Diffraction of Čerenkov radiation in a spatially periodic medium

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Some Čerenkov radiation features due to diffraction are considered for a particle moving in a medium with a spatially periodic distribution of the dielectric constant along planes with ϵ = const. It is shown that discontinuities due to diffraction occur in the Čerenkov cone. The positions of the discontinuities on the cone depend on the wavelength of the emitted light and on the particle velocity. In principle, the dependence of the position of the discontinuity on the particle velocity for a given wavelength can be used for mass separation of ultrarelativistic particles.

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1. It is known that the diffraction of light in a medium with a spatially periodic distribution of the refractive index is similar to the diffraction of x rays in a single crystal and is described by the dynamical theory of diffraction.[1-3] As such a medium it is possible to use, for example, a holographic lattice with a dielectric constant depending on the coordinates like

$$\varepsilon=\varepsilon_0+2\varepsilon_g\cos \mathbf{gr} \tag{1}$$

(**g** is the reciprocal lattice vector, $|\mathbf{g}|=2\pi/d$, and d is the spatial period), or a laminar medium consisting of layers with different dielectric constants. When a charged particle moves in such a medium, in addition to Čerenkov radiation, transition radiation is emitted, due to the fact that a periodic medium can either acquire or give up momentum in multiples of $\hbar\mathbf{g}$.

The theory of the radiation of a charged particle moving in a laminar medium in the direction perpendicular to the planes ε=const was developed by Faĭnberg and Khizhnyak.[4] Ter-Mikaélyan[5] considered Čerenkov and transition radiation in the eikonal approximation (i.e., far from the Bragg condition, neglecting diffraction) in a medium with a dielectric constant given by (1). The cylindrically symmetric problem of the radiation of a charged particle moving along the axis of a cylindrical lattice with a radial distribution of the dielectric constant $\varepsilon=\varepsilon_0+\varepsilon_g\cos g\rho$ was solved by Bolotovskiĭ and Chukhovskiĭ.[6] Belyakov, Dmitrienko, and Orlov[7] discussed the polarization characteristics and the ratio of the radiation intensities in the principal cone and the diffraction cone produced as a result of the diffraction of the Čerenkov radiation of a particle moving in a cholesteric crystal by the periodic structure of this crystal.

In the present study we consider the diffraction-related features of the radiation of a particle moving a periodic medium along the planes ε=const, i.e., perpendicular to the vector **g**. In this case there is no transition radiation (or, more precisely, when the recoil momentum is neglected the transition and Čerenkov radiations coincide) and the reflection cone coincides with the principal cone. It will be shown that diffraction leads to discontinuities in the surface of the Čerenkov cone, which attest to the presence of the two branches of the dispersion surface of a photon in a periodic lattice. For a given photon frequency the location of these discontinuities on the Čerenkov cone strongly depends on the particle velocity, which in principle can be used to mass-separate ultrarelativistic particles. The results are easily generalized to the case of arbitrary particle motion, in which the transition radiation, like the Čerenkov radiation, is due to the particular form of the refractive index for photons in the periodic medium.

2. According to the dynamical theory of diffraction in the two-ray approximation, propagation of light through a spatially periodic medium gives rise to two refractive indices related to the two branches of the dispersion surface:

$$(n^{(1,2)})^2=k^2/k_0^2=\varepsilon_0-\delta\pm(\delta^2+\varepsilon_g^2)^{1/2}, \tag{2}$$

where

$$\delta=\frac{k_g^2-k^2}{2k_0^2}=\frac{2\mathbf{kg}+g^2}{2k_0^2}\approx\varepsilon_0\left(\frac{\alpha^2}{2}-\alpha\sin\chi\right) \tag{3}$$

is the parameter of the deviation from the Bragg condition, χ is the angle between the photon wave vector in the medium **k** and the plane ε=const, $k_0=\omega/c$, ω is the photon frequency, and $\alpha=\lambda/d$, with λ the photon wavelength in the medium.

Stipulating henceforth that the photon field is expanded in the reciprocal lattice vectors $n\mathbf{g}$ ($n=0, \pm1, \pm2, \ldots$), the projection of the vector **k** on the direction of **g** varies from 0 to $-|\mathbf{g}|$ ($-g^2<\mathbf{k}\cdot\mathbf{g}\leqslant 0$). We also assume that at the end points of the $\mathbf{k}\cdot\mathbf{g}$ range the parameter δ is sufficiently large, so that there is no diffraction, i.e., $\varepsilon_0\alpha^2/2\gg\varepsilon_g$.

The operator for the interaction of a charged particle with the electromagnetic field responsible for the radiation in a periodic medium can be written as[8]

$$V=\frac{e}{mc}[\mathbf{A}_k^{(1)}\mathbf{p}\exp(i\omega_k^{(1)}t)+\mathbf{A}_k^{(2)}\mathbf{p}\exp(i\omega_k^{(2)}t)], \tag{4}$$

where

$$\mathbf{A}_k^{(1)}=A_0[\mathbf{e}_k\cos\gamma e^{-i\mathbf{kr}}+\mathbf{e}_{k_g}\sin\gamma\exp(-i\mathbf{k}_g\mathbf{r})], \tag{5}$$

$$\mathbf{A}_k^{(2)}=A_0[-\mathbf{e}_k\sin\gamma e^{-i\mathbf{kr}}+\mathbf{e}_{k_g}\cos\gamma\exp(-i\mathbf{k}_g\mathbf{r})] \tag{6}$$

are the Bloch functions of the photon in a periodic medi-

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um, $A_0=(2\pi\hbar c/kn_0)^{1/2}$, $n_0=\varepsilon_0^{1/2}$ (in normalizing the functions $\mathbf{A}_k^{(1)}$ and $\mathbf{A}_k^{(2)}$ we have neglected the difference arising from the small difference in the refractive indices), $\omega_k^{(\alpha)}=ck/n^{(\alpha)}$ $(\alpha=1,2)$, $\tan 2\gamma=\varepsilon_g/\delta$, and $\mathbf{p}=-i\hbar\nabla$ is the particle momentum operator. The intensities of the direct and reflected waves in the two branches (1) and (2) are determined by the quantity

$$\cos^2\gamma=\frac{1}{2}\left(1+\frac{\delta}{(\delta^2+\varepsilon_g^2)^{1/2}}\right). \qquad (7)$$

For positive values of δ far from the Bragg condition, i.e., for $\delta\gg\varepsilon_g$ ($\cos^2\gamma\approx 1$, $\sin^2\gamma\approx 0$), the first branch is a direct wave with wave vector $\mathbf{k}$ and frequency $\omega^{(1)}=ck/\varepsilon_0^{1/2}$ and the second branch is a reflected wave with wave vector $\mathbf{k}_g$ and frequency[1)]

$$\omega^{(2)}=ck/(\varepsilon_0-2\delta)^{1/2}=ck_g/\varepsilon_0^{1/2}.$$

On the other hand, for negative $\delta\ll-\varepsilon_g$ ($\cos^2\gamma\approx 0$, $\sin^2\gamma\approx 1$), the first branch describes a reflected wave and the second describes a direct wave. Therefore, far from the Bragg condition expressions (5) and (6) describe photons propagating in the direction of either $\mathbf{k}$ or $\mathbf{k}_g$, depending on the sign of δ, with a mean refractive index n_0, just as in a homogeneous medium. The transition from the direct wave to the reflected wave with changing δ occurs in the region $|\delta|\lesssim\varepsilon_g$. Here the direct and reflected waves are mixed, which causes the refractive index to change. At $\delta=0$ (exact satisfaction of the Bragg condition) the mixing is a maximum, which results in waves propagating along the planes $\varepsilon=$const and with a modulation in the direction of $\mathbf{g}$ having an intensity $\cos^2(\mathbf{gr}/2)$ for the first branch or $\sin^2(\mathbf{gr}/2)$ for the second branch, so that the intensity maxima of a wave belonging to the first branch occur on the plane $\varepsilon=\varepsilon_{max}=\varepsilon_0+2\varepsilon_g$, while those of a wave in the second branch occur on the plane $\varepsilon=\varepsilon_{min}=\varepsilon_0-2\varepsilon_g$. The mean refractive indices are given by $\varepsilon^{(1)}=\varepsilon_0+\varepsilon_g$ and $\varepsilon^{(2)}=\varepsilon_0-\varepsilon_g$, which also follows from (2) at $\delta=0$.

3. The amplitude for the transition of a particle from a state of momentum $\hbar\mathbf{k}_i$ to a state of momentum $\hbar\mathbf{k}_f$ with the emission of a quasiphoton of energy $\hbar\omega_k^{(1)}$ or $\hbar\omega_k^{(2)}$ is

$$a_{fi}(t)=-\frac{i}{\hbar}\int_0^t V_{fi}(t')\,dt'=\frac{e}{mc}\frac{A_0}{\hbar}\Big\{[\mathbf{e}_k\mathbf{p}\cos\gamma\,\delta(\mathbf{k}_f-\mathbf{k}_i+\mathbf{k})$$
$$+\mathbf{e}_{k_g}\mathbf{p}\sin\gamma\,\delta(\mathbf{k}_f-\mathbf{k}_i+\mathbf{k}_g)]\frac{\exp(i\omega_{fi}^{(1)}t)-1}{\omega_{fi}^{(1)}}+[-\mathbf{e}_k\mathbf{p}\sin\gamma\,\delta(\mathbf{k}_f-\mathbf{k}_i+\mathbf{k})$$
$$+\mathbf{e}_{k_g}\mathbf{p}\cos\gamma\,\delta(\mathbf{k}_f-\mathbf{k}_i+\mathbf{k}_g)]\frac{\exp(i\omega_{fi}^{(2)}t)-1}{\omega_{fi}^{(2)}}\Big\}. \qquad (8)$$

Here $\omega_{fi}^{(\alpha)}=\omega_f-\omega_i+\omega^{(\alpha)}$, where $\hbar\omega_f$ and $\hbar\omega_i$ are the final and initial particle energies, and $\mathbf{p}=\hbar\mathbf{k}_i=m\mathbf{v}$.

If the momentum conservation laws $\mathbf{k}_f-\mathbf{k}_i=\mathbf{k}$ are satisfied, we find

$$\omega_{fi}^{(\alpha)}=\frac{ck}{n^{(\alpha)}}\left(1-\frac{v}{c}n^{(\alpha)}\cos\theta\right). \qquad (9)$$

The equality $\omega_{fi}^{(\alpha)}=0$ in this case formally gives the usual condition for Vavilov-Čerenkov radiation of a particle moving with velocity v in a medium with refractive index $n^{(\alpha)}$.

At $\mathbf{k}_f-\mathbf{k}_i=\mathbf{k}_g$ we have

$$\omega_{fi}^{(\alpha)}=\frac{ck}{n^{(\alpha)}}\left[1-\frac{v}{c}n^{(\alpha)}\cos\theta+\frac{\mathbf{vg}}{\omega^{(\alpha)}}\right]=\omega^{(\alpha)}\left(1-\frac{v}{c}n^{(\alpha)}\cos\theta\right)+\mathbf{vg}. \qquad (10)$$

In this case $\omega_{fi}^{(\alpha)}=0$ is the formal condition for transition radiation.[5]

We note, however, that the subdivision of the particle radiation into Čerenkov and transition radiation is arbitrary in this analysis. This is because far from the Bragg condition one branch describes a direct wave and the other describes a reflected wave, whereas the refractive index $n^{(\alpha)}$ connects ω^{α} with the wave vector of the direct wave (see also footnote 1). It is therefore easily verified that, for example, $\omega_{fi}^{(1)}=0$ from (9) describes Čerenkov radiation at $\delta\gg\varepsilon_g$, but transition radiation at $\delta\ll\varepsilon_g$: These two types of radiation differ physically in that Čerenkov radiation has a particle-velocity threshold and a continuous frequency spectrum far from the Bragg condition, while the transient-radiation frequency is determined by the frequency at which the particles intersect the maxima of the lattice dielectric constant $\mathbf{v}\cdot\mathbf{g}$ and by the Doppler effect as a function of the direction of the radiation [see (10)]. However, when the Bragg condition is satisfied the transition radiation is actually reflected Čerenkov radiation, so that it is not possible to separate the two by a diffraction study.

4. Let us now consider the radiation of a particle moving along the planes $\varepsilon=$constant, i.e., $\mathbf{v}\perp\mathbf{g}$. In this case expressions (9) and (10) coincide and give identical radiation conditions. The emission probability in a time t, integrated over the final electron moments, will be

$$dP_{fi}(t)=\frac{e^2}{\hbar c}\frac{\pi v^2}{\varepsilon_0^{1/2}k}\Big\{(\mathbf{e}_k\mathbf{n})^2\Big[\Delta^2(\omega_{fi}^{(1)})\cos^2\gamma+\Delta^2(\omega_{fi}^{(2)})\sin^2\gamma$$
$$-2\sin\gamma\cos\gamma\cos\frac{\omega_{fi}^{(2)}-\omega_{fi}^{(1)}}{2}t\,\Delta(\omega_{fi}^{(1)})\Delta(\omega_{fi}^{(2)})\Big]$$
$$+(\mathbf{e}_{k_g}\mathbf{n})^2\Big[\Delta^2(\omega_{fi}^{(1)})\sin^2\gamma+\Delta^2(\omega_{fi}^{(2)})\cos^2\gamma$$
$$+2\sin\gamma\cos\gamma\cos\frac{\omega_{fi}^{(2)}-\omega_{fi}^{(1)}}{2}t\,\Delta(\omega_{fi}^{(1)})\Delta(\omega_{fi}^{(2)})\Big]\Big\}\frac{d^3k}{(2\pi)^3}, \qquad (11)$$

where $\Delta(\omega)=\sin(\omega t/2)(\omega/2)^{-1}$ and $\mathbf{n}$ is a unit vector in the direction of the particle momentum.

From (11) we see that when a particle moves along the planes $\varepsilon=$const two branches of Čerenkov radiation are emitted, corresponding to two refractive indices of refraction, and to each direction of direct radiation $\mathbf{k}$ there corresponds reflected radiation $\mathbf{k}+\mathbf{g}$. In each of these directions the two branches of Čerenkov radiation interfere with each other. At sufficiently large times the interference terms can be neglected and then, as $t\to\infty$ in (11), we find the following for the transition probability per unit time in the momentum interval d^3k:

$$dw=\frac{e^2}{\hbar c}\frac{v^2}{\varepsilon_0^{1/2}}\{(\mathbf{e}_k\mathbf{n})^2[\cos^2\gamma\,\delta(\omega_{fi}^{(1)})+\sin^2\gamma\,\delta(\omega_{fi}^{(2)})]$$
$$+(\mathbf{e}_{k_g}\mathbf{n})^2[\sin^2\gamma\,\delta(\omega_{fi}^{(1)})+\cos^2\gamma\,\delta(\omega_{fi}^{(2)})]\}k\,dk\frac{d\Omega}{2\pi}. \qquad (12)$$

In spherical coordinates (with the polar axis directed along the particle momentum and the angle φ measured from the direction of $\mathbf{g}$) the parameter δ has the form

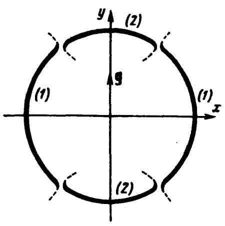

FIG. 1. Plots of $\tan\theta^{(1,2)}(\varphi)$ (in polar coordinates) on which the radiation condition $\cos\theta^{(1,2)}=1/\beta n^{(1,2)}(\theta,\varphi)$ is satisfied for particles with $\beta^2=0.99990$ (for example, 100-GeV protons). Here $\varepsilon_0=1.2$, $\varepsilon_g=0.01$, and $\alpha=0.66$. The length of the radius vector $r\propto\tan\theta$ and the angle φ is measured from the y axis. The line thickness is proportional to the radiation intensity (at the dashed lines the intensity is practically zero).

$$\delta=\varepsilon_0(\alpha^2/2-\alpha\sin\theta\cos\varphi); \tag{13}$$

here the angle θ and φ characterize the vector **k** $(0\leqslant|\varphi|<\pi/2)$. Changing in the second term in the square brackets in (12) to the variables characterizing the vector $\mathbf{k}_g$ and expanding the region of variation of the variables to all values of θ, φ, and k, expression (12) can be rewritten as

$$dw=\frac{e^2}{\hbar c}\frac{v^2}{\varepsilon_0^{1/2}}\frac{\sin^3\theta}{2}\left\{\left(1+\frac{\delta}{(\delta^2+\varepsilon_g^2)^{1/2}}\right)\delta\left(\omega^{(1)}\left[1-\frac{v}{c}n^{(1)}\cos\theta\right]\right)+\left(1-\frac{\delta}{(\delta^2+\varepsilon_g^2)^{1/2}}\right)\delta\left(\omega^{(2)}\left[1-\frac{v}{c}n^{(2)}\cos\theta\right]\right)\right\}k\,dk\,d\theta\frac{d\varphi}{2\pi}. \tag{14}$$

Here

$$\delta=\varepsilon_0(\alpha^2/2-\alpha\sin\theta|\cos\varphi|),\quad 0\leqslant\varphi<2\pi. \tag{13'}$$

The modulus of the cosine corresponds to the fact that for the reflected wave the Bragg condition is satisfied at $|\varphi|>\pi/2$ and came about because we expanded the region of variation of the variables.

5. From (14) we see that at $\beta n^{(2)}>1$ both branches of the Čerenkov radiation are excited ($\beta=v/c$). At $\delta\gg\varepsilon_g$ there is only the first branch that coincides with ordinary Čerenkov radiation in a homogeneous medium at an angle $\theta_0=\arccos(1/\beta\varepsilon_0^{1/2})$, since here $n^{(1)}\approx\varepsilon_0^{1/2}$. With decreasing δ the intensity of this branch begins to fall off, while the refractive index $n^{(1)}$ and, correspondingly, the angle θ increase. When the Bragg condition is exactly satisfied at some angle φ that depends on α, the intensity falls off by half and rapidly decreases with further decrease of δ to negative values. However, when the

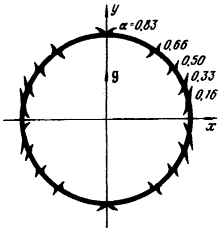

FIG. 2. The family of $\tan\theta^{(1,2)}(\varphi)$ plots for different values of α (the frequency resolution of the Čerenkov radiation spectrum against the angle φ). Here $\varepsilon_0=1.2$, $\varepsilon_g=0.01$, and $\beta^2=0.99990$.

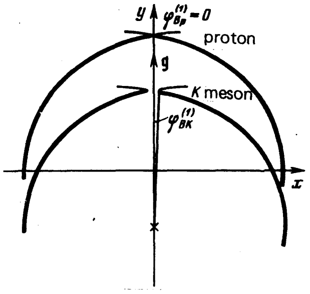

FIG. 3. Curves $\tan\theta^{(1)}(\varphi)$ for protons and K mesons of energy 100 GeV ($\beta_p^2=0.99990$, $\beta_K^2=0.999975$, $\varepsilon_0=1.05$, and $\varepsilon_g=0.01$) for a fixed proton frequency ($\alpha_0=n^{(1)}(\theta,\varphi)\ \alpha=0.49$). The centers of the curves are shifted arbitrarily.

first branch begins to fall off in intensity, the second branch appears at an angle $\theta<\theta_0$ since $n^{(2)}<\varepsilon_0^{1/2}$; its intensity increases with decreasing δ and at $\delta=0$ is comparable to that of the first branch; here the angle θ is determined by $n^{(2)}=(\varepsilon_0-\varepsilon_g)^{1/2}$. As δ decreases further to negative values, the second branch becomes ordinary Čerenkov radiation since $n^{(2)}\to\varepsilon_0^{1/2}$ and $\sin^2\gamma\to1$.

Thus, at certain values of the angle φ corresponding to satisfaction of the Bragg condition for photons with a given α, discontinuities appear on the Čerenkov cone $\theta=\theta_0$; part of the surface of the cone is turned inwards to the region $\theta<\theta_0$ (second branch), while the other part is turned outwards to the region $\theta>\theta_0$ (first branch) (Fig. 1). Since the location of the discontinuities (the angle φ_B) depends on the photon wavelength, the projection of the Čerenkov radiation on a plane perpendicular to the particle momentum acquires two colored rings in the angular ranges $\theta>\theta_0$ and $\theta<\theta_0$, each with its own wavelength resolution of the Čerenkov radiation as a function of the angle φ (Fig. 2). This resolution is given by the Bragg condition

$$|\cos\varphi_B^{(1,2)}|=\frac{\alpha}{2\sin\theta_B^{(1,2)}}=\frac{\lambda}{2d\sin\theta_B^{(1,2)}}, \tag{15}$$

where

$$\sin\theta_B^{(1,2)}=(1-1/\beta^2(\varepsilon_0\pm\varepsilon_g))^{1/2}. \tag{16}$$

The maximum wavelengths in both branches $\lambda_{max}^{(1,2)}$ are emitted at the angles $\varphi_B^{(1,2)}=0$ and $\varphi_B^{(1,2)}=\pi$ and are determined by the particle velocity:

$$\lambda_{max}^{(1,2)}=2d\sin\theta_B^{(1,2)}. \tag{17}$$

Since the arcsine depends very strongly on the argument near zero (the derivative tends to infinity), a small change in the particle velocity at fixed λ can lead to a significant change in the angle φ_B of the location of a discontinuity on the cone [see (15) and (16)]. This fact can in principle be used to mass-separate ultrarelativistic particles with a given momentum. In fact, for particles of energy E and mass m_1 ($E\gg m_1c^2$) assume that $\varphi_{1B}^{(1)}=0$ (at this wavelength the second branch is not excited, since $\alpha/2\sin\theta_B^{(2)}>1$). Then it is possible to satisfy the Bragg condition at this wavelength only for the radiation of lighter particles ($m_2\leqslant m_1$) of the same energy. Here it is easily shown that the angle $\varphi_{2B}^{(1)}$ is given by

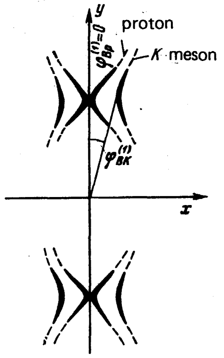

FIG. 4. Plots of $\tan\theta^{(1)}(\varphi)$ for protons and K mesons of energy 100 GeV for $\varepsilon_0=1$ and $\varepsilon_g=0.001$.

$$(\varphi_{2B}^{(1)})^2 \cong \frac{1}{\mathrm{tg}^2\,\theta_\infty^{(1)}}\left(\frac{1}{\mu_1^2}-\frac{1}{\mu_2^2}\right), \quad (18)$$

where $\cos\theta_\infty^{(1)}=1/(\varepsilon_0+\varepsilon_g)^{1/2}$ and $\mu_{1,2}=E/m_{1,2}c^2$ is the total energy of the particle in rest-mass units ($\beta^2=1-1/\mu^2$).

The corresponding change of the angle $\theta_B^{(1)}$, which is used in ordinary Čerenkov counters to separate particles according to their masses, is

$$\Delta\theta_B^{(1)} \approx \frac{1}{2\,\mathrm{tg}\,\theta_B^{(1)}}\left(\frac{1}{\mu_1^2}-\frac{1}{\mu_2^2}\right). \quad (19)$$

Comparing (18) and (19), it is easily seen that $\varphi_{2B}^{(1)} \gg \Delta\theta_B^{(1)}$. Figures 3 and 4 give the numerical calculation of the corresponding curves $\tan\theta^{(1)}(\varphi)$ for protons and K mesons of energy 100 GeV.

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1) Actually, $k_g^2/k_0^2 \equiv (n_g^{(1,2)})^2=(n^{(1,2)})^2+2\delta=\varepsilon_0+\delta\pm(\delta^2+\varepsilon_g^2)^{1/2}$.

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Two-step photoionization of helium via the $4p\,^1P_1$ state

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A new experimental method is proposed for the investigation of the photoionization of atoms from short-lived states, the excitation energy of which corresponds to the vacuum ultraviolet. The excited atoms are prepared by vacuum ultraviolet radiation from laser plasma. The technique has been used to investigate two-step photoionization of helium atoms. The ionization cross section σ_{4p} in the $4p\,^1P_1$ state has been measured. A theoretical analysis is given of two-step resonance photoionization in one-electron approximations, taking into account correlations between atomic electrons. Calculations of σ_{4p} in the Hartree-Fock approximation are in satisfactory agreement with the measured values, whereas the contribution of correlation processes to σ_{4p} is small.

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1. INTRODUCTION

Considerable progress in the photonionization of atoms from the ground state has been achieved as a result of the application of new experimental techniques and the development of theoretical methods capable of taking into account many-particle effects.[1] Ionization from excited states has not been studied to the same extent despite the fact that it is of considerable interest from the point of view of the properties of wave functions and the structure of complex atoms. Such studies are of considerable importance for applied purposes.

Direct experimental study of elementary interactions between photons and short-lived excited atoms require the use of high-intensity exciting and ionizing beams. Two-step photoionization of alkali atoms has been investigated in some isolated cases with the aid of gas-

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