Electron emission from highly charged ions: signatures of magnetic interactions and retardation in strong fields

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Abstract. The electron emission of highly charged ions has been reanalyzed with the goal of separating the magnetic and retardation contributions to the electron–electron (e–e) interaction from the static Coulomb repulsion in strong fields. A remarkable change in the electron angular distribution due to the relativistic terms in the e–e interaction is found, especially for the autoionization of beryllium-like projectiles, following a \(^1s^2 \rightarrow 2p_{3/2}\) Coulomb excitation in collision with some target nuclei. For low-energetic, high-\(Z\) projectiles with \(T_p \lesssim 10\) MeV u\(^{-1}\), a diminished (electron) emission in the forward direction as well as oscillations in the electron angular distribution due to the magnetic and retarded interactions are predicted for the autoionization of the \(^1s^2 2s^2 2p_{3/2} \ ^3P_2\) resonance into the \(^1s^2 2s^2 \ ^2S_{1/2}\) ground and the \(^1s^2 2p \ ^2P_{1/2}\) excited levels of the finally lithium-like ions, and in contrast to a pure Coulomb repulsion between the bound and emitted electrons. The proposed excitation–autoionization process can be observed at existing storage rings and will provide a novel insight into the dynamics of electrons in strong fields.

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1. Introduction

The electron–electron (e–e) interaction in (strong) Coulomb fields has attracted much interest since the early days of quantum mechanics [1, 2]. Relativistic calculations for neutral atoms and highly charged ions (HCI) soon showed that accurate level energies and fine-structure splittings can be obtained only if, in addition to the static Coulomb repulsion among the electrons, the magnetic and retardation contributions to the e–e interaction, also known as the Breit interaction in atomic physics, as well as the self-energy and vacuum polarization of (at least) the inner-shell electrons, are also taken into account ([3, 4] and references therein). Many computations on the electronic structure of few-electron HCI have been carried out since then, which helped advance atomic many-body techniques into the best-approved theories of physics today. More recently, the e–e interaction has been explored in different collision [5, 6], autoionization [7] and ion recombination processes [8–13]. In the dielectronic recombination of HCI, in particular, it was shown that the Breit interaction can modify individual (capture) cross sections and rates by 50% or more [14–18], or may even dominate the Coulomb repulsion and finally lead to a qualitative change in the expected x-ray emission pattern [19, 20]. These investigations demonstrated that ion and x-ray spectroscopies of HCI now provide a unique tool for improving our understanding of the e–e and electron–photon interactions in the presence of strong fields [21].

Until now, however, very little has been known about how the magnetic and retarded interactions among electrons affect the electron emission and hence the dynamics of electrons in strong fields. While shifts in the transition energies [22] and cross sections [23] of the recorded ion and x-ray spectra clearly illustrate the importance of the relativistic contributions to e–e interaction, when averaged over space and time, further insight into the dynamics of electrons in strong Coulomb fields can be obtained only by measuring the angular and spin properties of the emitted electrons. Indeed, angle-resolved measurements on the electron emission from few-electron ions may help explore how electron dynamics is affected by the relativistic motion of electrons and the presence of a negative continuum. Such measurements can be seen as complementary to accurate hyperfine and g-factor experiments [24] that explore the magnetic interaction among the bound electrons and their interplay with the charge and magnetization distribution of the nucleus.

In this paper, we analyze the electron emission from HCI following the Coulomb excitation of (lithium- and) beryllium-like ions. For the $1s^22s^2\ ^1S_0 + Z_T \rightarrow 1s2s2p\ ^3P_2$ excitation of beryllium-like projectiles and their subsequent autoionization into the $1s^22s\ ^2S_{1/2}$ and $1s^22p\ ^2P_{1/2,3/2}$ levels especially (figure 1), an unexpected strong variation in the angular distribution
Figure 1. Level scheme of the excitation–autoionization process (1). The Coulomb excitation of initially beryllium-like projectiles in the 1s\(^2\)2s\(^2\) 1S\(_0\) ground state gives rise to ions in the 1s2s2p\(^3\) \(^3\)P\(_2\) resonant state that may autoionize into the 1s2s \(^2\)S\(_{1/2}\) and 1s2p \(^2\)P\(_{1/2},3/2\) levels of the finally lithium-like projectiles. The resonant and final levels of this process can be easily identified by the kinetic energy of the emitted electrons as well as their relative intensity, if necessary.

of the electrons is found, caused by the magnetic interaction among the electrons. In addition to the details of the e--e interaction, the predicted change in angular distribution depends sensitively also on the velocity of the projectile ions, thus enabling one to ‘tune’ the (relativistic) interaction for every ion along the isoelectronic sequence separately. Angle-resolved measurements for different ions therefore provide quite a sensitive tool for analyzing the electron dynamics in strong fields, complementary to the well-established spectroscopy of hard x-rays.

While the Coulomb excitation of HCl and their subsequent electron emission can be readily described by the density matrix theory [25], care has to be taken to select an excitation–autoionization process for which the electron angular emission is affected by the mixing between different partial waves of the outgoing electrons, thus enabling one to separate the magnetic and retarded interactions from the instantaneous Coulomb repulsion. For HCl, in fact, the effects of multipole mixing on the (Auger) electron emission are most pronounced for the (single-electron) excitation of beryllium-like ions, while helium-like ions cannot autoionize after a single 1s \(\rightarrow\) 2l\(j\) excitation and lithium-like ions autoionize into the 1s\(^2\) \(^1\)S\(_0\) ground state (of the finally helium-like ion) via a single partial wave. Since the e--e interaction operator is scalar, it cannot affect the (shape of the) angular distribution of the electrons emitted for a single partial wave, irrespective of all the details of the relativistic terms. This remains true even for the Auger decay of the meta-stable 1s2s2p \(^4\)P\(_{3/2}\) level, which is forbidden in the non-relativistic limit and has thus been used widely as a test bed for studying magnetic and retardation effects over the last few decades [26–28]. The electron angular distribution for the autoionization of the 1s2s2p \(^4\)P\(_{3/2}\) level into the 1s\(^2\) \(^1\)S\(_0\) ground state is defined entirely by geometrical factors that depend on the parities and total angular momenta of the initial and final ionic states, but is independent of the transition amplitude; see equations (3.1.15)–(3.1.16) of [25].
To work out further details of the electron emission of HCI, let us consider the excitation–autoionization process

$$A^{q+}(α_0J_0M_0) + Z_T \rightarrow A^{q++}(α_tJ_tM_t) \rightarrow A^{(q+1)+}(α_tJ_tM_t) + e^-_A \quad (1)$$

of a $q$-fold charged projectile in collision with some target atom with nuclear charge $Z_T$. In this process, the projectile is initially assumed to be in its ground state $|α_0J_0M_0\rangle$ with well-defined total angular momentum $J_0$, $M_0$ and parity, and is excited by the target (nucleus) into some state $|α_tJ_tM_t\rangle$ that is embedded within the continuum of the next higher charge state of the ion, $A^{(q+1)+}$. Thus, the projectiles can also autoionize under emission of an electron to some final state $|α_tJ_tM_t\rangle$, in addition to the typical photon emission of inner-shell excited HCI. In the notation of the atomic states, $α_0$, $α_t$, . . . hereby denote all additional quantum numbers needed for a unique specification of these states. The autoionization of ions in the resonant state $|α_tJ_tM_t\rangle$ is caused by e–e coupling to the continuum [29] and leads to electrons with well-defined kinetic energy $ε = E_r - E_t$ as obtained from the total energies of the resonance and final ionic states, respectively.

2. Theory and computations

A (time-independent) density operator can be assigned to the ion at each step of process (1), by taking into account the finally emitted electron by this operator. In first-order perturbation theory, these density operators before and after a given ‘interaction’ are related to each other by

$$ρ_k = R_kρ_{k-1}R_k^T,$$

where $R_k$ refers to the transition operator of the $k$th step and $ρ_k$ to the corresponding density operator just after the interaction has occurred. Complete transformation of the density operators for both steps of process (1), the initial excitation and the subsequent (auto-)ionization, therefore enables one to explore how the angular distribution (and polarization) of the emitted electrons is affected by the various interactions that occur during the time evolution of the system. Here, we shall not provide a detailed derivation of the final-state density operator that is presented elsewhere [30], but display only those formulae that are needed for the further discussion below. For the Coulomb excitation of an initially unpolarized ion in level $|α_0J_0\rangle$ into some magnetic substate $|α_tJ_tM_t\rangle$ with a specified projection $M_t$, the partial cross section is given by

$$σ(α_tJ_tM_t) = \frac{2π}{2J_0 + 1} \sum_{M_0} \int_0^\infty dB \ |A(b; M_0, M_t)|^2 , \quad (2)$$

where

$$A(b; M_0, M_t) = i γ_p α Z_T \int dt e^{i(E_t - E_0)t} \left( α_tJ_tM_t \sum_{i=1}^N \frac{1 - β_p \hat{α}_3(i)}{r_i(t)} \right) |α_0J_0M_0\rangle \quad (3)$$

is the transition amplitude in first-order perturbation theory [31, 32]. These formulae for the projectile excitation in relativistic ion–atom collisions are displayed in natural units ($\hbar = m_e = c = 1$), and here $b$ denotes the impact parameter, $\hat{α}_3(i) = \hat{α}_3(i)$ the Dirac matrix for the $i$th particle and $E_0$ and $E_t$ refer to the total energies of the projectile ion in its initial and intermediate resonance states, respectively. Furthermore, integration over time and the impact parameter $b$ can be carried out analytically in momentum space, and this helps simplify the computation of the amplitudes (3).
The partial cross section (2) describes the excitation of the projectile electrons into some (magnetic) substate \( |\alpha_r J_r, M_r\rangle \) and, if the cascade processes due to the excitation of high-lying levels can be neglected, therefore also defines the relative population of the substates after Coulomb excitation has taken place. For a fixed direction of the ion beam, however, this population is in general not statistically distributed but typically aligned, i.e. with an equal population only of just the magnetic substates with the same modulus \(|M_r|\). In atomic theory, this alignment of the projectile ions can be described most conveniently in terms of one (or just a very few) parameters \([25, 33]\)

\[
A_k(\alpha_r J_r) = \frac{\sqrt{2J_r + 1}}{\sigma(\alpha_r J_r)} \sum_{M_r} (-1)^{J_r - M_r} \langle J_r M_r J_f - M_f | k0 \rangle \sigma(\alpha_r J_f M_f),
\]

that are directly related to the partial cross sections (2), and where \( \sigma(\alpha_r J_r) = \sum_{M_r} \sigma(\alpha_r J_f M_f) \) is the total excitation cross section and \( \langle J_r M_r J_f - M_f | k0 \rangle \) a Clebsch–Gordan coefficient. For a beam of initially unpolarized ions, as can be seen from expression (4) and the symmetry of the Clebsch–Gordan coefficients, the alignment parameters \( A_k(\alpha_r J_r) \) are non-zero only for even values of \( k \leq 2J_r \).

In the second step of process (1), the ion does stabilize itself under the emission of an (Auger) electron with characteristic angular distribution and polarization. Of course, both these properties are related to the (sublevel) population of the excited ion and hence to the alignment parameters \( A_k(\alpha_r J_r) \). For example, the angular distribution of the emitted electron is given by

\[
W(\theta) \propto 1 + \sum_{k=2J_r, \ldots} A_k(\alpha_r J_r) f_k(\alpha_r J_r, \alpha_l J_l) P_l(\cos \theta),
\]

where \( \theta \) is the polar angle with regard to the beam direction and

\[
f_k(\alpha_r J_r, \alpha_l J_l) = \frac{1}{N} \sum_{\epsilon \kappa} i^{l-l'} e^{i\Delta_\epsilon - i\Delta_{\kappa'}} (-1)^{J_r + J_l + k - 1/2} \times \langle \epsilon \kappa J_f || V || \alpha_r J_r \rangle \langle \epsilon \kappa J_f || V || \alpha_l J_l \rangle^*
\]

are characteristic functions that describe the dynamics of the autoionization, and including the normalization factor \( N = \sum_{\epsilon} |\langle \epsilon \kappa J_f || V || \alpha_r J_r \rangle|^2 \). Indeed, the angular distribution (5) nicely reflects the two steps of process (1): the excitation of the ions, characterized by the alignment parameters \( A_k(\alpha_r J_r) \), and the subsequent autoionization as described by \( f_k(\alpha_r J_r, \alpha_l J_l) \), although these steps are coupled via the alignment of the intermediate resonance (\( \sum_k \)). The function \( f_k \) in equation (6) merely depends on the (reduced) matrix elements \( \langle \epsilon \kappa J_f || V || \alpha_r J_r \rangle \) of the e–e interaction \( V \), where \( \epsilon \) and \( \kappa = \pm (j + 1/2) \) for \( l = j \pm 1/2 \) describe the energy, angular momentum and parity of the emitted electrons, and \( \Delta_\epsilon \) denotes the total phase of the outgoing electron.

In the relativistic atomic theory, as appropriate for medium and heavy elements, the (frequency-dependent) e–e interaction

\[
V = V^C + V^B = \sum_{i<j} \left( \frac{1}{r_{ij}} - (\alpha_i \cdot \alpha_j) \frac{\cos(\omega r_{ij})}{r_{ij}} + (\alpha_i \cdot \nabla_i)(\alpha_j \cdot \nabla_j) \frac{\cos(\omega r_{ij}) - 1}{\omega^2 r_{ij}} \right)
\]

comprises both, the instantaneous Coulomb repulsion (the first term) and the Breit interaction, i.e. the magnetic and retardation contributions (the second and third terms). In this

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representation of the interaction operator, moreover, \( \omega \) denotes the frequency of the virtual photon and \( \alpha_i \), the vector of Dirac matrices associated with the \( i \)th particle. The e–e interaction operator (7) has been derived rigorously within the framework of quantum electrodynamics (3) and applied to a large number of electronic structure calculations in recent years [34, 35].

To analyze the angular distribution (5) of the emitted electrons in further detail, the Coulomb excitation amplitudes (3) and alignment parameters (4) as well as the characteristic functions (6) for the autoionization of the inner-shell excited projectiles have been calculated with multiconfiguration Dirac–Fock wave functions [36, 37]. Wave function expansions of increasing size were tested in order to ensure that the calculated transition amplitudes and alignment parameters converge within 5%. This level of convergence is sufficient to reveal and demonstrate the influence of the (frequency-dependent) Breit interaction on the angular distribution (and polarization) of emitted electrons.

### 3. Results and discussion

Lithium-like ions are probably the simplest systems in which a \( 1s \rightarrow 2l j \) (single-electron) excitation leads to the emission of electrons. In particular, the \( 1s2s2p \ 4P \) term has received much interest in the literature [26, 38, 39] for the study of spin–orbit and relativistic interactions in (spin-) forbidden transitions, both by photon and electron spectroscopy. Even for lithium-like uranium, for example, the \( 1s2s2p \ 4P_{5/2} \) level has a lifetime that is roughly two orders of magnitude larger than that for the \( 2,4P_{1/2,3/2} \) levels from the same configuration. Nonetheless, the electron emission from all \( 2,4P_j \) levels of \( 1s \rightarrow 2p \) inner-shell excited lithium-like ions does not enable one (so easily) to distinguish different contributions to the e–e interaction since the \( 1s^2 \ 1S_0 \) ground state of the (finally) helium-like ion only allows emission of a single partial wave, i.e. \( \kappa = \kappa' \) in equation (6). For this reason, no interferences of different amplitudes due to the relativistic e–e interaction may arise in the angular distribution of the emitted electrons. This can be seen, for instance, from the properties of the \( 6j \) symbols in equation (6) which, for \( J_i = 0 \), are non-zero only for \( j = j' = J_f \) and \( l = l' \) and this reduces the characteristic function \( f_k \) to a purely geometrical factor, independent of any further detail with regard to the e–e interaction operator (7).

Inner-shell \( 1s \rightarrow 2p \) excited beryllium-like ions still have a reasonably simple level structure (figure 1) that helps distinguish individual Auger lines rather readily. For these ions, the \( 1s(2s^22p + 2s2p^2 + 2p^3) \) configurations with a single hole in the \( K \)-shell comprise in total 30 fine-structure levels \( 2S + 1L_J \) with \( J = 0, \ldots, 3 \). For medium- and high-\( Z \) ions along the beryllium sequence, all these levels have a comparable excitation energy that is well separated from the \( 1s \rightarrow 3l j \) excitations. Moreover, since the Coulomb excitation is caused by the one-electron operator \( \sum_j 1/r_j(t) \) in equation (3), mainly the \( 1s2s2p \) levels are excited, while the population of the \( 1s(2s2p^2 + 2p^3) \) levels is clearly suppressed or even negligible. In the following, we are especially interested in the Coulomb excitation and subsequent autoionization of the \( 1s2s^22p_{3/2} \) \( 3P_2 \) level, which does not decay by \( E1 \) dipole emission into the \( 1s2s^2 \ 1S_0 \) ground level but only by magnetic quadrupole (M2) radiation. For the Coulomb excitation of this \( J = 2 \) level, figure 2 displays the cross sections \( \sigma \) and alignment parameters \( A_2 \) and \( A_4 \) from equation (4) as a function of the projectile energy between 5 and 100 MeV u\(^{-1}\) and for the three beryllium-like ions with \( Z = 54, 79 \) and 92, respectively.

In fact, the strong energy dependence of both the cross sections and the alignment parameters, already enables one to ‘tune’ the importance of different terms in the angular
Figure 2. Excitation cross section and alignment parameters for the Coulomb excitation of the 1s2s^{2}2p_{3/2} \ ^3P_{2} resonance as a function of the projectile energy. Results are shown for the cross section \( \sigma \) (left panel), the alignment parameter \( A_2 \) (middle panel) and \( A_4 \) (right panel), and for three selected beryllium-like ions with \( Z = 54 \) (dashed lines), \( Z = 79 \) (dashed-dotted lines) and \( Z = 92 \) (solid lines), respectively.
Table 1. Characteristic functions $f_2$, $f_4$ and branching fraction (bf) for the autoionization of the $1s2s^22p_{3/2} J = 2$ resonance into the $1s^22s^22S_{1/2}$ ground and $1s^22p^23P_{1/2}$ first excited levels of lithium-like ions. Results are displayed without (Coulomb only) and with including the magnetic and retardation contributions (Coulomb + Breit) in the Auger amplitudes. See text for a further discussion.

| $Z$ | Coulomb only | Coulomb + Breit |
|-----|--------------|-----------------|
|     | $f_2$       | $f_4$          | $bf$         | $f_2$       | $f_4$          | $bf$         |
| 54  | -0.837      | < 0.01         | 0.238        | -0.849      | 0.053          | 0.242        |
| 79  | -0.837      | < 0.01         | 0.083        | -0.858      | 0.093          | 0.091        |
| 92  | -0.837      | < 0.01         | 0.060        | -0.862      | 0.112          | 0.071        |
|     | $1s^22s^22p_{3/2} \rightarrow 1s^22s^22S_{1/2}$ |
| 54  | -0.844      | 0.035          | 0.012        | -0.906      | 0.308          | 0.010        |
| 79  | -0.847      | 0.048          | 0.004        | -0.981      | 0.647          | 0.003        |
| 92  | -0.851      | 0.065          | 0.001        | -1.030      | 0.863          | 0.001        |

The outgoing electron leaves the ion predominantly in a $p_{3/2}$ partial wave with the amplitude $a_{p_{3/2}} \equiv \langle \ell S_{1/2}, \epsilon p_{3/2} J = 2 || V || 1s^22s^22p_{3/2} 3P_2 \rangle$ that is two orders of magnitude larger than the Auger amplitude $a_{f_{5/2}} \equiv \langle \ell S_{1/2}, \epsilon f_{5/2} J = 2 || V || 1s^22s^22p_{3/2} 3P_2 \rangle$ for an $f_{5/2}$ partial wave. The amplitude $a_{f_{5/2}}$ even vanishes to a very good approximation for a pure Coulomb repulsion among the electrons, $a_{f_{5/2}} \approx 0$, owing to the tensorial structure of this operator that requires a bound $f$-orbital to admix into the expansion of the $1s^22s^2p_{3/2} 3P_2$ wave function. With these definitions of $a_{p_{3/2}}$ and $a_{f_{5/2}}$, the characteristic functions for the autoionization into the $1s^22s^22S_{1/2}$ ground level can be simplified to

$$f_2 \approx -\frac{\sqrt{7}}{10} \left( 1 + 2\frac{\sqrt{6}}{7} \frac{a_{f_{5/2}}}{a_{p_{3/2}}} \cos(\Delta f_{5/2} - \Delta p_{3/2}) \right),$$

$$f_4 \approx 4\frac{3}{7} \frac{a_{f_{5/2}}}{a_{p_{3/2}}} \cos(\Delta f_{5/2} - \Delta p_{3/2}),$$

where $\Delta_e$ refers again to the phase of the outgoing electron. In single-configuration approximation and for a pure Coulomb interaction, $a_{f_{5/2}} = 0$, we therefore obtain $f_2 = -\sqrt{7/10} = -0.837$ and $f_4 = 0$ for all ions along the beryllium isoelectronic sequence, as seen from Table 1. If $a_{f_{5/2}}$ becomes non-zero because of the magnetic and retardation contributions to the e-e interaction, the two characteristic functions $f_2$ and $f_4$ then depend not only on the ratio between the $a_{f_{5/2}}$ and the (dominant) $a_{p_{3/2}}$ amplitudes but also on the phase difference of the involved partial waves, $\cos(\Delta f_{5/2} - \Delta p_{3/2})$. For the autoionization into the $1s^22s^22S_{1/2}$ ground state, the relativistic interactions lead to a 3% reduction of $f_2$ for $Z = 92$, while $f_4$ changes from nearly zero (for a pure Coulomb interaction) to about 0.11 if the complete e-e interaction is taken into account.

Since mainly the characteristic function $f_4$ is affected by the magnetic and retardation contributions to the e-e interaction, we need to choose projectile energies for which a sufficiently large product $A_4 \cdot f_4$ is predicted. A sizeable alignment parameter $A_4$ (in magnitude)
Figure 3. Angle-differential cross section for the electron emission of the 1s2s2p3/23P2 → 1s2s2S1/2 autoionization of beryllium-like ions with projectile energies \( T_p = 5, 10 \) and 100 MeV u\(^{-1}\), following the Coulomb excitation from their 1s2s2 \(^1S_0\) ground state. Results are shown in the rest frame of three beryllium-like projectiles with charges \( Z = 54, 79 \) and 92 and in two approximations. Angular distributions with only the Coulomb repulsion incorporated into the Auger amplitudes (blue dashed lines) in expression (6) are compared with those where the complete e–e interaction is taken into account (black solid lines).
increases steadily, owing to enhancement of the alignment parameter $A_2$, but with less effect than from the magnetic and retardation contributions to the e–e interaction, see figure 2.

A similar analysis as for the autoionization of the $1s^22s^22p_3/2\,^3P_2$ level into the $1s^22s\,^2S_{1/2}$ ground level of the lithium–like ion can be performed for its autoionization into the $1s^22p\,^2P_{1/2}$ lowest excited level of these ions. In this decay mode, the outgoing electron leaves the ion in either a $d_{3/2}$ or $d_{5/2}$ partial wave, with preference for a $d_{5/2}$ emission. In contrast to the autoionization into the $1s^22s\,^2S_{1/2}$ ground level, however, a non-zero $f_4$ value arises already for a pure Coulomb repulsion owing to the spin–orbit splitting of the relativistic 2p orbital. Nevertheless, the magnetic contributions to the e–e interaction have a remarkable effect on this electron line and enlarge the $f_4$ function by more than a factor of ten for all beryllium–like projectiles with nuclear charge $Z \gtrsim 67$ (table 1). This enhancement in the function $f_4$ gives rise, together with a small change of the $f_2$ function, to a strong alteration of the electron angular distribution, as shown in figure 4. Again, a double-peak structure and a reduced electron yield in the forward direction by up to a factor of 2 is predicted due to the magnetic and retarded interactions among the electrons. Let us note here, however, that autoionization into the $1s^22p\,^2P_{1/2}$ lowest-excited level is likely more difficult to observe because the branching fraction (relative intensity) of the $1s^22s^22p_3/2\,^3P_2 \rightarrow 1s^22p\,^2P_{1/2}$ line is much smaller than for the autoionization into the $1s^22s\,^2S_{1/2}$ ground level, see table 1.

The proposed excitation–autoionization scheme (1) facilitates, in contrast to electron capture, a clear distinction of the emitted electrons from the initial (Coulomb) excitation.
process. This scheme can be realized most easily at ion storage rings such as the experimental storage ring (ESR) in Darmstadt. The electrons from the projectiles are then measured in the laboratory frame and thus the predicted angular distributions need to be Lorentz transformed first before they can be compared with experiment. For example, figure 5 displays the angular distribution of the electrons emitted from beryllium-like U\(^{88+}\) projectiles at 5 MeV u\(^{-1}\) in the 1s2s\(^2\)2p\(^3/2\) \(^3\)P\(_2\)−1s2s \(^2\)S\(_{1/2}\) (left panel) and 1s2s\(^2\)2p\(^3/2\) \(^3\)P\(_2\)−1s2p \(^2\)P\(_{1/2}\) (right panel) autoionization. As seen from this figure, in particular the electron emission in the forward direction appears to be very sensitive with regard to relativistic contributions to the e–e interaction. Here, zero-angle electron spectrometry has been found sensitive in electron–photon coincidence measurements to resolve details at the high-energetic end of the bremsstrahlung spectrum [40]. Therefore, measurements of the electron emission in the forward direction, say between \(0 < \theta \lesssim 30^\circ\) in the laboratory frame, will enable one to resolve the predicted decrease in electron yield at low angles. As seen from figure 5, already a peak structure in the forward direction would provide a clear signature of the Breit interaction, and further details will be revealed if the range of measured angles is enlarged towards 90\(^\circ\) or even beyond.

In fact, the proposed measurements appear to be feasible with present-day available spectrometers at heavy-ion storage rings or those, for instance, planned at the future FAIR facility. The Coulomb excitation of high-Z projectile ions has been explored in good detail for initially hydrogen- and helium-like ions [23, 41]. In these measurements, it was found that the alignment of the inner-shell excited states itself is sensitive to magnetic interactions and, hence, to details of the Lienard–Wiechert potential. A relativistically modified alignment of such inner-shell excited states was also seen in the angular distribution of the emitted electrons, but these effects on the alignment can be separated (to a good extent) from the interaction of the emitted electrons and the bound-state density by varying the charge state of the projectiles. In beryllium-like ions, moreover, the 1s2s\(^2\)2p\(^3/2\) \(^3\)P\(_1\) resonance that lies quite close in energy to the \(^3\)P\(_2\) level (by about

Figure 5. Angle-differential cross section for the electron emission of the 1s2s\(^2\)2p\(^3/2\) \(^3\)P\(_2\)−1s2s \(^2\)S\(_{1/2}\) (left panel) and 1s2s\(^2\)2p\(^3/2\) \(^3\)P\(_2\)−1s2p \(^2\)P\(_{1/2}\) (right panel) autoionization of U\(^{88+}\) projectiles with energy \(T_p = 5\) MeV u\(^{-1}\). Results are shown in the laboratory frame and by incorporating only the Coulomb repulsion into the Auger amplitude (blue dashed lines) in expression (6) as well as for a full account of the e–e interaction (black solid lines).
32 eV for Xe$^{50+}$ and 60 eV for U$^{88+}$) and which also receives a substantial population due to the Coulomb excitation. The autoionization of the $^1P_1$, $^3P_1$, $^1S_0$, $^3P_1$, $^1S_0$, and $^3P_2$ levels may therefore partially blend the electron emission from the $1s2s^22p_{3/2}^3P_2$ resonance. Owing to the E1 dipole decay into the $1s^22s^21S_0$ ground state, however, the $^1P_1$ level has a 60 times shorter lifetime than $^3P_2$ and only a very tiny branching fraction for autoionization. In addition, cascade contributions to the population of the $^3P_2$ level mainly occur via the Coulomb excitations of the $1s2s^2(3d + 4d + 4f)$ levels but were found negligible at the present level of accuracy.

4. Conclusion

The electron emission of HCI has been reconsidered for unraveling the influence of the magnetic and retarded interaction among the electrons on the electron dynamics in strong fields. A remarkable change in the predicted angular distribution of the electrons is found for the Coulomb excitation of beryllium-like ions in the $1s2s^22p_{3/2}^3P_2$ resonance and the subsequent autoionization into the $1s^22s^21S_0$ and $1s2s^22p_{1/2}$ levels of the finally lithium-like ions. When compared with the radiative or dielectronic capture of electrons, the Coulomb excitation is particularly suitable for studying the electron emission in strong fields as it is not disturbed by free (incident) electrons. The present analysis of the relativistic contributions to the interaction extends a previous study of the photon emission of HCI in which the Breit terms also lead to a different behavior in the angular distribution and polarization of the emitted x-rays [19, 42] when compared with the Dirac–Coulomb theory, which was recently confirmed by experiment [20]. In contrast to photon emission, which only exhibits the integral effect of the relativistic motion of electrons, electron spectrometry enables one to access further details of the electron dynamics in strong Coulomb fields.

The simplest signatures of the relativistic contributions to the interaction in high-$Z$ ions are the reduced electron emission in the forward direction $(\theta \lesssim 5^\circ)$ as well as the double-peak structure in the expected angular distribution; these signatures arise especially at low projectile energies $\lesssim 10 \text{ MeV u}^{-1}$ and for beryllium-like ions with nuclear charge $Z \gtrsim 70$. The electron angular distribution from such projectiles can be analyzed with present-day electron spectrometers and provides complementary information about the electron dynamics in strong fields that is not accessible from x-ray spectra alone.

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