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ABSTRACT

We focus on the potentialities of nucleon-antinucleon ($\bar{N}N$) annihilation as a means of producing new mesonic states. The case for the existence of quasinuclear $\bar{N}N$ bound states is discussed in detail. Strong evidence for a $2^{++}(0^+)$ state of this type has been obtained at LEAR in annihilation from the p -wave ($L = 1$) $\bar{N}N$ system, in support of earlier sightings of this object in $L = 0$ annihilation at Brookhaven. In the next generation of LEAR experiments, the emphasis shifts to the search for mesons containing dynamical excitations of the gluonic field, namely glueballs and hybrids ($Q\bar{Q}g$). We discuss some features of the masses, decay branching ratios and production mechanisms for these states, and suggest particular $\bar{N}N$ annihilation channels which are optimal for their discovery.

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ABSTRACT

We focus on the potentialities of nucleon-antinucleon ($\bar{N}N$) annihilation as a means of producing new mesonic states. The case for the existence of quasinuclear $\bar{N}N$ bound states is discussed in detail. Strong evidence for a $2^{++}(0^+)$ state of this type has been obtained at LEAR in annihilation from the p -wave ($L = 1$) $\bar{N}N$ system, in support of earlier sightings of this object in $L = 0$ annihilation at Brookhaven. In the next generation of LEAR experiments, the emphasis shifts to the search for mesons containing dynamical excitations of the gluonic field, namely glueballs and hybrids ($Q\bar{Q}g$). We discuss some features of the masses, decay branching ratios and production mechanisms for these states, and suggest particular $\bar{N}N$ annihilation channels which are optimal for their discovery.

1. INTRODUCTION

Through the annihilation process, the $\bar{N}N$ system provides a rich source of mesons. Some mesonic states, for instance the $E(1420)$, were first seen in the debris of $\bar{N}N$ annihilation products¹. Although the final state is composed mostly of pions (multiplicity of about 5 at threshold), the main features of annihilation at low energies can be accounted for via a sequential process $\bar{N}N \rightarrow M_1 M_2$, followed by the strong decay of M_1 and M_2 into two or three pions². The quasi-two-body modes $M_1 M_2$ stand out prominently in Dalitz plots of the invariant mass distributions of such two and three meson sub-systems; for instance, the vector mesons (ρ, ω) occur frequently as annihilation products. To search for new mesons X , the two-body mode $\bar{N}N \rightarrow \pi X$ is of special interest. Accordingly, we focus attention on the selection rules for such reactions, as well as the very important kinematical factors which intervene for production of X from initial $L = 0, 1$ states.

These lectures are organized as follows: We first review (in Sect. 2) the general features of the $\bar{N}N$ interaction. The key point here is the existence of very strong tensor forces for isospin $I = 0$ $\bar{N}N$ configurations. This coherent tensor interaction registers its effect in cross section differences for $\bar{p}p$ and $\bar{n}p$ at low energies³, dramatic behavior in certain spin observables⁴, and a band of natural parity ($0^{++}, 1^{--}, 2^{++} \dots$) quasinuclear (QN) bound states with $I = 0$.⁵

In Sect. 3, we argue that at least one member of the QN band has in fact been seen by the ASTERIX collaboration at LEAR, namely the AX(1565) meson⁶. We present calculations of the mass and width, as well as the production and decay branching ratios of such objects. The preferential production of the AX(1565) from $L = 1$ rather than $L = 0$ $\bar{N}N$ states is naturally explained, using techniques developed⁷ for a general discussion of reactions of the type $\bar{N}N \rightarrow \gamma X, \pi X$. The AX(1565) is most likely the same object as seen earlier at Brookhaven in reactions $\bar{p}n \rightarrow 2\pi^- \pi^+$ ⁸ and $\bar{p}n \rightarrow 3\pi^- 2\pi^+$.⁹ In the $\bar{N}N$ potential model, the coherent tensor interaction provides the strong attraction necessary to achieve the rather substantial binding energy observed¹⁰. This is an essentially non-perturbative effect; in QCD, the tensor force obtained from one-gluon exchange is much weaker. We do not consider alternative explanations of the nature of the AX(1565), which must lie outside of the nonet of 2^{++} tensor mesons. Interpretations in terms of a glueball, four quark ($Q^2\bar{Q}^2$) or hybrid ($Q\bar{Q}g$) state are found to be unsatisfactory. Crucial experimental tests of the quasinuclear hypothesis are suggested (search for the $\eta\eta$ decay mode, etc.).

In Sec. 4, we review the theoretical predictions for the masses and preferred decay modes of the $Q\bar{Q}g \equiv X_g$ hybrid mesons. It is argued that the reaction $\bar{N}N \rightarrow \pi X_g$ (or $\gamma X_g, \pi\pi X_g$ in favorable cases) will proceed with measurable branching ratios. The case of *exotic* X_g (ex. 1^{-+}) is worthy of special attention; the candidate X_g state seen in the $\pi^0\eta$ decay channel¹¹ should be seen in the $\bar{p}p \rightarrow \pi^0\pi^0\eta$ reaction, for instance. We present arguments for identifying the E(1420) meson seen in the $\bar{p}p \rightarrow \pi^+\pi^-E$ reaction as the 0^{-+} (0^+) member of the $Q\bar{Q}g$ multiplet obtained by coupling a $Q\bar{Q}$ pair with the quantum numbers of the ω meson to a transverse electric (TE) gluon normal mode.

The $\bar{N}N$ annihilation modes involving ϕ mesons are of particular interest. Here the signal of new physics may be found in the order of magnitude violations of the Zweig rule which are found in certain channels, for instance $\bar{N}N \rightarrow \pi\phi$. It has been suggested¹², and disputed^{13,14}, that such Zweig rule violations signal a soft strange quark ($Q^3s\bar{s}$) component in the nucleon wave function. An alternate interpretation in terms of broad four quark ($Q\bar{Q}s\bar{s}$) resonances has also been suggested¹⁵.

A global conclusion of these lectures is that $\bar{N}N$ annihilation offers exciting prospects for the discovery of new mesonic states, including the elusive exotics. It forms an essential complement to $\pi p, pp, K^-p, e^+e^-$ and $\gamma\gamma$ channels, for instance, as a means of studying meson spectroscopy. The physics program of the new generation of LEAR detectors (Crystal Barrel, Obelix and Jetset) should be given high priority and vigorously pursued.

2. QUASINUCLEAR $\bar{N}N$ BOUND STATES

Models for the $\bar{N}N$ potential at low energies, as well as the nature of the bound state spectrum, have been discussed by numerous authors¹⁶⁻²⁵. The components of the NN and $\bar{N}N$ potentials due to t -channel meson exchange are related by the

Table 1: Signs of Meson Exchange Potentials for the NN System

	$I = 0$				$I = 1$			
	V_0	V_σ	V_{LS}	V_T	V_0	V_σ	V_{LS}	V_T
π	0	-	0	-	0	+	0	+
η	0	+	0	+	0	+	0	+
ρ	-	-	+	+	+	+	-	-
ω	+	+	-	-	+	+	-	-
δ	+	0	+	0	-	0	-	0
σ	-	0	-	0	-	0	-	0

G -parity transformation

$$\begin{aligned}
 V_{NN} &= \sum_i V_i \\
 V_{\bar{N}N} &= \sum_i G_i V_i
 \end{aligned} \tag{1}$$

where $i = \{\pi, \eta\}$, $\{\rho, \omega\}$, $\{\delta, \sigma\}$ correspond to exchanges of pseudoscalar (0^-), vector (1^-) and scalar (0^+) mesons, respectively. The real potential $V_{\bar{N}N}$ must be supplemented by a complex annihilation potential V_A to obtain a realistic optical potential V_{opt} to describe low and medium energy elastic scattering, charge exchange, annihilation cross sections, and spin observables. The potential $V_{\bar{N}N}$ is strongly attractive in certain channels, and supports a spectrum of bound states; these acquire a strong decay width Γ through the action of $\text{Im } V_A$. Generally, these quasinuclear (QN) bound states of the $\bar{N}N$ system will have typical hadronic widths of the order of $\Gamma \approx 100\text{--}200$ MeV, although in some cases narrower states are possible. These QN states are not to be confused with the putative long-lived "baryonium" states of $Q^2\bar{Q}^2$ character, whose narrow width had its origin in a topological selection rule²⁶.

The level order of the lowest-lying states of the $\bar{N}N$ spectrum is essentially model independent (not the absolute mass values!), and can be predicted¹⁸ on the basis of the *strong coherences* of the spin-isospin dependent components of the meson exchange potential $\sum_i G_i V_i$. Each component V_i can be decomposed into central, spin-spin, spin-orbit and tensor pieces:

$$V_i = \left(\begin{array}{c} 1 \\ \tau_1 \cdot \tau_2 \end{array} \right) (V_0 + V_\sigma \sigma_1 \cdot \sigma_2 + V_{LS} L \cdot S + V_T S_{12}) \tag{2}$$

where $S_{12} = 3\sigma_1 \cdot \hat{r}\sigma_2 \cdot \hat{r} - \sigma_1 \cdot \sigma_2$ is the tensor operator. The *signs* of the various terms V_0 , V_σ , V_{LS} and V_T are determined quite generally from the form of the interaction Lagrangian which couples mesons of different spin-parity to nucleons. The pattern of signs for pseudoscalar (π, η), vector (ρ, ω) and scalar (δ, ϵ) exchange is displayed in Tables 1 and 2 for NN and $\bar{N}N$ systems. A *coherence property*¹⁸ becomes immediately evident: for NN , V_σ and V_{LS} terms all have the same sign for $I = 1$, while for $\bar{N}N$ in $I = 0$ states, V_0 and V_T components are coherent. This striking contrast between NN and $\bar{N}N$ spin-isospin dependences is an effect of the G -parity transformation.

Table 2: Signs of Meson Exchange Potentials for the $\bar{N}N$ System

	$I = 0$				$I = 1$			
	V_0	V_σ	V_{LS}	V_T	V_0	V_σ	V_{LS}	V_T
π	0	+	0	+	0	-	0	-
η	0	+	0	+	0	+	0	+
ρ	-	-	+	+	+	+	-	-
ω	-	-	+	+	-	-	+	+
δ	-	0	-	0	+	0	+	0
σ	-	0	-	0	-	0	-	0

The influence of coherences is dramatic, both for NN scattering and for the spectrum of $\bar{N}N$ bound states. As an example, we display in Figure 1 the energy dependence of the $^{33}P_0$ and $^{13}D_2$ phase shifts δ for NN scattering. In both channels, the dominant tensor component of the π exchange potential is attractive, and generates a substantial $\delta > 0$. For $^{33}P_0$, the shorter range but coherent *repulsive* spin-orbit potentials due to scalar and vector exchange counteract the pion tensor force at higher energies, and produce the observed zero in δ . For $^{13}D_2$, on the other hand, the coherence is absent, and the exact δ remains close to the result for π exchange only.

In the $\bar{N}N$ system, the effects of spin dependent potentials are also dramatic, but now the tensor rather than spin-orbit components are most important. In Fig. 2, we illustrate how the predicted $\bar{p}p$ polarization is altered when various spin dependent terms are set to zero. $P(\theta)$ is seen to arise almost entirely from the coherent tensor forces, while $L \cdot S$ and $\sigma_1 \cdot \sigma_2$ terms play a minor role. The tensor effects also dominate for other $\bar{p}p$ spin observables⁴.

The spectrum of quasinuclear bound states is influenced in a qualitative way by the tensor force. For $I = 0$, natural parity ($L = J \pm 1$) states, the $\bar{N}N$ eigenfunctions⁵ are close to one of the configurations $|\alpha_J\rangle$ or $|\beta_J\rangle$ which diagonalize S_{12} , namely

$$\begin{aligned} |\alpha_J\rangle &= \left[(J+1)^{1/2} |L = J-1\rangle + J^{1/2} |L = J+1\rangle \right] / (2J+1)^{1/2} \\ |\beta_J\rangle &= \left[-J^{1/2} |L = J-1\rangle + (J+1)^{1/2} |L = J+1\rangle \right] / (2J+1)^{1/2} \end{aligned} \quad (3)$$

In this basis, the $\bar{N}N$ potential has the form

$$\begin{aligned} V_{\alpha\alpha} &= V_0 + V_\sigma - V_{LS} + 2V_T + \frac{\hbar^2}{m_N} \frac{J(J+1)}{r^2} \\ V_{\beta\beta} &= V_0 + V_\sigma - 2V_{LS} - 4V_T + \frac{\hbar^2}{m_N} \frac{J^2 + J + 2}{r^2} \\ V_{\alpha\beta} &= [J(J+1)]^{1/2} [-V_{LS} + 2\hbar^2/m_N r^2] \end{aligned} \quad (4)$$

For $I = 0$, the effective tensor potential is repulsive in channel α and attractive in channel β ; the effect of the coupling term $V_{\alpha\beta}$ is small⁵. In Fig. 3, we plot $V_{\beta\beta}(r)$

for the $^{13}P_2 - ^{13}F_2$ state of the $\bar{N}N$ system. The square of the radial wave function $b(r)$ for the bound state in this channel is also shown. The r.m.s. radius of the 2^{++} state is of normal hadronic size (0.8 fm); the centrifugal term in Eq. (4) prevents the probability $b^2(r)$ from leaking into the short distance regime $r < 0.5$ fm. In a realistic annihilation potential, such a configuration is expected to have a normal hadronic decay width, of order 100 – 200 MeV.

The presence of strong tensor forces favors the formation of an isoscalar band of $\bar{N}N$ states with quantum numbers

$$J^{\pi G}(I^G) = 0^{++}(0^+), 1^{--}(0^-), 2^{++}(0^+), 3^{--}(0^-) \dots \quad (5)$$

The structure of the low-lying $\bar{N}N$ spectrum is thus considerably simplified. For $I = 1$, there is no coherence of tensor forces, so bound states tend to cluster closer to the $\bar{N}N$ threshold. This pattern is evident in Fig. 4, which shows an $\bar{N}N$ spectrum²² as a function of a parameter λ , the strength of a complex annihilation potential. Other calculations^{16,17,18,21} yield a qualitatively similar picture. The absolute binding energies are very sensitive to the prescription used to regularize the $\bar{N}N$ potential at short distances, but the level order of the lowest lying states is the same in most of the existing calculations.

We emphasize again that the dynamical mechanism for producing relatively high spin (3^{--} in Fig. 4) $\bar{N}N$ bound states is the attractive coherence for $I = 0$ of tensor potentials from pseudoscalar (π, η) and vector (ρ, ω) meson exchange. The off-diagonal tensor force induces virtually complete mixing of orbital states with $L = J \pm 1$, and is particularly important in bringing $J \geq 2$ states down to energies below the $\bar{N}N$ threshold.

The spectrum of $\bar{N}N$ states is quite distinct from that for a four quark ($Q^2\bar{Q}^2$) system, as calculated by Jaffe²⁷ or Badalyan and Kitoroage²⁸, among others. A $Q^2\bar{Q}^2$ spectrum, with Q^2 in a color $\{\bar{3}\}$ and $\bar{Q}^2$ in $\{3\}$, is shown in Fig. 5. We can have the diquark in a spin-isospin state α with $I = S = 0$ or β with $I = S = 1$. This gives trajectories $A = \alpha\bar{\alpha}$, $B = \alpha\bar{\beta} \pm \bar{\alpha}\beta$ and $C = \beta\bar{\beta}$. The trajectory A , which couples most strongly to $\bar{N}N$, has the same quantum numbers as the lowest-lying band of $\bar{N}N$ states given in Eq. (5). Trajectory B displays G -parity doublets, while C has degenerate $I = 0, 1, 2$ states. These features are not shared by the spectrum of $\bar{N}N$ quasinuclear states.

The considerations of this section are speculative. In the next section, we argue that at least some of the predicted $\bar{N}N$ bound states may be realized in nature.

3. EXPERIMENTAL EVIDENCE FOR $\bar{N}N$ QUASINUCLEAR BOUND STATES

The ASTERIX group at LEAR studied the $\bar{p}p \rightarrow \pi^+\pi^-\pi^0$ reaction from an initial $L = 1$ atomic state⁶. They saw a peak in the $\pi^+\pi^-$ mass spectrum in the vicinity of 1565 MeV, which they denoted as the AX(1565). This data is shown in Fig. 6. From an analysis of decay angular distributions and the non-observation of a peak in the $I = 1$ $\pi^\pm\pi^0$ mass spectrum, an assignment of $2^{++}(0^+)$ quantum numbers

was given. The branching ratio is

$$B(\bar{p}p(L=1) \rightarrow \pi^0 + \text{AX}(1565)) \cdot B(\text{AX}(1565) \rightarrow \pi\pi) = (5.6 \pm 0.9) \times 10^{-3} \quad (6)$$

There is good evidence that the AX(1565) was seen earlier by Gray *et al.*⁸ in bubble chamber experiments (essentially pure $L=0$) in the $\bar{p}n \rightarrow 2\pi^-\pi^+$ and $\bar{p}p \rightarrow 3\pi^0$ channels. Some of this data is plotted in Fig. 7. The quoted⁸ branching ratio is

$$B(\bar{p}n(L=0) \rightarrow \pi^- + \text{AX}(1565)) \cdot B(\text{AX}(1565) \rightarrow \pi\pi) \approx (3.9 \pm 0.6) \times 10^{-3} \quad (7)$$

Correcting for an isospin factor. Eqs. (6) and (7) imply the ratio

$$\frac{B(\bar{p}p(L=0) \rightarrow \pi^0 + \text{AX}(1565))}{B(\bar{p}p(L=1) \rightarrow \pi^0 + \text{AX}(1565))} \approx 0.35_{-0.10}^{+0.13} \quad (8)$$

Another broad structure in the same mass region has been seen in the $\bar{p}n \rightarrow 3\pi^-2\pi^+$ reaction at Brookhaven by Bridges *et al.*⁹ and at LEAR by Ahmad *et al.*²⁹. The quantum number assignment preferred by Bridges *et al.*⁹ is $2^{++}(0^+)$, identical to the AX(1565). The four pion invariant mass plots are shown in Figs. 8 and 9. The peak in the 4π spectra from the $\bar{p}d \rightarrow 3\pi^-2\pi^+p$ reaction lies somewhat lower than the AX(1565), and its apparent mass decreases with increasing spectator momentum, as seen in Fig. 9. This effect has been explained by Kolybasov *et al.*³⁰ in terms of final state interactions of the pions with the proton. A plausible hypothesis is that the AX(1565) has been seen by Bridges *et al.*⁹ in its $\rho\rho$ decay mode; if this is tenable, it is clear that $B(\text{AX}(1565) \rightarrow \rho\rho) > B(\text{AX}(1565) \rightarrow \pi\pi)$. Furthermore, Gray *et al.*⁸ indicate that the $\bar{K}K$ decay mode is strongly suppressed, so that the structure at about 1530 MeV in Fig. 7 cannot be identified with the $f'(1525)$, i.e., the well known $2^{++}(0^+) s\bar{s}$ meson.

The arguments for interpreting the AX(1565) as the $^{13}P_2 - ^{13}F_2$ quasinuclear bound state of the $\bar{N}N$ system have been developed in detail by Dover *et al.*³¹. The mechanism which gives strong binding for this particular $I=0$ configuration is the coherent tensor potential. Calculations employing a complex $\bar{N}N$ optical potential fit to scattering data predicted a quasinuclear (QN) state in the $2^{++}(0^+)$ channel with roughly the right mass and width: for instance, Vinh Mau²¹ gives $M \approx 1500$ MeV, $\Gamma \approx 65$ MeV. Long ago, Dover *et al.*⁷ predicted branching ratios for reactions of the type $(\bar{N}N)_{\text{atom}} \rightarrow \pi + (\bar{N}N)_{\text{QN}}$; a sample of these is shown in Fig. 10. When these estimates are adjusted to obtain the observed width $\Gamma \approx 170$ MeV of the AX(1565) and to incorporate the correct phase space factor for the reaction $\bar{N}N \rightarrow \pi + \text{AX}(1565)$, a branching ratio consistent with Eq. (6) is obtained³¹. The preferential production of AX(1565) from an initial $L=1$ $\bar{N}N$ state, as indicated by Eq. (8), is understood as essentially a kinematical effect.

An estimate of the relative branching ratios for the two-meson decays $\bar{N}N \rightarrow M_1 M_2$ has been provided in Ref. 31, assuming the mechanism shown schematically in Fig. 11. The $Q\bar{Q}$ vertices are taken to have vacuum quantum numbers (the 3P_0 model). Such a model has been successfully applied to a description of $\bar{p}p$ annihilation at rest and in flight^{32,33}. Meson and baryon resonance decay are well accounted for using an effective 3P_0 vertex. This approximation is motivated by flux tube models of QCD in the strong coupling limit^{34,35}. The strength of the 3P_0 vertex has been determined from the $Q\bar{Q}$ and gluon condensates by Weber³⁶.

The transition amplitude T corresponding to Fig. 3 is given by³³

$$T = \int d^3k_1 \dots d^3k_6 d^3k'_1 d^3k'_4 d^3k'_7 d^3k'_8 \psi_{M_1 M_2}^\dagger \mathcal{O} \psi_{\bar{N} N} \quad (9)$$

where $\mathcal{O}$ is the operator

$$\mathcal{O} = \lambda \delta(k'_1 - k_1) \delta(k'_4 - k_4) \delta(k_3 + k_6) \delta(k_2 + k_5) \delta(k'_7 + k'_8) \\ \sigma_{-\mu}^{36} y_{1\mu} (k_3 - k_6) \sigma_{-\nu}^{25} y_{1\nu} (k_2 - k_5) \sigma_{-\gamma}^{7'8'} y_{1\gamma} (k'_7 - k'_8) (-)^{\mu+\nu+\gamma} / 3\sqrt{3} \quad (10)$$

where the factor $(-)^{\mu} / \sqrt{3} = (1\mu 1 - \mu | 00)$ ensures that each vertex (with z -component of spin $-\mu$) is coupled to 0^{++} quantum numbers. We now sum over allowed spin-flavor assignments for the Q and $\bar{Q}$ lines in Fig. 11, and assume $SU(6)$ wave functions for N , $\bar{N}$, M_1 and M_2 . For each transition $\bar{N}N \rightarrow M_1 M_2$ with final c.m. momentum q , we obtain

$$\Gamma = CqF(q)W_{12} \quad (11)$$

where $F(q)$ is a kinematical form factor proportional to $|T|^2$ and W_{12} is a spin-flavor weight. These have been tabulated by Maruyama *et al.*³³.

The predicted³¹ branching ratios for the decay of a $^{13}P_2$ quasinuclear $\bar{N}N$ state are displayed in Fig. 12. We note that decay is likely to occur at a measurable rate to several final states. The $\rho\rho$ mode is largest, which is consistent with the large signal in this channel seen by Bridges *et al.*⁹, assuming that the X(1480) is in fact the same object as the AX(1565). A $\pi\pi$ branching ratio of order 0.2, together with the calculated production rate^{7,31} for $\bar{N}N(L=1) \rightarrow \pi + \text{AX}(1565)$, is consistent with Eq. (6). The $\bar{K}K$ decay mode of the $^{13}P_2$ state (not shown in Fig. 12) is predicted to be strongly suppressed, in agreement with the absence of a signal in the experiment of Gray *et al.*⁸. Such a suppression is also seen in the recent data of Doser *et al.*³⁷, who give

$$B(\bar{p}p(L=1) \rightarrow K^+ K^-) / B(\bar{p}p(L=1) \rightarrow \pi^+ \pi^-) \approx 0.06 \pm 0.012 \quad (12)$$

From Fig. 12, we note that $B(\text{AX}(1565) \rightarrow \eta\eta) \approx 0.05$. This would be a clear experimental signature, which could be looked for in the $\bar{p}p \rightarrow \pi^0 \eta\eta$ channel with the Crystal Barrel detector at LEAR.

The picture of $\bar{N}N$ quasinuclear states presented here is somewhat similar to the interpretation³⁸ of the S^* and δ scalar mesons as $I = 0, 1$ S -wave “ $\bar{K}K$ molecules” rather than 3P_0 $Q\bar{Q}$ states. One could also attempt to interpret the AX(1565) in terms of a $Q^2\bar{Q}^2$ states^{39,40}. Note, however, that in the Bag Model calculation of Jaffe²⁷, as seen in Fig. 5, the lowest-lying $2^{++}(0^+)$ state at 1550 MeV is accompanied by degenerate $I = 1, 2$ partners, for which there is no evidence in the $\bar{p}p \rightarrow \pi^+ \pi^- \pi^0$ data.

4. PRODUCTION OF HYBRID MESONS IN $\bar{N}N$ ANNIHILATION

We consider in this section the utility of the $\bar{N}N$ annihilation process as a means of producing hybrid mesons, namely $Q\bar{Q}g$ composites in which the gluon field g is excited to a transverse electric (TE) or transverse magnetic (TM) normal

mode. Among these are found “exotic” mesons, which have J^{PC} quantum numbers forbidden for a non-relativistic $Q\bar{Q}$ system, for example 0^{--} , 0^{+-} or 1^{-+} .

4.1 The Spectrum of $Q\bar{Q}g$ Hybrids

There are estimates of the masses of hybrids in a variety of models, for instance the Bag Model^{41–43}, the flux tube model⁴⁴, the technique of QCD sum rules^{45–47}, the constituent gluon model^{48–50}, and lattice QCD⁵¹. In the latter approach, a static $Q\bar{Q}$ potential is calculated in the presence of a gluon field with non-vacuum quantum numbers. A recent review of the phenomenology of hybrid mesons is due to Godfrey⁵².

In the Bag Model, the lowest lying $Q\bar{Q}g$ states are obtained by coupling an S -wave $Q\bar{Q}$ pair (0^{-+} (1S_0) or 1^{--} (3S_1)) to a TE (1^{+-}) or TM (1^{--}) mode of the gluon field. In a spherical cavity, the TM mode lies about 500 MeV above the TE mode, so the lowest-lying $Q\bar{Q}g$ hybrids are 1^{--} ($^1S_0 \otimes TE$) and 0^{-+} , 1^{-+} , 2^{-+} ($^3S_1 \otimes TE$). The 1^{-+} exotic member comes in four flavor combinations ρ_g , ω_g , ϕ_g , and K_g^* , having the quark structure $(u\bar{u} - d\bar{d})/\sqrt{2}$, $(u\bar{u} + d\bar{d})/\sqrt{2}$, $s\bar{s}$ and $\bar{s}d$, respectively. According to Barnes *et al.*⁴¹, the masses are approximately 1.4, 1.55, 1.52 and 1.75 GeV/ c^2 , respectively. The mass scale depends on the ratio of TE and TM self energies, and the values given by Chanowitz and Sharpe⁴² are somewhat larger. The near degeneracy of ω_g and K_g^* is a signature of hybrid structure in the Bag Model.

In the flux tube model⁴⁴, the ρ_g and ω_g lie at about 1.9 GeV/ c^2 , and 0^{+-} ($0^{-}, 1^{+}$) and 2^{+-} ($0^{-}, 1^{+}$) exotics are predicted at the same mass, a degeneracy not present in the Bag Model. In this case, it would not be possible to produce exotics in $\bar{N}N$ annihilation at rest. In the sum rule approach⁴⁵, the ρ_g is predicted to lie in the range 1.6 – 2.1 GeV/ c^2 , with a most probable mass at the upper end of this range.

4.2 Decays of $Q\bar{Q}g$ Hybrid Mesons

Decay branching ratios of hybrids have been estimated in the flux tube model⁴⁴, using QCD sum rules³⁹, and from the constituent gluon picture^{53,54}. In all of these calculations, *strong selection rules* restrict the possible decay modes of hybrids. In the $SU(6)$ approximation, for $Q\bar{Q}g$ hybrids constructed with a $1^{+-} TE$ gluon, we have

$$\begin{aligned} (Q\bar{Q})_{\ell=0} \otimes TE &\rightarrow sp (\ell_f = 0) \\ &\not\rightarrow sp (\ell_f = 2) \\ &\not\rightarrow ss (\ell_f = 1) \end{aligned} \quad (13)$$

where s denotes a $(Q\bar{Q})_{\ell=0}$ meson and p is a $(Q\bar{Q})_{\ell=1}$ meson, and ℓ_f is the meson-meson relative orbital angular momentum. The decay mechanism which leads to these selection rules is depicted in Fig. 13. In the allowed decays of $Q\bar{Q}g$, the spin of the gluon field appears as an internal orbital excitation of a $Q\bar{Q}$ meson rather than as a relative P -wave between two s mesons. These selection rules emerge in both the constituent gluon and flux tube models.

4.3. Candidates for $Q\bar{Q}g$ Mesons

Alde *et al.*¹¹ have presented evidence for the existence of a 1^{-+} (1^{-}) meson (the ρ_g) with a mass of 1406 ± 20 MeV/c² and a width of 180 ± 30 MeV. The putative ρ_g was seen in its $\pi^0\eta$ decay mode in the $\pi^-p \rightarrow \pi^0\eta n$ reaction at 100 GeV. The observed mass lies close to the Bag Model prediction, but according to Eq. (13), $\rho_g \rightarrow \pi^0\eta$ ($\ell_f = 1$) should be essentially a forbidden decay mode, except for $SU(3)$ breaking contributions. This is a puzzle. There are a number of other candidates for mesons which fall outside of the usual $Q\bar{Q}$ nonets. A recent comprehensive review of the experimental situation is due to Landsberg⁵⁵. Further studies of the $\pi^-p \rightarrow \rho_g^-p$ reaction are underway at Brookhaven and KEK.

4.4 Hybrid Production in $\bar{N}N$ Annihilation

We expect that the $\bar{N}N$ annihilation process is a copious source of gluons through the successive destruction of $Q\bar{Q}$ pairs from the initial state. Non-perturbative $Q\bar{Q}$ annihilation can lead to glue in its ground state (with vacuum quantum numbers) or, in some cases, an excited gluonic configuration. Such excitations could combine with each other to form glueballs, or fuse with a $Q\bar{Q}$ pair to form a hybrid meson. A possible mechanism for $Q\bar{Q}g$ formation in $\bar{N}N$ annihilation is shown in Fig. 14. Calculations are underway⁵⁶ to estimate the relative branching ratios for annihilation reactions $\bar{N}N \rightarrow (Q\bar{Q}) + (Q\bar{Q}g)$, for instance $\bar{p}p \rightarrow \pi + \rho_g$.

The principal advantage of $\bar{N}N$ annihilation at rest is that one can tag the formation of an $L = 1$ atomic state³⁷ using a gaseous hydrogen target by observing a $3d \rightarrow 2p$ gamma ray in coincidence with the mesonic decay products. In liquid hydrogen, on the other hand, the annihilation occurs dominantly from the $L = 0$ state. This ability to filter the initial L is crucial. An example is provided by Eq. (8): the production of AX(1565) is predicted and observed to be enhanced from the $L = 1$ state, thus strengthening the interpretation of AX(1565) as a quasinuclear $\bar{N}N$ bound state. By comparing $\bar{p}n$ to $\bar{p}p$, one can filter the initial isospin. One further selects G or C by choosing certain exclusive final states. To illustrate the power of quantum number filtration, we consider the three reactions $\bar{p}p \rightarrow \pi^0 X^0$, $\bar{p}p \rightarrow \pi^\pm X^\mp$ and $\bar{p}n \rightarrow \pi^- X^0$. For $\bar{p}p \rightarrow \pi^0 X^0$, C parity conservation considerably diminishes the number of allowed transitions. In $\bar{p}p \rightarrow \pi^\pm X^\mp$, we select only isovector and isotensor mesons $X^\mp$, whereas in $\bar{p}n \rightarrow \pi^- X^0$, we prepare an $I = 1$ initial state. Among the $J^{\pi C}$ exotic mesons with $J = 0, 1$, only the ρ_g can be produced in all three of the above reactions for both $L = 0, 1$.

Consider the production of ρ_g . If ρ_g decays predominantly in sp modes like πB or πD , as predicted by Eq. (13), then a promising sequence of reactions is

$$\bar{p}p \rightarrow \pi^\pm \rho_g^\mp, \quad \rho_g^\mp \rightarrow \pi^\mp B^0, \quad B^0 \rightarrow \pi^0 \omega, \quad \omega \rightarrow \pi^0 \gamma \quad (14)$$

as proposed for study by the Crystal Barrel collaboration at LEAR. The $\omega \rightarrow \pi^0 \gamma$ radiative decay is chosen to minimize combinatorial background. If the observation of the ρ_g by Alde *et al.*¹¹ is correct, on the other hand, then the neutral channels $\pi^0 \pi^0 \eta$ and $\pi^0 \pi^0 f$ are worthy of study. For $L = 0$, these correspond to the decay

chains

$$\begin{aligned} \bar{p}p (^{11}S_0) &\rightarrow \pi^0 \rho_g^0 (\ell_f = 1), \quad \rho_g^0 \rightarrow \pi^0 \eta (\ell_f = 1), \quad \eta \rightarrow 2\gamma \\ \bar{p}p (^{11}S_0) &\rightarrow \pi^0 \rho_g^0 (\ell_f = 1) \quad \rho_g^0 \rightarrow \pi^0 f (\ell_f = 2), \quad f \rightarrow 2\pi^0 \end{aligned} \quad (15)$$

These are optimum channels for study by the Crystal Barrel detector. The $\pi^0 \pi^0 \eta$ system would be detected as six photons.

As another example for ρ_g , consider $\bar{p}p \rightarrow 2\pi^+ 2\pi^-$. The “signal” is $\bar{p}p (^{11}S_0, ^{33}S_1) \rightarrow \pi^\pm \rho_g^\mp (\ell_f = 1)$, $\rho_g^\mp \rightarrow \pi^\mp \rho^0 (\ell_f = 1)$, and the background due to ordinary two-body channels arises from $\bar{p}p (^{11}S_0, ^{33}S_1) \rightarrow \pi^\pm A_2^\mp (\ell_f = 2)$, $A_2^\mp \rightarrow \pi^\mp \rho^0 (\ell_f = 2)$, for instance. As with the analysis of Alde *et al.*¹¹ for the $\pi^0 \eta$ system, one could look for an interference of $\ell_f = 0$ and $\ell_f = 2$ $\pi \rho$ amplitudes.

The 1^-+ isoscalar partner of the ρ_g , the ω_g , could be searched for in the neutral channels $\bar{p}p \rightarrow 4\pi^0$, $\pi^0 \eta \eta$, $\pi^0 \eta \eta'$. In $\bar{p}p \rightarrow 3\pi^0$, there is no background from $\pi \rho$, and $\pi^0 f$ is known to be small, so $\bar{p}p \rightarrow \pi^0 \omega_g$ might be observable. Earlier experiments^{9,57} show very puzzling differences in the Dalitz plots for $\bar{p}p \rightarrow 3\pi^0$, $\bar{p}p \rightarrow \pi^+ \pi^- \pi^0$ and $\bar{p}n \rightarrow 2\pi^- \pi^+$, which are still unexplained.

4.5. The E(1420) as a $Q\bar{Q}g$ Hybrid?

The ASTERIX group has recently presented results⁵⁸ for the annihilation channel $\bar{p}p \rightarrow \pi^+ \pi^- K^\pm \pi^\mp K^0$ with a gaseous target. In the invariant mass spectrum of the $K^\pm \pi^\mp K^0$ system, evidence is seen for the D(1285) and E(1420) mesons. In earlier experiments¹ in liquid hydrogen (mostly $L = 0$ annihilation), the E(1420) was seen with a larger branching ratio, indicating a formation rate which depends strongly on L . The dependence of the E(1420) branching ratio on the fraction of $L = 1$ annihilation [$61 \pm 6\%$ for ASTERIX in this channel] is shown in Fig. 15. Baillon *et al.*¹ saw no charged E 's, so $I = 0$. Also, $E \rightarrow K_s K_s \pi^0$ was seen, but $E \rightarrow K_s K_L \pi^0$ was not, so $C = 1$ $G = 1$ for the E . The possible quantum number assignments for the E are then $J^{\pi C}(I^G) = 0^{-+}(0^+)$ or $1^{++}(0^+)$. On the basis of the observed suppression of E production from $L = 1$ initial states (see Fig. 15), Duch *et al.*⁵⁸ prefer $E(1420) = 0^{-+}(0^+)$. The argument is a simple kinematical one: if the E is a $0^{-+}(0^+)$ pseudoscalar meson, the allowed transitions from $L = 0, 1$ are

$$\begin{aligned} ^{11}S_0 &\rightarrow [(\pi\pi)_{\ell=0} \otimes E]_{\ell_f=0} \\ ^{33}S_1 &\rightarrow [(\pi\pi)_{\ell=1} \otimes E]_{\ell_f=1} \\ ^{13}P_1 &\rightarrow [(\pi\pi)_{\ell=0} \otimes E]_{\ell_f=1} \\ ^{31}P_1 &\rightarrow [(\pi\pi)_{\ell=1} \otimes E]_{\ell_f=0} \end{aligned} \quad (16)$$

Since relatively little phase space is available, the $\ell = \ell_f = 0$ case is favored kinematically, implying dominant production of the E(1420) from the $^{11}S_0$ initial state. Similarly, for a $1^{++}(0^+)$ axial vector meson like the D(1285), we have

$$^{11}S_0 \rightarrow [(\pi\pi)_{\ell=0} \otimes D]_{\ell_f=1}$$

$$\begin{aligned}
^{33}S_1 &\rightarrow [(\pi\pi)_{\ell=1} \otimes D]_{\ell_f=0} \\
^{13}P_1 &\rightarrow [(\pi\pi)_{\ell=0} \otimes D]_{\ell_f=0} \\
^{31}P_1 &\rightarrow [(\pi\pi)_{\ell=1} \otimes D]_{\ell_f=1}
\end{aligned} \tag{17}$$

The same phase space argument favors production of the D(1825) from the $^{13}P_1$ initial state. This is what is observed experimentally: the D(1285) is not seen by Baillon *et al.*¹ ($L = 0$) but is quite distinct in the ASTERIX data⁵⁸ (60% $L = 1$).

In liquid hydrogen, the branching ratios for $\bar{p}p \rightarrow \pi^+\pi^-E(1420)$ are given¹ as $(2.0 \pm 0.2) \times 10^{-3}$ for the 1S_0 state and $(0.6 \pm 0.06) \times 10^{-3}$ from the 3S_1 state. Removing the $(2S+1)$ factor for an assumed statistical population of 1S_0 and 3S_1 initial states, we find

$$\frac{\Gamma(^{33}S_1 \rightarrow \pi^+\pi^-E)}{\Gamma(^{11}S_0 \rightarrow \pi^+\pi^-E)} \approx \frac{1}{10} \tag{18}$$

This seems to support the kinematical argument based on Eq. (16), since transitions from $^{33}S_1$ are suppressed because $\ell = \ell_f = 1$. However, then it is not clear that the $L = 1$ production of E is as strongly suppressed as indicated by the solid line in Fig. 1. Indeed, transitions $^{13}P_1, ^{31}P_1 \rightarrow \pi^+\pi^-E$ should be less suppressed than $^{33}S_1 \rightarrow \pi^+\pi^-E$, since only one relative P -wave is required (rather than two) in the final state. Note that a branching ratio for $L = 1 \rightarrow \pi^+\pi^-E$ up to 2×10^{-4} can easily be accommodated by keeping the straight line within the errors in Fig. 15.

The allowed two body decay modes of a $0^{-+}(0^+)$ meson are

$$\begin{aligned}
0^{-+}(0^+) &\rightarrow \pi\delta (\ell_f = 0), \eta\sigma (\ell_f = 0), \bar{K}K^* + c.c (\ell_f = 1), \\
&\rho\rho (\ell_f = 1), \omega\omega (\ell_f = 1), \pi A_2 (\ell_f = 2)
\end{aligned} \tag{19}$$

where $\sigma = (\pi\pi)_{\ell=0}$. For E(1420), the decay analysis of Duch *et al.*⁵⁸ prefers $E \rightarrow \pi\delta$ as the dominant $K^\pm\pi^\mp K^0$ decay mode, but $\bar{K}K^* + c.c$ is not ruled out at the 20 - 30% level.

The E(1420) has the quantum numbers of the η and η' . It lies too low in mass to be a radial excitation. Thus it is a prime candidate for a non- $Q\bar{Q}$ meson. It is not unreasonable to interpret the E(1420) as a $Q\bar{Q}g$ hybrid meson of the type $[\omega \otimes TE]_{0^{-+}}$, the pseudoscalar partner of the exotic 1^{-+} meson ω_g discussed earlier. From Eq. (13), we then have the approximate selection rules

$$\begin{aligned}
E(1420) &\rightarrow \pi\delta, \eta\sigma \\
E(1420) &\not\rightarrow \bar{K}K^*, \rho\rho, \omega\omega, \pi A_2
\end{aligned} \tag{20}$$

The mass of the E is close to the value predicted for a $0^{-+}(0^+) Q\bar{Q}g$ meson in the Bag Model⁴¹. Eq. (20) would explain naturally the preference for $\pi\delta$ over $\bar{K}K^*$ in the decay of the E . A puzzle concerns the decay $E \rightarrow \eta\sigma \rightarrow \eta\pi\pi$, which was not seen by Ando *et al.*⁵⁹. It is possible that this is caused by a destructive interference between the two quasi-two-body modes $\pi\delta$ and $\eta\sigma$, which lead to the same $\pi\pi\eta$ final state. The consistency of the hypothesis that E(1420) is a $Q\bar{Q}g$ meson is being studied⁵⁶ in detail, based on the production and decay mechanisms of Figs. 13 and 14.

5. CONCLUSIONS

Studies of $\bar{N}N$ annihilation offer exciting prospects for the discovery of mesonic states beyond the usual $Q\bar{Q}$ nonets. We have presented plausible interpretations of the new objects (for instance, the AX(1565) and the E(1420)) already seen in the debris of $\bar{N}N$ annihilation. The $\bar{N}N$ system forms an essential complement to pp , e^+e^- , K^-p , $\gamma\gamma$ and other channels as a means of shedding new light on the meson spectrum. The LEAR facility is unique for these studies and its physics program should be vigorously pursued in the 1990's, employing the full range of new detectors coming on line.

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Figure 1: Nucleon-nucleon scattering phase shifts δ as a function of kinetic energy E . The dashed curves include only pion exchange, while the solid curves labeled "EXACT" correspond to a full calculation with pseudoscalar, scalar and vector meson exchange. The strong effect of the coherent repulsive spin-orbit potentials is evident in the ${}^{33}P_0$ channel.

Figure 2: Angular distribution of elastic $\bar{p}p$ polarization $P(\theta)$. The solid curve is the result of an optical model calculation⁴, and the dashed curves show the effect of turning off various spin-dependent components of the meson exchange potential. The effect of removing the tensor (S_{12}) term is dramatic.

Figure 3: The potential $V_{\beta\beta}(r)$ of Eq. (4) as a function of r for the $2^{++}(0^+)$ state of the $\bar{N}N$ system, together with $b^2(r)$, the square of the bound state wave function; from Ref. (5).

Figure 4: Binding energies (left-hand scale) and masses (right-hand scale) for $\bar{N}N$ bound states, as a function of the strength λ of the annihilation potential; from Niskänen and Green, Ref. (22).

Figure 5: Masses of $Q^2 \bar{Q}^2$ states, from Jaffe, Ref. (27).

Figure 6: The $\pi^+ \pi^-$ mass spectrum for $\bar{p}p \rightarrow \pi^+ \pi^- \pi^0$ annihilation from atomic P -states; from May *et al.*, Ref. (6).

Figure 7: The $\pi^+\pi^-$ mass spectrum for $\bar{p}n \rightarrow 2\pi^-\pi^+$ annihilation from atomic S -states; from Gray *et al.*, Ref. (8).

Figure 8: In (A), the mass spectra of 4π systems recoiling against a π^+ or π^- are shown. In (B), the difference spectrum is plotted. The solid curve in (B) represents a fit to an assumed reaction $\bar{p}n \rightarrow \pi^- + X(1480)$; from Bridges *et al.*, Ref. (9).

Figure 9: Four pion difference spectra, as in Fig. 8(B), for the reaction $\bar{p}d \rightarrow 3\pi^- 2\pi^+ p$. The dashed curve corresponds to recoil protons in the low momentum “spectator” region, while the solid curve represents the data for “non-spectator” protons of higher momentum; from Ahmad *et al.*, Ref. (29).

Figure 10: Widths Γ for π^0 emission from quasi-atomic (QA) $\bar{N}N$ bound states to deeply bound quasi-nuclear (QN) states. Selected transitions between QN states are also shown.

Figure 11: Mechanism for $\bar{N}N$ annihilation into two mesons M_1, M_2 via quark-antiquark ($Q\bar{Q}$) creation and destruction.

Figure 12: Predicted relative branching ratios for the decay of the $2^{++}(0^+)$ $\bar{N}N$ bound state into two meson channels, as a function of mass; from Dover *et al.*, Ref. (31).

Figure 13: Mechanism for the decay of a $Q\bar{Q}g$ hybrid meson into two $Q\bar{Q}$ mesons.

Figure 14: Mechanism for the production of $Q\bar{Q}g$ hybrid mesons in $N\bar{N}$ annihilation.

Figure 15: Branching ratio for $\bar{p}p \rightarrow \pi^+\pi^- E$ as a function of P wave annihilation fraction; from Duch *et al.*, Ref. (58).