BARYONIUMS — THE S-MATRIX APPROACH*

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Physics of baryonium is discussed within dual S-matrix framework. Some difficulties within the context of Standard FESR are pointed out. The experimental situation concerning the baryonium spectroscopy and phenomenology of baryonium exchange in two body and inclusive reactions is reviewed. Phenomenological evidence for ω-baryonium mixing is presented. A model for this mixing based on dual unitarisation scheme is discussed.

Baryoniums are a set of mesons which couple primarily to baryon-antibaryon channels

$$\begin{aligned} B &\to b\bar{b} \\ &\not\to MM, \end{aligned}$$

just as the charmoniums couple primarily to pairs of charmed particles

$$\begin{aligned} \psi &\to D\bar{D} \\ &\not\to MM, \end{aligned}$$

or the ϕ, f′ couple primarily to pairs of strange particles

$$\begin{aligned} \phi &\to K\bar{K} \\ &\not\to MM. \end{aligned}$$

These baryoniums are fairly old objects — they were invented by Dualists over a decade back [1]. Only the christening is new. The name was inspired no doubt by their rich cousins — the charmoniums, and following the same inspiration, the ϕ, f′ particles are now called strangoniums.

How are the baryoniums related to charmoniums and strangoniums — why are they cousins? In the S-Matrix framework, they all follow from the same pair of hypotheses, Duality and No Exotics. Let us see how? Here one does not have to invoke any underlying quark structure, except that inherent in the assumption of No Exotics. This assumption

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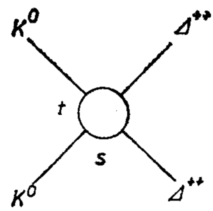
means that there are no mesons outside the singlet and octet representation of SU(3) and no baryons outside the singlet, octet and decuplet. In other words all mesons occur within the quantum number of a $q\bar{q}$ system and all baryons within those of qqq. This seems to be an experimental fact, which has no natural explanation within the *S*-Matrix framework except that it is the minimal non-zero solution to the Duality constraints. The approach in the past has been to take it as an experimental input and build up a phenomenological *S*-Matrix framework. Lately it has been realised that the answer may come from the colour dynamics of quarks. If true this would provide an important link between the fundamental but invisible field theory of quarks and gluons and the phenomenological but visible *S*-Matrix theory overlying it.

Strangonium and charmonium

Anyway, given the no exotics hypothesis duality implies not only exchange degeneracy but also the existence of strangoniums and charmoniums and the so called Zweig rule [2]. We can see this from the following example [3]. Consider



	Q	S
s	2	1
u	2	-1

where both the *s* and *u* channels are exotic. Duality implies

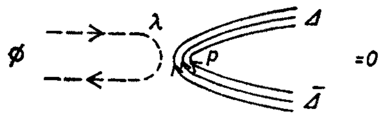
$$\sum_{T\pm V} \gamma_R s^{\alpha_R} = \langle \mathrm{Res} \rangle_{s,u} = 0,$$

i.e. the vector contributions (ω, ρ, ϕ) and the tensor contribution (f, A_2, f′) have to separately vanish and since each must vanish over a wide range of energy, the cancellation has to occur amongst a set of degenerate trajectories. The minimal solution is of course the null solution where the trajectories are not degenerate, in which case all the couplings vanish and there is no scattering to talk of. But the minimal non-zero solution corresponds to the physical situation, where ρ, ω are degenerate and ϕ is not. Hence ρ and ω couplings are required to cancel and ϕ is required to decouple. These 2 constraints determine the 2 parameters — the singlet to octet coupling ratio and the singlet-octet mixing angle in $\omega = V_1 \cos\theta + V_8 \sin\theta; \phi = V_8 \cos\theta - V_1 \sin\theta$. Thus one has a completely determined system apart from the overall normalization, which can not be determined by the duality constraints, since they are linear.

Although no quarks were involved in deriving these results, these can be translated into the quark language by defining

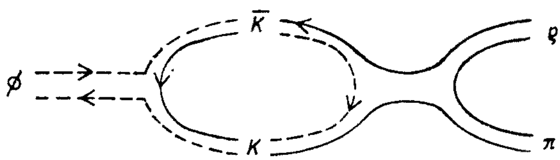
$$V_1 = \frac{p\bar{p} + n\bar{n} + \lambda\bar{\lambda}}{\sqrt{3}}, \quad V_8 = \frac{p\bar{p} + n\bar{n} - 2\lambda\bar{\lambda}}{\sqrt{6}}.$$

The mixing angle results ($\tan\theta = \sqrt{\frac{1}{2}}$) implies $\phi = \lambda\bar{\lambda}$ i.e. a purely strange quark–antiquark state (or strangonium). The second results — the decoupling of ϕ from $\Delta\bar{\Delta}$ — implies the vanishing of the disconnected diagram



which is the Zweig rule. Similarily one gets $\phi \nrightarrow \rho\pi$ and $\psi \nrightarrow \rho\pi$.

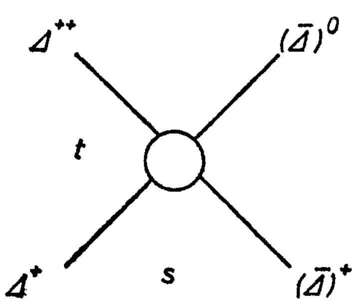
Of course these decays are not completely forbidden. They can go through via the higher orders diagram line



where the 2 halfs are Zweig allowed dual amplitudes and the stitching is done via unitarity. There are, of course, unknown things like the relative phase of the 2 amplitudes and their dependence on the pair of internal legs which is to be summed over, and the result will depend on one's model assumptions for these things. Nonetheless there are some general predictions — e.g. the Zweig disallowed couplings are much smaller than the corresponding allowed couplings (i.e. left half of the diagram); and this relative suppression increases with the mass of the decaying particle.

Baryonium

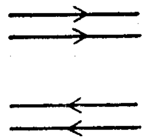
The same duality constraints, when applied to baryon-antibaryon scattering, implies the existence of baryoniums. We can again see this from a simple example [3]. Consider the charge exchange process



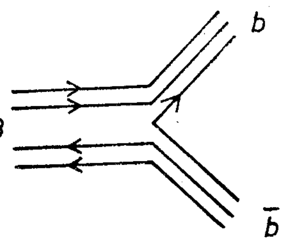
Where again the s-channel is exotic (double charge) and the u channel is exotic (di-baryon). But the vector and tensor exchange are only ρ and A_2 which cannot decouple — apart from contradicting data, this would require via factorization that all the meson-baryon amplitudes vanish. Hence the only way out is to have meson resonances in exotic baryon--antibaryon channels[1]. Now for consistency with the earlier solution, these exotic mesons must decouple from the meson-meson channels. Hence they are called baryoniums in analogy with the strangoniums or charmoniums.

One can again translate the whole thing into a quark language.

1. Since these mesons (or some of them at least) have exotic quantum numbers, they must correspond to 4 quark lines

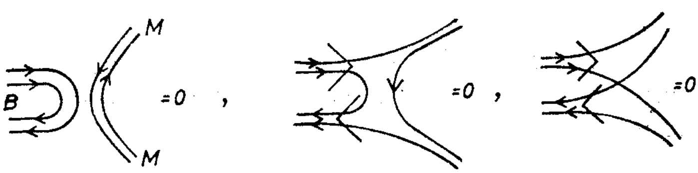


2. Their coupling to the baryon-antibaryon channel goes through by the normal connected graph



like the ϕ coupling to $K\bar{K}$.

3. Their decoupling from the meson-meson channels, i.e.

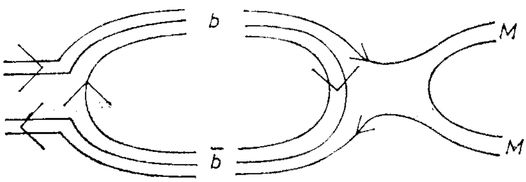


is ensured by a slight generalization of the Zweig rule which disallows both the disconnected graphs and the splitting of the diquark boundary[2]. (In the case of strangonium or charmonium the boundary was a single quark line and hence the question of splitting did not arise.)

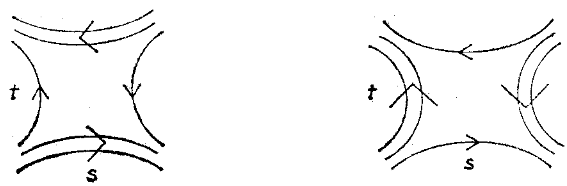
[1] The other alternative of having dibaryon resonances is ruled out by the factorization constraints. The meson-baryon scattering example considered earlier fixes the relative sign of the ρ and A_2 contribution in KΔ scattering, and similar considerations for meson-meson scattering fixes the relative sign in KK. By factorization, then, the relative sign of the ρ and A_2 coupling to $\Delta\bar{\Delta}$ is fixed, which ensures that they cancel in the ΔΔ channel and add up in $\Delta\bar{\Delta}$.

[2] More precisely one requires that each quark line is shared by 2 mesons and each pair of mesons share at least one quark line (Freund, Rosner and Waltz, *Nucl. Phys.* **B13**, 237 (1969)).

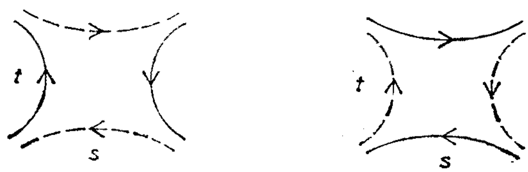
4. The baryonium can decay into ordinary mesonic channels via the higher order graphs



similar to the decay of strangoniums into nonstrange mesons. But again, such couplings are expected to be much smaller than the Zweig allowed couplings, and the suppression to increase with the baryonium mass.
5. In baryon-antibaryon scattering, the normal $(q\bar{q})$ resonances are dual to the baryonium exchange and vice versa,



analogous to the $K\bar{K}$ scattering case, where the ω, ρ resonances are dual to strangonium exchange and vice versa.



This means that knowing the resonance and Regge parameters of ordinary mesons in baryon-antibaryon scattering provides significant constraints on the baryonium parameters. One cannot push their mass arbitrarily high or their couplings arbitrarily low, and as we shall see later, the baryonium resonance data seems to have a serious quantitative discrepency with such predictions.

Finally, the duality constraint applied to the baryonium-baryonium scattering has been shown [3] to imply still higher exotics. These would correspond to $3q3\bar{q}$-states and their couplings to baryon-antibaryon and meson-meson channels would both be forbidden in the lowest order. These may be relevant for describing some narrow meson states, reportedly seen way above the baryon-antibaryon threshold, as we shall discuss later.

Spectroscopy

Of course it is not possible to predict the precise mass of the baryonium states in specific quantum number channels without going to a dynamical model like the various potential models, which have been proposed both in the *S*-Matrix [12a] and the QCD [21] frameworks. Nonetheless one can predict some significant systematics in the baryonium spectroscopy, from fairly general considerations coupled with the duality constraints [4].

Let us denote the diquark system by Q and let us first look at its isospin I. In particular one has

$$I_{\mathrm{Q}} = 1 \quad \text{since} \quad I_{\Delta} = I_{\mathrm{Qq}} = 3/2.$$

The corresponding baryoniums have isospin

$$I_{\mathrm{B}} = I_{\mathrm{Q\bar{Q}}} = 2, 1, 0 \quad (I_{\mathrm{M}} = I_{\mathrm{q\bar{q}}} = 0, 1),$$

i.e. they occur in both exotic and non-exotic quantum numbers. Moreover duality predicts all the 3 isospin states to be degenerate, analogous to the ρ, ω degeneracy prediction for ordinary ($\mathrm{q\bar{q}}$) mesons. Experimentally there seems to be some indications of at least $I = 0$ and 1 degeneracy for some of the baryonium states.

Let us next consider spin S. Again one expects a $S_{\mathrm{Q}} = 1$ system, since $S_{\Delta} = S_{\mathrm{Qq}} = 3/2$ in the $L = 0$ state. Hence

$$S_{\mathrm{B}} = S_{\mathrm{Q\bar{Q}}} = 2, 1, 0 \quad (S_{\mathrm{q\bar{q}}} = 1, 0),$$

i.e. the baryoniums are expected to occur in spin quintet, triplet and singlet states analogous to the spin triplet and singlet states of the ordinary mesons. In analogy with the ordinary mesons, one may expect the lowest mass states for $L = 0$. One then gets the following pattern

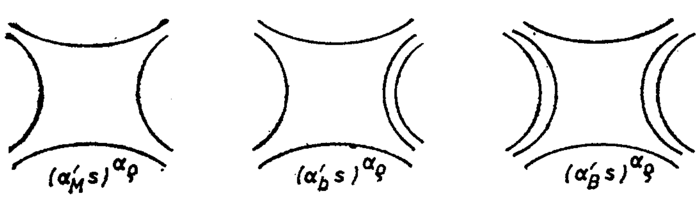
$$\mathrm{Q\bar{Q}} \quad (L = 0) \quad \mathrm{q\bar{q}}$$

$$J = 2, 1, 0 \qquad\qquad J = 1, 0$$

$$P = (-1)^L = \mathrm{NUN} \qquad\qquad P = -(-1)^L = \mathrm{NU}.$$

Of course the different spin-parity states would not be degenerate and one does not know the spacing between them. But again in analogy with the ordinary mesons one may speculate the spacing between the adjacent spin state to be $\frac{1}{2}$ GeV^2 in mass [2].

Finally the slope of the baryonium trajectory is expected to be similar to that of ordinary meson ($\simeq 0.9$ GeV^{-2}) from the following duality argument [5]. Consider ρ exchange in the meson-meson, meson-baryon and baryon-antibaryon scattering.

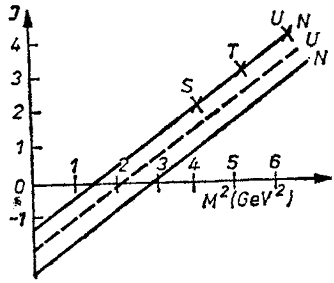


The ρ exchange contribution in a general Veneziano model is shown below each amplitude. Factorization constraint for the ρ residue then implies

$$\alpha'_{\mathrm{B}} = (\alpha'_{\mathrm{b}})^2/\alpha'_{\mathrm{M}},$$

and since the ordinary meson and baryon slopes are both around 0.9 GeV^{-2}, so should be the baryonium slope.

Thus one expects the following Chew–Frautschi plot for the baryonium trajectory.



Of course the trajectory intercept is still arbitrary, but assuming the favoured spin--parity assignment 2^+ for the S(1936) state one expects a leading baryonium intercept of ~ -1, which seems to be consistent with the rough estimate one has from baryonium exchange phenomenology.

Experimental Evidence — over the past 5 years there have been several experimental evidences of baryonium type resonances and several phenomenological evidences for baryonium exchange.

Baryonium resonances

The most reliable ones — i.e. those which have made their way to the Rosenfeld diary — are the S, T and U resonances. They are all seen in the formation channel — $\bar{p}p$ scattering (Fig. 1 and 2). The essential properties of these resonances are summarised in Table I.

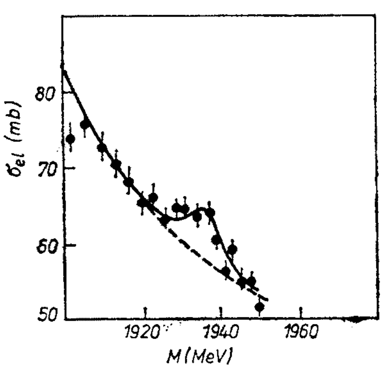


Fig. 1

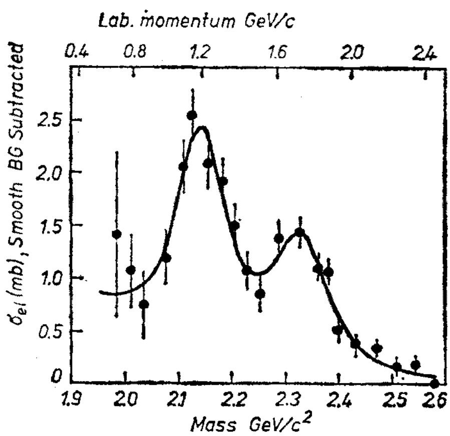


Fig. 2

Fig. 1. The S signal in elastic $\bar{p}p$ cross-section [7]
Fig. 2. The T and U signals in elastic $\bar{p}p$ cross-section [10]

TABLE I

Properties of the S, T, U resonances

Total and Elastic Scatt. [6-10] $M(\Gamma)I$ σ_T, σ_{el} (mb)	$\bar{p}p \to \pi^+\pi^-$ [13] $M(\Gamma)I$ J^{PC}, x	$\bar{p}p \to K^+K^-$ [13]	$\bar{p}p \to \pi^0\pi^0$ [14] $M(\Gamma)I$ J^{PC}	$\pi^-p \to \bar{p}pn$ [15] $M(\Gamma)$ J^P
	2480 ± 30 ($\simeq$ 280)1 5^{--}			
U: 2350 ± 15 ($\simeq$ 160) 1, 0, 3, 2	2310 ± 30 ($\simeq$ 200)0 4^{++}, 0.14	also $I = 0$		$\simeq$ 2300 ($\simeq$ 250) 4^+
T: 2190 ± 10 ($\simeq$ 90)1 5, 2.5	2150 ± 30 ($\simeq$ 200)1 3^{--}, 0.2	also $I = 0$	2150 ± 10 ($\simeq$ 210)0 2^{++}	$\simeq$ 2100 ($\simeq$ 250) 3^-
S: $1936 \pm 1(<10)1(0)$ 10.6 ± 2.4, 7 ± 1.4	$\bar{p}p$ elastic data favours [12, 13] 2^{++}, 0.15			$\simeq$ 1950 ($\simeq$ 250) 1^-

The S(1936) state has all the trappings of a textbook baryonium. It is narrow, being close to the $N\bar{N}$ threshold, and has a large elasticity X (branching ratio into $\bar{p}p$) — all the features are in perfect analogy with the strangonium ϕ. The T and U states are broad presumably for the same reason as f′ is broad — i.e. having a large phase space into the allowed channel $N\bar{N}$. Again like f′ they are characterised by a large elasticity.

The isospin assignments come from comparing the $\bar{p}p$ and $\bar{p}d$ total cross-sections. Carrol et al. [6] find $I = 1$ for S, but do not rule out a $I = 0$ component[3]. Abrams et al. [9] find $I = 1$ for T, and both $I = 0$ and 1 in the U region. Spin parity assignment for S come from fitting the elastic differential cross-section, which favours 2^+ but cannot rule out higher spins [12, 13]. From the phase shift analysis of $p\bar{p} \to \pi^+\pi^-$ Carter et al. [13] have found a 3^- and 4^+ resonances in the T and U region respectively, which can be tentatively identified with these states[4]. They have also found a 5^- state at 2480 MeV. One may note that the M^2 spacing between the adjacent spin states are roughly around 1 GeV^2, which makes their assignment on the same baryonium trajectory quite plausible.

Carter et al. have also found $I = 0$ components in their 3^{--} and 5^{--} states by extending their analysis to $\bar{p}p \to K^+K^-$, and from $\bar{p}p \to \pi^0\pi^0$ Dulude et al. [14, 17], have found a 2^{++} ($I = 0$) state at 2150, which may be a daughter (or a spin singlet state). Finally 3 broad

[3] A recent deuteron bubble chamber experiment [11], does not seem to see an S signal in the annihilation cross-section ($\bar{p}n \to \pi$'s), which may cast some doubt on the $I = 1$ assignment.

[4] With these spin assignments for S, T, U the elasticity X turn out to be smaller than those suggested by the ratio σ_{el}/σ_T. This can be accounted for by interference between the resonance and a BG contribution [12] in the same channel. Similarily the non-observance of these resonance signals in the charge exchange cross-section can be accounted for by interference between a resonance and a BG contribution [12] (or between 2 resonance contributions) in opposite isospin channels.

$\bar{p}p$ resonances have been seen by the CERN Omega spectrometer [15, 17] in $\pi^- p \rightarrow \bar{p}pn$. They are at 2300, 2100 and 1950 with spin parity 4^+, 3^- and 1^-.

A number of narrow resonance candidates have been reported from the production experiments.

1. In addition to the S, 2 narrow $\bar{p}p$ resonances have been observed at 2020 and 2200 MeV by the CERN Omega spectrometer [16] in the backward production experiment (Fig. 3)

$$\pi^- p \rightarrow \frac{p_f \pi^-}{\Delta, N^*} \frac{p\bar{p}}{2020{,}2200}.$$

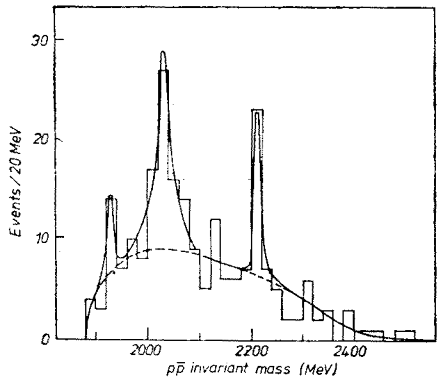


Fig. 3. Narrow peaks in the backward ($\bar{p}p$) mass distribution in $\pi^- p \rightarrow \frac{\pi^- p_f}{\Delta} p\bar{p}$ at 9 and 12 GeV/c [16]

Some independent support for these 2 states have been presented at the Tokyo Conference [17]. The Pittsburg–Massachusetts-Collaboration see a signal at 2200 in the annihilation reaction

$$\bar{p}p \rightarrow \pi_f^- \frac{\pi^+ K^+ K^-}{2200}$$

and a signal at 2020 has been seen in the virtual photoproduction reaction

$$\gamma_V p \rightarrow \frac{\bar{p}p}{2020} p$$

at Cornell. One should note that, since these 2 states have substantial phase space into ahe $\bar{p}p$ channel there is no natural justification for their narrow width ($\Gamma \sim 20$ MeV) in · standard baryonium framework.

2t A narrow resonance signal (6σ) was seen at 2950 MeV by the CERN Omega spectrometer [18], in the forward production reaction

$$\pi^- p \rightarrow \frac{\pi^- p_f \bar{p}}{2950} X.$$

However a repeat of this experiment with higher statistics by the same group, does not seem to show this signal [17].

3. Recently a narrow peak at 2460 MeV have been reported from the CERN Omega spectrometer [18a] in the forward production process

$$K^+p \rightarrow \frac{\bar{\Lambda}p\pi^+}{2460} n$$

$$\hookrightarrow \bar{\Lambda}\Delta^{++} \quad \text{or} \quad \bar{\Sigma}^+p.$$

This is the first resonance candidate with exotic quantum numbers, but there is again no natural justification for the narrow width or for the forward production, within the standard baryonium framework.

4. A narrow peak at 2600 was reported from BEBC [19] in $K^0\pi^\pm\ \pi^+\pi^-$ in the annihilation reaction

$$\bar{p}p \rightarrow K^0\pi^+\pi^+\pi^+\pi^-\pi^-\pi^- + \text{neutrals}.$$

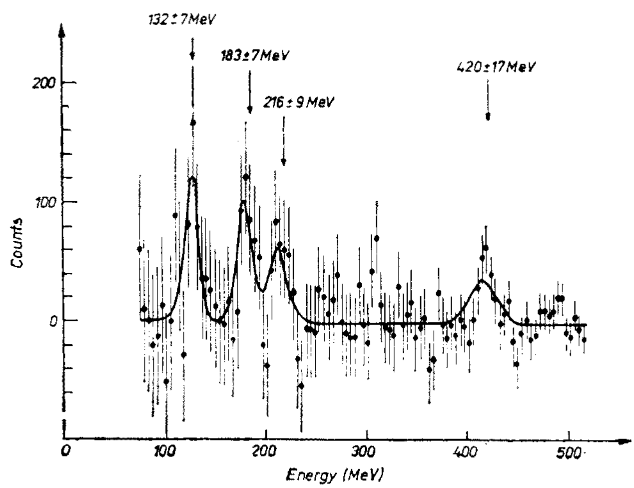
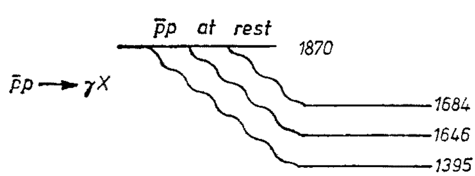


Fig. 4. Narrow peaks in E_γ in the radiative decay of $\bar{p}p$ at rest [20]

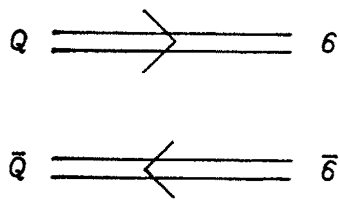
But 2 other bubble chamber searches for this channel have drawn a blank [19a].

5. Finally 3 narrow peaks below the $\bar{p}p$ threshold have been reported [20] from the radiative decay of $\bar{p}p$ at rest (Fig. 4)

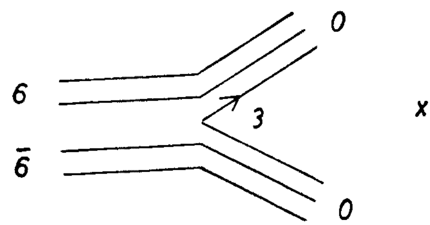


The first peak occuring at $E_\gamma \simeq 132$ MeV is associated with the radiative capture of π^- coming from annihilation. The 3 other peaks at $E_\gamma = 183$, 216 and 420 MeV correspond to peaks at 1684, 1686 and 1395 MeV in the recoiling system.

The resonances below the $\bar{p}p$ threshold are, of course, expected to be narrow from duality arguments. But the narrow resonance candidates, occuring substantially above the $\bar{p}p$ threshold, cannot be interpreted as standard dual baryonium. One interpretation is in terms of a hidden colour state [21]



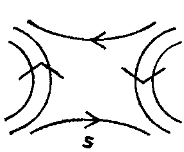
where the diquark supposedly belongs to a colour sextet ($3 \times 3 = 6 + \bar{3}$). As a result their couplings to the baryon-antibaryon channels are inhibited since the latter are colour singlet states.



Alternatively they could be interpreted as the higher exotic ($3q\ 3\bar{q}$) states implied by duality, as discussed earlier. In the S-Matrix framework, at least, the latter interpretation seems more natural. One should bear in mind, of course, that each of these candidates is essentially based on a single experiment.

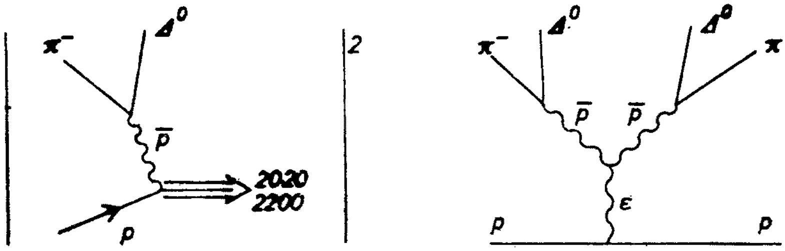
FESR constraint

Next let us look at the duality constraint for the couplings of the baryonium resonances to the baryon-antibaryon channels, i.e.



$$\Rightarrow \langle g^2_{Bb\bar{b}} \rangle \sim \langle \gamma^2_{\omega b \bar{b}}\, s^{\alpha_\omega} \rangle$$

Pennington [22] has drawn attention to a violent contradiction between the two sides, for backward production of the 2020 and 2200 states, i.e. between



$$\left.\frac{d\sigma}{du}\right|_{\pi^- p \to \Delta^0 B} \quad \text{and} \quad \gamma^2_{N\pi\Delta}\gamma_{\omega NN}\Gamma_{\omega NN}s^{2\alpha_N - 2}(M_B^2)^{\alpha_\omega - 2\alpha_N}$$

since the rough magnitudes of the Regge couplings (γ's) and the triple Regge coupling Γ on RHS are known one can estimate the RHS. This turns out to be roughly a factor of 100 larger than the production cross-section for the 2020 and 2200 states (LHS) measured at the Omega spectrometer [16].

However these 2 narrow peaks, even if they are confirmed, may not have anything to do with a normal dual baryonium resonance, as remarked earlier. A more appropriate test of this duality constraint should therefore be the one involving the S, T, U resonances. Since the only reliable signals for these resonances are in the formation channel, we have to look at the old fashioned FESR for the elastic $\bar{p}p$ amplitude. This gives

$$\int_{s_1}^{s_2} \text{Im}\, A_{\text{Bar.Res.}} = \sum_{i=S,T,U} \frac{s_i}{2q_i}(2J_i+1)x_i\Gamma_i = 2\beta_{\omega pp}\frac{s_2^{\alpha_\omega+1} - s_1^{\alpha_\omega+1}}{(\alpha_\omega+1)},$$

where A is the kinematic singularity free amplitude,

$$\text{Im}\, A = \frac{q\sqrt{s}}{2\pi^2}\sigma_T.$$

Note that the individual resonance contribution to σ_T, quoted in Table I, is simply

$$\left.\frac{\pi(2J_i+1)x_i}{q^2\left[\left(\frac{E-E_i}{\Gamma/2}\right)^2+1\right]}\right|_{\max} = \frac{\pi}{q_i^2}(2J_i+1)x_i.$$

The individual resonance contribution to the LHS, from Abrams et al. [9], are compared in Table II against the ω Regge contribution to the RHS[5].

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[5] Change of variable from s to $\nu\left(=\frac{s}{m}-2m\right)$ makes ~5% difference.

TABLE II

Baryonium Resonance and ω Regge Contribution to the FESR $\int_{s_2}^{s_1} ds\, \mathrm{Im}A_{\bar{p}p}$ ($s_1 - s_2 = 3.5 - 6.7\ \mathrm{GeV}^2$; the error shown on the Reggeon Contribution comes from changing these limits by 0.5 GeV^2)

Res.	Cont.	(GeV^2)			Regge Cont. (GeV^2)
S	T	U(I = 1)	U(I = 0)	Total	
0.05 ± 0.04	0.52 ± 0.17	0.7 ± 0.24	0.8 ± 0.27	2.07 ± 0.4 (1.67 ± 0.33)	28.3 ± 5.0

TABLE III

Strangonium Resonance and ω, ρ Regge Contribution to the FESR $\int_{s_1}^{s_2} ds\, \mathrm{Im}\, A_{\bar{K}K}$ ($s_1 - s_2 = 1 - 2.5\ \mathrm{GeV}^2$; the error shown on the Reggeon Contribution comes from changing these limits by 0.2 GeV^2)

Res.	Cont.	(GeV^2)	Regge Cont. (GeV^2)
ϕ	f′	Total	
0.102 ± 0.006	1.3 ± 0.2	1.4 ± 0.2	1.73 ± 0.44

One first notes that the S contribution is negligible compared to T, U because of its narrow width. This is analogous to the ϕ contribution to elastic $\bar{K}K$ being negligible compared to the f′ (Table III). For the same reason, the other narrow states, even if they are confirmed, would make negligible contribution to the FESR.

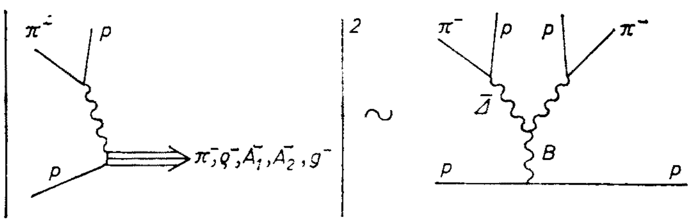
We see a yawning gap of an order of magnitude between the S, T, U. resonance contributions and the ω Regge exchange. This may be contrasted with the situation in $K\bar{K}$ scattering (Table III), where the strangonium (ϕ, f′) resonance contributions seem to be roughly equal to the ω-Regge exchange. There are two possible explanations of this discrepancy: I. The first involves pulling down the ω-Regge contribution — i.e. to assume a large renormalization of the planar ωNN Regge coupling (by a factor of 2–3, say) from unitarity effects. The FESR supposedly builds up only the planar Regge coupling, since the renormalization effects are negligible in the resonance region. Such a large renormalization for the $\omega N\bar{N}$ Regge coupling has been suggested before in the context of a specific ω-baryonium mixing model [23]. However, the only phenomenological estimate of the planar $\omega N\bar{N}$ coupling available is through the $\bar{K}N$ FESR, where two earlier analyses found the FESR estimate smaller than the physical ω coupling by a factor 1.4–1.8 [24] and 2–2.4 [25], depending on the choice of resonance parameters. In view of this large uncertainty, associated with the large resonance region below the $\bar{K}N$ threshold, it gives no significant constraint either in favour or against a large renormalization of the $\omega N\bar{N}$ coupling. We shall see later however, that in the simplest (one parameter) mixing prescription a large renormalization of $\omega N\bar{N}$ coupling ($\sim$2–3) would imply an even larger renormalization of the ω $\Delta\bar{\Delta}$ coupling ($\sim$10–20) as measured from the inclusive cross-section difference. This is certainly ruled out by the inclusive data, but of course the one parameter mixing prescription may be totally wrong.

II. The other alternative is to jack up the resonance contribution by an order of magnitude. This would mean that S, T, U peaks are only the tip of the ice-berg — there has to be a very large number of other baryonium resonances hiding underneath. This may seem inplausible in terms of our experience with the ordinary ($q\bar{q}$) meson spectrum; but as we have seen earlier, one expects a much higher density of baryonium states compared to the ordinary ($q\bar{q}$) mesons. These resonances have to be broad and highly elastic to make a significant contribution to the FESR. It would be extremely interesting therefore, to search for such states in the $\bar{p}p$ elastic differential cross-section over the S, T, U region, together with the $\bar{p}p \to \pi\pi$ and $K\bar{K}$ phase shift analyses.

I will simply state my personal prejudice, which is in favour of the second alternative.

Baryonium exchange

1. The first phenomenological evidence for baryonium exchange came from studying the backward production cross-section of ordinary ($q\bar{q}$) mesons which should be dual to baryonium exchange.



$$\left.\frac{d\sigma}{du}\right|_{\pi^- p \to pR} \sim G s^{2\alpha_\Delta - 2} (M_R^2)^{\alpha_B - 2\alpha_\Delta}.$$

This is complimentary to the backward production case considered above where the production of baryonium resonances was dual to ordinary meson exchange. Anyway, plotting the backward production cross-section of the ordinary meson resonances (π, ρ, A_1, A_2, g) against the resonance mass M_R, one found that the dual trajectory, indeed, has a low intercept $\alpha_B \lesssim -0.5$. This was in striking contrast to the forward production cross-section of the same set of resonances, which had shown that the dual trajectory has the normal intercept $\simeq +0.5$.

2. The second evidence came from comparing a set of exotic exchange reactions in baryon-baryon channels with the corresponding ones in meson baryon channels [27]. These are

$$\sigma(pn \to \Delta^-\Delta^{++}) \overset{3.7}{=} 550 \pm 200, \qquad \sigma(\pi^- p \to \pi^+\Delta^-) \overset{3.2}{=} 50 \pm 12,$$

$$\sigma(\bar{p}p \to \bar{Y}^{*+}Y^{*-}) \overset{3.7}{=} 8 \pm 3, \qquad \sigma(\pi^- p \to K^+Y^{*-}) \overset{4}{\simeq} 1,$$

$$\sigma(\bar{p}p \to \bar{\Sigma}^+\Sigma) \overset{3.6}{=} 6 \pm 1, \qquad \sigma(K^- p \to \pi^+\Sigma^-) \overset{3.8}{=} 0.6 \pm 0.1,$$

where all the cross-sections are in μb and the quantities in the middle refer to the incident momenta in GeV/c. Apart from the baryonium, of course, there would be other exchanges with exotic quantum numbers (e.g. Regge-Regge cuts) — but the latter ones have no particular reason to favour the baryon-baryon channel over the meson-baryon channel. The fact that the baryon-baryon cross-sections seem to be systematically higher than the corresponding meson-baryon ones by an order of magnitude, therefore, suggests the baryonium exchange to be the dominant mechanism at least for the baryon-baryon cross-sections.

One may also note that the strangeness exchange processes are significantly suppressed, but there is no significant difference between octet and decuplet production. These systematics are purely empirical but they should be of help in comparing different baryonium exchange processes.

There is unfortunately very limited data on these exotic exchange processes, which is restricted moreover to a very limited range of low energy. Consequently the estimate of the baryonium trajectory is quite sensitive to the choice of variable. Comparing the $pn \to \Delta^-\Delta^{++}$ cross-sections available at the two incident momenta 3.7 and 7 GeV/c, gives a baryonium intercept anywhere in the range 0 to -1 depending on the choice of variable. More precisely one gets [28]

$$\alpha_B(0) = \begin{array}{lll} -0.1 \pm 0.1 & \text{for} & P_L \\ -0.3 \pm 0.4 & & \nu \\ -0.5 \pm 0.5 & & s. \end{array}$$

It seems to me that the clearest estimate of the baryonium trajectory can come from the measurement of these exotic exchange cross-sections over a reasonably wide energy range (up to 20 GeV/c, say).

3. There also seems to be an evidence for baryonium exchange in the good old proton-antiproton cross-section difference Δ_{pp}. In fact as soon as the first set of total cross-section data came out from the Fermilab, it was realised that the Δ_{pp} does not quite behave the same way as Δ_{Kp} or $\Delta_{\pi p}$. Where as the latter two follow a single power law (characteristic of ω, ρ exchange) from 250 down to 2 GeV/c, the Δ_{pp} shows a distinct upward curvature at the low energy end [29]. But since at that time the baryonium episode had been forgotten, this was looked upon as somewhat of an embarrassment and accommodated through an empirical change of variable. Following the revival of interest in the baryonium, we have reanalysed these cross-section differences [28]. At least for the 3 conventional choices of variables P_L, ν and s one finds that the $\Delta_{\pi p}$ and Δ_{Kp} show a single power law behaviour but Δ_{pp} does not. This is illustrated in Fig. 5 for the P_L variable and in Fig. 6 for the s variable (the situation for the ν variable is inbetween these two). Instead of going into details of the fit let me just mention the essential points. Although the departure from linearity looks small, the low energy data points have so tiny errors that a simple ω, ρ fit gives a $\chi^2/\nu \sim 100$ (even for the P_L variable). Adding a single low lying trajectory brings down χ^2/ν to $\simeq 1$ in each case; although this trajectory parameter is sensitive to the choice of

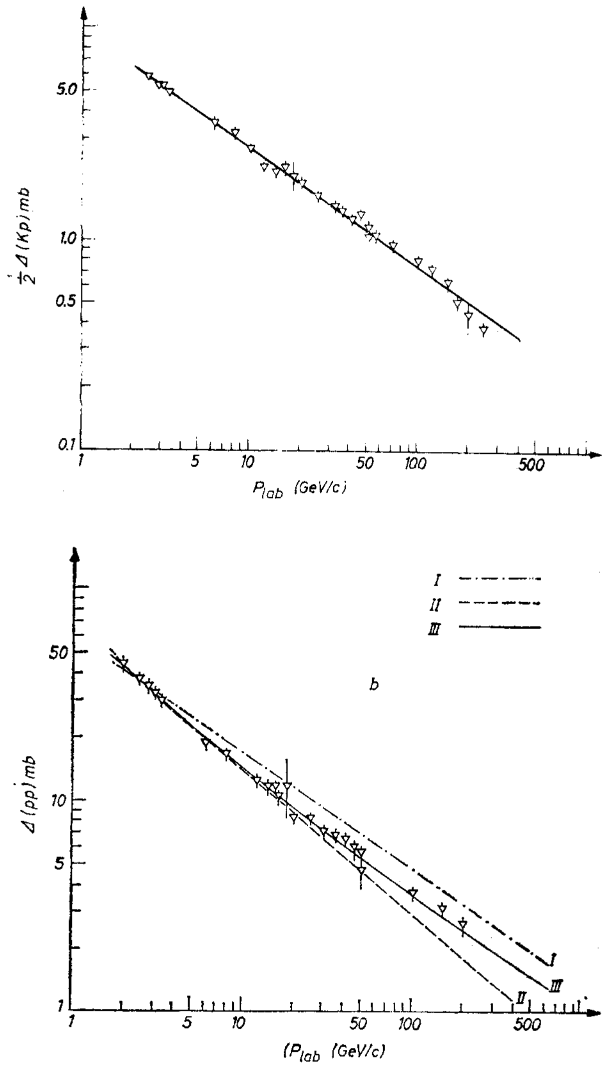


Fig. 5. The total cross-section difference a) Δ(Kp), b) Δ(pp) versus $p_{\rm lab}$, showing a departure from linearity for the latter. The fits are from Ref. [28], which also gives references to the data points

parameter. One gets

$$\alpha_B(0) = \begin{array}{lll} -0.24 & \text{for} & P_L \\ -0.64 & & \nu \\ -1.1 & & s, \end{array}$$

which agree with the corresponding estimates from the exotic exchange reaction $pn \to \Delta^-\Delta^{++}$. Moreover the baryonium residues estimated from the 2 reactions have roughly similar magnitude as well.

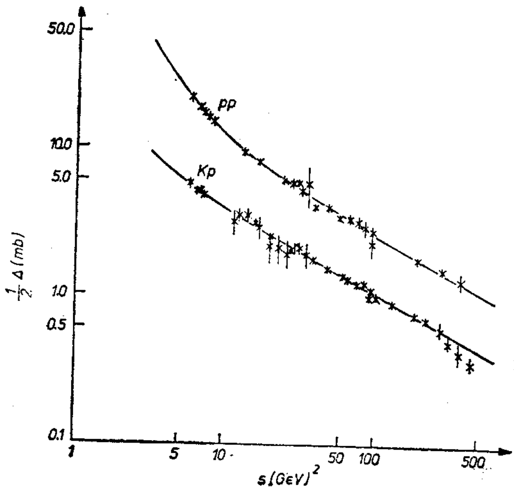
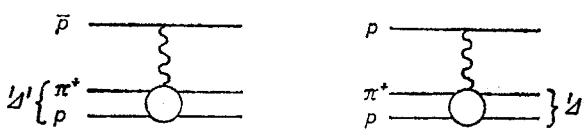


Fig. 6. Same quantities as Fig. 5, plotted against s. The fits are from Ref. [30]

4. By far the most convincing signal of baryonium exchange is seen in the difference of inclusive cross-section [30].

$$\Delta_{pp}(\pi^-) = \sigma(p \xrightarrow{\bar{p}} \pi^-) - \sigma(p \xrightarrow{p} \pi^-) = \sum_{i=\rho,\omega,B} \gamma^i_p \gamma^i_{'\Delta'} s^{\alpha_i - 1}$$



let us recall that

$$\Delta_{Kp}(\pi^-) = \sum_{i=\rho,\omega} \gamma^i_K \gamma^i_{'\omega'} s^{\alpha_i - 1},$$

$$\Delta_{\pi p}(\pi^-) = \gamma^\rho_\pi \gamma^\rho_{'\omega'} s^{\alpha_\rho - 1}.$$

Moreover the ρ and ω coupling to 'Δ' are degenerate

$$\gamma^\rho_{'\Delta'} = \gamma^\omega_{'\omega'},$$

where as the 2-body couplings are related by

$$\gamma^\rho_K = \gamma^\omega_K = \tfrac{1}{2}\gamma^\rho_\pi = \gamma^\rho_p = \tfrac{1}{3}\gamma^\omega_p.$$

As a result one expects

$$\Delta_{Kp}(\pi^-) \simeq \Delta_{\pi p}(\pi^-),$$

and the $\Delta_{pp}(\pi^-)$, minus the baryonium contribution,

$$\Delta^M_{pp}(\pi^-) \simeq 2\Delta_{\pi p}(\pi^-).$$

It was realised quite early in the inclusive game that where as $\Delta_{\pi p}(\pi^-)$ and $\Delta_{Kp}(\pi^-)$ both fall like $s^{-\frac{1}{2}}$ and are roughly equal, the $\Delta_{pp}(\pi^-)$ neither falls like $s^{-\frac{1}{2}}$ nor follows the above systematics. But a meaningful quantitative analysis has become possible only after the BNL–Pennsylvania data [31], which came out last year. It is an incredible piece of data which covers both the BNL and Fermilab energy ranges ($P_L = 8-150$ GeV/c) with the

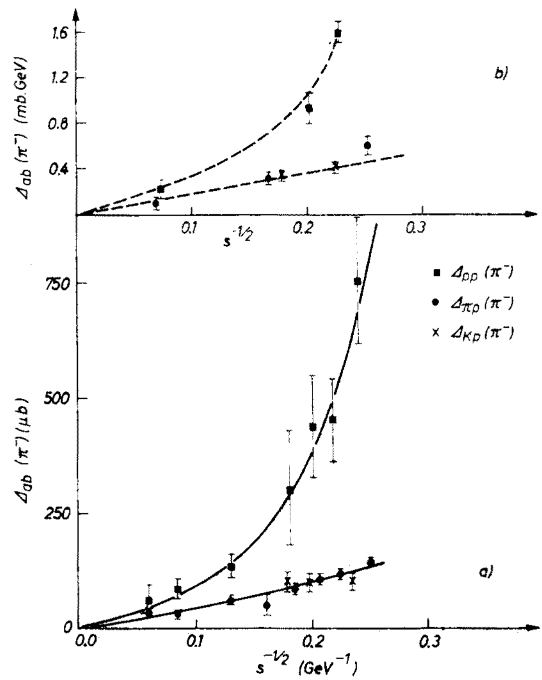


Fig. 7. a) The inclusive cross-section difference $\Delta_{pp,\,\pi p,\,Kp}(\pi^-)$ plotted against $s^{-\frac{1}{2}}$. b) The same plot for a different kinematic cut. For details see the text and Ref. [30]

same spectrometer and the same kinematic cut. In fact, this is the only piece of data which covers a large enough energy range with a low point-to-point systematic error to do a quantitative Mueller Regge Analysis. Fig. 7 shows their result for $\Delta_{\pi p}(\pi^-)$ and $\Delta_{pp}(\pi^-)$ against $s^{-\frac{1}{2}}$. Unfortunately their Kp data has rather poor statistics — the $\Delta_{Kp}(\pi^-)$ shown are interpolated from Bubble Chamber data for their kinematic cut.

One first notices that the $\Delta_{\pi p}$ is approximately linear over the entire energy range — one gets $\alpha_\varrho(0) \simeq 0.4$. The $\Delta_{Kp}(\pi^-)$ is equal to $\Delta_{\pi p}(\pi^-)$ within the errors. Over the Fermilab range ($P_L > 30$ GeV/c) the $\Delta_{pp}(\pi^-)$ is roughly twice the $\Delta_{\pi p}(\pi^-)$. Over the low energy range however there is a clear excess in $\Delta_{pp}(\pi^-)$. Fitting this excess in terms of a single baryonium exchange gives

$$\alpha_B(0) \sim -1.3.$$

In fact, the information content in this single figure is so large that one can get from this a pretty tight bound on the ω–baryonium mixing, which we shall discuss next.

Phenomenological estimate of ω–baryonium mixing at $t = 0$

If we have 2 states with the same quantum numbers like ω and B, then unitarity effects will induce some amount of mixing between the 2. In a specific unitarization model, of course, the magnitude of mixing can be predicted, and two such models have been studied in detail by Chew and Rosenzweig [32] (CR) and Chan and Tsou [23] (CT).

But let us see first if we can estimate the mixing parameter from the data. For this we have to assume that the ω–B mixing is described by a single parameter

$$\omega = \cos\theta\,\omega^P + \sin\theta\,B^P, \qquad B = \cos\theta\,B^P - \sin\theta\,\omega^P,$$

where P refers to the pure quark (planar) state. This prescription is built into the CR model through the assumption of a J independent kernel. On the other hand, in the CT model

$$\omega = \varrho_\omega \cos\theta_\omega\,\omega^P + \varrho_\omega \sin\theta_\omega\,B^P, \qquad B = \varrho_B \cos\theta_B\,B^P - \varrho_B \sin\theta_B\,\omega^P,$$

i.e. a priori there is no relation between the planar and the physical couplings, but even in this model the output values of θ_ω and θ_B are roughly equal and $\rho_\omega \sim \rho_B \sim 1$. Therefore the one parameter prescription given above may be reasonable. In any case since the data can determine essentially a single quantity let us stick to this one parameter prescription.

We first note that the physical baryonium coupling to the $K\bar{K}$ channel is proportional to $\sin\theta$. Since both the inclusive and the total cross-section data show a clear distinction between the $\bar{p}p$ and Kp with respect to the baryonium signal, one expects the mixing angle to be small[6]. Thus we can work to the first order in θ. Then the extraction of the mixing angle is quite transparent.

I. Since $\gamma^{\varrho}_{`\Delta'} = \gamma^{\varrho P}_{`\Delta'} = \gamma^{\omega P}_{`\Delta'}$, the deviation of $\Delta_{pp}(\pi^-)$ from $2\Delta_{\pi p}(\pi^-)$ at the high energy end essentially measures the deviation of $\gamma^{\omega}_{`\Delta'}$ from its planar value. Taking account of the errors in these 2 measurements and using the actual 2-body Regge couplings instead of the approximate relation given earlier, we get

$$\frac{\gamma^{\omega}_{`\Delta'} - \gamma^{\rho}_{`\Delta'}}{\gamma^{\rho}_{`\Delta'}} = \sin\theta\,\frac{\gamma^{B}_{`\Delta'}}{\gamma^{\omega}_{`\Delta'}} = 0.15 \pm 0.39. \tag{1}$$

II. The relative size of the baryonium contribution in Kp and pp gives

$$\frac{\Delta^{B}_{Kp}(\pi^-)}{\Delta^{B}_{pp}(\pi^-)} = -\sin\theta\,\frac{\gamma^{\omega}_{K}}{\gamma^{B}_{p}} = -0.05 \pm 0.15. \tag{2}$$

Since the $\Delta_{Kp}(\pi^-)$ data was available over a limited energy range, the limit on the baryonium contribution was obtained from the difference of $\Delta_{Kp}(\pi^-)$ and $\Delta_{\pi p}(\pi^-)$. There is also a contribution to this difference from $\gamma^{\omega}_{`\Delta'} - \gamma^{\omega P}_{`\Delta'}$, which has been taken into account using Eq. (1). Alternatively we can get this quantity from the total cross-section data

$$\frac{\Delta^{B}_{Kp}}{\Delta^{B}_{pp}} = -\sin\theta\,\frac{\gamma^{\omega}_{K}}{\gamma^{B}_{p}} = 0.05 \pm 0.01, \tag{2'}$$

[6] If the signal were comparable for the 2 cases, it would, in fact, be impossible to distinguish the baryonium from other low-lying exchanges like Regge-Regge cuts.

where the baryonium contribution to Δ_{Kp} is estimated from the energy dependence (all we need of course is only an upper bound).

III. Finally

$$\frac{\gamma^{B}_{\text{`}\Delta\text{'}}}{\gamma^{\omega}_{\text{`}\Delta\text{'}}} \sim 10 \frac{\gamma^{B}_{p}}{\gamma^{\omega}_{p}}, \tag{3}$$

which simply comes from the fact that the baryonium exchange signal relative to ω exchange is an order of magnitude stronger in the inclusive cross-section difference than in the total cross-section difference.

Moreover

$$\frac{\gamma^{\omega}_{p}}{\gamma^{\omega}_{K}} \sim 3. \tag{4}$$

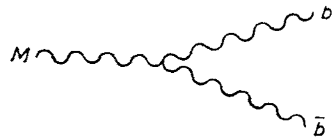
From equations (1), (2), (3) and (4) one gets a 1 standard deviation bound $\theta \lesssim 15°$ which essentially comes from the inclusive data alone. If one uses the baryonium signal in the total cross-section data as well — i.e. Eq. (2′) in place of (2) — then one gets even a stronger bound, $\theta \lesssim 5°$.

I should add that although the $\omega-\phi$ mixing was ignored above its inclusion up to the first order in mixing angle, would not affect our result, as we shall see later.

Models for ω–baryonium mixing

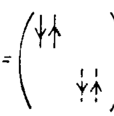
Let us compare this phenomenological bound, with some of the model estimates of ω–B mixing. Chan and Tsou have estimated ω–B mixing in the SU(3) symmetry limit (no ω–ϕ mixing). Their estimate of the mixing angle is $\sim 45°$, i.e. several times higher than even the more liberal bound $\lesssim 15°$ following from inclusive data alone. Although the bound was obtained in the 1-parameter mixing prescription (i.e. $\varrho_{\omega B} \simeq 1$), substituting the CT model values for $\varrho_{\omega B}$ pushes up the bound by only 50%, which is still significantly lower than the model value. Recently this model has been extended by Hansson [33] to incorporate ϕ mixing as well — i.e. to describe ω–ϕ–B mixing. But the resulting ω–B mixing angle is still close to the CT estimate, i.e. several times higher than the phenomenological bound. Of course, in obtaining the phenomenological bound, one has neglected things like mixing in the Isospin 1 channel (ρ) and second order effects in ω–ϕ mixing. Each of these would contribute to the difference $\gamma^{\omega}_{\text{`}\Delta\text{'}} - \gamma^{\rho}_{\text{`}\Delta\text{'}}$ along with the ω–baryonium mixing effects considered above. In case the former effects are not negligible one may possibly arrange to cancel them with a large ω–baryonium mixing contribution, so as to reproduce a small breaking of ω–ρ degeneracy, and hence small a effective angle θ. Barring such accidental cancellation however, I do not see anyway of reconciling a large ω–baryonium mixing model with the above data. In particular it may be emphasised that there are only 2 places where the ω–baryonium mixing effect shown up experimentally (1) ω, ρ degeneracy breaking $(\gamma^{\omega}_{\text{`}\Delta\text{'}} - \gamma^{\rho}_{\text{`}\Delta\text{'}})/\gamma^{\rho}_{\text{`}\Delta\text{'}}$, and (2) the relative size of the baryonium signal in Kp and $\bar{p}p$ $\Delta^{B}_{kp}(\pi^-)/\Delta^{B}_{pp}(\pi^-)$, and both the effects seem to be rather small.

I shall not be able to describe the CT model. Let me just mention that it is a highly ambitious scheme, with no free parameters. Not only are the physical ω and B trajectories and the mixing parameters all determined in terms of the planar trajectories, but the planar trajectories are also determined by the planar bootstrap constraints. In fact apart from a single parameter, the meson–baryon–antibaryon triple Regge coupling

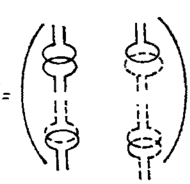


which is taken from phenomenological fits, all other parameters are determined by the dynamical constraints. Of course the price you have to pay for such a zero parameter model is that one has to make a number of detailed dynamical assumptions and approximations on the way.

On the other hand the CR model is a less ambitious programme. But it is simpler and more phenomenological. It was developed by Chew and Rosenzweig for ω–ϕ mixing. Recently it has been extended by Gavai [34] to the ω–ϕ–baryonium mixing case. Let me summarize the essential feature. In the ω–ϕ mixing case the planar propagator is

$P^{\pm} =$  $= \begin{pmatrix} \frac{1}{J-\alpha_0} & 0 \\ 0 & \frac{1}{J-\alpha_3} \end{pmatrix}$

where α_0 and α_3 are the planar ω, f and ϕ, f′ intercepts and the superscript refers to even and odd charge conjugation. The mixing is done by a kernel C, which is assumed to be SU(3) symmetric and J independent.

$C^{\pm} =$  $= \pm K \begin{pmatrix} 2 & \sqrt{2} \\ \sqrt{2} & 1 \end{pmatrix}$

i.e. it is specified at a given t, by a single parameter K. In the J plane the unitarity integral equation can be written as

$$P' = P + PCP + PCPC + \ldots = (P^{-1} - C)^{-1}$$

and the renormalised pole position are obtained by solving the equation.

$$\mathrm{Det}\,(P^{-1} - C) = \mathrm{Det}\,(JI - C') = 0,$$

where

$$C' = C - P^{-1} + JI = (C) + \begin{pmatrix} \alpha_0 & 0 \\ 0 & \alpha_3 \end{pmatrix}.$$

Now the J independence of C implies C' is also J independent, which makes life simple. It means firstly that the renormalised pole positions are simply the eigen values of the real symmetric matrix C' and hence real. Secondly it means that the mixing matrix i.e. the diagonalising matrix of P' is also the diagonalising matrix for C', which is a J independent real orthogonal matrix. In 2 dimension this can be written in terms of a single parameter, as

$$M = \begin{pmatrix} \cos\theta & \sin\theta \\ -\sin\theta & \cos\theta \end{pmatrix}.$$

The renormalised poles and the mixing angle are, of course, functions of $\alpha_{0,3}$ and the kernel C. For the SU(3) symmetric kernel, they are

$$\alpha_{f,f'} = \tfrac{1}{2}[\alpha_0+\alpha_3+3K\pm\{(\alpha_0-\alpha_3+K)^2+8K^2\}^{1/2}],$$

$$\alpha_{\omega,\phi} = (K \to -K)$$

and

$$\text{Tan } 2\theta^+ = \frac{\sqrt{8}\,K}{\alpha_0-\alpha_3+K}, \quad \text{Tan } 2\theta^- = -\frac{\sqrt{8}\,K}{\alpha_0-\alpha_3-K}.$$

Thus given the planar intercepts α_0, α_3 and one of the physical intercepts α_f say, the model predicts the remaining physical intercepts and the mixing angles. With a $\alpha_0 = 0.58$ (taken from the phenomenological ρ intercept), and a physical f (which is the Pomeron) intercept $\alpha_f = 0.96$ they get a reasonable $\alpha_\omega = 0.41$. The resulting mixing angles are $\theta^+ = 20.3°$ and $\theta^- = -33.7°$.

However, there are 2 points where one may need to improve upon this model.

I. The physical ω intercept is very sensitive to the choice of α_0. Some authors [23, 33] have preferred the value $\alpha_0 = 0.5$, as obtained from the planar bootstrap condition or also from the Chew-Frautschi plot to the value 0.58, in view of possible non-planar effects in the $I = 1$ channel. If one accepts this then for the same α_f, the α_ω goes down to a rather low value of 0.34. The mixing angles remain roughly unchanged.

II. The mixing angle $\theta^- = -34°$ seems to be too large for the K* production data. The reason will be clear from the following relations

$$\gamma^{\omega}_{\pi\rho} = 2g\cos\theta^-,$$

$$\gamma^{f}_{KK^*} = g(\cos\theta^+ - \sqrt{2}\sin\theta^+),$$

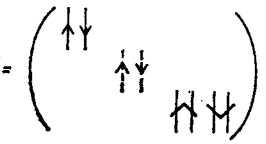
$$\gamma^{\omega}_{KK^*} = g(\cos\theta^- + \sqrt{2}\sin\theta^-).$$

Since the couplings are normalised to the ω production data ($\gamma^{\omega}_{\pi\rho}$), it is evident that a large negative θ^- would enhance the Pomeron contribution and reduce the ω contribution. The K* data shows very little Pomeron contribution; and from a detailed fit Tan et al. [35] have concluded that it can at most accommodate a $\theta^- \lesssim -15°$.

It would be interesting therefore to see if the incorporation of baryonium mixing within its phenomenological bound can give enough flexibility to the model to account for the above discrepancies. This was the phenomenological motivation for Gavai's work.

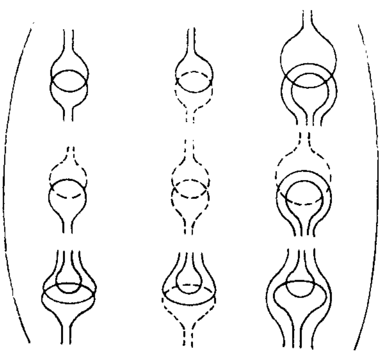
I should emphasise that since we have only an upperbound on the ω-baryonium mixing but no lower bound, the general solution will contain the CR solution, but show what are the maximal variations allowed.

For the ω–ϕ baryonium mixing the planar propagator

$P^{\pm} =$ 

$$= \begin{pmatrix} (J-\alpha_0)^{-1} & & 0 \\ & (J-\alpha_3)^{-1} & \\ 0 & & (J-\alpha_E)^{-1} \end{pmatrix}$$

where α_E is the planar baryonium intercept. With some approximations the cylinder kernel can be written as

$C^{\pm} =$ 

$$= \begin{pmatrix} \pm 2K_0 & \pm\sqrt{2}\,K_0 & 4K_1 \\ \pm\sqrt{2}\,K_0 & \pm K_0 & 2\sqrt{2}\,K_1 \\ 4K_1 & 2\sqrt{2}\,K_1 & 4K_2 \end{pmatrix}$$

where again SU(3) and J independence are assumed for the kernel C. The physical intercepts are obtained by solving the cubic eigen value equation

$$\text{Det}\,[JI-C'] = 0, \qquad C' = C-P^{-1}+JI,$$

and the mixing matrix M is the diagonalising matrix of C'. Being a real orthogonal matrix in 3 dimension, M can be described in terms of the 3 Euler angles, i.e.

$$M = \begin{bmatrix} \cos\xi & \cos\eta\sin\xi & \sin\eta\sin\xi \\ -\sin\xi\cos\zeta & \cos\xi\cos\eta\cos\zeta-\sin\eta\sin\zeta & \cos\xi\cos\zeta\sin\eta+\cos\eta\sin\zeta \\ \sin\xi\sin\zeta & -\cos\xi\cos\eta\sin\zeta-\sin\eta\cos\zeta & -\cos\xi\sin\eta\sin\zeta+\cos\eta\cos\zeta \end{bmatrix}.$$

Without going into details let me discuss the essential features of this model. Where as the original CR model had 3 parameters (α_0, α_3, K) the present one has six (α_0, α_3, α_E, K_0, K_1, K_2). These are determined in terms of the 6 intercepts,

$$\alpha_0 = 0.5 \text{ and } 0.58,$$

$$\alpha_3 = 0.2,$$

$$\alpha_E = -0.5\pm 0.3,$$

$$\alpha_f = 0.95 \pm 0.03,$$

$$\alpha_\omega = 0.43 \pm 0.03,$$

$$\alpha_{B^-} = -1 \pm 0.5.$$

Here the planar baryonium intercept α_E has been chosen from the analyses of Exotic exchange reactions [27, 28] and from the duality analysis [26], where as the renormalised baryonium intercept α_{B^-} has been chosen from those of the inclusive and the total cross--section differences. In terms of these all the mixing angles are fixed. Because of the large uncertainties in the baryonium intercepts, of course, one gets a range of solutions. It is remarkable, however, that the whole range of the ω–B mixing angle is within the phenomenological bound both for $\alpha_0 = 0.50$ and 0.58. The corresponding range for the ω–ϕ mixing angle is from $-33°$ to $-20°$.

More precisely the quantity corresponding to the $\sin^2\theta$ of reference [28] is[7]

$$S = \sin^2 \text{`}\theta\text{'} = M_{13}M_{31} - \sqrt{2}\, M_{32}$$

for $\alpha_0 = 0.58$, $0 < S < 0.04$ $(0 < \text{`}\theta\text{'} < 12°)$ and for $\alpha_0 = 0.50$, $0.009 < S < 0.074$ $(5° < \text{`}\theta\text{'} < 15°)$.

Thus it seems incorporation of ω–B mixing within the phenomenological bound can give a reasonable ω intercept for both $\alpha_0 = 0.58$ and 0.5. Moreover it can bring down the ω–ϕ mixing to a phenomenologically acceptable value.

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[7] In contrast to the SU (3) symmetry assumption here, Ref. [30] had assumed the ϕ–B mixing to be small relative to ω–B ($M_{32} \ll M_{31}$). The truth maybe somewhat inbetween. Under this assumption ($M_{32} \ll M_{31}$) one can see that sin 'θ' = sin ξ up to the first order term in ω–ϕ mixing — i.e. $\sin\eta = -\sin\zeta = 1$ (pure ω–B mixing corresponds to $\eta = -\zeta = \pi/2$).

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