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Relativistic electron-ion recombination in the presence of an intense laser field

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Abstract
Radiative recombination of a relativistic electron with a highly charged ion in the presence of an intense laser field is considered. The high electron energy and its strong coupling with the laser field lead to a very broad energy spectrum of the emitted γ-photons featuring very pronounced side wings and showing a saturation at extremely relativistic electron energies. Moreover, characteristic shifts in the angular distribution occur due to the impact of the laser photon momentum.

Radiative recombination (RR) of a free electron with an ion is a fundamental process which plays an important role in all kinds of astrophysical and laboratory plasmas. It represents the inverse of the photoeffect and has been studied for a long time [1].

More recently, the availability of intense laser sources has triggered a growing interest in RR in the presence of an external laser field which can lead to strong modifications of the field-free properties of the process. When, for example, the laser photon frequency is resonant with the electron transition energy, the field stimulates the RR process and substantially enhances its rate. Corresponding experimental investigations rely on merged beams of ions, electrons and photons in a storage ring [2]. Very recently, multiphoton-assisted RR of low-energy electrons into barium ions in a microwave cavity has been observed in a weak-coupling regime [3].

Modern powerful laser devices are capable of producing field intensities in excess of 10^{20} W cm^{-2} in the near-optical frequency domain. Electrons exposed to such strong fields may be accelerated to highly relativistic energies, e.g. through laser–plasma interactions [4] or strong-field photoionization [5]. Modifications of many atomic collision processes such as Mott and Møller scattering, or Bethe–Heitler pair creation due to the presence of a relativistically strong laser field have been theoretically studied during the last decade [6–8].

The theoretical investigation of laser-assisted radiative recombination (LARR), however, has so far been restricted to the nonrelativistic regime where the dipole-approximation to the field applies (see [9–11, 12, 13] and references therein). The main emphasis was placed upon the x-ray energy spectrum. Effects from the intensity profile of a focused laser pulse have also been addressed [11] and the possibility of phase control has been shown [12]. The laser field may also affect the x-ray polarization [13].

Note also that LARR represents the last step in the coherent process of high-harmonic generation from gas targets which has intensively been studied both theoretically and experimentally for 20 years. In this process, RR of field-ionized atomic electrons occurs through laser-driven recollisions, with the highest harmonic photon energies achieved of about 1 keV [14]. The relativistic domain of high-harmonic generation, however, is inaccessible until now since at optical laser intensities exceeding $\sim 10^{17}$ W cm^{-2}, the magnetic-field-induced Lorentz drift motion of the electron prevents efficient recollisions [7, 15]. This circumstance provides additional motivation for studies of relativistic LARR utilizing free electron and ion beams.

In this communication, we extend the consideration of LARR into the relativistic domain where various characteristic modifications of the process are demonstrated to occur. Relativistic effects arise from the high initial electron energy (in the MeV range) and the large laser field strength-to-frequency ratio leading to strong electron-field coupling. In particular, a very large number of low-frequency laser photons participate in the process whose total energy and momentum are imprinted on the emitted high-frequency (γ-ray) photon influencing its energy and angular distributions in a distinctive non-dipole manner.

In the following, atomic units are used unless otherwise indicated.

The laser field is assumed to be a classical monochromatic plane wave of circular polarization switched on and off adiabatically at $t \to \mp\infty$, respectively. We shall describe this field by the four-potential $(A_0, \mathbf{A})$:

$$A_0 = F_0(\mathbf{e}_2 \cdot \mathbf{r} \cos\varphi - \mathbf{e}_1 \cdot \mathbf{r} \sin\varphi), \qquad \mathbf{A} = A_0 c \mathbf{k}_0/\omega_0. \tag{1}$$

Here $\mathbf{r}$ and t are the space-time coordinates, ω_0 and $\mathbf{k}_0$ are the frequency and wave vector, $\varphi = \omega_0 t - \mathbf{k}_0 \cdot \mathbf{r}$ is the phase, $\mathbf{e}_{1,2}$ are the polarization vectors ($\mathbf{e}_i \cdot \mathbf{e}_j = \delta_{ij}$, $\mathbf{e}_i \cdot \mathbf{k}_0 = 0$) and F_0 and c are the strength of the laser field and the speed of light, respectively. Note that in the dipole approximation ($\mathbf{k}_0 \cdot \mathbf{r} \to 0$) the choice (1) for the field potentials corresponds to the so-called length gauge.

The transition amplitude for the radiative recombination reads

$$S_{fi} = -\mathrm{i} \int_{-\infty}^{+\infty} \mathrm{d}t \langle \Psi_f(t) | \hat{W} | \Psi_i(t) \rangle, \tag{2}$$

where $|\Psi_i(t)\rangle$ and $|\Psi_f(t)\rangle$ are the initial and final states of the system 'electron+radiation field' and $\hat{W}$ is the interaction with the radiation field. The latter is chosen in the form $\hat{W} = \boldsymbol{\alpha} \cdot \hat{\mathbf{A}}_\gamma$ where $\boldsymbol{\alpha} = (\alpha_x, \alpha_y, \alpha_z)$ are the Dirac matrices and

$$\hat{\mathbf{A}}_\gamma(\mathbf{r}, t) = \sum_{\mathbf{k},\rho} \sqrt{\frac{2\pi c^2}{V\omega_k}} \mathbf{e}_{\mathbf{k},\rho} \left(c^+_{\mathbf{k},\rho} \, \mathrm{e}^{\mathrm{i}(\omega_k t - \mathbf{k}\cdot\mathbf{r})} + \mathrm{C.C.} \right). \tag{3}$$

In (3), $c^+_{\mathbf{k}\rho}$ is the creation operator for a photon with momentum $\mathbf{k}$, polarization vectors $\mathbf{e}_{\mathbf{k}\rho}$ ($\rho = 1, 2$) and frequency $\omega_k = c|\mathbf{k}|$, V is the normalization volume for the radiation field and the sum runs over all photon modes.

The states $\Psi_i(t)$ and $\Psi_f(t)$ are given by

$$|\Psi_i(t)\rangle = \psi_i(t)\, |0\rangle, \qquad |\Psi_f(t)\rangle = \psi_f(t)\, |\mathbf{k}\rho\rangle, \tag{4}$$

where $|0\rangle$ and $|\mathbf{k}\rho\rangle$ denote states of the radiation field with no photons and with one photon having momentum $\mathbf{k}$ and polarization ρ. The initial and final states of the electron, ψ_i and ψ_f, are solutions of the Dirac equation

$$\mathrm{i}\frac{\partial\psi}{\partial t} = c\boldsymbol{\alpha} \cdot \left(\hat{\mathbf{p}} + \frac{1}{c}\mathbf{A} \right) \psi + V_0\psi - A_0\psi + \beta c^2 \psi, \tag{5}$$

where $V_0 = -Z/r$ is the interaction between the electron and the ionic nucleus and β is the Dirac matrix. Both in the initial and final states, the electron is subject to the simultaneous presence of two fields: the field of the nucleus and the laser field. Since an exact solution of such a problem is not known, suitable approximations to describe these states are needed.

The incident electron is supposed to have relativistic asymptotic momentum p_i and energy E_i. Assuming that the charge z of the nucleus satisfies the condition $Z \ll c$, one can neglect the Coulomb interaction in the initial state. At the same time, the low-frequency laser field can very substantially affect the motion of the incident electron and its interaction with this field should be taken into account to all orders. Therefore, we shall approximate the initial state of the electron using the Gordon-Volkov solution of [16] multiplied by the factor $\exp(\mathrm{i}F_0(\mathbf{e}_1 \cdot \mathbf{r}\cos\varphi + \mathbf{e}_2 \cdot \mathbf{r} \sin\varphi)/\omega_0)$ to adapt it to the gauge (1).

Let us now assume that although the laser field is quite strong, its strength F_0 nevertheless remains much smaller than the typical nuclear field $F_a \sim Z^3$ acting on the electron bound in the ground state of the ion[1]. Then, taking into account that in the gauge (1) the ratio between the terms in equation (5), describing the interactions of the bound electron with the laser and nuclear fields, is $\sim F_0/F_a \ll 1$, the final state ψ_f can simply be taken as

$$\psi_f = \phi_0(\mathbf{r}) \exp(-\mathrm{i}\varepsilon_0 t), \tag{6}$$

where ϕ_0 is the wave function of the electron in the ground state and ε_0 is the corresponding electron energy.

Inserting the approximate electron states into (2) and assuming that the condition $v_{i\perp}/\omega_0 \gg r_0$ is fulfilled (where $v_{i\perp}$ is the part of the initial electron velocity v_i perpendicular to $\mathbf{k}_0$ and $r_0 \sim 1/Z$ is the size of the final bound state of the electron) we obtain

$$S_{fi} = -\mathrm{i}2\pi \sqrt{\frac{2\pi c^2}{V\omega_k}} \sqrt{\frac{mc^2}{E_i}} \times \sum_{n=-\infty}^{+\infty} G_n \, \delta(n\omega_0 + \varepsilon_0 + \omega_k - \xi\omega_0 - E_i). \tag{7}$$

In the above expression

$$G_n = \int \mathrm{d}^3\mathbf{r} \exp(\mathrm{i}(\mathbf{p}_i - \mathbf{k} - n\mathbf{k}_0 + \xi\mathbf{k}_0)\cdot\mathbf{r}) \times \mathbf{e}_{\mathbf{k}\rho} \cdot [\mathbf{f}_+ J_{n-1}(\Lambda) + \mathbf{f}\, J_n(\Lambda) + \mathbf{f}_- J_{n+1}(\Lambda)], \tag{8}$$

where $\mathbf{f}_\pm = \phi_0^\dagger \boldsymbol{\alpha}\, \eta_\pm\, u(\mathbf{p}_i, s) \exp[\pm\mathrm{i}(\chi_{\mathbf{p}} + \mathbf{b}\cdot\mathbf{r})]$ and $\mathbf{f} = \phi_0^\dagger \boldsymbol{\alpha}\, u(\mathbf{p}_i, s)$, with $u(\mathbf{p}, s)$ being the free Dirac spinor for an electron with momentum $\mathbf{p}$ and spin s,

$$\eta_\pm = \frac{F_0}{4\omega_0} \frac{\omega_0/c + \boldsymbol{\alpha}\cdot\mathbf{k}_0}{\omega_0 E_i/c^2 - \mathbf{k}_0\cdot\mathbf{p}_i} \boldsymbol{\alpha}\cdot(\mathbf{e}_1 \mp \mathrm{i}\mathbf{e}_2), \tag{9}$$

J_n is the Bessel function with argument

$$\Lambda = \lambda - \mathbf{d}\cdot\mathbf{r}, \qquad \mathbf{d} = \frac{F_0^2}{\lambda\,\omega_0^2} \frac{(\mathbf{e}_1\cdot\mathbf{p}_i)\,\mathbf{e}_2 - (\mathbf{e}_2\cdot\mathbf{p}_i)\,\mathbf{e}_1}{\omega_0 E_i/c^2 - \mathbf{k}_0\cdot\mathbf{p}_i},$$

$$\lambda = \frac{F_0/\omega_0}{\omega_0 E_i/c^2 - \mathbf{k}_0\cdot\mathbf{p}_i} \sqrt{(\mathbf{e}_1\cdot\mathbf{p}_i)^2 + (\mathbf{e}_2\cdot\mathbf{p}_i)^2} \tag{10}$$

and

$$\xi = \frac{c^2 F_0^2/\omega_0^2}{2(\omega_0 E_i - c^2\mathbf{k}_0\cdot\mathbf{p}_i)}, \qquad \chi_{\mathbf{p}} = \arctan\left(\frac{\mathbf{e}_2\cdot\mathbf{p}_i}{\mathbf{e}_1\cdot\mathbf{p}_i}\right),$$

$$\mathbf{b} = \frac{F_0^2}{\lambda^2\,\omega_0^2} \frac{(\mathbf{e}_1\cdot\mathbf{p}_i)\,\mathbf{e}_1 + (\mathbf{e}_2\cdot\mathbf{p}_i)\,\mathbf{e}_2}{\omega_0 E_i/c^2 - \mathbf{k}_0\cdot\mathbf{p}_i}. \tag{11}$$

Note that the matrix element in equation (8) can be evaluated by analytical means (see [17] for a similar calculation).

One should mention that had we used the 'conventional' radiation gauge for the laser field potentials ($A_0 = 0, \mathbf{A} = cF_0/\omega_0(\mathbf{e}_1\cos\varphi + \mathbf{e}_2\sin\varphi)$), the expression (7) for the transition amplitude would still be the same. Indeed, in the radiation gauge, the Gordon–Volkov solution would not possess the factor $\exp(\mathrm{i}F_0(\mathbf{e}_1\cdot\mathbf{r}\cos\varphi + \mathbf{e}_2\cdot\mathbf{r}\sin\varphi)/\omega_0)$ but then an additional factor $\exp(-\mathrm{i}F_0(\mathbf{e}_1\cdot\mathbf{r}\cos\varphi +$

[1] Provided $Z \gg 1$, the condition $F_0 \ll F_a$ is not very restrictive. For $Z = 50$, e.g. one has $F_a \sim 10^5$ au so that even fields produced by the most powerful modern lasers ($F_0 \lesssim 3\times10^2$ au) still fit very well into this condition.

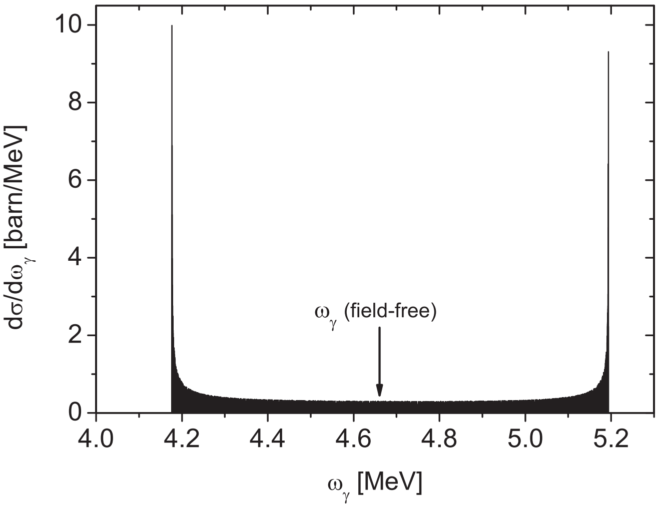

Figure 1. Energy spectrum of the γ-ray photons emitted in LARR of a relativistic electron ($p_i = 10mc$) with a bare Sn nucleus ($Z = 50$) in a circularly polarized laser field with $F_0 = 7.5$ au and $\omega_0 = 0.055$ au crossing the beam under 90°.

$\mathbf{e}_2 \cdot \mathbf{r} \sin\varphi)/\omega_0)$ would appear in the expression (6) for the final electron state (see e.g. [18]).

In what follows we consider, based on equations (7)–(11), the reaction $e^- + \mathrm{Sn}^{50+} \to \mathrm{Sn}^{49+}(1s) + \gamma$ occurring in a laser field with $F_0 = 7.5$ au ($I \approx 4 \times 10^{18}$ W cm^{-2}), $\omega_0 = 0.055$ au (1.5 eV), $\mathbf{e}_1 = (1, 0, 0)$ and $\mathbf{e}_2 = (0, 1, 0)$ in which the initial electron momentum $p_i = 10$ mc is directed along the x-axis.

Figure 1 shows the energy spectrum of the γ-ray photons. The spectrum consists of a multitude of lines each separated by a laser photon energy, thus forming a quasi-continuous distribution with a total width of $\Delta\omega_\gamma \sim 1$ MeV and very pronounced side wings. Without the laser field, the spectrum would comprise a single line at $\omega_\gamma = E_i - \varepsilon_0 \approx 4.66$ MeV.

The motion of an electron in the presence of a strong laser field has quasi-classical character. Therefore, similarly as in the nonrelativistic case (see e.g. [10]), certain features of the relativistic LARR energy spectrum can be understood in classical terms: in the presence of the field, the instantaneous electron energy is modulated,

$$E(\varphi) = E_i + U_p + \lambda\omega_0 \cos\varphi$$
$$= E_i + \frac{c^2 F_0^2}{2\omega_0^2 E_i} + \frac{c p_i}{E_i}\frac{c F_0}{\omega_0}\cos\varphi, \qquad (12)$$

so that the available kinetic energy depends on the instant phase of recombination. While the ponderomotive energy $U_p = \xi\omega_0 \approx 25$ keV leads to a relatively small shift of the centre of the spectrum, the oscillating term in equation (12) causes a spectral width of $\Delta\omega_\gamma = 2\lambda\omega_0 \sim 1$ MeV between the minimum and maximum energies corresponding to recombination occurring at $\cos\varphi = -1$ and $\cos\varphi = 1$, respectively.

In the quantum picture, the broad spectrum results from the emission or absorption of n laser photons during the RR process, with the boundaries $|n| \lesssim n_{\max} \simeq \lambda$ determined by the properties of the Bessel functions. While already the classical consideration gives the range of the energy spectrum, of course only a quantum consideration can predict the shape of the spectrum.

It is known that in the nonrelativistic domain, the spectral width grows with the kinetic energy of the incoming electron [9, 10]. In the present case, however, we find that the width becomes practically independent of the incoming electron energy, i.e. it saturates. This occurs because $\lambda \propto v_i$ (see equation (10)) and the electron velocity v_i cannot exceed the speed of light. This implies in fact that the width and shape of the spectrum in figure 1 (where $v_i \approx c$) are 'universal' in the sense that they will remain practically unchanged when E_i is further increased. Enhancement of the electron energy would shift the centre of the spectrum to correspondingly higher energies but, even in the limit $E_i \to \infty$, the number of laser photons participating in the process would not increase noticeably.

In the quasi-classical case ($\lambda \gg 1$), the energy spectrum of the γ-ray photons reaches large maxima at the side wings. This can be qualitatively understood on the basis of expression (12): the incident electron spends more time in states, where the energy (12) is close to its maximum or minimum value than in states with intermediate energies ($\cos\phi$ varies more slowly when it is close to its extremal points). With the further increase of λ, this feature becomes even more pronounced because the electron spends relatively more and more time in states with the extreme energies. In fact, the ratio of the spectral intensity at the side wings as compared with the plateau region in between scales with $\lambda^{1/3}$ and amounts to $\approx$30 in figure 1 which is significantly larger than in the nonrelativistic regime.

Characteristic relativistic signatures also arise in the angular emission spectrum, shown in figure 2. In the nonrelativistic domain, the photons produced by RR into the ground state are emitted mainly in the direction perpendicular to the momentum of the incident electron, forming a typical dipole pattern. In contrast to that, the γ-rays are emitted essentially along the direction of the incoming electron. This feature is also known from relativistic field-free RR [19]. However, as compared with the field-free case, the LARR spectrum in figure 2 is shifted to larger angles and broadened.

The shift is most pronounced in the energetic side wings where it is almost solely caused by the momentum carried by the laser photons (and thus would be absent in the nonrelativistic domain) and where it may again be understood in classical terms by inspecting the instantaneous electron momentum $\mathbf{p}(\varphi)$ with the components

$$p_x(\varphi) = p_i + \frac{F_0}{\omega_0}\cos\varphi, \qquad p_y(\varphi) = \frac{F_0}{\omega_0}\sin\varphi,$$
$$p_z(\varphi) = \frac{U_p}{c} + \frac{p_i c}{E_i}\frac{F_0}{\omega_0}\cos\varphi. \qquad (13)$$

The x and y components lie in the polarization plane of the field and show a corresponding modulation due to the coupling with the electric component of the field, whereas the electron momentum along the z-axis arises solely from the light pressure exerted by the field momentum. By equation (13), the typical polar angle θ_γ under which the electron impinges on the ion before recombination is given by $\tan\theta_\gamma = p_\perp/p_x$

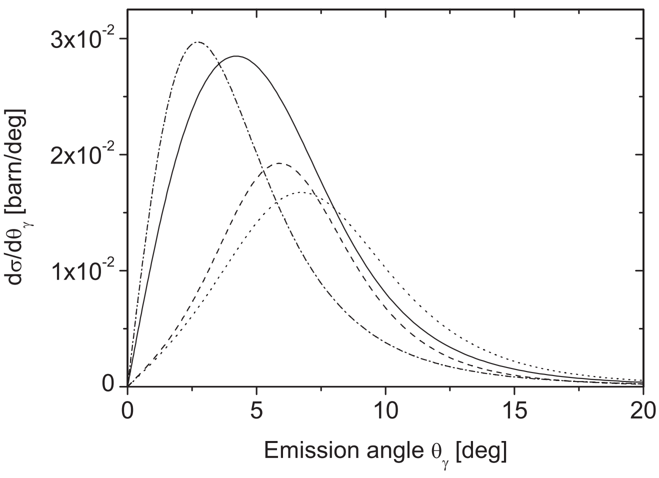

Figure 2. Angular distributions of the emitted γ-ray photons with respect to the incoming electron momentum. The dash-dotted curve shows the field-free case, whereas the solid line refers to LARR with the laser parameters of figure 1. The dotted and dashed lines show the spectra for the energy side wings around $\omega_\gamma \approx 4.2$ MeV and $\omega_\gamma \approx 5.2$ MeV (see figure 1), respectively, within an energy interval of $\delta\omega_\gamma = 1$ keV; the latter two curves are multiplied by a factor of 30.

with $p_\perp = \sqrt{p_y^2 + p_z^2}$ and depends on the phase φ. For the minimum instantaneous electron energies ($\varphi = \pi$), we have $p_\perp = p_z \approx F_0/\omega_0 \approx mc$, $p_x \approx 9mc$ and $\theta_\gamma \approx 6.3°$, whereas for $\varphi = 0$ where the maximum electron energy occurs, we obtain $p_x \approx 11mc$, $p_\perp \approx mc$ and $\theta_\gamma \approx 5.2°$. Both θ values are in agreement with the side-wing distributions shown in figure 2.

A similar, but somewhat more complicated, explanation can also be found for the position of the maxima of medium-energy γ-rays. Let us consider, for example, the centre of the energy plateau in figure 1 where $\cos\varphi = 0$ and thus a net number of $n = 0$ laser photons participate in the process. The corresponding angular spectrum (not shown) peaks at small angles around $\theta \approx 2.5°$ which coincides with the field-free case (see figure 2). This can be understood by noting that now, within a single laser cycle, there are two possible phases $\varphi_+ = \pi/2$ and $\varphi_- = 3\pi/2$, both with $\cos\varphi_\pm = 0$ but opposite directions of $p_y(\varphi_\pm) \approx \pm mc$, whose contributions to the LARR amplitude interfere. The peak for $n = 0$ may thus be considered as arising from the 'average' electron momentum with $\frac{1}{2}[p_y(\varphi_+) + p_y(\varphi_-)] = 0$ and $p_x = p_i$.

For increasing photon numbers n, with $0 \ll |n| \ll n_{\max}$, the interference effect still exists but is becoming less pronounced. When the side wings are eventually reached ($n \approx \pm n_{\max}$), both quantum paths merge into a single one and the classical picture becomes applicable. The peak of the total LARR angular spectrum, which represents a sum over the contributions from all laser photon numbers, hence lies in between the field-free peak and the side-wing peaks.

During the electron–ion interaction, high-energy photons may also be produced by bremsstrahlung. Similar as for LARR, the endpoint of the bremsstrahlung spectrum is determined by the maximum difference $E_i(\varphi = 0) - E_f$ of the classical instantaneous electron energies in the field [20],

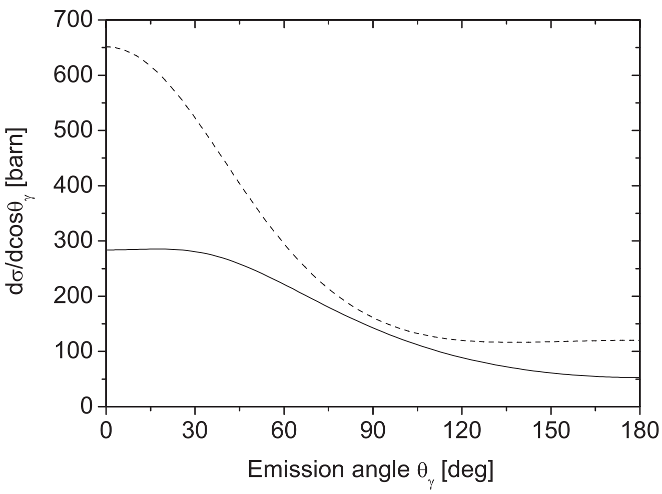

Figure 3. Angular photon spectra $d\sigma/d\cos\theta_\gamma$ for LARR of a weakly relativistic electron ($p = 0.3$ mc) with a bare Zn nucleus ($Z = 30$) assisted by a circularly polarized laser field with $F_0 = 2.3$ au and $\omega_0 = 0.055$ au. The solid and dash curves show the results of the relativistic and nonrelativistic theories, respectively.

where $E_f = mc^2 + F_0^2/(2\omega_0^2 m)$ is the minimum energy of the final-state continuum electron. For the parameters assumed in figure 1, the bremsstrahlung cutoff thus lies around 4.9 MeV which is considerably lower than the LARR endpoint; the difference is $F_0^2/(2\omega_0^2 m) + mc^2 - \varepsilon_0 \approx 0.3$ MeV. Consequently, the high-energy part of the LARR spectrum is not overlaid by bremsstrahlung photons.

In our case, the total LARR cross section of about 0.2 b is comparable with the relevant high-energy tail of bremsstrahlung. The comparable absolute signal strengths are to some extend fortuitous since LARR scales with Z^5 while bremsstrahlung as a free-free transition raises like Z^2. In particular, in a laser-generated plasma involving very highly charged ions LARR is expected to dominate the production of energetic photons.

A nonrelativistic description of the LARR process clearly does not apply to the electron and field parameters of figures 1 and 2; in particular, the nonrelativistic total cross section comes out by four orders of magnitude too large. Therefore, in order to compare the relativistic and nonrelativistic theories, we consider in figure 3 a parameter set at the borderline between the respective regimes. It shows that already at modest energies and field strengths, the proper inclusion of relativistic effects leads to strong modifications of the angular spectrum.

Before we proceed to conclusions, two more remarks seem to be appropriate. First, the beam geometry we have considered, in which the asymptotic electron momentum $\mathbf{p}_i$ lies in the polarization plane $(\mathbf{e}_1, \mathbf{e}_2)$ of the laser field, maximizes the value of the parameter λ (see equation (10)) and thus the coupling between the incident electron and the field. When the angle δ between $\mathbf{p}_i$ and the plane $(\mathbf{e}_1, \mathbf{e}_2)$ increases, the coupling becomes weaker leading to smaller numbers of photons exchanged between the laser field and the recombining electron-ion system. In particular, both the quantum considerations based on equations (7)–(11) and the classical expression (12) predict that the width of the energy distribution of the emitted γ photon reduces when λ decreases. Moreover, based on expression (12), one would even expect that the width of the spectrum becomes zero when $\delta = 90^\circ$. While equations (7)–(11) do not enable us to access this particular geometry[2], a more general consideration shows, however, that even in such a case, the width does not reduce to zero but is of the order of $F_0 r_0 \sim F_0/Z$.

Second, in our consideration, we have neglected the influence of the Coulomb field on the incident electron. Based on the very numerous previous studies of various atomic processes in the presence of a strong laser field, one may expect that the account of the Coulomb effects would not lead to qualitative changes in the calculated photon spectra. However, it is very difficult to predict how this account would change the spectra quantitatively. A rough estimate of the Coulomb effects can be obtained by considering the so-called Gamov factor, which in the case of an ultrarelativistic electron is given by $2\pi Z/c(1 - \exp(-2\pi Z/c))$. Taking $Z = 50$, we obtain that the calculated cross sections would increase by a factor of 2.5 without changing their shape.

In conclusion, we have considered radiative recombination of a relativistic electron with a highly charged ion assisted by an intense laser field of circular polarization. We have shown that the field substantially modifies the shape of the spectra of the emitted γ photons. We have discussed relativistic effects caused by the high energy of the incoming electron and its strong coupling with the laser field resulting in very large energy-momentum exchanges between the recombining electron-ion system and the field.

An observation of the predicted effects in principle is feasible in many high-power laser laboratories worldwide because such lasers can be used to generate the relativistic electron beams and highly charged ions required [4, 7].

Acknowledgments

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[2] The condition $v_{i\perp}/\omega_0 \gg 1/Z$ (where $v_{i\perp}$ is the component of the initial electron velocity perpendicular to $\mathbf{k}_0$), used to derive equations (7)–(11), implies that these expressions cannot be applied when $\lambda \to 0$.