ИНСТИТУТ ЯДЕРНОЙ ФИЗИКИ СО АН СССР

25

V.N. Baier, V.M. Katkov, V.M. Strakhovenko

TO RADIATION THEORY OF HIGH ENERGY ELECTRONS IN ALIGNED SINGLE CRYSTALS

PREPRINT 86-43

НОВОСИБИРСК

1986

TO RADIATION THEORY OF HIGH ENERGY ELECTRONS IN ALIGNED SINGLE CRYSTALS

V.N.Baier, V.M.Katkov and V.M.Strakhovenko

Institute of Nuclear Physics,
630090, Novosibirsk, U S S R

Abstract

A general approach is developed to a description of the radiation emitted by high energy particles in aligned single crystals. At small angles of incidence, this theory describes the radiation in the field of separate axes and, at large angles of incidence, go over to the coherent bremsstrahlung theory.

In the case when the initial particle is incident at a small angle ϑ_0 to the direction of the axes or the planes of a single crystal, the processes of photon emission by a charged particle and of electron-positron pair creation by a photon are essentially modified compared with amorphous medium. Recently the authors have developed an approach to describe pair creation by photons of ultrahigh energy in single crystals, which is valid at any energies and angles of incidence ϑ_0 /1,2/. This approach is based on a general quasi-classical formalism suited for describing the phenomena occuring in external fields /3/. At small angles of incidence $\vartheta_0 \ll V_0/m$ (V_0 is the scale of the axis (plane) potential and m is the electron mass) the approach describes the pair photoproduction in the field of an axis (plane), and at $\vartheta_0 \gg V_0/m$ it transforms into the theory of coherent pair production. A similar approach is possible to develop for the radiation problem, too. However, in this case these appears a number of complications mainly relating to the range of small ϑ_0 . One of them is associated with the necessity to take into account a redistribution of the flux of incident particles.

Let us introduce the parameters determining the radiation properties*:

$$\varrho = 2\gamma^2\left[\langle \underline{v}^2\rangle - \langle \underline{v}\rangle^2\right] \simeq \frac{2\langle \underline{p}_\perp^2\rangle}{m^2}, \quad \chi = \frac{V_0 \varepsilon}{m^3 a} \qquad (1)$$

where $\gamma = \varepsilon/m$, ε is the energy of a particle, a is the radius of potential action and $\langle \underline{v}\rangle (\langle \underline{v}^2\rangle)$ is the average va-

---

* The system of units $\hbar = c = 1$ is used.

lue of the (squared) velocity of the particle. The parameter ϱ determines the radiation multipolarity; at $\varrho \ll 1$ ($\rho_\perp \ll m$) the radiation is dipole, while at $\varrho \gg 1$ it is of magnetic bremsstrahlung nature. For rather small angles of incidence, ϑ_0, $\varrho \simeq \varrho_c = \frac{2V_0\varepsilon}{m^2}$. The parameter χ determines the quantum properties of the radiation. Note that already at χ = 0.1 the classical intensity is roughly 1.5 times higher than the correct result. Here we will be mainly interested in the energy range where the value of χ is not too low. Then, since $ma \gg 1$ we have $\varrho_c \gg 1$, and at fairly small ϑ_0 the magnetic bremsstrahlung limit holds for the description of the radiation. For the angles of incidence $\vartheta_0 > \vartheta_c \equiv \sqrt{\frac{2U_0}{\varepsilon}}$ (U_0 is the depth of the potential well) we have $\varrho \sim \left(\frac{V_0}{m\vartheta_0}\right)^2$ [1,2]. This means that at $\vartheta_0 \gg V_0/m$ the parameter $\varrho \ll 1$ and the dipole description is valid, which coincides with the coherent bremsstrahlung theory, for this case /4/. It is this range of energies for which we are here developing the approach intended to describe the radiation, which holds at any angles of incidence ϑ_0 of the electrons and positrons on the crystal.

Making use of the formalism of Ref. /3/, we have got a convenient representation for the spectral distribution of radiation intensity on a given trajectory (see Sect. 2 in /5/). This expression needs to be summed up over all of the particle trajectories in the crystal. In what follows we will consider mainly the thin crystals where, by definition, the distribution function over the transverse coordinates of particles is determined by the initial conditions of incidence the crystal by the particles. For definiteness, we will analyse the axial case. For this purpose, the amount of particles

within $d^2\boldsymbol{\varrho}$ is given by the expression /6/

$$dN(\boldsymbol{\varrho}) = N\frac{d^2\varrho}{S}\int \frac{d^2\varrho_0\,\vartheta\left(\varepsilon_\perp(\boldsymbol{\varrho}_0) - U(\boldsymbol{\varrho})\right)}{S\left(\varepsilon_\perp(\boldsymbol{\varrho}_0)\right)} \qquad (2)$$

where N is the total number of particles, S is the area of the transverse cross section per one axis, $U(\boldsymbol{\varrho})$ is the potential dependent on the transverse coordinate $\boldsymbol{\varrho}$, $S(\varepsilon_\perp(\boldsymbol{\varrho}_0)) = \int d^2\varrho\,\vartheta(\varepsilon_\perp(\boldsymbol{\varrho}_0) - U(\boldsymbol{\varrho}))$, $\varepsilon_\perp(\boldsymbol{\varrho}_0) = \frac{\varepsilon\vartheta_0^2}{2} + U(\boldsymbol{\varrho}_0)$, and ϑ_0 is the angle of incidence of the particle (with respect to the axis under consideration). Note that at $\vartheta_0 > \vartheta_c = \sqrt{\frac{2U_0}{\varepsilon}}$ we have $dN(\boldsymbol{\varrho}) = N\frac{d^2\varrho}{S}$.

The crystal potential may be represented as follows:

$$U(\mathbf{r}) = \sum_{\mathbf{q}} G(\mathbf{q})\, e^{-i\mathbf{q}\mathbf{r}} \qquad (3)$$

At $\varrho_c \gg 1$, for any ϑ_0 one can make use of the straight-path approximation (see /1,2/). This is associated with the fact that there exist ϑ_0 such that $\frac{V_0}{m} \gg \vartheta_0 \gg \vartheta_c$, i.e. both the bremsstrahlung description and the straight-path approximation are valid for the particles moving high above the barrier (the approach accuracy is $\sim 1/\varrho_c$). In this approximation let us represent $\mathbf{v}(t) = \mathbf{v}_0 + \Delta\mathbf{v}(t)$ (cf. formula (2.5) in /2/) in this approximation:

$$\Delta\mathbf{v}(t) = -\frac{1}{\varepsilon}\sum_{\mathbf{q}} G(\mathbf{q})\, e^{-i\mathbf{q}\mathbf{r}}\left(e^{-iq_\parallel t} - 1\right)\frac{\mathbf{q}_\perp}{q_\parallel} \qquad (4)$$

where $q_\parallel = \mathbf{q}\mathbf{v}_0$ and $\mathbf{q}_\perp = \mathbf{q} - \mathbf{v}_0(\mathbf{q}\mathbf{v}_0)$. Substituting eqs. (1) and

(4) into the quasiclassical expression for intensity (formulae (2.1)-(2.5) in /5/) we have (cf. eq. (2.8) in /2/):

$$I=\frac{i\alpha m^2}{2\pi\varepsilon^2}\int\omega d\omega\int\frac{d^3r}{V}\int\frac{d^2\boldsymbol{\rho}_0\,\vartheta(\varepsilon_\perp(\boldsymbol{\rho}_0)-U(\boldsymbol{\rho}))}{S(\varepsilon_\perp(\boldsymbol{\rho}_0))}\int_{-\infty}^{+\infty}\frac{d\tau}{\tau-i0}\Bigg\{1-$$

$$-\varphi(\varepsilon)\Big(\sum_{\mathbf{q}}\frac{G(\mathbf{q})\mathbf{q}_\perp}{mq_\parallel}e^{-i\mathbf{q}\mathbf{r}}\sin q_\parallel\tau\Big)^2\Bigg\}\exp\Bigg\{-\frac{im^2\omega\tau}{\varepsilon\varepsilon'}\Big[1+\sum_{\mathbf{q},\mathbf{q}'}\frac{e^{-i(\mathbf{q}+\mathbf{q}')\mathbf{r}}}{m^2q_\parallel q'_\parallel}\cdot$$

$$\cdot G(\mathbf{q})G(\mathbf{q}')(\mathbf{q}_\perp\mathbf{q}'_\perp)\Big(\frac{\sin(q_\parallel+q'_\parallel)\tau}{(q_\parallel+q'_\parallel)\tau}-\frac{\sin q_\parallel\tau}{q_\parallel\tau}\,\frac{\sin q'_\parallel\tau}{q'_\parallel\tau}\Big)\Big]\Bigg\}, \qquad (5)$$

where $\alpha=\left(\frac{e^2}{\hbar c}\right)=\frac{1}{137}$, $\varepsilon'=\varepsilon-\omega$, ω is the photon energy, $\varphi(\varepsilon)=\frac{\varepsilon}{\varepsilon'}+\frac{\varepsilon'}{\varepsilon}$, and V is the crystal volume. Formula (5) describes all the radiation properties including its orientation dependence. Its applicability region coincides, as will be demonstrated below, with that of the quasiclassical approximation. As far the spectral distribution derived from (5) is concerned, if we omit the integration over ω in it, it valid only at $\rho_c\gg1$. A comprehensive analysis of the spectrum for this case is made in Ref. /7/. For the case considered in Ref. [1,2] the radiation and pair production have been considered also in Refs. /8,9/, using the original general expressions of Ref. /3/ and making further calculations only numerically for the particular cases.

Let us choose now the z-axis along the chain of atoms forming the axis. In the case when $\vartheta_0\ll V_0/m$, we have $q_\parallel\tau\sim\sim\frac{\vartheta_0 m}{V_0}\ll1$ for the vectors $\mathbf{q}$ lying in the (x,y) plane and $q_\parallel\tau\sim\frac{m}{V_0}\gg1$ for the rest of $\mathbf{q}$ (with $q_z\neq0$). By virtue of this, one should remain only the terms with $q_z=0$ in the double sum in eq. (5) (the contributions with $q_z\neq0$ are sup-

pressed in the powers of (V_0/m)). Expanding in powers $q_\parallel\tau$ in eq. (5), we find

$$I=\frac{\alpha m^2}{\sqrt{3}\pi\varepsilon^2}\int\omega d\omega\int\frac{d^2\boldsymbol{\rho}}{S}\Bigg\{\int\frac{d^2\boldsymbol{\rho}_0\,\vartheta(\varepsilon_\perp(\boldsymbol{\rho}_0)-U(\boldsymbol{\rho}))}{S(\varepsilon_\perp(\boldsymbol{\rho}_0))}\Big[\varphi(\varepsilon)K_{2/3}(\lambda)-$$

$$-\int_\lambda^\infty K_{1/3}(y)dy\Big]-\frac{1}{3}\frac{(\mathbf{b}(\nabla\mathbf{n})^2\mathbf{b})}{\mathbf{b}^4}\varphi(\varepsilon)\Big[K_{2/3}(\lambda)-\frac{2}{3\lambda}K_{1/3}(\lambda)\Big]+ \qquad (6)$$

$$+\frac{\lambda}{30\,\mathbf{b}^4}\Big[((\mathbf{n}\nabla)\mathbf{b})^2+3(\mathbf{b}(\mathbf{n}\nabla)^2\mathbf{b})\Big]\Big[K_{1/3}(\lambda)-\frac{4}{3\lambda}K_{2/3}(\lambda)+$$

$$+\varphi(\varepsilon)\Big(\frac{4}{\lambda}K_{2/3}(\lambda)-K_{1/3}(\lambda)-\frac{16}{9\lambda^2}K_{1/3}(\lambda)\Big)\Big]\Bigg\}$$

where $\lambda=2m^2\omega/3\varepsilon\varepsilon'|\mathbf{b}|$, K_ν is the McDonald function and $\mathbf{b}=\nabla U(\boldsymbol{\rho})/m$; here the potential $U(\rho)=\sum_{q_z=0}G(q)e^{-i\mathbf{q}\boldsymbol{\rho}}$ is dependent only on the transverse coordinate $\boldsymbol{\rho}$, i.e. $\nabla=\partial/\partial\boldsymbol{\rho}$. The first two terms in formula (6) (which contain no $\mathbf{b}$) is independent on ϑ_0 and, due regard for a redistribution of the flux of charged particles (the coefficient at $[\ldots]$) give the spectral intensity of radiation in the magnetic bremsstrahlung approximation (if the integral over ω is omitted). The remaining terms in eq. (6) $\propto\left(\frac{\vartheta_0 m}{V_0}\right)^2$ and are the corrections of order to this approximation (see formula (3.6) in Ref. /2/). In eq. (6) we have beared in mind the fact that since at $\vartheta_0\sim\vartheta_c$ we have $\left(\frac{\vartheta_0 m}{V_0}\right)^2\sim\frac{1}{\rho_c}$ and this accuracy is ensured by the validity of the initial formula (5), the corrections of the order $\sim1/\rho$ should be taken into account only at $\vartheta_0>\vartheta_c$ when the distribution of the flux does not change.

The expression (6) can be further simplified if the potential of the axis is axisymmetrical and one can perform integration over the angles of the vector $\boldsymbol{\rho}$. Introducing the variable $u=\frac{\omega}{\varepsilon-\omega}$ and integrating by parts in (6), we obtain for the total radiation intensity (cf. eq. (3.7) in /2/):

$$I=\frac{\alpha m^2}{3\sqrt{3}\pi x_0}\int_0^\infty \frac{u\,du}{(1+u)^4}\int dx\left\{\int d^2\varrho_0 \frac{\vartheta(\varepsilon_\perp(\varrho_0)-U(\varrho))}{S(\varepsilon_\perp(\varrho_0))}(4u^2+5u+4)K_{2/3}(\lambda)\right.$$

$$\left.-\left(\frac{\vartheta_0 m}{V_0}\right)^2\frac{3\chi}{10}\frac{\sqrt{x}}{(1+u)^2}K_{1/3}(\lambda)\left[5A(u)+\frac{1+2\eta}{x}B(u)+\frac{1}{2}g(x)C(u)\right]\right\}, \tag{7}$$

where

$$A(u)=4u^3-u^2+14u-1\,,\quad B(u)=-4u^3+27u^2-2u+27,$$

$$C(u)=16u^3-43u^2+38u-43$$

Here we have proceed to the variable $x=\varrho^2/a_s^2$, a_s is the screening radius, $x_0^{-1}=\pi a_s^2 n d$, d is the average distance between the atoms in the chain forming the axis, and n is the density of atoms in the crystal. We define $U(x)=V_0 g(x)$ for the potential of the axis we are using /6/ $g(x)=(\eta+x)^{-1}-(1+\eta+x)^{-1}$ and $\eta\simeq 2u_1^2/a_s^2$, u_1 is the amplitude of thermal vibrations), then

$$\lambda=\frac{u}{3\chi_s\sqrt{x}\,g(x)},\quad \chi_s=\frac{V_0\varepsilon}{m^3 a_s}$$

Just as in the pair creation theory /2/, as χ_s increases the contribution to the radiation comes from the larger region of transverse coordinates whose boundary is determined by the relation $\chi(x)=\frac{2V_0\varepsilon}{m^3a_s}\sqrt{x}\,g(x)=2\chi_s\sqrt{x}\,g(x)\sim 1$. So we have that at $\chi_s\gg 1$ the contribution is given by the region $x\sim\chi_s^{2/3}\gg 1$. On the other hand, at such x the magnetic bremsstrahlung approximation must be applicable, i.e. the condition $\frac{U(x)}{m\vartheta_0}\simeq\frac{V_0}{m\vartheta_0 x}\gg 1$ must be satisfied. Thus, at $\chi_s\gg 1$ the validity region of the bremsstrahlung radiation is limited by the angles of incidence

$V_0/m\chi_s^{2/3}\gtrsim\vartheta_0$. In addition, note that for large χ when the inequality $\chi^{2/3}\ll x_0$ is not satisfied, the axial symmetry can be violated and formula (6) should be used.

Similarly to the pair creation theory /2/, for the angles $\vartheta_0\gtrsim V_0/m$ one can derive, from the general expression (5), approximate formulae for the spectral intensity of radiation whose applicability region is far wider than that of the standard coherent bremsstrahlung theory, at high energies /4/:

$$dI^m_{coh}=\frac{\alpha\omega\,d\omega}{4\varepsilon^2}\sum_{\underline{q}}|G(\underline{q})|^2\frac{\underline{q}_\perp^2}{q_\parallel^2}\left\{\varphi(\varepsilon)-\frac{2\omega m^2}{\varepsilon\varepsilon' q_\parallel^2}\left(|q_\parallel|-\right.\right.$$

$$\left.\left.-\frac{\omega m_*^2}{2\varepsilon\varepsilon'}\right)\right\}\vartheta\left(|q_\parallel|-\frac{\omega m_*^2}{2\varepsilon\varepsilon'}\right) \tag{8}$$

where

$$m_*^2=m^2\left(1+\frac{\varrho}{2}\right),\quad \frac{\varrho}{2}=\sum_{q_\parallel\neq 0}|G(\underline{q})|^2\frac{\underline{q}_\perp^2}{q_\parallel^2 m^2}$$

In the standard theory, $m_*^2\to m^2$. Integrating in (8) over ω we have for the total radiation intensity:

$$I^m_{coh}=\frac{\alpha}{4}\sum_{\underline{q}}|G(\underline{q})|^2\frac{\underline{q}_\perp^2}{q_\parallel^2}F(z),\quad z=\frac{2\varepsilon q_\parallel}{m_*^2} \tag{9}$$

where $$F(z)=\left(\ln(1+z)-\frac{z(2+3z)}{2(1+z)^2}\right)\left(1-\frac{8(3+z)}{z^2(2+\varrho)}\right)+$$

$$+\frac{z}{(1+z)^2}\left(\frac{8}{2+\varrho}-\frac{z(3+2z)}{3(1+z)}\right)$$

Equation (5) has been derived assuming that $\varrho_c\gg 1$. However, for the total radiation intensity it holds in a much broader range and its validity region is the same as that of the quasiclassical approximation. This is due to the fact that at the energies when $\varrho_c \lesssim 1$ and the description of the radiation using the magnetic bremsstrahlung formulae fails, for the angles of incidence $\vartheta_0 \lesssim \vartheta_c$ the classical description of the radiation ($\chi \ll 1$) is valid, the total intensity of which is expressed via local characteristics /6/. In this case, the applicability region of the classical theory is bounded by the angles ϑ_0 when the characteristic energy of the emitted photon is comparable with the electron energy: $\omega \sim \vartheta_0 \gamma^2/a \sim \varepsilon(z-1)$. In this case $\vartheta_0 \sim V_0/m\chi \gg V_0/m$ and the formulae of coherent bremsstrahlung theory (9), directly following from (5), are applicable ($\varrho \ll 1$).

Thus, at $\chi_s \ll 1$ in thin crystals and for the angles of incidence $\vartheta_0 \ll V_0/m\chi_s$, the results of Ref. /6/ are valid, where at $\vartheta_0 = 0$ the electron radiation intensity has been shown to have a maximum and the positron radiation intensity to have a minimum: the former and the latter are associated with a redistribution of the flux of charged particles at $\vartheta_0 \lesssim \vartheta_c$. At $\vartheta_0 \gtrsim \vartheta_c$ the radiation intensity reaches the plateau, while at $\vartheta_0 \sim V_0/m\chi_s$ its decreases, which is described by eq. (9). When $\vartheta_0 \gtrsim \vartheta_c$ the radiation intensity is equal for electrons and positrons. The applicability regions of eqs. (7) and (9) are overlapped at $\chi_s \ll 1$ so that at $\chi_s \ll 1$ the orientation dependence of the radiation intensity is completely described by these simple formulae*. With increasing the parameter χ_s this plateau narrows, and disappears when $\chi_s \sim 1$. At $\chi_s \gtrsim 1$ the intensity of electron radiation is maximum, as previously, at $\vartheta_0 = 0$, and decreases monotonously as ϑ_0 grows. For posi-

* For the case of planar channeling, the orientation dependence was traced in the authors' paper /10/ for this situation.

trons, the intensity also is minimum at $\vartheta_0 = 0$; then with increasing ϑ_0 the intensity grows, achieves its maximum at $\vartheta_0 \simeq \vartheta_c$ and coincides with that of electron radiation when $\vartheta_0 > \vartheta_c$. Fig. 1 demonstrates the typical behaviour of the angular dependence of the electron and positron radiation intensity in a thin crystal, for $\chi_s \sim 1$. Curve I has been obtained according to the formulae of the present paper. It is seen that at $\vartheta_0 < \vartheta_c$ the intensity depends considerably on the kind of particles and on the angular width of the beam $\Delta\vartheta_0$: curves 1 and 3 represent, correspondingly, the positron and electron radiation ($\Delta\vartheta_0 = 0$). Curve 2 shows the case of a uniform distribution with respect to the coordinates and this corresponds to a wide, with respect to the angle, incident beam ($\Delta\vartheta_0 \sim \vartheta_c$). At all these curves coincide at $\vartheta_0 > \vartheta_c$ (curve I). The curves in Fig. 2 have been obtained after calculations on the basis of the derived results. The experimental data are taken from Refs. /11,12/ where the orientation dependence of a relative energy loss $\Delta\varepsilon/\varepsilon_0$ has been measured (Ge crystal: T = 100 K, $\langle 110 \rangle$ axis, initial energy ε_0 = 150 GeV, crystal thickness L = 1.4 mm). The crystal used in the experiment cannot be considered now to be thin because in an appropriate amorphous medium the angle of multiple scattering $\langle \vartheta_c \rangle \simeq 35\,\mu rad \simeq 0.5\,\vartheta_c$. Besides, the angular width of the incident beam was $\Delta\vartheta_0 \simeq \pm 30\,\mu rad$ and the quality of the crystal could correspond to the overall curvature less than 10 μrad (Ref. /11/). One should bear in mind that the multiple scattering of electrons intensifies while the same process for positrons weakens at $\vartheta_0 < \vartheta_c$, what is due to the flux redistribution mentioned above. Since (cf. /7/) for the indicated effective divergence $\Delta\vartheta_0 \sim (30 \div 40)\,\mu$rad most of the electrons populate immediately the above-barrier states and the intensified multiple scattering makes the rest of electrons leave the channel, the electron distribution $dN/d\varrho^2$ has taken uniform at all angles of incidences. On the contrary, for positrons in the first approximation one can neglect multiple scattering and the calculation has been performed for a thin crystal with the effective width of the incident beam taken into account. Note that in thinner crystals and for a smaller angular divergence of the incident beam the orientation dependence of the intensity at $\vartheta_0 < \vartheta_c$ can be observed for electrons as well, while for positrons the effect becomes more "contrast". Further calculations of the energy losses has been performed as in /7/. The solid parts of the curves are constructed according to formulae (7) and (9). The dashed part is the result of interpolation made in the same way as in Ref. /2/.

References

1. V.N.Baier, V.M.Strakhovenko and V.M.Katkov. Phys. Lett. 109A (1985) 179.
2. V.N.Baier, V.M.Katkov and V.M.Strakhovenko, Zh. Eksp. Theor. Fiz. 90 (1986) 801.
3. V.N.Baier, V.M.Katkov and V.S.Fadin. Radiation from Ultrarelativistic Electrons (Atomizdat, Moscow 1973).
4. M.A.Ter-Mikaelian. High-Energy Electromagnetic Processes in Condensed Media, (Wiley-Interscience, New York, 1972).
5. V.N.Baier, V.M.Katkov and V.M.Strakhovenko. Zh. Eksp. Theor. Fiz. 80 (1981) 1348, Soviet Phys. JETP 53 (1981) 688.
6. V.N.Baier, V.M.Katkov and V.M.Strakhovenko. Nucl. Instrum. and Meth. B4 (1984) 346.
7. V.N.Baier, V.M.Katkov and V.M.Strakhovenko. Quantum effects in radiation emitted by ultrahigh energy electrons in aligned crystals. Preprint INP 85-115, Novosibirsk (1985).
8. A.Belkacem, N.Cue and J.C.Kimball. Phys. Lett. 111A (1985) 86.
9. J.C.Kimball, N.Cue and A.Belkacem. 11th International Conference on Atomic Collisions in Solids, Washington (USA) 1985.
10. V.N.Baier, V.M.Katkov and V.M.Strakhovenko "Radiation of Ultrarelativistic Particles at Planar Channeling". Preprint INP 80-03, Novosibirsk 1980.
11. A.Belkacem, G.Bologna, M.Chevallier et al. Phys. Rev. Lett. 54 (1985) 2667.
12. A.Belkacem, G.Bologna, M.Chevallier et al. Preprint CERN-EP/85-150 (1985).

Figure Captions

Fig. 1. The orientation dependence of the radiation intensity in the thin germanium single crystal ($\langle 110 \rangle$ axis, T = = 100 K, $\mathcal{E}$ = 150 GeV). Curve I is obtained according to the formulae of the present paper: according to eq. (9) to the right from the dashed line, to eq. (7) to the left and the dashed part is interpolation; branching of curve I at small angles: 1 - positrons, 3 - - electrons, 2 - uniform distribution. Curve II is calculation according to the standard coherent bremsstrahlung theory.

Fig. 2. The relative energy losses vs. the angle of incidence with respect to the $\langle 110 \rangle$ axis (Ge crystal, T = 100 K, initial energy $\mathcal{E}_0$ = 150 GeV, L = 1.4 mm). The numbers of the curves denotes the same as in Fig. 1. The experimental results were taken from /11,12/; full circles: electrons, open circles: positrons.

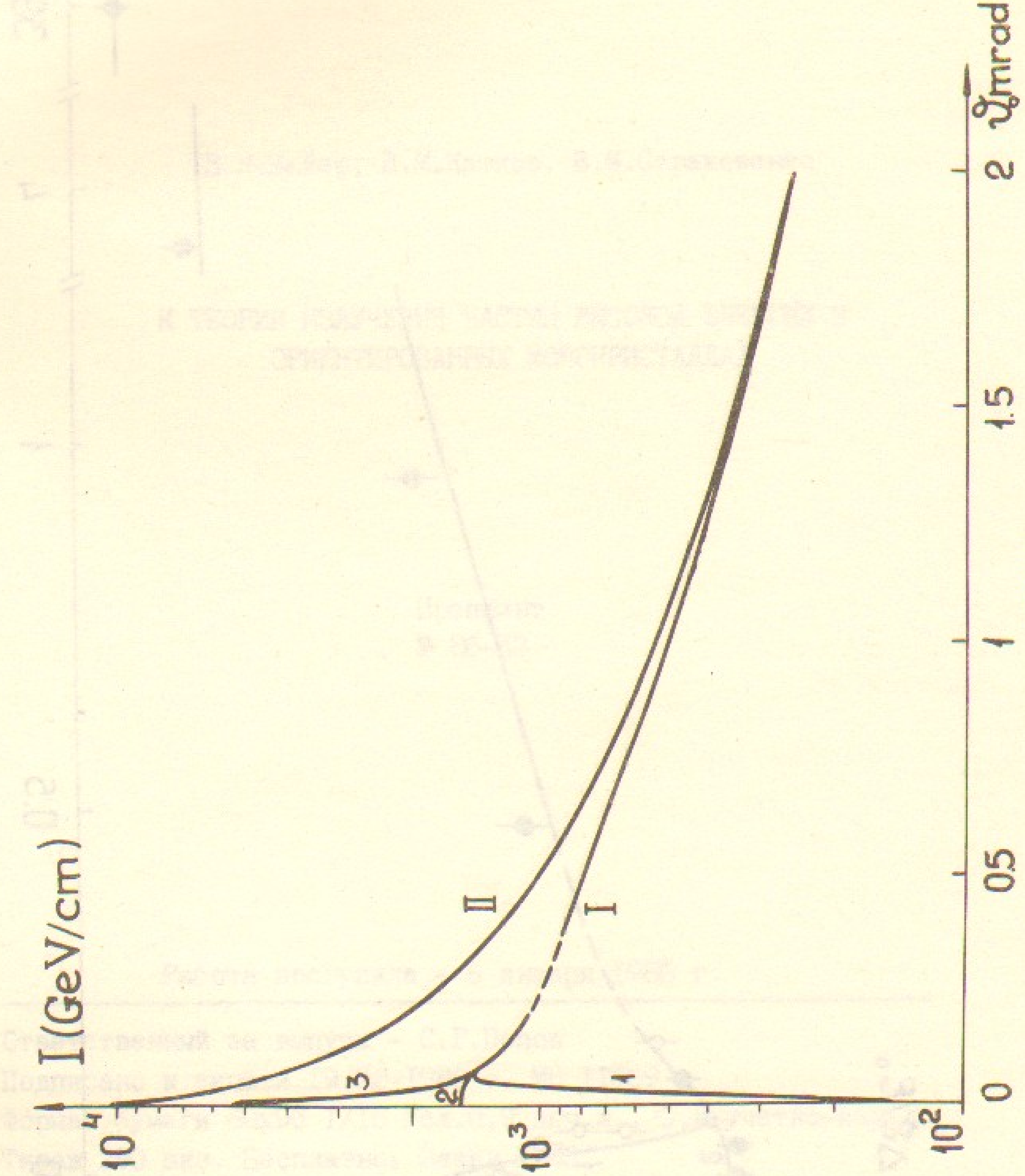

Fig.1

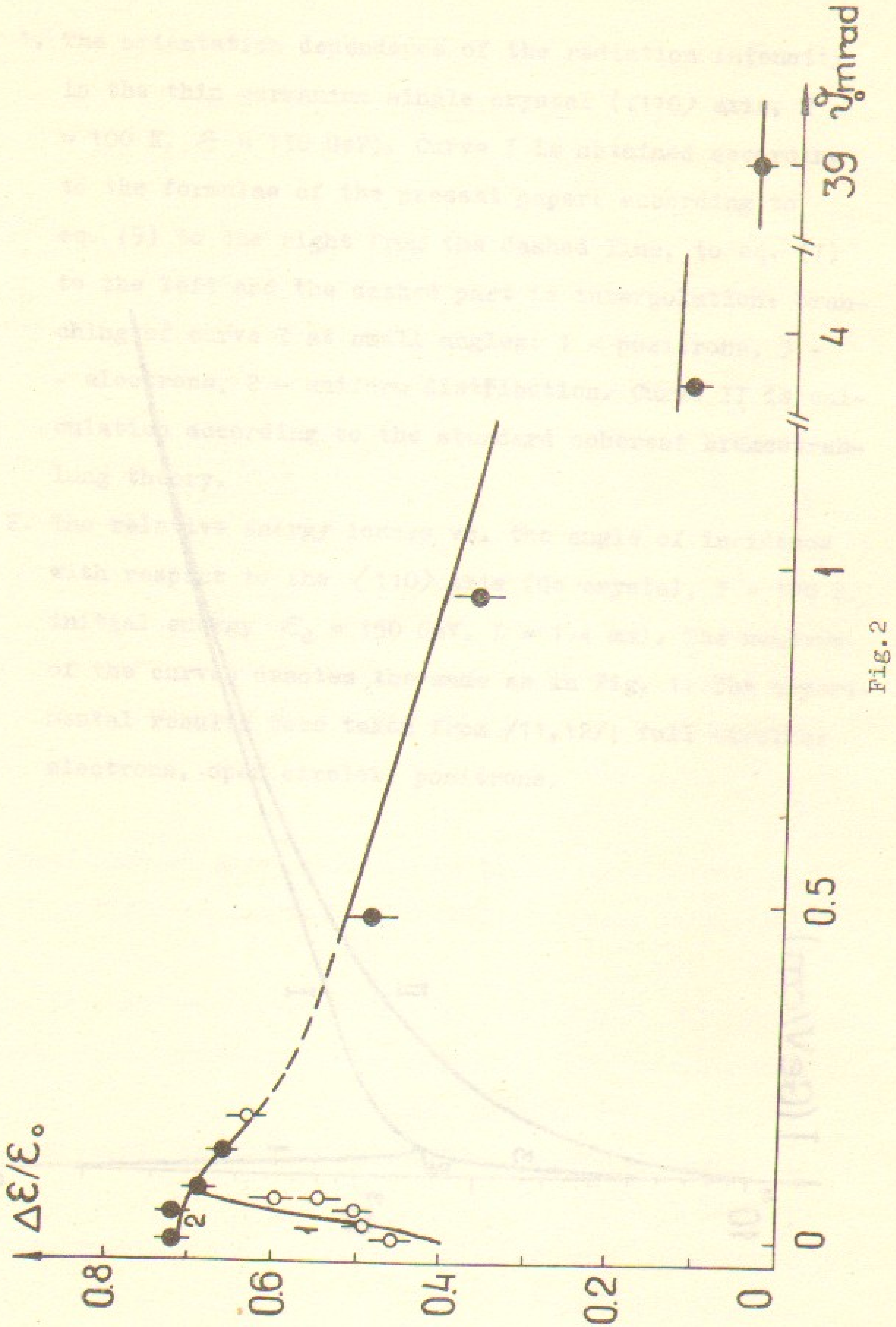

Fig.2

В.Н.Байер, В.М.Катков, В.М.Страховенко

К ТЕОРИИ ИЗЛУЧЕНИЯ ЧАСТИЦ ВЫСОКОЙ ЭНЕРГИИ В ОРИЕНТИРОВАННЫХ МОНОКРИСТАЛЛАХ

Препринт
№ 86-43

Работа поступила - 6 января I986 г.

Ответственный за выпуск - С.Г.Попов
Подписано к печати I9.02-I986 г. МН II669
Формат бумаги 60x90 I/I6 Усл.0,9 печ.л., 0,8 учетно-изд.л.
Тираж 290 экз. Бесплатно. Заказ № 43.

Ротапринт ИЯФ СО АН СССР, г.Новосибирск, 90