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PRODUCTION OF A POSITRON AND A BOUND ELECTRON BY HIGH-ENERGY PHOTON IN A STRONG COULOMB FIELD

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НОВОСИБИРСК

Production of a positron and a bound electron by high-energy photon in a strong Coulomb field.

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Abstract

The cross section of e^+e^- pair creation by a high-energy photon with capture of the produced e^- in arbitrary bound state of arising hydrogen-like atom is found using the quasiclassical approach with exact regard for a Coulomb field. Formulae are essentially simplified for large quantum numbers of the bound state. That permits us to find the total cross section of the process.

1 Introduction

During recent years QED processes in a strong Coulomb field attracted significant attention, in particularly, connected with construction of relativistic heavy-ion colliders. Besides other important problems, these colliders will allow to check QED predictions for the processes in extremely strong external fields. The process of bound-electron-positron pairs production is of peculiar interest, because it limits essentially the beam lifetime.

The consideration of the phenomena may be conveniently performed in the rest frame of an ion (target ion) capturing an electron. For very high ions energies, the equivalent-photon method is valid. Within this method the cross section of the process at ion collision can be expressed via the cross section of the pair production by the photon in target ion field. Then the cross section of the process under consideration is proportional to $Z_p^2 \ln\gamma$, where Z_p is the projectile ion charge (in units of the electron charge) and γ is its Lorentz-factor. Recently [1, 2] it was pointed out that for moderate ions energies the nonperturbative mechanism turns out to be important. This mechanism gives essentially different Z dependence of the process cross section as compared to that obtained within the equivalent photon method, but does not contain the factor $\ln\gamma$ that is large in ultrarelativistic case.

To apply the equivalent photon method we must know the corresponding photo-process cross section. The total cross section of bound-electron-positron pair production is determined by equivalent photon energies of the order of a few electron masses (we use in this paper the system of units $\hbar = 1$, $c = 1$). This cross section was calculated for K and L shells in a series of papers (see reviews [3, 4] and references cited there). The cross section differential with respect to energies of outgoing positrons is also of obvious interest. We shall consider the case of ultrarelativistic positrons. Hence, we must find the cross section of bound-electron-positron pairs production by a high energy photon in a Coulomb field ($\omega \gg m$, m is the electron mass, ω is the photon energy). This cross section was investigated in details for K and L shells in [5, 6], where earlier published papers are cited as well. As it was noted in [5], the bounded pair production cross section coincides for $\omega \gg m$ with those of the photoeffect and radiative recombination (up to the factor connected with summing or averaging over polarizations of involved particles). The results in [5, 6] were obtained by direct calculation of matrix elements using Sommerfeld-Maue approximation [7] for wave functions.

The present paper is devoted to the calculation of the cross section of e^+e^- pair creation by a high-energy photon when a produced electron is captured into arbitrary state of discrete spectrum with exact regard for a Coulomb field. The consideration is based on the use of the integral representation for the electron Green function in a Coulomb field obtained by us in [8]. The quasiclassical limit of this function [9, 10] is used as well. The case of large electron-state quantum numbers is analyzed in details. The asymptotic expressions obtained are of a high accuracy that permits us to find the total (summed up over all discrete electron states) cross section of the process.

2 Green functions and general expression for the cross section of the process

According to the common Feynman-rules, the cross section of bound-electron-positron pair production by a photon in a Coulomb field of a nucleus is

$$\sigma_\gamma = \frac{\alpha}{2\pi\omega}\int d\vec{p}\,\delta(\omega - \epsilon_p - \epsilon_n)|M|^2, \qquad (1)$$

where $\alpha = e^2 = 1/137$ is the fine structure constant, summing over the electron and positron polarization states is assumed. The matrix element M

is

$$M = \int d\vec{r}\,\bar{\psi}_n^{(+)}(\vec{r})\,\hat{e}\,\psi_p^{(-)}(\vec{r})\exp(i\vec{q}\vec{r}).$$

Here $\psi_n^{(+)}(\vec{r})$ is the bound state wave function; $\psi_p^{(-)}(\vec{r})$ is negative frequency wave function for the continuum state corresponding to positron, e_μ is the photon polarization 4-vector, $\vec{q}$ is the photon wave-vector. According to the usual definition (see, e.g., [11], (109.19)), the electron Green function in an external field can be represented in the form

$$G^{\pm}(\vec{r},\vec{r}'|\,\epsilon) = \sum_n \frac{\psi_n^{(+)}(\vec{r})\bar{\psi}_n^{(+)}(\vec{r}')}{\epsilon - \epsilon_n \pm i0} +$$

$$\int\frac{d\vec{p}}{(2\pi)^3}\left[\frac{\psi_p^{(+)}(\vec{r})\bar{\psi}_p^{(+)}(\vec{r}')}{\epsilon - \epsilon_p \pm i0} + \frac{\psi_p^{(-)}(\vec{r})\bar{\psi}_p^{(-)}(\vec{r}')}{\epsilon + \epsilon_p \mp i0}\right], \qquad (2)$$

where G^+ is the Green function in the upper and G^- in the lower half-plane of the complex variable ϵ . We find from (2)

$$\delta G\,(\vec{r},\vec{r}'\,|-|\epsilon|) = G^+ - G^- = 2i\pi\int\frac{d\vec{p}}{(2\pi)^3}\,\psi_p^{(-)}(\vec{r})\,\bar{\psi}_p^{(-)}(\vec{r}')\delta(|\epsilon| - \epsilon_p),$$

$$\psi_n^{(+)}(\vec{r})\,\bar{\psi}_n^{(+)}(\vec{r}') \equiv \rho_n(\vec{r},\vec{r}') = \lim_{\epsilon\to\epsilon_n}(\epsilon - \epsilon_n)G(\vec{r},\vec{r}'|\,\epsilon). \qquad (3)$$

Substituting (3) into (1) we get a formula for the cross section expressed via Green function only:

$$\sigma_\gamma = \frac{i\pi\alpha}{\omega}\int\!\!\int d\vec{r}\,d\vec{r}'\,Sp\,\{\rho_n(\vec{r},\vec{r}')\gamma_\mu\delta G(\vec{r}',\vec{r}|\,\epsilon_n - \omega)\gamma^\mu\}\exp(i\vec{q}(\vec{r} - \vec{r}'))\,. \qquad (4)$$

The averaging over photon polarization ($e_\mu e_\nu \to -g_{\mu\nu}/2$)) is carried out in (4) since the cross section is independent of it. Note that in the case of degeneracy the matrix $\rho_n(\vec{r},\vec{r}')$ and correspondingly the cross section σ_γ (4) is the sum over degenerate states with definite electron energy.

In our paper [8] the convenient integral representation for the electron Green function in a Coulomb field was obtained. Using formulae (19)–(22) of this paper we find for $\rho_n(\vec{r},\vec{r}')$:

$$\rho_n(\vec{r},\vec{r}') = \frac{Z\alpha(-1)^n}{8\pi^2 rr'[(\gamma + n)^2 + (Z\alpha)^2]}\times$$

$$\int_{-\pi/2}^{\pi/2} d\tau \exp\left(i[k(r+r')\tan\tau - 2Z\alpha\tau E_n/k]\right)T, \qquad (5)$$

where n is the radial quantum number, Z is the nucleus charge (in units of $|e|$), $\gamma = (l^2-(Z\alpha)^2)^{1/2}$, $l = j+1/2$, j is the total angular momentum of the state, E_n is the energy of the bound state

$$E_n = m(\gamma+n)/\left((\gamma+n)^2+(Z\alpha)^2\right)^{1/2},$$

$k = \sqrt{m^2-E_n^2}$, and the matrix T has the form:

$$T = \left[1+\vec{n}\vec{n}\,'+i\vec{\Sigma}[n\times\vec{n}\,']\right]\left[\frac{y}{2}J'_{2\gamma}(y)(\gamma^0 E_n+m)-\right.$$

$$\left. iZ\alpha J_{2\gamma}(y)\gamma^0 k\tan(\tau)\right]B + \left[1-\vec{n}\vec{n}\,'-i\vec{\Sigma}[n\times\vec{n}\,']\right]\times$$

$$(\gamma^0 E_n+m)J_{2\gamma}(y)A + imZ\alpha\gamma^0(\vec{\gamma},\vec{n}+\vec{n}\,')J_{2\gamma}(y)B+$$

$$\left[\frac{ik^2(r-r')}{2\cos^2\tau}(\vec{\gamma},\vec{n}+\vec{n}\,')B - k\tan(\tau)(\vec{\gamma},\vec{n}-\vec{n}\,')A\right]J_{2\gamma}(y), \qquad (6)$$

$$A = l\frac{d}{dx}\left[P_l(x)+P_{l-1}(x)\right], \quad B = \frac{d}{dx}\left[P_l(x)-P_{l-1}(x)\right],$$

$$x = \vec{n}\vec{n}\,', \quad \vec{n} = \vec{r}/r, \quad \vec{n}\,' = \vec{r}\,'/r', \quad y = \frac{2k\sqrt{rr'}}{\cos\tau}.$$

Here $J_{2\gamma}$ is the Bessel function. The integration with respect to τ in (5) can be easily carried out by means of the residue theory if we expand the Bessel function in the series, go over to the variable $v = \tan\tau$ and close the integration contour in the upper v half-plane. We have:

$$\int_{-\pi/2}^{\pi/2} d\tau \exp\left(i[k(r+r')\tan\tau - 2Z\alpha\tau E_n/k]\right)J_{2\gamma}\left(\frac{2k\sqrt{rr'}}{\cos\tau}\right) =$$

$$2\pi(-1)^n\sum_{m=0}^{n}\frac{(k^2rr')^{\gamma+M}}{M!\,(n-M)!\,\Gamma(2\gamma+M+1)}\left(\frac{d}{dv}\right)^{n-M}\times$$

$$\left[(1+v)^{2\gamma+n+M-1}\exp\left(-k(r+r')v\right)\right]_{v=1}. \qquad (7)$$

The main contribution to the cross section (4) comes from the distances r, r' of the order of electron Compton wave length $1/m$ and the angles between vectors $\vec{k}$, $\vec{r}$, $\vec{r}\,'$ of the order of unity. The contribution from small distances

$r, r' \sim \omega^{-1}$ as wall as from small angles ($\sim m/\omega$) is suppressed as some power of the small parameter m/ω. For $\omega \gg m$, the energy of a produced positron is approximately equal to ω. Hence, the characteristic values of the positron angular momentum, giving the main contribution to the cross section are large: $l_p \sim \omega r \sim \omega/m \gg 1$. Therefore in the expression for the quantity $\delta G(\vec{r}\,',\vec{r}\,|\,\epsilon)$ (4) corresponding to positron we can use quasiclassical limit of the Green function which is much more simply than exact one. This function was obtained in [9, 10]. Making use of Eq. (5) from [10], we have

$$\delta G(\vec{r}\,',\vec{r}\,|\,\epsilon) = \frac{i\nu}{4\pi}\int_{-\infty}^{\infty} ds \exp\left(i[\nu(r+r')\coth s + 2Z\alpha s\epsilon/\nu]\right)\times$$

$$\left\{J_0(w)\left[\gamma^0\epsilon + m + \frac{\nu}{2}(\vec{\gamma},\vec{n}-\vec{n}\,')\coth s\right] + \frac{iJ_1(w)}{w}\times\right.$$

$$\left.\left[(\vec{\gamma},\vec{n}+\vec{n}\,')\left(\frac{\nu^2(r'-r)}{2\sinh^2 s}+Z\alpha m\gamma^0\right) - Z\alpha\nu\gamma^0(1-(\vec{\gamma}\vec{n}\,')(\vec{\gamma}\vec{n}))\coth s\right]\right\}, \qquad (8)$$

where

$$\nu = (\epsilon^2-m^2)^{1/2}, \quad w = [2rr'(1+\vec{n}\vec{n}\,')]^{1/2}/\sinh s.$$

Substituting (8) and (5) into (4) we can take the trace of γ-matrices in (4). Now some transformations simplifying further calculations can be conveniently done. As the cross section is independent of the momentum direction of an incoming photon, we can average over these directions in (4). Then

$$\exp\left(i\vec{q}(\vec{r}-\vec{r}\,')\right) \to \frac{\sin(\omega|\vec{r}-\vec{r}\,'|)}{\omega|\vec{r}-\vec{r}\,'|}.$$

After that the integrand in (4) turns out to be dependent only on r, r' and $x = (\vec{n}\vec{n}\,')$. Let us introduce new variables: $r = \rho(1+t)/2$, $r' = \rho(1-t)/2$. From what was said above it follows that the main contribution to the cross section is given by $\rho \sim 1/m$, $x \sim 1$, where the arguments of exponential functions in (4), (8) and Bessel functions in (8) are large ($\sim \omega/m$). Therefore we can use asymptotics of the Bessel functions and the stationary phase method to take the integral with respect to t. The stationary point t_0 is

$$t_0 = \left(1 - \frac{2\nu^2}{(1+x)(\nu^2+\omega^2\sinh^2 s)}\right)^{1/2}.$$

Using (7) we can now take the integral over ρ that reduces to Euler's Γ-function. In addition, it is convenient to go over from the variable x to the variable $y = \left(\cosh^2 s - 2/(1+x)\right)^{1/2} / \sinh s$ and to shift the integration contour with respect to variable s: $s \to s - i\pi/2$, after which to go over to the variable $u = \tanh s$. Carrying out the indicated transformations, we obtain the final expression for the bound-electron-positron pair production cross section at $\omega \gg m$ for the electron with definite energy belonging to any state of discrete spectrum:

$$\sigma_\gamma = \frac{2\pi\alpha l\, \beta^{2\gamma+1}\, e^{-\pi Z\alpha}}{\omega m(\gamma+n)\sqrt{1+\beta^2}} \sum_{m=0}^{n} \frac{\Gamma(2\gamma+2M+2)\beta^{2M}}{M!\,(n-M)!\,\Gamma(2\gamma+M+1)2^{2\gamma+2M}} \times$$
$$\left(\frac{d}{dv}\right)^{n-M} (1+v)^{2\gamma+n+M-1} \int_0^1 dy \int_{-1}^{1} du \left(\frac{1-u}{u+1}\right)^{iZ\alpha} \frac{(y^2-u^2)^{\gamma+M}}{g^{2(\gamma+M+1)}(v)} \times$$
$$\left\{\left[\beta^2(1-v^2)(\gamma+M+1)y^2 \frac{1}{g(v)} + \gamma + M + \beta Z\alpha v\right] b + lag(v)\right\}_{v=1}, \quad (9)$$

In this expression

$$a = P_{l-1}(w) - P_l(w), \quad b = P_l(w) + P_{l-1}(w), \quad w = \frac{2-u^2-y^2}{y^2-u^2},$$

$$g(v) = 1 - ivu\beta, \quad \beta = \frac{k}{E_n} = \frac{Z\alpha}{n+\gamma}.$$

The integration contour over u in (9) goes over the real axis. Differentiation with respect to v in (9) can be easily carried out in the explicit form. We keep this differentiation to make the notation shorter. Note else that for any values of l and n with the help of the relation

$$P_l(w) = (-1)^l F\left(-l, l+1; 1; \frac{1-u^2}{y^2-u^2}\right),$$

where $F(a,b;c;x)$ is the hypergeometric function, the integral over y in (9) can be reduced to the integral

$$\int_0^1 dy\,(y^2-u^2)^\lambda = \frac{1}{2}\exp\left(-i\pi\lambda\,\mathrm{sgn}(u)\right)(u^2)^{\lambda+1/2} B(\lambda+1, 1/2) +$$
$$\frac{(1-u^2)^{\lambda+1}}{2(\lambda+1)} F(1/2, 1; \lambda+2; 1-u^2), \quad (10)$$

where bypass rule over the variable u is taken into account. For $n=0$ the expression (9) for the cross section is essentially simplified:

$$\sigma_\gamma(n=0) = \frac{\pi\alpha(2\gamma+1)le^{-\pi Z\alpha}}{\omega m\gamma} \left(\frac{Z\alpha}{\gamma}\right)^{2\gamma+1} \int_0^1 dy \int_{-1}^{1} du \left(\frac{1-u}{u+1}\right)^{iZ\alpha} \times$$
$$\frac{(y^2-u^2)^\gamma}{[1-iZ\alpha u/\gamma]^{2\gamma+2}} \left[(\gamma - iZ\alpha u)a + lb\right]. \quad (11)$$

For $l = 1$ ($1s_{1/2}$-state) it goes over into that obtained in [5] if we make change of variables $u = 1/x$ and deform the contour of integration over x to the closed interval [0, 1]. Eq. (9) is very convenient for performing numerical calculations and obtaining asymptotics, which will be considered in the next section.

3 Asymptotics of the cross section and discussion

Consider the dependence of the cross section on the radial quantum number n, the total angular momentum j (remind that the parameter l in (9) is equal to $j+1/2$) and the quantity $Z\alpha$. Let us start with the case $l \gg 1$, $n \sim 1$. Making use of the integral representation for the Legendre polynomials at the argument larger than unity (see [13]), we easily find that for $l \gg 1$

$$P_l(w) \approx \frac{2^{2l} e^{-(u^2+y^2)l/2}}{\sqrt{\pi l}\,(y^2-u^2)^l}. \quad (12)$$

It follows from (12) that the main contribution to the integral (9) for $l \gg 1$ is given by the region of variables $y \sim u \sim 1/\sqrt{l}$. Substituting (12) into (9), expanding the integrand at small u and taking the integrals, we obtain

$$\sigma_\gamma(l \gg 1) = \frac{\pi\alpha(2l)^n\sqrt{l\pi}}{8\omega m n!} \left(\frac{2Z\alpha}{l+n}\right)^{2l+3} e^{-\pi Z\alpha}. \quad (13)$$

It is seen that even at $Z\alpha \sim 1$ the cross section is numerically suppressed for $l \gg 1$. Comparing (13) with the results of numerical calculations performed in [6] for $l=2$, $n=0$ ($2p_{3/2}$-state), we find that already at $l=2$ an accuracy of (13) in the region $Z\alpha \leq 0.7$ ($Z < 96$) is better than 5%. If we additionally take into account that the contribution of $2p_{3/2}$-state to the cross section when an electron is captured on the L shell ($\sigma_L = \sigma_\gamma(l=1, n=1) + \sigma_\gamma(l=2, n=0)$) does not exceed 11%, than it is clear that the use of (13) provides an accuracy in obtaining σ_L better than 1%.

Consider now the case $n \gg 1$, $l = 1$. The leading term of expansion (9) in powers of $1/n$ has the form

$$\sigma_\gamma = \sigma_0 \frac{F(Z\alpha)}{(\gamma_1 + n)^3}, \qquad (14)$$

where $\sigma_0 = 4\pi\alpha(Z\alpha)^5/m\omega$ and the function $F(Z\alpha)$ reads

$$F(Z\alpha) = e^{-\pi Z\alpha} \int_{-1}^{1} du \left(\frac{1-u}{1+u}\right)^{iZ\alpha} e^{2iZ\alpha u} \sum_{m=0}^{\infty} \frac{(Z\alpha)^{2(d-2)}}{M!\,\Gamma(\gamma+d+1)} \times$$

$$\int_0^1 dy\,(y^2-u^2)^{d-1}\{R(2d+1)\left[R(1-u^2)+(y^2-1)\right] +$$

$$(Z\alpha)^2(1-u^2)\left[(1-y^2)(2R+1)+u^2-y^2\right]\} \qquad (15)$$

where

$$d = M + \gamma_1, \quad \gamma_1 = \sqrt{1-(Z\alpha)^2}, \quad R = d + iZ\alpha u.$$

The function $F(Z\alpha)$ tends to unity at $Z\alpha \to 0$ in accordance with calculations performed in Born-approximation [12]. Note that, as it follows from (13), when an electron is captured into the state with $l \geq 2$ the process cross section contains at $Z\alpha \ll 1$ an additional (as compared with $l = 1$ case) suppressing factor $\propto (Z\alpha)^{2(l-1)}$ (see also discussion in [6]). The result of summation in (15) can be expressed via the Bessel function $J_{2\gamma_1}(2Z\alpha\sqrt{u^2-y^2})$ and its derivative. Such form of notation can be directly obtained from (4) if we note that for $n \gg 1$ small values of $\pi/2 - |\tau| \sim k/E_n = Z\alpha/(n+\gamma)$ contribute to the integral in (5). At such derivation of the asymptotics it can be easily proved that in evaluating corrections to (14) the quantity $(k/E_n)^2$ is the parameter of expansion. With account for the first correction in this parameter the process cross section for $n \gg 1$ takes the form

$$\sigma_\gamma = \sigma_0 \frac{F(Z\alpha)}{(\gamma_1+n)^3}\left[1 - \left(\frac{Z\alpha}{\gamma_1+n}\right)^2 f(Z\alpha)\right]. \qquad (16)$$

In view of awkwardness of the function $f(Z\alpha)$ we don't give here its explicit form. Estimations show that for $0.1 < Z\alpha < 0.9$ the function $f(Z\alpha)$ changes

(not monotonically) in the interval $0.9 < f(Z\alpha) < 2$. This permits one to estimate an accuracy of (14). It is interesting that if we simply put $f(Z\alpha) = 4/3$ in (16) than this formula fits for $Z\alpha < 0.7$ the results of numerical calculations performed in [6] within 1.6%.

The function $F(Z\alpha)$ is plotted in Fig.1. At calculating $F(Z\alpha)$ the integral over y was taken by means of (10) after which the integral over u was

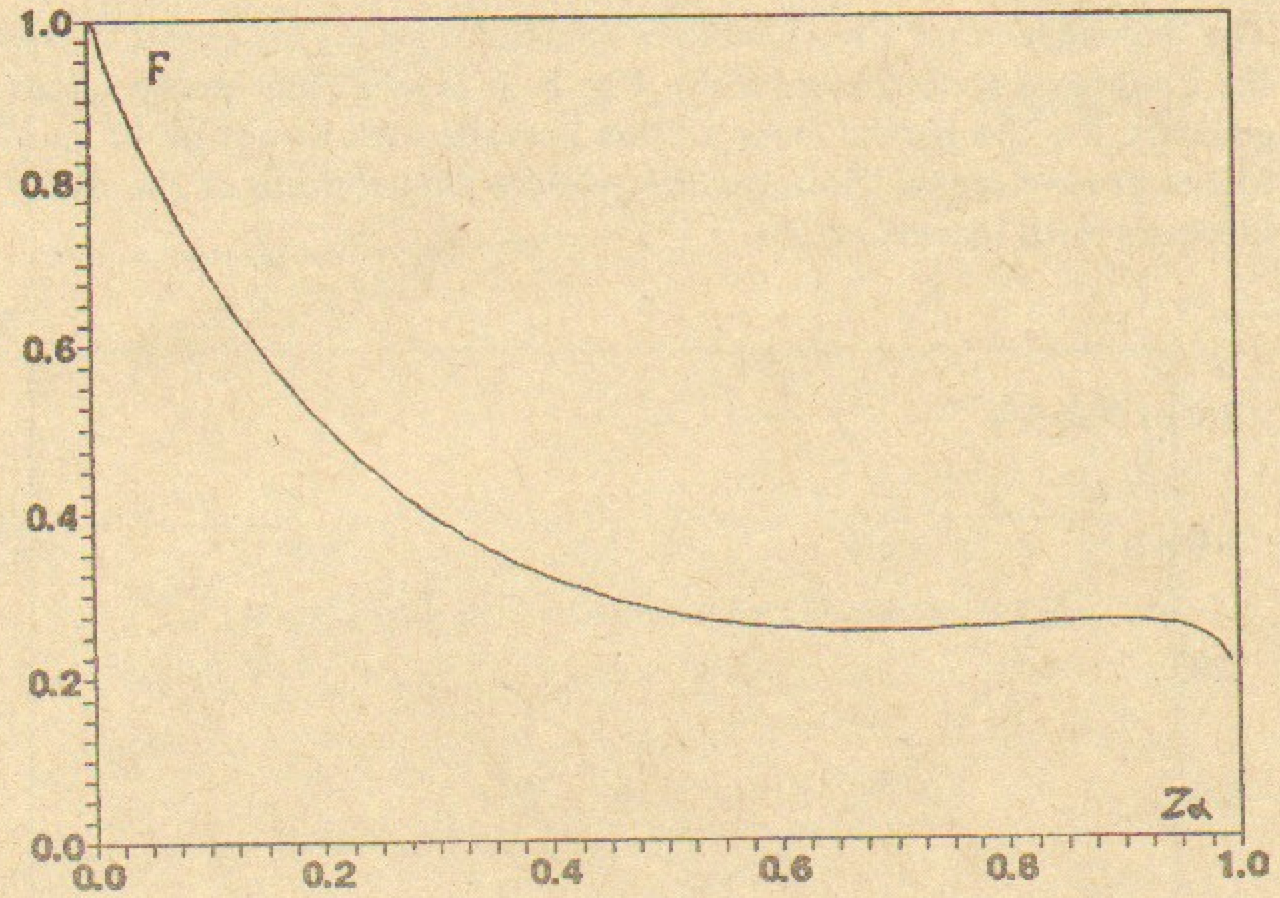

Fig. 1. Function $F(Z\alpha)$ determining dependence of the cross section σ_γ on $Z\alpha$ for $n \gg 1$ (see (15)).

expressed in terms of hypergeometric functions, so that the result was represented as a double sum containing these functions. Note that with increasing $Z\alpha$ the function $F(Z\alpha)$ first decreases rapidly owing to the factor $\exp(-\pi Z\alpha)$ in (15) while in the region $0.55 < Z\alpha < 0.95$ it is almost constant.

Obtained approximate expressions (13) and (16) for partial (with given values of l and n) cross sections $\sigma_\gamma(l, n)$ provide finding within an accuracy of 1% the total cross section of the process

$$\sigma_{tot} = \sum_{n=0}^{\infty}\sum_{l=1}^{\infty} \sigma_\gamma(l, n).$$

For contributions to this sum from K shell ($1s_{1/2}$-state, $\sigma_K = \sigma_\gamma(1,0)$) and L shell ($2s_{1/2}$, $2p_{1/2}$ and $2p_{3/2}$ states, $\sigma_L = \sigma_\gamma(1,1) + \sigma_\gamma(2,0)$) we use the numerical results obtained in [6]. For the sum of all the other contributions $\sigma_{M'} = \sigma_{\rm tot} - \sigma_K - \sigma_L$ we use (13) and (16). The dependence on $Z\alpha$ of the ratio $\sigma_{\rm tot}/\sigma_0$ is shown in Fig. 2 in the interval $Z\alpha < 0.7$. It is seen that an exact account for a Coulomb field drastically changes the result as compared to Born approximation.

The dependence on $Z\alpha$ of relative contributions $\sigma_L/\sigma_{\rm tot}$ and $\sigma_{M'}/\sigma_{\rm tot}$ to the total cross section is shown in Fig. 3. It is seen that at any values of the quantity $Z\alpha$ the partial cross section prevails with a capture of an electron to the ground state. However the relative contribution of the other states increases with increasing $Z\alpha$.

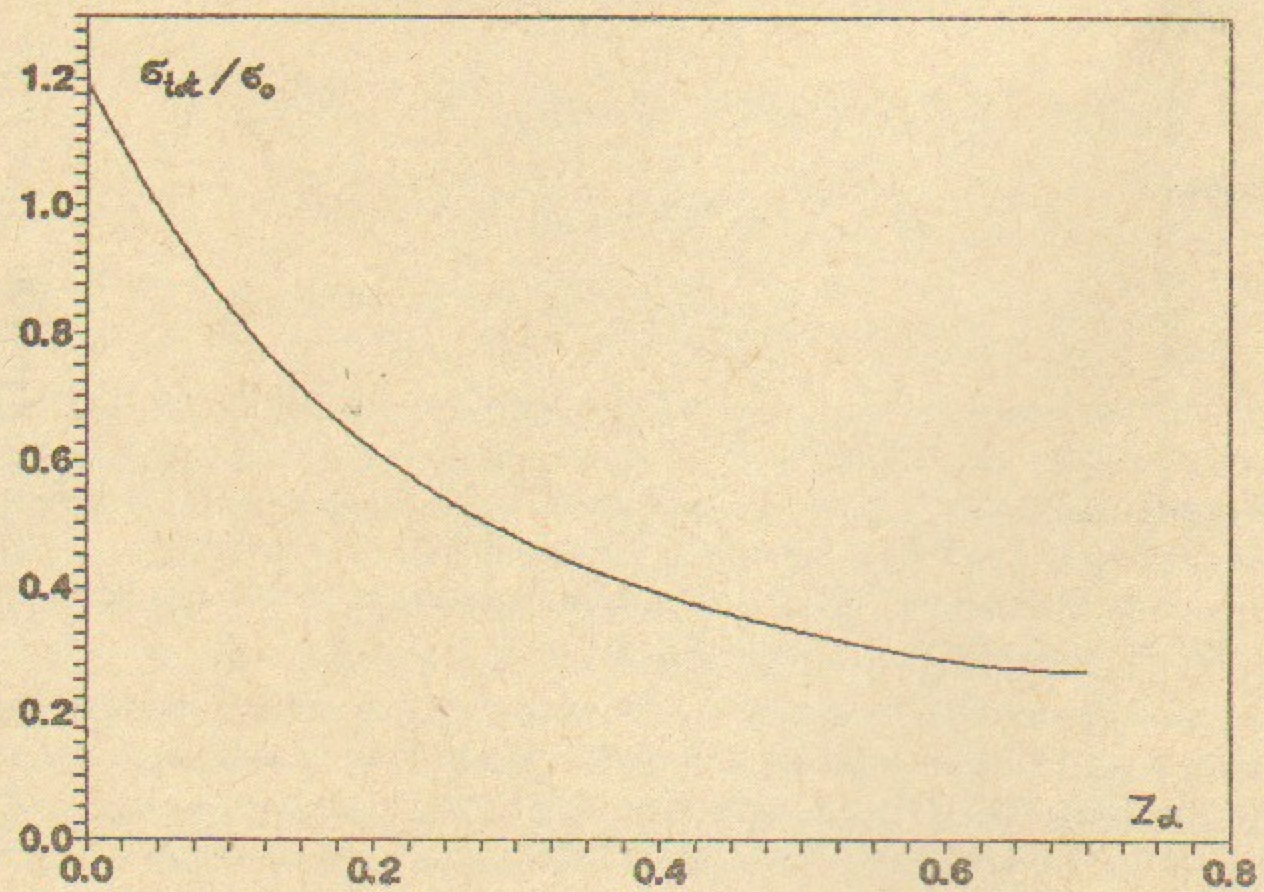

Fig. 2. Total cross section of the process in units of σ_0.

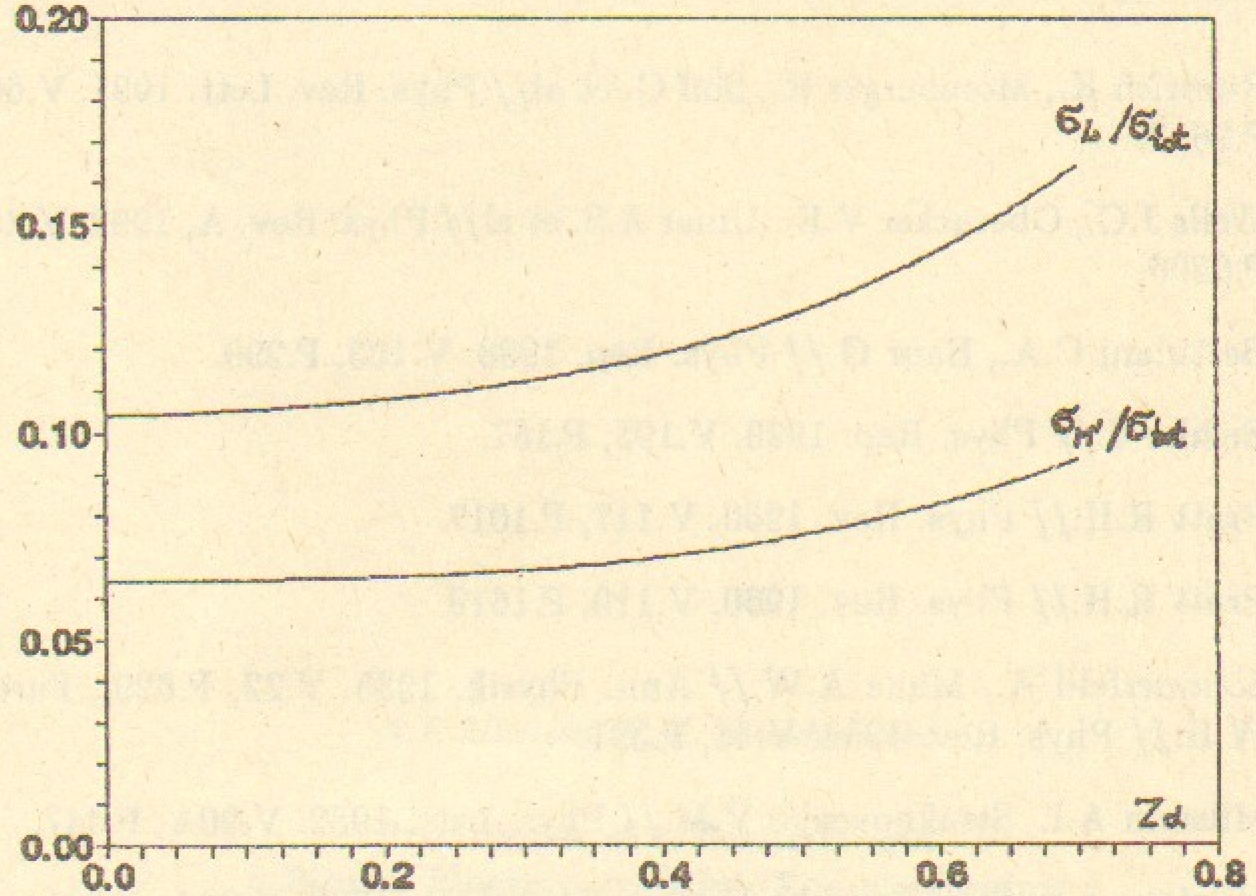

Fig. 3. Dependence of the relative contributions $\sigma_L/\sigma_{\rm tot}$ and $\sigma_{M'}/\sigma_{\rm tot}$ on $Z\alpha$.

References

[1] Rumrich K., Momberger K., Soff G. et al// Phys. Rev. Lett. 1991. V.66, P.2613.

[2] Wells J.C., Oberacker V.E., Umar A.S. et al// Phys. Rev. A. 1992. V.45, P.6296.

[3] Bertulani C.A., Baur G.// Phys. Rep. 1988. V.163, P.299.

[4] Eichler J.// Phys. Rep. 1990. V.193, P.167.

[5] Pratt R.H.// Phys. Rev. 1960. V.117, P.1017.

[6] Pratt R.H.// Phys. Rev. 1960. V.119, P.1619.

[7] Sommerfeld A., Maue A.W.// Ann. Physik. 1935. V.22, P.629.; Furry W.H.// Phys. Rev. 1934. V.46, P.391.

[8] Milstein A.I., Strakhovenko V.M.// Phys. Lett. 1982. V.90A, P.447.

[9] Milstein A.I., Strakhovenko V.M.// Phys. Lett. 1982. V.92A, P.381.

[10] Milstein A.I., Strakhovenko V.M.// Sov. Phys. -JETP 1983. V.58, P.8.

[11] Berestetsky V.B., Lifshitz E.M., Pitaevsky L.P.// Quantum Electrodynamics. Pergamon Press, Oxford, 1982.

[12] Sauter F.// Ann. Physik. 1931. V.11, P.454.

[13] Gradstein I.S. and Ryzhik I.W.// Tables of Integrals, Series and Products. Academic Press, New York, 1965

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Production of a Positron and a Bound Electron by High-Energy Photon in a Strong Coulomb Field

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Рождение позитрона и связанного электрона фотоном большой энергии в сильном кулоновском поле

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