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ANTINUCLEON–NUCLEON ANNIHILATION DYNAMICS*

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The antinucleon–nucleon annihilation is predominantly described by a hot-fireball process where the many final quantum numbers are distributed in a statistical fashion. It is argued that caution must be used in employing the long-range meson-exchange forces to describe the protonium atomic states. The simplest processes of two final mesons do show puzzling behavior which might be a reflection of quark dynamics, but no guiding principles for these quark calculations have been established yet.

1. INTRODUCTION

The measured antinucleon–nucleon ($\bar{N}N$) annihilation processes reflect the fact that there are many degrees of freedom available and these degrees of freedom are distributed statistically among most of the final mesonic states. No other spin or isospin effects are seen experimentally in most of the annihilation channels. If there are signatures reflecting the underlying quark–gluon dynamics we have to search for these in the above mentioned statistical "noise" as will be discussed.

First, I will concentrate on the dominant, statistical, "hot plasma (fireball)" nature of $\bar{N}N$ annihilation and compare this description with experimental data. In addition, we know that for a given $\bar{N}N$ energy there are many nearby meson thresholds. These thresholds seem to be important in describing the measured $N\bar{N}$ annihilation into three-, four-, five- *etc.* mesons cross sections at LEAR energies. After having presented the experimental evidence for the above statements, I hope to have convinced you that we have to focus on a few two-meson annihilation channels in order to search for signals of the underlying quark-gluon "dynamics". This will be the second part of my presentation, where I will give a short, critical review of these quark model calculations. We should keep in mind that what I will present are model ideas and calculations. *A priori*, we do not have any good theoretical guiding principles from which to start these specific calculations. With the cooperation of our experimental colleagues we hope to establish these guiding principles (if they can be found) and try to link these to QCD inspired models in order to enhance our understanding of $\bar{N}N$ annihilation. We expect (hope) that the LEAR facility at CERN will contribute significantly toward resolving the question: "What are the signatures of the underlying quark-gluon dynamics in $\bar{N}N$ annihilation?" The search we are faced with resembles the search for signals of the quark gluon plasma in relativistic heavy-ion collisions. We are lucky in the sense that we can concentrate on a few simple annihilation channels to look for our signatures.

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To begin with we should keep in mind that the ratio of the total annihilation to the total cross section at LEAR energies is

$$\sigma_{\text{ann}}/\sigma_{\text{tot}} \simeq 2/3 \ . \tag{1}$$

This says we do not have a black sphere situation where we cannot ask what takes place inside the black sphere itself.

2. THE DOMINANT ANNIHILATION CHANNELS

2.1 Statistical Distributions

We are here concerned about $\bar{N}N$ annihilating into many pions or into $\bar{K}K$ plus pions, reactions which make up the bulk of the annihilation cross section. At rest the $\bar{p}p$ center-of-mass energy is

$$E_{\text{c.m.}} = 2M \simeq 1876\,\text{MeV} \tag{2}$$

which allows for $13\,\pi^0$ in the final state. On the average $\bar{N}N$ annihilates into five pions ($\langle n\rangle \simeq 5$) at LEAR energies. In Fig. 1 is shown the measured yield[1] for different number (n) of pions detected. For a given number (n) of pions in the final state, Pais[2] proposed that the different charge combination of the final pions should show a statistical distribution in isospin, for example, for $n = 5$ we have the possible charge combinations $5\pi^0$, $3\pi^0\pi^+\pi^-$ and $\pi^0 2\pi^+ 2\pi^-$. Pais assumed all charge channels have equal moduli and this statistical distribution of charges is confirmed by experiments for $n = 3$ to 7.[1]

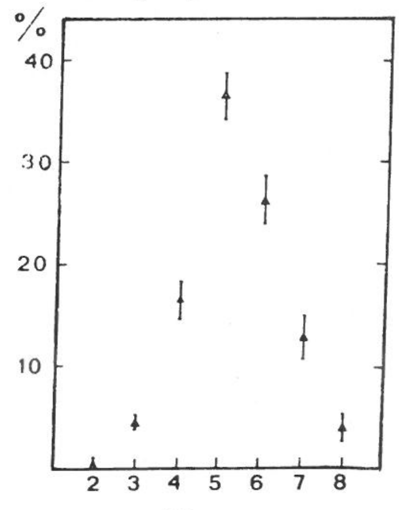

Figure. 1
Total pion multiplicity distribution for $p_{\text{lab}} \simeq 1.6$ *GeV/c.*

2.2 The Hot Plasma and Meson Evaporation

In the hot plasma models again all information about the incoming spin and isospin states are lost (averaged). We can use this model to, for example, describe the energy spectrum of the pions emitted in the reaction

$$\bar{p}p \rightarrow \text{“hot plasma”} \rightarrow \pi + \text{anything} \ . \tag{3}$$

The model assumes that $\bar{N}N$ annihilate into fragments (pions, meson resonances, ...) which are in thermal equilibrium.[3] Further, it is assumed that the pions evaporated from this "hot plasma" have the same energy distribution as the fragments. The cross section for the reaction, Eq. (3), is according to Kimura and Saito[4] (see Ref. [5] for later discussions, as well as earlier work by Hamer[3])

$$\frac{d^3\sigma}{dk^3} \sim \frac{k^2}{\omega^3} e^{-\omega/T} \tag{4}$$

where the pion energy is $\omega = \left(m_\pi^2 + k^2\right)^{1/2}$ and the factor k^2/ω^2 is required by Adler's soft pion consistency condition.[6] The temperature T is determined by the measured pion energy spectrum[7], Fig. 2, and Kimura and Saito found $T = 117$ MeV (compare to Hagedorn's temperature[8] $T = 160$ MeV). This temperature is not too different from what is found in relativistic heavy-ion collisions. This cross section (4) measures the momentum distribution[4] $\Phi_\pi(k)$

$$\frac{d^3\sigma}{dk^3} = \langle n \rangle \, \sigma_{\text{ann}} \left|\Phi_\pi(k)\right|^2 \tag{5}$$

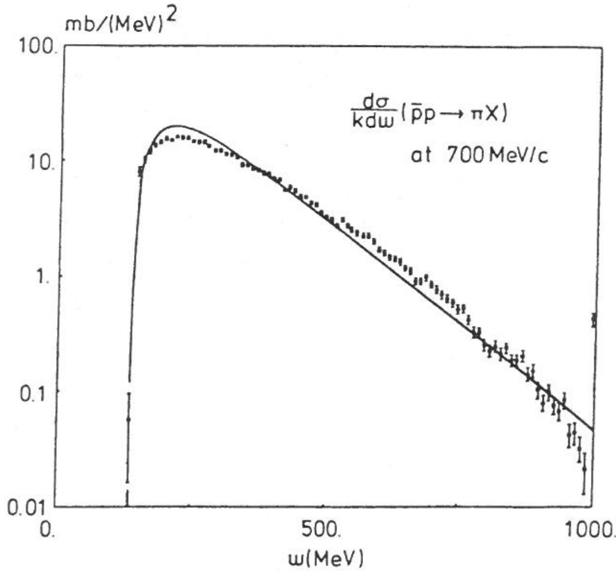

Figure 2
The pion energy distribution from Ref. [4]. This line is from Eq. (4).

where as before $\langle n \rangle = 5$ and $\sigma_{\text{ann}} \simeq 75$ mb for $p_{\text{lab}} \sim 700$ MeV/c. The Fourier transform of $\Phi_\pi(k)$ is assumed to reflect the $\bar{N}N$ annihilation region and Kimura and Saito[4] then postulated an imaginary potential $W(r) = \sigma_{\text{ann}} \tilde{\Phi}_\pi(r)$ which when added to a meson-exchange potential describes measured $d\sigma/d\Omega$ for $\bar{N}N$ elastic and charge-exchange just as good as Bryan and Phillips[9] or later potential models.

2.3 A Statistical Model with Some "Dynamics"

This is a third "minimal" annihilation model which describes experimental data very well. Vandermeulen[10] assumes that annihilation proceeds via a two-meson $(a + b)$ intermediate state and that the nearest meson-thresholds dominate.[11] The reaction is

$$\bar{N}N \to a + b \to \text{pions} + \text{kaons} \tag{6}$$

where a and b are S- and P-wave $\bar{q}q$ mesons. This is an old idea (early sixties) (based on baryon exchange, Fig. 3) which states that the dominant annihilation channels are those

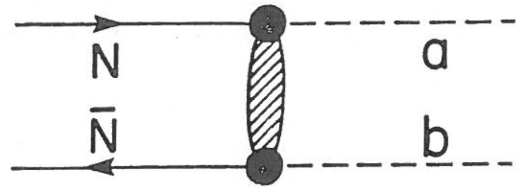

Figure 3
The old baryon exchange picture of $\bar{N}N$ annihilation into two mesons $a + b$.

which have the minimal energy/momentum transfer at the two vertices (NBa and NBb). This means a and b should be as heavy as possible for a given $\bar{N}N$ energy to get the least vertex suppression.

(Here I want to insert a short comment on the baryon-exchange calculations of, *e.g.* Moussallam[12] who studied $p\bar{p} \to \pi^+\pi^-$ and $p\bar{p} \to K^+K^-$. Due to the very high *energy* and momentum transfer in the $NN\pi$, $N\Delta\pi$, *etc.* vertices, his parametrization of these vertices are unrealistic because these vertices are timelike and, for example, intermediate states like $\pi\rho \to \pi$ can be on-shell). Vandermeulen[10] assumes in addition that the *different* $a + b$ channels leading to the *same* final states add incoherently. The branching ratio (or cross section) B for the process $p\bar{p} \to a + b$ is

$$B = C_{\bar{s}s}(2J_a + 1)(2J_b + 1)W_{ab}2^{-\delta_{ab}} k\, e^{-A\left(E_{\rm cm}^2 - (m_a + m_b)^2\right)^{1/2}} \quad . \tag{7}$$

Here the arbitrary constant, $C_{\bar{s}s}$ equals one for non-strange final states, J_a and J_b are the meson spins, W_{ab} *is the isospin-averaged weight* (Clebsch–Gordan coefficients), the factor 1/2 takes care of statistics for identical mesons ($a = b$) and k is the cm momentum. The exponential contains the threshold dominance postulate where m_a and m_b are the meson masses and A is a universal constant determined ($A = 1.2\,\mathrm{GeV}^{-1}$) by fitting the $p\bar{p} \to \pi^+\pi^-$ total cross section as a function of energy. For strange final states ($K\bar{K}$+pions) $C_{\bar{s}s}$ is adjusted to be 0.15. Vandermeulen's publication[10] shows a very impressive description of the total cross sections as a function of energy ($p_{\rm lab} < 3$ GeV/c) for the annihilation reactions, $p\bar{p} \to \pi^+\pi^-$, $\pi^+\pi^-\pi^0$, $2\pi^+2\pi^-$, $2\pi^+2\pi^-\pi^0$, $3\pi^+3\pi^-$, $3\pi^+3\pi^-\pi^0$, $4\pi^+4\pi^-$, $K^+K^-\pi^+\pi^-$, $K^+K^-\pi^+\pi^-\pi^0$, $K_sK^+\pi^+2\pi^-$(+c.c.), $K_sK^+\pi^+2\pi^-\pi^0$(+c.c.) and $2K_s2\pi^+2\pi^-\pi^0$. Only $p\bar{p} \to K^+K^-$ (and $\bar{p}p \to \pi^+\pi^-$ in an energy interval) is measured to be ***larger*** than this model description.

A few comments will summarize the above findings:

(i) The spin and isospin quantum numbers are averaged and then distributed statistically in most annihilation channels. This means these annihilation channels with several final mesons give little information about the underlying dynamics. Another consequence is that the imaginary part of optical models for $\bar{N}N$ elastic scattering should have no strong, explicit spin or isospin dependence.

(ii) Some "dynamics" seem to favor the nearest meson thresholds.

(iii) The strange final channels are suppressed relative to the non-strange ones ($C_{\bar{s}s} \simeq 1/7$). This suppression is supported by the theoretical model analysis of Dosch and Gromes[13] of $\bar{q}q$ creation in a "background" chomoelectric field (similar to Schwinger's e^+e^- creation in an electric field). This strange channel suppression (exception: $\bar{p}p \to K^+K^-$) says that $SU(3)$ breaking should be included into the general quark line rule analysis of Hartman *et al.*[14] and Genz *et al.*[15]

(iv) We have to concentrate on specific, two-meson final state channels which deviate strongly from the above expectations to learn something (we hope) about the underlying dynamics.

3. ANNIHILATION INTO TWO MESONS

In this last part I will limit myself to $\bar{p}p$ annihilation at rest where experimentally it is possible to specify the initial $\bar{N}N$ state via atomic X-ray coincidence measurements. However, as mentioned in Landua's presentation[16] there are some puzzling and possibly interesting differential cross section measurements[17] of $\bar{p}p \to K^+K^-$ and $\pi^+\pi^-$. In the $\bar{p}p$ atom the K_α-line has been observed,[18] and the inferred strong interaction energy-shift and width of the atomic $1S$-state are both in agreement with theoretical predictions of Bryan and Phillips[9] and later works.[19] The measured energy is[18]

$$\Delta E + i\Gamma/2 = [-0.70(15) + i0.80(20)]\,\text{keV} \quad . \tag{8}$$

3.1 The Atomic Wave Function

The protonium annihilation into $\pi^+\pi^-$ and K^+K^- has recently been measured by the ASTERIX collaboration[20] and together with earlier measurements of the $K^0\bar{K}^0$ channels[21,22] they find that the rate of $\bar{p}p \to \pi^+\pi^-$ is slightly larger from the $\bar{p}p$ P-state than from the S-state *whereas* the $K\bar{K}$ channel gets suppressed by a factor 5 from the protonium P-state relative to the S-state.[16] Due to symmetry requirements the $\pi^+\pi^-$ channel has isospin zero for decay from $\bar{p}p$ in the P-state and one for $\bar{p}p$ in the S-state. However, the $K\bar{K}$ is found to come from a mixture of both isospin one and zero in the atomic $\bar{p}p$ P-state[16,20] whereas only one isospin channel dominates the $\bar{p}p$ annihilation into $K\bar{K}$ from the atomic S-state.[20] Which one is a delicate question to be discussed next in connection with a critical presentation of a possible solution to the $\pi\rho$ "puzzle."

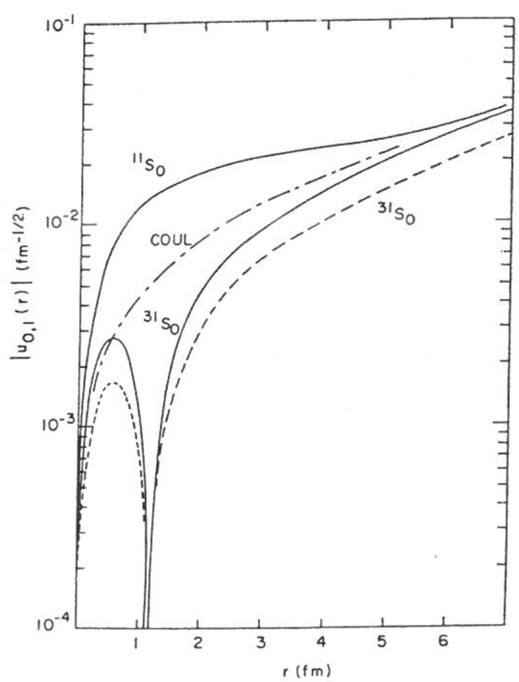

Figure 4

The $^1S_0$ wave function taken from Ref. [24] shows the two isospin states where $I = 1$ has a node. The dashed-dotted curve is the pure Coulombic wave function.

We know that the strong interaction affects both the S- and P-wave $\bar{p}p$ atomic states. We also know that the annihilation compares favorably with the electromagnetic K_α transition since $\Gamma\,(\text{radiation}) < 0.1\,\Gamma\,(\text{annihilation})$ in the lowest atomic P-state.[16] Furthermore, the $\bar{p}p$ atomic wave function at distances, r, of the order $1-2$ fm are modified by $\bar{p}p$ strong interactions.[23,24] In addition, we cannot exclude the possibility that the atomic $1S$-state is not the ground state of the $\bar{p}p$ system.[25] To illustrate this latter point, examine Fig. 7 in Buck, Dover and Richard[26] which shows the $\bar{N}N$ spectrum for several $\bar{N}N$ potential models all of which are G-parity transforms of NN potentials which describe NN scattering data. Please note that these energy spectra are sensitive to the arbitrary, strong interaction cut-offs used in these models. Furthermore, the annihilation is neglected when these spectra were calculated. In Fig. 4 is shown the $\bar{p}p\ {}^1S_0$ atomic wave functions taken from Dover, Richard and Zabek[24] (who neglect annihilation in calculating this atomic wave function). The isospin $I = 1$ wave function has a node at $r \simeq 1.2$ fm which makes it orthogonal to a $\bar{p}p$ strongly bound state of quantum numbers $I = 1$, $J^{PC} = 0^{-+}$ $({}^1S_0)$ generated by the $\bar{N}N$ meson-exchange potential in this calculation. The maximal $\bar{p}p$ annihilation strength is also around $r \simeq 1$ fm.[27] This means when we evaluate the overlap integral of this atomic wave function and the annihilation matrix element of for example $\bar{p}p \to \rho\pi$ we can get a very small number for $\bar{p}p({}^1S_0) \to \rho\pi$ when $I = 1$. However, as stated, the size of the strong interactions effect on the atomic wave function (or position of a possible node) is *extremely sensitive* to the arbitrary cut-off procedures used in the meson exchange $\bar{p}p$ potentials. Furthermore how annihilation processes will change these "predictions" are not really known. (A bound state pole at energy $E = E_0$ will move into the complex energy plane and the position of the complex bound state pole will move toward more positive $\mathrm{Re}\,(E)$ or more negative $\mathrm{Re}\,(E)$ on *a* Riemann-sheet depending on the position of the meson thresholds relative to E_0.[28])

After this exercise in "explaining" why the observed rate for $\bar{p}p({}^1S_0) \to \rho\pi << \bar{p}p({}^3S_1) \to \rho\pi$, we turn to the explanation of Maruyama *et al.*[29] They postulate that $\bar{p}p({}^3S_1) \to \rho\pi >> \bar{p}p({}^1S_0) \to \rho\pi$ because the strong interaction tensor force produce a sizeable ${}^3D_1$ atomic wave function component for $r \sim 1-2$ fm and they find,[29] for $\bar{p}p \to \rho\pi$, that $\Gamma({}^3D_1)/\Gamma({}^3S_1) \simeq 3$ (the overall atomic D-state probability is tiny). My statement above about the extreme sensitivity of the *arbitrary* strong interaction cut-off is not discussed by Maruyama *et al.*[29] (Their total $1S$ atomic width is also much larger than measured.) The problem with these type of "explanations" is the following. We know that the $\bar{N}N$ meson-exchange potential is in general more attractive for $I = 0$ than for $I = 1$. Kaufmann and Pilkuhn[23] found using basically $\bar{p}p$ one-pion-exchange that the atomic wave functions of the P-states $({}^1P_1,\ {}^3P_0,\ {}^3P_2)$ are mainly $I = 0$ whereas ${}^3P_1$ is mainly $I = 1$ for $r \le 1-2$ fm. However, ASTERIX find from the measured $K\bar{K}$[20] and $\pi\rho$[16] final states that the atomic P is *both* $I = 0$ and $I = 1$ indicating the meson-exchange analysis above has to be carefully re-examined. We should also keep in mind that the meson-exchange models did not predict $\bar{p}p$ polarization data[30] for $p_{\text{lab}} < 1$ GeV/c very well, so we should not be surprised if it fails to predict specific $\bar{N}N$ partial waves like the atomic states discussed above.

However, the meson-exchange models do offer a couple of speculative possibilities. If we have *no* strongly bound states then the $\bar{N}N$ atomic wave function will have no nodes and atomic $\bar{p}p$ wave function will be dominantly $I = 0$ at shorter distances for *many* $\bar{p}p$ quantum numbers. This seems to be in conflict with some ASTERIX measurements. Here a way out is that selective annihilation "dynamics" is important. The other possibility is that we do have strongly bound $\bar{N}N$ states (not experimentally observed) *i.e.* we will have nodes in the atomic wave functions and the positions of the node(s) will determine if an annihilation channel is suppressed or not. At this point it seems to me more fruitful to extend the nice analysis of Oades *et al.*[31] who used analytic extrapolations of "known" off-shell $\bar{N}N \to \pi^+\pi^-$ and $\bar{N}N \to K^+K^-$ helicity amplitudes and $\bar{p}p$ threshold arguments to predict what we can expect from the atomic measurements.

3.2 Quark Model "Dynamics"

In discussing this topic, the question I have in mind is: do the experimental data indicate something unusual that warrants dynamics on the quark level? Let me first present some critical comments on quark calculations relating to $\bar{N}N \to$ two mesons in which at least one $\bar{q}q$ pair has to annihilate. The so-called $^3P_0$ model, where $\bar{q}q$ annihilate with vacuum quantum numbers, does describe some aspects of meson-decay.[32] This model has been studied extensively in $\bar{N}N$ annihilation[29,33] (some disagreements in its prediction have to be settled[29,33].) The $^3P_0$ vertex is an effective operator and one has to add different quark diagrams leading to the same final states. The question is what are the relative strength and phases of these diagrams? And how does one account for the large internal energy-flow in these diagrams? To make progress with this model one has to make some "connection" to QCD in order to get some understanding for the approximations employed in these effective operator calculations. Several groups[13,29,33] have stated that the $\bar{q}q$ annihilating into a "flux-tube" simulates the $^3P_0$ model but these arguments have to be put on firmer ground.

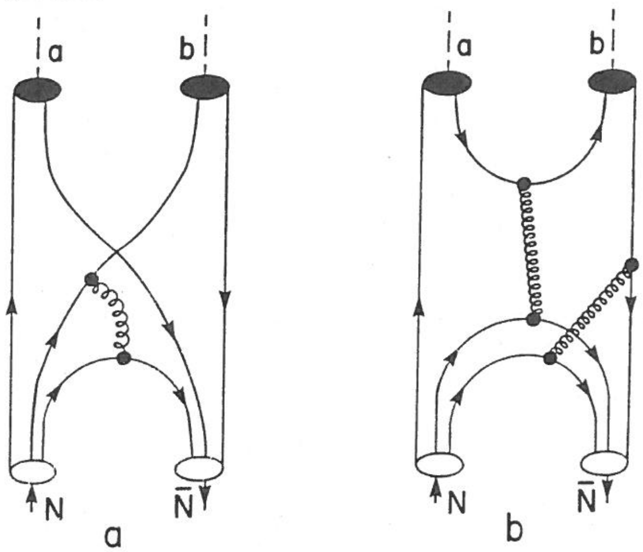

Figure 5
Two of the many effective quark-gluon diagrams for $\bar{N}N \to a + b$.

More obvious QCD related diagrams have been proposed by many.[29,33,34] Since the energy/momentum flow is relatively small (< 1 GeV) this is not a weak coupling limit ($\alpha_s \to 0$ for $Q^2 \to \infty$). On the average the annihilating $\bar{q}$ and q has a small relative momentum and since the gluons are confined (effectively heavy) the spatial part of the gluon propagator can be approximated by a delta-function. However, the gluon in Fig. 5a carries an energy $E_G \simeq 2E_0$ ($E_0 = m_Q = 330$ MeV in the non-relativistic quark model). This means the quark absorbing this energy becomes highly relativistic (and is far off-shell). If the final mesons are pointlike, Fig. 5a is zero (similar to pion absorption of a pointlike deuteron). Only a $\bar{q}q \to$ meson form factor makes this diagram non-zero and then this process probes the very high momentum components of this form factor. Furthermore the virtual gluon carries both transverse and longitudinal components. The non-relativistic quark model with $m_Q \simeq 330$ MeV makes sense in static spectroscopy where $m_Q = E_0 =$ energy of a confined quark. We know that a massless confined quark has a magnetic moment $\mu \simeq 1/2E_0$[35] ($E_0 \simeq m_Q$) and for the relativistic processes of $\bar{p}p$ annihilation it is more reasonable to assume that the dynamical quark masses are zero. The advantage is that the massless quark's *helicity* is conserved in quark-gluon interactions (no quark spin-flip). The confined quark wave function will be of the form $\begin{pmatrix} iF \\ G\underline{\sigma}\cdot\hat{\underline{r}} \end{pmatrix}$ where one automatically gets some extra spin dependence from the lower components not considered in non-relativistic calculations. In $\bar{N}N$ annihilation the focus *can be* on a few timelike, effective gluon diagrams since if one calculates the quark propagators in Fig. 6 in the closure approximation (sum over all states, but keep a fixed excited quark energy) then the extra gluons will be spacelike and contribute to the renormalization of the $N(\bar{N})$ wave functions only.[36] These calculations have of course to be applied to meson decay to test the approximations involved.

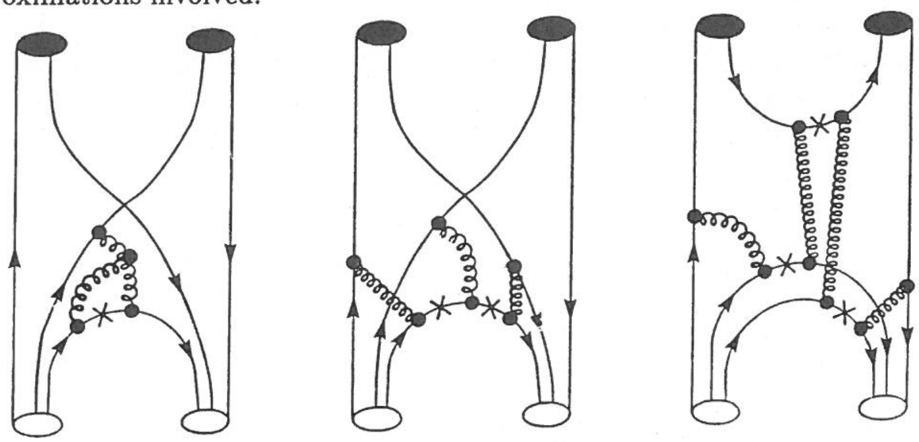

Figure 6
A few effective quark-gluon diagrams where the quark propagators with a cross is treated in the closure approximation.

4. CONCLUSIONS

As should be evident, most of the $\bar{N}N$ annihilation channels *do not* come from specific quark-gluon processes but resemble more the decay of some "hot plasma." Why the threshold dominance model[10] works so well is not understood. For $\bar{p}p$ annihilation at rest Kaufmann and Pilkuhn's "dynamical" isospin selection rules[23] due to meson exchange modifications of the $\bar{p}p$ atomic wave functions at short distances, seem to conflict with the *measured* isospin content of $\bar{p}p \to \bar{K}K$ and $\bar{p}p \to \pi\rho$ from atomic P-states[16] which indicate that we do not have strong initial state selection rules. We should always keep in mind that many of these $\bar{N}N$ mesons exchange models are extremely sensitive to the arbitrary cut-off procedures employed. However, a success of these meson-exchange ideas is the prediction of the strong energy-shift and width of the atomic $1S$-state.

To make progress in the effective quark-operator calculations in our search for some guiding principles in understanding annihilation we have to focus on $\bar{N}N \to$ two mesons and strive to derive and thereby improve our understanding of the $^3P_0$-model (connection to flux-tubes?) and the effective gluon diagrams. We can hope that in collaboration with our experimental colleagues we will reach some understanding of the underlying quark-gluon processes. Aspects of $\bar{p}p$ at rest decaying into one vector meson plus one pseudoscalar meson as well as $\bar{p}p \to \bar{K}K$ and $\pi\pi$ are puzzling and require some *careful* studies and in my opinion these reactions do reflect some quarkish aspects of the annihilation process.

5. DISCUSSION

Furui (Comment): If you adopt the so-called A3 model, *i.e.* $p\bar{p}$ annihilation into three mesons via two $q\bar{q}$ pair annihilations and two $q\bar{q}$ pair creations to form three mesons, I find that the S-wave annihilation into three pions with two pions correlated to make a ρ occurs only from $^3S_1$. The amplitude for $^1S_0$ channel vanishes due to the symmetry among the pions. The relation between the $^3P_0$ model and strong coupling lattice QCD theory is discussed by Isgur and Kokoski, *Phys. Rev.* **D35** (1987).

Answer: We have to understand better the theoretical foundation for the quark annihilation models before we calculate more diagrams.

Klempt (Question): You claim that the suppression of strangeness in $\bar{p}p$ annihilation is incompatible with $SU(3)$ analyses. But this suppression can also be described by the ratio of rearrangement and annihilation graphs, so I do not understand on which basis you refuse $SU(3)$ models.

Answer: In general, data says the strange channels are suppressed, but there are exceptions. I do not see how you can compare directly the *amplitudes* of "rearrangement" and "annihilation" graphs just based on symmetries.

REFERENCES

1. E. Fett *et al.*, *Nucl. Phys.* **B130** (1977) 1.
2. A. Pais, *Ann. Phys.* (NY) **9** (1960) 548.
3. C. J. Hamer, *Nuovo Cim.* **12A** (1972) 162.
4. M. Kimura and S. Saito, *Nucl. Phys.* **B178** (1981) 477.
5. J. H. Kim and H. Toki, *Prog. Theor. Phys.* **78** (1987) 616.
6. S. J. Orfandis and V. Rittenberg, *Nucl. Phys.* **B56** (1973) 561.
7. R. Hamatsu *et al. Nucl. Phys.* **B123** (1977) 189.
8. R. Hagedorn, in *Proc. Quark Matter 1984* Lecture Notes in Physics **221**, ed. K. Kajantie, (Springer, New York, 1985), p. 53.
9. R. A. Bryan and R. J. N. Phillips, *Nucl. Phys.* **B5** (1968) 201.
10. J. Vandermeulen, *Z. Phys.* **C37** (1988) 563.
11. G. F. Chew, in *Fundamental Problems in Elementary Particle Physics* (Wiley, New York, 1968), p. 72.
12. B. Moussallam, *Nucl. Phys.* **A429** (1984) 429.
13. H. G. Dosch and D. Gromes, *Phys. Rev.* **D33** (1986) 1378; *Z. Phys.* **C34** (1987) 139.
14. U. Hartmann, E. Klempt and J. Körner, *Phys. Lett.* **B155** (1985) 163; *Z. Phys.* **A331** (1988) 217; E. Klempt, *Z. Phys.* **A331** (1988) 211.
15. H. Genz, *Phys. Rev.* **D28** (1983) 1094; **D31** (1985) 1136 and in *Third Conf. on the Intersection Between Particle and Nuclear Physics* (Rockport, ME, 1988).
16. R. Landua, "What's New in $\bar{p}p$ Annihilation at Rest?," this conference.
17. T. Tanimori *et al.*, *Phys. Rev. Lett.* **55** (1985) 1835.
18. ASTERIX Collaboration, M. Zeigler *et al. Phys. Lett.* **B206** (1988) 151.

19. O. D. Dalkarov and V. M. Samoilov, *JETP Lett.* **16** (1972) 249; W. B. Kaufmann, *Phys. Rev.* **C19** (1979) 440; J. M. Richard and M. E. Sainio, *Phys. Lett.* **B110** (1982) 349; A. M. Green and S. Wycech, *Nucl. Phys.* **A377** (1982) 441; M. A. Albert *et al.*, *Phys. Rev.* **D27** (1983) 536; B. Moussallam, *Z. Phys.* **A325** (1986) 1.

20. ASTERIX Collaboration, M. Doser *et al.* CERN preprint CERN/EP 88-42 (1988).

21. ASTERIX Collaboration, S. Ahmad *et al.* *Proc. III LEAR Workshop* Tignes, France (Editions Frontières, Paris, 1985), p. 353.

22. R. Armenteros and B. French, *High Energy Physics*, Vol. IV, ed. E. H. S. Burhop (Academic Press, London, 1969), p. 284.

23. W. B. Kaufmann and H. Pilkuhn, *Phys. Rev.* **C17** (1978) 215.

24. C. B. Dover, J. M. Richard and M. C. Zabek, *Ann. Phys.* (NY) **130** (1980) 70.

25. O. D. Dalkarov, V. B. Mandelzweig and I. S. Shapiro, *JETP Lett.* **10** (1967) 257; I. S. Shapiro, *Phys. Rep.* **C35** (1978) 129.

26. W. W. Buck, C. B. Dover and J. M. Richard, *Ann. Phys.* (NY) **121** (1979) 47.

27. R. Tegen, T. Mizutani and F. Myhrer, *Phys. Rev.* **D32** (1985) 1672; F. Myhrer, in *AIP Conf. Proc.* **133** (1985), p. 96.

28. B. R. Karlsson and B. Kerbikov, *Nucl. Phys.* **B141** (1978) 241.

29. M. Maruyama, T. Gutsche, G. Strobel, A. Faessler and E. M. Henley, University Tübingen preprint (1988).

30. R. A. Kunne *et al.*, CERN preprint CERN-EP/88-22 (1988).

31. G. C. Oades, A. Miranda, L. Mandrup, A. S. Jensen and B. I. Deutch, *Nucl. Phys.* **A464** (1987) 538.

32. F. Close, *An Introduction to Quarks and Partons* (Academic Press, London, 1979).

33. C. B. Dover, P. M. Fishbane and S. Furui, *Phys. Rev. Lett.* **57** (1986) 1538; C. B. Dover, in *Third Conf. Intersection of Particle and Nuclear Physics* (Rockport, ME, 1988); M. Maruyama, S. Furui and A. Faessler, *Nucl. Phys.* **472** (1987) 643; S. Furui and A. Faessler, *Nucl. Phys.* **A468** (1987) 669; M. Maruyama and T. Ueda, *Prog. Theor. Phys.* **78** (1987) 841; A. M. Green and J. A. Niskanen, *Prog. Part. Nucl. Phys.* **18** (1987) 93.

34. E. M. Henley, T. Oka and J. Vergados, *Phys. Lett.* **166B** (1986) 274, E. M. Henley, L. Wilets and M. A. Alberg, University of Washington preprint 40048-40-NZ (1987); J. A. Niskanen and F. Myhrer, *Phys. Lett.* **157B** (1985) 247; M. Kohno and W. Weise, *Phys. Lett.* **B152** (1985) 303; *Nucl. Phys.* **A454** (1986) 429.

35. T. DeGrand, R. L. Jaffe, K. Johnson and J. Kiskis, *Phys. Rev.* **D12** (1975) 2060.

36. F. Myhrer and R. Tegen, *Phys. Lett.* **162B** (1985) 237 and in preparation.