Nucleon-antinucleon annihilation dynamics[†]

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NUCLEON-ANTINUCLEON ANNIHILATION DYNAMICS*

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ABSTRACT

Recent progress towards a microscopic understanding of nucleon-antinucleon $N\bar{N}$ annihilation is reviewed. We consider statistical models, the constraints imposed by $SU(2)$ or $SU(3)$ flavor symmetry and the spin-flavor-color dependence of effective operators for the creation/destruction of quark-antiquark ($Q\bar{Q}$) pairs in $N\bar{N}$ annihilation processes. The importance of dynamical selection rules is emphasized as a means of revealing the underlying reaction mechanism.

INTRODUCTION

Nucleon-antinucleon $N\bar{N}$ annihilation provides a testing ground for models of non-perturbative quantum chromodynamics (QCD). An *ab initio* QCD calculation of $N\bar{N}$ annihilation is beyond our reach at present, so a more phenomenological approach has been adopted, in which one tries various hypotheses for the dependence of the effective $Q\bar{Q}$ creation/destruction operator $\mathcal{O}$ on spin, flavor and color. Even at threshold, the $N\bar{N}$ system offers a rich ensemble of annihilation channels, whose relative branching ratios B offer significant constraints on dynamical models. Recently, the technique of γ -ray tagging at the LEAR facility at CERN has enabled one to prepare the $N\bar{N}$ atomic system in a state of definite orbital angular momentum L . The other quantum numbers such as spin S , isospin I , and charge parity C can also be filtered by comparing particular mesonic final states. The observed ratios B display a non-trivial dependence on L , suggesting dynamical content beyond simple statistical or flavor symmetry models. The key signatures of the underlying annihilation mechanism appear in the guise of dynamical selection rules (DSR), i.e., the suppression of transitions $N\bar{N} \rightarrow M_1 M_2$ which are allowed by the conservation of $\{J, \pi, C, G\}$. The emergence of DSR depends on the interplay between the color-spin-flavor structure of the effective $Q\bar{Q}$ operators $\mathcal{O}$ and the topology of the annihilation process. We focus on these features in the present note.

STATISTICAL MODELS

Given the sizable pion multiplicity $\langle N_\pi \rangle \approx 5$ in $N\bar{N}$ annihilation at threshold, it is natural to try to understand the data in terms of statistical and phase space factors, without any dynamical input. A recent analysis along these lines by Vandermeulen¹ was successful in understanding most features of the data, particularly at higher momenta where numerous partial waves contribute.

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In this approach, the branching ratios B (or cross sections) for the processes $p\bar{p} \rightarrow a + b$ are written as

$$B(p\bar{p} \rightarrow a + b) = C_{s\bar{s}} (2J_a + 1) (2J_b + 1) W_{ab} 2^{-\delta_{ab}} q e^{-A(s-s_{ab})^{1/2}} \quad (1)$$

where $J_{a,b}$ are the meson spins, W_{ab} is an isospin Clebsch-Gordon coefficient, $2^{-\delta_{ab}}$ is a Bose-Einstein factor, q is the c.m. meson momentum, $\sqrt{s}$ is the total c.m. energy, and $s_{ab} = (m_a + m_b)^2$. It is assumed that quasi-two-body (QTB) channels $a + b$ act as "doorway states." The subsequent decay of $\{a, b\}$ into pions builds up the observed value of $\langle N_\pi \rangle$. The exponential factor in Eq. (1) ensures that each channel $a + b$ peaks close to its threshold $\sqrt{s_{ab}}$; this is related to the constraints of multi-channel unitarity. The constant $A = 1.2 \text{ GeV}^{-1}$ is chosen to reproduce the energy dependence of the $p\bar{p} \rightarrow \pi^+\pi^-$ cross section. The constant $C_{s\bar{s}}$ is taken¹ to be 0.15 for modes containing $K\bar{K}$ pairs and unity for non-strange modes. The strangeness suppression factor $C_{s\bar{s}}$ expresses the fact that $K\bar{K} + n\pi$ modes represent only about 5% of the low energy annihilation cross section, whereas in a purely statistical picture they should constitute about 30%. This $s\bar{s}$ suppression is also observed² in other hadronic reactions. In the context of strong coupling QCD, $C_{s\bar{s}}$ can be understood³ in terms of the quantum tunneling of massive $s\bar{s}$ quark pairs, in analogy to the Schwinger formula for e^+e^- pair creation in an electric field. The need for a factor $C_{s\bar{s}}$ in meson decay has been detailed by Dosch and Gromes⁴.

The exponential factor in Eq. (1) may be viewed as an effective penetrability $P \approx \exp(-\tilde{R}q)$, where $\tilde{R} = 2A \approx 0.5 \text{ fm}$. In quark bag models for $N\bar{N} \rightarrow a + b$, Gaussian wave function overlaps⁵ yield

$$P \approx x^{2\ell_f} e^{-\beta x^2}, \quad (2)$$

where $x = qR$ and ℓ_f = meson orbital angular momentum. Although the functional form of (2) differs from that used in Eq. (1), the numerical values for typical $\ell_f = 0$ channels are rather similar if the bag radius $R \approx 0.5 \text{ fm}$. Note that $x > 1$ for all channels of interest, so the naive approximation $P \approx x^{2\ell_f}$, as often used in the past, is inadequate.

The statistical approach¹ describes most global features of $N\bar{N}$ annihilation, even at threshold, where one might expect the selectivity associated with particular $L \rightarrow \ell_f$ transitions to be important. Nevertheless, there are interesting discrepancies with the data, which allow some dynamical aspects to shine through the dull background of statistics and phase space. For instance, Eq. (1) contains no dependence on the quantum numbers of the $N\bar{N}$ initial state, so it does not allow for the possibility of dynamical selection rules which depend on the intrinsic spin S and isospin I . For example, experiment⁶ yields $B(\pi^\pm \rho^\mp) \approx 2B(\pi^0 \rho^0)$ for both $L = 0, 1$, implying that $\pi\rho$ production is dominated by the $^{2I+1, 2S+1}L_J = ^{13}\text{S}_1$ and $^{11}\text{P}_1$ channels (this is the " $\pi\rho$ puzzle"). Equation (1) predicts $B(\pi^\pm \rho^\mp) = 5B(\pi^0 \rho^0)$, which disagrees with experiment.

For the $\rho\rho$ channel, Eq. (1) implies $B(\rho^+\rho^-) = B(\rho^0\rho^0)$, while dynamical models, for instance Ref. 5, lead to $B(\rho^+\rho^-) \gg B(\rho^0\rho^0)$, because of the dominance of the $^{33}\text{S}_1 \rightarrow \rho^+\rho^-$ transition over $^{11}\text{S}_0 \rightarrow \rho^0\rho^0$. A similar situation occurs for the πA_2 channel: data⁹ imply $B(^{11}\text{S}_0 \rightarrow \pi A_2) \approx (3.6 \pm 0.9) B(^{33}\text{S}_1 \rightarrow \pi A_2)$, in disagreement with the statistical model¹.

Although the global fit requires $C_{33} \ll 1$, some ratios do not display the expected $(2J+1)$ factors. For instance, for $N\bar{N}$ at rest, we have

$$\frac{B(K^\pm K^{\mp})}{B(K^+ K^-)} \approx \begin{cases} 3.6 & \text{(Eq. (1))} \\ 0.8 - 1.4 & \text{(Exper, Ref. 7)} \end{cases} \quad (3)$$

This discrepancy may reflect the delicate isospin interference effects which occur in strange particle channels (no G -parity restriction). These are sensitive to the state dependent admixtures⁸ of $p\bar{p}$ and $n\bar{n}$ (and possible $\bar{\Delta}N \pm \Delta\bar{N}$, $\bar{\Delta}\Delta$) which occur at short distances in the atomic state.

In sum, although the statistical-phase space approach is consistent with many features of $N\bar{N}$ annihilation, it fails to account for the observation of dynamical selection rules, as seen for example in $N\bar{N} \rightarrow \pi\rho$, πA_2 . These are then the key to extracting non-trivial information from the data on $N\bar{N} \rightarrow$ mesons.

SU(2)/SU(3) FLAVOR SYMMETRY

The implications of flavor symmetry for $N\bar{N}$ annihilation have been examined by Genz¹⁰ and Klempt et al¹¹. These authors consider the quark line diagrams A2 and R2 of Fig. 1 corresponding to transitions to two meson final states. By considering non-strange modes only (SU(2)), Genz¹⁰ concludes that the planar topology A2 dominates. Klempt et al¹¹ include $K\bar{K}$, $\bar{K}K^*$ and $\bar{K}^*K^*$ modes in the analysis (SU(3)) and conclude that the rearrangement graph R2 is dominant, in apparent contradiction to Genz¹⁰. The argument of Genz is based largely on triangle inequalities for R2, such as

$$|B^{\frac{1}{2}}(\eta\eta') - \frac{1}{2}B^{\frac{1}{2}}(\pi^0\pi^0)|^2 \leq 2B(\pi^0\eta') \leq |B^{\frac{1}{2}}(\eta\eta') + \frac{1}{2}B^{\frac{1}{2}}(\pi^0\pi^0)|^2 \quad (4a)$$

$$|B^{\frac{1}{2}}(\rho^0\rho^0) - 2B^{\frac{1}{2}}(\rho^0\omega)|^2 \leq B(\omega\omega) \leq |B^{\frac{1}{2}}(\rho^0\rho^0) + 2B^{\frac{1}{2}}(\rho^0\omega)|^2 \quad (4b)$$

Using recent values $B(\pi^0\eta') \approx 0.5 \times 10^{-3}$ (Ref. 12) and $B(\pi^0\pi^0) \approx 0.2 - 0.25 \times 10^{-3}$ (Refs. 12, 13), Eq. (4a) could only be satisfied if $B(\eta\eta') \approx 10^{-3}$, but Tauscher¹⁴ gives a limit $B(\eta\eta') < 10^{-3}$. Similarly, since $B(\rho^0\rho^0)$ is very small (a value $(0.12 \pm 0.12) \times 10^{-2}$ is given by Diaz et al⁹), Eq. (4b) implies $B(\omega\omega) \approx 4B(\rho^0\omega)$. However, the data¹⁵ indicate $B(\omega\omega) \leq B(\rho^0\omega)$. Thus, the constraints (4a,b) from graph R2 appear to be violated. The situation is not so clear, however, since the A2 topology implies $B(\omega\omega) = B(\rho^0\rho^0)$, which is also not well satisfied by the data, since¹⁶ $B(\omega\omega) \approx (1.4 \pm 0.6) \times 10^{-2}$.

If one assumes that SU(3) symmetry is unbroken, we obtain for the planar A2 graph $B(^{13}\text{S}_1 \rightarrow \bar{K}K^* + \text{c.c.}) = B(^{13}\text{S}_1 \rightarrow \pi\rho)$. The observed $\bar{K}K^*/\pi\rho$

Fig. 1. Quark line diagrams for $N\bar{N}$ annihilation into two ($R2$, $A2$) or three ($R3$, $A3$) mesons. The dot stands for an effective operator for $Q\bar{Q}$ pair creation/destruction.

ratio is small. Since graph $R2$ produces $\pi\rho$ but not $\bar{K}K^*$, Klempt et al¹¹ conclude that $R2$ dominates. A problem with this argument is that it ignores a known mechanism of $SU(3)$ breaking, namely the $s\bar{s}$ suppression factor $C_{s\bar{s}}$ of Eq. (1). In addition, isospin interference is crucial for strange meson modes, i.e., the atomic system is not pure $p\bar{p}$ at short distances where annihilation occurs, as assumed in Ref. 11. The $\bar{n}n$ admixture is important in $S = 1$ channels, where the tensor potential plays an important role¹⁷ in providing a coherent $^{13}S_1 - ^{13}D_1$ mixture¹⁷ at short distances which diagonalizes the tensor operator S_{12} . Thus, the $\bar{K}K^*$ mode, for instance, is suppressed by $C_{s\bar{s}} \approx 1/6$ and also possibly by interference effects. And so it seems inappropriate to attribute the suppression of strange modes to the dominance of $R2$ over $A2$ processes.

For $SU(2)$, certain relations hold for any relative weight of $R2$ and $A2$. These are

$$B(^{13}S_1 \rightarrow \eta\omega) = B(^{11}S_0 \rightarrow \omega\omega) = B(^{11}S_0 \rightarrow \eta f) \quad (5a)$$

$$B(^{13}S_1 \rightarrow \pi\rho) = B(^{11}S_0 \rightarrow \pi A_2) \quad (5b)$$

The values $B(\eta\omega) \approx 1.3 \times 10^{-2}$ given by Tauscher¹⁴ and $B(\omega\omega) \approx 1.4 \times 10^{-2}$ of Bloch et al¹⁶ are consistent with Eq. (5a), while Eq. (5b) is also approximately satisfied. If $R2$ is absent, the above five B values become equal, and proportional to $|A_1^+ + A_2^+|^2$, using the $SU(3)$ invariant amplitudes of Klempt¹¹. If the experimental B values in Eq. (5) are corrected for a phase space factor, assumed¹¹ to be $q^{2\ell_f+1} \exp(-R^2 q^2/2)$, they are consistent with this prediction. Thus, if one restricts attention to non-strange modes ($SU(2)$) then the topology $A2$ is more or less consistent with the data.

DYNAMICAL MODELS

Thus far we have not assigned any operator structure to the dots in Fig. 1, and relied on general statistical or symmetry considerations. Although one may perform a fit to the data in order to obtain SU(3) amplitudes, this procedure gives no insight into the origin of dynamical selection rules, or the strong L dependence observed for certain modes (examples are πf , which is more prominent for $L = 1$ than for $L = 0$, and $K^0 \bar{K}^0$, which is strongly suppressed for $L = 1$). If the effective operator $\mathcal{O}$ corresponding to $Q\bar{Q}$ pair formation depends on spin and orbital angular momentum, the $N\bar{N} \rightarrow a + b$ amplitude will not factorize in general into orbital, spin, and flavor dependent terms (à la $0(3) \otimes \text{SU}(2)_o \otimes \text{SU}(3)_f$).

Recently, there has been a great deal of phenomenological work, in which one combines various prescriptions for the dots in Fig. 1, and tries to determine which topologies among $\{R2, A2, R3, A3\}$ are dominant. Major efforts have been devoted to this problem by the Helsinki¹⁸, Osaka¹⁹ and Tübingen²⁰ groups, as well as others^{5,21-24}.

The simplest dynamical approximation consists of using $R2$ only, with the dot replaced by an effective operator $\mathcal{O}$ with the quantum numbers of a single gluon²²⁻²⁴ (3S_1 , $I = 0$, color $\{8\}$). For meson decay, this approximation fails to account for either relative widths²⁵ or polarization properties²⁶ in the decays $B \rightarrow \omega\pi$ and $A_1 \rightarrow \rho\pi$. Also for $N\bar{N} \rightarrow a + b$, the one gluon approximation fails qualitatively: for example²², in $p\bar{p}$ annihilation at very low momentum, we predict the ratios $\eta\rho^0/\eta\omega \approx 0.1$ and $\eta\rho^0/\pi^0\rho^0 \approx 0.0075$, while the experimental values at rest are much larger, i.e., 0.6 and 0.55, respectively. The transitions $N\bar{N} (L=0) \rightarrow a + b$ ($\ell_f = 2$) vanish in one-gluon approximation, while πA_2 ($\ell_f = 2$) is observed to be one of the largest single two-body modes (3-5%).

One can try to repair the deficiencies of the effective one-gluon approximation by adding graph $A2$ of Fig. 1 with two-gluon exchanges^{19,23}. This can be made to give reasonable results for some ratios such as $K^+K^-/\pi^+\pi^-$, $K_L K_S/K^+K^-$, and $\pi^0\omega/\pi^+\pi^-$, but it fails for others such as $K^{*+}K^- + \text{c.c.}/K^+K^-$ (prediction 10 versus exper. $\sim 1 - 1.5$). The above-mentioned ratios $\eta\rho^0/\eta\omega$ and $\eta\rho^0/\pi^0\rho^0$ increase to about $1/4$ and $1/8$, respectively, but remain too small. In a global sense, the use of effective 3S_1 $Q\bar{Q}$ operators does not seem promising, even when graph $A2$ is included.

Meson and baryon resonance decay are well described^{27,28} phenomenologically using an effective $Q\bar{Q}$ operator with the quantum numbers of the vacuum (3P_0 , $I = 0$, color $\{1\}$). This approximation has been given a basis in the context of QCD in the strong coupling limit^{3,29}, where $Q\bar{Q}$ pairs are created by breaking a flux tube (formed as a normal (collective) mode of the gluon field). Since the 3P_0 model works so well for meson decays, it is reasonable to apply the

same technique to $N\bar{N}$ annihilation, although its success is by no means guaranteed. Indeed, the $N\bar{N}$ system is a "heavy meson", which shares the quantum numbers of the ordinary $Q\bar{Q}$ mesons, but differs in its internal quark composition and number of energetically allowed decay modes. Because of competition between different decay channels and topologies (Fig. 1), $N\bar{N}$ annihilation offers a richer dynamics than ordinary meson decay. As an example, $A_1 \rightarrow \rho\pi$ is the dominant decay, while $A_1 \rightarrow f\pi$ is a closed channel. In contrast, $N\bar{N}$ (3P_1), which has the same quantum numbers as the A_1 , can decay into both $\rho\pi$ and $f\pi$ (the latter dominates). This channel competition gives rise to dynamical selection rules, which are the most novel aspect of the $N\bar{N}$ annihilation process.

In Ref. 5, the 3P_0 model was applied to two-body decays, and it was concluded that the planar graph A_2 was more important than the rearrangement graph R_2 . Further support for this conclusion is provided by the recent AS-TERIX results³⁰

$$\frac{B(L=1 \rightarrow \pi^+\pi^-)}{B(L=0 \rightarrow \pi^+\pi^-)} \approx 1.44 \pm 0.16 \quad (6a)$$

$$\frac{B(L=1 \rightarrow K_s K_s)}{B(L=1 \rightarrow K^+ K^-)} \approx (1.46 \pm 0.3) \times 10^{-1} \quad (6b)$$

For A_2 , the 3P_0 model predicts³¹

$$\begin{aligned} \frac{B(L=1 \rightarrow \pi^+\pi^- (\ell_f=0))}{B(L=0 \rightarrow \pi^+\pi^- (\ell_f=1))} &\approx \left(\frac{45}{13}\right)^2 \frac{1}{(qR)^2} \approx \frac{3}{2} \\ \frac{B(L=1 \rightarrow K_s K_s)}{B(L=1 \rightarrow K^+ K^-)} &\approx \frac{1}{8} \end{aligned} \quad (7)$$

in excellent agreement with Eq. (6a,b). For R_2 , on the other hand, the transition $L=0 \rightarrow \pi^+\pi^- (\ell_f=1)$ is forbidden and $K\bar{K}$ modes also vanish.

The Tübingen group²⁰ has done the most extensive analysis of data on $N\bar{N} \rightarrow M_1 M_2$, $M_1 M_2 M_3$, comparing the topologies R_2 , A_2 , R_3 , A_3 of Fig. 1. They conclude that the use of 3P_0 $Q\bar{Q}$ operators together with the graphs A_2+A_3 offer the most successful phenomenological description. The rearrangement model R_3 yields too small a summed cross section and also leads to the selection rule $N\bar{N}(L=0) \not\rightarrow s\bar{s}$ for which there seems to be little evidence. In these calculations, initial state distortions are included and are important.

OUTLOOK

On the experimental side, one needs to obtain spin-separated rates for $L=0, 1$ $N\bar{N}$ states and measurement of modes with more than one neutral meson, in order to search for further dynamical selection rules (ex. ${}^3S_1 \rightarrow \rho^+\rho^- \rightarrow \pi^+\pi^-\pi^0\pi^0$). It would also be valuable to measure the polarization of mesons produced in $N\bar{N}$ annihilation; as for meson decays ($B \rightarrow \pi\omega$), this is a delicate

test of the reaction mechanism through the interference of amplitudes with different ℓ_f .

On the theoretical side, one needs a "guiding principle" in order to sort out the many somewhat confusing and *ad hoc* phenomenological prescriptions. The "flux tube" picture is successful for meson decays, but needs to be developed into a calculational scheme for $N\bar{N}$ annihilation. One should try to understand the relation of the 3P_0 vertex used in models of meson and $N\bar{N}$ decay to the $Q\bar{Q}$ and/or gluon condensates obtained from QCD sum rules. The $N\bar{N}$ system is also potentially rich as a means of producing new "hybrid" mesons containing dynamical gluon excitations, a possibility which merits theoretical and experimental study.

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