" at Argonne National Laboratory, August, 1975.
20. Particle Data Group, Phys. Letters 75B, 1 (1978).

Figure Captions

Fig.1 Tetrahedron constructed by two quarks and two antiquarks.

Fig.2 The predicted spectrum of the S-wave diquonium states. The superallowed decay modes are shown.
(a) These decays which would be energetically forbidden actually occur because of the widths of ρ and ω .

Fig.3 The predicted spectrum of the P-wave diquonium states.

Fig.4 The predicted spectrum of the D-wave diquonium states. The cascade decay modes of the $\beta\bar{\beta}$ and $\beta^*\bar{\beta}^*$ states are shown.

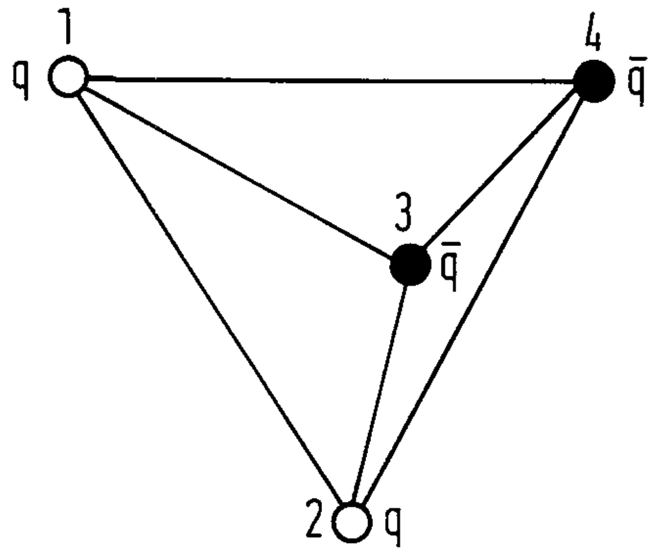


Fig.1

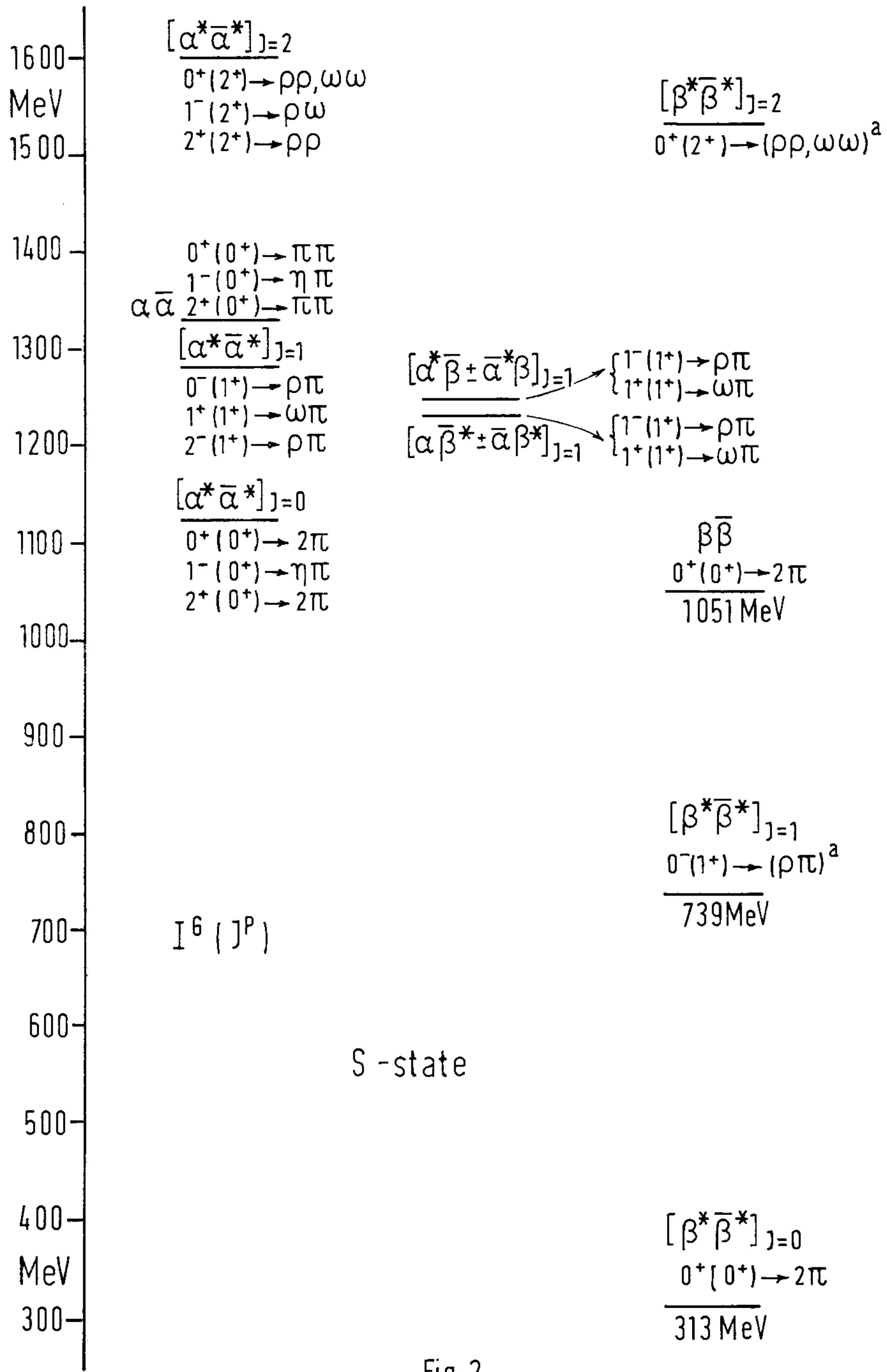
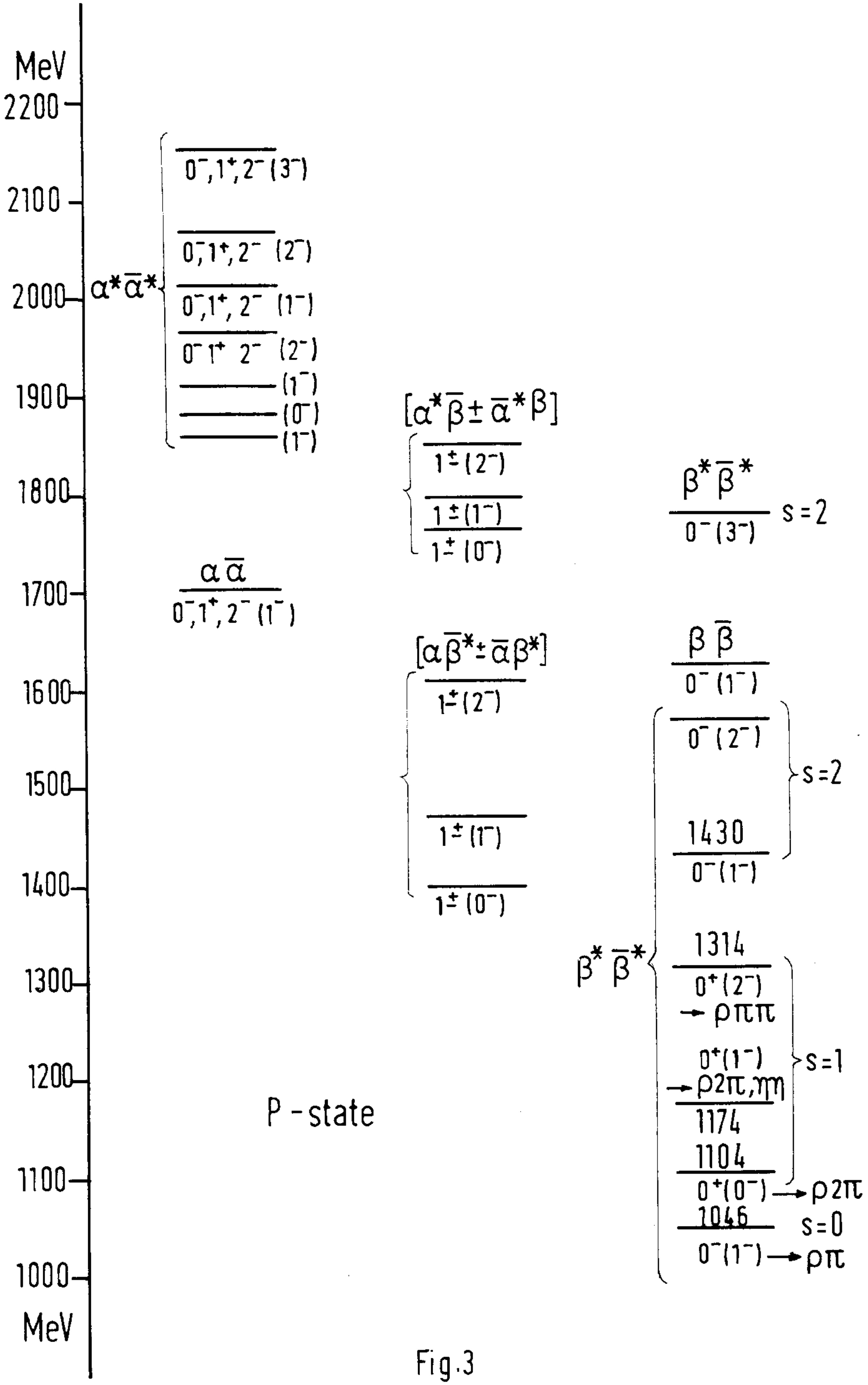

1600 MeV
1500
1400
1300
1200
1100
1000
900
800
700
600
500
400 MeV
300
$[\alpha^*\bar{\alpha}^*]_{J=2}$
$0^+(2^+) \to \rho\rho, \omega\omega$
$1^-(2^+) \to \rho\omega$
$2^+(2^+) \to \rho\rho$
$[\beta^*\bar{\beta}^*]_{J=2}$
$0^+(2^+) \to (\rho\rho, \omega\omega)^a$
$0^+(0^+) \to \pi\pi$
$1^-(0^+) \to \eta\pi$
$\alpha\bar{\alpha}$ $2^+(0^+) \to \pi\pi$
$[\alpha^*\bar{\alpha}^*]_{J=1}$
$0^-(1^+) \to \rho\pi$
$1^+(1^+) \to \omega\pi$
$2^-(1^+) \to \rho\pi$
$[\alpha^*\bar{\beta} \pm \bar{\alpha}^*\beta]_{J=1}$
$1^-(1^+) \to \rho\pi$
$1^+(1^+) \to \omega\pi$
$[\alpha\bar{\beta}^* \pm \bar{\alpha}\beta^*]_{J=1}$
$1^-(1^+) \to \rho\pi$
$1^+(1^+) \to \omega\pi$
$[\alpha^*\bar{\alpha}^*]_{J=0}$
$0^+(0^+) \to 2\pi$
$1^-(0^+) \to \eta\pi$
$2^+(0^+) \to 2\pi$
$\beta\bar{\beta}$
$0^+(0^+) \to 2\pi$
1051 MeV
$[\beta^*\bar{\beta}^*]_{J=1}$
$0^-(1^+) \to (\rho\pi)^a$
739MeV
$I^G(J^P)$
S -state
$[\beta^*\bar{\beta}^*]_{J=0}$
$0^+(0^+) \to 2\pi$
313 MeV

Fig. 2



MeV
2200
2100
2000
1900
1800
1700
1600
1500
1400
1300
1200
1100
1000
MeV
α*ᾱ*
0⁻,1⁺,2⁻ (3⁻)
0⁻,1⁺,2⁻ (2⁻)
0⁻,1⁺, 2⁻ (1⁻)
0⁻ 1⁺ 2⁻ (2⁻)
(1⁻)
(0⁻)
(1⁻)
αᾱ
0⁻,1⁺,2⁻ (1⁻)
[α*β̄ ± ᾱ*β]
1± (2⁻)
1± (1⁻)
1± (0⁻)
[αβ̄* ± ᾱβ*]
1± (2⁻)
1± (1⁻)
1± (0⁻)
β*β̄*
0⁻ (3⁻)
s = 2
β β̄
0⁻ (1⁻)
0⁻ (2⁻)
s = 2
1430
0⁻ (1⁻)
β* β̄*
1314
0⁺ (2⁻)
→ ρππ
0⁺ (1⁻)
s = 1
→ ρ2π, ηη
1174
1104
0⁺ (0⁻) → ρ2π
1046
s = 0
0⁻ (1⁻) → ρπ
P - state

Fig.3

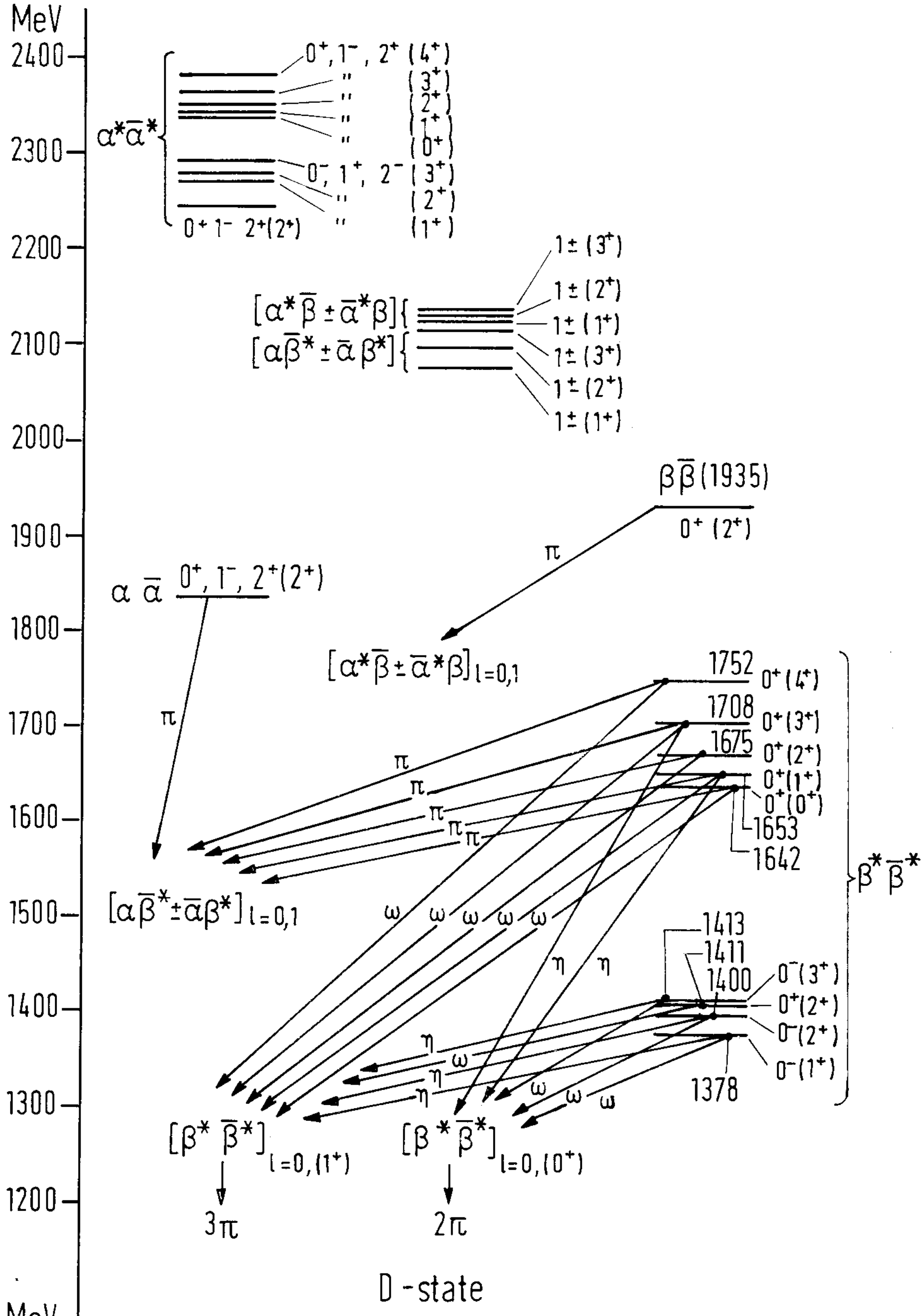

MeV
2400
2300
2200
2100
2000
1900
1800
1700
1600
1500
1400
1300
1200
α*ᾱ*
0⁺, 1⁻, 2⁺ (4⁺)
(3⁺)
(2⁺)
(1⁺)
(0⁺)
0⁻, 1⁺, 2⁻ (3⁺)
(2⁺)
(1⁺)
0⁺ 1⁻ 2⁺(2⁺)
[α*β̄ ± ᾱ*β]
[αβ̄* ± ᾱβ*]
1± (3⁺)
1± (2⁺)
1± (1⁺)
1± (3⁺)
1± (2⁺)
1± (1⁺)
ββ̄ (1935)
0⁺ (2⁺)
α ᾱ 0⁺, 1⁻, 2⁺(2⁺)
π
[α*β̄ ± ᾱ*β] I=0,1
1752 0⁺ (4⁺)
1708 0⁺ (3⁺)
1675 0⁺ (2⁺)
0⁺ (1⁺)
0⁺ (0⁺)
1653
1642
β*β̄*
[αβ̄* ± ᾱβ*] I=0,1
ω
η
1413
1411
1400 0⁻ (3⁺)
0⁺ (2⁺)
0⁻ (2⁺)
0⁻ (1⁺)
1378
[β* β̄*] I=0,(1⁺)
[β* β̄*] I=0,(0⁺)
3π
2π
D-state
MeV

Fig. 4