ISSN: 2639-0108

Editorial

Advances in Theoretical & Computational Physics

A Novel Super Current Based on Dichromatic Potential Surface Deformation in Two-Electron Atoms

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Submitted: 26 Nov 2019; **Accepted**: 02 Dec 2019; **Published**: 02 Jan 2020

Abstract

Slow electron scattering from a highly excited Rydberg atom is shown to be controlled by deformed potential surfaces. More over, that deformation shows a dichroism. The potentials are different for shrinking and expanding modes of motion, respectively. We attribute this effect to a novel fictitious force in atomic structure..

Subject of this paper is the study of a two-electron atom near its threshold for double escape; i. e. we have to treat the zero-energy wave equation

$$\left[-\tfrac{1}{2}(\Delta_1+\Delta_2)-\frac{Z}{r_1}-\frac{Z}{r_2}+\frac{1}{r_{12}}\right]\Psi=0 \qquad (1)$$

where Z is the nuclear charge, r_1 and r_2 are the electron-nucleus distances, r_{12} is the electron-electron separation, and the Δ's are Laplacians.

A solution of (1) by direct integration is unknown; a look onto the potential surface of the three-body-system provides, however, a helpful hint. We observre two attractive Coulomb valleys separted by a ridge. The motion across the ridge is unstable in consequence of radial correlation. The importance of that ridge was emphasized long ago by Wannier [1]. The bending motion, however is stable, and is not relevant for the following. Hypersherical coordinates are very useful to treat the present situation. For simplicity we focuss here only to S states. Then we need only three coordinates to describe electron pair configurations. In the collinear configuration we choose the two relevant coordinates following Sommerfeld as follows [2]

We map the set of zero-eenrgy collinear configurations onto the hypersherical equator plane of a six-dimensional hypershere parametrized by the hyperradius [3].

$$R=\sqrt{r_1^2+r_2^2} \qquad (2)$$

and the azimuth angle defined by [2]

$$\varphi=tan^{-1}\frac{r_1^2-r_2^2}{2r_1r_2} \qquad (3)$$

In terms of these coordinates the wave equation (1) rewrites as

$$\left\{-\tfrac{1}{2}\frac{\partial^2}{\partial R^2}-\frac{2}{R^2}\frac{\partial^2}{\varphi^2}-\frac{C_0}{R}+\frac{15}{8R^2}-\frac{C_2}{2R}(\varphi-\pi)^2\right\}R^{-5/2}\Psi=0 \qquad (4)$$

where the potential has been approximated in quadratic order around the top of the hill.

In terms of these coordinates the top of the ridge is located at the position φ=π, and the ridge itself runs along R. The constant

$$C_0=\frac{4Z-1}{\sqrt{2}} \qquad (5)$$

is the net charge on the ridge top where is

$$C_2=\frac{12Z-1}{4\sqrt{2}} \qquad (6)$$

is the curvature across the ridge. The Ansatz

$$\Psi=R^{5/2}e^{i\int K(R')dR'}e^{i\kappa(\varphi-\pi)^2} \qquad (7)$$

for the wave function leads then directly to he Riccati equation [3]

$$\frac{8\kappa^2}{R^2}=\frac{C_2}{2R}-K\frac{d\kappa}{dR}. \qquad (8)$$

with the squared radial wave number

$$K^2=2\left(E+\frac{C_0}{R}-\frac{15/8+4i\kappa}{R^2}\right) \qquad (9)$$

where the hypercentrifugal barrier $^{15}/_{8R^2}$ prevents the collaps of the whole atom.The second exponential in (7) constitutes an eigen function of the antioscillator across the ridge; i.e. its wave number κ is determined by (8) [4]. In the derivation of the Ricati equation only the first derivative of κ with respect to R has been taken into account. The present article shows below that the second derivative would deliver a short-range contribution not relevant here.

The rhs of (8) contains the electrostatic curvature plus a velocity-dependent term to be analyzed now. To this end we remark that at threshold E=0 the Coulomb zone extends to infinity. Therefore the centrifugal term in (9) does not contribute to the leading term of the radial wave number K now given by

$$K = \pm\sqrt{\frac{2C_0}{R}}. \qquad (10)$$

We put (10) into (8) and conclude that the Ansatz

$$\kappa = \bar{\kappa}\sqrt{R} \qquad (11)$$

solves the Ricati equation in the Coulomb zone. The coupling term K (dκ)/dR delivers therefore a contribution $\propto {}^{1}/_{\mathrm{R}}$ and may be combined with the electrostatic curvature. Thus we arrive at a dynamical curvature given by

$$C_2^{dyn} = C_2 - \bar{\kappa}\sqrt{2C_0} \qquad (12)$$

We stress that both wave numbers K and κ carry a sign depending on the mode of breathing motion. That may be either a shrinking motion or an expansion of the whole atom what constituts a three-body generalization of incoming or outgoing waves in the two-body case. We remark that in the shrinking mode (K< 0, κ < 0) occurs an electron-electron attraction.

We find for the shrinking mode the dynamic ridge curvature

$$C_{2shrinking}^{dyn} = C_2 - \frac{1}{8}\left\{\sqrt{C_0^2 + 32C_0C_2} + C_0\right\} \qquad (13)$$

In contrastto (13) the expanding system (K>0 ,κ>0) manifests itself by a different dynamic curvature givenb y

$$C_{2expanding}^{dyn} = C_2 - \frac{1}{8}\left\{\sqrt{C_0^2 + 32C_0C_2} - C_0\right\} \qquad (14)$$

At a first glance the dichromatic potential surface deformation presented above appears surprising. On the other hand our effect is not entirely unexpected because we desribe the electron-electron interaction in a moving frame. The 7 motion proceed here along the hyperradius R. It is well known that in moving frames fictitious forces occur.

The interpretation in terms of a fictitious force is supported by the structure of our wavee quation (8) because the rhs term K ${}^{\mathrm{d}\kappa}/_{\mathrm{dR}}$ constitues a velocity-dependent contribution to the electrostatic curvature.

In summary, the stationary state in two-electron atoms near threshold of double escape may be described as follows. Using He^{**} as prototype target atom we have scattered a slow electron from He^{+*}. The electron pair moves then on the potential surface whose curvature is given by (13). That curvature is largely reduced in comparison to its static value C_2 and allows the electron pair to jump towards the top of the ridge. Physically this implies a reduction of the electrron-electron separation corresponding to a mutual electron attraction. The pair on the ridge moves now towards the nucleus; i. e. the whole complex is reduced in size. That process was above denoted as shrinking. This shrinking is controlled by the motiona long the evolution coordinate R, the hyperradius. The classically allowed interval for that motionis R∈[R(0,),∞) where R_0 is the turning point determined by the repulsive hypercentrifugal barrier given by $15/8R^2$. Reflection of the two-electron wave prevents a collaps of the whole three-body complex. We stress that an incoming electron scattered from the ion He^{+*} causes a shrinking motion of the whole complex since that motion occurs on the potential surface whose curvature is given by (13). This dxnamical curvature is strongly reduced compared to its static value. This reduction makes it easier for he electrons to jump onto the ridge top.

After the reflection from the turning point R_0 the two-electron wave moves on the potential surface whose curvaruture remains close to its static value (14). Therefore this step of the motion is highly unstable, and the pair decays. At the total energy E=0 one electron is trapped into a high Rydberg state of the ion He^{+*}. with a small negative binding energy -ε < 0, and the other electron escapes slowly with a small kinetic energy of +ε > 0 into the continuum.

There it may be scattered from another He^{+*} ion, and the whole process repeats. Along these lines we obtain a series of pair creations and annihilations. That process compairs favorably with Cooper pairs. There is, however, a hugh difference between Coopeer pairs and our pairs. The binding mechanism of a Cooper pair stems from an electron-phonon interaction which is strongly dependent on the temperature. In contrast to that our pairs stem from a dominant three-body correlation independent of the temperature, and may be the basis for a high temperature super current [5, 6].

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