ANTIPROTON INDUCED REACTIONS*

J.-M. Richard
Institut des Sciences Nucléaires, Université Joseph Fourier-IN2P3-CNRS,
53, avenue des Martyrs, F-38026 Grenoble, France

February 5, 2022

Abstract

The various aspects of antiproton physics are shortly reviewed, and its relevance for the possible discovery of new particles and effects is pointed out. Then a survey of the nucleon-antinucleon interactions is given. In the nucleon-antinucleon annihilations, there is a big amount of experimental data that call for theoretical explanation. Importance of specific spin and isospin channels for our understanding of antiproton physics is stressed.

1 A survey of \( \bar{p} \) physics

1.1 Some references

The results obtained in experiments with low-energy antiprotons and their interpretation are summarized in the Proceedings of numerous Workshops and Symposia. One should first mention a series of “European Antiproton Symposia” held at Chexbres [1], Liblice [2], Stockholm [3], Barr [4], Brixen [5], Santiago [6], Durham [7], Thessaloniki [8] and finally Mainz [9]. The Low-Energy Antiproton (LEAR) machine at CERN, and the proposals and experiments there were discussed at the LEAR Workshops: Karlsruhe [10], Erice [11], Tignes [12], and Villars-sur-Ollon [13]. The two series merged into the Low-Energy Antiproton Physics (LEAP) conferences: Stockholm (LEAP 90) [14], Courmayeur (LEAP 92) [15]. The next one is planned to be held in Slovenia (LEAP 94).

More pedagogical (in principle) are the Schools held in Erice on specialized topics: fundamental symmetries [16], meson spectroscopy [17], nucleon-antinucleon (NN) and antinucleon-nucleus (NA) scattering [18], medium-energy physics with antiprotons [19].

*Lecture presented at the Indian-Summer School on Interaction in Hadronic Systems, Praha (The Czech Republic), 25–31 August 1993.
Among the review articles, one should quote at least the recent ones by C. Amsler and F. Myhrer \cite{20}, and by C.B. Dover et al. \cite{21}, which provide a very comprehensive survey of this field.

The data on $\Lambda N$ scattering and their interpretation are also discussed in the Proceedings of the Telluride \cite{22} and Bad Honnef \cite{23} workshops.

1.2 Antiproton beams

The first evidence for antiprotons was obtained at Berkeley in 1955. Shortly after this discovery, $\bar{p}$ cross sections were measured and antineutrons were produced in the charge-exchange reaction $p\bar{p} \rightarrow n\bar{n}$.

Secondary beams of antiprotons were also used at CERN, BNL, KEK,... The $\bar{p}$ are used immediately after being produced with the consequence of many impurities ($\pi^-, K^-, \ldots$) and a wide momentum distribution.

Stochastic cooling offers the possibility of storing the antiprotons that are produced in proton-nucleus collisions. This leads to beams of much higher intensity, 100% purity, and with a very sharp momentum resolution. These $\bar{p}$ facilities were designed for high energy physics ($W^\pm$ and $Z^0$ production, in particular), but applications at intermediate (charmonium physics) or low energies (antiprotonic atoms, symmetries, ...) were immediately considered as interesting by-products \cite{10}.

Antiproton beams of this type are now available at CERN and Fermilab. One can dream of even better $\bar{p}$ beams if a kaon factory is ever built, at Vancouver or elsewhere \cite{24, 25}. At the same time, some improvements are proposed at CERN as well, e.g. the Super LEAR project \cite{26, 27}.

1.3 CP violation

The violation of parity ($P$) conservation was suspected in the early 50’s and established in 1956. A simultaneous violation of the charge-conjugation ($C$) symmetry was discovered, but the combined operator $CP$ looked as being conserved. For instance, the neutrino is always left-handed, an obvious violation of $P$, but the antineutrino is always right-handed, an indication that $CP$ is a good symmetry.

In 1964, a violation of $CP$ symmetry was detected in the $K^0, \bar{K}^0$ system \cite{29}. If one defines $|\bar{K}^0\rangle = CP | K^0\rangle$, then the combinations

$$K_1^0 = \frac{1}{\sqrt{2}} (K^0 + \bar{K}^0)$$

$$K_2^0 = \frac{1}{\sqrt{2}} (K^0 - \bar{K}^0)$$

\footnote{See also, for instance, \cite{28}.}
correspond to the eigenvalues $CP = +1$ and $CP = -1$, respectively. If $CP$ is conserved, $K^0_1$ can decay into $\pi\pi$ (and also into $\pi\pi\pi$), and has to be identified with the short-living component $K^0_S$ of neutral K beams. On the other hand, $K^0_2$ cannot decay into $\pi\pi$ and coincides with the long-living component $K^0_L$.

However, the experiment of Christenson et al. in 1964 [29] has shown evidence for $K^0_0$ decaying into $\pi\pi$. This can be attributed either to the impurity of the eigenstates

$$K^0_S = K^0_1 + \epsilon K^0_2$$

$$K^0_L = -\epsilon K^0_1 + K^0_2$$

(2)

or to the $CP$-violating amplitude $K^0_2 \rightarrow \pi\pi$, measured by a parameter $\epsilon'$. The $CP$-violation is of primordial importance for fundamental physics and cosmology. The 1964 experiment has been repeated several times on both sides of Atlantic. The accuracy and statistics were significantly improved by using high energy beams.

The $CP$ experiment performed at LEAR [12, 13] offers an interesting alternative, with hopefully a better control of systematic errors, and an access to new fundamental quantities, besides $\epsilon$ and $\epsilon'$. The $K^0$ particles are tagged via the reaction

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

(3)

and the $\bar{K}^0$ by the conjugate reaction.

Low-energy kaons will also be produced in the DAΦNE facility at Frascati, in the reaction

$$e^+e^- \rightarrow \phi \rightarrow K\bar{K}.$$  

(4)

The study of $CP$ violation is too important to be restricted to a single system, namely $K^0$-$\bar{K}^0$. Similar effects should be observed in the beauty sector, i.e. $B^0$-$\bar{B}^0$ and $B_s$-$\bar{B}_s$ systems, where $B^0 = (\bar{b}d)$ and $B_s = (\bar{b}s)$ in terms of quarks. These B mesons can be produced either in very-high energy colliders or in B-factories that are to be constructed in near future.

There are also some speculations about $CP$-violation in hyperon decays [30,27]. If $CPT$ symmetry is exact (see next Section), then $\Lambda$ and $\bar{\Lambda}$ should have the same lifetime. Nevertheless, the angular correlation factors in $\Lambda \rightarrow p + \pi$ and in the charge conjugate decay $\bar{\Lambda} \rightarrow \bar{p} + \pi$ may differ. It is proposed to use the reactions

$$p\bar{p} \rightarrow \Lambda + \bar{\Lambda}, \quad \Lambda \rightarrow p\pi, \quad \bar{\Lambda} \rightarrow \bar{p}\pi$$

(5)

slightly above the threshold, and to compare the decays of $\Lambda$ and $\bar{\Lambda}$ in similar kinematical conditions. More precisely, the observable one aims to measure in experiment is

$$A \propto \vec{P}(\Lambda) \cdot (\vec{q}(p) \times \vec{q}(\pi))$$

(6)

defined in the rest frame of $\Lambda$, and its counterpart for $\bar{\Lambda}$. It involves the momenta $\vec{q}$ of the decay product and the polarization of the $\Lambda$. Thanks to the efforts of
the PS 185 collaboration at LEAR\textsuperscript{\textdagger}, we know that $\langle \vec{P}(\Lambda) \rangle$ is sizeable in a wide angular range.

It is believed that $CP$ violating effects could be even more pronounced in the $\Xi^-\Xi^-$ system, i.e. for baryons with two units of strangeness. One should first measure whether or not the production reaction

$$p + \bar{p} \rightarrow \Xi + \Xi$$

provides the hyperons with an important polarization. This question is interesting by itself for our understanding the dynamics of strangeness exchange reactions, as we shall see in Subsection 2.7.

### 1.4 CPT tests

$CPT$ symmetry has not been seriously questioned so far. It implies for the inertial masses and magnetic moments

$$m(\bar{p}) = m(p), \quad \mu(\bar{p}) = -\mu(p) \quad (8)$$

Accurate measurements have been performed in Penning trap experiments\textsuperscript{32} and are presently done using a Cyclotron trap [11–13]. The goal is

$$\frac{\Delta m}{m} = \left| \frac{m(\bar{p}) - m(p)}{m(p)} \right| \leq 10^{-8} \quad (9)$$

but lower limits could perhaps be reached. Earlier bounds on $\Delta m/m$ were obtained from antiprotonic atoms by observing the energies of the transitions between the high orbits, where strong interaction effects are negligible. These experiments with antiprotonic atoms give access to $\mu(\bar{p})$, as well.

Note that the equality of the inertial masses also holds for the imaginary parts, i.e. for the lifetimes. Proton decay experiments have provided bounds of the order of

$$\tau(p) \gtrsim 10^{31–33} \text{ years.} \quad (10)$$

(The value depends a little on whether one believes that $e^+\pi^0$ should be the dominant decay mode, if any.)

As already mentioned, the early antiproton experiments used antiprotons immediately after their production. With stochastic cooling, $\bar{p}$ are stored for several days, so that

$$\tau(\bar{p}) \gtrsim \text{several days.} \quad (11)$$

In Penning traps, one routinely stores electrons for months. Thus, reaching a limit $\tau(\bar{p}) \gtrsim 1$ month or 1 year seems feasible, if needed.

\textsuperscript{32}See the contributions of the PS185 collaboration in [31, 13–15].
1.5 Gravity experiments

If CPT symmetry holds, an “anti-Earth” would attract an antiproton with the same strength as Earth attracts a proton. The problem of antimatter gravity is whether Earth attracts protons and antiprotons at the same rate.

As pointed out e.g. by Hughes[3], inertial mass measurements provide an indirect answer to that question. Imagine the extreme situation where \( \bar{p} \) does not feel gravity, i.e., \( m_g(\bar{p}) = 0 \). From the equivalence principle, this 100% difference in \( m_g \) implies a relative difference

\[
\delta \nu = \frac{MG}{Rc^2} \approx 10^{-9}
\]

between the eigenfrequencies of \( p \) and \( \bar{p} \) in the same electromagnetic device. Here \( M \) and \( R \) are the mass and the radius of Earth. In other words, a \( 10^{-15} \) measurement of the \( \bar{p} \) inertial mass would test its gravitational mass to \( 10^{-9} \).

A direct measurement of \( m_g(\bar{p}) \) is, however, desirable. An experiment is planned at LEAR, by a team from Los Alamos [11–13]. They are presently testing their equipment by launching protons.

1.6 Very cold antiprotons

The measurement of the inertial and gravitational masses \( m(\bar{p}) \) and \( m_g(\bar{p}) \) requires slowing down the antiprotons extracted from LEAR when the machine is operating at its lowest momentum.

In recent years, some other applications of very low energy antiprotons have been proposed. In particular, one could combine \( \bar{p} \) and \( e^+ \) to form antihydrogen atoms. It is probably easier to measure the gravity of the neutral \( \bar{H} = (\bar{p}e^+) \) than that of the charged \( \bar{p} \). One could also measure with a very high accuracy the frequency of some electromagnetic transitions in \( \bar{H} \), and thus perform a sensitive test of matter-antimatter symmetry.

Protonium \((p\bar{p})\) and antihydrogen \((\bar{p}e^+)\) are the first examples coming to our mind when considering atomic physics with antiprotons. Some other configurations are stable, as far as one keeps the Coulomb interaction only and neglects annihilation and strong interactions. In the positron sector, one also knows the positronium ion \((e^+e^-)\), the positronium hybride \((pe^-e^+e^-)\) etc., which cannot undergo spontaneous dissociation. With antiprotons, one expects \((pp\bar{p})\), \((pppe^-)\) or \((ppp\bar{p})\), for instance, to be stable [34].

There may already be some indication for metastable exotic configurations. When one studies annihilation at rest, there are events with more time than expected between \( \bar{p} \) capture and its annihilation[4]. In the current picture, the \( \bar{p} \) is captured in some high orbit and quickly decays toward low-lying states (see Section 2.8.), while the electrons either remain outside or are ejected by Auger emissions.

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3See, for instance, [33] and references therein.
4See [35] and references therein to earlier works.
To explain the events with delayed annihilation, it is suggested that in rare circumstances the $\bar{p}$ reaches alternative intermediate orbits, where it remains trapped for some time. There are already some calculations of $A\bar{p}e^-$ system ($A =$ nucleus), indicating that the $m$-th radial excitation in the $n$-th Born-Oppenheimer potential, with typically $m = 89$ and $n = 7$, corresponds to $\bar{p}$ well localized outside the peak of the $e^-$ distribution, and having small probability of decaying into lower $(m,n)$ states. Other metastable orbits involve orbital instead of radial excitations. This new atomic spectroscopy remains to be studied in detail.

1.7 The proton form factor

The proton form factor is usually studied in the reaction $e^- + p \to e^- + p$, which corresponds to the domain $t < 0$ of the Mandelstam variable $t = (\tilde{p}_f - \tilde{p}_i)^2$, where $\tilde{p}_i$ and $\tilde{p}_f$ are the four-momenta of the initial and final proton, respectively.

The reaction $e^+e^- \to p\bar{p}$ gives information on the $t > 4m^2$ domain of the form factor. It was seen in electron-positron colliders, with however little statistics.

The reversed reaction $p\bar{p} \to e^+e^-$ has been measured by the PS170 collaboration at LEAR, from the threshold to $t \approx 4.2 \text{ GeV}^2$. An interesting structure was seen. It might be related to structures observed in $p\bar{p}$ elastic scattering. The reaction $p\bar{p} \to e^+e^-$ can be considered as the ultimate form of annihilation, where all incoming quarks disappear. We shall come back in Subsection 3.4 on the relative importance of diagrams where all, some or none of the initial quarks annihilate.

1.8 Low energy strong interactions

Section 2 will be devoted to the $NN$ scattering, and to protonium, while the dynamics of annihilation will be discussed in Section 3. Unfortunately, we shall not have enough time to review the antiproton-nucleus physics ($\bar{p}A$). The first results of LEAR dealt with some $\bar{p}A$ differential cross sections, in elastic or inelastic channels. This motivated many theoretical investigations.

The $\bar{p}A$ annihilation was also measured in several experiments, and compared to the “elementary” $\bar{p}N$ annihilation, in search for new phenomena, e.g. the annihilation of a $\bar{p}$ on two nucleons simultaneously, excess of strangeness production, etc. [13–15].

Heavy hypernuclei have been produced by shooting some $\bar{p}$ on a Uranium target, and the lifetime of these hypernuclei has been measured and compared to that of the free $\Lambda$ [14, 15].
1.9 Charmonium spectroscopy

The direct formation of charmonium states in electron machines proceeds via the reaction \( e^+e^- \rightarrow c\bar{c} \) and is restricted to \( J^{PC} = 1^{--} \) levels since there is a virtual photon in the intermediate state. One easily gets the \( 3S_1 \) state. The notation is \( ^{2S+1}L_J \). Some \( ^3D_1 \) states, like \( \psi''(3.772) \) are also seen thanks to some S-D mixing at short \( c\bar{c} \) separation, due to tensor forces or coupling to decay channels.

These \( 1^{--} \) states give access to \( 3P_0, 3P_1 \) and \( 3P_2 \) states via the dominant \( E1 \) radiative transition. The \( M1 \) signal \( 3S_1 \rightarrow 1S_0 + \gamma \) is less clear. A wrong value for the mass of the \( \eta_c \) was published in the 70’s. The true \( \eta_c \) seems now established \( 117 \) MeV below the \( J/\psi \) \([36]\), but the present candidate for \( \eta_c' \) is far from being firmly established.

We note that there is no access to \( ^1P_1 \) and to \( D \) states in electron machines. Another problem is that the masses and widths are not accurately measured, due to the limitations on the energy resolution of electron beams and \( \gamma \)-ray detectors.

The alternative reaction \( p\bar{p} \rightarrow c\bar{c} \) was successfully used by the R704 collaboration at CERN \([12]\) and later in the E760 experiment at Fermilab \([37]\). The widths of the \( \chi_2(3P_2) \) and \( \chi_1(3P_1) \) have been measured with great accuracy. These widths are very important quantities in QCD, where \( c\bar{c} \) decay is described in terms of 2 or 3 intermediate gluons.

The \( ^1P_1 \) state has been seen in these experiments. One can now analyse the P-state multiplet in terms of a central potential \( V_C \) supplemented by spin corrections with spin-spin, spin-orbit, tensor, and other components.

\[
\delta V = V_{SS}\hat{\sigma}_1\cdot\hat{\sigma}_2 + V_{LS}\vec{L}\cdot\vec{S} + V_T S_{12} + \cdots
\]

If they are treated at first order, then

\[
\begin{align*}
M(^1P_1) &= M_0 - 3\langle V_{SS} \rangle, \\
M(^3P_0) &= M_0 + \langle V_{SS} \rangle - 2\langle V_{LS} \rangle - 4\langle V_T \rangle, \\
M(^3P_1) &= M_0 + \langle V_{SS} \rangle - \langle V_{LS} \rangle + 2\langle V_T \rangle, \\
M(^3P_2) &= M_0 + \langle V_{SS} \rangle + \langle V_{LS} \rangle - \frac{2}{5}\langle V_T \rangle.
\end{align*}
\]

Experimentally, the \( ^1P_1 \) almost coincides with the centre-of-gravity of the triplet, i.e.,

\[
\delta M = \frac{1}{9}[M(^3P_0) + 3M(^3P_1) + 5M(^3P_2)] - M(^1P_1)
\]

is very small. Since \( \delta M = 4\langle V_{SS} \rangle_{l=1} \) at first order, we conclude that the spin-spin potential does not act on P-waves. In contrast, a value

\[
\langle V_{SS} \rangle_{l=0} = \frac{1}{4}[M(J/\psi) - M(\eta_c)] \approx 30 \text{ MeV}
\]

is very small.
is observed in S-waves.

The short range character of $V_{SS}$ is confirmed in lattice QCD calculations [38]. It agrees with phenomenological models, where $V_{SS}$ is of Breit-Fermi type, i.e. essentially of zero range.

At this stage of the analysis, the actual value $\delta M = -1$ MeV cannot be taken too seriously. Second order contributions of $V_{LS}$ and $V_T$, or a small quadratic spin-orbit term in Eq. (13) would easily generate such a $\delta M$ and it would be premature to conclude that $\langle V_{SS} \rangle_{l=1} < 0$ supports higher-order corrections in $\alpha_s$ or other fancy effects. On the contrary, an analysis based on a non-perturbative account for spin forces would probably lead to a slightly positive value for $\langle V_{SS} \rangle_{l=1}$ [37].

1.10 High energy physics

As already said, the commissioning of improved $\bar{p}$ beams was motivated by high-energy particle physics. The $p\bar{p}$ colliders at CERN and Fermilab have produced many results of basic importance on $W^\pm$ and $Z^0$ bosons, jets, heavy quarks, etc.

The diffractive part of the interaction was studied as well. One observes a rise of the $p\bar{p}$ cross section as a function of the c.m. energy $\sqrt{s}$, as $\sigma \propto (\ln s)^2$, i.e. the maximal behaviour allowed by the Froissart-Martin bound (the constant in front of $(\ln s)^2$ is, however, far from saturation). A measurement of $pp$ cross-sections in this energy region $\sqrt{s} \approx 1$ TeV is badly needed. There are speculations that the total $pp$ cross section could become larger than the $p\bar{p}$ one at very high energy.

It is not sure, however, that $p\bar{p}$ collisions would ever be performed at future colliders such as LHC. At very high energy, $pp$ and $p\bar{p}$ have comparable (very small) cross-sections for rare events producing Higgs bosons or supersymmetric particles. The choice is dictated by intensity considerations and proton beams are much better than antiproton beams in this respect.

1.11 A broad field of physics

To summarize, there are many important investigations done with antiproton beams. One could also mention some aspects of atomic or solid-state physics: ionization, channeling, wake-riding electrons induced by $\bar{p}$ collisions, etc.

The present improvements on the Fermilab collider and associated detectors might well lead to the discovery of the top quark. Several charmonium states await discovery, and the spectroscopy of hybrids, glueballs, radial excitation, etc. near or above 2 GeV/$c^2$ requires $\bar{p}$ beams such as these of the SuperLEAR proposal [26–28]. In the low-energy sector, one would much benefit from polarized $\bar{p}$ beams, and several methods of polarizing antiprotons are presently under investigation. Finally, a source of very cold antiprotons is required for symmetry tests, gravity experiments, $\bar{p}$ chemistry and antihydrogen formation.
2 Nucleon-antinucleon interaction

2.1 The nucleon-nucleon forces

There are rather good models available for describing the nucleon-nucleon (NN) interaction at low energy: Paris, Bonn, Nijmegen potentials, etc. We refer to the lectures given by K. Holinde at this Summer School [39].

The long-range part (LR) is described in terms of mesons which are exchanged: $\pi$, $\pi\pi$ (including $\pi\pi$ resonances such as $\rho$, $\sigma$) etc., and in terms of excitation of resonances ($\Delta, N^*, \ldots$) in the intermediate states. There is a solid piece of conventional strong-interaction physics, with many connections to the physics of mesons and baryons.

In potential models used in nuclear-structure calculations, the short-range part (SR) of the NN interaction is treated phenomenologically. The meson-exchange potential is regularized at short distances by ad-hoc form factors and supplemented by an empirical core, whose parameters are adjusted to fit the NN data.

There are some attempts to understand the SR part of NN in terms of quarks. When two nucleons come close together, the Pauli principle and the chromomagnetic force start acting between the quarks. This explains semi-quantitatively the observed repulsion.

The situation is, however, far from being fully satisfactory. We have a theory for the LR and another one for the SR forces. We badly need a unified treatment. The Skyrmion model, for instance, is an attempt to re-express some aspects of QCD in a way that is compatible with the Yukawa model of meson exchanges. Perhaps more promising is the use of effective Lagrangians adjusted to reproduce low-energy data on pions and nucleons.

Even at the phenomenological level, one can hardly draw any conclusion from a particular model, with quarks and gluons on the one side, and mesons and baryons on the other. The data are very sensitive to the transition region, and there are too many ambiguities in designing a matching between LR and SR potentials.

2.2 The G-parity rule

If a meson (or a set of mesons) $m$ is exchanged between two nucleons, and thus contributes to nuclear forces, it can also be exchanged between a nucleon and an antinucleon. In QED, we know that since the photon has charge conjugation $C = -1$, the repulsive Coulomb potential between two electrons becomes attractive between $e^-$ and $e^+$. This is the “$C$-conjugation rule”. A similar result holds for strong interactions, since they are invariant under $C$. For instance,
the LR potential between p and p mediated by $\pi^0$ exchange is identical to the Yukawa potential between two protons, since the neutral pion has $C(\pi^0) = +1$.

There is, however, another symmetry of strong interaction, isospin. It is convenient to analyse the data with potentials $V(I = 0)$ and $V(I = 1)$ acting on isospin eigenstates rather than with the linearly dependent $V(p\bar{p})$, $V(p\bar{n})$, $V(n\bar{n})$ and $V(p\bar{p} \to n\bar{n})$. Isospin symmetry and the $C$-conjugation rule are conveniently combined in the "$G$-parity rule" which links the NN and $\bar{N}\bar{N}$ potentials in a given isospin state.

$$V^I(\text{NN}) = \sum_m V_m \quad \Rightarrow \quad V^I(\text{N}\bar{N}) = \sum_m G(m) V_m \quad (17)$$

where $G = C \exp(-i\pi I_2)$, as usual. In particular, the potentials mediated by pion and omega exchanges flip sign, since $G(\pi) = G(\omega) = -1$.

Note that this $G$-parity rule should not be confused with crossing symmetry. This latter principle states that the same analytic function $F(s, t)$ describes the reaction $a + b \to c + d$ and its crossed reaction $a + \bar{c} \to b + \bar{d}$. For $a = b = c = d = N$, one would, indeed, recover a relation between $\text{NN} \to \text{NN}$ and $\text{NN} \to \text{NN}$, but one would have to perform an analytic continuation in the Mandelstam variables, from $(s > 4m^2, t < 0)$ to $(s < 0, t > 4m^2)$. Such analytic continuation would be unreliable, unlike for instance the case of $\gamma e^- \to \gamma e^-$ and $e^+ e^- \to \gamma\gamma$ in QED, for which one has in hand an analytic expression that is exact, or exact to some order in the coupling constant.

In the $G$-parity rule, one compares $\text{NN}$ and $\bar{N}\bar{N}$ elastic reactions in the same kinematical conditions, and exploits the fact that they share the same crossed channel, namely $\bar{N}\bar{N} \to m \to \text{NN}$. This is illustrated in Fig. 1. There is no approximation. In particular, the rule holds for both a single pole, corresponding to a stable meson $m$, and for unstable resonances, like $\rho$ which consists of a pair of correlated pions.

$$\begin{array}{ccc}
\text{N} & \text{m} & \text{N} \\
\hline
\text{N} & \text{m} & \text{N} \\
\end{array} \quad \begin{array}{ccc}
\bar{N} & \text{m} & \bar{N} \\
\hline
\text{N} & \text{m} & \text{N} \\
\end{array}$$

Figure 1: Exchange of a meson state $m$ in NN (left) and $\bar{N}\bar{N}$ (right) interactions. They have in common the same $t$-channel reaction $\bar{N}\bar{N} \to m \to \text{NN}$.

As a first consequence of the $G$-parity rule, the $\bar{N}\bar{N}$ potential is, on the average, more attractive than the NN one. In the past, this led to speculations about "quasi-nuclear baryonia", i.e. bound states of N and $\bar{N}$ with a binding energy larger than for the deuteron. The situation concerning the AX and other baryonium candidates will be reviewed by C. Amsler [41].

The LR attraction is also crucial for understanding the observed cross-sections, as we shall see in Subsection 2.4. Another consequence of the $G$-parity
rule is a strong spin and isospin dependence of the LR forces. We shall discuss this point in Subsection 2.5.

2.3 Empirical optical models

Meson exchanges tentatively account for the long and medium range part of the $NN$ potential. At short distances, the interaction is dominated by annihilation. In Section 3 we shall discuss the slow progress in our theoretical understanding of annihilation.

For a phenomenological description of the whole $NN$ interaction, one does not need the detailed knowledge of all branching ratios. What essentially matters is the cumulated strength of the coupling to all mesonic channels. The situation is reminiscent of nuclear reactions in which many final states are accessible and one describes the distortion of the initial state by means of an optical potential. Ultimately, the optical potential can be derived from the microscopic dynamics. Nevertheless, it is simply deduced by fitting the experimental data in most cases.

It is well known from elementary scattering theory that a real potential always produces real phase shifts and thus a purely elastic cross section. A complex potential $V$ with $\text{Im} \, V < 0$ provides inelasticity.

Several optical potentials have been designed to reproduce the early $NN$ data [42–44]. They have in common a meson-exchange tail deduced by means of the $G$-parity rule from the current NN potentials (and regularized at short distances) and a very simple parametrization of the core, typically taken in the Wood-Saxon form [42–44]

$$V_{\text{core}}(r) = \frac{V_0}{1 + \exp[(r - R)/a]}.$$

Note that there is no reason to believe that $V_{\text{core}}$ should be local. For instance, simple quark models lead to separable forms [13, 14]. The parametrization (18) is dictated by simplicity. Even so, without spin or isospin dependence, the parameters $V_0$ (complex), $R$ and $a$ were not determined unambiguously: one can arrange either a strong $|V_0|$ and $R = 0$ [14], corresponding to a sharply decreasing potential, or a moderate $|V_0|$ and $R \approx 0.8$ fm [17], with a shoulder shape. (It is surprising that some authors were able to determine tens of parameters for the core on the basis of the same pre-LEAR data and even include a complicated spin-isospin dependence [18].)

The situation with $V(r)$ is a little similar to low-energy heavy-ion scattering: the long range is dominated by the Coulomb potential, and the inner part of the inter-ion potential is never seen; everything comes from the surface interaction. Here, one needs a strong absorption near $r \approx 0.8–1.0$ fm, and all possible fits give similar values of $\text{Im} \, V$ in this region [14].

The absorption range of 0.8–1.0 fm was a bit surprising: one would expect a very short range, as e.g. in the $e^+e^-$ annihilation. We shall come back to this point in Subsection 3.4.
Note that the range of absorption is clearly read off only for local models. If you introduce energy or angular momentum dependence in your model (a perfectly reasonable strategy with regards to the underlying microscopic mechanisms), you can increase the rate of absorption in high partial waves and mimic the amplitudes generated by local potentials with the size 0.8–1.0 fm. The wave functions obtained from both the local potentials and the non-local ones are rather similar, indeed.

2.4 Integrated cross-sections

The integrated cross-sections which have been measured in experiments and fitted in optical models are rather large. Their order of magnitude is 100 mb and the most striking features are:

i) The large ratio of inelastic to elastic cross-sections \( \sigma_{\text{inel}}/\sigma_{\text{el}} \approx 2 \). A simple black sphere would give \( \sigma_{\text{inel}}/\sigma_{\text{el}} \approx 1 \), typically. The departure is usually understood as an effect of LR attraction, which pulls out the wave function into the annihilation region [49].

ii) The smallness of the integrated charge-exchange cross-section \( \sigma_{\text{ce}} \). In a pure one-pion-exchange model, \( \sigma_{\text{ce}} \) would be comparable to \( \sigma_{\text{el}} \) or even larger, since the isospin Clebsch-Gordan coefficients are more favourable. When absorption is considered, both \( I = 0 \) and \( I = 1 \) amplitudes are suppressed (\( I \) is the isospin in the direct channel) and nearly equal at short distances. The charge-exchange amplitude

\[
M_{\text{ce}} \propto M_{I=0} - M_{I=1}
\]

(19)

becomes extremely small in the central region. It was pointed out [44] that the smallness of \( \sigma_{\text{ce}} \) was the most constraining property of pre-LEAR data when adjusting the parameters of the optical models. The authors who insist on having a short-range absorption get a too large charge-exchange cross-section.

iii) Isospin \( I = 1 \) cross-sections are usually well reproduced by the same optical models that work well for \( \bar{p}p \) and charge-exchange cross-sections. An interesting study of the possible isospin dependence in the core region was carried out in Ref. [50]. Recently, the \( \bar{n} \) cross-sections were measured at very low energy [52] by the OBELIX collaboration, with rather surprising results, which contradict current potential models. One should wait, however, for the final analysis of this delicate experiment.

2.5 Differential cross-sections

The integrated cross-sections are dominated by the low partial waves, at least in the low energy region relevant to LEAR experiments. The role of higher-\( L \)
waves is better seen in angular distributions. Differential cross-sections have been measured in several experiments with following results:

i) Even at very low energy, the angular distributions are far from being flat. This means a dominance of P-waves and the other $L \leq 2$ partial waves.

ii) The differential cross-section for $p\bar{p} \rightarrow n\bar{n}$ exhibits a structure in the forward hemisphere. This structure is sensitive to the interference between $\pi$-exchange, $\rho$-exchange and absorption, which makes it interesting for testing the models. I had several discussions on this point with the late Helmut Poth, who was a pioneer in all aspects of LEAR: machine design, physics experiments, and also their theoretical interpretation.

iii) The Coulomb-nuclear interference in the forward hemisphere gives access to the so-called ”$\rho$ parameter” defined as

$$\rho(s) = \frac{\text{Re} M(s,t)}{\text{Im} M(s,t)} igg|_{t=0}. \quad (20)$$

There is a straightforward generalization for scattering of particles with spin, where more than one amplitude contribute. $\rho(s)$ exhibits rapid variations near threshold in the region $s \gtrsim 4m^2$. Determining whether there are genuine resonances would require further experimental investigations with very low energy antiprotons. Note that $\rho(s)$ is available at threshold, thanks to protonium experiments: from the Truemann formula that will be written in Subsection 2.8, $\rho(4m^2) = \frac{\text{Re} (\delta E)}{\text{Im} (\delta E)}$, where $\delta E$ is the complex energy shift of the 1S state of protonium, with respect to the pure Coulomb binding energy. $\rho(s)$ has also been measured at high energy. So one can use dispersion relations to try to understand the energy dependence [51]. Again, the sharp variations at low energy raise difficulties.

### 2.6 Spin forces and spin observables

We just recall that differential cross-sections are more sensitive to high partial waves than the integrated ones. Still, they are not sufficient to reconstruct the interaction and to test the validity of models in this way. One should also measure a sufficient number of spin observables to get some real insight into the dynamics.

Let us give some examples to illustrate this point. In atomic physics, the very precise measurements of fine and hyperfine structures provide unambiguous tests of the vector character of QED. We mentioned in Section 1 the efforts in improving the spectroscopy of charmonium and in comparing the spin dependence of the heavy-quark potential with QCD predictions. In both the nuclear physics and the low-energy hadron physics [39,52], the pseudoscalar character of the pion reflects the fundamental properties of the underlying theory and leads to clear consequences, such as the quadrupole deformation of deuteron.
It took many years to achieve a comprehensive experimental investigation of the NN interaction, with a delicate machinery of polarized beams and polarized targets, and an obstinate phenomenological analysis of the data, mostly by phase-shift analysis [53]. Clearly more experimental efforts should be devoted to study the $\overline{N}N$ interaction.

Current optical models can be used to make rough simulations of spin observables. The results are rather dramatic, with some parameters nearly saturating the limits allowed by unitarity [54,55]: spin transfers close to 100%, very large depolarization effect etc. Unfortunately, the spin dependence that is expected is more complicated than in the NN case. In the latter case, we have mostly spin-orbit contributions, so that polarization (or analysing power) is the first observable to look at.

In $N\overline{N}$ interaction, the tensor force is the dominant component [56]. This is a rank-2 operator in the language of specialists. This means that the best signatures are seen in observables where at least two spins are measured: spin correlation or spin transfer.

If one treats the tensor operator

$$S_{12} = 3(\vec{\sigma}_1 \cdot \hat{r})(\vec{\sigma}_2 \cdot \hat{r}) - (\vec{\sigma}_1 \cdot \vec{\sigma}_2)$$

in DWBA approximation starting from the wave function generated by the central potential, there is no polarization. The observed polarization [57,58] might well be due to tensor forces acting at second order or beyond it [55]. This is why it is crucial to measure quantities that are sensitive to tensor forces at first order.

This dominance of the tensor force can be understood on rather general grounds. In $N\overline{N}$, the $\pi$- and $\rho$-exchange contributions to the tensor forces add up coherently. In NN, they tend to cancel each other, while other coherent effects show up in the spin-orbit component [56]. We note that in theoretical predictions, the most striking effects are expected in the charge-exchange reaction $p\overline{p} \rightarrow n\overline{n}$. It guided the choice of the PS199 collaboration [12,57,58].

Once a decent amount of spin observables will be accumulated, one will be able to get rid of the various components of $N\overline{N}$ forces by fitting the data with an optical model that contains detailed spin and isospin dependence. Some recent attempts in that direction [59] are clearly premature. One cannot fix both spin-orbit and tensor forces from polarization data alone. Similarly, if one returns to charmonium spectroscopy and looks at Eq. (14), one cannot determine $\langle V_{LS} \rangle$ and $\langle V_T \rangle$ without knowing both $^3P_2$-$^3P_0$ and $^3P_1$-$^3P_0$ splittings. This is not a profound physics statement, simply the same counting of the degrees of freedom. It is also clear that one would need many more data to determine the N$\overline{N}$ phase-shifts. One may exhibit a particular set of phase-shifts [60] that is compatible with some of the available data, but there are certainly many other possible solutions.
Finally a word on integrated spin observables. The quantity

$$\Delta \sigma_T = \sigma(\uparrow \uparrow) - \sigma(\uparrow \downarrow)$$

is crucial for one of the proposed tools for polarizing antiprotons [13]. It roughly consists of letting a $\bar{p}$ beam pass many times on a transversally polarized proton target, to eliminate the $\bar{p}$ with the transverse spin corresponding to the largest cross-section. Unfortunately, the theoretical expectations on $\Delta \sigma_T$ are not too encouraging. It might be, however, that

$$\Delta \sigma_L = \sigma(\rightarrow \rightarrow) - \sigma(\rightarrow \leftarrow)$$

is slightly larger. Since only transverse polarization can safely travel in accelerator rings, this would mean some additional magnets in the device to flip the spin of $\bar{p}$ before and after the filtering target.

A possible way of producing polarized antinucleons is provided by the charge-exchange reaction, with a (unpolarized) $\bar{p}$ beam shooting protons with longitudinal polarization [54]. The $\bar{n}$ produced in the forward angles are expected to be highly polarized. This is essentially an effect of $\pi$ exchange and thus seems rather safe.

### 2.7 Strangeness-exchange reactions

We have seen in the previous subsections that the $p\bar{p} \rightarrow n\bar{n}$ charge-exchange cross-section deserves a special attention when studying $NN$ scattering. It is strongly suppressed at short distances, but enjoys coherent contributions from meson exchanges at larger distances.

A natural generalization is flavour exchange. The reaction $p\bar{p} \rightarrow \Lambda\bar{\Lambda}$ was carefully measured by the PS185 collaboration at CERN [31,13–15]. (An outstanding member of this collaboration was Nikolaus Hamann, who died recently, and also was the driving force of the JETSET experiment that will be mentioned when discussing annihilation.) PS185 has found very striking results, which motivated tens of theoretical papers. Attention is paid in particular to

i) a possible structure in the cross-section close above the threshold

ii) the very large $P$-wave, and even $L \leq 2$ contributions even at very low values of $\Lambda$ and $\bar{\Lambda}$ momenta in the c.m.s.

iii) the rather large polarization of $\Lambda$ and $\bar{\Lambda}$ in final state. This is a striking feature that hyperons produced at various energies in various reactions have similar polarizations.

iv) the complete suppression of the singlet fraction. Without spin-dependent forces, one would expect a fraction $F_0 = 1/4$ of spin $S = 0$ events and $F_1 = 3/4$ for $S = 1$ if both the $\bar{p}$ beam and the $p$ target are unpolarized. The observed $F_0 \approx 0$ is very puzzling.
The above results stimulated a continuation of the program to study $\Lambda \Sigma + \text{c.c.}$ and $\Sigma \Sigma$ production. The production of strangeness $-2$ and $-3$ hyperons or even of charmed baryons could be studied with higher-energy machines. It would also be rather interesting to push further the study of the spin dependence by producing hyperons on a polarized proton target and analyse the correlations between the spin of the proton and that of $\Lambda$ and $\bar{\Lambda}$.

As for the theoretical predictions, there are essentially two classes of models [61–66] which are schematically represented in Fig. 2 (see also [67]):

- a) nuclear-physics type of models, where the transition is described in terms of $K, K^*, \ldots$ exchanges, with proper account for initial and final state interaction, in particular the strong absorption.
- b) quark models, with typically, for $p\bar{p} \to \Lambda \bar{\Lambda}$, a $u\bar{u}$ pair annihilated in the initial state and an $s\bar{s}$ pair created in the final state.

The models give similar results, illustrating once more that both hadron and quark basis might be used to describe low-energy physics. The choice is a matter of convenience, but one should avoid double counting when trying to combine the two pictures.

The complete suppression of the spin-singlet fraction has, however, never been understood in simple terms. See, for instance, Ref. [68] for a recent discussion.

### 2.8 Protonium

Most spectroscopy experiments at LEAR, so far, use annihilation at rest, i.e. the initial state is a protonium or another antiprotonic atom. It is thus important to understand the physics of these atoms to analyse the results of annihilation experiments. Antiprotonic atoms are interesting by themselves since they provide information on the $NN$ interaction at zero relative energy. We shall restrict here to protonium and say a word on antiprotonic deuterium in the next subsection. However, some of the statements would also hold for more complicated antiprotonic atoms or other exotic atoms.
When a low-energy antiproton is sent in a hydrogen target, it is slowed down by electromagnetic interaction. It is then captured by the proton, while the electron is expelled by Auger emission. The most probable state corresponds to a radius that matches the Bohr radius of the initial electron. This gives a principal quantum number \( n = \sqrt{\frac{2m_e}{m}} = 30 \). The protonium then decays via radiative transitions, populating circular state (with maximal angular momentum \( L = n - 1 \)) preferentially.

In a very dilute gas, the scenario is rather simple

i) as long as \( L \geq 2 \), strong interactions are negligible (with a marginal exception for \(^3\text{D}_1\) states which have a small mixing with \(^3\text{S}_1\)). The spectroscopy is dominated by QED. Note that pp radius being much smaller than that of \( \text{pe}^- \), the average electric field is much larger and vacuum polarization effects more important.

ii) for \( L = 1 \) (2P level), the real part of the energy shift is still very small in comparison with the Coulomb energy, but there is already a contribution of strong interactions to fine and hyperfine splittings \([69, 47]\). A measurement is planned at CERN \([70]\). If successful, it will provide an estimate of spin-spin, spin-orbit and tensor forces at rest. The hadronic width for \( L = 1 \) exceeds by a large factor (which depends on the particular \(^{2S+1}\text{P}_J\) state one considers) the electromagnetic width for \( 2\text{P} \rightarrow 1\text{S} \) transition. This means 2P states mostly decay by annihilation.

iii) a few \( L = 0 \) states (1S) levels are formed. Here the real and imaginary shifts are large, of the order of 1 keV, to be compared to the pure Coulomb energy 12.5 keV. These states of course cannot do anything but annihilate. The hyperfine structure, i.e. the separation between \(^3\text{S}_1\) and \(^1\text{S}_0\), has not been clearly seen.

When one increases the pressure in the gaseous target, or uses a liquid target, the cascade processes become more involved. A protonium, thanks to its small radius, can travel inside an ordinary hydrogen atom and experiences Stark mixing in the corresponding electric field \([71]\). This means that high-\( L \) states are mixed with high-\( n \) S-states, from which annihilation can take place. In short, the ratio of S-wave to P-wave annihilation is very sensitive to the pressure of the hydrogen target.

The quantum mechanics of protonium is easily formulated, but requires a lot of care when one carries out the calculations. The basic equation is

\[
u''(r) - \frac{L(L+1)}{r^2}u(r) + m(E - V)u(r) = 0,
\]

where \( u(r) \) is the reduced radial wave function, with \( u(0) = 0 \), and the Coulomb potential \( V^c \) is replaced by the total potential \( V = V^c + V^n \) with a (complex) nuclear piece \( V^n \). For \( L = 0 \), one typically gets shifts of the order of 1 keV for
the ground state (with principal quantum number \( n = 1 \)), small compared to
the Bohr energy \( E^c = -12.49 \text{ keV} \). This does not mean, however, that \( V^n \)
can be treated as a perturbation. A first order estimate

\[
\delta E = E - E^c = \int_0^\infty u^c(r)^2 V^n dr
\]

(25)

would overestimate \(|\delta E|\) by orders of magnitude. The ordinary expansion
in powers of the additional potential is not applicable here. What is appropriate
is “radius perturbation theory” \([72, 73]\), where the expansion parameter is the
ratio \( a/R_0 \) of the scattering length \( a \) in the nuclear potential to the Bohr radius
of the atom. At first order, one gets the famous Trueman formula, which reads

\[
\delta E = \frac{4\pi m}{4\pi} |\Psi^c(0)|^2 \frac{a}{R_0}
\]

(26)

for S-waves. There is an analogue for P-waves where the first derivative of
the radial wave function, \( du^c(r)/dr|_{r=0} = \sqrt{4\pi}\Psi^c(0) \) is replaced by the second
derivative, and \( a \) by the scattering volume.

There is some confusion in the literature about the validity of the Trueman
formula, with a tentative clarification\(^7\). We first note a lack of unified conven-
tions for defining \( a \), \( E \) and \( \delta E \). Secondly, there are claims for the Trueman
formula being inaccurate. The problem comes in fact from the Coulomb cor-
rections to the scattering length, which are often omitted or badly computed.
Once the Coulomb corrected scattering length is properly estimated, the True-
man formula turns out to be very precise. The only exception is the situation
where \( a \) is large, i.e., where the nuclear potential supports a bound state or a
resonance very close to the threshold \([49]\).

The physics is actually more involved than suggested in Eq. (24). The
Coulomb potential acts between \( p \) and \( \bar{p} \) only, while the nuclear piece is di-
agonal in the isospin basis. One has to consider a two-component wave function

\[
\Psi = \frac{u(r)}{r} |pp\rangle + \frac{w(r)}{r} |nn\rangle
\]

(27)

resulting into coupled equations, with the neutron-to-proton mass difference
taken into account \([38]\). In the natural-parity sector, there is also some orbital
mixing \((L = J - 1 \leftrightarrow L = J + 1)\), and 4 coupled equations altogether. At
this point, the Trueman formula still holds, provided one accounts for these
couplings when computing the scattering length \( a \).

Another method to estimate \( \delta E \) consists of solving numerically the coupled
radial equations. This is a little difficult if only \( \delta E \) is needed, but it has
the advantage of providing the wave functions. The results are rather dramatic.
They were perhaps not paid enough attention when first obtained by Kaufmann

\(^7\)For a review, see [74].
and Pilkuhn [8] and confirmed in subsequent calculations [47, 75]. Nowadays they are taken into account in most serious analyses of annihilation experiments. In particular

i) protonium is far from being pure $p\bar{p}$ in the annihilation region. There is a copious $n\bar{n}$ mixing, and when one projects out on the isospin eigenstates, one finds one of the isospin components dominating, $I = 0$ or $I = 1$, depending on the partial wave. For instance, the decay of the scalar $^3P_0$ is mostly $I = 0$.

ii) there is also some S-D or P-F mixing in states of natural parity.

iii) the various P-states ($^1P_1$, $^3P_0$, $^3P_1$, $^3P_2$) have very different hadronic widths.

iv) there are some oscillations in the density probabilities and this might influence the branching ratios [76].

2.9 Antiprotonic deuterium

Antiprotons have been stopped on a variety of nuclear targets. For studying $\bar{p}A$ atoms with large $A$ nuclei, one usually derives an optical potential, by folding the elementary $\bar{p}p$ and $\bar{p}n$ amplitudes with empirical nuclear wave functions [77,22,23]. The case of antiprotonic deuterium ($\bar{p}d$) is somewhat intermediate [78–80].

If one only aims at estimating the energy shift $\delta E$, then simple approximations seem to work quite well. However, none of the simple methods provides reliable wave functions in the annihilation region. One would have to perform a detailed 3-body calculation.

These wave functions should be useful to analyse in detail $\bar{p}d$ annihilation at rest, and to compare it with $\bar{p}p$ and $\bar{p}n$ annihilations. It is possible that the spin, isospin and angular momentum content of each $NN$ sub-system is modified by the presence of the other nucleon.

3 Annihilation

3.1 Why annihilation?

We have at least two good reasons for studying annihilation extensively. At first, annihilation is a fascinating process, where matter undergoes a kind of phase transition, from a baryonic structure to mesonic states. It was first (and it is still, by some authors) thought in analogy with $e^+e^-$ annihilation in QED. The nucleons play the role of the electrons, and the mesons that of the photons. This is the baryon-exchange mechanism. Now, the quark model offers a dramatic alternative where annihilation into three ordinary meson resonances can proceed
via a simple rearrangement of the constituents, similar to the rearrangement
of atoms and ions in molecular collisions. Presently, there is much activity
in analysing whether annihilation consists mostly of quark rearrangement, or
involves several quark-antiquark pairs being absorbed into or created out of the
gluon field.

The second reason deals with meson spectroscopy. In the past, several
mesons have been revealed by $\bar{p}p$ annihilation. The experiments presently run-
ning at LEAR are very useful for clarifying the situation of meson resonances in
the mass range 1.0–1.5 GeV/$c^2$. There are already indications for exotic mesons.
Hopefully, some hybrids, glueballs, multiquarks or quasi-nuclear bound states
will show up in the spectrum.

3.2 Quantum numbers

Annihilation at rest takes place in the S- or P-waves of protonium. Annihilation
in flight can involve an angular momentum $L \geq 2$ between $N$ and $\overline{N}$, but the
algebra of quantum numbers is the same, and we can restrict ourselves to $L \leq 1$
in this subsection.

The partial waves are denoted by the standard spectroscopic notation $^{2S+1}L_J$
or $^{2I+1, 2S+1}L_J$, i.e. the same notation as for charmonium, supplemented when
needed by the isospin multiplicity. Orbital mixing such as $^3S_1 \leftrightarrow ^3D_1$ does not
change the selection rules and is thus omitted in this subsection. For each par-
tial wave, one can compute the parity, $C$-conjugation and $G$-parity. The results
are listed in Table 1.

| $^{2I+1, 2S+1}L_J$ | $^{2I}S^0_0$ | $^{3I}S^0_0$ | $^{3I}S^1_1$ | $^{3I}S^1_1$ | $^{1I}P^0_1$ | $^{3I}P^1_1$ |
|-------------------|-------------|-------------|-------------|-------------|-------------|-------------|
| $J^{PC}$ ($I^G$)  | $0^+ (0^+)$ | $0^+ (1^-)$ | $1^- (0^-)$ | $1^- (1^+)$ | $1^+ (0^-)$ | $1^+ (1^-)$ |

| $^{2I+1, 2S+1}L_J$ | $^{13}P^0_0$ | $^{33}P^0_0$ | $^{13}P^1_1$ | $^{33}P^1_1$ | $^{13}P^2_2$ | $^{33}P^2_2$ |
|-------------------|-------------|-------------|-------------|-------------|-------------|-------------|
| $J^{PC}$ ($I^G$)  | $0^+ (0^+)$ | $0^+ (1^-)$ | $1^+ (0^+)$ | $1^+ (1^-)$ | $2^+ (0^+)$ | $2^+ (1^-)$ |

Let us consider a few mesonic states as pedagogical examples.

i) $p\bar{p} \rightarrow \pi^+\pi^-$. It selects the natural-parity partial waves with $P = C = (-1)^J$.
Since the $G$-parity of $\pi\pi$ is $G = +1$, we end with only $^{33}S_1$, $^{13}P_0$ and $^{13}P_2$
as candidates with $L \leq 1$. Alternatively, one can use a generalized Pauli
principle: if the spatial wave function of the two pions is even, as in the case
of $J = 0$ or 2, the isospin wave function should be symmetric, implying
$I = 0$; an antisymmetric space wave function ($J = 1$) is associated with
an antisymmetric coupling of the isospins, i.e., $I = 1$.

ii) $p\bar{p} \rightarrow K^+K^-$. As in the previous case, only $J^{PC} = 0^{++}, 1^{--}, 2^{++}, \ldots$ are
possible. But there is no further restriction. $K^+$ and $K^-$ do not belong
to the same isospin multiplet, and thus need not obey a generalized Bose

\[ S_1, \quad S_3, \quad S_{1/2}, \quad S_{3/2}, \quad P_0, \quad P_{1/2}, \quad P_{3/2} \]

contribute.

iii) \( p\bar{p} \rightarrow K^0 \bar{K}^0 \). This seems at first to be the same situation as for \( K^+K^- \).

In fact, the experimentalists do not detect \( K_0^+ \) or \( K_0^- \), but \( K_0^0 \) or \( K^0_L \). The
reaction \( p\bar{p} \rightarrow K_0^0K_0^0 \) selects \( CP = (-)^J \), i.e. the intrinsic \( CP \) and the orbital
contribution, and thus arises from \( 3\,P_0 \) or \( 3\,P_2 \), while \( p\bar{p} \rightarrow K_0^0K^0_L \)
comes from \( 3\,S_1 \).

iv) \( p\bar{p} \rightarrow \pi^0\pi^0 \) selects \( I = 0, \quad G = +1, \quad \text{natural parity}, \quad \text{and} \quad J \) even (Bose
statistics). Thus only \( 1\,P_0 \) and \( 1\,P_2 \) are possible.

v) \( p\bar{p} \rightarrow \rho^0\rho^0 \) requires \( I = 0, \quad G = +1 \) and a symmetric final state. For \( 3\,P_0 \), \( 1\,P_1 \) and \( 1\,P_2 \), one can combine a symmetric space wave function
(\( \ell \rho \) even) and symmetric spin wave function \( (S = 0, \quad 2) \). For \( 1\,S_0 \), one
should choose the antisymmetric coupling \( S = 1 \) and an orbital momentum
\( \ell \rho \) = 1.

vi) \( p\bar{p} \rightarrow \pi^0\pi^0\pi^0 \) implies \( C = +1, \quad G = -1 \) and thus \( I = 1 \). To implement Bose
statistics, it is convenient to use the Jacobi variables already introduced
by G. Karl in his lectures on baryon structure [81]

\[
\vec{\rho} = \vec{r}_2 - \vec{r}_1 ,
\]

\[
\vec{\lambda} = (2\vec{r}_3 - \vec{r}_1 - \vec{r}_2)/\sqrt{3} .
\]

A wave function is symmetric if it survives \((1 \leftrightarrow 2)\) exchange \( P_{12} \) and

circular permutation \( P_- \). One easily checks that

\[
P_{12}(\vec{\lambda} + i\vec{\rho}) = (\vec{\lambda} + i\vec{\rho})^* ,
\]

\[
P_-(\vec{\lambda} + i\vec{\rho}) = j(\vec{\lambda} + i\vec{\rho}) ,
\]

where \( j = \exp(2i\pi/3) \), as usual. A constant is symmetric and this con-
vinces us that \( 3\,S_0 \) is allowed. In the same \( J = 0 \), we cannot accommodate
\( 3\,P_0 \), because it has the wrong parity. For \( 3\,P_1 \), we can exhibit

\[
(\lambda^2 - \rho^2)\vec{\lambda} - (2\vec{\lambda} \cdot \vec{\rho})\vec{\rho} \quad (30)
\]

as having the right \( J^P \) and being fully symmetric. In [81], one indeed

reminds that the pair \((\lambda^2 - \rho^2), (2\vec{\lambda} \cdot \vec{\rho})\) behaves like the pair \( \vec{\lambda}, -\vec{\rho} \). The
above expression is thus a generalization of the well-known symmetric
polynomial

\[
\text{Re} [(\vec{\lambda} + i\vec{\rho})(\vec{\lambda} - i\vec{\rho})] \quad (31)
\]

which is proportional to \( \sum r_{ij}^2 \). We note some structure in [81]. There are
altogether three units of internal angular momenta to end with \( J = 1 \). So
one might expect some suppression, or at least some structure in the Dalitz plot. One can also find explicit wave-functions for $^{3}P_{2}$. In summary, $^{3}S_{0}, ^{3}P_{1}$ and $^{3}P_{2}$ are possible initial states.

vii) The case of $p\bar{p} \rightarrow \pi^{0}\pi^{0}\pi^{0}\pi^{0}$ is even more delicate. Unfortunately, modern detectors such as Crystal Barrel are able to record events with 8 photons and even more, and this channel has to be considered. We have $C = +1$, $G = +1$, $I = 0$. A constant, or say an overall S-wave is symmetric under all permutations, so one easily conceives that $^{13}P_{0}$ is a possible source of $4\pi^{0}$.

For constructing explicit examples of wave-functions with full permutation symmetry and a $J^{P}$ less obvious than $0^{+}$, it is convenient to use the relative variables

$$
\vec{u} = \vec{r}_{4} + \vec{r}_{1} - \vec{r}_{2} - \vec{r}_{3}, \\
\vec{v} = \vec{r}_{4} + \vec{r}_{2} - \vec{r}_{1} - \vec{r}_{3}, \\
\vec{w} = \vec{r}_{4} + \vec{r}_{3} - \vec{r}_{1} - \vec{r}_{2},
$$

(32)

in terms of which the permutations are easily translated: $(1 \leftrightarrow 2)$ exchange results into $(\vec{u} \leftrightarrow \vec{v})$, $(1 \leftrightarrow 4)$ exchange into $(\vec{v} \leftrightarrow -\vec{w})$, etc.

For $J^{P} = 2^{+}$, and for instance $J_{z} = 2$, one can exhibit

$$
u^{2}_{+} + v^{2}_{+} + w^{2}_{+}
$$

(33)

which is fully symmetric.

For $J^{P} = 1^{+}$, we have for instance

$$[(\vec{u} \times \vec{v}) \cdot (\vec{v} \times \vec{w})] \vec{u} \times \vec{v} + [(\vec{v} \times \vec{u}) \cdot (\vec{v} \times \vec{u})] \vec{v} \times \vec{u}

+ [(\vec{w} \times \vec{v}) \cdot (\vec{u} \times \vec{v})] \vec{v} \times \vec{u}.
$$

(34)

And for $J^{P} = 0^{-}$

$$(u^{2} - v^{2}) (v^{2} - w^{2}) (w^{2} - u^{2}) [(\vec{u} \times \vec{v}) \cdot \vec{w}].
$$

(35)

In the later case, we have 9 units of internal orbital excitation to end with $J = 0$. So the transition is likely to be suppressed in the $0^{-}$ channel. From the structure of the polynomial and of its analogue in momentum space, one expects many holes in the Dalitz plot.

In summary, $^{11}S_{0}, ^{13}P_{0}, ^{13}P_{1}$ and $^{13}P_{2}$ can decay into four neutral pions.

### 3.3 Phase-space considerations

Annihilation at rest produces an average of 5 pions in the final state. Many channels are open, and the results of annihilation experiments are often expressed in terms of branching ratios $B_{\alpha} = BR(pp \rightarrow \alpha)$. Each $B_{\alpha}$ is small, and one needs many contributions to achieve $\sum B_{\alpha} = 1$. 

22
As reviewed by Ecker [52], the pion is very light since it is the Goldstone boson associated with the chiral symmetry of the QCD Lagrangian. Indeed, the pion mass is \( m_\pi = 140 \text{ MeV} \), while \( m_N = 940 \text{ MeV} \), so \( N\bar{N} \) annihilation can involve up to 13 pions.

In the heavy quark limit, the inequality

\[
Q\bar{Q} + \bar{Q}Q > 3(Q\bar{Q})
\]

is believed to hold [82], so that annihilation can proceed into three mesons. On the other hand, if the baryon and the antibaryon are built of very different quarks, one might get the inverted inequality

\[
\bar{Q}Q + q\bar{q} > 3(\bar{Q}q),
\]

provided the mass ratio is large enough. This means that a very heavy antibaryon such as \( \Omega_{ccc} \) would not annihilate on ordinary matter.

Back to the real word of \( NN \) annihilation. Even at rest, the phase-space is comfortable and meson resonances can be produced in the primary process and then these resonances decay into the pions that are eventually detected.

### 3.4 Range of annihilation

In QED, \( e^+e^- \) annihilation into photons proceeds via the exchange of an electron, a very short-range process on the atomic scale. The naive analogue for \( NN \) would be the exchange of a baryon. The range is given by the inverse of the nucleon mass, of the order of 0.1 fm. To understand the observed cross-sections, in particular the large ratio of inelastic to elastic, one needs an absorption acting up to 0.8–1 fm. This is a clear conclusion of the optical model analysis.

In the baryon-exchange picture, one has to introduce large form-factor corrections to account for the finite size of \( N \) and \( \bar{N} \).

At this point, since the structure of the nucleon is involved, it becomes natural to think in terms of quarks. The simplest mechanism is quark rearrangement (see Fig. 3). It is similar to rearrangement collisions in molecular physics. Its range is governed by the size of the incoming and outgoing clusters. This leads rather naturally to the desired order of magnitude, 1 fm.

One may even push the reasoning a little further [83]. Annihilation into pions receives a contribution of quark rearrangement which has the largest possible range. To produce a \( K\bar{K} \) pair, one needs to create and annihilate some quark pairs, as in the second diagram in Fig. 3. This makes the range shorter.

The extreme case is annihilation into \( \phi\phi \), slightly above the threshold and currently under investigation in the JETSET experiment at CERN. One should get rid of all incoming quarks and antiquarks. This is a genuine annihilation at the level of elementary constituents and one expects a rather short range.

These ideas are supported by the observation [41] that the ratio of branching ratios \( BR(K\bar{K})/BR(\pi\pi) \) is larger in S-wave than in P-wave: the production of strangeness is more central.
3.5 Strength of annihilation

The Jülich group, among others, has tried to obtain a realistic annihilation potential in terms of baryon exchanges, with the same coupling constants as for meson-nucleon and nucleon-nucleon scattering [39]. Note, however, that such models heavily rely on the form factors, so that the baryon-exchange mechanism is not tested in detail.

The contribution of the quark rearrangement diagram (on the right of Fig. 3) was calculated within simple quark models. It gives a significant fraction of the observed annihilation [46, 84]. So one can reasonably hope that when this rearrangement process is supplemented by the other contributions, one ends with a realistic strength for annihilation.

Let us repeat our warning. In such model calculations, the results are rather sensitive to the matching between the Yukawa potential of LR forces and the SR annihilation described in terms of quarks, and there is no safe guideline on how to arrange the transition.

3.6 Microscopic mechanisms

Several models reproduce the observed annihilation cross-section $\sigma_{\text{inel}}$, which measures the cumulated strength of annihilation. It is much more difficult and constraining to describe the detailed results of annihilation, namely the branching ratios $BR(p\bar{p} \rightarrow \alpha)$ into the various mesonic channels $\alpha$ and for any channel $\alpha$ with more than 2 particles, the momentum distribution of the mesons.

The goal is to find a leading mechanism that explains the main features of the data and then to work out minor improvements involving higher-order mechanisms, rescattering corrections, etc. So far, however, no such dominant mechanism has been identified and there are animete debates to stimulate the searches.

We have no time here to examine the details of the available models and to gauge their success and shortcomings. We shall restrict ourselves to a survey
of the various approaches and refer to the recent reviews [20, 21] for further
discussions and references.

The baryon-exchange model will be discussed by Holinde [39] who wrote
several papers on the subject. Once a baryon is exchanged between N and \( \bar{N} \),
you can produced two mesons or two meson resonances. You then enter the
club of “quasi-two-body” annihilation, on which we shall come back shortly.
The problem is first to describe the relative abundance of the different channels
with two mesons.

There is another school whose members analyse annihilation with consider-
ations based on flavour SU(3) symmetry. The first of these attempts, to my
knowledge, is in a paper by Rubinstein and Stern [85]. More recent works were
done by Genz, Körner, Klempt, etc. The literature can be traced back from the
most recent articles [86]. If this approach proves successful, it will allow us to
make predictions for other baryon-antibaryon channels.

In another branch of studies, annihilation is discussed in terms of quark
diagrams, those of Fig. 3 and others. One should first notice that these pictures
are not genuine Feynmann diagrams. Some theory should be included, or at
least some model, to associate these diagrams with actual numbers. One has
to consider that each diagram corresponds to many QCD diagrams where the
 gluons, not shown, are exchanged in all possible ways between the quarks and
antiquarks.

There is a widespread belief that QCD favours planar diagrams, with amni-
hilation of quark pairs in the initial state, and creation of new quark pairs in the
final state. The arguments based on \( 1/N_c \) expansion, where \( N_c \) is the number of
 colours, have been revised by Pirner [87] who concluded that there is no reason
to eliminate the non-planar diagrams.

These \( 1/N_c \) arguments are usually associated with very high energy physics.
At low energy, we have empirical models which work remarkably well, such as
the constituent quark model. Its most attractive property is that the number
of constituents is frozen: mesons contain a quark and an antiquark, baryons
are made out of three quarks and higher configurations with additional quark-
antiquark pairs do not seem to play a very important role. One is thus tempted
to describe hadron-hadron interactions by keeping as much as possible this sim-
licity, i.e. by minimizing the number of pair creations and annihilations. For
instance, the decay of a resonance involves a single pair creation. \( K^+ N \) scatter-
ing is understood by the interaction and exchange of the constituents. \( K^- N \)
 involves one creation and one annihilation, this providing the usual s-channel
resonances. We mentioned in Subsection 2.1 the semi-quantitative success of
quark model to account for the short-range repulsion in NN interaction. These
observations suggested for NN a scenario, where quark rearrangement domi-
nates, with corrections due to diagrams with a few creations or annihilations.
The corresponding calculations are summarized in [10].

At this point, it is rather difficult to draw conclusions, the theoretical argu-
ments being a little empirical or, say, handwaving. So one may try to answer the
question of the leading mechanism by a fair phenomenological analysis. There are here two categories of contributions:

i) global fits of all branching ratios, with a complicated superposition of planar and non-planar diagrams. Unfortunately, the answer seems not unique owing to the large number of parameters.

ii) detailed investigations of selected branching ratios. For instance, there are interesting papers on the πρ channels or on the decay into two pseudoscalars: ππ, πη, KK. Of particular interest are the channels involving the φ meson, more precisely the ratio BR(φ + X)/BR(ω + X). A clear violation of the OZI rule is observed, but the ratio puzzlingly depends very much on the partner X associated with φ or ω. These selected channels give direct insight into the physics, but there is a risk to promote a new mechanism for each peculiar set of BR, without reaching a global understanding.

May be we will never converge towards a simple quark mechanism. We have been reminded that the main features of annihilation can be understood by production of two meson resonances, followed by their decay into stable mesons \[8\]. The rate for producing the two primary resonances seems governed by simple phase-space considerations \[8\]. On the other hand, the non-resonating part of pion production can be viewed as a radiation of the incoming nucleon and antinucleon, which are strongly accelerated by their mutual interaction \[8\].

### 3.7 Annihilation in flight

There are not too many data on annihilation in flight, but the Crystal Barrel collaboration will take some data during the next runs \[41\]. The E760 experiment at Fermilab has also collected interesting results, as a side product of their study of charmonium \[37\].

As the phase-space increases, higher meson resonances can be produced, and higher N\(N\) partial wave contribute. According to theoretical calculations, several glueballs and hybrids are expected near 2 GeV/c\(^2\) and thus are not accessible in annihilation at rest.

Let us mention now the beautiful results obtained in studying the reactions p\(\bar{p}\) → \(\pi^+\pi^-\) and p\(\bar{p}\) → K\(^+\)K\(^-\) with a target which is polarized in the transverse direction. One measures the angular distribution \(d(\vartheta, \varphi)\), whose average over the azimuthal angle \(\varphi\) is the usual differential cross-section \(\sigma(\vartheta)\). The azimuthal dependence reflects the correlation with the spin of the proton, and provides the asymmetry parameter \(A\). It is found that \(A\) is very large, nearly saturating the unitarity bound \(|A = 1|\) in a wide domain of energy and scattering angle \(\vartheta\).

This result was analysed in several papers \[90\,93,94\]. To get a large \(A\), one needs altogether the strong tensor force in the initial state, and a specific

\[8\] See also the contributions by these authors in \[91,92\,26\].
spin dependence of the annihilation operator. The comparison of $\pi\pi$ and $K\bar{K}$ reactions is instructive. The latter contains a larger fraction of initial S-wave, in full agreement with the analysis of annihilation at rest.

4 Conclusions

Nucleon-antinucleon interaction is clearly a very dense field. At present, we have many data, but no clear view of the basic dynamics has been completed yet. Analysing the complex optical images, one has to apply filters to separate the various contributions.

The first filter is the spin degree of freedom. Spin observables allow us to study long-range forces, and their interplay with annihilation. After all, the meson-exchange picture of LR forces need not be postulated. It has to be checked and the crucial experiments remain to be done.

The second filter is isospin, and more generally, flavour. Much information has been obtained from the charge-exchange reaction and from the production of hyperon-antihyperon pairs. We hope that the antineutron program will be resumed at OBELIX. The hyperon program requires either a polarized target, or higher energies.

In annihilation experiments, one has cleverly used X-ray detectors or the pressure of the target to filter S-wave from P-wave annihilation, this providing very useful information. The new detectors enable us to measure new channels, with neutral mesons or strange particles and to analyse all possible correlations between the mesons in the final state. Another development is annihilation in flight and again polarization might well be a useful tool, as it proved to be for the channels with two pions or two kaons.

My participation in this Summer School was supported in part by the PARCECO program of the Ministère de l’Enseignement Supérieur et de la Recherche, Paris. I am very grateful to A. Cieply for his help in preparing the lecture notes for the Proceedings. The manuscript was completed during a visit at the University of South Carolina. I thank F. Myhrer and K. Kubodera for their warm hospitality and for several useful discussions.

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