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ON IONIZATION LOSS DISTRIBUTION OF RELATIVISTIC PARTICLES IN THIN DETECTORS

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Ionization energy loss distribution of high-energy charged particles in thin detectors is considered. Its modification due to the fact that a part of δ-electrons, kicked out from the atoms by the incident particles, escape from the detector and do not dispose of all their kinetic energy inside it, is estimated and shown to be negligibly small. The physical reasons for this are discussed. The parameters responsible for the magnitude of this modification are obtained.

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INTRODUCTION

The value of a fast charged particle ionization energy loss in a thin target is stochastic. It is distributed according to the function, firstly derived in [1]. This function (ionization loss spectrum) resembles a Gaussian, but is asymmetrical with respect to its maximum, steeply increasing on the left of it and more slowly decreasing on the right. The maximum corresponds to the most probable value E_{MP} of the particle ionization loss in the target. It is smaller than the average ionization loss value defined by the Bethe-Bloch formula [2]. After the seminal paper [1] a series of theoretical and experimental works have been devoted to investigation of ionization loss spectra (ILS) both in amorphous and oriented crystalline targets (see, e. g., [3–5] and refs. therein).

Experimentally, ILS is often registered with the use of thin semiconductor detectors. In this case the detector plays the role of the target itself, and the ionization loss is measured as the value, proportional to the current of electrons kicked out from the atoms by the incident particle and holes, which appear as a result of such ionization. The electron-hole pairs are created not only by the incident particle itself, but also appear in the processes of secondary, tertiary etc. ionization, as a result of energy deposition by δ-electrons kicked out from the atoms by the incident particle. In thin detectors a δ-electron with a relatively low kinetic energy completely loses it inside the detector, while an electron of higher energy can escape the detector transforming just a part of its energy into electron-hole pairs. Due to this fact, with the use of thin detectors one usually measures not the total value of the particle ionization energy loss in the detector, but the resticted one [2]. It is a part of the total ionization loss associated with collisions in which the energy transfer from the particle to the atomic electron does not exceed some maximum value ε_0. It is of the order of the energy at which δ-electrons cease to be absorbed inside the detector (the electron range becomes of the order of the target thickness). For ultrarelativistic incident particles in thin detectors ε_0 is usually much smaller than the kinematically maximum energy $\varepsilon_{\max}$, which can be transferred in a single particle-electron collision. In the theory of average ionization loss this fact is taken into account by restriction of the interval of integration with respect to transferred energies by the value ε_0 instead of $\varepsilon_{\max}$, which noticeably changes both the ionization loss value and its dependence on the particle energy. It is also natural to expect some influence of this restriction on ILS. Indeed, acting in the analogous way as in the case of average loss and removing from the consideration energy transfers exceeding ε_0 results in 'cutting' the ILS tail, which corresponds to large energy transfers. It should automatically lead to some increase of the ILS values in the region of its maximum, since the total area under the ILS curve should remain equal to one (ILS normalization condition). Nevertheless, the influence of the discussed restriction on ILS has not been considered. Even in sufficiently thin detectors ILS is always calculated as if all δ-electrons with the energies up to $\varepsilon_{\max}$, leave all their energy inside the detector. The results obtained thereby well coinside with the results of experimental measurements of ILS, at least with the precision achieved in these measurements. In the present work we theoretically investigate the reason for such a small impact of the discussed restriction on ILS, unlike in the case of average ionization loss. For this aim we approximately estimate the corrections to ILS associated with the fact that high-energy δ-electrons leave the detector losing just a part of their energy inside it and analyse the parameters, responsible for the value of such corrections.

1. ANALYTICAL CONSIDERATION

Let us consider normal incidence of a monoenergetic beam consisting of n high-energy particles on a target (detector) of thickness L (Fig. 1). The quantity $nf(x,\Delta)d\Delta$, where $f(x,\Delta)$ is ionization loss probability distribution function or ILS, defines the number of particles which transmit energy Δ (i. e., energy in the interval from Δ to $\Delta + d\Delta$) to the target after passing the layer of thickness x. The change of this quantity in a thin layer dx is defined by the difference between the number of particles dN_+ which, initially having a smaller energy loss $\Delta - \varepsilon$, lose energy ε inside dx, becoming the particles with the energy loss Δ, and the number of particles dN_- with the energy loss Δ which lose any energy inside dx.

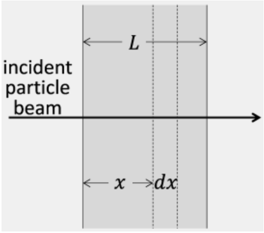

Fig. 1. Passage of a particle beam through the detector

Let us apply a simplified model in which δ-electrons with the kinetic energy $\varepsilon < \varepsilon_0$ do not exit the detector and lose all their energy inside it, while the ones with the energy $\varepsilon > \varepsilon_0$ exit the detector losing just a part of their energy inside it. The value ε' of the energy loss of the latter is defined by the probability distribution function (or ILS), which we denote as $F(\varepsilon, L-x, \varepsilon')$. Generally, it depends on the electron energy ε and the thickness $L-x$ of the layer in front of it. In this case the number dN_+ can be written as

$$dN_+ = ndxd\Delta\left\{\int_0^{\varepsilon_0} d\varepsilon f(x,\Delta-\varepsilon)w(\varepsilon) + \right.$$
$$\left. + \int_{\varepsilon_0}^{\varepsilon_{\max}} d\varepsilon w(\varepsilon)\int_0^{\varepsilon} d\varepsilon' F(\varepsilon, L-x, \varepsilon') f(x, \Delta-\varepsilon')\right\}, \quad (1)$$

where $w(\varepsilon)d\varepsilon dx$ is the probablility for the incident particle to lose energy ε (i. e., in the interval from ε to $\varepsilon + d\varepsilon$) in dx. The layer dx is considered to be thin enough, so that the probability of a single particle collision inside it is much less than unity. The first term in (1) is associated with the particles which create δ-electrons of energy $\varepsilon < \varepsilon_0$, while the second one is due to production of δ-electrons of higher energies, which escape the target. Presently, $nf(x, \Delta-\varepsilon')w(\varepsilon)d\varepsilon dx d\Delta$ stands for the number of incident particles which, having the energy loss $\Delta - \varepsilon'$ before dx, lose energy ε inside this layer, while $F(\varepsilon, L-x, \varepsilon')d\varepsilon'$ is the probability for a δ-electron of energy ε to lose energy ε' before leaving the target. The presence of the second term in (1), as well as integration up to ε_0, instead of $\varepsilon_{\max}$, in the first term, differs the expression for dN_+ in the present case from the expression for this quantity in the case when the δ-electron escape from the target is not taken into account. The expression for dN_-, on the contrary, has the same form as in this case:

$$dN_- = ndxd\Delta\int_0^{\varepsilon_{\max}} d\varepsilon f(x,\Delta)w(\varepsilon)\,, \quad (2)$$

since arbitrary energy loss in dx removes the particles from the company with the loss Δ. It is due to the fact that every energy transfer up to $\varepsilon_{\max}$, results in deposition of some part of the δ-electron energy in the detector and is not remained 'unnoticed' by it. Taking into account (1) and (2) we obtain the following kinetic equation for $f(x, \Delta)$:

$$\frac{df(x,\Delta)}{dx} = \int_0^{\varepsilon_{\max}} [f(x,\Delta-\varepsilon) - f(x,\Delta)]w(\varepsilon)d\varepsilon +$$
$$+\int_{\varepsilon_0}^{\varepsilon_{\max}} d\varepsilon w(\varepsilon)\left[\int_0^{\varepsilon} d\varepsilon' F(\varepsilon, L-x, \varepsilon') f(x,\Delta-\varepsilon') - \right.$$
$$\left. - f(x,\Delta-\varepsilon)\right]. \quad (3)$$

The first line here presents the conventional kinetic equation for $f(x, \Delta)$ without taking into accout the δ-electron escape from the target, while the second and third lines describe the correction associated with it. Equation (3) can be solved with the use of Laplace transform by the procedure analogous to the one applied in [1]. As a result, we obtain:

$$f(L,\Delta) = \frac{1}{2\pi i}\int_{-i\infty}^{+i\infty} dp \exp\{p\Delta -$$
$$-L\int_0^{\varepsilon_{\max}} d\varepsilon w(\varepsilon)(1-e^{-p\varepsilon}) - \quad (4)$$
$$\left. -L\int_{\varepsilon_0}^{\varepsilon_{\max}} d\varepsilon w(\varepsilon)\left[e^{-p\varepsilon} - \int_0^{\varepsilon} d\varepsilon' \overline{F}(\varepsilon, L, \varepsilon')e^{-p\varepsilon'}\right]\right\},$$

where the integration with respect to p is performed along the imaginary axis and

$$\overline{F}(\varepsilon, L, \varepsilon') = L^{-1}\int_0^L F(\varepsilon, L-x, \varepsilon')dx \quad (5)$$

is the ionization loss distribution function for δ-electrons averaged over the target thickness. The first two terms in the exponent in (5) have the same form as in the conventional ILS and can be presented in the form [1]:

$$p\Delta - L\int_0^{\varepsilon_{\max}} d\varepsilon w(\varepsilon)(1-e^{-p\varepsilon}) =$$
$$= p\Delta - p\xi[1 - \Gamma - \ln(p\varepsilon_1)], \quad (6)$$

where

$$\xi = \frac{2\pi Z_1^2 Z_2 N e^4}{mv^2}L,$$
$$\ln\varepsilon_1 = \ln\frac{I^2}{2mv^2\gamma^2} + \frac{v^2}{c^2} + \delta(\gamma),$$

$Z_1 e$ is the incident particle charge, v and γ – are its velocity and Lorenz-factor, $\Gamma \approx 0.577$ is the Euler's constant, Z_2 and N are the atomic number and atomic density of the target, I is its mean ionization potential, m is the electron mass. We also take into account the density effect correction $\delta(\gamma)$ in the definition of $\ln\varepsilon_1$ which is important for high-energy particles [2].

Making the substitution $p = iy/\xi$ in (4) and performing some transformations, it is possible to present this expression in the explicitly real form:

$$f(L,\Delta) = \frac{1}{\pi\xi}\int_0^{+\infty} dy \exp\left\{-y\frac{\pi}{2} - L\int_{\varepsilon_0}^{\varepsilon_{\max}} d\varepsilon w(\varepsilon) \times\right.$$
$$\left.\times\left[\cos(y\varepsilon/\xi) - h_C(y,\varepsilon,L)\right]\right\} \times$$
$$\times \cos\left\{y\lambda + y\ln y + L\int_{\varepsilon_0}^{\varepsilon_{\max}} d\varepsilon w(\varepsilon) \times\right. \quad (7)$$
$$\left.\times\left[\sin(y\varepsilon/\xi) - h_S(y,\varepsilon,L)\right]\right\},$$

where

$$h_C(y,\varepsilon,L) = \int_0^{\varepsilon} d\varepsilon' \overline{F}(\varepsilon, L, \varepsilon')\cos\left(\frac{y\varepsilon'}{\xi}\right), \quad (8)$$

$$h_S(y,\varepsilon,L) = \int_0^{\varepsilon} d\varepsilon' \overline{F}(\varepsilon, L, \varepsilon')\sin(y\varepsilon'/\xi) \quad (9)$$

and $\lambda = \Delta/\xi - \ln(\xi/\varepsilon_1) - 1 + \Gamma$.

Let us analyse the function $\overline{F}(\varepsilon, L, \varepsilon')$ in the integrals (8) and (9). In order to estimate it, it is necessary to know the ionization loss probability distribution function $F(\varepsilon, t, \varepsilon')$ (i. e., ILS) for δ-electrons with $\varepsilon > \varepsilon_0$, since $\varepsilon = \varepsilon_0$ is the lower limit of integration with respect to ε in (7). The function F should be known for various thicknesses t of the layer which the electrons penetrate before leaving the target. For definiteness, we will consider silicon detectors of thickness larger than about several hundred microns, for which at $\varepsilon\sim\varepsilon_0$ δ-electrons are moderately relativistic (e. g., if ε_0 is considered as the energy at which the electron range in the target equals its thickness, for 1 mm silicon target $\varepsilon_0 \approx 740$ keV). In this case, for estmation of $\overline{F}$ we will apply the function F in the form typical for ultrarelativistic electrons, which does not depend on the electron energy ε due to the density effect. This approximation is, certainly, violated for ε in vicinity of ε_0 in the case when the electron has to penetrate almost all the target thickness L, since in this case it is not possible to use the approximation of a constant electron velocity in the target (the one which is applied for calculation of ultrarelativistic particle ILS in thin targets). Though, one the one hand, this approximation rather quickly regains its applicability with the increase of ε, and, on the other hand, a significant part of δ-electrons penetrate through layers thinner than L before leaving the target. The ILS for moderately relativistic electrons with $\gamma\sim 1$ differs from the ILS at $\gamma \gg 1$. Though, for not very thin detectors, which we presently consider (with $\varepsilon_0 \sim mc^2$ or larger) this is not expected to be a source of a noticeable error. For instance, Fig. 2 shows the functions F for δ-electrons in a silicon layer of 200 μm thickness for several values of ε. We see that, although at $\varepsilon \approx 740$ keV (which is the value of ε_0 for $L = 1$ mm) this function is noticeably different from the one typical for ultrarelativistic electrons (for definiteness, their energy is chosen to equal 200 MeV), already at $\varepsilon = 1$ MeV it becomes much closer to it.

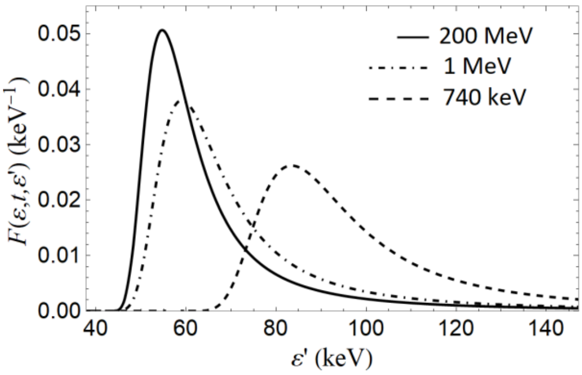

Fig. 2. Ionization energy loss probability distribution functions for electrons of various energies (indicated in the legend) in a 200 μm silicon layer

The function $\overline{F}(\varepsilon, L, \varepsilon')$, obtained with the use of the discussed approximation, is presented in Fig. 3 for the cases of $L = 1$ mm and $L = 300$ μm. Presently we approximately assume that the value $L - x$ defines the path of a δ-electron, emitted in the point x, inside the target, neglecting both the electron scattering and angular distribution. Note that due to the applied approximation for F as a function independent on ε, the function $\overline{F}$ also becomes independent on this quantity and will be further written as $\overline{F}(L, \varepsilon')$. We see that this function is almost flat in the most part of the region of ε', where it is noticeably different from zero. It begins to rapidly decrease at $\varepsilon' \sim E_{MP}$, which is the most probably value of a high-energy electron ionization energy loss in the target of thickness L. At very small ε' the value of $\overline{F}$ is dominated by the contribution from $F(t, \varepsilon')$ in the layers of very small thickness t. In this case F can be noticeably different (e. g., broader) from the conventional Landau distribution and the applied approximation for F fails (thus the curves in this region are dotted). Though, this region makes a small contribution to the integrals in (8) and (9), which allows to neglect this fact.

The analysis of the data on the electron ranges and the values of E_{MP} in thin silicon targets shows that the value of ε_0 in this case is usually several times larger than E_{MP}. Hence, the integration in (8) and (9), which is performed up to $\varepsilon > \varepsilon_0$ can in fact be performed up to infinity, since the integrand vanishes at $\varepsilon' \sim E_{MP}$ due to the discussed behavior of $\overline{F}(L, \varepsilon')$. This makes the functions h_C and h_S independent on ε and we will further write them as $h_C(y, L)$ and $h_S(y, L)$. This allows to calculate the integrals with respect to ε in (7). For this aim we will take the energy loss probability $w(\varepsilon)$ in the form [1]

$$w(\varepsilon) = \frac{\xi}{L\varepsilon^2}, \tag{10}$$

which corresponds to the incident particle Rutherford scattering on free atomic electrons. Evidently, this expression does not take into account the effects of atomic electron binding. These effects are manifested in sufficiently thin detectors (for silicon targets, up to $L\sim 1$ mm), leading to some broadening of the ILS without a noticeable change of the value of E_{MP}. As shown [6], these effects can be taken into account by convolution of the Landau distribution, calculated without taking them into account, with a Gaussian depending on the atomic electron binding energies. Therefore, presently we will not take into account these effects [in the estimation of $\overline{F}(L, \varepsilon')$ as well], concentrating on the considered effect of δ-electron escape from the target.

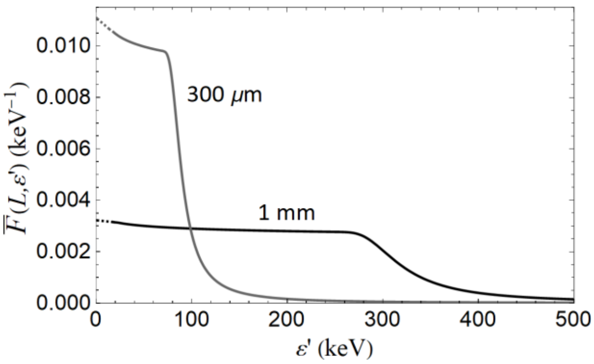

Fig. 3. Averaged ionization energy loss probability distribution functions for electrons in the targets of different thicknesses L (indicated near the curves)

Taking into account the above, the integrals with respect to ε in (7) can be calculated as follows:

$$L \int_{\varepsilon_0}^{\varepsilon_{\max}} d\varepsilon w(\varepsilon)[\cos(y\varepsilon/\xi) - h_C(y, L)] =$$

$$= \frac{\xi}{\varepsilon_0}[\cos(y\varepsilon_0/\xi) - h_C(y, L)] + y[\mathrm{Si}(y\varepsilon_0/\xi) - \pi/2], \tag{11}$$

$$L\int_{\varepsilon_0}^{\varepsilon_{\max}} d\varepsilon w(\varepsilon)[\sin(y\varepsilon/\xi)-h_S(y,L)] =$$
$$=\frac{\xi}{\varepsilon_0}[\sin(y\varepsilon_0/\xi)-h_S(y,L)] - y\mathrm{Ci}(y\varepsilon_0/\xi), \quad (12)$$

where $\mathrm{Si}(x)$ and $\mathrm{Ci}(x)$ are the integral sine and cosine functions defined in such way that at $x\to+\infty$, $\mathrm{Si}(x)\to\pi/2$ and $\mathrm{Ci}(x)\to 0$. Presently we took into account that $y\varepsilon_{\max}/\xi \gg 1$. It is associated with the fact, that the values of y which make the main contribution to the integral in (7) are $y{\sim}1$, while $\varepsilon_{\max}/\xi$ for high-energy incident particles in thin detectors is very large. For instance, for 200 MeV electrons in 1 mm silicon target $\varepsilon_{\max}/\xi{\sim}6000$ and grows with the decrease of the target thickness. Substituting (11) and (12) into (7) we obtain:

$$\begin{aligned} f(L,\Delta) = \frac{1}{\pi\xi}\int_0^{+\infty} & dy\exp\{-y\pi/2 - \\ & -y[\mathrm{Si}(y\varepsilon_0/\xi)-\pi/2] - \\ & -\frac{\xi}{\varepsilon_0}[\cos(y\varepsilon_0/\xi)-h_C(y,L)]\Big\}\times \\ & \times\cos\{y\lambda + y\ln y - \\ & - y\mathrm{Ci}(y\varepsilon_0/\xi)+\frac{\xi}{\varepsilon_0}[\sin(y\varepsilon_0/\xi)-h_S(y,L)]\Big\} \end{aligned} \quad (13)$$

with

$$h_C(y,L)\approx\int_0^{+\infty} d\varepsilon'\overline{F}(L,\varepsilon')\cos\left(\frac{y\varepsilon'}{\xi}\right), \quad (14)$$

$$h_S(y,L)\approx\int_0^{+\infty} d\varepsilon'\overline{F}(L,\varepsilon')\sin(y\varepsilon'/\xi). \quad (15)$$

The function (13) differs from the corresponding function [1], which does not take into account the effect of δ-electron escape from the target, by the presence of the second and third lines in the exponent, as well as by the presence of the second line in the cosine argument.

2. NUMERICAL ESTIMATIONS AND DISCUSSION

The function $\overline{F}(L,\varepsilon')$, as well as the nonaveraged function $F(t,\varepsilon')$, is normalized to unity. Thus, its almost constant value on the 'plato' can be estimated as $\sim 1/E_{MP}$. Substituting this value to the integrals (14) and (15) and changing the upper integration limit to E_{MP}, it is possible estimate the values h_C and h_S. Thereby we get

$$h_C(y,L)\sim\frac{\xi}{yE_{MP}}\sin\left[\frac{yE_{MP}}{\xi}\right], \quad (16)$$

$$h_S(y,L)\sim\frac{\xi}{yE_{MP}}\{1-\cos[yE_{MP}/\xi]\}. \quad (17)$$

The restriction of the integration interval by E_{MP} results in the presence of some additional oscillations in the functions (16) and (17) around the 'average' values (14) and (15), though the amplitude $\xi/(yE_{MP})$ provides a reasonable estimation of h_C and h_S magnitudes.

Generally, the value ε_0 is not a strictly defined one. As noted, it is usually considered as the energy at which the electron range equals the target thickness L. In this case, as mentioned above, for a 1 mm silicon target $\varepsilon_0\approx 740$ keV. Meanwhile, the quantities ξ and E_{MP} are: $\xi\approx 20$ keV (for the electron energy $\varepsilon=\varepsilon_0$, when ξ is maximal), $E_{MP}\approx 300$ keV. For comparison, in a 300 µm target these values are: $\varepsilon_0\approx 315$ keV, $\xi\approx 9$ keV and $E_{MP}\approx 105$ keV. In both cases the parameter ξ/ε_0 is almost the same and is slightly smaller than 0.03. The parameter ξ/E_{MP} is also much smaller than unity for both cases, but slightly changes with L due nonlinear dependence of E_{MP} on L. As expressions (13), (16) and (17) show, these two parameters define the magnitude of the considered impact of δ-electron escape from the target on $f(L,\Delta)$. Indeed, in the exponent in (13) the function $\mathrm{Si}(y\varepsilon_0/\xi)$ is close to $\pi/2$ due to a large value of its argument for $y{\sim}1$, which mainly contribute to the integral. The term in the third line in the exponent is proportional to the small parameter ξ/ε_0, while the function h_C is additionally suppressed being of the order of ξ/E_{MP} (again, for $y{\sim}1$). This makes the contribution from the second and third lines to the value of the exponent small. The analogous considerations indicate the smallness of contribution from the second line in the cosine. These conclusions are supported by Fig. 4 which presents the results of numerical calculation of $f(L,\Delta)$ on the basis of (13) and their comparison with the conventional Landau distributions [1], which neglect the effect of δ-electron escape from the target. The figure shows that for the considered detector thicknesses the function (13) is almost undistinguishable from it.

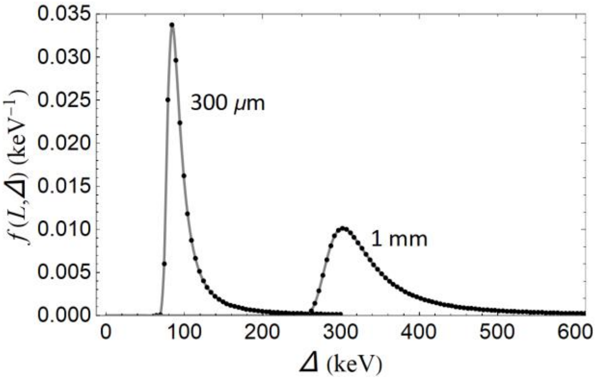

Fig. 4. Ionization loss spectra for high-energy particles in the targets of different thicknesses L (indicated near the curves). Points – calculation on the basis of (13), solid lines – calculation without taking into account the effect of δ-electron escape from the target

The obtained result indicates that, calculating ILS in thin detectors, one should not substitute $\varepsilon_{\max}$, by ε_0 unlike in the case of the average ionization loss calculation. This can be explained by the fact that δ-electrons with $\varepsilon>\varepsilon_0$, though leaving the detector, still lose a considerable amount of energy inside it and do not remain 'unnoticed' by the detector. Since ε_0 is several times larger than E_{MP}, which corresponds to the ILS maximum, such electrons lose energy of the order of E_{MP} (though, electrons emitted from the regions of $x{\sim}L$ lose smaller amounts of energy). Nevertheless, the region in the vicinity of the ILS maximum is mainly formed by a large amount of energy transfers much smaller than ε_0, which are much more numerous than the ones, resulting in the production of δ-electrons with $\varepsilon>\varepsilon_0$. The latter may rather influence on the high-energy tail of ILS, which is though might be rather problematic to register experimentally. For the average ionization loss the situation is different, since every consecutive interval of energy transfers ε in a single collision, on which ε increases by a factor of 10, makes the same contribution to the integral with respect to ε due to the $1/\varepsilon$ dependence of the corresponding integrand [the energy transfer probability (10) scales as $1/\varepsilon^2$, while the transferred energy as ε for $\varepsilon < \varepsilon_0$]. With the increase of the energy transfer above ε_0, the part of the corresponding δ-electron energy, left inside the detector, saturates to a constant [precisely, it is destributed according to the function $\overline{F}(L,\varepsilon')$, which does not depend on ε due to the density effect]. This makes the mentioned integrand to scale as $1/\varepsilon^2$ and the integral over ε to converge at $\varepsilon \sim \varepsilon_0$. Effectively, this corresponds to the substitution $\varepsilon_{\max} \to \varepsilon_0$ of the upper integration limit.

CONCLUSIONS

In the present work we investigated the modification of ILS due to the fact that a certain part of δ-electrons, kicked out from the atoms by the incident particles, escape from the detector and do not dispose of all their kinetic energy inside it. For detectors of thickness exceeding several hundred microns, an approximate estimation of ILS is performed taking into account this fact. It is shown that the corresponding magnitude of ILS modification is defined by the parameters ξ/ε_0 and ξ/E_{MP}, which are usually rather small for high-energy particles in thin detectors. The numerical estimation of ILS for the considered conditions showed that it is almost undistinguishable from the conventional Landau distribution [1]. This indicates that, calculating ILS in thin detectors, one should not substitute $\varepsilon_{\max}$, by ε_0, unlike in the case of the average ionization loss calculation.

The fact that $\varepsilon_{\max}$,should be considered as a maximum energy transferred to the detector in a single collision (though δ-electrons with such an energy take the most part of it out of the detector with them) is of interest, particularly, for numerical simulations of ILS. For instance, the procedure of ILS simulation for particles which move in oriented crystals, applied in [5, 9], demonstrated some sensitivity to the formal substitution $\varepsilon_{\max} \to \varepsilon_0$, which raised the question about the correct value of the energy transfer upper limit and triggered a more detailed investigation of this issue, reported in the present paper.

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ЩОДО РОЗПОДІЛУ ІОНІЗАЦІЙНИХ ВТРАТ РЕЛЯТИВІСТСЬКИХ ЧАСТИНОК У ТОНКИХ ДЕТЕКТОРАХ

С.В. Трофименко

Розглянуто розподіл іонізаційних втрат енергії високоенергетичних заряджених частинок у тонких детекторах. Оцінено його модифікацію внаслідок того, що певна частина δ-електронів, вибитих з атомів налітаючою частинкою, вилітає з детектора і не втрачає в ньому всю свою кінетичну енергію, та показано, що така модифікація є дуже малою. Обговорено фізичні причини цього. Отримано параметри, відповідальні за величину цієї модифікації.