New State of Binding of Antiprotons in Atoms

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It is suggested that there exists a novel "molecular" state of binding of antiprotons to atoms in which an excited highly polarized electron is located between the nucleus and the antiproton. The lifetime of this state is expected to be decided by the radiative lifetime of the excited electronic state.

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The behavior of massive negatively charged particles $(\pi^-, K^-, \mu^-, \bar{p})$ in matter, of interest since the early work of Fermi and Teller,[1] has attracted renewed attention[2] due to the availability of beams of low-energy antiprotons of increased intensity.[3] Such negatively charged particles eventually stop in matter by being captured in bound states of atoms. Capture occurs by ejection of an atomic electron and the exotic particle initially occupies a Rydberg orbit of high principal quantum number n. Decay of the high-lying state occurs by cascades of Auger processes in which the atomic electrons are "boiled off." Indeed, it is known that complete ionization of noble-gas atoms occurs.[3] Where Auger transitions are not possible the antiprotonic states (here we will concentrate principally on $\bar{p}$) decay radiatively.

The $\bar{p}$ states of most interest to date have involved radii of a few atomic units (a.u.) or much less and, particularly in the case of helium,[4] it has usually been sufficient to consider them to be essentially hydrogenic with weak perturbation by the atomic electrons. Recently, Yamazaki *et al.*[5] have observed long-lived states of negative kaons bound to He^+. The presence of these states was suggested by Condo.[6] They are essentially hydrogenic with principal quantum number $n \gtrsim 30$ and very high $l \lesssim n-1$. For these states Auger transitions are suppressed and radiative lifetimes of $\sim 10^{-8}$ sec are predicted. Here we suggest novel states of binding of massive negatively charged particles in matter in which Auger transitions are similarly suppressed. However, these states are extremely nonhydrogenic, and have large radii (typically 100 a.u.) with the $\bar{p}$ strongly localized in space. In contrast to hydrogenic states their existence depends upon an intimate interaction with one of the atomic electrons. In many ways they resemble "molecules" formed from two heavy particles, the positive nucleus and negative $\bar{p}$, and an electron bound to them.

The existence of such states can be inferred from an extremely simple argument based on the static potential energy of an array of point charged particles. Indeed, the first hint of the existence of such states arose from a study of the *classical* mechanics of two electrons in the field of a point nucleus.[7] Naively, one would expect that the optimally stable configuration would be one in which the two negatively charged particles stay out of each others' way by arranging themselves on opposite sides of the nucleus. Such states of course exist, for example, the doubly excited symmetric resonant states of helium.[8] However, states also exist in which both electrons are essentially on the same side of the nucleus but in orbits of very different principal quantum number ("planetary" atoms[9]). It is the antiprotonic analog of such electronic states that we suggest here as giving rise to a novel form of binding of $\bar{p}$ in atoms.

The states occur in three-body Coulomb systems of almost collinear form comprised of a nucleus of charge Z, an electron, and a negatively charged particle N. In the static situation the potential energy of N consists of the sum of an attractive potential due to the nucleus (taken as the origin) and a repulsive potential due to the electron (taken as fixed at unit distance). In this configuration the potential $V(R)$ as a function of the distance R between N and Z shows a minimum at $R=\sqrt{Z}/(\sqrt{Z}-1)>1$ for $Z>1$, i.e., all nuclei other than hydrogen (or an atomic core with effective charge >1) should bind an antiproton in this configuration. This potential is shown in Fig. 1. It is also easily estimated that, for $\bar{p}$, the quantum of vibration in this outer potential well is much smaller than the well depth. Consequently the well can support many bound states for $Z \geq 2$. If the electron is moved off the internuclear axis, the well depth diminishes as the $\bar{p}$ sees more and more of the nuclear charge.

Although the simple arguments above suggest the existence of bound states in the outer potential, the confirmation requires a quantum-mechanical calculation. This is simplified by the large mass of the nucleus and of

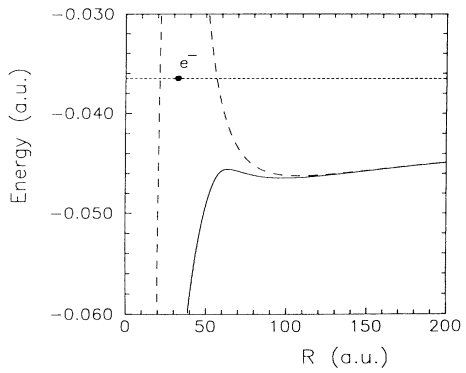

FIG. 1. The continuous curve is the BO potential energy of a singly negatively charged particle N located a distance R from a nucleus of charge $Z=2$ when the electron has MO quantum numbers $n_\lambda=6$, $n_\mu=0$, $m=0$. The dashed curve is the corresponding potential for a collinear static point-charge configuration with the electron located as shown.

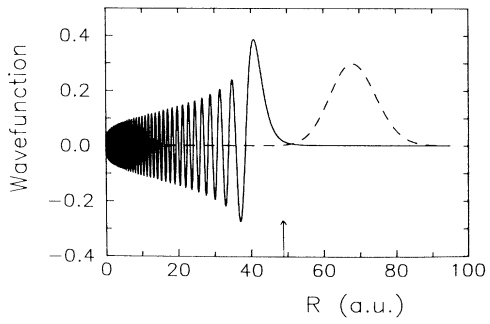

FIG. 2. The wave function of a μ^- in the BO potential for an electron with $n=6$. The continuous curve is for the μ^- state with $n=89$. The dashed curve is for the state next higher in energy. The arrow indicates the location of the maximum of the potential barrier.

the antiproton compared to the mass of the electron. This allows the Born-Oppenheimer (BO) approximation to be used. Fortunately, programs to calculate the bound states of an electron in the field of two stationary charged particles are readily available.[10] The quantum analog of the collinear configuration is an electron in a hydrogenic state of maximum polarization. For principal quantum number $n=n_1+n_2+m+1$, where n_1, n_2, and m are parabolic and magnetic quantum numbers, this implies $n_1=n-1$, $n_2=0$, and $m=0$. This is the state produced at $R=\infty$ from a molecular orbital (MO) with prolate spheroidal quantum numbers $n_\lambda=n-1$, $n_\mu=0$, $m=0$. The BO potential for this state first shows a minimum for $Z=2$ at $n=5$ (below this value the radial off-axis extent of this wave function is too great). The potential is shown in Fig. 1 for the rotationless case of total S symmetry in comparison with the point-charge static potential. It is seen that the major effect of quantum delocalization is to smear out the potential singularity into a potential hump.

The heavy-particle vibrational eigenstates in the BO potential have been calculated and the results confirm fully the simplified picture. In Fig. 2 are shown the wave functions for μ^- states (those for $\bar{p}$ show too many nodes to be drawn easily). For $n<89$ the states are confined to the inner well and are essentially hydrogenic. The $n=90$ wave function shows a dramatic change. It has almost zero amplitude in the inner well and is clearly the ground-state oscillator function in the outer well. The wave function for $n=91$ is then the first excited state in this well. The calculated oscillator spacing for an antiproton is 1.3×10^{-3} eV which agrees qualitatively with 4×10^{-3} eV from the collinear static model.

To estimate the lifetime of these states it is necessary to consider the possible decay processes. For this we consider the $\bar{p}$ states since the lifetime against decay by non-BO interactions increases with the reduced mass of the (ZN) system. Let us first consider a $\bar{p}$ particle in the ground state of the outer potential well. It can decay by three processes: (a) tunneling into the inner potential well, where it decays by normal channels, i.e., the Auger and radiative decays so far observed; (b) nonradiative transitions due to non-BO couplings in which the electron transfers to a lower state whose potential energy curve has no outer potential well (see Fig. 3) and the $\bar{p}$ then transfers to a Rydberg orbit in the inner well of this potential curve; (c) radiative processes in which the inner electron decays radiatively and the $\bar{p}$ makes a Franck-Condon transition to bound or unbound states in a lower potential well.

In case (a) the tunneling is estimated from the depth and width of the potential barrier separating inner and outer wells. For the electron in $n=5$ only one bound state is possible and this has a tunneling lifetime of only 10^{-11} sec since it is located near the top of the barrier.

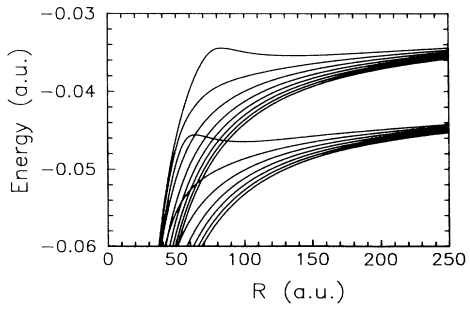

FIG. 3. The BO potential curves for all MO states asymptotic to given parabolic states of the $n=7$ and $n=8$ manifolds. Only the highest state with maximum polarization along the internuclear line has a potential well. Potential curves of lower energy are for MO in which the nucleus increasingly lies between the electron and the antiproton.

However, for an electron in $n=6$ or 8 the ground $\bar{p}$ state has a tunneling lifetime of 10^5 and 10^{40} sec, respectively. Hence if an electron can be produced in a high enough Rydberg state the tunneling decay is negligible. This confirms the results of classical mechanical studies.[7] Since the electron has maximum probability to be located near the barrier, tunneling will almost always give rise to Auger emission of the electron.

Decay of the type (b) occurs by non-BO coupling terms and corresponds to the electron transferring to a lower-energy less-polarized state corresponding to an attractive $\bar{p}$ potential as shown in Fig. 3. In other words, the electron moves off the internuclear line and the antiproton accelerates towards the nucleus. The rate for this to occur can only be ascertained by a detailed calculation of non-BO couplings. However, in the case of a two-electron planetary-atom state in which the electrons have very different principal quantum numbers the lifetime is known to be $\sim 10^{-12}$ sec. The same decay process operates for the $\bar{p}$ state except that due to the inertia of the $\bar{p}$ the lifetime is increased by a factor $(M_p/m_e)^2$. This gives an estimated lifetime of $\sim 10^{-6}$ sec. Indirect evidence for the stability of these states against decay via non-BO couplings has been obtained from studies of the classical motion of the three-body Coulomb system.[7] In the classical calculation the full dynamics is solved, i.e., the non-BO couplings are taken completely into account. An example of the stable classical motion corresponding to the quantum state described here is depicted in Fig. 4. While the electron rapidly orbits the nucleus in a highly elliptical orbit, the $\bar{p}$ oscillates back and forth close to the internuclear line. We have shown that this motion is stable in that it occupies a torus in phase space whose center is a completely stable periodic orbit corresponding to a collinear configuration of nucleus, electron, and antiproton. That the regular phase-space volume in the neighborhood of this orbit corresponds to very long-lived resonances is supported by the fact that in classical motion of the type shown in Fig. 4 the bounded oscillation of the $\bar{p}$ can be followed for hundreds of periods.

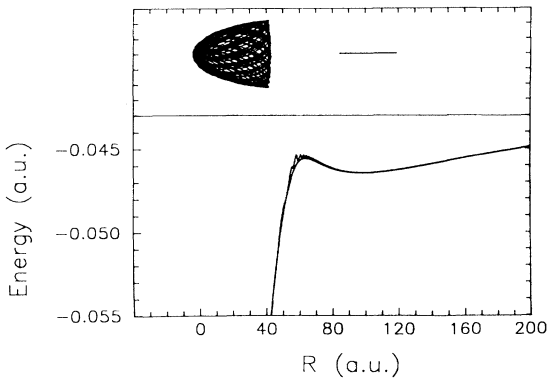

FIG. 4. The upper panel shows the classical trajectory of an electron orbiting the nucleus at the origin while a $\bar{p}$ executes a stable linear oscillation. In the lower panel the BO potential of Fig. 1 is shown as a smooth curve. The corresponding classical adiabatic potential is coincident with this curve apart from small structure in the peak region.

The validity of the quantum-mechanical BO approximation has also been tested by constructing the adiabatic classical potential shown in Fig. 4. This has been calculated by solving the full three-body problem with initial conditions such that the $\bar{p}$ moves extremely slowly in the direction of the nucleus. At each separation R the total (kinetic plus potential) energy of the electron is calculated. The result shown in Fig. 4 is in remarkable agreement with the quantum-mechanical BO potential which is also shown. The only difference is in the peak region where electron and $\bar{p}$ are close together and can interchange energy. Note that the center of the $\bar{p}$ oscillation is coincident with the minimum of the adiabatic potential. That the exact $\bar{p}$ motion shown in the upper part of Fig. 4 clearly corresponds to oscillation in the potential well of the lower part confirms the validity of the classical adiabatic approximation and hence the corresponding quantum BO approximation.

The lifetime against decay of the type (c), direct radiative transition of the electron to a lower state, is easy to estimate if it is assumed that the electronic state is well approximated by a hydrogenic parabolic state with $n_1=n-1$, $n_2=0$, and $m=0$. The lifetime of such an excited state increases as n^3 but for $n\approx 6$ is estimated to be $>10^{-8}$ sec. However, we consider this to be a pessimistic estimate of the lifetime. It should also be noted that the electron's radiative lifetime and the stability of the $\bar{p}$ state increases strongly with the principal quantum number of the parabolic state occupied by the electron.

Although such states should exist, we do not underestimate the difficulty of their production and detection. They have very large radii; nevertheless, the similar states which exist in "planetary atoms" can be produced by laser pumping[9] and are readily observed. Very low-energy $\bar{p}$ can also be obtained.[11] A marrying of these two techniques could result in a collision of a neutral doubly excited atom (a bare nucleus as core is not required, only a core with effective charge in excess of unity) with a cold $\bar{p}$ in which the $\bar{p}$ is captured by ejection of the outer electron. Alternatively, a singly charged ion with an electron excited to a suitable parabolic orbit could be viewed as a trap for cold antiprotons. Although the well is quite shallow ($\sim$20–50 meV), we have shown that the oscillator level spacing ($\sim$1–2 meV) is such that very many levels are available in the well.

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