Dynamic polarization of light halo nuclei in strong fields: ^{6}He + ^{209}Bi elastic scattering below and close to the Coulomb barrier

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The elastic scattering of light halo nuclei in the field of heavy targets has been studied for collision energies below the Coulomb barrier. Based on the assumption that the neutron halo follows the projectile adiabatically along its classical trajectory, a dynamic polarization potential is derived which describes both the (electrical) polarization as well as the breakup of the projectile in the field of the target. Detailed computations have been carried out for the elastic scattering of ^{6}He + ^{209}Bi at energies between 14.7 MeV and 19.1 MeV near to the Coulomb barrier. It is demonstrated that the polarization of the halo nucleus leads to a clear decrease of the (elastic) scattering cross section in excellent agreement with a recent measurement by Aguilera *et al.* [Phys. Rev. Lett. **84**, 5058 (2000)].

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I. INTRODUCTION

The existence of (neutron) halos in light and midrange nuclei has been one of the most fascinating discoveries in nuclear and astrophysics during the last two decades [1–8]. Today, such neutron halos are well confirmed for a number of light nuclei, including ^{6}He, ^{11}Li, ^{14}Be as well as for several others. Usually, these halos are formed by two (or more) loosely attached neutrons which surround a deeply bound core with binding energies of less than 1 MeV, in contrast to the 6–8 MeV of stable nuclei. Therefore, by studying the properties of halo nuclei, one may obtain information not only about the nuclear forces but also on neutron-neutron correlations [2,3] or (so-called) astrophysical S-factors [9] for describing various nuclear reactions. Aside from experimental investigations on the β-delayed deuteron emission [10] and the scattering and breakup of halo nuclei [11,12], several theoretical models have been suggested in order to understand the structure of halo nuclei including, for example, the adiabatic treatment of the internal coordinates of the neutron halo [4], a microscopic model [13], the two-frequency shell model [14], continuum-discretized coupled channels (CDCC) [15–21], or the tree-body continuum theory [22–26].

In order to explore the (shape and) structure of neutron halos, perhaps the simplest process is the elastic scattering of halo nuclei in the field of heavy targets at energies below and near the Coulomb barrier. For a pure Coulomb field, for example, the angle-differential scattering cross section of a point-like projectile along a classical (Rutherford) trajectory is well known to follow a $1/\sin^4(\theta/2)$ dependence with θ being the scattering angle of the projectile. For heavy (and nearly magic) targets, moreover, one can often neglect the effects of the Coulomb excitation as the first excited level has a large excitation energy from the ground state. During the last decade, therefore, the elastic scattering and breakup of halo nuclei has attracted a lot of recent interest [27–29] and may provide new insights into the structure of halo nuclei. The Coulomb breakup of $2n$ halo nuclei (^{6}He, ^{11}Li) exhibits a large dipole transition strength just close to the breakup threshold [30,31]. Owing to the elastic scattering experiments for ^{6}He and ^{11}Li, in addition, an extended neutron density distribution was established recently [11,12].

In this article, we present an adiabatic model for describing the elastic scattering of light halo nuclei in the Coulomb field of heavy targets. In this model, the neutron halo is assumed to be polarized by the Coulomb field of the target while moving along a classical (Rutherford) trajectory as summarized in Sec. II A. In Sec. II B, we then present the Schrödinger equation for this motion of the halo nucleus in the Coulomb field of the heavy target. Since the internal motion of the (halo) projectile is supposed to follow the field adiabatically, we are able to divide this Schrödinger equation into two equations: One for the center-of-mass motion along the Rutherford trajectory and a second one to describe the internal motion of the halo nucleus. Based on the adiabatic and the zero-range approximations, we derive in particular an expression for the relative wave function in terms of the Greens propagator. In Sec. II E, moreover, an expression is obtained for the dynamic polarization potential which applies for any arbitrary system with a dineutron configuration. This polarization potential is later utilized in Sec. III in order to calculate the elastic angle-differential cross sections for the collision of low-energy ^{6}He ions by ^{209}Bi. A remarkable reduction of the elastic scattering cross sections is found due to the asymmetry of the halo nucleus, and in very good agreement with the experiment [27,32]. Finally, a brief summary is given in Sec. IV.

II. THEORY

To analyze the cross sections for the elastic scattering of light, neutron-rich nuclei by heavy targets at sub-barrier energies, i.e., below and near the Coulomb barrier, we may assume a *deuteron-like* structure for the projectile. This means

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that the halo nucleus ('deuteron') with mass $m_{\rm d}$ consists of a charged core ('proton') with mass $m_{\rm p}$ and charge $Z_{\rm p}$, and the neutral halo ('neutron') with mass $m_{\rm n}$: $m_{\rm d} = m_{\rm n} + m_{\rm p}$. For such a structure of the projectile, the effective two-body potential between the neutral halo and charged core, $V_{\rm np}(r)$, can be approximated by the well-known Hulthen potential [33] and expressed in terms of the binding energy $\varepsilon_0 = -\hbar^2\alpha^2/2\mu$ and the reduced mass $\mu = m_n m_p/m_d$. The (binding) energy ε_0 is those needed in order to break the halo nucleus into its 'proton' and 'neutron' clusters. For the target, moreover, we assume a heavy and (nearly) magic nucleus with mass $M_T \gg m_d$ and charge $Z_T \gg 1$, for which Coulomb excitations are considered to be negligible.

A. Geometry and coordinates of the neutron halo nucleus in the Coulomb field of the heavy target

To describe the scattering of the halo projectile, let us introduce the coordinates

$$\mathbf{R} = \frac{\mu}{m_p}\mathbf{r}_n + \frac{\mu}{m_n}\mathbf{r}_p, \quad \mathbf{r} = \mathbf{r}_n - \mathbf{r}_p, \tag{1}$$

where the target nucleus is taken as the origin, and $\mathbf{r}_n$ and $\mathbf{r}_p$ are the position vectors of the neutral halo and the charged core, respectively, while $\mathbf{R}$ denotes the center-of-mass coordinate and $\mathbf{r}$ the relative position (vector) between the core and the halo. These coordinates are displayed in Fig. 1; they allow us to describe both the polarization as well as a breakup of the projectile.

B. Schrödinger equation for the motion of the halo nucleus in the Coulomb field of the heavy target

In the Coulomb field of the target, the motion of the halo nucleus is described by the time-independent Hamiltonian

$$\begin{aligned} H &= \hat{T}_{\mathbf{R}} + \hat{T}_{\mathbf{r}} + V_{\rm np}(r) + \frac{Z_p Z_T e^2}{R} - \Delta V(\mathbf{r}, \mathbf{R}) \\ &= H_0 - \Delta V(\mathbf{r}, \mathbf{R}). \end{aligned}$$

This total Hamiltonian accounts for the kinetic energy (operators) $\hat{T}_{\mathbf{R}} = -\hbar^2\Delta_{\mathbf{R}}/2m_d$ for the center-of-mass and

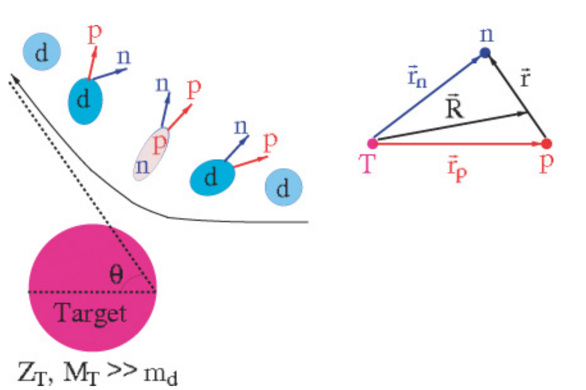

FIG. 1. (Color online) Coordinates to describe the elastic scattering of a halo nucleus in the Coulomb field of heavy targets. The deuteron-like projectile with mass m_d is formed by a charged core p ('proton') and a neutral halo n ('neutron'). See text for further details.

$\hat{T}_{\mathbf{r}} = -\hbar^2\Delta_{\mathbf{r}}/2\mu$ for the relative motion of the charged core and the neutral halo as well as the potential energy. Apart from the Hulthen potential $V_{\rm np}(r)$ of the deuteron-like halo and the Coulomb repulsion $Z_p Z_T e^2/R$, here we have to include, owing to the internal structure of the halo nucleus, the additional potential

$$\Delta V(\mathbf{r}, \mathbf{R}) = Z_p Z_T e^2[1/R - 1/|\mathbf{R} - \mu/m_p\mathbf{r}|] \tag{2}$$

in order to describe the deviations from a (pure) Coulomb scattering of projectile in the field of heavy target.

Of course, the (total) wave function of the projectile 'core + halo' must obey the Schrödinger equation [34]

$$(H_0 - E)\,\Psi(\mathbf{r}, \mathbf{R}) = \Delta V(\mathbf{r}, \mathbf{R})\Psi(\mathbf{r}, \mathbf{R}), \tag{3}$$

where the total energy $E = E_d + \varepsilon_0$ is the sum of the (asymptotic) kinetic energy E_d and the binding energy ε_0 of the halo projectile.

C. Approximation to the Schrödinger equation for describing the motion of the halo nucleus in the Coulomb field of the heavy target

1. Adiabatic approximation

An (approximate) solution to this equation can be found in the adiabatic approach, in which we assume that the internal motion of the halo nucleus is fast compared to the motion of the projectile through the Coulomb field of the target [35–38]. Mathematically, this adiabatic approximation can be expressed

$$\Psi(\mathbf{r}, \mathbf{R}) \approx \chi(\mathbf{R})\,\varphi^+(\mathbf{r}, \mathbf{R}), \tag{4}$$

where $\chi(\mathbf{R})$ refers to the wave function of the center of mass and $\varphi^+(\mathbf{r}, \mathbf{R})$ to those of the relative motion of the projectile, and in which $\mathbf{R}$ now occurs as a parameter.

Substituting Eq. (4) into the Schrödinger equation (3), we obtain the two (ordinary) differential equations

$$\left(\hat{T}_{\mathbf{R}} + \frac{Z_p Z_T e^2}{R}\right)\chi(\mathbf{R}) = (E - \varepsilon_0 - \delta V(R))\,\chi(\mathbf{R}), \tag{5}$$

and

$$(\hat{T}_{\mathbf{r}} - \Delta V(\mathbf{r}, \mathbf{R}) + V_{\rm np}(r))\varphi^+(\mathbf{r}, \mathbf{R}) = (\varepsilon_0 + \delta V(R))\varphi^+(\mathbf{r}, \mathbf{R}), \tag{6}$$

where the first equation describes the motion of the halo nucleus along the Rutherford trajectory and the second, Eq. (6), the relative motion of the neutron halo and the charged core of the projectile. In these equations, in addition, we modified the binding energy ε_0 of the projectile at the radius R by the (complex) potential $\delta V(R)$ ($|\delta V(R)/\varepsilon_0| \ll 1$) to account for the polarization of the projectile in the field of the target. Below, we shall refer to this potential as the *dynamic (Coulomb) polarization* potential. In fact, the Coulomb field of the target affects the halo nucleus in two different ways in moving along the trajectory: Aside from a polarization of the projectile with regard to its mass and charge distribution, the Coulomb field may even lead to a *breakup* of the halo nucleus into the charged core and the neutral halo. Both effects are incorporated into the theoretical description by using the dynamic polarization potential, leading again however to a parametric dependence on the center-of-mass coordinate R.

2. Zero-range approximation

We need to solve Eq. (6) in order to understand the internal motion of the (halo) projectile while moving along its classical trajectory. A solution to this equation can be found in the *zero-range* approximation [39] in which a point-like two-body interaction [36,40]

$$V_{\rm np}(r)\varphi^+(\mathbf{r},\mathbf{R}) \approx -\frac{2\pi\hbar^2}{\mu}\left(\frac{\alpha}{2\pi}\right)^{\frac{1}{2}}\delta(\mathbf{r}) \tag{7}$$

is assumed between the charged core and the neutral halo, and with $1/\alpha$ being the mean distance between the core and the neutrons for a *free* halo nucleus [41]. In the zero-range approximation, then Eq. (6) simplifies to

$$(\hat{T}_{\mathbf{r}} - \Delta V(\mathbf{r},\mathbf{R}) - \varepsilon_0 - \delta V(R))\varphi^+(\mathbf{r},\mathbf{R}) = -\frac{2\pi\hbar^2}{\mu}\left(\frac{\alpha}{2\pi}\right)^{\frac{1}{2}}\delta(\mathbf{r}). \tag{8}$$

As shown by Demkov and Drukarev [42], who studied the breakup and polarizability of charged particles in an external electric field, any solution of Eq. (8) must behave for small values of r as $\varphi^+(\mathbf{r},\mathbf{R})\xrightarrow[r\to 0]{}\frac{1}{r}-\alpha$. We may utilize this result to obtain an 'initial' value for Eq. (8) by evaluating the logarithmic derivative of $r\varphi^+(\mathbf{r},\mathbf{R})$,

$$\lim_{r\to 0}\{(\partial/\partial r)\ln[r\varphi^+(\mathbf{r},\mathbf{R})]\} = -\alpha. \tag{9}$$

Having the condition (9), we can solve Eq. (8) uniquely for any value R along the trajectory of the projectile.

D. Coulomb Green function and wave function of the internal motion

Instead of the relative coordinate $\mathbf{r}$, it is more convenient below to use the coordinate $\mathbf{r}_p$ of the charged core together with the center-of-mass coordinate R. For this choice, we then have $\mathbf{r} = (m_p/\mu)(\mathbf{R}-\mathbf{r}_p)$ and $\delta(\mathbf{r}) = (\mu/m_p)^3\delta(\mathbf{R}-\mathbf{r}_p)$, while the operator of the kinetic energy can be expressed as $\hat{T}_{\mathbf{r}} = \frac{\mu}{m_p}\hat{T}_{\mathbf{r}_p} = -\frac{\hbar^2\mu}{2m_p^2}\Delta_{\mathbf{r}_p}$ in terms of the kinetic energy of the charged core in the Coulomb field of the target. With these substitutions, the relative wave Eq. (8) becomes

$$\left(\hat{T}_{\mathbf{r}_p} + \frac{m_p}{\mu}\frac{Z_pZ_Te^2}{r_p} - \tilde{E}_p\right)\frac{m_p}{\mu}\left(-\frac{m_p}{2\pi\hbar^2}\right)\left(\frac{2\pi}{\alpha}\right)^{\frac{1}{2}}\varphi^+(\mathbf{r}_p,\mathbf{R}) = \delta(\mathbf{r}_p-\mathbf{R}), \tag{10}$$

and where we use (again) the notation φ^+ to represent the wave function of the internal motion. From Eq. (10), we see that the function

$$\varphi^+(\mathbf{r}_p,\mathbf{R}) \sim G_C(\tilde{E}_p,\mathbf{R},\mathbf{r}_p) \tag{11}$$

is proportional to the Coulomb Green function $G_C(\tilde{E}_p,\mathbf{R},\mathbf{r}_p)$ as known, for example, from the scattering theory [40] for a projectile with energy

$$\tilde{E}_p = \frac{m_p}{\mu}\left(\frac{Z_pZ_Te^2}{R} + \varepsilon_0 + \delta V(R)\right) = \frac{\hbar^2}{2m_p}k^2(R), \tag{12}$$

moving in the Coulomb potential $Z_pZ_Te^2/r_p$, and where $k(R)$ on the right hand side of Eq. (12) is used in order to denote the wave number of the charged core in the field of the target.

An explicit expression for the Coulomb Green function of a charged particle has been worked out originally by Hostler [43], but can be written also in the form [44]

$$G_C \sim \frac{1}{r}\left(\frac{\partial}{\partial\rho_+} - \frac{\partial}{\partial\rho_-}\right)\{H_0^+(\rho_+)F_0(\rho_-)\}, \tag{13}$$

where $\rho_\pm$ are two radial coordinates given by

$$\rho_\pm = \frac{k(R)R}{2}\left(1 \pm \frac{r_p}{R} + \frac{1}{R}|\mathbf{R}-\mathbf{r}_p|\right) = \frac{k(R)R}{2}\left(1 \pm \frac{\mu}{m_p}\frac{r}{R} + \frac{1}{R}|\mathbf{R}-\frac{\mu}{m_p}\mathbf{r}|\right). \tag{14}$$

In this representation of the Coulomb Green function, we have $H_0^+ = G_0 + iF_0$ and with F_0 and G_0 being the (well-known) regular and irregular Coulomb functions at the origin. The form (11) for the wave function of the internal motion of the halo nucleus in the Coulomb field of the target is one of the main results of this work, together with the expressions (13) and (14). From the viewpoint of scattering theory, this function $\varphi^+(\mathbf{r},\mathbf{R})$ describes the (quasi)stationary state of the halo nucleus while moving along the Rutherford trajectory. It is *quasistationary* because the projectile may *break* into its charged core and neutral halo, with a width determined by the imaginary part of the dynamic polarization potential $\delta V(R)$.

E. Dynamic Coulomb polarization potential

In order to analyze the probability (cross section) for a breakup of the projectile, we can apply expression (13) for evaluating the polarization potential $\delta V(R)$. Substituting this expression into the logarithmic derivative (9), and taking the limit $r\to 0$, we find that the potential $\delta V(R)$ must obey the implicit equation

$$\frac{\mu}{m_p}k(R)\left\{H_0^{+\prime}(\rho)F_0'(\rho) - \left(\frac{2\eta}{\rho}-1\right)H_0^+(\rho)F_0(\rho)\right\} = -\alpha, \tag{15}$$

where $\eta = (m_p^2/\mu)Z_TZ_pe^2/k(R)$ denotes the Sommerfeld parameter and $\rho = \lim_{r\to 0}\rho_\pm = k(R)R$. This equation is another important result of this work; indeed, its solution enables us to describe both the polarization of the projectile as well as the breakup into different parts. Note that Eq. (15) is free of any parameter other than needed to describe the scattering process and, hence, supports a direct solution of the elastic scattering of halo nuclei in the Coulomb field of heavy targets. Unfortunately, Eq. (15) cannot be solved analytically but has to be considered numerically of R. However, we can obtain the analytical expression for the dynamic polarization potential $\delta V(R)$ by applying, in addition, a quasiclassical approach to the internal motion of the projectile.

1. *Quasiclassical analytical expression for the dynamic polarization potential $\delta V(R)$*

In order to obtain the analytical expression for the dynamic polarization potential $\delta V(R)$, we must use the second order Wentzel-Kramers-Brillouin (WKB) approximation for Coulomb functions in Eq. (15). The WKB approximation for the regular Coulomb function can be written as [45]

$$F_0(\rho) \simeq \tfrac{1}{2} B e^{A}. \tag{16}$$

Its derivative is defined as

$$F_0'(\rho) \simeq \left(B^{-2} + \frac{1}{8\eta} t^{-2} B^4\right) F_0(\rho). \tag{17}$$

Here,

$$B = \left(\frac{t}{1-t}\right)^{1/4}, \tag{18}$$

$$A = 2\eta\left([t(1-t)]^{1/2} + \arcsin t^{1/2} - \frac{\pi}{2}\right), \tag{19}$$

$$t = \frac{\rho}{2\eta}. \tag{20}$$

The irregular Coulomb function is given by [45]

$$G_0(\rho) \simeq \tfrac{1}{2} B e^{-A} \tag{21}$$

within the WKB approximation. The derivative of this irregular Coulomb function can be expressed as

$$G_0'(\rho) \simeq \left(-B^{-2} + \frac{1}{8\eta} t^{-2} B^4\right) G_0(\rho). \tag{22}$$

Substituting expressions (16)–(20) into the implicit equation (15), reflecting $k(R), \rho, \eta$ in the variables of the $\delta V(R), \varepsilon_0, R$ and expanding in the small parameter $\delta V(R)/\varepsilon_0$, we obtain the quasiclassical analytical expression for the dynamic polarization potential $\delta V(R)$

$$\mathtt{Re}\,\delta V(R) = -\frac{(Z_p Z_T e^2)^2}{16 R^4 \alpha^2}\left(\frac{\mu}{m_p}\right)^2 \frac{1}{\varepsilon_0}, \tag{23}$$

$$\mathtt{Im}\,\delta V(R) = -\frac{Z_p Z_T e^2}{4R^2\alpha}\frac{\mu}{m_p}\left(1 + \frac{Z_p Z_T e^2}{8R^2\alpha}\frac{\mu}{m_p}\frac{1}{\varepsilon_0}\right) e^{2A}. \tag{24}$$

Here, the coefficient A which represents Eq. (19) is the tunneling coefficient of charged particle through a Coulomb potential barrier within a quasiclassical approach. The quasiclassical analytical expressions for the real (23) and imaginary (24) parts of the dynamic polarization potential $\delta V(R)$ are obtained from the first principles for the first time. As we can see from Eq. (23), the real part of the potential $\delta V(R)$ is inversely proportional to the center-of-mass coordinate R. This result agrees with other similar articles [15,36,37], where the adiabatic approximation has been used, as well as with the continuum-discretized coupled channels method of Ref. [17]. The imaginary part of the $\delta V(R)$ potential is proportional to the probability of the tunneling of charged particle through a Coulomb potential barrier. The imaginary part of the $\delta V(R)$ [Eq. (24)] is larger than in Refs. [15,17]. The main reason for this is that in our case the halo nucleus decays from the ground state which is treated as a quasistationary state and which is often called the 'direct' decay. In Refs. [15,17], in contrast, first a dipole resonance must be excited before the nucleus may decay. Although there is a finite probably for this 'excitation-and-decay' mechanism, the probably was found smaller than for the direct decay [33].

2. *Illustration of the dynamic polarization potential for ^{6}He scattering on ^{209}Bi*

The polarization potential of Eqs. (23) and (24) as well as Eq. (15) has been calculated for the case of ^{6}He scattering on ^{209}Bi. In Figs. 2 and 3 we present the real and imaginary parts of the dynamic polarization potential $\delta V(R)/\varepsilon_0$ for the elastic scattering of ^{6}He on ^{209}Bi nuclei as function of the center-of-mass coordinate, taken relative to the biding energy for ^{6}He. The binding energy between an α- and dineutron particles was observed experimentally and was equal to -0.975 MeV [30]. The calculations were obtained for the projectile energy $E_d = 17.8$ MeV, i.e., below of the Coulomb barrier.

The results are presented for the exact value $\delta V(R)$ (solid line) from Eq. (15) as well as for the quasiclassical analytical expressions for the dynamic polarization potential $\delta V(R)$ (dashed lines) from Eqs. (23) and (24). We note that the exact and analytical results basically coincide for $R > R_t$ where R_t is the classical Coulomb turning point. This classical point is $R_t = 15$ fm for the ^{6}He ions in the Coulomb field of the ^{209}Bi. Therefore, our computations show that small by the absolute value real and imaginary parts of the polarization potential have a long range.

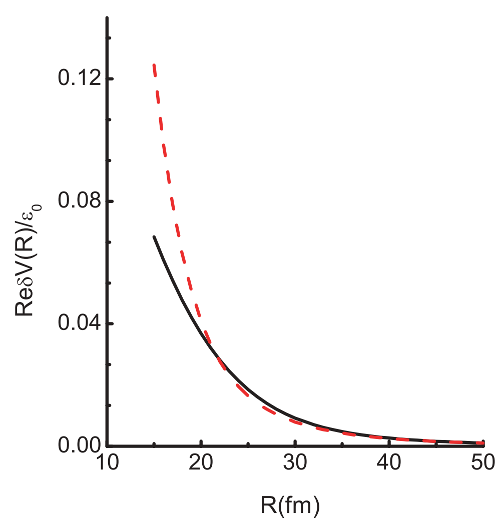

FIG. 2. (Color online) Real parts of the dynamic polarization potential $\delta V(R)$ as functions of the projectile-target separation R. Calculations have been performed for the scattering of ^{6}He ions in the Coulomb field of a ^{209}Bi for the collision energy $E_{^6\mathrm{He}} = 17.8$ MeV. Results are presented for the exact value $\delta V(R)$ (—) from Eq. (15) as well as by applying the quasiclassical analytical expression for the dynamic polarization potential $\delta V(R)$ (- -) from Eq. (23). The potentials are taken relative to the binding energy of ^{6}He.

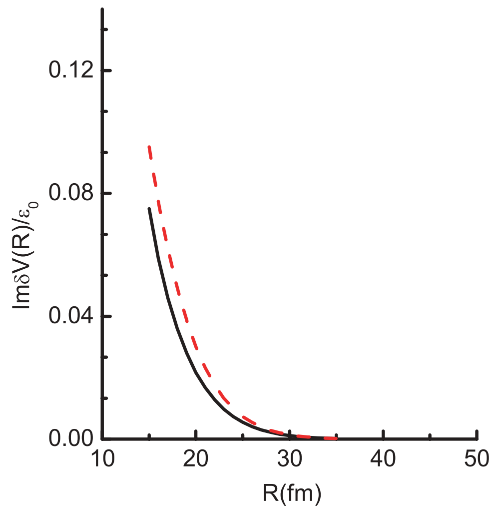

FIG. 3. (Color online) Imaginary parts of the dynamic polarization potential $\delta V(R)$ as functions of the projectile-target separation R. Calculations have been performed for the scattering of ^{6}He ions in the Coulomb field of a ^{209}Bi for the collision energy $E_{^6\mathrm{He}} =$ 17.8 MeV. Results are presented for the exact value $\delta V(R)$ (—) from Eq. (15) as well as by applying the quasiclassical analytical expression for the dynamic polarization potential $\delta V(R)$ (- -) from Eq. (24). The potentials are taken relative to the binding energy of ^{6}He.

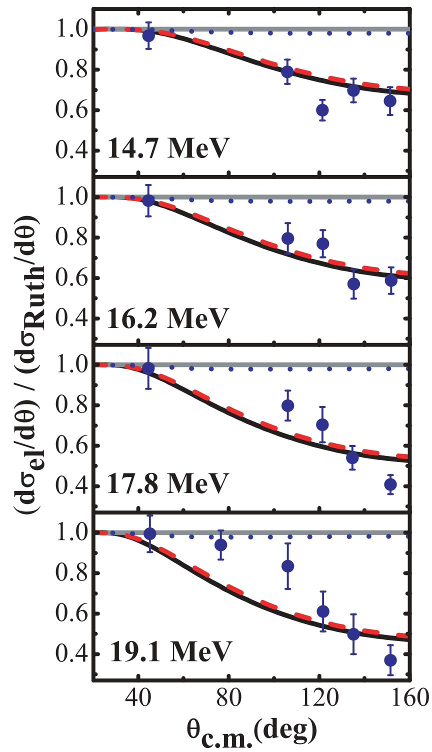

FIG. 4. (Color online) Elastic angle-differential cross sections for the collision of low-energy ^{6}He ions by ^{209}Bi nuclei as function of the scattering angle θ in the center-of-mass frame. The cross sections are taken relative to the cross sections of a pure Rutherford scattering of a projectile with mass m_d and charge Z_p and are shown for the four collision energies as observed experimentally [27,32].

III. CALCULATION OF THE ELASTIC CROSS SECTION FOR ^{6}He

By calculating the dynamic polarization potential $\delta V(R)$ from Eq. (15), we can apply the theory developed so far to describe the elastic scattering of $^6\mathrm{He} + {}^{209}\mathrm{Bi}$ at energies below the Coulomb barrier (∼20 MeV). Using $\delta V(R)$ as an optical potential without any parameter variation, the differential cross section for the elastic scattering of a ^{6}He has been calculated, assuming the ^{6}He halo nucleus to be formed by an α- and dineutron particles.

Figure 4 displays the cross section $(d\sigma_{\mathrm{el}}/d\theta)/(d\sigma_{\mathrm{Ruth}}/d\theta)$ for the elastic scattering of ^{6}He on ^{209}Bi target nuclei as a function of the scattering angle, taken relative to the cross section for a pure Rutherford scattering of a projectile with mass m_d and charge Z_p. Theoretical predictions, derived in the center-of-mass frame, are compared with a recent measurement by Aguilera *et al.* [27,32] for the four projectile energies $E_d = 14.7$ MeV, 16.2 MeV, 17.8 MeV, and 19.1 MeV, i.e., near and just below of the Coulomb barrier. In the experiments by Aguilera and coworkers, the ^{6}He beam was produced by the TwinSol radioactive nuclear beam facility at the University of Notre Dame [46]. Apart from the Rutherford cross section (straight lines), Fig. 4 also shows the relative cross sections, if only the polarization (dotted lines) or the breakup of the projectile (dashed lines) has been taken into account in the computations. These approximations correspond to the case where only the real or the imaginary part of the dynamic polarization potential $\delta V(R)$ is considered in Eq. (15). As seen from Fig. 4, the full (elastic) scattering cross sections are in very good agreement with experiment, except for the energy $E_d = 19.1$ MeV for which the cross sections are slightly underestimated, especially for small scattering angles.

The elastic cross sections decrease systematically by about 50% at the scattering angle $\theta_{\mathrm{c.m.}} = 160^\circ$ if the collision energy is enlarged from 14.7 MeV to 19.1 MeV. This decrease is due to the Coulomb field of the target since, at 19.1 MeV, the ^{6}He projectile comes closer to the ^{209}Bi target nuclei. While the deviations from the Rutherford scattering for small angles can be explained by the (so-called) 'electric' polarization of the projectile [36], the decrease at larger angles (back scattering) arise from the Coulomb breakup of the ^{6}He halo nuclei which increases because of the strong Coulomb forces.

IV. CONCLUSION

In summary, the elastic scattering of light halo nuclei in the (Coulomb) field of heavy targets has been investigated for energies below and near the Coulomb barrier. It is shown that the asymmetry of the halo nuclei, which is assumed to be formed by a charged core and a dineutron, leads to a considerable decrease of the (elastic) scattering cross section. To understand this effect quantitatively, emphasis was placed on the polarization (of the mass and charge distribution) of the projectiles which, for strong enough fields, are associated with a breakup of the halo nuclei. Assuming an adiabatic motion of the projectile along a classical Rutherford trajectory, we were able to derive an equation for the (dynamic) polarization potential which describes the internal dynamics of the halo nucleus, including its (electrical) polarization and breakup, and which leads to a clear reduction of the cross section at large scattering angles. Detailed computations have been carried out for the elastic scattering of ^{6}He + ^{209}Bi at energies between 14.7 MeV and 19.1 MeV, i.e., near the Coulomb barrier ∼20 MeV. The (theoretical) cross sections obtained from our adiabatic model agree very well with the measurements by Aguilera and coworkers (cf. Fig. 4) and, hence, make a dineutron configuration very likely for the ground state of the ^{6}He halo nucleus. No additional parameters were needed in order to explain the deviations from Rutherford's scattering cross sections.

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