ON THE PARAMETRIC X-RAYS ALONG THE VELOCITY OF AN EMITTING PARTICLE

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The forward PXR from relativistic electrons crossing a thick crystalline target (the thickness of the target exceeds an absorption length) is studied theoretically with account of the contributions of PXR, transition radiation and bremsstrahlung to the formation of a total emission yield. The peculiarities in an emission spectrum observed in the recent Tomsk experiment are explained on the basis of the developed theory.

1. INTRODUCTION

When a fast charged particle crosses a crystalline target it can emit X-rays due to the coherent Bragg scattering of its Coulomb field by system of parallel atomic planes of the crystal. This emission mechanism well known as the parametric X-rays (PXR) has been studied theoretically and experimentally in detail [1—6], but this statement is true for the emission flux propagating to the Bragg scattering direction only. An additional PXR flux propagating along an emitting particle velocity has been observed last time only [7, 8].

Two different experimental procedures have been used in experiments [7, 8] in order to separate a low-power PXR signal on the background of broadband bremsstrahlung and transition radiation. A thin light crystal has been used in [7] to suppress an influence of photoabsorption and multiple scattering (and bremsstrahlung consequently). The main contribution to background in the experiment [7] was formed by transition radiation. The negative interference between transition radiation waves emitted from in and out surfaces of the target was used to suppress this contribution. Mentioned advantages of the discussed approach was attended by one substantial disadvantage consisting in the occurrence of an additional X-ray flux due to dynamical diffraction effect in the transition radiation emitted from a crystal [9].

An alternative experimental procedure was used in the experiment [8]. The used heavy crystalline target (tungsten) allowed to increase substantially the spectral width of PXR signal. Because of this the effect of averaging of the narrow spectrally PXR peak on background of broadband bremsstrahlung and transition radiation (this effect was caused by a finite energy resolution of X-ray detector) was suppressed. In addition to this there were no dynamical effects in the transition radiation contribution since a photoabsorption length was less then the thickness of the target. On the other hand the contribution of bremsstrahlung was of fundamental importance in the emission process [8]. It has been shown in the experiment that the structure of total emission spectrum in the vicinity of Bragg frequency where PXR yield was concentrated depends strongly on the value of orientation angle between the velocity of emitting particle and reflecting crystallographic plane. Among other things the transformation of PXR peak into the dip hole in emission spectrum was observed in [8] under changes of the orientation angle.

The aim of this work is to present the theoretical explanation of the peculiarities of experimental results [8]. Two wave approximation of the dynamical diffraction theory [10] and the system of units $\hbar = c = 1$ are used in the work.

2. THE INTENSITY OF X-RAYS EMITTED FROM RELATIVISTIC ELECTRONS MOVING THROUGH AN UNBOUNDED CRYSTAL

It is well known that the formulae describing the forward PXR from the absorbing crystal with a finite thickness are rather complicated even without regard to bremsstrahlung contribution [11, 12], Since these formulae are not convenient for our purposes we derive the simple formula for the spectral-angular distribution of the intensity of photon flux emitted from relativistic electrons moving in an unbounded non-absorbing crystal.

Based on Maxwell equations for Fourier-transforms of an electrical field $\mathbf{E}_{\omega\mathbf{k}} = (2\pi)^{-4}\int dtd^3r\mathbf{E}(\mathbf{r},t)e^{i\omega t - i\mathbf{kr}}$ excited by a relativistic electron in a crystal with dielectric susceptibility $\chi(\omega,\mathbf{r}) = \chi_0(\omega) + \sum_{\mathbf{g}} \chi_{\mathbf{g}} e^{i\mathbf{gr}}$

$$(k^2 - \omega^2(1+\chi_0))\mathbf{E}_{\omega\mathbf{k}} - \mathbf{k}(\mathbf{k}\mathbf{E}_{\omega\mathbf{k}}) - -\omega^2 \sum_{\mathbf{g}} \chi_{-\mathbf{g}}(\omega)\mathbf{E}_{\omega\mathbf{k}+\mathbf{g}} = 4\pi i\omega \mathbf{J}_{\omega\mathbf{k}} \quad (1)$$

One can obtain within the frame of two-wave approximation of the dynamical diffraction theory [10] the following expression for the field $\mathbf{E}_{\omega\mathbf{k}} \approx \mathbf{E}^{tr}_{\omega\mathbf{k}} = \sum_{\lambda=1}^{2} \mathbf{e}_{\lambda\mathbf{k}} E_{\lambda 0}$ ($\mathbf{e}_{\lambda 0}$ are the polarization vectors, $\mathbf{k}\mathbf{e}_{\lambda 0} = 0$):

$$\mathbf{E}_{\lambda 0} = 4\pi i\omega \frac{\mathbf{e}_{\lambda\mathbf{k}}\mathbf{J}_{\omega\mathbf{k}}}{k^2 - \omega^2\varepsilon_{eff}(\omega,\mathbf{k})}, \quad (2.1)$$

$$\varepsilon_{eff} = 1 + \chi_0 + \frac{\omega^2\chi_g\chi_{-g}\alpha_\lambda^2}{(\mathbf{k}+\mathbf{g})^2 - \omega^2(1+\chi_0)}, \quad (2.2)$$

where $\alpha_1 = 1, \alpha_2 = (\mathbf{k},\mathbf{k}+\mathbf{y})/k\,|\mathbf{k}+\mathbf{g}|$.

It is of interest to note [12] that the expression (2.1) coincides formally with that describing the Cherenkov radiation from a fast charged particle moving with a constant velocity trough an uniform medium. This property establishes the quasicherenkov nature of the forward PXR.

In order to describe the spectral-angular distribution of emitted photons one should calculate the Fourier-integral $E^{Rad}_{\lambda} = \int d^3k e^{i\mathbf{kn}r} E_{\lambda 0}$ where $\mathbf{n}$ is the unit vector to the direction of emitted photon propagation. The result of integration obtained by the stationary phase method can be presented in the form

$$\mathbf{E}^{Rad}_{\lambda} = \sum A^{(\pm)}_{\lambda\mathbf{n}} \cdot \frac{e^{ik_\pm r}}{2},$$
$$A^{(\pm)}_{\lambda\mathbf{n}} = 4\pi^3 i\omega\left(1 \pm \frac{\Delta}{\chi_\lambda}\right)\mathbf{e}_{\lambda\mathbf{k}_\pm}\mathbf{J}_{\omega\mathbf{k}_\pm},$$
$$\mathbf{k}_\pm = (\omega + \xi_\pm)\mathbf{n},$$
$$\Delta = \frac{g^2}{2\omega} + \mathbf{ng}(1 + \frac{1}{2}\chi_0),$$
$$\xi_\pm = \frac{1}{2(1+\mathbf{ng}/\omega)}(-\Delta \pm \chi_\lambda) + \frac{\omega}{2}\chi_0,$$
$$\chi_\lambda = \sqrt{\Delta^2 + \omega^2\chi_g\chi_{-g}\alpha_\lambda^2(1+\mathbf{ng}/\omega)} \quad (3)$$

Obtained formulas allow to determine the spectral-angular distribution of the number of emitted photons

$$\omega\frac{dN}{d\omega d\Omega} = \frac{e^2\omega^2}{16\pi^2}\left\langle\left(1+\frac{\Delta}{\chi_\lambda}\right)^2\left|\int dt\mathbf{e}_{\lambda\mathbf{k}_+}\mathbf{V}(t)e^{i\omega t - i\mathbf{k}_+\mathbf{r}(t)}\right|^2 + \right.$$
$$\left. + \left(1-\frac{\Delta}{\chi_\lambda}\right)^2\left|\int dt\mathbf{e}_{\lambda\mathbf{k}_-}\mathbf{V}(t)e^{i\omega t - i\mathbf{k}_-\mathbf{r}(t)}\right|^2\right\rangle, \quad (4)$$

where the brackets $\langle\ \rangle$ mean averaging over all possible trajectories of emitting electrons $\mathbf{r}(\mathbf{t})$, $\mathbf{V}(t) = \frac{d}{dt}\mathbf{r}(t)$. It should be noted that the expression (4) is reduced in the case of large values of so-called resonance defect Δ to well known formula [1] describing the spectral-angular distribution of ordinary bremsstrahlung.

For a further analysis it is convenient to introduce the angular variables $\vec{\theta}$ and $\vec{\psi}_t$ in accordance with formulas

$$\mathbf{n} = \mathbf{e}(1 - \frac{1}{2}\theta^2) + \vec{\theta}, \quad \mathbf{e}\vec{\theta} = 0,$$
$$\mathbf{V}(t) = \mathbf{e}(1 - \frac{1}{2}\gamma^{-1} - \frac{1}{2}\psi_t^2) + \vec{\psi}_t, \quad \mathbf{e}\vec{\psi}_t = 0 \quad (5)$$

The angles $\vec{\theta}$ and $\vec{\psi}_t$ are shown in Fig. 1. Using (5) and assuming that the thickness of the target L exceeds substantially the value of emission formation length $l_{coh} \approx 2\gamma^2/\omega$ one can reduce (4) to more simple formula

$$\omega\frac{dN}{dtd\omega d^2\theta} = \frac{e^2}{\gamma\pi^2} Re\left\langle\int_0^\infty d\tau\mathbf{\Omega}_{\lambda t}\mathbf{\Omega}_{\lambda t+\tau}e^{-i\omega\tau} + \right.$$
$$\left. + \left[\left(1+\frac{\Delta}{\chi_\lambda}\right)^2 e^{i\mathbf{k}_+(\mathbf{r}_{t+\tau}-\mathbf{r}_t)} + \left(1-\frac{\Delta}{\chi_\lambda}\right)^2 e^{i\mathbf{k}_-(\mathbf{r}_{t+\tau}-\mathbf{r}_t)}\right]\right\rangle \quad (6)$$

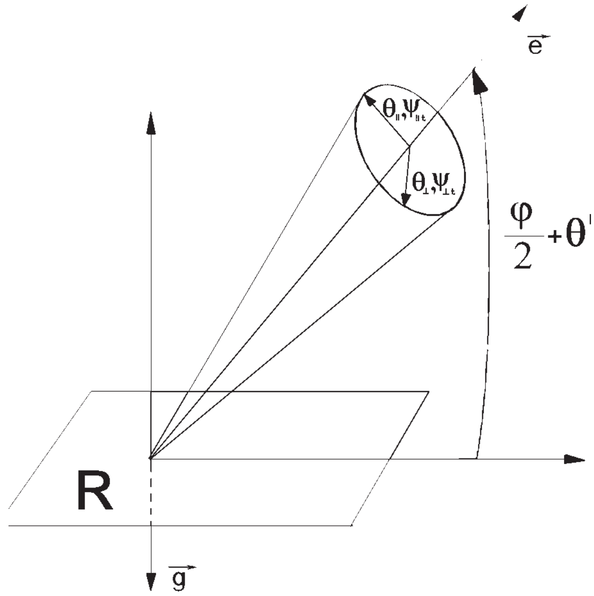

Fig. 1. The geometry of the emission process: R is the reflecting plane, **e** is the axis of the X-ray detector, $\bar{\theta}$ and $\bar{\varphi}_t$ are observation orientation angle, counted from the position of reflecting plane corresponding to exact Bragg resonance

describing the distribution of the emission intensity. Here $\Omega_{1t} = \psi_{\perp t} - \theta_{\perp}$, $\Omega_{2t} = \psi_{\| t} - \theta_{\|}$, $\Omega^2 = \Omega_{+}^2 + \Omega_{L}^2$, $\mathbf{r}_t \equiv \mathbf{r}(t)$.

like that depicted in (6) is described in [1] in detail. Using results [1] one can obtain from (6) the following final formula

$$\omega \frac{dN}{dt d\omega d^2\theta} = - \frac{e^2\omega^2}{\gamma\pi^2} \frac{\Omega_{\lambda t}^2}{\Omega_t^2} Re\int_0^\infty d\tau \exp\left[-\sqrt{\frac{i\omega\gamma^2 L_{Sc}}{2}} \Omega_t^2 \sqrt{\frac{i\omega}{2\gamma^2 L_{Sc}}} \tau \times \left(1 + \frac{\Delta}{\chi_\lambda}\right)^2 \nu_{(-)} e^{-\frac{i\omega}{2}\nu_{(-)}\tau} + \left(1 - \frac{\Delta}{\chi_\lambda}\right)^2 \nu_{(+)} e^{-\frac{i\omega}{2}\nu_{(+)}\tau} \right], \tag{7.1}$$

$$\nu_{(\pm)} = \gamma^{-2} + \frac{\omega_0^2}{\omega^2} + \frac{\Delta \pm \chi_\lambda}{\omega \cos\varphi}, \tag{7.2}$$

where $L_{Sc} = \frac{e^2}{4\pi} L_R$, L_R is the radiation length, ω_0 is the plasma frequency of the target ($\chi_0 = -\omega_0^2/\omega^2$). The obtained result (7) provided a basis for the further analysis.

3. THE TOTAL EMISSION YIELD FROM A THICK ABSORBING TARGET

Characteristics of emitted photons on the output of the thick absorbing crystal is of prime interests for the explanation of experimental results [8]. The simple formula

$$\frac{dN_\lambda}{d\omega d^2\theta} = \int_0^L dt \exp\left[-\frac{L-t}{l_{ab}} \right] \int d^2\psi_t \frac{\gamma^2 L_{Sc}}{\pi t} \times \exp\left[-\frac{L_{Sc}}{t} \gamma^2 \psi_t^2 \right] \frac{dN_\lambda}{dt d\omega d^2\theta} \tag{8}$$

is used in this work to solve this task. Here l_{ab} is the photoabsorption length.

It is of first importance that X-ray detector with small acceptance $\theta_d < \gamma^{-1}$ was used in experiment [8]. In addition to this the absorption length l_{ab} was substantially less than the thickness of the target L and the characteristic angular scale of PXR angular distribution $\theta_{eff} \approx \sqrt{\gamma^{-2} + \omega_0^2/\omega^2}$ was small relative to multiple scattering angle $\psi_{Sc} \approx \gamma^{-1}\sqrt{L/L_{Sc}}$. In conditions under consideration the integration in (8) can be performed easy. The result can be presented in the form

$$\omega \frac{dN_\lambda}{d\omega d^2\theta} = \frac{e^2}{4\pi^2} \frac{1}{\chi_0} \left[\left(1 - \frac{\Delta}{\lambda}\right)^2 \int_0^\infty \frac{d\tau e^{-\tau}\left(\sin\tau \left(1 + \sqrt{\frac{\omega}{\gamma^2 L_{Sc}}} L \cdot \mathrm{th}\left(\frac{2}{\nu_{(+)}} \frac{\tau}{\sqrt{\omega\gamma^2 L_{Sc}}} \right)\right) - \cos\tau \right)}{1 + \left(1 + \sqrt{\frac{\omega}{\gamma^2 L_{Sc}}} L \cdot \mathrm{th}\left(\frac{2}{\nu_{(+)}} \frac{\tau}{\sqrt{\omega\gamma^2 L_{Sc}}} \right)\right)^2} + \left(1 + \frac{\Delta}{\chi_\lambda}\right)^2 \int_0^\infty \frac{d\tau e^{-\tau}\left(\sin\tau \left(1 + \sqrt{\frac{\omega}{\gamma^2 L_{Sc}}} L \cdot \mathrm{th}\left(\frac{2}{\nu_{(-)}} \frac{\tau}{\sqrt{\omega\gamma^2 L_{Sc}}} \right)\right) - \cos\tau \right)}{1 + \left(1 + \sqrt{\frac{\omega}{\gamma^2 L_{Sc}}} L \cdot \mathrm{th}\left(\frac{2}{\nu_{(-)}} \frac{\tau}{\sqrt{\omega\gamma^2 L_{Sc}}} \right)\right)^2} + \pi \cdot \left(1 + \frac{\Delta}{\chi_\lambda}\right)^2 \frac{\gamma^2 L_{Sc}}{L} |\nu_{(-)}| \cdot \sigma\left(-\nu_{(-)}\right) \left(\exp\left(-\frac{\gamma^2 L_{Sc}}{L} |\nu_{(-)}| \right) + \frac{\exp\left(-\frac{\pi}{2}\sqrt{\omega\gamma^2 L_{Sc}} \left|\nu_{(-)}\right|\right)\left(\cos\left(\frac{\pi}{2}\sqrt{\omega\gamma^2 L_{Sc}} \left|\nu_{(-)}\right|\right) - \exp\left(-\frac{\pi}{2}\sqrt{\omega\gamma^2 L_{Sc}} \left|\nu_{(-)}\right|\right)\right)}{1 - 2\exp\left(-\frac{\pi}{2}\sqrt{\omega\gamma^2 L_{Sc}} \left|\nu_{(-)}\right|\right)\cos\left(\frac{\pi}{2}\sqrt{\omega\gamma^2 L_{Sc}} \left|\nu_{(-)}\right|\right) + \exp\left(-\pi\sqrt{\omega\gamma^2 L_{Sc}} \left|\nu_{(-)}\right|\right)} \right) \right] \tag{10}$$

$$\omega \frac{dN_\lambda}{d\omega d^2\theta} = -\frac{e^2\omega^2 l_{ab}}{16\pi^2} Re\left[\left(1+\frac{\Delta}{\chi_\lambda}\right)^2 \times \right.$$
$$\times \nu_{(-)} \int_0^\infty \frac{d\tau e^{-\frac{i\omega}{2}\nu_{(-)}\tau}}{1+\sqrt{\frac{i\omega}{2\gamma^2 L_{Sc}}}\cdot L\cdot \mathrm{th}\left(\sqrt{\frac{i\omega}{2\gamma^2 L_{Sc}}}\tau\right)} +$$
$$\left. + \left(1-\frac{\Delta}{\chi_\lambda}\right)^2 \cdot \nu_{(+)} \int_0^\infty \frac{d\tau e^{-\frac{i\omega}{2}\nu_{(+)}\tau}}{1+\sqrt{\frac{i\omega}{2\gamma^2 L_{Sc}}}\cdot L\cdot \mathrm{th}\left(\sqrt{\frac{i\omega}{2\gamma^2 L_{Sc}}}\tau\right)}\right]. \quad (9)$$

Formula (9) is not convenient for the analysis. Based on the methods of contour integration one should put this result in a more convenient form

The correlation $l_{ab} = 1/\omega\chi_0^{''}$ has been used in (10), $\chi_0^{''}$ is the imaginary part of dielectric susceptibility $\chi_0(\omega)$.

The formula (10) describes the contribution of the forward PXR and bremsstrahlung modified because of the Bragg diffraction of emitted bremsstrahlung photons. In order to determine the total emission to above yield must be added the contribution of transition radiation. This contribution is calculated on the basis of ordinary transition radiation theory [13]. In conditions analogous to that used above the transition radiation yield is described by the formula

$$\omega\frac{dN^{tr}}{d\omega d^2\theta} = \frac{e^2}{\pi^2}\frac{\gamma^2 L_{Sc}}{L}\left[\left(1+2\frac{\omega^2}{\gamma^2\omega_0^2}+\frac{L_{Sc}}{L}\right)\exp\left(\frac{L_{Sc}}{L}\right)+\right.$$
$$+E_1\left(\frac{L_{Sc}}{L}\right)+\left(1+2\frac{\omega^2}{\gamma^2\omega_0^2}-\frac{L_{Sc}}{L}\left(1+\frac{\gamma^2\omega_0^2}{\omega^2}\right)\right)\times$$
$$\left.\times\exp\left(\frac{L_{Sc}}{L}\left(1+\frac{\gamma^2\omega_0^2}{\omega^2}\right)\right)\cdot E_1\left(\frac{L_{Sc}}{L}\left(1+\frac{\gamma^2\omega_0^2}{\omega^2}\right)\right)-2\right]\approx$$
$$\approx\frac{e^2}{\pi^2}\frac{\gamma^2 L_{Sc}}{L}\cdot\left[\left(1+2\frac{\omega^2}{\gamma^2\omega_0^2}\right)\ln\left(1+\frac{\gamma^2\omega_0^2}{\omega^2}\right)-2\right]. \quad (11)$$

The obtained results (10) and (11) give a total description of the discussed emission from relativistic electrons crossing a thick aligned crystal.

4. DISCUSSION OF THE OBTAINED RESULTS

Before proceeding to properties of the emitted photon flux it should be noted that the spectrum in the vicinity of Bragg resonance $\Delta = 0$ is of prime interest for our purpose to explain experimental results [8].

Obviously, the spectral-angular distribution (11) describing a contribution of transition radiation varies lightly in the frequently range under consideration. Because of this, one may set approximately $\omega \approx \omega_B$ in the formula (11). Bragg frequently $\omega_B = g/\sin(\varphi/2)$ follows from the formula

$$\Delta = g\sin\left(\frac{\varphi}{2}\right)\left[\frac{\omega_B}{\omega}-1-(\theta'+\theta_{||})\mathrm{ctg}\left(\frac{\varphi}{2}\right)+\right.$$
$$\left.+\frac{1}{2}\left(\frac{\omega_0^2}{\omega^2}+(\theta'+\theta_{||})^2\right)\right]\approx\frac{\omega_g^2}{\omega_B}\delta,$$
$$\delta = \frac{g^2}{2\omega_g^2}\left(\frac{\omega_B}{\omega}-1-(\theta'+\theta_{||})\mathrm{ctg}\left(\frac{\varphi}{2}\right)\right),$$

$$\chi_g\chi_{-g} = \omega_g^4/\omega^4 =$$
$$= (\omega_0/\omega)^4 (F(g)/Z)^2 (|\,S(\mathbf{g})\,|/N_0)^2 e^{-g^2u_T^2}, \quad (12)$$

determining the resonance defect $\Delta(\omega)$. Here $F(g)$ is the atom formfactor, $S(\mathbf{g})$ is the structure factor of an elementary cell containing N_0 atoms, Z is the number of electrons in a single atom, u_T is the amplitude of formal vibrations of atom in a crystalline lattice.

Since $g^2 >> \omega_g^2$, the quantity δ in (12) is the fast function on ω, $\theta^{'}$ and $\theta_{||}$. Therefore the frequency ω is approximately constant and equal to ω_B in the formula (10) except for the variable $\delta(\omega)$. Thus, the characteristic of the total emission $\omega dN^{tot}/d\omega d^2\theta$ given by the sum of (15) and (11) are described by the variable $\delta(\omega,\theta^{'},\theta_{||})$ in the vicinity of Bragg frequently.

It should be noted that the quantify $\nu_{(\pm)}$ in (10) can be presented in the form

$$\nu_{(\pm)} = \gamma^{-2}+\frac{\omega_0^2}{\omega_B^2}+$$
$$+\frac{\omega_g^2}{\omega_B^2\cos\varphi}\left(\delta\pm\sqrt{\delta^2+\alpha_\lambda^2\cos\varphi}\right)\equiv\gamma^{-2}-\chi_{eff}^{(\pm)} \quad (13)$$

where χ_{eff} is the effective dielectric susceptibility, so that condition $\nu_{(-)} < 0$ in (10) coincides with that of Cherenkov effect. Therefore last two terms in the formula (10) may be interpreted as the contribution of Cherenkov radiation (it is easily to show that the contribution of last term is small). On the other hand first two terms in (10) contain not only bremsstrahlung contribution, but the contribution of Cherenkov radiation as well (such a contribution is included into the term, proportional to the coefficient $(1+\Delta/\chi_\lambda)^2$). In accordance with (10) and (11) the spectral-angular distribution of total emission outside the vicinity of Bragg frequency $|\delta|\leq 1$, or $|\omega-\omega_B|\leq 2\omega_g^2/g^2 \ll 1$ coincides with that of ordinary bremsstrahlung and transition radiation. This quantity is the background in the experiment [8]. Obviously, relative contributions of bremsstrahlung and transition radiation to background depend strongly on the value of emitted photon energy $\omega\approx\omega_B$. The contribution of bremsstrahlung decreases with decreasing ω because of two reasons: increase of a photoabsorption and suppression due to Ter-Mikaelian effect [1] manifested in the frequency range $\omega\approx\omega_B<\gamma\omega_0$. On the other hand, the contribution of transition radiation increases with decreasing ω in accordance with (11).

Two new effects appear inside the vicinity of Bragg frequency $|\delta|\leq 1$. The strong Cherenkov peak (the forward PXR peak) appears in this range and bremsstrahlung contribution decreases due to the diffraction of bremsstrahlung photons on the crystallographic plane responsible for the forward PXR generation. An important point is that the hole in bremsstrahlung spectrum and the Cherenkov peak appear in the same frequency range. Therefore the form of total emission spectrum is determined by the competition between the discussed effects. The result of this competition depends strongly on the value of Bragg frequency. Indeed, the contribution of bremsstrahlung is suppressed if $\omega_B<\gamma\omega_0$ due to Ter-Mikaelian effect. As a consequence, the Cherenkov peak can manifest more clearly in the range $\omega_B<\gamma\omega_0$.

In order to explain experimental data [8] on the base of obtained theoretical results on should take into account that the orientation dependence of the yield of strongly collimated emission has been measured in [8]. Such yield is described by the formula

$$N(\theta',\omega_B,\omega_*)=\int\limits_{-\frac{1}{2}\Delta\theta_\perp}^{\frac{1}{2}\Delta\theta_\perp} d\theta_\perp \int\limits_{-\frac{1}{2}\Delta\theta_\parallel}^{\frac{1}{2}\Delta\theta_\parallel} d\theta_\parallel \int\limits_{\omega_*-\frac{1}{2}\Delta\omega}^{\omega_*+\frac{1}{2}\Delta\omega} d\omega F(\omega)\frac{dN}{d\omega d^2\theta}, \quad (14)$$

where $\Delta\theta_\perp\cdot\Delta\theta_\parallel$ is the angular acceptance of crystal-diffractometer, $\Delta\omega$ is the frequency range of photons scattered by the crystal-diffractometer, $F(\omega)$ is the reflection coefficient, centered in ω_*. The value of ω_* is determined by the corresponding orientation of crystal-diffractometer relative to the axis of emitted photon beam. The value of ω_* has been chosen to be equal to ω_B for the position of the target corresponding to exact Bragg resonance.

In order to show most prominently peculiarities in the emission spectrum, observed in the experiment [8] the function $N(\theta',\omega_B)$ calculated by (10) and (11) with account of experimental conditions [8] (the spectral-angular distribution of emitted photons have been measured practically in [8] because of small angular size of crystal-diffractometer) is illustrated in Fig. 2—4. Obviously, the presented results of performed calculations allow to explain all peculiarities of the experiment [8] discussed above. Particularly, the presented orientation curves explain one of the most unexpected results observed in the experiment [8] which consists in the strong dependence of the form of observed orientation curve on the value of Bragg frequency ω_B.

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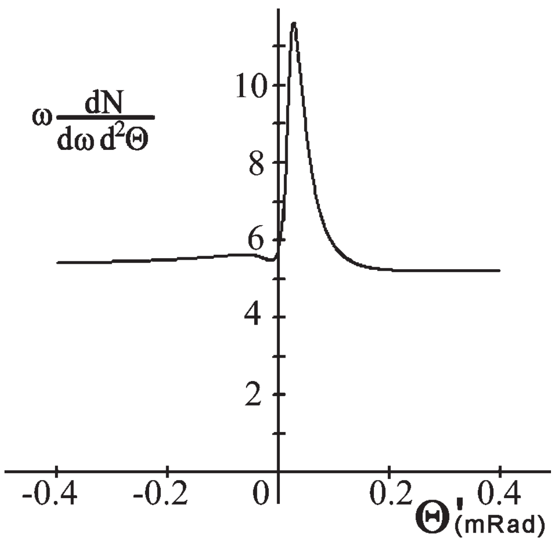

Fig. 2. The orientation dependence of the emission spectral-angular distribution, calculated for $\omega_B=28$ keV

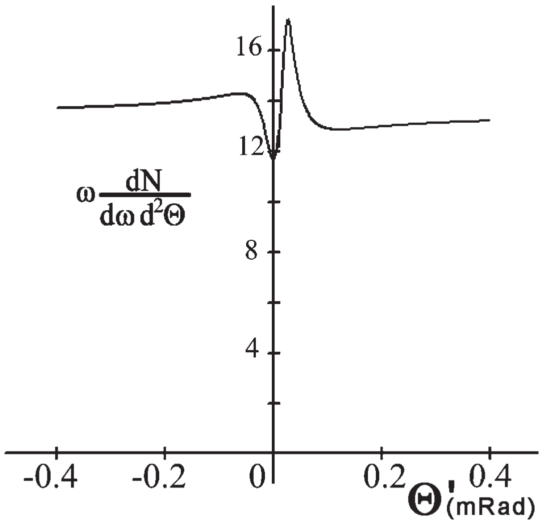

Fig. 3. The same, but for $\omega_B = 40$ keV

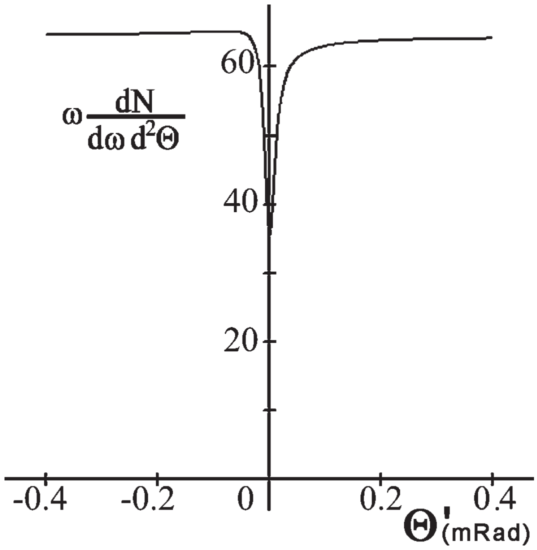

Fig. 4. The same, but for $\omega_B = 95$ keV

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