Resonance Phenomena in Antiproton-Nucleus Scattering

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This Letter suggests that $\bar{p}$-nucleus scattering might display resonant phenomena at certain energies, arising from the orbiting of the $\bar{p}$ in a pocket of nuclear attraction extending beyond the region of strong absorption. This results in resonances of typical widths 15–150 MeV in the 180° excitation function and large polarizations for large-angle elastic scattering.

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There has been a great deal of experimental and theoretical interest[1] in the interactions of antinucleons ($\bar{N}$) with nucleons (N). Experimental work has thus far concentrated on the $\bar{p}p$ and $\bar{p}d$ systems: Elastic, charge exchange, and total cross sections have been measured in the low- and medium-energy region. At various times, there has been evidence for narrow structures in these elementary cross sections, although the experimental situation remains confused.[1] It is hoped that such structures may correspond to the long-sought "baryonium" states predicted in both quark and potential models.[1,2]

The corresponding interactions of antinucleons with nuclei [$\bar{p}$-$\mathfrak{N}$ for $A(\mathfrak{N}) \geq 3$] are essentially unexplored, either theoretically or experimentally. With use of currently available $\bar{p}$ beams at either CERN (Centre Européen de Recherches Nucléaires) or Brookhaven National Laboratory, it should be possible to investigate $\bar{p}$-induced nuclear reactions experimentally in the near future. The proposed Low-Energy Antiproton Ring (LEAR) facility at CERN will greatly increase the potential scope of such experiments. In view of these prospects, it is important to pose some theoretical questions. Among these are the following: (i) If there are narrow two-body resonances or bound states, how are they reflected in $\bar{p}$-$\mathfrak{N}$ scattering? (ii) Are there any new phenomena which could be anticipated in $\bar{p}$-$\mathfrak{N}$ scattering because of the unique features of the many-body system? We present here some preliminary results along these lines.

The first question also arises in π-nucleus scattering, where one looks for the effects of the Δ and N^* resonances inside a nucleus. One observes an energy shift and a considerable broadening of the Δ, even in light nuclei. Similarly, a two-body $\bar{N}N$ resonance such as the $S(1930)$, if it exists,[1] is Fermi and collision broadened and will not remain as a narrow peak in $\bar{p}$-$\mathfrak{N}$ total or annihilation cross sections. Difficult as this structure has been to observe in the $\bar{N}N$ system, it may be virtually unobservable in a nucleus. An interesting effect of broad $\bar{N}N$ resonances still remains, however, in that the $\bar{N}$ annihilation potential will be very strong, much deeper than for nucleon scattering at comparable energies. The $\bar{p}$-$\mathfrak{N}$ total cross section σ_{tot}, although varying smoothly with energy, can considerably exceed the naive geometrical limit $\sigma_G = 2\pi(R + \lambdabar)^2$, where $R = (\frac{5}{3}\langle r^2\rangle)^{1/2}$ is the equivalent spherical radius. The difference $\sigma_{tot} - \sigma_G$ is sensitive to the shape of the nuclear density in the extreme surface region[3]; even in the absence of narrow structures, measurements of σ_{tot} and the reaction cross section σ_R should prove very revealing.

It was hoped that despite the strong absorption resulting from annihilation even in the $\bar{N}N$ configuration, narrow bound or resonant molecular states[2] would arise at least for sufficiently large orbital angular momentum l. If, as now appears increasingly likely, sharp resonances are absent in the elementary $\bar{N}N$ system, one may still look for resonant structure in $\bar{p}$-$\mathfrak{N}$ cross sections (not arising from embedding a two-body resonance in the medium).

The real part of the average $\bar{p}$-$\mathfrak{N}$ potential is generated by meson exchanges of effectively longer range than the processes which gives rise to annihilation. Thus the $\bar{N}$-nucleus interaction will possess a pocket of attraction extending beyond the edge of the strongly absorbing region, and many-body shape resonances can arise. We now explore the likelihood of such a phenomenon, using a simple optical-model picture for the $\bar{p}$-nucleus interaction.

In principle, the $\bar{N}$ optical potential is derivable from a folding of the nuclear density $\rho(r)$ with an effective $\bar{N}N$ interaction V_{eff} inside nuclear matter. The qualitative features of this optical potential are apparent from the start. The short-ranged two-body annihilation will create an absorptive part for the optical potential which more or less follows the nucleon density distribution $\rho(r)$. Thus if we write

$$V_{\bar{N}}^{\mathrm{opt}}(\vec{r}) = V(\vec{r}) + iW(\vec{r}), \tag{1}$$

then we expect $W(\vec{r}) = -W_0\rho(r)/\rho(0)$ to be a reasonable first approximation, with the value of W_0 somewhat uncertain, but in principle consistent with known two-body annihilation and $\bar{p}$ atomic data.[4,5] On the other hand, the real part obtained from

$$V_{\bar{N}}^{\mathrm{opt}}(\vec{r}) = \int d^3r' V_{\mathrm{eff}}(\vec{r} - \vec{r})\rho(\vec{r}') \tag{2}$$

will extend appreciably beyond $\rho(\vec{r})$. For $\bar{N}$ interactions with a spin-isospin saturated core, $V(r)$ will obtain contributions from attractive ω- and ϵ-meson exchanges as well as from a considerably longer-ranged nonresonant 2π exchange (second-order tensor force). In addition, $V(r)$ will contain an annihilation-generated dispersive term.[6]

The folded potential of Eq. (2) can be converted to an equivalent Woods-Saxon potential

$$V_{\bar{N}}^{\mathrm{opt}}(r) = -V_0 f_R(r) - iW_0 f_I(r), \tag{3}$$

where $f_{R,I}(r) = \{1 + \exp[(r - R_{R,I})/a_{R,I}]\}^{-1}$. We take $R_{R,I} = r_{0R,I}A^{1/3}$. The central point is that we have $r_{0R} > r_{0I}$, i.e., a longer-ranged real potential. For a Yukawa two-body interaction, the folding model gives $R_{R,I} \cong (\frac{5}{3}\langle r^2\rangle_{\mathfrak{N}} + 6\mu_{R,I}^{-2})^{1/2}$, where $\mu_I^{-1} \approx m_N^{-1} \approx 0.2$ fm and $\mu_R^{-1} \approx 0.5$–0.7 fm are the ranges associated with annihilation (baryon exchange) and meson exchange (ϵ or uncorrelated 2π exchange). Writing $r_{0R} = r_{0I} + \delta$, we have $\delta \approx 0.9A^{-2/3}$, i.e., $\delta = 0.3$ ($A = 16$) or 0.1 fm ($A = 27$). For convenience, we take an average value $\delta \approx 0.2$ fm here, and so we use $r_{0I} = 1.1$ fm and $r_{0R} = 1.3$ fm. In principle, the diffusivities a_R and a_I can also differ (we expect $a_R > a_I$); in the absence of a detailed calculation of $V_{\bar{N}}^{\mathrm{opt}}$, we simply assume $a_R = a_I = 0.52$ fm, a value used in the analysis of antiprotonic-atom data.[5] Note that a reasonable difference $a_R - a_I > 0$ is also effective in producing the type of orbiting phenomenon discussed here, although the resulting resonances are somewhat narrower and more closely spaced in energy than those generated by having $r_R - r_I > 0$.

The depth parameters V_0 and W_0 can be estimated from meson-exchange models[6,7] or from fits to energy shifts and widths in antiprotonic atoms.[4,5] From Barnes *et al.*[4] we obtain $V_0 \approx 240$ MeV and $W_0 \approx 120$ MeV, while from Roberson *et al.*[5] we get $V_0 \approx 70$ MeV and $W_0 \approx 210$ MeV. The depths are not well determined by the somewhat crude experimental antiprotonic-atom data, and also depend strongly on the choice of radius and diffuseness parameter for $\rho(r)$. The values of V_0 and W_0 given above serve only as rough guides. Other estimates are obtained by taking $V_{\mathrm{eff}}^{\bar{N}N}$ to be either (a) the free-space two-body potential $\sum_{i=\epsilon,\omega} v_i$ plus the annihilation potential of Ref. 6, or (b) the $\bar{N}N$ t matrix.[7] From (a) we obtain extremely deep potentials with $V_0 \approx 1$ GeV and $W_0 \approx 700$ MeV, and from (b) we obtain a *repulsive* real part with $V_0 \approx -70$ MeV and $W_0 \approx 65$ MeV. Neither of these methods yields a potential remotely consistent with the antiprotonic-atom data, indicating a very strong modification of the $\bar{N}N$ interaction in the nuclear medium. The calculation of V_{eff} is in itself a very interesting problem, which we will not address here. For the sample calculations shown below, we have chosen $\{V_0, W_0\} = \{300, 100\}$ or $\{100, 200\}$ MeV as representative values, corresponding roughly to Refs. 4 and 5.

The most natural way for resonances to arise is through a radius difference ($r_R > r_I$). One can also force the appearance of backward enhancements with $r_R = r_I$, if W_0 is unrealistically small. However, such extreme values are not consistent with the antiprotonic-atom data. Such peaks generally correspond to smaller backward cross sections and larger widths than those which occur for $r_R > r_I$.

Some of our results for $\bar{p} + {}^{16}$O and $\bar{p} + {}^{40}$Ca are shown in Fig. 1. The fact that the real $\bar{p}$-nucleus potential extends beyond the imaginary part has dramatic consequences: This longer-range "pocket" of attraction can produce a strong radial localization of the $\bar{p}$ wave function, which is also focused by the centrifugal potential. If the absorptive potential is relatively weak in this "pocket," one can get *orbiting* of the nucleus by an antiproton. The signature for this phenomenon is a strong enhancement of the backward $\bar{p}$-nucleus elastic scattering cross section for characteristic energies. The excitation function for 180° scattering in Fig. 1 displays the most noticeable resonances for parameter set (a), which corresponds to $r_R > r_I$, a deep real potential, and an absorption which is sufficiently weak to permit orbiting. If we remove the radius difference (i.e., take $r_R = r_I$), but keep a deep real potential [set (b)], backward peaking still occurs, but is considerably reduced in magnitude. A "background" curve is generated by taking stronger absorption, a weaker real part, and $r_R = r_I$ [set (c)]. Although the excitation function still exhibits diffractive oscillations for set (c), the strong maxima which

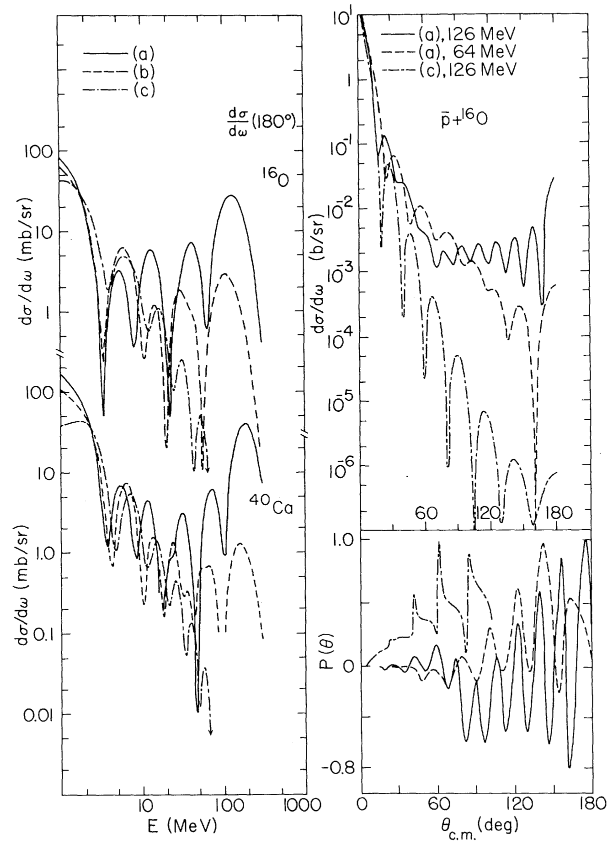

FIG. 1. On the left-hand side, we show excitation functions $d\sigma/d\omega$ (180°) for $\bar{p}$ + ^{16}O (upper half) and $\bar{p}$ + ^{40}Ca (lower half). The three curves for each target are for the parameter sets (a) V_0 = 300 MeV, W_0 = 100 MeV, r_{0R} = 1.3 fm, r_{0I} = 1.1 fm, and $a_R = a_I$ = 0.52 fm; (b) V_0 = 300 MeV, W_0 = 100 MeV, $r_{0I} = r_{0R}$ = 1.2 fm, and $a_R = a_I$ = 0.52 fm; and (c) V_0 = 100 MeV, W_0 = 200 MeV, $r_{0R} = r_{0I}$ = 1.20 fm, and $a_R = a_I$ = 0.52 fm (background). A Coulomb potential, if we assume a uniform charge distribution in the target, has been included. Higher-energy cross sections for (c) are below 0.01 mb/sr. On the right-hand side (upper half), we display angular distributions $d\sigma/d\omega$ for $\bar{p}$ + ^{16}O at the highest peak of the excitation curve (126 MeV) and its preceding minimum (64 MeV) for parameter set (a). Note the relative constancy of the average value between 60° and 170° at the "peak" energy. For comparison, the angular distribution for the "background" parameters (c) at 126 MeV is also given. In the lower half, we show polarization angular distributions $P(\theta)$ for parameter sets (a) and (c) with the addition of a spin-orbit potential of the Thomas form with $V_{s.o.} = -6$ MeV (sign opposite to that for proton). The addition of the spin-orbit term shifts the maximum to 116 MeV; for parameter set (a), P is also plotted for 64 MeV. Note the predominantly opposite signs for the 116-MeV curves; thus, even with poor resolution (~10°) the two can be distinguished. The $P(\theta)$ curve for the "background" parameters has been terminated at 90° for clarity.

occur above 10 MeV for set (a) are removed. Note that backward peaks can occur in ordinary diffractive scattering, without any resonance phenomenon. The solid curve in Fig. 1, describing orbiting, differs in two major ways from the nonorbiting curves. First, the average backward-hemisphere cross section is much larger and therefore easily measured. Second, the orbiting phenomenon is signalled by an order-of-magnitude rise in $\sigma(180°)$ above an essentially flat plateau. The predicted resonance cross sections at 180° are $\gtrsim$ 10 mb/sr and should certainly be observable.

On the right-hand side of Fig. 1, we show angular distributions for $\bar{p}$ + ^{16}O scattering at energies corresponding to a maximum or a minimum in the 180° excitation function. At the resonance (126 MeV) for parameter set (a), the cross section is enhanced over background throughout the entire backward hemisphere, compared with the "background" calculation. The resonances for $\bar{p}$ + ^{16}O shown in Fig. 1 have widths varying from about $\Gamma \approx 15$ MeV (E_{res} = 13 MeV) to $\Gamma \approx 150$ MeV ($E_{res} \approx$ 126 MeV). Their effects are thus more readily seen in elastic angular distributions than in total cross sections. The situation is reminiscent of backward-angle heavy-ion scattering[8]; for instance, the ^{16}O-^{28}Si system at c.m. energies of 20–60 MeV, where enhancements interpretable as orbiting phenomena are seen. For the $\bar{p}$ case, the 180° rise is generally due to one (or sometimes two) partial waves, so one can speak of a genuine single-particle resonance (very sharp l localization).

Large $\bar{p}$ polarizations $P(\theta)$ are also associated with the orbiting phenomenon. Some typical results for $P(\theta)$ are shown in Fig. 1. One sees that $|P(\theta)|$ attains values approaching unity for $\theta > 90°$ in the region of a peak in the 180° excitation function. On the other hand, the values of $|P(\theta)|$ are much smaller in the forward hemisphere ($\theta < 60°$) than for "background" scattering. These results may have some implications for proposals[9] to produce a polarized $\bar{p}$ beam at LEAR.

The polarization arises from a one-body antinucleon spin-orbit term $V_{\bar{N}}^{LS}(r)$, which should be added to Eq. (2). It is not clear how large this term should be. For nucleons, V_N^{LS} arises mostly from the underlying two-body spin-orbit potential, which receives coherent contributions from ϵ and ω exchange. There is in addition a significant correction to V_N^{LS} in heavy nuclei from the second-order iteration of the two-body tensor force. For the $\bar{N}$, the contribution of ω exchange is of opposite sign, and now largely cancels against the ϵ contribution. However, the $I=0$ tensor force for the $N\bar{N}$ system is *coherently attractive*,[2] unlike the NN case, and thus second-order tensor contributions to $V_{\bar{N}}^{LS}$ could be very large. Thus it is not even clear what the *sign* of $V_{\bar{N}}^{LS}$ should be at low energy, and measurements of $d\sigma/d\omega$ and $P(\theta)$ for $\bar{p}$-$\mathfrak{N}$ scattering should be encouraged. If the orbiting phenomenon is found experimentally, one may obtain some useful information on both the real and imaginary parts of the $\bar{N}N$ effective interaction. The appearance or nonappearance of $\bar{p}$ orbiting is rather delicate; it depends sensitively on the choice of well depths and the difference in the radii and diffusenesses of the real and imaginary potentials. The results presented in Fig. 1 are only representative of a number of calculations we have performed. In the presence of *extremely* strong $\bar{p}$ absorption in the nuclear interior, the existence of narrow surface localized resonances would be somewhat remarkable, and well worth more detailed study. Finally, it should be noted that the $\bar{N}$ absorptive strength should be roughly energy independent below the pion production threshold, since it arises predominantly from two-body annihilation rather than $\bar{N}$-nucleus inelastic effects. Hence if resonances appear in the 10–60-MeV region, they should also be present at higher energies.

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