Hadron physics in low-energy antiproton beams

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The present state of low-energy nucleon-antinucleon physics is reviewed. Some recent experimental results, described here, cannot be explained in the standard potential models without appealing to a quark-gluon structure of the hadrons. Various theoretical approaches which are being taken in the physics of the nucleon-antinucleon interaction are discussed and compared. The outlook for further research in antiproton beams is outlined.

INTRODUCTION

Antiprotons were first produced at the Betatron accelerator in Berkeley in 1955.[1] Up to the early 1980s, the beams of low-energy antiprotons were of low intensity ($\lesssim 10^3 \bar{p}/s$), and most measurements were carried out with bubble chambers (see the review by Armenteros and French[2]). Among the most important achievements of this initial stage of the experiments one should include the discoveries of the ω, D, E, and Q_1 mesons[3–6]; measurements of the widths of the ω and φ mesons[7]; analysis of the quantum numbers of the D, E, A_2, and K* mesons[8]; observation of the decay[9] $\delta^+ \to \eta\pi^\pm$; and the experimental confirmation of the Day–Snow–Sucher mechanism,[10] which governs the annihilation of antiprotons which are stopped in liquid hydrogen.

A new stage in low-energy antinucleon physics began in 1983, when the LEAR (Low Energy Antiproton Ring) antiproton complex came on line at CERN. The basic characteristics of this installation are listed in Table I. About 20 experiments were carried out at LEAR, involving the efforts of about 300 physicists from 50 countries, including the USSR. Several experiments in low-energy antiproton beams were also carried out over the same period at Brookhaven National Laboratory (BNL, in the US) and at the National Laboratory of High-Energy Physics (KEK, in Japan) (the characteristics of the corresponding beams are shown in Table I). Some experiments with unseparated secondary beams of antiprotons and antineutrons were carried out in the USSR (at the Institute of High-Energy Physics and the Institute of Theoretical and Experimental Physics).

The LEAR experiments were devoted to the structure of the amplitude of the $\bar{N}N$ interaction at very low energies ($T \leqslant 50$ Mev). The properties of this amplitude turned out to be extremely unusual: The higher-order partial waves "turn on" even near the $\bar{N}N$ threshold; and a puzzling irregularity is observed in the energy dependence of the ratio of the real and imaginary parts of the amplitude for forward elastic $\bar{p}p$ scattering. A systematic study was made of the nuclear shifts and annihilation widths of the levels of the $\bar{p}p$ atom at LEAR; these studies yielded important information about the amplitude for the $\bar{N}N$ interaction at zero energy. Unexpected results emerged from a study of exclusive annihilation channels, in particular, the channel $\bar{p}p \to \Lambda\bar{\Lambda}$. Important information was obtained on the interaction of antiprotons with nuclei. A search for narrow resonances near the $\bar{N}N$ threshold did not meet with success. The large statistical base of these experiments, forces the conclusion that the numerous earlier indications of the existence of narrow resonances in the $\bar{N}N$ system were unreliable.

In 1988, after a reconstruction of the CERN antiproton storage ring at LEAR, experiments were begun with a much-improved beam (Table I). A program of experiments is being performed at the new generation of installations, with a high energy resolution. At these installations it is possible to detect completely and to identify both the charged and neutral particles and to study exclusive annihilation channels with a large statistical base. In addition to the research on the dynamics of annihilation and meson spectroscopy, which are standard questions in antiproton physics, the new experiments are making tests of the fundamental symmetries (CP, T, CPT) and searching for exotic states and interactions.

Low-energy antinucleon physics is presently developing at a very fast pace. In a review of limited size we obviously cannot discuss every single result of significance. Our basic goals are to present the most interesting experimental results of recent years and to discuss briefly the existing theories and the major directions of future research.

The material of this review is organized in five sections. Section 1 deals with the basic properties of the nucleon-antinucleon interaction. Section 2 deals with the dynamics of annihilation. In Sec. 3 we discuss the present state of the search for baryonium and exotic mesons in $\bar{p}p$ annihilation. In Sec. 4 we discuss some interesting effects in antinucleon-nucleus interactions. In Sec. 5 we discuss the outlook for further research in antiproton beams.

Earlier stages in the experimental and theoretical research on low-energy antinucleon physics are reflected in some previous reviews in *This Journal*.[11,12]

Diverse additional information on these questions can be found in the proceedings of conferences on the LEAR research programs.[14–17]

1. BASIC PROPERTIES OF THE NUCLEON-ANTINUCLEON INTERACTION

1.1. Elastic $\bar{N}N$ scattering, charge exchange, and annihilation

The interaction in the $\bar{N}N$ system is quite different from the well-studied NN interaction. The $N\bar{N}$ system has a baryon number zero, and it easily annihilates into mesons. At low energies the annihilation cross section is more than twice the elastic cross section (Fig. 1). The Pauli principle imposes limitations on the wave function of the $N\bar{N}$ system which are less stringent than those imposed on the wave function of the NN system.[19] For this reason, the number of partial nucleon-nucleon scattering amplitudes which are allowed is roughly twice the number of amplitudes in NN scattering.[20] Furthermore, annihilation renders each of the $N\bar{N}$-scattering phase shifts complex. A phase-shift analysis of the $N\bar{N}$ interaction is thus a very difficult problem. For exam-

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TABLE I.

a) Comparative characteristics of low-energy antiproton beams.

	Pre-LEAR	LEAR (1983)	LEAR (after reconstruction)
Intensity, $\bar{p}$/s	10^3	$6\cdot10^5$	$3\cdot10^6$
Momentum, GeV/c		0.1—1.5	0.06—2
Momentum spread	1%	$(2—3)\cdot10^{-3}$	10^{-4}
Impurity of secondary particles	$10^4—10^5$	None	None
Beam emittance, mm·mrad	100	3—5	3
Extension	0.1	1	1

b) Separated antiproton beams, BNL[13]

Beam	Momentum p, GeV/c	$\Delta p/p$, %	Particle	Flux per 10^{12} primary protons
B2, B4	1.5—9	3	$\bar{p}$	10^5 ($\pi/\bar{p}\approx 3/4$)
C2, C4	1.1	2	$\bar{p}$	$2\cdot10^3$
C6, C8	0.8	2.5	$\bar{p}$	$1.4\cdot10^3$

c) Antiproton beams, KEK[13]

Beam	Momentum p, GeV/c	$\Delta p/p$, %	Particle	Flux per 10^{12} primary protons
π2	2—4.3	1	p/$\bar{p}$	$10^6/10^3$ at 3 GeV/c
K2	1—2	3	p/$\bar{p}$	$5\cdot10^6/3\cdot10^3$ (2 GeV/c)
K3 — L	0.5—1	2	p/$\bar{p}$	$2\cdot10^7/90$ (0.8 GeV/c)
K4	0.4—0.8	2	$\bar{p}$	175 at 0.6 GeV/c

ple, while np scattering at a fixed energy is described by five parameters in the case $J\neq0$ and two in the case $J=0$, $\bar{p}p$ scattering is characterized by 20 parameters at $J\neq0$ and five at $J=0$ (Ref. 20).

The isospins of the nucleon and the antinucleon combine to give either 0 or 1. The expansion of physical states in states of the isospin basis $|I,I_3\rangle$ thus takes the form

$$|\bar{n}p\rangle=|1,+1\rangle,\quad |\bar{n}n\rangle=\frac{1}{\sqrt{2}}(|0,0\rangle+|1,0\rangle),$$
$$|\bar{p}n\rangle=|1,-1\rangle,\quad |\bar{p}p\rangle=\frac{1}{\sqrt{2}}(|0,0\rangle-|1,0\rangle). \tag{1.1}$$

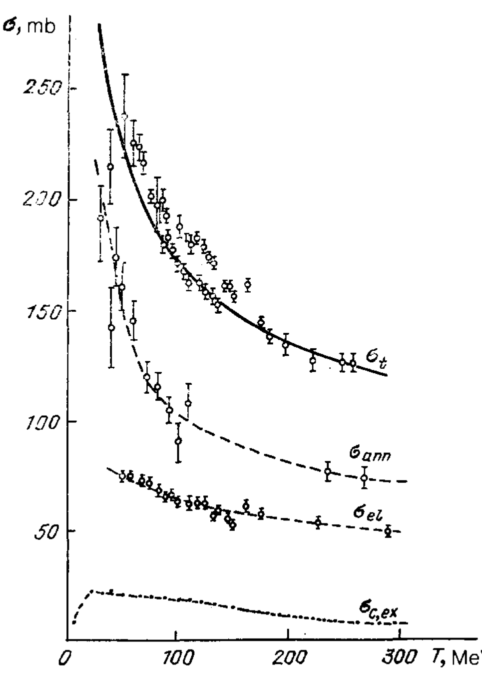

FIG. 1. Cross sections for the $\bar{p}p$ interaction. The experimental data are taken from the compilation of Ref. 18. The curves are drawn to aid the eye.

If we ignore the Coulomb interaction (see Ref. 21, for example, for refinements), we can express the cross sections for elastic $\bar{p}p$ and $\bar{n}p$ scattering and for the charge exchange $\bar{p}p\to\bar{n}n$ in terms of the amplitudes $f_{ll'}{}^{JS}(I)$ in the isospin basis:

$$\sigma_{el}(\bar{p}p)=\pi\sum_{J=0}(2J+1)\sum_{S}\sum_{l,l'}|f_{ll'}^{JS}(I=0)+f_{ll'}^{JS}(I=1)|^2,$$
$$\sigma_{el}(\bar{n}p)=4\pi\sum_{J=0}(2J+1)\sum_{S}\sum_{l,l'}|f_{ll'}^{JS}(I=1)|^2, \tag{1.2}$$
$$\sigma_{ce}(\bar{p}p\to\bar{n}n)$$
$$=\pi\frac{k_n}{k_p}\sum_{J=0}(2J+1)\sum_{S}\sum_{l,l'}|f_{ll'}^{JS}(I=0)-f_{ll'}^{JS}(I=1)|^2,$$

where $l'=l$ or $l\pm2$, $|J\text{-}S|\leqslant l$, $l'\leqslant(J+S)$, and k_p are the momenta in the $\bar{p}p$ and $\bar{n}n$ channels in the c.m. frame, which are related by

$$k_p=\frac{p_L}{2},\quad k_n^2-k_p^2=m_p\delta,$$
$$\delta=2(m_n-m_p)=2.59\text{ MeV}. \tag{1.3}$$

Figure 1 illustrates the relations among the cross sections for annihilation, elastic scattering, and charge exchange. At low energies the total cross section σ_t and the annihilation cross section σ_a for the $\overline{N}N$ interaction are customarily approximated by

TABLE II. The coefficients A_i and B_i in parametrization (1.4) of the energy dependence of the $\overline{\rm N}$N interaction cross sections. Shown for comparison are theoretical results[25] on the $\bar{\rm p}$n cross section according to $\bar{\rm p}$d data.

Quantity	A_i, mb	B_i, mb·GeV/c	Range of p_L, MeV/c	Ref.
$\sigma_t(\bar{p}p)$	65.55	53.84	220—600	[22]
$\sigma_a(\bar{p}p)$	29.43±1.83	32.00±0.92	400—600	[23]
$\sigma_t(\bar{n}p)$	94.4±9.0	36.0±2.9	100—500	[24]
$\sigma_a(\bar{n}p)$	41.4±9.0	29.0±2.9	100—500	[24]
$\sigma_t(\bar{p}n)$	60.59	44.26	300—800	[25]
$\sigma_a(\bar{p}n)$	37.20	27.97	300—800	[25]

$$\sigma_i = A_i + \frac{B_i}{p_L} \quad (i = \text{t, a}), \tag{1.4}$$

where A_i and B_i are empirical constants.

Table II shows sets of the parameters A_i and B_i which have been found through an analysis of LEAR data[22,23] and BNL data.[24] It follows from Table II that in the energy interval in which the measurements were taken the total cross section of the $\bar{\rm n}$p direction is about 20–30% smaller than the corresponding values for the $\bar{\rm p}$p interaction. It can thus be concluded that the amplitude of the $\overline{\rm N}$N interaction has a significant isospin dependence (see also Subsec. 2.1).

In the annihilation of antiprotons it is primarily π mesons which are formed. The average multiplicity in $\bar{\rm p}$p annihilation at rest is $\bar{n}_\pi = 5.0 \pm 0.2$ (Ref. 26). Table III shows the relative probabilities for the various channels for $\bar{\rm p}$p annihilation at rest. The pion energy spectrum has a maximum near $T_\pi \approx 200$ MeV, i.e., in the region of the excitation of the Δ_{33} resonance (Fig. 2). This feature results in an intensification of pion rescattering effects in the annihilation of antiprotons with nuclei. The yield of kaons in $\bar{\rm p}$p annihilation at low energies is about 5%.

The total worldwide statistical base on $\bar{\rm p}$p annihilation at rest has reached $\approx 10^7$ events. Nevertheless, we are still a long way from having comprehensive information about the annihilation process. The uncertainty stems primarily from the channels which contain several neutral particles. These channels account for about 60% of the annihilation events.[28–30]

The cross section for the charge exchange $\bar{\rm p}{\rm p}\to\bar{\rm n}{\rm n}$ is only 6–7% of the total cross section. This process has attracted interest for two reasons. First, the differential cross section for charge exchange is far more sensitive than the differential cross section for elastic scattering to the model of the N$\overline{\rm N}$ interaction.[31,32] Second, charge exchange has recently been used to generate antineutron beams.[24,33]

1.2. Nucleon-antinucleon interaction near the threshold; antiprotonic atoms

Information about the $\overline{\rm N}$N scattering amplitude at the threshold can be found by measuring the nuclear shifts and annihilation widths of protonium: the atom consisting of a proton and as antiproton. The $\bar{\rm p}$p atom has a first Bohr radius $a_B \approx 2\alpha m_p = 57.6$ fm, and the binding energy of the 1S level in the Coulomb potential is $E_{1S} = 1/(m_p a_B^2) = 12.5$ keV. Three experimental groups have measured the nuclear shifts Re ΔE_{1S} and the annihilation width of the 1S level of

TABLE III. Branching ratios of certain channels for the annihilation of antiprotons at rest in liquid hydrogen (B is given as a percentage).

Channel	B	Channel	B
$\pi^+\pi^-$	0.37±0.03	$\pi^0\varphi$	$(3.3\pm1.5)\cdot10^{-2}$
$\pi^0\pi^0$	$(2.06\pm0.14)\cdot10^{-2}$	$\pi^0\eta'$	0.05±0.019
$\pi^+\pi^-\pi^0$	6.9±0.35	$\pi^0\omega$	0.52±0.05
$3\pi^0$	$(7.6\pm2.3)\cdot10^{-1}$	$\pi^0\rho^0$	1.6±0.10
$\pi^+\pi^-2\pi^0$	9.3±3.0	$\pi^0\eta$	0.0133±0.003
$\pi^+\pi^-3\pi^0$	23.3±3.0	π^0a_2	1.07±0.13
$\pi^+\pi^-4\pi^0$	2.8±0.7	$\eta\rho^0$	0.5±0.17
$\pi^+\pi^-X$	35.8±0.8	$\eta\omega$	0.4±0.1
$2\pi^+2\pi^-$	6.9±0.6	$\eta\eta$	$(8.1\pm3.1)\cdot10^{-3}$
$2\pi^+2\pi^-\pi^0$	19.6±0.7	$\pi^0\gamma$	$(1.74\pm0.22)\cdot10^{-3}$
$2\pi^+2\pi^-2\pi^0$	16.6±1.0		
$2\pi^+2\pi^-3\pi^0$	4.2±1.0	$\gamma\gamma$	$<1.7\cdot10^{-4}$
$2\pi^+2\pi^-X$	20.8±0.7	$\pi^\pm a_0^\mp$	0.76±0.12±0.08
$3\pi^+3\pi^-$	2.1±0.2	$\pi^\pm b_1^\mp$	1.76±0.2±0.18
$3\pi^+3\pi^-\pi^0$	1.9±0.2	$\pi^\pm a_2^\mp$	2.64±0.2±0.26
Neutrals:	4.1±0.3	$\rho^\pm\pi^\mp$	3.014±0.41±0.3
$\rho^0\pi^+\pi^-\pi^0$	13.7±0.6	ρ^0f^0	0.9±0.2
$\rho^\pm\pi^\mp\pi^+\pi^-$	6.4±1.8	$\rho^0\pi^+\pi^-$	1.5±0.3
		$\omega\pi^+\pi^-$	3.0±0.3
		$\omega\rho^0$	3.49±0.56
		ωf^0	1.7±0.2

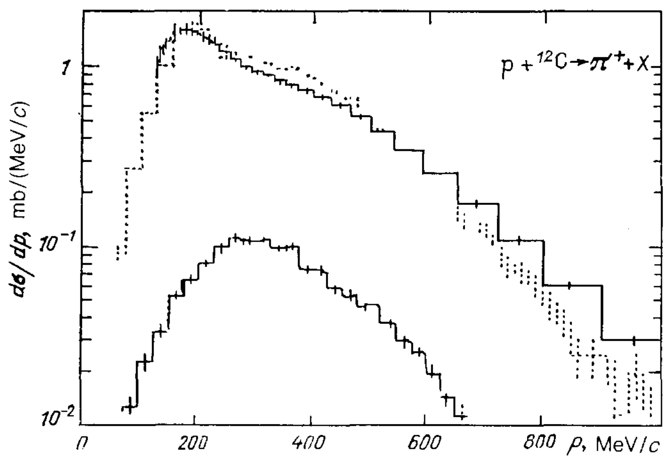

FIG. 2. The spectrum of π^+ mesons in $\bar{p}p$ and $\bar{p}^{12}C$ annihilation at $p_L = 608$ MeV/c. The solid and dotted histograms in the upper part of the figure show the inclusive pion spectra in $\bar{p}^{12}C$ annihilation which correspond to data and calculations based on the cascade model.[27]

the protonium atom at LEAR[34]; they found similar results (Table IV).

According to these data, the 1S level of the $\bar{p}p$ atom has an annihilation width $\Gamma_{1S} \approx 1$ keV, while the nuclear shift is $\mathrm{Re}(\Delta E_{1S}) \approx 0.7$ keV. The level is shifted upward; i.e., the binding energy is smaller than the unperturbed Coulomb value. The nuclear shift is small in comparison with the distance between the 1S and 2P levels, so the well-known perturbation-theory formula can be used[35]:

$$\mathrm{Re}\,\Delta E_{1S} - i\,\frac{\Gamma_{1S}}{2} = -\frac{4}{m_p a_B^3}\,a_{CS}, \qquad (1.5)$$

where a_{CS} is the nuclear scattering length distorted by the Coulomb interaction.

It might appear that we could find two parameters from (1.5): $(a_0 + a_1)/2$, i.e., half the sum of the nucleon-antinucleon scattering lengths in the states with $I = 0$, 1; and $\rho(p_L = 0) = \mathrm{Re} f_{el}/\mathrm{Im} f_{el}|_{\theta=0}$, i.e., the ratio of the real part of the amplitude for forward elastic $\bar{p}p$ scattering to the imaginary part at zero energy. It would appear at first glance to be natural[33,36–38] to identify a_{CS} with $(a_0 + a_1)/2$ and to identify the parameter $\rho(p_L = 0)$ with the quantity

$$\tilde{\rho} = -\frac{2\,\mathrm{Re}\,\Delta E_{1S}}{\Gamma_{1S}}. \qquad (1.6)$$

That conclusion, however, would be a bit hasty. In the first place, we are interested in the scattering amplitudes and in the quantity ρ in the limit $p_L \to 0$, while (1.5) refers to the 1S level below the threshold. Furthermore, it is necessary to take into account the Coulomb corrections and the circumstance that the problem is not a single-channel problem (there is a difference between the thresholds for the $\bar{p}p$ and $n\bar{n}$ channels). As was shown in Refs. 21, 39, and 40, the correct expressions relating the atomic shifts and widths with the scattering lengths are

$$a_{CS} = \frac{[(a_0 + a_1)/2] + a_0 a_1 \varkappa}{1 + (1/2)(a_0 + a_1)(\Delta_C + \varkappa) + a_0 a_1 \varkappa \Delta_C}, \qquad (1.7)$$

$$\rho(p_L \to 0) = \tilde{\rho}\left[1 + \frac{\pi}{2\alpha}\,a_B \Gamma_{1S}\,(1 + \tilde{\rho}^2)\right]^{-1}; \qquad (1.8)$$

Here $\varkappa = [2m_n(m_n - m_p)]^{1/2} \approx 0.25$, and $\Delta_C \approx -0.08$ fm^{-1} is the Coulomb correction to the scattering lengths.[41,42]

It can be seen from (1.7) that information on the shifts and widths of the $\bar{p}p$ atom is generally not sufficient in itself for a correct determination of the $\bar{p}p$ scattering lengths. Information about the $p\bar{n}$ interaction near the threshold is also necessary. The first experiments with slow antineutrons, recently carried out at BNL,[33] have revealed the imaginary part of the $\bar{N}N$ scattering length in the state with an isospin $I = 1$:

$$\mathrm{Im}\, a_1 = 0.83 \pm 0.07 \text{ fm}.$$

If we make use of additional information about the relationship between a_0 and a_1 which follows from the potential models (Table V), we can find the following average value of the $\bar{p}p$ scattering length, averaged over all existing data on atomic shifts[33]:

$$a_{\bar{p}p} = -(0.93 \pm 0.09) + i\,(0.95 \pm 0.12) \text{ fm}.$$

This value corresponds to

$$\mathrm{Im}\, a_0 = 1.07 \pm 0.16 \text{ fm}.$$

In the derivation of these values of $a_{\bar{p}p}$ and $\mathrm{Im}\, a_0$ in Ref. 33, the simplified formula $a_{CS} = a_{\bar{p}p}$ was used. Nevertheless, the errors which are introduced when we use these numerical values of the parameters do not exceed the experimental uncertainties.

The use of (1.8) to determine ρ from data obtained by the PS-174 collaboration (Table IV) leads to

$$\tilde{\rho} = -1.29 \pm 0.14, \quad \rho(p_L \to 0) = -1.08 \pm 0.14. \qquad (1.9)$$

The value of the parameter ρ at the threshold thus turns out to be large and negative (i.e., the sign of ρ corresponds to an effective repulsion in the S wave). A corresponding conclusion was reached in an analysis of data[47,48] on isotopic effects in the $\bar{p}^{16}O$, ^{17}O, and ^{18}O atoms. That analysis showed that the parameter ρ for the $\bar{p}n$ interaction is also large and negative[49]: $\rho_n(p_L \to 0) \approx -1$.

Because of the large annihilation width of the 1S level, its hyperfine structure—i.e., the splitting of the levels $^3S_1$

TABLE IV. The hadron shift $\mathrm{Re}E_{1S}$ and the annihilation width Γ_{1S} of the ground level of the antiprotonic-hydrogen atom.[34]

$\mathrm{Re}\Delta E_{1S}$, keV	Γ_{1S}, keV	Experiment
0.70 ± 0.15	1.60 ± 0.40	PS-171
0.73 ± 0.05	1.13 ± 0.09	PS-174
0.66 ± 0.13	1.13 ± 0.23	PS-175

TABLE V. S-wave NN scattering lengths in potential models and in the effective-radius approximation (ERA).

	a_0, fm, for isospin $I=0$	a_1, fm, for isospin $I=1$
Paris potential[43,44]	$-1.02+i0.68$	$-1.04+i0.91$
Dover–Richard potential[45]	$-0.86+i0.88$	$-0.94+i0.63$
ERA[39b]	$-1.1+i0.6$	$-0.1+i0.8$
ERA[21]	$-1.0+i1.1$	$-0.3+i0.2$
ERA[46]	$-0.6+i0.9$	$0.4+i0.5$

and $^1S_0$, has yet to be observed. Resolving these levels will require simultaneously detecting x rays of the Lyman series and annihilation mesons in certain exclusive channels. For example, an annihilation through the channel $\pi^0\pi^0\eta$ (in a state with quantum numbers $J^{PC}=0^{-+}$) could go only from the $^1S_0$ level.

A phenomenological analysis of data on $\bar{N}N$ scattering at low energies, based on the approximation of an effective radius, has been carried out in several places.[21,39,40,46,50,51] That formalism makes possible a model-independent description of the experimental data on the basis of the principles of analytically and unitarity. The multichannel nature of the $\bar{N}N$ interaction complicates the problem substantially. Unfortunately, the results presently available on the low-energy parameters, found in Refs. 21, 39, 40, and 46, are ambiguous. The primary reason for the ambiguity is the absence of experimental data at very low energies ($T\leqslant 20$ MeV), where the effective-radius expansion is most justified. Nevertheless, the solutions which have been found give a satisfactory description of the existing data on the differential cross sections and on the parameter ρ (more of this below).

1.3. Unusual energy dependence of the parameters of the $\bar{N}N$ interaction

1.3.1. Intensification of higher-order partial waves

In the LEAR experiments it was possible to move a good distance down the energy scale and to measure the differential cross sections for elastic $\bar{p}p$ scattering[38,52] and for the charge exchange [31]$\bar{p}p\to\bar{n}n$. These cross sections were measured down to an antiproton momentum $p_L=181$ MeV/c ($T=18$ MeV). An interesting effect was observed: The differential cross sections for elastic scattering and for charge exchange are sharply anisotropic down to the lowest momentum values (Fig. 3), in stark contrast with the picture seen in proton-proton scattering. The anisotropy of the angular distributions is evidence of a large component from waves with nonzero orbital angular momentum, primarily, an intense P wave. At $P_L=287$ MeV/c, for example, the P-wave component of the elastic $\bar{p}p$ cross section is 40%, while the D-wave component is 10% (Ref. 52). By way of comparison, 90% of the cross section for elastic pp scattering is in the S wave at the same energy.

Some very interesting results have come from measurements of the cross section and polarization of the $\Lambda(\bar{\Lambda})$ in the reaction $\bar{p}p\to\Lambda\bar{\Lambda}$ near the threshold[53] ($P_{th}=1435$ MeV/c). A surprising fact was observed: The angular distribution of the Λ hyperons remains anisotropic, while the polarization is large even at an energy of the $\Lambda\bar{\Lambda}$ system below 1 MeV (Fig. 4). Various approaches have been taken to describe the reaction $\bar{p}p\to\Lambda\bar{\Lambda}$: the exchange of K and K* mesons,[54] one-gluon exchange,[55] the Jaffe–Low P-matrix method,[56] the scattering-length approximation,[57] and a reasonance model.[58] Two facts—unrelated to specific models—can be judged solidly established. First, the interaction in the initial and final states, in particular, annihilation, plays an important role in the $\bar{p}p\to\Lambda\bar{\Lambda}$ reaction. Second, transitions $l_{\Lambda\bar{\Lambda}}=l_{p\bar{p}}-2$, in particular, $^3F_3\to{}^3P_2$, which are caused by the tensor interaction, are important. The absorption in the initial and final states is at its strongest for the S wave. A weakening of the S wave may be the main reason for the anisotropy of the angular distribution.[54b,56,57].

1.3.2. ρ oscillations

The behavior of the quantity $\rho(p_L) = \mathrm{Re}\, f_{el}/\mathrm{Im}\, f_{el}|_{\theta=0}$, i.e., the ratio of the real and imaginary parts of the amplitude for elastic $\bar{p}p$ forward scattering, has turned out to be very unusual. Figure 5 shows ρ as a function of the momentum p_L according to data on the nuclear-Coulomb interference in small-angle elastic $\bar{p}p$ scattering.[36–38] We see that ρ varies rapidly over the narrow energy interval from the $\bar{p}p$ threshold, where it is large and negative, to $p_L\approx 200$ MeV/c, where ρ is close to zero and may in fact be oscillating. The rapid increase in ρ near the threshold and its vanishing near $p_L\approx 200$ MeV/c are not described by the potential models of the $\bar{N}N$ interaction,[32,37,59,60] and they do not agree with predictions based on dispersion relations.[61,62]

Note that the parameter ρ which was measured in Refs. 36–38 is related to the amplitudes averaged over spin projections. At $p_L\leqslant 300$ MeV/c, for example, where the S and P

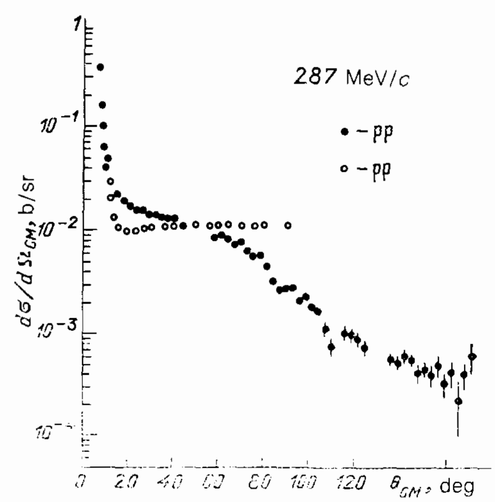

FIG. 3. Differential elastic $\bar{p}p$ scattering cross sections at 287 MeV/c (Ref. 52). Shown for comparison are data on elastic pp scattering (the open circles).

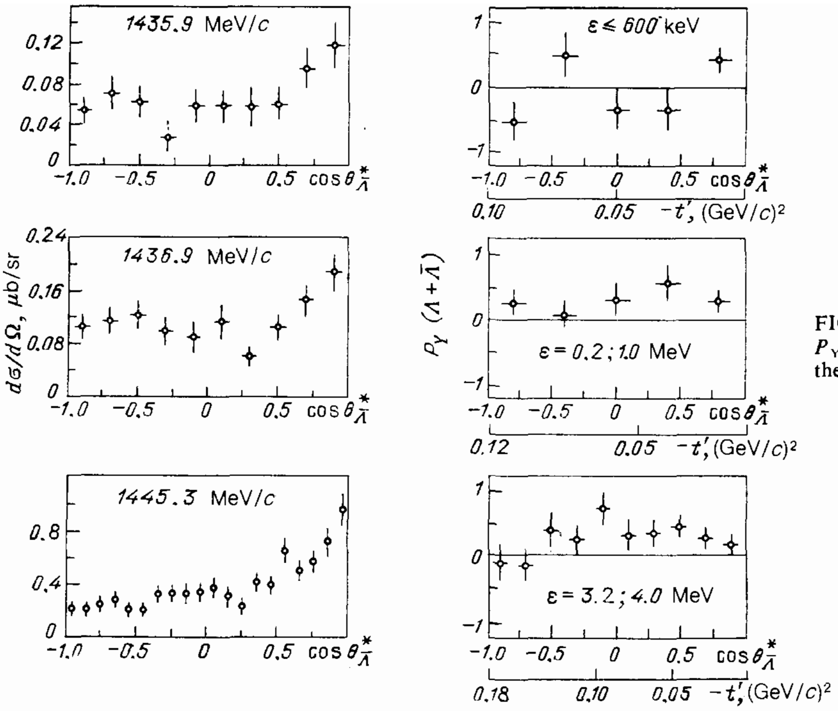

FIG. 4. Differential cross sections and average polarization $P_Y(\Lambda + \bar{\Lambda})$ in the reaction $\bar{p}p \to \bar{\Lambda}\Lambda$ at various energies near the threshold.[53]

waves are dominant, the ratio ρ can be expressed in terms of the amplitude f^{JS}_{ll} [see (1.2)] in the following way:

$$\rho = \frac{\mathrm{Re}\,(f^{00}_{00} + 3f^{11}_{00} + 3f^{10}_{11} + f^{01}_{11} + 3f^{11}_{11} + 5f^{21}_{11})}{\mathrm{Im}\,(f^{00}_{00} + 3f^{11}_{00} + 3f^{10}_{11} + f^{01}_{11} + 3f^{11}_{11} + 5f^{21}_{11})}. \qquad (1.10)$$

It is thus an extremely nontrivial fact that even the value of ρ averaged over spins is such a strong function of the energy. As a rule it turns out that if structural features are present in an amplitude averaged over spins then they will be even more prominent in the individual terms. This is the situation, for example, in pp scattering in the energy interval 2.2–2.5 GeV (Ref. 63).

Factors which have been suggested as possible reasons for the sharp energy dependence of ρ are a threshold structural feature associated with the opening of the $\bar{p}p \to \bar{n}n$ channel at $p_L = 98$ MeV/c (Refs. 36 and 64) and the existence of resonances in the P and D waves.[60,62,65–67] With regard to this threshold structural feature, we note that a systematic analysis[59,66] showed that it is smoothed out almost entirely by strong annihilation. The resonance explanation also runs into serious problems. It is very difficult to achieve the ρ behavior required even if the masses and widths of the resonances are left as completely adjustable parameters.[60,67] It was shown in Ref. 60 that the best fit of ρ is achieved by introducing contributions of two narrow ($\Gamma \approx 10$ MeV) resonances near the threshold in the $^{11}P_1$ and $^{31}P_1$ channels. The dynamics by which these resonances arise, however, is not clear. In the coupled-channel model[67] states with such quantum numbers should lie below the $\bar{N}N$ threshold and should have a width of about 70 MeV. Experimental searches for narrow resonances states in the $N\bar{N}$ system have not been successful (Sec. 3), but the energy region in the immediate vicinity of the threshold (± 20 MeV) has not yet been studied adequately. Nevertheless, there are serious theoretical arguments against the existence of narrow

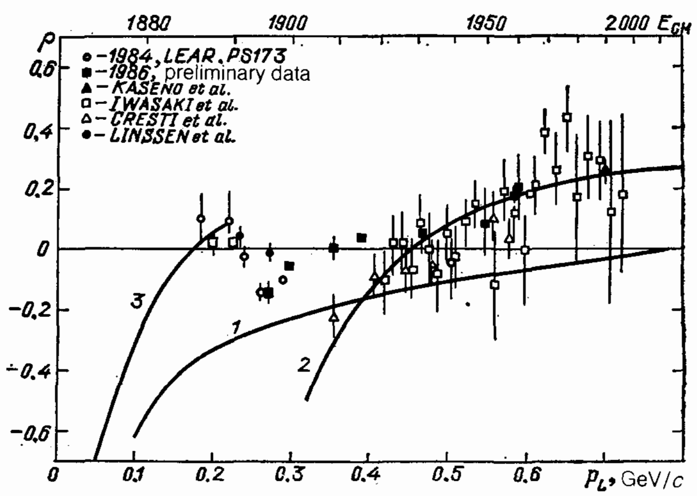

FIG. 5. The ratio $\rho = \mathrm{Re}\, f_{\bar{p}p}(0)/\mathrm{Im}\, f_{\bar{p}p}(0)$ as a function of the energy.[36–38] The solid lines show theoretical results found in the Dover-Richard potential model[45] (1), from dispersion relations[61,62] (2), and in the effective-radius approximation[21] (3).

$\overline{\mathrm{N}}$N states (Sec. 3). Resonances near the threshold are also absent from the Argand diagrams plotted on the basis of an analysis of data in the effective-radius approximation.[46,66]

In the effective-radius approximation, the rapid change in ρ from the threshold to 200 MeV/c is explained in terms of a destructive interference between an S-wave amplitude corresponding to an effective repulsion and an intense P-wave amplitude corresponding to an attraction (see the solid lines in Fig. 5).

The hypothesis of Ref. 68 regarding a relationship between the unusual behavior of ρ and the dynamics of various quark annihilation mechanisms is interesting. There are at least two quark mechanisms which are strongly dependent on the energy at $p_L \leqslant 2$ GeV/c. According to the unitarity condition, these mechanisms also contribute substantially to the cross section for elastic $\bar{p}$p scattering. One of these processes corresponds to a quark diagram with an exchange of diquark-antidiquark states in the t channel. Another mechanism is the annihilation which stems from a restructuring of valence quarks. Since the impact-parameter method which was used in Ref. 68 yields very crude estimates at low energies ($p_L \leqslant 500$ MeV/c), it would be useful to examine this hypothesis in more detail by other approaches, e.g., by the method of multichannel dispersion relations.

1.3.3. Expansion of the diffraction cone in elastic $\bar{p}$p scattering (the "antishrinkage effect")

The elastic scattering of antiprotons at low energies has yet another interesting aspect: the slope of the differential cross section for small-angle elastic scattering,

$$\frac{d\sigma}{dt} = \frac{\sigma_t^2}{16\pi} |\rho + i|^2 \exp(-bt), \qquad (1.11)$$

i.e., the parameter b in (1.11), is large. Figure 6 shows b as a function of the momentum of the antiproton. We see that the slope increases with decreasing p_L, reaching $b = 60$–80 (GeV/c)$^{-2}$ at $p_L = 200$ MeV/c. This behavior of b as a function of p_L is called "cone antishrinkage." We recall that in $\bar{p}$p scattering $T \approx 1$ GeV the slope of the cone is far smaller, $b = 6$ (GeV/c)$^{-2}$, and it increases with increasing energy.

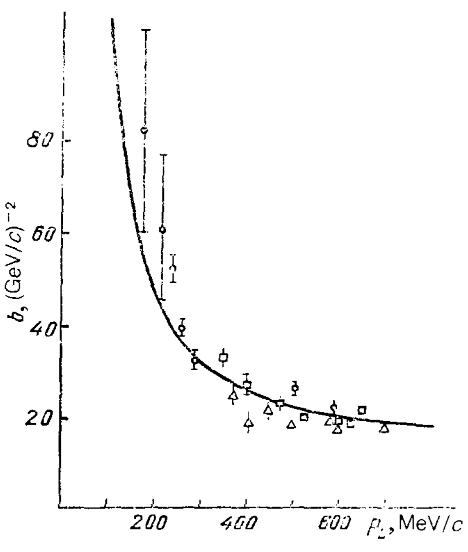

FIG. 6. Energy dependence of the cone slope parameter b for elastic $\bar{p}$p scattering. The solid line is a fit by the formula $b = 0.25[C + (A/k)]^2$, where $A = 0.691$, $C = 1.34$, and k is the momentum in the c.m. frame.[25,69]

It should of course be recalled that the differential cross section for elastic $\bar{p}$p scattering at small energies can be approximated by a single exponential function of the type in (1.11) only in a limited interval of the momentum transfer. At $p_L \leqslant 200$ MeV/c the concept of a diffraction cone becomes entirely meaningless.

A possible explanation of the antishrinkage effect was proposed in Ref. 68. If one assumes that the seed amplitudes associated with the various t-channel exchanges and with the annihilation mechanism fall off smoothly in impact-parameter space, then "unitarization" gives rise to an amplitude which corresponds to scattering by a black disk with an effective radius which increases with decreasing energy. Physically, this effect stems from the decrease in the transparency of the edges of the disk due to the rapid growth of the seed amplitude for annihilation, which sets in as the energy decreases (i.e., it is actually caused by the strong energy dependence of the annihilation amplitude, associated with quark dynamics).

Experiments on low-energy scattering of antiprotons have thus revealed some significant effects which stem from higher-order partial waves. Although the dynamic reason for the appearance of these effects is not yet completely clear, it appears to involve a nontrivial energy dependence of quark annihilation mechanisms (see the following section of this paper).

2. THE ANNIHILATION PROBLEM

The annihilation problem is the most complicated and most interesting one in the physics of the N$\overline{\mathrm{N}}$ interaction. Attempts to derive a theory for annihilation from the basic principles of quantum chromodynamics have not yet been successful. The primary reason is that the characteristic distances for annihilation are ≈ 0.6–1 fm, and this is the region in which nonperturbative effects are manifested. As a result, the experimental data are analyzed on the basis of various models, the most popular of which are the potential and quark models. In this section of the paper we will briefly discuss some common approaches to the annihilation problem and the most important experimental results.

2.1. Annihilation radius and potential models

Attempts have been made for many years now to describe the N$\overline{\mathrm{N}}$ system by a potential approach. Several of these models are based on Martin's argument[70] that the annihilation radius is small. Martin analyzed some very simple annihilation diagrams with baryon exchange (Fig. 7). The nearest singularity of these diagrams in terms of the variable t is at $t = 4m^2$ (m is the mass of a nucleon), so it can be concluded[70,71] that the annihilation radius is smaller than half the Compton wavelength of a nucleon:

$$r_a \leqslant \frac{1}{2m} \approx 0.1 \text{ fm}. \qquad (2.1)$$

Martin's arguments, however, refer to point particles. In the cases of the nucleon and the antinucleon, with nonzero dimensions, the concept of an annihilation radius requires further definition. It is natural to expect that the annihilation radius is of the order of the size of a nucleon. In an analysis of the annihilation $\overline{\mathrm{N}}$N interaction on the basis of

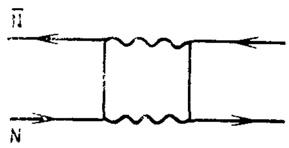

FIG. 7. Diagram of annihilation scattering with baryon exchange.

nonrelativistic quark models, for example, the value $R_{ann} = 0.7$ fm has been found.[32]

Nevertheless, potential models based on the assumption that the annihilation radius is small yield a fairly good description of the cross sections for elastic scattering and annilation. What is the explanation?

To analyze the situation we consider the example of the Dover-Richard potential[45]

$$V_{ann} = -\frac{V_0 + iW_0}{1 + \exp(r/a)}, \tag{2.2}$$

where $a = 0.2$ fm, $v_0 = 21$ GeV, and $W_0 = 20$ GeV. The value of the parameter a agrees with Martin's argument, but the constant W_0 is so large that the potential turns out to be strong, Im $V_{ann} \approx 150$ MeV, out to a distance $r \approx 1$ fm. Solving the scattering problem with potential (2.2), Dover and Richard[45] concluded that the absorption of antiprotons is concentrated in a narrow shell $1 \leqslant r \leqslant 1.1$ fm. A similar situation arises in the "Paris-potential" model,[43,44] where the imaginary part of the potential is determined by a set of diagrams with an exchange of two mesons ($\pi,\rho,\varepsilon,\omega$) in an intermediate state (Fig. 7). In other words, the imaginary part of the potential falls off as $\exp(-2m_N r)$. Again in this model, however, it is necessary to choose huge values for the constant in front of the exponential function, ≈ 10–100 GeV.

The list of examples of this type could be continued, but the general picture is already clear. By choosing an annihilation potential with a smaller radius, substituting it into the Schrödinger equation, and fitting the potential to the experimental data, we perform a "unitarization," as a result of which the effective annihilation radius reaches a normal value ≈ 1 fm. The behavior of the potential at small distances is thus important for describing the data at low energies. All the annihilation potentials which have been described in the literature—no matter how greatly they differ at small distances—are characterized by a value Im$V_{ann} \approx 100$ MeV at $r \approx 1$ fm (Refs. 32, 72, and 73).

Other phenomenological approaches to the description of annihilation have also been developed. For example, there are the coupled-channel model[74–77a,67,20] and the geometric approaches,[72,78] along which the nucleon and the antinucleon are treated as extended objects of the bag type, and the quantity ImV_{ann} is assumed to be proportional to the volume of the region in which these bags overlap. The most comprehensive description of the entire set of experimental data in the coupled-channel model is given by Timmers *et al.*[20] They also estimate the range of the effective nonlocal potential corresponding to this model. As expected, this range turns out to be large, $R \approx 1.5$ fm. The bag overlap model predicts Im$V_{ann} \approx 100$ MeV at $r = 1$ fm (Ref. 72). The model of boundary conditions also leads to a large annihilation radius.[77a]

Most of the potential models are successful in describing both the total and differential cross sections for elastic $\bar{p}p$ scattering.[32,52] The angular distributions of the charge-exchange reaction are not taken into account in all the models; the Paris potential,[43] for example, is noticeably less successful than the Nijmegen potential in describing the data.[20,32] A simultaneous description of elastic scattering and charge exchange has been achieved on the basis of optical potentials fitted especially to the LEAR data, which have an annihilation radius $r_a \approx 1$ fm (Refs. 32, 79, and 80). As a rule, however, the potential models run into difficulties in attempts to describe more-detailed characteristics of the $\bar{N}N$ interaction, e.g., the energy dependence of ρ (Subsec. 1.3) or data on the polarization.[17] The spin dependence and the isospin dependence of the annihilation potential are also unclear.[81]

The PS-179 experiment[82,83] yielded an indication of a strong isospin dependence of the P-wave amplitude for $\bar{N}N$ scattering. It was found that in the annihilation of antiprotons at rest with ^{4}He and ^{3}He the cross-section ratio

$$R = \frac{\sigma(\bar{p}n)_{ann}}{\sigma(\bar{p}p)_{ann}} \tag{2.3}$$

is $R = 0.48 \pm 0.1$, i.e., smaller by a factor of nearly two than the same ratio in the case of annihilation with deuterium, $R(d) = 0.82 = 0.82 \pm 0.03$ (Ref. 84).

We know that the capture of an antiproton in liquid deuterium occurs from high-lying S states, while 92% of the annihilations in helium come from P and D levels.[85] The potential models predict values of the ratio R in the S wave in the interval 0.75–0.9; roughly the same value is predicted for R in the P wave; $R(P) = 0.76$ (Ref. 45; see also Fig. 8 of the present paper). In other words, models which ignore the isospin dependence of ImV_{ann} cannot explain the pronounced increase in the annihilation interaction in a state with an isospin $I = 0$ in the P wave—an increase which follows from the experimental data.[82–83]

In summary, the phenomenological models tell us that the characteristic distances for annihilation are the same as those for confinement. We thus clearly see just how compli-

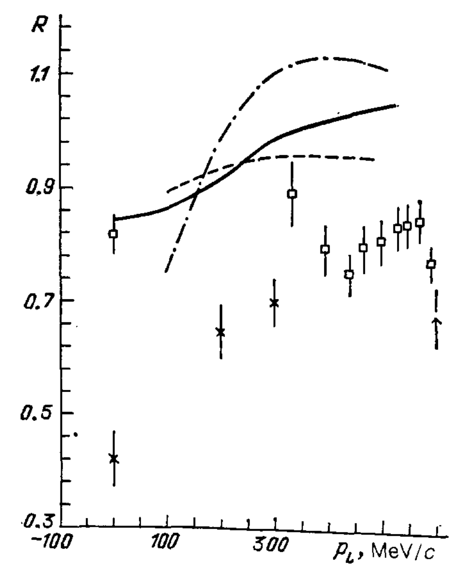

FIG. 8. The ratio $R = \sigma_a(\bar{p}n)/\sigma_a(\bar{p}p)$ as a function of the momentum according to data on the $\bar{p}^3$He and $\bar{p}^4$He interactions.[82,83] The curves are calculations from the potential model of Ref. 45 for (the dashed curve) the S-wave component (the dot-dashed curve), the P-wave component, and (the solid line) the combination S + P.

cated this problem of drawing a quark-gluon picture of annihilation actually is.

2.2. Exclusive annihilation channels

Table III shows data on the branching ratios of various exclusive channels for the annihilation of antiprotons at rest. We see that meson resonances are readily produced in annihilation: The branching ratio for the production of ρ mesons is $\approx 32\%$, that for ω mesons is 11%, and that for η mesons is $\approx 7\%$. Unfortunately, it is difficult to distinguish reliably certain channels because of large uncertainties regarding the subtraction of the background in the case of broad resonances (e.g., ρ and ε). The neutral modes for $\bar{p}p$ annihilation and for the annihilation of antineutrons have not been studied adequately; there are discrepancies between the results reported by different experimental groups.

Several general properties of annihilation can be described in comparatively simple phenomenological models.[20,32,44,45,86] For example, the following expression was proposed in Ref. 86 for calculations on reactions $\bar{p}p \to a + b$:

$$\bar{B}\,(pp \Rightarrow a + b) = C_{\bar{s}s}\,(2J_a + 1)(2J_b + 1)\,W_{ab}\,2^{-\delta_{ab}}q \times \exp[-A(s - s_{ab})^{1/2}]; \qquad (2.4)$$

here $J_{a,b}$ are the spins of the mesons, W_{ab} are the Clebsch–Gordan coefficients in terms of the isospin indices, $2^{-\delta_{ab}}$ is the Bose–Einstein factor, q is the momentum of the mesons in the c.m. frame, $s_{ab} = (m_a + m_b)^2$, and the coefficient $C_{\bar{s}s} = 0.15$ reflects the suppression of the production of pairs of strange quarks (this factor is absent in the case of channels which do not have strange quarks). Under the assumption that the annihilation occurs only through two-particle channels, expression (2.4) can reproduce the multiplicities of the various pion channels and their energy dependence. In order to reach an understanding of the annihilation mechanism, however, it is particularly important to study the dependence of the annihilation probability on the quantum numbers of the initial state. Several experiments in this direction have recently been carried out at LEAR.

From the methodological standpoint, it is possible to determine the quantum numbers of the initial state because the annihilation of antiprotons at rest goes through a stage of the production of an antiprotonic atom. Through detection in coincidence of the x rays emitted in transitions between levels in the $\bar{p}p$ atom and the charged particles from the annihilation, the experimentalists determine the orbital angular momentum of the state from which the annihilation occurs. Systematic studies with antiprotonic atoms "tagged" in terms of orbital angular momentum have become possible only in recent years in intense antiproton beams. It turns out that the annihilation probability in certain exclusive channels is highly dependent on the quantum numbers of the initial state of the $\bar{N}N$ system. We wish to emphasize that we are talking about a dependence which stems from the dynamics of the process, not some trivial consequences of selection rules. The clearest example comes from annihilation in the $\pi^+\pi^-$ and K^+K^- channels. The ratio of the probabilities for the annihilation of antiprotons at rest by these channels, $R = B(K^+K^-)/B(\pi^+\pi^-)$ has turned out to be quite different for the S and P states of the $\bar{p}p$ system. According to LEAR data,[87] the ratio $R_l = B_l(K^+K^-)/B_l(\pi^+\pi^-)$, where l is the orbital angular momentum of the $\bar{p}p$ system, is

$$R_0 = 34 \pm 3\%,\quad R_1 = 6.0 \pm 1.2\%. \qquad (2.5)$$

We thus have $R_1 \ll R_0$. The difference between R_1 and R_0 is due primarily to the fourfold suppression of the reaction $\bar{p}p \to K^+K^-$ as we go from the S state to the P state and also to some increase in the probability for the process $\bar{p}p \to \pi^+\pi^-$.

Here is the simplest, but not the only possible, explanation of the suppression of K^+K^- production for the P state.[88–91] In the reaction $\bar{p}p \to K^+K^-$, an additional $s\bar{s}$ pair—a strange quark and an antiquark— is produced. According to the standard $^3P_0$ model[92] (Subsec. 2.3), the $s\bar{s}$ pair is produced in the P state. It follows from parity conservation, however, that in annihilation from a P initial state the K^+K^- system should have an orbital angular momentum $l = 0, 2$. Accordingly, the annihilation $\bar{p}p \to K^+K^-$ is suppressed from P states.

A selectivity in terms of quantum numbers of this type, which does not stem from conservation laws or discrete symmetries, is called a "dynamic selection rule." Yet another interesting manifestation of these rules is the "$\pi\rho$ puzzle."[93,94] It has been observed that in the S state the reaction $\bar{p}p \to \pi\rho$ comes from the $^{13}S_0$ state in 95% of the cases[93] (here we are using the standard spectroscopic notation $^{2I+1,\,2S+1}L_J$).

The experimental data thus tells us that the production of the $\pi\rho$ system is dominated by initial states with an isospin $I = 0$ and a charge parity $C = -1$. As a result, the observed relation $W(\pi^\pm\rho^\mp) \approx 2W(\pi^0\rho^0)$ between the branching ratios for the charged and neutral modes of the $\pi\rho$ system differs by a factor of nearly two from the predictions of the phenomenological statistical model of annihilation,[86] according to which we have $W(\pi^\pm\rho^\pm) = 5W(\pi^0\rho^0)$.

The dynamic selection rules are not exhausted by these examples; a similar contradiction of the predictions of the statistical models is observed in other exclusive channels, e.g.,[95]

$$W(N\bar{N}(^{11}S_0) \to \pi A_2) = (3.6 \pm 0.9)\,W(N\bar{N}(^{33}S_1) \to \pi A_2). \qquad (2.6)$$

Data on the annihilation of stopped antiprotons into two neutral particles ($\pi^0\pi^0$, $\pi\rho$, $\gamma\pi^0$, etc.) were recently obtained at LEAR and KEK. An interesting result was found in measurements of the probability of the annihilation[30,96] $\bar{p}p \to \pi^0\pi^0$, which can occur only from states with an odd angular momentum. It follows from the data of Refs. 30 and 96 that the ratio of the probability for annihilation from the P state to all two-pion channels is

$$R_P(\pi\pi) = \frac{3W(\pi^0\pi^0)}{W(\pi^+\pi^-) + W(\pi^0\pi^0)} = 18 \pm 2\%\ (\text{Ref. } 30), \quad = 23 \pm 4\%\ (\text{Ref. } 96). \qquad (2.7)$$

The value $R_p \approx 20\%$ seems unusually high in comparison with estimates of R_p found in other experiments [e.g., $R_P = 8.6 \pm 1\%$ (Ref. 87) or $R_p \leqslant 6\%$ (Ref. 97)] or if we work from arguments concerning the similarity of the cascade processes in the $\bar{p}p$, π^-p, and K^-p atoms.[98] The π^- and K^- mesons stopped in liquid hydrogen annihilate from S states in more than 90% of the cases.[10,99] Annihilation from states with $l > 0$ is strongly suppressed by virtue of the

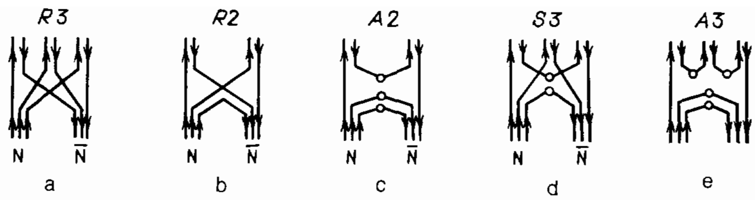

FIG. 9. Quark diagrams of $\bar{\mathrm{N}}\mathrm{N}$ annihilation.

Stark mixing of levels with different orbital angular momenta in the electric field produced by neighboring hydrogen atoms. As a result of the mixing, the π^- and K^- mesons which fill states with $l>0$ are absorbed through wide S states. In the antiprotonic hydrogen atom the picture is slightly different: The P levels themselves have a fairly large annihilation width,[34] and cascade calculations[100] are capable of reproducing the value $R_p \approx 20\%$. It must be kept in mind, of course, that this value of $R_p \approx 20\%$. It must be kept in mind, of course, that this value of R_p refers to the rare two-pion channel and could not be the same as the values for other annihiliation channels.

We have presented some new experimental results on the annihilation process. We turn now to attempts to derive a theory for this phenomenon.

2.3. Quark models of annihilation

Figure 9 shows the basic quark diagrams which describe annihilation. Two interrelated questions arise in an analysis of these diagrams: (a) What is the hierarchy of the various diagrams? In other words, what are their relative contributions to the annihilation cross section? (b) What is the effective operator corresponding to the annihilation of a quark-antiquark pair? In the original papers[101–103] it was suggested that the annihilation cross section was dominated (to an extent $\approx 80\%$) by a quark restructuring diagram and that the diagram with a single $q\bar{q}$ annihilation vertex and with the production of two mesons [Fig. 9(b)] made a correction of 10–20%. That conclusion, however, leans heavily on some assumptions regarding the form of the annihilation operator, the choice of the wave functions of the baryons and the mesons, and the incorporation of the interaction in the initial and final states (see the reviews of Refs. 73,88,104,105). The experimental data do not permit the unambiguous conclusion that three-particle annihilation plays a dominant role, since two-particle processes involving the production and subsequent decay of broad resonances (of the type $\varepsilon \Rightarrow \pi\pi$) are difficult to distinguish from genuine three-particle reactions. It is thus not possible to refute the suggestion that two-particle processes are actually dominant.[65,86]

Two basic types of the vertex operator $\hat{0}$ for the $q\bar{q}$ annihilation are being discussed: the $^3\mathrm{P}_0$ and $^3\mathrm{S}_1$ models (Fig. 10). The $^3\mathrm{P}_0$ model corresponds to the annihilation of a $q\bar{q}$ pair into gluon states with the quantum numbers of the vacuum, while the $^3\mathrm{S}_1$ vertex describes the conversion of a $\bar{q}q$ pair into one gluon and its subsequent absorption by another quark or antiquark:

$$\hat{O}(^3\mathrm{S}_1) = \lambda_S \chi_f \chi_C^{(8)} \mathbf{S}_{12} [\boldsymbol{\sigma}_3 \mathbf{q}]\, \delta(\mathbf{k}_1 + \mathbf{k}_2 - \mathbf{q}), \tag{2.8}$$

$$\hat{O}(^3\mathrm{P}_0) = \lambda_P \chi_C^{(1)} \chi_m\, (lm\, l\, -m\,|\,00)\, Y_{l,-m}(\mathbf{k}_1 - \mathbf{k}_2)\, \delta(\mathbf{k}_1 + \mathbf{k}_2),$$

where λ_S and λ_P are coupling constants, $\chi_f = (u\bar{u} + d\bar{d} + s\bar{s})\sqrt{3}$, χ_C is a color function, χ_m is a spin function, and $\mathbf{S}_{12}$ and $\boldsymbol{\sigma}_3$ are spin operators.

Just which vertex operator is dominant is not clear. Arguments in favor of the $^3\mathrm{P}_0$ model are successful in describing the decays of mesons and baryons[92, 106b] and the fact that such a model can be derived from lattice quantum chromodynamics (QCD) in the strong-coupling limit.[106] Adherents of the $^3\mathrm{S}_1$ model are oriented toward perturbative QCD, but it would be a bit optimistic to expect perturbation theory to be applicable at characteristic momenta ≈ 500 MeV/c.

A very simple calculation using the first approximation in terms of the vertex operators [the diagram in Fig. 9(b)] leads to substantially different predictions for the $^3\mathrm{S}_1$ and $^3\mathrm{P}_0$ models. In the $^3\mathrm{P}_0$ model, for example, the S-wave annihilation into two S-wave mesons is forbidden; i.e., the channel $\bar{\mathrm{N}}\mathrm{N}(^3\mathrm{S}_1) \to \pi\rho$ should be suppressed. That conclusion, however, contradicts experimental data (Subsec. 2.2). In the $^3\mathrm{S}_1$ model, this decay is allowed; in addition, in that model we find a natural explanation of why the reaction $\bar{p}p \to \pi\rho$ in the S state comes from the $^{13}\mathrm{S}_0$ state in 95% of the cases.[90,107] The $^3\mathrm{S}_1$ model, however, forbids the annihilation $\bar{\mathrm{N}}\mathrm{N}$ $(^{13}\mathrm{S}_1) \to \pi\,\mathrm{B}$ (1233 MeV), which is reliably seen experimentally. Furthermore, the ratios $W(\eta\rho^0)/W(\eta\omega) \approx 0.1$ and $W(\eta\rho^0)/W(\pi^0\rho^0) \approx 0.0075$ predicted by this model are greatly at odds with the experimental values, which are 1 and 0.5, respectively. The model of one-gluon exchange also forbids transitions from the S wave of $\bar{\mathrm{N}}\mathrm{N}$ to a two-pion state in the D wave, of the type $\bar{\mathrm{N}}\mathrm{N}(L=0) \to \pi A_2(l_f = 2)$ while in fact the channel $\bar{p}p \to \pi A_2$ is one of the dominant channels in the two-particle mode.

Attempts to describe annihilation by means of a hybrid model based on a superposition of $^3\mathrm{P}_0$ and $^3\mathrm{S}_1$ mechanisms in the first approximation in the annihilation interaction [Fig. 9(b)] have also been unsuccessful.[109]

The simplest dynamic approximations are thus inapplicable. It is necessary to move up to the next higher orders in the annihilation interaction, e.g., diagrams of the types in Figs. 9(c) and 9(d). Furthermore, since it is important to consider also the interaction in the initial and final states, as was pointed out in Refs. 110–112, the situation becomes extremely tangled. This entanglement is reflected in the fact that the $\pi\rho$ puzzle can be explained in either the $^3\mathrm{P}_0$ or $^3\mathrm{S}_1$ model.[107,113,114]

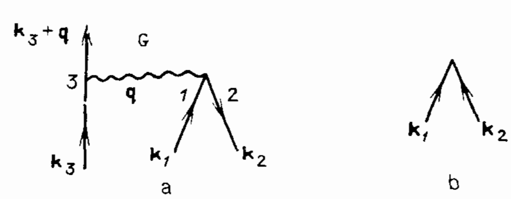

FIG. 10. Vertices of quark-antiquark annihilation in the (a) $^3\mathrm{S}_1$ model and (b) the $^3\mathrm{P}_0$ model.

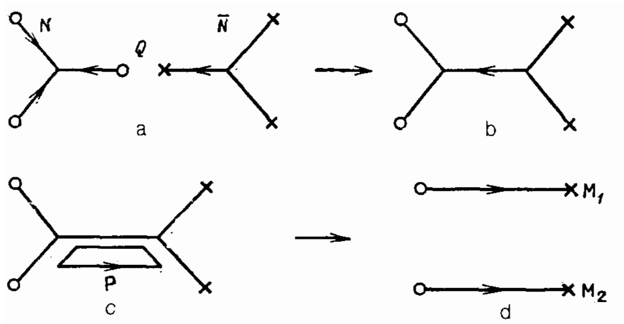

FIG. 11. Diagrams representing various stages of the annihilation $\bar{N}N \to M_1M_2$ in a string model (see the text proper) (a–d).

The most detailed analysis of the two- and three-particle annihilation channels has been carried out by a Tübingen group,[111] who considered all diagrams in Fig. 9. The best phenomenological description of the data was achieved on the basis of diagrams A2 and R3 with the $^3P_0$ annihilation vertex.

The hierarchy of the various annihilation diagrams was analyzed in the large-N_c limit in Ref. 115. The result

$$\sigma_{k,m}^{\mathrm{ann}} \xrightarrow[N_c \to \infty]{} e^{-AN_c} (\Upsilon^2)^{k+m} N_c^{2m} \qquad (2.9)$$

was found there; here k and m correspond to the number of pairs which are produced and which have annihilated, Υ^2 is the probability for the production or annihilation of a $q\bar{q}$ pair, N_c is the number of colors, and A is a numerical parameter.

It can be seen from (2.9) that if N_c is large then planar diagrams (diagrams with no intersections of quark lines) of types A2 and A3 are dominant. With $N_c = 3$ and $\Upsilon^2 = 0.23$ (Ref. 116), however, we find the opposite situation: $\sigma(\mathrm{R3})\cdot\sigma^{-1}(\mathrm{A}^3) = 3$ and $\sigma(\mathrm{R2}) \approx 2\sigma(\mathrm{R3}) \approx 2\sigma(\mathrm{A2})$ (Ref. 115). These results should be regarded as simply estimates, since they ignore the spin and isospin relations. The results may also change greatly when the interactions in the initial and final states are taken into account.

An interesting attempt to explain the $^3P_0$ mechanism on the basis of a string model was made in Refs. 117 and 118. Figure 11 shows the scenario proposed in Ref. 117 for the R2 annihilation mechanism. A nucleon (or antinucleon) is represented as a system of three quarks (or antiquarks) which are connected by strings at a central point: a node. As they move closer together, there is a coalescence of one string due to the annihilation of $q\bar{q}$ in the $^3P_0$ state [the process which is the inverse of the breaking of a string during the decay of mesons; Fig. 11(b)]. A perturbation of the vacuum generates a gluon loop [a little plaquette; Fig. 11(c)], which leads to a redistribution of the strings and to the formation of a pair of mesons [Fig. 11(c)].

Models of other annihilation mechanisms were also discussed in Refs. 117 and 118. Clearly, these attempts are only a first step toward incorporating gluon degrees of freedom. The presence of strong gluon fields can lead to the production of hybrid (qqq) mesons and glueballs in $N\bar{N}$ annihilation (Subsec. 3.4).

3. BARYONIUM AND HADRON SPECTROSCOPY

3.1. Mass spectrum and widths of baryonium

The term "baryonium" was first used by Chew[119] to label resonance states in q^2q^{-2} four-quark systems.[120–122] In Refs. 123–125 there was an independent discussion of the possible existence of bound states in the $\bar{N}N$ system near the threshold; in today's terminology they would be called "$q^3 q^{-3}$ six-quark systems with separate white clusters." Such states were called "quasinuclear mesons" in Refs. 124 and 125. Frequently, however, no distinction is drawn between these two types of states, and the term "baryonium" is applied to both.[71]

In principle, the $q^2 q^{-2}$ and $q^3 q^{-3}$ configuration may mix. There has been essentially no discussion of this mixing problem in the literature. All that we find are semiquantitative arguments, based on dual models, that this mixing is apparently small (see Ref. 126, for example). Specifically, if we take the $q^2 q^{-2}$ state as a zeroth approximation, then the energy shift due to an admixture of the six-quark state $q^3 q^{-3}$ should be at the level of the loop corrections in order of magnitude. In other words, it should be on the scale of the width of the $q^2 q^{-2}$ states. It is natural to assume that this shift would be of the same order of magnitude as the shift of the $\bar{q}q$ mesonic state due to $q^2 q^{-2}$ loop corrections. These corrections were estimated in Ref. 126 on the basis of a quark-gluon model and were shown to be small (see also Ref. 127).

3.1.1. Quasinuclear $\bar{N}N$ states

The starting point for the quasinuclear model of baryonium is the representation of the nucleon-nucleon interaction potential as a sum of one-meson exchanges,

$$V_{\mathrm{NN}}(r) = \sum_i V_i(r), \qquad (3.1)$$

where the index i corresponds to the exchange of different mesons ($i = \{\pi,\eta,\rho,\omega,\delta,\varepsilon\}$). The transition to the $N\bar{N}$ channel is made through a G transformation[128,129]

$$V_{\mathrm{ME}}^{\bar{N}N}(r) = \sum_i G_i V_i(r), \qquad (3.2)$$

where G_i is the G-parity of meson i. Because of the factors G_i, the real part of the $N\bar{N}$ potential is radically different from the nucleon-nucleon potential. The exchange of an ω meson leads to a strong short-range attraction in the $N\bar{N}$ system. Futhermore, the contributions of the π and ρ mesons add coherently in the tensor forces and also give rise to an attraction. The ultimate result is that about ten bound states arise in the $\bar{N}N$ system, with binding energies ranging up to hundreds of MeV, and there are several resonances above the threshold.[71,130,131]

These discussions, however, have not incorporated the strong annihilation in the $\bar{N}N$ channel. How does it affect quasinuclear levels?

The simplest estimate of the annihilation width of an S level is given by the well-known perturbation-theory formula.[35,132]

$$\Gamma_a^S = (v\tilde{\sigma}_a)_{v\to 0} |\psi_S(0)|^2, \qquad (3.3)$$

where $\tilde{\sigma}_a$ is the purely annihilation cross section with the one-boson-exchange potential switched off. Lacking a microscopic theory for annihilation, we cannot find a reliable estimate of $\tilde{\sigma}_a$. The observed annihilation cross section $v\tilde{\sigma}_a \approx 50$ mb could evidently serve as an upper limit on $v\tilde{\sigma}_a$, while a lower estimate would be $v\tilde{\sigma}_a \approx 4\pi r_a^2 \approx 1.2$ mb, if we assume $r_a = 1/2m \approx 0.1$ fm (Refs. 70 and 71). As was shown in subsec. 2.1, however, the annihilation radius is definitely greater than $1/2m$, so the latter estimate is far too low. It should be kept in mind that the cross section $\tilde{\sigma}_a$ is determined not by the annihilation radius itself but by the scattering length corresponding to the annihilation interaction. The scattering length is not necessarily small, even if $r_a \to 0$.

Estimating the weight function at the origin of coordinates is considerably simpler. For example, on could use the known expression[133]

$$|\psi_S(0)|^2 = \frac{m}{4\pi}\left\langle\frac{dV}{dr}\right\rangle \approx \frac{m}{4\pi}\frac{1\,\text{GeV}}{1\,\text{fm}}. \qquad (3.4)$$

Substituting (3.4) into (3.3), and setting $v\tilde{\sigma}_a = 1.2$ mb, we find the estimate $\Gamma_a^S \approx 50$ MeV, which is definitely on the low side. As a more reliable estimate of the annihilation width we could use

$$\Gamma_a^S \geqslant 100 \text{ MeV}. \qquad (3.5)$$

Another way to derive (3.5) is described in the review by Badalyan *et al.*[76] They also cite the papers in which (3.5) was derived on the basis of specific annihilation models. Estimate (3.5) is strikingly different from the value $\Gamma_a^S \approx 2$–7 MeV given in the review by Shapiro.[71] It is clear from this discussion that it is not possible to find a value $\Gamma_a \lesssim 10$ MeV under realistic assumptions regarding the annihilation cross section and the wave function. Estimate (3.5) remains valid for levels with $l > 0$, as was shown in Ref. 76. Consequently, if levels of quasinuclear baryonium do exist, they must have large annihilation widths.

3.1.2. Diquonium

We turn now to the properties of four-quark baryonium: q^2q^{-2} diquonium.[134] The conclusion that such states may exist may be regarded as a consequence of the existence of mesonic and baryon Regge trajectories with identical slopes α':

$$M_l^2 = M_0^2 + \frac{l}{\alpha'}. \qquad (3.6)$$

At large l the mesonic and baryon resonances lying on Regge trajectories have a string-like configuration with color charges at the ends of the string. The slope of the trajectory is inversely proportional to the string tension ν:

$$\alpha' = \frac{1}{2\pi\nu}. \qquad (3.7)$$

This tension is in turn proportional to the color charge:

$$\nu = \text{const}\cdot(F_c^2)^{1/2}. \qquad (3.8)$$

If the meson resonances correspond to a rotating string with triplet color charges at its ends which carry a quark and an antiquark, the baryon resonances should then correspond to the qq^{-2} configuration, where the diquark has the same color charge as the antiquark. If the quark in this system is replaced by an antiquark, we obtain the diquonium state with color-triplet diquarks ($F_c^2 = 4/3$).

Such states are called "T-baryonium"[122,134] or "T-diquonium." States which contain color-sextet diquarks ($F_c^2 = 10/3$) are called "M-baryonium" or "M-diquonium." Various models have been used to describe the spectrum of diquonium "a dual model"[134] (D), a quark-gluon model[126] (QG), extended rotating bags[135,106a] (ERB), and a relativistic model of a QCD string with a spin-orbit coupling[136] (RS).

States which contain color-triplet diquarks $(D)_{00}(I=0,\ S=0)$ and $D_{11}(I=1,\ S=1)$ are strongly coupled with the $B\bar{B}$ channel. The rupture of a triplet string results most probably from the production of a quark and an antiquark, which, combining with the diquarks q^2 and q^{-2}, can easily form a baryon-antibaryon system. M-diquonium is weakly coupled with the $\bar{N}N$ channel, since the breaking of a sextet string could result only from the production of several $q\bar{q}$ pairs from vacuum.

As an example, Table VI compares predictions of the masses of T-baryonium, q^2q^{-2}, derived on the basis of various models. The states which are obtained in the ERB model and the RS model when $\mathbf{l}$, $\mathbf{S}_1$, and $\mathbf{S}_2$ are combined into the angular momentum $\mathbf{J} \leqslant \mathbf{J}_{max} = \mathbf{l} + \mathbf{S}_1 + \mathbf{S}_2$ must be placed on daughter Regge trajectories in the D and QG models. Consequently, such resonances should have the masses as states lying on the main trajectory for a given l in the D and QG models. In the RS model, the masses of such systems will be different because of the spin-orbit coupling.

The spectrum of q^2q^{-2} states is thus a very rich one in all versions of quark models: There should be at least ten levels near the $\bar{N}N$ threshold.

Among the q^2q^{-2} states there must be some exotic ones with an isospin $I = 2$. According to the predictions of the RS model, the lightest $I = 2$ state has a mass of about 1.8 GeV (Ref. 136).

The states of T-Diquonium which lie above the $\bar{N}N$ threshold easily decay by the $B\bar{B}$ channel through the rupture of a triplet string. It has thus been assumed from the outset that such states should be broad ($\Gamma \geqslant 100$ MeV). On the other hand, it has been assumed that states which lie below the $\bar{N}N$ threshold may be narrow. The reason is that the rupture of a string in this case should result primarily from the formation of a q^2q^{-2} pair and the production of two diquonia. As follows from a description of hadron production processes in string models,[137] the probability for the rupture of a string accompanied by the formation of a pair of diquarks is about an order of magnitude lower than the probability for the corresponding rupture accompanied by the production of a $q\bar{q}$ pair. Consequently, if T-diquonium decayed only by virtue of the rupture of a string then the width of "light" diquonium ($M < 2m_N$) would be about an order of magnitude smaller than the widths of ordinary $q\bar{q}$ mesons.

There is, however, a cascade decay of the higher-lying orbital states of diquonium to low-lying states, through the emission of a pion by a diquark. To estimate the typical possibility for this process, we consider the decay $\Delta \to \pi N$ in a diquark model, in which it can be associated with a $D_{11} \to \pi D_{00}$ transition. Since the width of the Δ resonance is $\Gamma_\Delta = 120$ MeV, a transition between diquarks occurs over a time characteristic of the strong interaction ($\tau \approx m_\pi^{-1}$). The difference between the energies of the levels of T-diquonium with orbital angular momenta differing by one is, for a mass M = 1.5–1.9 GeV,

$$\Delta M \approx \frac{1}{2\alpha' M} \approx 0.25 - 0.35 \text{ GeV}$$

(α' is the slope of the Regge trajectory of T-diquonium). In

TABLE VI.

a) Spectra of the D_{00}-D_{00} system in various models $[D_{IS} = (q^2)_{I=S=0}]$ (the masses are in GeV).

l^P	$I^G(J^P)$	DRM	ERB	QG	RS
0^+	$0^+(0^+)$	—	—	1.26	—
1^-	$0^-(1^-)$	1.68	1.50	1.63	1.28
2^+	$0^+(2^+)$	1.92	1.76	1.93	1.70
3^-	$0^-(3^-)$	2.13	2.05	2.19	2.02
4^+	$0^+(4^+)$	2.33	2.29	2.42	2.30
5^-	$0^-(5^-)$	2.51	2.51	2.63	2.54

b) Spectrum of the $|D_{00}-\bar{D}_{11}\rangle \pm |\bar{D}_{00}-D_{11}\rangle$ system.

l^P	$I^G(J^P)$	RS	QG	DRM	ERB
0^+	$1^\pm(1^+)$	—	1.46	—	—
1^-	$1^\pm 0^-$	1.72	—	—	
	1^-	1.64	—	—	1.72
	2^-	1.46	1.79	1.89	
2^+	$1^\pm 1^+$	2.08	—	—	
	2^+	2.00	—	—	2.01
	3^+	1.84	2.06	2.13	
3^-	$1^\pm 2^-$	2.39	—	—	
	3^-	2.31	—	—	2.28
	4^-	2.15	2.31	2.34	
4^+	$1^\pm 3^+$	2.65	—	—	
	4^+	2.57	—	—	2.52
	5^+	2.41	2.53	2.54	

c) Spectrum of the $D_{11}-\bar{D}_{11}$ $(I=0,1,2)$ system.

S_u	J^P	RS	QG	DRM	ERB
$l^P=0^+$ 0	0^+	—	0.705	—	—
1	1^+	—	1.26	—	—
2	2^+	—	1.63	—	—
$l^P=1^-$ 0	1^-	1.66	—	—	1.86
1	0^-	1.90	—	—	
	1^-	1.81	—	—	1.90
	2^-	1.64	—	—	
2	1^-	2.11	—	—	
	2^-	1.94	—	—	1.94
	3^-	1.68	1.93	2.11	

other words, it is approximately the same as the difference between the masses of the Δ and the nucleon. We would thus naturally expect that cascade transitions of T-diquonium would also occur over a time $\tau \approx m_\pi^{-1}$, until the $l=0$ ground state forms. In this ground state, diquarks (not separated by a centrifugal barrier) easily convert into a pair of ordinary mesons. Consequently, one could hardly expect narrow states of T-diquonium even below the $\bar{N}N$ threshold.

The expected widths of the levels of T-baryonium near the threshold are thus $\geqslant 100$ MeV. In addition, it has been predicted that there are a large number of such states (Table VI). For this reason, four-quark states could hardly lead to narrow structural features in the cross sections. Nevertheless, the presence of states near the threshold with large spins (up to $J=4$) might be one possible reason for the strengthening of the $\bar{N}N$ interaction in the higher-order partial waves (Sec. 1). A P-matrix analysis may prove useful for observing q^2q^{-2} baryonium.[138]

3.2. Search for narrow states in the $N\bar{N}$ system

The history of the search for narrow states in the $N\bar{N}$ system could legitimately be called dramatic. By the mid-1970s, research had revealed many indications of the existence of a large number of $N\bar{N}$ levels with small widths (see the reviews of Refs. 71, 125). As new experimental data accumulated, however, the faith in the existence of narrow baryonium gradually crumbled, and by now all the narrow peaks observed previously can be regarded as disproved (see Refs. 94, 139, and 140 for a review of the current situation regarding the search for baryonium).

The most obvious candidate for the role of baryonium has been the so-called S meson, with a mass of 1936 MeV, which was first observed in 1974 (Ref. 141) as a narrow peak ($\Gamma \approx 10$ MeV) in the total cross section for the $\bar{p}p$ interaction, which stood out from the background by (5–7) σ. This resonance was subsequently observed in other experiments,[142–145], although its characteristics have differed slightly according to the results of the different groups. Structure near 1936 MeV has also been seen in the annihilation cross sections and in certain exclusive channels.[142,143,145]

Since 1980, however, experimentalists have not seen any structural features in either the total or annihilation cross sections.[146–149] Furthermore, no structural features have been seen in recent LEAR experiments designed especially for precise measurements of σ_t and σ_{ann} in the vicinity of the S meson at steps of 5 MeV/c (Refs. 22, 37, and 150). It is asserted that at a width $\Gamma = 3.5$ MeV the value of $(\sigma\Gamma)$ does not exceed 2 mb·MeV (Ref. 150).

Unconditional proof of the existence of bound $N\bar{N}$ states would be the observation of narrow lines in the spectrum of γ rays accompanying the annihilation of stopped antiprotons in hydrogen and deuterium.[71,151,152] The source of monochromatic γ rays would be radiative transitions from the states of a $p\bar{p}$ atom to baryonium levels: $(\bar{p}p)_{atom} \to \gamma + B$. The typical energies of the γ rays in such transitions should be of the order of 100 MeV; theoretical estimates of the transition intensity $B_\gamma = \Gamma_\gamma/\Gamma_a$ yield 10^{-4}–10^{-2} of the total annihilation probability.[71,151,152] A search for monochromatic γ lines is exceedingly complicated, since it is necessary to distinguish a weak signal against the background of a huge number of γ rays from the decay of π^0 mesons. Over the years, the statistical base from such experiments has grown. In the experiments of Refs. 153–155, which were carried out in 1978–1984, four narrow lines corresponding to bound-state masses of 1210, 1638, 1894, and 1771 MeV were observed. This result was perceived as nearly conclusive proof of the existence of baryonium, although the statistical significance of the observed signals was low (at the 3σ level).

The faith in these results eroded after experiments[156,157] at the KEK accelerator (in Japan), with a statistical base an order of magnitude greater than that of the experiments of Ref. 153, yielded no narrow lines with a statistical significance greater than 4σ. It is true that some narrow peaks with intensities $\approx 10^{-4}$ were found,[157] but their statistical significance was definitely low: at the level of $(2$–$3)\sigma$.

Negative results also emerged from an experimental search for discrete γ lines carried out at LEAR.[158–159] Negative results were also found by the Bakenstoss group,[159] which had previously reported the observation of monochromatic γ lines.

Studies of the inclusive spectra of π and K mesons also failed to reveal discrete lines.[160,161] Such lines would have corresponded to the production of baryonium in the reaction

$$(\bar{p}p)_{atom} \to \pi^{\pm}(K^{\pm}) + B^{\mp}. \tag{3.9}$$

The limitations on the transition intensities are[139] $B_{\pi^\pm B^\mp} < 8\times10^{-4}$ for the mass region $1000 \leqslant m_x \leqslant 1670$ MeV and $B_{K^\pm B^\mp} < 1{,}9\times10^{-4}$ for $1040 \leqslant m_x \leqslant 1280$ MeV.

In Ref. 162 there was a report of the experimental observation of a sharp peak in the cross sections for the reaction $\bar{p}p \to K^+K^-$, with a mass of 1940 ± 20 MeV and a width $\Gamma \approx 40$ MeV. Measurements by another group,[163] however, revealed a completely smooth energy dependence of the cross section for this annihilation channel.

The present status of experiments carried out to search for narrow baryons can be summarized as follows:

1. No bound states or narrow resonances have been found in the $N\bar{N}$ system.
2. If such states do exist, they would have to satisfy one of the following conditions: (a) They are wide ($\Gamma > 50$ MeV). (b) They lie very close to the $N\bar{N}$ threshold (within ± 20 MeV). (c) They have a small production probability ($\leqslant 10^{-4}$).

3.3. Wide $N\bar{N}$ resonances

In the shadow of the dramatic events associated with the discovery and dismissal of narrow baryonia, the fate of broad resonances in the $N\bar{N}$ system went comparatively unnoticed. These resonances were first observed in 1970 in measurements of the total cross sections of the $\bar{p}p$ and $\bar{p}d$ interactions.[164] Simultaneous measurements of these two cross sections made it possible to determine the cross sections for interactions in states with isospins $I = 0$ and $I = 1$. As can be seen in Fig. 12, two broad resonances are observed in the $I = 1$ channel, with masses of 2190 ($\Gamma = 85$ MeV) and 2350 MeV ($\Gamma = 140$ MeV), while in the $I = 0$ channel one observes a single resonance with a mass of 2375 MeV ($\Gamma = 190$ MeV). Similar structural features were subsequently observed in the annihilation cross sections,[165] in elastic scattering,[166] and in a charge-exchange reaction.[167]

Broad structural features ($\Gamma \approx 150$–300 MeV) have also been seen in the region 2100–2300 MeV in a phase-shift analysis of the two-particle reactions $\bar{p}p \to \pi^+\pi^-$, $\pi^0\pi^0$ (Refs. 168 and 169) and in a phase-shift analysis of the $(\bar{p}p)$ system produced in the reaction $\pi^-p \to (\bar{p}p)n$ (Ref. 170).

Although these resonances are not presently included in the main part of Rosenfeld's tables, so far no one dismisses them.

A recent study of the annihilation of antiprotons with deuterium revealed two new wide states: $\xi(1480)$, with the quantum numbers $J^{PC} = 2^{++}$, $I^G = 0^+$, $\Gamma = 116 \pm 9$ MeV; and X(1110), with $J^{PC} = 0^{++}$ or 2^{++}, $I = 0$, and a width $\Gamma = 111 \pm 8$ MeV (Refs. 171–173). There is particular interest in the state $\xi(1480)$. It has been observed both in the difference between the energy spectra of π^+ and π^- mesons from $\bar{p}d$ annihilation[171] and in the effective-mass spectrum of the four-pion system in the reaction[172]

$$\bar{p} + d \to \pi^- + \xi(1480) + p_S, \quad \xi(1480) \to 2\pi^-2\pi^+ \tag{3.10}$$

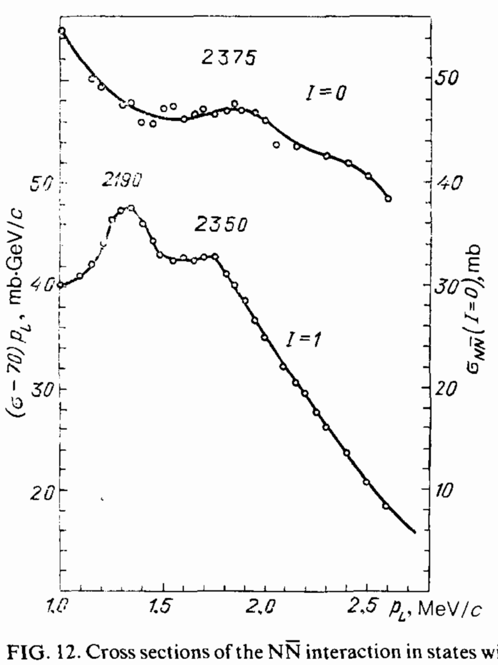

FIG. 12. Cross sections of the $N\bar{N}$ interaction in states with isospins 0 and 1 measured in the experiment of Ref. 164. For a clearer separation of the resonance structural features, the quantity $(\sigma - 70)P_L$ (in mb·GeV/c) is plotted along the ordinate. The ordinate scale at the right is for $\sigma_{N\bar{N}}$ $(I = 0)$.

The feature $\xi(1480)$ coincides with a clearly defined and broad peak which has been observed[175,176] in the reaction $\gamma\gamma\rightarrow\rho^0\rho^0$.

Several hypotheses have been advanced regarding the nature of $\xi(1480)$: (a) quasinuclear baryonium[177]; (b) an exotic q^2–$\bar{q}^2$ meson[178,179]; (c) a resonance in the $\rho^0\rho^0$ system.[180] Before we can seriously analyze any hypotheses concerning the nature of $\xi(1480)$, however, it must be clearly proved that this structural feature is not a kinematic feature stemming from an interference and rescattering of π^- and ρ^- mesons in the reaction $\bar{p}n\rightarrow 3\pi^-2\pi^+$ (Refs. 181 and 182).

In summary, although several candidates for broad resonances have been seen experimentally, their status has not been finally resolved. We do not rule out the possibility that some of these features may be exotic ($\bar{q}^2 - q^2$) mesons. Solving this problem is one of the most important tasks for antiproton complexes of the SUPERLEAR type with antiproton energies ≈ 10 GeV (see also Sec. 5).

3.4. Annihilation and the search for glueballs and exotic mesons

3.4.1. The problem of $E/f_1(1420)$ and $\iota/\eta(1440)$

When one or several $q\bar{q}$ pairs annihilate in an $\bar{N}N$ interaction, hard gluons can form in intermediate states. In the course of their bleaching, purely gluon states may arise: glueballs or mixed (hybrid) states of the $\bar{q}qg$ type (see, for example, the reviews of Refs. 183–185). Glueballs are flavor singlets, so it is expected that they should interact in an identical way with quarks of various types. They would preferably be sought in the mass spectra of hadronic states containing pairs of strange quarks [i.e., in channels with $K(\bar{K})$, $K^*(\bar{K}^*)$, η, and η' mesons]. The following considerations are of importance here: First, several models predict a large admixture of a gluon component in the η' meson. Second, SU(3) symmetry leads to certain relations among various two-particle decays of glueballs; and these relations could conveniently be tested in channels in which strange quarks are produced, because of the relatively favorable background conditions.

Although there are already a number of candidates for glueballs, their status is still unclear. One of the leading candidates for such states is the $\iota/\eta(1440)$ meson, with the quantum numbers $J^{PC} = 0^{-+}$. The reason for the interest in this entity is obvious: The nonet of pseudoscalar mesons is filled, and any new state with $J^{PC} = 0^{-+}$ would have to be some exotic state, possibly a radial excitation of one term of the nonet, a glueball, or a hybrid meson.

The history of the observation of this state dates back to early experiments with bubble chambers in antiproton beams. In a study of the annihilation of antiprotons at rest with protons, Armenteros *et al.*[5] observed a peak at an energy of 1420 MeV in the effective-mass spectrum of the $K_S^0K^{\pm}\pi^{\mp}$ system. A subsequent analysis[8] of these results revealed the quantum numbers of this resonance to be $J^{PC} = 0^{-+}$. In 1980, a peak was found in the reaction $\pi^-p\rightarrow K_S^0K^{\pm}\pi^{\mp}$ near 1420 MeV, but with the quantum numbers $J^{PC} = 1^{++}$. It was called[186] an "E meson" [$f_1(1420)$ in the new designation scheme]. The situation became more complicated when a state with a similar mass, $\iota(1440)$ [$\eta(1440)$ in the new scheme], but with quantum numbers 0^{-+}, was found in radiative decays of J/ψ particles.[187,188] Is the ι meson the same particle as the E meson which has been observed in collisions of hadrons, or are these two different states?

In an effort to resolve this E/ι problem, a numer of recent experiments have been carried out on the production of this resonance in both hadron-nucleon and e^+e^- collisions (see Refs. 183–185 and 189–199 for reviews). On the basis of all the evidence, it was concluded that E and ι are two different particles; E is an ordinary $s\bar{s}$ meson, while ι is a good candidate for a glueball.

Furthermore, in a recent experiment carried out by the MARK III group[199] on the $KK\pi$ and $\eta\pi\pi$ states in the radiative decays of the J/ψ meson, two structural features with quantum numbers 0^{-+} were observed in the mass interval 1/4–1.5 GeV: $M_1 = 1409 \pm 5$ MeV, $\Gamma_1 = 69 \pm 11$ MeV; $M_2 = 1499 \pm 9$ MeV, $\Gamma_2 = 138 \pm 25$ MeV. There are also indications that the ι/η structure is masking a 1^{++} state with a mass of 1420 MeV (Refs. 190 and 191).

The results of the early experiments, which seemed contradictory at first glance, can thus be explained in terms of the presence of several close-lying 1^{++} and 0^{-+} states in the mass region 1.3–1.5 GeV. These states are formed with different probabilities and different phases in different reactions at different energies.

The complexity of the problem obviously requires new experiments. The production of the ι/η meson is being analyzed in all the meson-spectroscopy experiments in the new phase of operation of LEAR. Apparently of greatest interest are the quasinuclear channels of the type $\bar{p}p\rightarrow M + \iota/\eta$, where $M = \pi,\eta,\rho,\omega$, or φ meson, and the ι/η resonance is in the $\eta\pi\pi$ or $K\bar{K}\pi$ channel. In particular, a study of these reactions would make it possible to test whether there actually are two pseudoscalar states in this region, which may differ is the magnitude of an $s\bar{s}$ admixture.[200]

3.4.2. Unusual $C/\rho(1470)$ and $U/M(3100)$ states

In a study of the mass spectrum of the $\varphi\pi^0$ system in the reaction $\pi^-p\rightarrow\varphi + \pi^0 + n$ at 32.5 GeV/c on the accelerator of the Institute of High-Energy Physics, the Lepton-Φ group observed a new resonance, C/ρ (1470), with the quantum numbers $I = 1$, $J^{RC} = 1^{--}$ and a width $\Gamma = 130 \pm 60$ MeV (Ref. 201). This resonance could not be an ordinary $\bar{q}q$ meson since it decays by a mechanism which is forbidden by the Okubo–Zweig–Iizuki rule and is not produced in the $\pi^-p\rightarrow\omega\pi^0n$ reaction. It is a good candidate for a (ds)–$(\bar{d}\bar{s})$ T-diquonium state with a latent strangeness, constituting a P-wave excitation of two strange spin-0 diquarks coupled by a triplet string.[202]

There is, however, an alternative interpretation of the $C/\rho(1470)$ peak: as a kinematic feature stemming from the chain[203] $\pi^-p\rightarrow\rho'n\rightarrow\bar{K}K\pi^0n\rightarrow\varphi\pi^0n$. In order to confirm the resonance nature of $C/\rho(1470)$ as diquonium with a latent strangeness it would be important to observe its strange partners $(qs) - \bar{D}_{00}$ and $(\bar{q}s) - D_{00}$ ($q = u,d$) with a mass in the vicinity of 1380 MeV and also the nonstrange partner D_{00}–$\bar{D}_{00}$, with a mass of about 1280 MeV (Table VI).

Recent CERN experiments in a hyperon beam[204] revealed three charge states of a narrow ($\Gamma \leqslant 30$ MeV) resonance with a negative strangeness (U^0, U^+, and U^-) and a mass of 3.1 GeV, which decayed into Λ, $\bar{p}$, and charged pions. The same states were found by the BIS-2 group from the Joint Institute for Nuclear Research, working at the accelerator of the Institute of High-Energy Physics.[205] That group also observed a state of this resonance with a double negative charge. In the classification of Ref. 134, this state would be designated M_S. Further research by the BIS-2 group[206] revealed narrow states with a positive strangeness, $\overline{U}^0/\overline{M}_s$, $\overline{U}^+/\overline{M}_s$, and $\overline{U}^{++}/\overline{M}_s$, in this region (they decay into $\overline{\Lambda}$, p and pions). In addition, a narrow resonance with a latent strangeness M_φ and a mass of $3255 \pm 10 \pm 30$ MeV, which decays into $\Lambda\bar{p}K^+ + \pi$ or $\overline{\Lambda}pK^- + \pi$, was observed.[207] The CERN group also observed a resonance with a mass of 3.4 GeV in the $\Lambda\bar{p}\pi^+\pi^-\pi^-$ channel, and the BIS-2 group observed states with a mass of 2.4 GeV in the $\overline{\Lambda}p$ system. It seems quite likely that these resonances could not be weakly decaying states.[208]

The presence of doubly charged states of a U/M_S meson is clear evidence that this meson has an isospin $I \geqslant 3/2$; i.e., it could not be an ordinary meson consisting of a $q\bar{q}$ pair. Furthermore, these resonances could hardly be regarded as states of T-baryonium,[208] $(q^2)_{3^*} - (\bar{q}^2)_3$, since such high-lying excitations could not be this narrow (see also Ref. 209 and Subsec. 3.1). Possibly a more plausible interpretation of the U/M mesons would be states of M-baryonium, in which qs and $\bar{q}^2$ diquarks have a color charge of 6 (Refs. 209–211). A color charge of 6 cannot be neutralized by the production of a single $q\bar{q}$ pair, so we have an explanation for the production of one or several mesons instead of a $B\overline{B}$ pair in the decay of the U/M resonance.

Chan and Tsou[209] have suggested identifying the states at 3.1 and 3.4 GeV as resonances with $J = 4^-$ and 5^+ which lie on a trajectory of M-baryonium with an isospin $I = 3/2$ and with $S = -1$. At the same time, there are arguments against that interpretation. First, why is it that only M-diquonium states with a spin of 4–5 tend to form in these experiments, and why do we not see states with a smaller spin? Second, how can we explain such a small width of the observed states, despite the large probability for their decay to similar states with smaller spins through pion emission? (See also Subsec. 3.1.)

There is also an interesting possibility for identifying the U/M resonances with triquonium: the 6q state $(q^2s)_8 - (\bar{q}^3)_{8'}$ in which there is a color octet tube between two octet 3q clusters[212,213] ($F_c^2 = 3$). In principle, a model of this sort would make it possible to achieve small spins for the U/M_S particles, in contrast with the model of M-diquonium.

It would of course be very interesting to see confirmation of the existence of these resonances and subsequent measurements of their quantum numbers. Antiproton beams with an energy $\approx$6–10 GeV would open up some new possibilities for studying such states, since in $\bar{p}p$ interactions the 3q and $3\bar{q}$ clusters are present even in the initial state, and they might convert into a $(3q)_8$–$(3\bar{q})_8$ color dipole through the exchange of a gluon. In this case, U/M resonances could form in the direct channel without strange quarks or with a latent strangeness. Strange U/M_S resonances might form, for example, in a reaction of the type $\bar{p}p \to K + U/M_S$.

4. QUARK-GLUON ASPECTS OF THE INTERACTIONS OF ANTINUCLEONS WITH NUCLEI

In research on the interactions of antiprotons with nuclei one can distinguish two groups of problems. The first includes such traditional questions for nuclear physics as research on the dynamics of the interaction and tests of the applicability of various models for hadron-nucleus scattering in describing antiproton data. The problems of the second group involve specific phenomena which arise in the annihilation of antiprotons in nuclear matter. Since a large amount of energy (≈ 2 GeV) is released in a comparatively small volume (≈ 1 fm^3) in the process, we are immediately faced with the interesting question of whether "hot droplets" of nuclear or hadronic matter or a quark-gluon plasma can form. Quark-gluon degrees of freedom may also turn out to be important in describing reactions of such a nature that they cannot occur in the case of free nucleons but do occur in interactions with a nucleus. Examples of such reactions might be the one-meson and zero-meson annihilations of an antiproton with a nucleus, of the type $\bar{p} + d \to \pi^- + p$ and $\bar{p} + {}^3He \to p + n$, which were first studied by Pontecorvo.[214]

Since the traditional set of problems has been discussed quite thoroughly in the existing reviews (e.g., Refs. 73 and 215–219), we will discuss in this section of the paper only certain questions in which the appearance of nontrivial effects is expected in $\bar{p}$ annihilation with a nucleus.

4.1. Strange particles as signals of the possible production of "hot droplets" of nuclear matter or of a quark-gluon plasma

The annihilation of antiprotons with nuclei occurs primarily at the periphery of the nucleus, where a "beam" of five or six π mesons with energies near the Δ_{33} resonance is formed. The mesons which are produced penetrate into the nucleus. Calculations based on cascade models[220,221] show that 40–50 % of the π mesons undergo a secondary interaction. Reactions involving the absorption and quasielastic scattering of annihilation π mesons lead to the transfer of a substantial energy to the nucleus. In the annihilation of 50-MeV antiprotons with ^{12}C and ^{208}Pb nuclei, for example, the values of the average energy transfer are respectively 350 and 730 MeV (Ref. 220).

How is this substantial amount of energy distributed in the nucleus? Does the entire nucleus undergo a "thermalization," or is there just a "local heating" of some small regions near the point of annihilation? Opinion is divided on this matter. Let us examine one possible scenario, which is discussed in Refs. 222–224.

It is assumed that the annihilation leads to the production of a fireball: a blob of highly excited nuclear matter with a baryon charge $B = 0$. It is the decay of this fireball into mesons which primarily determines the multiplicity of the mesons and their spectra. There is, however, a certain probability that the annihilation fireball will be able to absorb one or several nucleons in the vicinity of the annihilation point before the decay. Quark-gluon bubbles with a nonzero baryon charge form in the process. These bubbles can be identified with a highly excited multiquark bag. Inside the bag, quarks and gluons form a liquid of fermions and bosons. Beginning at $B = 3$, one can ignore the effects of the finite dimensions of the bag and assume that the parameters of the bubbles are those of a quark-gluon plasma. The temperature of the quark-gluon droplet is estimated to be $T = 160$ MeV, and it should depend weakly on B. The lifetime of such a fireball is estimated to be $\tau \approx 4 \cdot 10^{-23}$ s. In this fireball, the reactions $q\bar{q} \to s\bar{s}$ and $gg \to s\bar{s}$ form an admixture of strange quarks, which is some two to five times as large as the admixture of the sea of strange quarks in a nucleon. The intensification of the yield of strange particles in annihilation with nuclei in comparison with annihilation with free nucleons is thus regarded as one of the basic indicators of the formation of a quark-gluon droplet. It was accordingly predicted in Ref. 223 that the ratio K^+/π would increase by a factor of three, that the yield of strange baryons would reach 10% of the total annihilation probability, etc. Nevertheless, one should note that in this model there is an uncertainty which stems from the impossibility of reliably evaluating the fireball production probability.

How well do the conclusions of the fireball model correspond to experiment? The LEAR results do not support the predicted growth of the ration K^+/π; it remains essentially constant in annihilation with various nuclei up to lead.[225] The PS 179 experiment, however, detected an unusually large production of Λ hyperons in the annihilation of antiprotons with neon and helium nuclei at low energies[226] ($p_L \leqslant 600$ MeV/c). At these energies, the production of a Λ at a single nucleon would not be possible, since the threshold for the reaction $\bar{p}p \to \Lambda\Lambda$ is $p_L = 1435$ MeV/c.

It nevertheless turns out that in annihilation in neon the number of Λ particles produced is nearly twice the number of K_S^0 mesons. The ratio of the cross sections for Λ and K_S^0 production is $R = 2.3 \pm 0.7$ at 600 MeV/c (Ref. 226). Even stopped antiprotons have a large probability for producing Λ hyperons [the value of R at the threshold is $R(\Lambda/K_S^0) = 1.1 \pm 0.2$; Ref. 226]. A corresponding intensification of the production of Λ hyperons has been observed in the annihilation of antiprotons in tantalum at 4 GeV/c (Ref. 227).

An analysis[228] of these results showed that the elevated Λ yield could be explained under the assumption that these particles are formed in the rescattering of annihilation mesons.

If, on the other hand, we assume that all the Λ particles in the experiments of Ref. 226 were produced exclusively through the evaporation of fireballs with $B=0$ then we could find an upper limit on the probability for the production of such fireballs. In the case of $\bar{p}^4$He annihilation, this limit turns out to be[229] $\approx 18\%$.

Another anomaly associated with the production of strange particles in $\bar{p}A$ interactions was discovered back in 1973 (Ref. 230). It turns out that the momentum spectrum of the spectator protons in the reaction $\bar{p} + d \to p_s + KK + N\pi$ differs quite noticeably from a purely Hulthén distribution (Fig. 13). Remarkably, this signal disappears almost entirely if the $K\bar{K}$ trigger is not used. A theoretical analysis[231] has shown that this effect cannot be explained in terms of a rescattering of pions and kaons (see also Ref. 232). Unfortunately, the very fact that there is such a pronounced distortion of the momentum spectrum of the spectator nucleons has not been checked in other experiments.

The pronounced rescattering of the annihilation products in the nucleus thus leads to several interesting effects, e.g., an abundant production of Λ particles. This aspect of the $\bar{p}A$ annihilation has been utilized in experiments[233] on the production of heavy hypernuclei. It turns out that heavy hypernuclei (Bi, U) can be produced with a probability up to 10^{-3}–10^{-4} per annihilation event.

In principle, similar effects could occur in the annihilation of antiprotons with a higher energy, sufficient for the production of, for example, D mesons or J/ψ particles. The rescattering of charmed particles should lead to the production of a Λ_C^+ hyperon; the latter could become bound in a nucleus and form a supernucleus,[234] in the manner in which a Λ hyperon produced in the rescattering of annihilation mesons becomes bound in a hypernucleus. A search for hypernuclei in antiproton-nucleus annihilation is one of the interesting problems for complexes of the SUPERLEAR type.

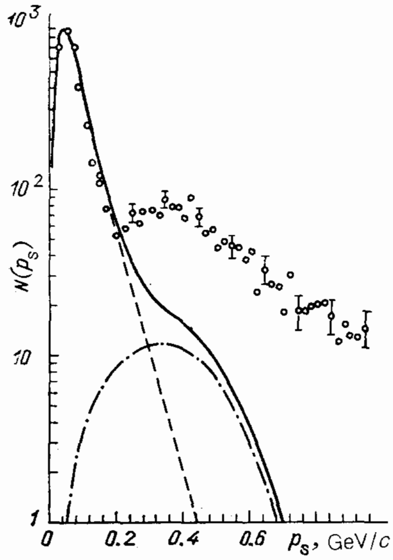

FIG. 13. Momentum spectrum of spectator protons in the reaction $\bar{p} + d \to p_s + K\bar{K} + N\pi$ (Ref. 230). Dashed and dot-dashed lines—Contributions of the impulse approximation (pole diagram) and of the mechanism with a final-state rescattering of mesons (triangle diagram); solid line—the coherent sum of these contributions.[231]

4.2. One-meson and zero-meson annihilation

Back in 1956, just half a year after the discovery of the antiproton, Pontecorvo[214] called attention to the possible existence of some unusual annihilation reactions which were forbidden at a free nucleon but which could occur in nucleons bound in nuclei. Among these processes are (1) an annihilation accompanied by the appearance of only one meson in the final state, e.g.,

$$\bar{p} + d \to \pi^- + p, \quad (4.1)$$

$$\bar{p} + d \to K^+ + \Sigma^-, \quad (4.2)$$

$$\bar{p} + d \to K^0 + \Lambda, \quad (4.3)$$

and (2) an annihilation with no mesons at all in the final state,

$$\bar{p} + {}^3He \to p + n, \quad (4.4)$$

$$\bar{n} + {}^3He \to p + p. \quad (4.5)$$

Unfortunately, we are obliged to state that Pontecorvo processes have not received much study so far. All that we have are data on the branching ratios for the annihilation of stopped antiprotons: $W(\pi^- p) = (0.9 \pm 0.4) \cdot 10^{-5}$ (Ref. 232) and $W(\pi^- p) = (2.8 \pm 0.3) \cdot 10^{-5}$ (Ref. 235). For the branching ratio for reaction (4.2) we have only an upper limit: $W(K + \Sigma^-) < 8 \cdot 10^{-6}$. Reactions (4.4)–(4.5), of zero-meson annihilation, have yet to be observed.

The Pontecorvo processes have attracted interest pri-

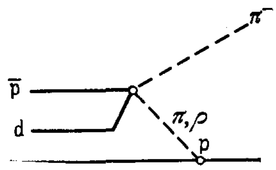

FIG. 14. Two-step mechanism (triangle diagram) of the reaction $\bar{p}d \to \pi^- p$.

marily because they should be sensitive to the high-momentum component ($k > m_N$) of the nuclear wave function $\Psi(k)$, in which quark-gluon degrees of freedom can play an important role (Ref. 236, for example). We will illustrate this assertion in the example of the two-step mechanism described by the triangle diagram in Fig. 14. In the annihilation of an antiproton with a nucleon of a deuteron, two energetic π mesons are produced with momenta $k \approx m_N$. One is then absorbed by the second nucleon. Energy-momentum cannot be conserved in each step of this process; the virtuality of the particles in the intermediate state is large ($> m_N$), so the amplitude for the reaction must be sensitive to small nucleon-nucleon distances $r \leqslant 1/m_N$ in the dueteron. Specific calculations[236] provide quantitative support for these arguments.

The probability for reaction (4.1) calculated for realistic wave functions of the deuteron and for a dipole form factor of the πNN vertex turns out to be of the order of 10^{-7}–10^{-6}, i.e., lower than the experimental value. The contributions from the absorption of ρ and ω mesons turn out to be small in comparison with that from the diagram with π-meson absorption and could not correct the situation.

One possible hypothesis is that the existence of quark degrees of freedom in the deuteron should lead to an increase in $\psi_d(k)$ at large k because of a tunneling of quarks through nucleon bags, resulting in a nonzero $\psi_d(r)$ at short range.[236] If this is the case, then the branching ratios for reactions (4.2) and (4.3) would have to be at the level $W(\bar{p}d \to K + \Sigma^-) = 8 \times 10^{-9}$ and $W(\bar{p}d \to K^\circ \Lambda) = 3 \times 10^{-7}$. These results are very different from the prediction of Ref. 224, where the branching ratio for reactions (4.2) and (4.3) was calculated in the model of the evaporation of a fireball with a nonzero baryon charge, and where the results $W(\bar{p}d \to K^+ \Sigma^-) = (7.8 \pm 0.8) \cdot 10^{-6}$ and $W(pd \to K^\circ \Lambda) = (8.1 \pm 0.9) \cdot 10^{-6}$ were found.

Brodsky[237] has suggested using Pontecorvo reaction (4.1) to test the predictions of perturbative QCD. Under the assumption that amplitude (4.1) can be represented as the product of the amplitude for a hard collision of quarks and of corresponding structure functions, which determine the distribution of valence quarks in the colliding nuclei, one can derive the following result from the quark counting rules:

$$\frac{d\sigma}{dt}(\bar{p}d \to \pi^- p) = \frac{1}{(p_T^2)^{12}} f(\theta_{c.m.}), \tag{4.6}$$

where $f(\theta_{c.m.})$ is some function of the scattering angle. Since it would be difficult to expect the quark counting rules to remain applicable at energies ≈ 1 GeV, it is suggested that the formalism of reduced QCD amplitudes be used. Brodsky[237] argues that formalism gives a correct description of the quark-gluon dynamics of the process beginning at $p_T \geqslant 1$ GeV.

The Pontecorvo reactions should be thought of as a useful tool for studying the high-momentum component of nuclear wave functions. An extensive experimental program has been planned for studying Pontecorvo reactions at LEAR and the OBELIX apparatus.[29]

5. PROSPECTS FOR PHYSICS RESEARCH IN LOW- AND INTERMEDIATE-ENERGY ANTIPROTON BEAMS

The prospects for the development of antiproton rings and possible programs for physics research at them are being discussed widely.[14–17] Since 1988, a second stage of experiments with an improved antiproton beam (see the Introduction) has been under way at CERN. The research is being carried out with apparatus of a new generation: large 4π detectors, designed for precise measurements of various exclusive (including multiparticle) channels for the interaction of antiprotons with protons and nuclei. Let us take a brief look at the basic aspects of the program for these new experiments and also some suggestions for further research.[14–17,238–241]

5.1. Experiments with tagged kaons for testing discrete symmetries (CP, CPT, T, and $\Delta S = \Delta Q$)

The high antiproton-beam intensity ($\approx 3 \times 10^6$ $\bar{p}$/s) which has been achieved at LEAR since its reconstruction is making possible precise measurements of the characteristics of the breaking of CP and T invariance in decays of kaons produced in the annihilation reactions

$$\bar{p}p \to \bar{K}^0 K^+ \pi^-, \ K^0 K^- \pi^+. \tag{5.1}$$

The branching ratio of these reactions is $\approx 0.2\%$ of the total annihilation probability. The sign of the charged kaon tells us unambiguously whether a K^0 or a $\bar{K}^0$ is produced. Measurements of the momenta and energies of the $K^\pm$ and $\pi^\pm$ mesons make possible a reliable identification of the reaction of the determination of the momenta of the K^0 and $\bar{K}^0$ mesons.[242]

This procedure of "tagging" kaons is exceedingly convenient for experiments carried out to measure the characteristics of CP breaking, since the K^0 and $\bar{K}^0$ contain identical K_L and K_S admixtures. In ordinary experiments, K^0 mesons are produced in hadron-nucleus collisions, and it is primarily K_L mesons which reach the detector.

Knowing the behavior of the branching ratios for the decays $K^0/\bar{K}^0 \to \pi^+\pi^-$ and $K^0/\bar{K}^0 \to \pi^0\pi^0$ as a function of the transit time of the neutral kaon, one can determine the parameters

$$\eta_{+-} = K_L/K_S \to \pi^+\pi^- = \varepsilon + \varepsilon'$$

and

$$\eta_{00} = K^0/\bar{K}^0 \to \pi^0\pi^0 = \varepsilon - 2\varepsilon'.$$

According to an estimate[242] based on an analysis of 10^{13} events in which antiprotons at rest underwent annihilation, there is hope that it will be possible even to improve somewhat the accuracy of the determination of the ratio $|\varepsilon'/\varepsilon|$ and the magnitude of the phase difference $\Phi_{+-} - \Phi_{00}$ (Table VII).

CPT invariance makes it possible to relate the phase shifts ε and ε' with other observable parameters,[243,244] ($\tau_+ \equiv \tau(K^+), \tau_- \equiv \tau(K^-)$). Consequently, this experiment

TABLE VII. Accuracies attainable in the testing of discrete symmetries in experiments with "tagged" kaons[239,245] ($\tau_+ \equiv \tau(K^+)$, $\tau_- \equiv \tau(K^-)$).

Parameter	Breaking	Limit	
		(a) existing	(b) expected
$\lvert\tau_+ - \tau_-\rvert / \lvert\tau_+ + \tau_-\rvert$	CPT	$(1.1 \pm 0.9) \cdot 10^{-3}$	$1.5 \cdot 10^{-4}$
$\delta(M_L - M_S)/\lvert M_L - M_S\rvert$	CPT	$4.1 \cdot 10^{-3}$	$1.2 \cdot 10^{-5}$
$\Phi_{00} - \Phi_{+-}$	CPT	$(10 \pm 5)^\circ$	<1
ε'/ε	CP	$(3.2 \pm 1.0) \cdot 10^{-3}$	$1.4 \cdot 10^{-3}$
$\lvert\eta_{+-0}\rvert^2$	CP	$<1.2 \cdot 10^{-1}$	$<10^{-3}$
$\lvert\eta_{000}\rvert^2$	CP	$<10^{-1}$	$<10^{-5}$
Re X	$\Delta S = \Delta Q$	$<2 \cdot 10^{-2}$	$6 \cdot 10^{-4}$
Im X	$\Delta S = \Delta Q$	$<2.6 \cdot 10^{-2}$	$7 \cdot 10^{-4}$
K^+e^+/K^-e^- ($\bar{p}p \to K^+e^+X$; K^-e^-X)	T, CPT or $\Delta S = \Delta Q$	—	10^{-3}

will also make it possible to find some new limits of the characteristics of the breaking of CPT invariance. We know (Ref. 239, for example) that regardless of whether CPT is broken the phase shifts measured in the decay $K_L \to 2\pi$ indicate a breaking of T invariance in the K^0–$\bar{K}^0$ system by virtue of unitarity.

It has been suggested that T breaking be tested directly by comparing the branching ratios for the transitions $\bar{K}^0$–K^0 and K^0–$\bar{K}$ in the reactions

$$\bar{p}p \to \bar{K}^0 K^+ \pi^-, \quad \bar{K}^0 \to K^0 \to (\pi^- e^+ \nu) K^+ \pi^-, \tag{5.2}$$

$$\bar{p}p \to K^0 K^- \pi^+, \quad K^0 \to \bar{K}^0 \to (\pi^+ e^- \bar{\nu}) K^- \pi^+. \tag{5.3}$$

The idea is based on the rule $\Delta S = \Delta Q$, which forbids the decays $K^0 \to (\pi^+ e - \nu)$ and $\bar{K}^0 \to (\pi^+ e^- \bar{\nu})$. A nonzero difference between the branching ratios for (5.2) and (5.3) would be a direct measure of T breaking. Interestingly, one can work from the magnitude of the effect and its time dependence also to determine whether CPT invariance or the selection rule $\Delta S = \Delta Q$ is broken.

5.2. Test of CPT invariance in the p-$\bar{p}$ system

The best limitation on the scale of a possible breaking of CPT invariance follows from the difference between the K_L and K_S mesons: The K^0–$\bar{K}^0$ mass eigenvalues may differ by an amount[245] $\Delta m/m \leqslant 6 \times 10^{-19}$. Nevertheless, the fundamental importance of CPT invariance requires an improved accuracy of tests of this invariance also, in baryon systems in which there is an upper limit on the CPT breaking parameter. This upper limit, which follows from the mass difference between p and $\bar{p}$, is[245] 6×10^{-5}. Three experiments have been proposed for improving this limit at LEAR: PS–189, PS–196, and PS–200 (Ref. 246).

In two experiments of an rf determination of the $\bar{p}$/p inertial-mass ratio, the cyclotron frequencies of antiprotons and protons (or of H^- atoms) revolving in an ultrauniform magnetic field will be compared. The PS–189 experiment uses an rf mass spectrometer about 1 m in diameter. It is believed that the natural linewidth is no greater than $10^{-6}\nu_C$. At the error level expected in the determination of the center of the peak, $\approx 10^{-3}$, and $\bar{p}$ and p masses could be compared within an error as small as 10^{-9}.

It is hoped that the same limit on $\Delta m/m$ will be achieved in the PS–196 experiment, which uses new magnetic-trap technology. In previous experiments with ultracold protons, a natural linewidth $\sim 10^{-9}\nu_C$ was achieved by this method. The success of this experiment will depend on whether it is possible to develop a procedure for cooling $\bar{p}$'s to energies $\sim 10^{-3}$ eV and capturing them in a trap.

An experiment (PS–200) to determine the gravitational mass of the antiproton uses a magnetic trap. Antiprotons with an energy $\sim 10^{-3}$ eV captured in the trap will subsequently be "launched" vertically upward into a drift tube. The duration of the upward flight of the $\bar{p}$'s, which depends on the initial velocity and the gravitational force, will be compared with the corresponding time for H^- ions. At the expected comparison error $\sim 1\%$, it will be necessary to launch $\sim 10^6$–10^7 $\bar{p}$'s and H^- ions into the tube. This accuracy will be sufficient to derive new limitations on the mass and interactions constant of the "graviphoton" - a hypothetical particle which is the supersymmetric counterpart of the graviton. The existence of such a graviphoton would result in an additional short-range attraction of a $\bar{p}$ toward the earth and a repulsion of a p away from the earth (see also Ref. 240).

5.3. Antihydrogen

There has recently been a widespread discussion[248–250] of Budker and Skrinskiĭ's idea[247] of producing, and studying under laboratory conditions, the simplest atom of antimatter: the antihydrogen atom ($e^+\bar{p}$). The primary difficulty along the path to the realization of this idea is that the antiparticle densities which have been achieved are still low. Among the several schemes[248] for producing $\bar{H}$ atoms, that which offers the most advantages is analogous to one which is used for electron cooling of proton beams (with a replacement of e^- by e^+, and of p by $\bar{p}$). At the existing intensities of antiproton and positron beams, the rate at which $\bar{H}$ atoms would be produced through spontaneous radiative capture, $e^+\bar{p} \to \bar{H} + \gamma$, would be only a few atoms per hour.[251] It has accordingly been suggested[252] that capture be induced with a laser beam. The rates at which $\bar{H}$ would be formed could be increased in this manner to thousands of atoms per second[248,250] or, according to the most optimistic estimates, to 10^6 atoms per second.[251]

An extensive program of physics research involving antihydrogen is being discussed.[248–250] One idea is to measure fundamental spectroscopic characteristics: the Rydberg constant, the $2S_{1/2}$–$2P_{1/2}$ Lamb shift, the hyperfine splitting of the 2S and 2P states, and the lifetimes of levels. The goal of these studies would be to test fundamental symmetries. In addition, there are plans for gravitational experiments with antihydrogen, and there are plans to use antihydrogen beams to produce polarized antiprotons.[248-250] A more remote possibility is that of studying the H-$\overline{\mathrm{H}}$ interaction and, in particular, finding an answer to the question of whether a metastable state exists in this system.[253,254] Progress in experiments with antihydrogen will depend greatly on whether it is possible to achieve the required probability for the formation of $\overline{\mathrm{H}}$ atoms. It has been suggested that the rate of $\overline{\mathrm{H}}$ formation at LEAR be studied by investigating the radiative capture of positrons by antiprotons.[251]

5.4. Search for and study of exotic hadrons

Antiproton beams open up several new opportunities for studying exotic hadrons, with masses ≈ 1–2 GeV: glueballs, mixed states of the $\bar{q}qg$ type, multiquark exotic entities of the q^2–$\bar{q}^2$ type, etc.[239] One of the primary advantages of the $\bar{p}p$ channel is that it is possible to study such states with known quantum numbers of the initial state, with the annihilation starting from the S or P level of the $\bar{p}p$ atom (Section 2). By selecting various final states, one can vary the background conditions to a great extent. Let us examine, for example, the S-wave $\bar{p}p$ annihilation by the (1) $\pi^0\pi^0\eta$ and (2) $\pi^0\pi^0\Phi$ channels.[239] In reaction (1), the quantum number of the final states should be $J^{PC}=0^{-+}$, $I^G=0^+$ or 1^- (the $\pi^0\pi^0\eta$ state could have quantum numbers 1^{--}). If there is an exotic resonance with quantum numbers $J^{PC}=1^{-+}$ in the $\pi^0\eta$ system, then it will be relatively easy to distinguish it from the background. The reason is that 2^{++} and 3^{-+} states in the $\pi\pi$ and $\pi\eta$ channels (where there are the resonances f, f', A_2, etc.) should form in the D wave with respect to the third meson and would probably be suppressed. The contribution of the 0^{++} state to the mass spectrum, where there is a comparatively narrow $\delta(960)$ resonance, could easily be distinguished from the contribution of the 1^{-+} resonance, which is expected to be heavier (one resonance in the $\pi^0\Phi$ system has already been found[201]).

Searches for resonances in the $\pi^0\Phi$ system in reaction (2) are facilitated by the circumstance that it is possible to avoid a large background from the $\pi\rho$ state. A study of the $\bar{p}p\to\pi^0\eta\Phi$ channel is extremely interesting in connection with searches for glueballs in the $\eta\Phi$ mass spectrum with quantum numbers $I^G=0^-$ and $J^{PC}=1^{+-}$, 0^{--}, 1^{--}, 2^{--}, etc.

In $\bar{p}p$ annihilation one can expect a large probability for the formation of exotic q^2–$\bar{q}^2$ states, if only because both diquarks will have been prepared beforehand. Such resonances should have a substantial effect onthe $\bar{p}p$ interaction in the S channel.[136]

It is also expected that $\bar{p}p$ annihilation should be a good source of glueballs and $\bar{q}qg$ hybrids (Refs. 185 and 239; see also Subsec. 3.4).

Several experiments will be devoted to spin effects in the $\overline{\mathrm{N}}\mathrm{N}$ interaction[255,256] and research in an antineutron beam.[257]

5.5. Some more-remote possibilities involving intense $\bar{p}$ beams in the momentum range 2–60 GeV/c

The outlook for further research in antiproton beams is linked with increasing the beam energy above the threshold for the production of charmonium ($p\geqslant 4$–5 GeV/c) or bottomonium ($p\geqslant 60$ GeV/c). Several suggestions for corresponding installations have been discussed in the literature:

1. SUPERLEAR: a storage ring for antiprotons with momenta from 2 to 10 GeV/c, a luminosity $L\approx 10^{32}$ cm$^{-2}\cdot$s^{-1}, and $\Delta p/p\approx 10^{-3}-10^{-4}$ (Ref. 241).

2. SUPERLEAR-2-$\bar{p}p$: a collider with an energy up to 10×10 GeV and $L\approx 10^{32}$ cm$^{-2}\cdot$s^{-1} (Ref. 241).

3. A Fermilab plan with $E(\bar{p})=2$–8 GeV, $N_{\bar{p}}\approx 5\cdot 10^7$ $\bar{p}$/s, and $\Delta p/p\approx 10^{-5}$ (Ref. 16).

4. The suggestion of a European Hadron Factory (EHF): an extracted beam of antiprotons with momenta of 3–10 GeV/c, $N_{\bar{p}}\sim 3\cdot 10^9$ $\bar{p}$/s, and $\Delta p/p\approx 5\cdot 10^{-2}$ (Ref. 258).

5. A suggestion of the Institute of Theoretical and Experimental Physics: an accelerating ring for accelerating protons and antiprotons to 40–50 GeV, with $N_{\bar{p}}\approx 10^7$ $\bar{p}$/s and $\Delta p/p\approx 10^{-5}$ (Ref. 238).

In installations of this type, with a high energy resolution (10^{-4}–10^{-5} and better), were to come on line it would become possible to take a new step toward the development of a spectroscopy of heavy quarkonia $\bar{c}c$ and $\bar{b}b$ (Refs. 238, 239, and 241). For example, such installations would make it possible to do the following:

a) Observe and study narrow charmonium levels $1^{+-}(^1P_1)$, $2^{--}(^3D_2)$, and $2^{-+}(^1D_2)$. Studies of these levels are beyond the capabilities of e^+e^- colliders.

b) Measure more accurately the total widths of all the narrow states of charmonium ($\Delta M<50$ keV).

c) Measure the probabilities for various multipole transitions in radiative decays of charmonium. In the decay $2^{-+}\to 1^{--}$, for example, one could work from the measured angular distributions of γ rays to distinguish the contributions of multipolarities M1, E2, and M3.

d) Study the spectra of mesons and baryons with heavy quarks.

e) Search for and study heavy narrow exotic resonances of the $\xi(2220)$ and U/M_S (3100) type and also glue balls and hybrid states.

New possibilities would also be opened up in research on discrete symmetries. At the SUPERLEAR-2, for example, it might be possible to study the characteristics of CP and CPT breaking in the Λ-$\overline{\Lambda}$ system, produced in the reaction $\bar{p}p\to\overline{\Lambda}\Lambda$ (to compare the lifetimes of Λ and $\overline{\Lambda}$ and also to compare the asymmetry parameters of their decays [α,β, and γ) through the $p\pi^-$ and $\bar{p}\pi^+$ channels]. The symmetric production of Λ and $\overline{\Lambda}$ in a $\bar{p}p$ collider would make it possible to reduce the systematic errors and to improve the measurement error to a level of 10^{-4}. The optimum beam momenta for these experiments would be in the interval 2–3 GeV/c (Ref. 241). If the installation had a luminosity $L\approx 10^{32}$ cm$^{-2}\cdot$s^{-1}, one could produce 10^6 events of the reaction $\bar{p}p\to\Lambda+\overline{\Lambda}$ per day. An experiment carried out in the existing beam at 1.546 GeV/c has demonstrated the feasibility in principle of this suggestion and has yielded the result $\overline{A}=(\alpha+\overline{\alpha})/(\alpha-\overline{\alpha})=-0.023\pm 0.05$ (Ref. 259).

As was mentioned in Ref. 260, measurement of the asymmetry $A=k([\mathbf{pq}]-[\bar{\mathbf{p}}\bar{\mathbf{q}}])$ in the reaction $p(\mathbf{k})+\bar{p}(-\mathbf{k})\to\Lambda+\overline{\Lambda}\to\pi^-+p+\pi^++\bar{p}$, where $\mathbf{p}$ and $\mathbf{q}$ are the momenta of the proton and the pion, at a level of 10^{-4}, might turn out to be important for testing models for the breaking of CP invariance (but see the more cautious discussion in Ref. 238).

Important progress might also be achieved toward the solution of several other QCD problems. In addition to studying the properties of heavy quarkonia and exotic hadrons, it might be possible to study such dynamic QCD predictions as the power-law falling-off with momentum of the cross sections for hard exclusive processes, the presence of amplitudes which break chiral symmetry, etc.[238–241]

Even this brief list, of by no means all the possibilities which would be opened up by the implementation of the suggestions for new intermediate-energy antiproton accelerators, demonstrates the exceedingly great research potential here. Since the antiproton rings which have been proposed would make it possible to achieve a mass resolution at least 100 times better than that at existing accelerators, we do not rule out the possibility that experiments at such installations might yield some unexpected results which would prove important for a next stage in the development of particle physics and nuclear physics.

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