APHYSTHE

Theory for Atomic Physics with Heavy Ions

Radiative electron capture into heavy, lithium–like ions

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During the last decade, relativistic collisions of highly–charged ions with electrons and low–Z target atoms have been studied intensively at the GSI storage ring in Darmstadt. In such collisions, one of the dominant processes is the radiative electron capture (REC) of either a target or cooler electron into a bound state of the projectiles which is accompanied by the emission of photons. So far, however, most experimental REC studies concerned the capture of electrons by *bare* projectile ions, including measurements of the total and angular–differential cross sections [1]. Usually, good agreement is found when the cross sections are compared with computations based on Dirac's theory [2]. Less attention, in contrast, was paid to the electron capture by *few–electron* projectiles owing to the fine–structure effects for ions with (initially) open shells which arise from the interaction among the electrons. A first step towards the study of *interelectronic* effects has been done only recently by Bednarz and co–workers [3], who measured the angular distributions of the x–ray photons following the REC by *hydrogen–like*, *helium–like* and *lithium–like* uranium ions. In this experiment, rather high collision energies ($T_p > 200$ MeV/u) was employed for which the radiative recombination is not sensitive to the electron–electron repulsion. However, further studies on the REC angular distributions are planned at the GSI storage ring, and will be focused on the use of *decelerated* ions, for which one expects more pronounced effects from the interaction of the electrons.

In this contribution, we report about detailed calculations of the angle–differential REC cross sections for the capture into *few–electron* heavy ions, performed for a wide range of collision energies. In order to explore the influence of the interelectronic–interaction on the angular distributions of the emitted photons, two theoretical approaches were applied: In the independent–particle model (IPM), we first made use of hydrogenic functions and proper set of Slater determinants to account for the Pauli principle in the capture of (additional) electron. Apart from this simple "one–electron" model, the angular distributions of the recombination photons have been also calculated by means of the Multi–configuration Dirac–Fock (MCDF) approach which allows to incorporate the electron–electron interaction effects in a more systematic fashion. Within the MCDF approach, both the initial and the final ionic states with angular momentum and parity $\left(J^P\right)$ are approximated by a linear combination of the so–called configuration state functions (CFS) of the same symmetry:

$$|\psi_\alpha(PJM)\rangle \;=\; \sum_{r}^{n_c} c_\gamma(\alpha)\;|\gamma_r(PJM)\rangle\;, \qquad (1)$$

where n_c is the number of CSF and $c_r(\alpha)$ denotes the representation of the ionic state in this basis. As usual, the CFS are antisymmetrized products of a common set of *orthonormal* orbitals which are optimized on the basis of the Dirac–Coulomb Hamiltonian.

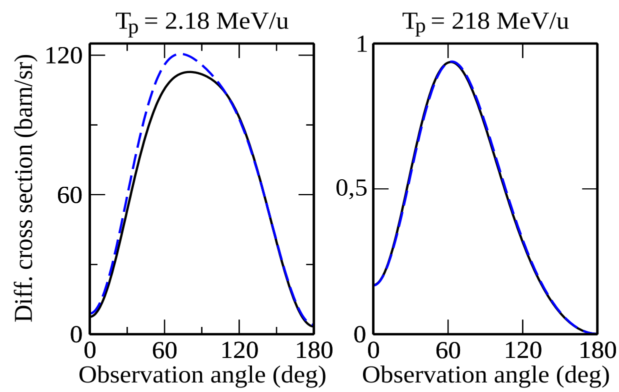

Figure 1: Angle–differential cross sections for the radiative electron capture into the $1s^2\ 2s^2$ state of (initially) lithium–like uranium ions U^{89+} with projectile energies of 2.18 and 218 MeV/u. Results are presented within the independent particle model (- - -), and Multi–configuration Dirac–Fock approach (—).

The MCDF method is implemented, for instance, in the RATIP [4] program and has been utilized to calculate the *bound–free* transition amplitude and, hence, the angular distributions of the recombination x–ray photons [5]. Figure 1 shows the angle–differential cross sections for the radiative capture of electrons into the $1s^2\ 2s^2$ state of (initially) lithium–like uranium ions U^{89+} with projectile energies $T_p = 2.18$ and 218 MeV/u and compares them with the cross section data as obtained in the IPM. For sufficiently large projectile energies, such as $T_p = 218$ MeV/u, both, the "one–electron" and "many–electron" approximations yield virtually identical results, which shows that the electron–electron interaction effects become negligible for high collision energies. For *decelerated* ions with an energy $T_p = 2.18$ MeV/u, in contrast, the electron–electron correlation alters the angular distribution of the recombination photons considerably with a rather strong effect around the angle $\theta = 70°$ in the laboratory system [5].

In the present work, we explored the influence of the electron–electron correlations on the *angle–differential* REC cross sections. More pronounced effects due to the interelectronic–interactions are expected, however, for the (linear) *polarization* of the emitted radiation for which a first analysis is currently performed by combining the MCDF approach with the density matrix theory [6].

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Excitation of heavy nuclei by electron capture

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In the resonant process of nuclear excitation by electron capture (NEEC) a free electron is captured into a bound shell of an ion with the simultaneous excitation of the nucleus. This is the time-reversed process of the internal conversion. The excited nuclear state can then decay radiatively by the emission of a photon or by internal conversion. The process is similar to the dielectronic recombination (DR) in which the captured electron transfers its energy to the excitation of a bound electron. This electronic excitation is now replaced by the excitation of the nucleus.

The investigation of NEEC opens the possibility to explore the spectral properties of heavy nuclei through atomic physics experiments.

Considering the analogy with DR we have extended the formalism developed in [1] for calculating the DR cross section, so that it accounts for nuclear transitions. Zimmerer [1] has applied a relativistic formalism for expanding the transition operator introducing Feshbach projector operators to clearly separate the different subspaces involved in this process.

In the following we consider that the electron is captured to a bound state in the Coulomb field of a bare nucleus, so that no competition with the DR process occurs. The cross section for the NEEC process, followed by the radiative deexcitation of the nucleus, reads:

$$\sigma = \frac{2\pi}{p^2}\frac{A_r Y_n}{\Gamma_d} L_d(E-E_d). \qquad (1)$$

Here, $L_d(E-E_d)$ is the Lorentz resonance profile, centered on the resonance energy E_d and having the total natural width Γ_d of the excited nuclear state. Further, p and E are the momentum and the energy of the continuum electron, respectively, A_r is the radiative decay rate of the excited nuclear state and Y_n the NEEC rate, which was calculated in the Coulomb gauge according to the formula

$$Y_n = \frac{8\pi^2 e^2}{(2L+1)^3} R_0^{-2(L+2)}$$

$$\times \sum_\kappa |R_{L,j_f,j_i}|^2 B(EL)(2j_f+1) C\left(j_f\ L\ j_i; \frac{1}{2}\ 0\ \frac{1}{2}\right)^2. \qquad (2)$$

For describing the nucleus we have used a nuclear collective model [3] which assumes irrotational flow of the nuclear currents. This model accounts only for the excitation of low lying nuclear levels by electric multipole transitions. Here, L denotes the multipolarity of the nuclear transition, $B(EL)$ is the reduced electric nuclear transition probability and R_0 the nuclear radius. In the numerical calculations we have used experimental values for $B(E2)$ and A_r taken from [2]. The electronic state is labelled by $\kappa = (-1)^{j+l+1/2}(j+1/2)$, with the total angular momentum j and the orbital angular momentum l. The indices i and f stand for the initial, respectively final electronic states. R_{L,j_f,j_i} is the radial integral given by

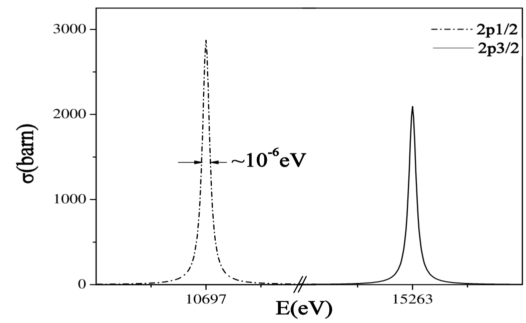

Figure 1: NEEC cross section for ^{238}U with excitation of the first nuclear 2^+ state

$$R_{L,j_f,j_i} = \frac{1}{R_0^{L-1}}\int_0^{R_0} dr r^{L+2}\left[f_f(r)f_i(r)+g_f(r)g_i(r)\right]$$

$$+R_0^{L+2}\int_{R_0}^{\infty} dr r^{-L+1}\left[f_f(r)f_i(r)+g_f(r)g_i(r)\right], \qquad (3)$$

where $g(r)$ and $f(r)$ are the large and small radial components of the relativistic electron wavefunction, respectively. These integrals were calculated numerically.

For the numerical calculation we have considered the case of ^{238}U with its $E2$ transition from the ground state to the first 2^+ state at the excitation energy of 44.913 keV. This energy value allows the excitation of the nucleus by capturing a continuum electron into the L atomic shell. The transition rate of capturing the free electron into the $2p_{1/2}$ or $2p_{3/2}$ orbitals is $\sim 10^{10}s^{-1}$, two orders of magnitude larger than that for the capture in the $2s_{1/2}$ orbital, which is $\sim 5\cdot 10^8 s^{-1}$. Calculations for ^{236}U produced very similar results.

Figure 1 shows the NEEC cross sections for the $E2$ transitions in ^{238}U, with the capture of the free electron in the $2p_{1/2}$ and $2p_{3/2}$ orbitals of the L shell, followed by the radiative decay of the excited nuclear state. The resonance energies of the continuum electron are 10.697 and 15.263 keV, respectively. The resonance strength is 10^{-2} barn·eV in the case of the $2p_{1/2}$ capture and $7\cdot 10^{-3}$ barn·eV in the case of the $2p_{3/2}$ capture. These values are 5-6 orders of magnitude smaller than the resonance strength for DR, making NEEC challenging to be observed experimentally. Calculations of magnetic nuclear transitions induced by the current of the electron are in progress.

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Angular correlations in the two–photon decay of highly–charged ions

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During recent years, a number of experiments have been performed at the GSI facilities in Darmstadt in order to explore the two–photon decay in medium and heavy ions [1, 2]. When compared to theoretical predictions, the outcome of these measurements provides a unique possibility to probe the influence of relativistic as well as quantum electrodynamical (QED) effects on the electronic structure of ions (atoms). Until now, most of the experimental and theoretical studies focused on the total decay rates as well as the energy distributions of the emitted photons. In contrast, less attention, has been paid so far to the effects of relativity on the photon–photon *angular correlation* function, that is if the angular distribution of one of the photons is observed relative to the emission of the second one in a two–photon transition. However, with the recent progress in x–ray detector design, measurements on the angular correlations in the two–photon decay of heavy atomic systems are likely to be become possible within the next few years.

To understand the details of the two–photon decay, we analyzed the photon–photon angular correlation function for highly–charged, hydrogen–like ions in the framework of second–order perturbation theory and Dirac's equation [3, 4]. Within this framework, the photon–photon angular correlation function can be traced back to the evaluation and computation of the second–order transition amplitude:

$$M_{fi} = \sum\!\!\!\!\!\!\int_\nu \frac{\langle\psi_f\,|\mathbf{A}_1^*|\,\psi_\nu\rangle\,\langle\psi_\nu\,|\mathbf{A}_2^*|\,\psi_i\rangle}{E_\nu - E_i + E_{\gamma_2}} + \frac{\langle\psi_f\,|\mathbf{A}_2^*|\,\psi_\nu\rangle\,\langle\psi_\nu\,|\mathbf{A}_1^*|\,\psi_i\rangle}{E_\nu - E_i + E_{\gamma_1}}\,, \qquad (1)$$

where $(\psi_i,\,E_i)$, $(\psi_\nu,\,E_\nu)$ and $(\psi_f,\,E_f)$ refer to the *relativistic* wavefunctions and the energies of the initial, intermediate and final atomic states, respectively. In this expression, as usual, the electron–photon interaction is described in terms of the transition operator $\mathbf{A} = \boldsymbol{\alpha}\,\mathbf{u}_\lambda\,e^{i\boldsymbol{k}\cdot\boldsymbol{r}}$ which includes the unit vector $\mathbf{u}_\lambda$ in order to describe the polarization of the emitted photons.

As indicated in equation (1), the summation over the intermediate states refers to the *complete* one–particle spectrum (ψ_ν, E_ν), including a summation over the discrete part of the spectrum as well as the integration over the continuum. Since such a summation over the *complete* spectrum is not feasible in practice, a number of alternative methods has been proposed for calculating second–order amplitudes. In our work, for example, we first express the transition amplitude (1) by means of relativistic Coulomb–Green's functions [3] in order to avoid an explicit summation over the spectrum. For the hydrogen–like ions, the Coulomb–Green's functions are known *analytically* and allow for an accurate and fast computation of the transition amplitude (1) and, hence, of the photon–photon angular correlation function.

Apart from the summation over the complete hydrogen spectrum (ψ_ν, E_ν), the transition amplitude (1) includes formally also the – infinite – summation over (the products of) the multipole photon fields, E1E1, E1M2, M1M1, E2M2, ... which arise from the decomposition of the photon plane wave $\mathbf{u}_\lambda\,e^{i\boldsymbol{k}\cdot\boldsymbol{r}}$ into its multipole components. However, often, this summation can be restricted to the 'leading' term which usually dominates the two–photon transitions for any given pair of initial and final bound states. For example, the $2s_{1/2}$ and $3d_{5/2}$ levels decay into the $1s_{1/2}$ ground state primarily by the emission of two electric–dipole (E1E1) photons, while all the higher multipoles contribute with less than 0.5 % to the total decay rate. Nevertheless, a significant effect from the higher multipoles is expected for the *angular* distributions of the emitted photons. Figure 1 displays this (higher multipole) effect for the $2s_{1/2} \rightarrow 1s_{1/2}$ and $3d_{5/2} \rightarrow 1s_{1/2}$ two–photon decay of hydrogen–like uranium U^{91+} ions. As seen from this Figure, the photon–photon angular correlation function, which is found to be symmetric with respect to the angle $\theta = 90°$ in the electric dipole (E1E1) approximation, becomes asymmetric due to the higher (non–dipole) components of the radiation filed, and this effect is even more pronounced for the decay of the $3d_{5/2}$ state.

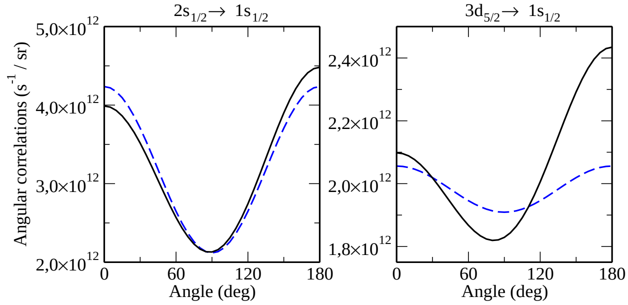

Figure 1: Photon–photon angular correlations in the $2s_{1/2} \rightarrow 1s_{1/2}$ and $3d_{5/2} \rightarrow 1s_{1/2}$ (two–photon) decay of hydrogen–like uranium U^{91+} ions. Results are presented for the exact relativistic theory (—) and the relativistic electric dipole approach (– –) for two emitted photons with energies $E_{\gamma_1} = E_{\gamma_2}$.

Until now, we have restricted our calculations of the photon–photon angular correlations to *hydrogen–like* ions. However, the Green's function approach which has been successfully applied for these (one–electron) ions, can also be extended for the study of the two–photon decay in many–electron heavy ions. First investigations of the $2^1S_0 \rightarrow 1^1S_0$ (two–photon) decay in the helium–like ions are currently under way.

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Angular distribution of the cascade photons in the radiative electron capture

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Recent experiments on the cascade photon distribution in the collisions of uranium ions with light targets for the bare uranium nuclei U^{92+} and the one-electron ions U^{91+} in the ground states provided different results. In the former case distributions of the photons emitted in the decay of intermediate $2p_{3/2}$ state exhibit strong angular dependence. In the latter case the angular dependence of the superposition of the emission of the corresponding 2^1P_1 and 2^3P_2 states is very weak [1, 2]. In this note I present a theoretical explanation of this phenomena.

In the case of bare nucleus U^{92+} all the electron states can be described by the relativistic Coulomb field functions (RCFF). There are different forms for presentation of the RCFF. In [3] the RCFF are presented as a fast converging $(\alpha Z)^2$ series. We carried out the calculations by using several lowest terms of the expansion [3]. This confirmed previously obtained experimental and theoretical results on the angular distribution. The latter is shown to vary by about 50% in the whole interval of the values of angles.

In the case of capture by a single-electron ion U^{91+} we must build the two-electron wave functions. Since the electromagnetic interaction between the electrons is $1/Z \sim 10^{-2}$ times smaller then their interaction with the nucleus, the asymmetric combination of the single-electron RCFF functions is a reasonable approximation. In other words, we neglect the electromagnetic interaction between the electrons, while the spin correlations are included (Pauli principle is respected).

By using presentation [3] for RCFF we obtain an analytical expression for the amplitude. The distribution of the photons emitted in the decays of the intermediate states 2^1P_1 and 2^3P_2 treated separately still exhibit strong angular dependence. However there is a large cancellation of their angular dependent contributions to the superposition of the distributions which is detected in the experiment. The final result is that the angular dependence should not vary by more then 10% in the observable superposition of the decays of the two states.

Thus we conclude that the observed difference of the angular distributions for the capture by bare nucleus of U and by its single-electron ion is due to the Pauli principle.

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QED Calculations of Bound-Electron g Factors in Highly-Charged Ions

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Highly-charged ions or simple atoms offer an opportunity for high-accuracy calculations of atomic properties within the ab-initio framework of bound-state QED. Such systems possess relatively simple spectra and some of their transitions can be measured with very high precision allowing for sensitive tests of both the magnetic and the electric sector of QED. The remarkable progress made in experimental determinations of Lamb shifts and bound-electron g factors (see e.g. Refs. [1, 2] and references therein) together with recent extensions and further developments of the calculation methods in bound-state QED of atomic systems [3, 4] facilitates independent precision determinations of phenomenological parameters, such as the electron mass m_{e}, the fine-structure constant α and electron g factors, or nuclear magnetic moments, nuclear radii and ion masses.

The most recent experimental determination of the magnetic moment of the bound electron in hydrogen-like oxygen [2] may exemplify the present situation: The experimental value for the g factor of the bound $1s$-electron in an $^{16}\mathrm{O}^{7+}$ ion, $g_{1s}^{\mathrm{exp}} = 2.000\,047\,025\,4(15)(44)$, agrees within 1.1σ with the predicted value $g_{1s}^{\mathrm{theory}} = 2.000\,047\,020\,2(6)$ and thus provides a stringent test of bound-state QED to the level of 0.25%. Moreover, assuming the validity of the underlying theory allows for the determination of the most precise value for the electron mass $m_{\mathrm{e}} = 0.000\,548\,579\,909\,6(4)$ u [2]. The experiments were performed with single H-like ions (with spin-zero nuclei) confined in a Penning trap measuring the ratio of the Larmor frequency ω_{L} of the electron and the cyclotron frequency ω_{c} of the ion, which is related to the electron g factor via

$$\frac{\omega_{\mathrm{L}}}{\omega_{\mathrm{c}}} = g\,\frac{|e|}{m_{\mathrm{e}}}\,\frac{m_{\mathrm{ion}}}{q_{\mathrm{ion}}}\,. \qquad (1)$$

Here e denotes the elementary charge unite, m_{ion} and q_{ion} is the mass and the charge of the ion, respectively.

From experimental side further progress and extensions of the measurements, e.g., to Li-like ions and to H-like ions with nonzero nuclear spin is anticipated in the near future. A major advantage of comparing Li-like and H-like ions with the same nucleus lies in the fact that uncertainties due to nuclear effects can be significantly reduced in the specific difference g' between the corresponding g factors

$$g' = g_{(1s)^2 2s} - \xi g_{1s}\,. \qquad (2)$$

The parameter ξ is calculated numerically. In Ref. [5] we have provided most accurate calculations of various relativistic and QED corrections to the g factor of Li-like ions in the region of nuclear charge numbers $Z = 6 - 92$. The interelectronic-interaction correction of order $1/Z^2$ and higher obtained within the configuration-interaction Dirac-Fock method is combined with the $1/Z$-correction term derived rigorously from QED. Improved results for the one-electron QED corrections of first and second order in α together with screened QED corrections have been taken into account. The Mainz-GSI collaboration is aiming for the determination of the g factor, in particular, for Li-like calcium with anticipated experimental accuracy at the level of ppb. Corresponding data could be compared with recent predictions provided in [5] (see Table 1 as an example). Extensions of direct measurements of the electron g factor of heavy H-like ions with odd nuclei in the HITRAP facility at GSI seems very promising for tests of the magnetic sector of QED. Moreover, such results can be utilized for independent, precise determinations of nuclear properties, such as magnetic moments. Accordingly, a fully relativistic theory of the g factor of H-like ions with nuclear spin $I > 0$ has been developed [6]. The total magnetic moment of an ion state with total angular momentum F thus contains informations about the electronic and nuclear g factors, respectively. In [6] calculations for the g factor of ions in the ground state ($F = I \pm 1/2$) have been performed for a variety of ions. They account for corrections resulting from QED, nuclear recoil, finite-nuclear size, magnetic dipole, electric quadrupole and hyperfine interactions. To provide accurate theoretical results suitable for the analysis of corresponding data obtained in experiments with highly-charged ions at the GSI facilities represents an important goal of future theoretical investigations at the frontier of atomic and nuclear physics.

Table 1: Contributions to the ground-state g factor of Li-like calcium.

Dirac value (point nucleus)	1.996 426 011
Finite-nuclear size	0.000 000 014
Interelectronic interaction	0.000 454 45(14)
one-loop QED $\sim \alpha$	0.002 325 555(5)
one-loop QED $\sim \alpha^2$	-0.000 003 517(0)
Screened QED	-0.000 000 33(10)
Nuclear recoil	0.000 000 61(2)
Total	1.999 202 24(17)

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Dynamical Schwinger effect in the focus of optical laser beams

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The Schwinger effect of vacuum pair creation by a quasi-classical electric field although predicted long time ago has still not been experimentally verified since critical field strengths $E_{\rm cr} = m^2/e = 10^{16}$ V/m could not be reached. With the planned construction of X-ray free electron lasers producing field strengths close to the critical one, the possibility to prove the Schwinger effect attracts attention again. Here we consider the region of parameters already achievable with the operating SLAC and Jena Terawatt laser systems [1, 2] as well as the PHELIX laser [3] under construction at GSI, where $m\omega \ll eE \ll m^2$, with the result that optical lasers can generate a number of e^+e^- pairs large enough to allow their experimental observation.

For the theoretical analysis we use the kinetic equation approach, which allows to consider the dynamics of the creation process while taking into account the initial conditions [4]. This approach has been applied already to the periodical field case [5] with near-critical values of the field strength and X-ray frequencies. We consider here a simple model of the laser field, which can be formed in the focus of two counter-propagating laser beams: a harmonic linearly polarized field which acts during N periods. We solve the quantum kinetic equation numerically with the initial condition of a vanishing field and obtain the particle number density as a moment of the distribution function.

For a weak field $E \ll E_{cr}$, the e$^+$e$^-$ plasma density oscillates with twice the field frequency. We obtain that the electric field $E \approx 7 \times 10^{-6} E_{\rm cr}$ produces $\approx 3 \cdot 10^5$ pairs in a volume of λ^3 averaged over one period of the laser field what corresponds to a particle density of 10^{18} cm^{-3}. This dense enough electron-positron plasma exists during a laser pulse but vanishes almost completely after switching off the field. The residual density n_r corresponding to an integer number of field periods, is negligible in comparison with the mean density $\langle n\rangle$ per period. Under the above conditions the ratio of $\langle n\rangle/n_r$ is approximately $3\cdot 10^{11}$. As a consequence, in spite of the fact that the residual density for the X-ray laser exceeds the one for the optical laser by a large factor, the situation is different regarding the mean density: the optical laser can produce more pairs in the volume of λ^3 than the X-ray one [6].

In order to prove the suggested effect experimentally, we discuss here two schemes. The first one is the two-photon annihilation of e$^+$e$^-$ within the plasma volume. The corresponding γ quanta with the total energy ≈ 1 MeV can be registered in coincidence outside the focus of the counter-propagating laser beams. We have estimated the number of annihilation events per laser pulse $d\nu$ with the following parameters: pulse intensity $I = 10^{18}$ W/cm^2, pulse duration $\tau_L \sim 80$ fs, wavelength $\lambda = 795$ nm, cross size of laser beams $z_0 = 9$ μm [2]. The estimate results in ≈ 1 annihilation event per laser pulse, see Fig. 1. Another possible method to detect the formation of the quasiparticle plasma consists in the observation of the nonlinear plasma response on a weak monochromatic probe signal with intensity $I_{\rm pr}$. For example, the intensity of the third harmonic is of the order $I_3 \sim 10^{-4} I_{\rm pr}$ [7].

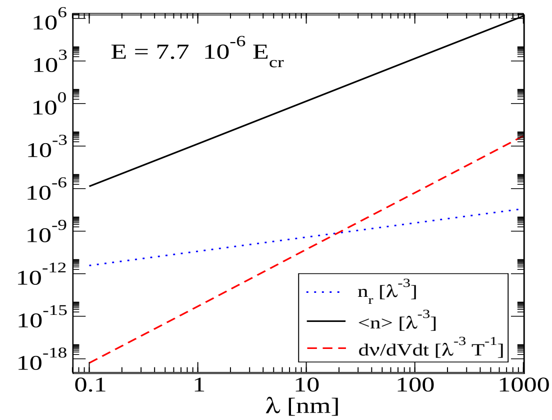

Figure 1: The laser wavelength dependence of the mean ($\langle n\rangle$) and residual (n_r) e$^+$e$^-$ pair densities as well as the γ pair production rate per volume for the field strength $E = 7.7 \cdot 10^{-6} E_{\rm cr}$.

In conclusion, we have shown that the simplest model of the laser field predicts the creation of a dense quasiparticle plasma in the foci of the counter-propagating optical laser beams with parameters corresponding to presently operating systems [2, 1]. The plasma lives during a laser pulse and vanishes almost completely after switching off the field. The mean density does not depend on the frequency and reaches the values $10^{17} - 10^{18}$ cm^{-3} for field strengths $10^{10} - 10^{11}$ V/cm. One possible manifestation of this picture can be the pair of ≈ 1 MeV γ quanta emitted per laser pulse, another is the modulation of a probe laser beam with the time-dependent plasma density.

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Electron Cooling of Highly Charged Ions in Traps

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Electron cooling in HITRAP [1] is an essential tool for providing cold ions required for planned precision tests of QED and experiments with low energetic highly charged ions or antiprotons. This method is well established in storage rings and is based on the energy loss of the ions due to Coulomb collisions with cold electrons. In Penning traps, like in HITRAP, the presence of a strong magnetic field presents a theoretical challenge. We thus developed several complementary analytical and numerical methods to calculate the cooling force [2], comprising linearized dielectric theory, a binary collision model [3], classical trajectory Monte Carlo (CTMC) calculations and particle-in-cell (PIC) simulations. Based on these results for the cooling force $\vec{F}(\vec{v}_i, T_e)$ or energy loss rate $dE_i/dt(\vec{v}_i, T_e) = \vec{F} \cdot \vec{v}_i$ we estimate cooling times from calculating the time evolution of the energy of the ions. The actual energy loss now depends on the actual ion velocity $\vec{v}_i$ and the electron temperature T_e. Since ions and electrons are trapped together, the temporal changes of T_e are important here. Hence three effects are taken into account for these calculations. The deceleration of the ion and its related energy loss, the heating of the electrons due to the energy transfer from the ions, and the cooling of the electrons by emission of synchrotron radiation [4]. This leads to the coupled differential equations for $\vec{v}_i(t)$ and $T_e(t)$ which are solved numerically:

$$\frac{d\vec{v}_i}{dt} = \frac{1}{M}\vec{F}(\vec{v}_i, T_e)$$

$$\frac{dT_e}{dt} = -\frac{2}{3k_B}\frac{n_i}{n_e}\frac{dE_i}{dt}(\vec{v}_i, T_e) - \frac{1}{\tau_e}(T_e - T_{e,0}) \quad .$$

M is the ion mass, n_i the ion density, n_e the electron density, $T_{e,0}$ the temperature to which the trap is cooled and τ_e the time constant for cooling of the electrons by the emission of synchrotron radiation.
Fig. 1 shows the ion energy $E_i = \frac{1}{2}Mv_i^2$ for U^{92+} ions as a function of time with $n_e = 10^7$ cm^{-3}, $B = 6$ T at a fixed electron temperature $T_e = T_{e,0} = 4$ K (that is for $n_i/n_e \to 0$) and different $\alpha := \sphericalangle(\vec{v}_i, \vec{B})$. The large variation of the cooling time with α results from the strong anisotropy of the cooling force $\vec{F}$ in the presence of the strong magnetic field. This is compared with the case $B = 0$, which yields the fastest cooling. If the heating of the electrons is taken into account (see Fig. 2 with $n_i/n_e = 10^{-4}$) this anisotropy is averaged out as a result of an intricate feedback between the electron temperature T_e (shown in the lower part of Fig. 2) and the force $\vec{F}(\vec{v}_i, T_e)$ on the ion. A similar qualitative behavior is also found for $n_i/n_e = 10^{-5}$ and 10^{-3} with cooling times 0.5 s and 4 s, respectively. Such a dependence on n_i/n_e and typical cooling times of the order of 1 s agree well with the previous findings of Ref.[4] where the influence of the magnetic field on dE_i/dt had been neglected. However, more quantitative comparison of the results of [4] with our calculations shows lower electron temperatures and cooling times which are longer by a factor $2-4$. - This work was supported by the BMBF (06ER128) and a GSI collaboration contract.

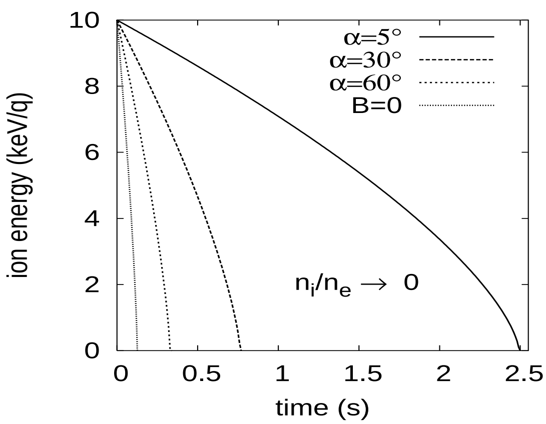

Fig.1: Temporal evolution of the energy of U^{92+} ions at a fixed electron temperature $T_e = 4$ K.

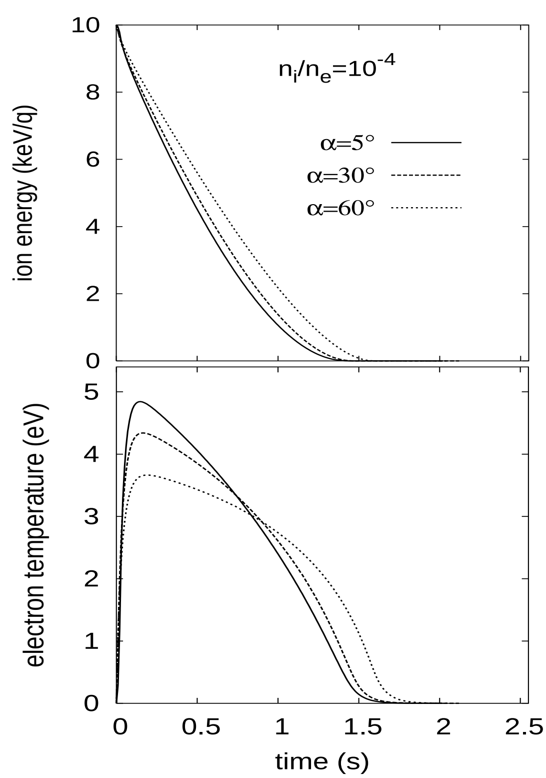

Fig.2: Time evolution of the ion energy and the electron temperature if electron heating is taken into account.

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