The right-hand side of (18) actually contains G^0, the sum of the first $\sum_{n=0}^{n_0}(2n+1)$ (in the one-dimensional case n_0+1)] spherical harmonics of the Green's functions. Therefore, the obtaining of G is reduced to the obtaining of G^0 by the method of spherical harmonics, that is, to the solution of a small set of integral equations which can be solved with consideration of the dependence of λ and λ_α on u. The Green's function G and the functions $\psi^{(0)}$ and $\psi^{(n)}$ are then determined according to (18) and (16) in finite form.

This is especially important for calculating neutron distributions in media which contain hydrogen (water, petroleum, etc.) since the usual expansion of the distribution in spherical functions[1-6] converges poorly in this case.

In conclusion, I wish to express my sincere thanks to S. A. Kantor for his assistance.

[1] A. I. Akhiezer and I. Ia. Pomeranchuk, *Some Problems of Nuclear Theory*, GITTL, Moscow-Leningrad, 1950.

[2] R. E. Marshak, Rev. Mod. Phys. **19**, 185 (1947).

[3] Marshak, Brooks and Hurwitz, Nucleonics **4**, 5, 6 (1949); **5**, 1, 2 (1949).

[4] M. Verde and G. C. Wick, Phys. Rev. **71**, 852 (1947).

[5] G. C. Wick, Phys. Rev. **75**, 738 (1949).

[6] G. Holte, Ark. f. Fysik **2**, 523 (1951); **3**, 209 (1951); **6**, 243 (1953); **8**, 165 (1954).

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The Effect of Finiteness of Target-Nucleus Mass on Angular Distribution in (d, p) and (d, n) Reactions

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(Submitted to JETP editor June 12, 1956)
J. Exptl. Theoret. Phys. (U.S.S.R.) 31, 895-897
(November, 1956)

In analyzing the angular distribution of the products of (d, p) and (d, n) reactions, it is customary to use the formulas derived by Butler[1] for an infinitely heavy target-nucleus. We have calculated the angular distribution with allowance for the finiteness of the nuclear mass by the same method of smooth connection of functions and subject to the same assumptions as in Ref. 1 without considering the question of the legitimacy of these assumptions. The results of the calculation are given below.

The angular distribution is given by the following formula:

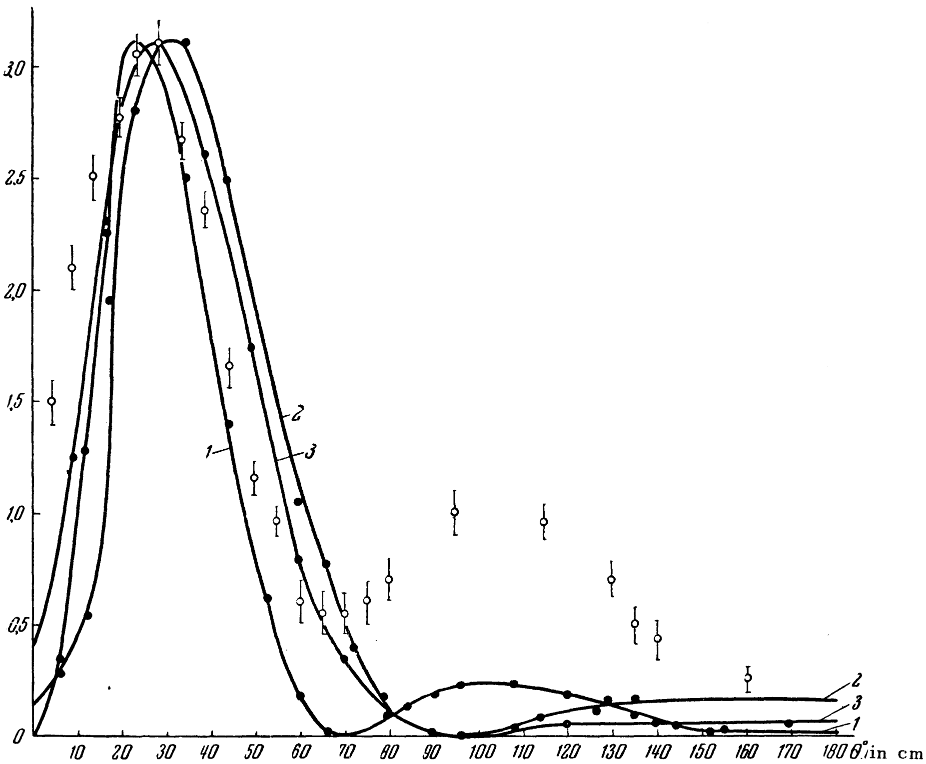
$$S(\theta) = |(Z_1^2+\alpha^2)^{-1} - (Z_1^2+(\alpha+\beta)^2)^{-1}|^2 \times \sum_l N^J_{l k_n^s} \Bigg| \left(\sum_{q=0}^{l} \frac{(l+q)!}{(l-q)!\,q!\,(2k_n^sR_0)^q} \times [l+q+k_n^sR_0]\right) f_l(Z_2R_0) - \left(\sum_{q=0}^{l}\frac{(l+q)!}{(l-q)!\,q!\,(2k_n^sR_0)^q}\right)(Z_2R_0)\, f_{l+1}(Z_2R_0)\Bigg|^2, \quad (1)$$

where l is the orbital angular momentum of a nucleon captured in an orbit of the ultimare nucleus; k_n^s is the wave number of this nucleon in the final state. The values of N are the same as in Ref. 1. R_0 is the "reaction radius" and is obtained by setting in agreement the theoretical and experimental curves for the angular distribution of the products of stripping. The wave numbers α and β which describe the state of the deuteron, are taken to be $\alpha = 0.23\times10^{13}$ cm^{-1} and $\beta = 1.4\times10^{13}$ cm^{-1}; $f_l(x)$ are spherical Bessel functions: $f_l(x) = \sqrt{\pi/2x}\,J_{l+1/2}(x)$. This formula differs from that of Butler only in that K is replaced by Z_1, Z by Z_2 (and r_0 by R_0). Here Z_1 and Z_2 are given by

$$Z_1 = \frac{1}{\hbar}\left|\left(\frac{M_a}{M}\right)^{1/2}\left(\frac{M_b}{M_b+M_c}\right)\sqrt{2(M_b+M_c)W_1}\,\mathbf{n}_1 - \left(\frac{M_a+M_c}{M}\right)^{1/2}\sqrt{2M_bW_2}\,\mathbf{n}_2\right|, \quad (2)$$

$$Z_2 = \frac{1}{\hbar}\left|\left(\frac{M_a}{M}\right)^{1/2}\sqrt{2(M_b+M_c)W_1}\,\mathbf{n}_1 - \left(\frac{M_a^2}{M(M_a+M_c)}\right)^{1/2}\sqrt{2M_bW_2}\,\mathbf{n}_2\right|,$$

where $(\mathbf{n}_1, \mathbf{n}_2) = \cos\theta$, θ is the angle between the directions of the escaping nucleon and the incident deuteron in the center-of-mass system, W_1 and W_2 are the kinetic energy of the deuteron and nucleon with respect to the initial and final nucleus; M_a, M_b and M_c are the masses of the target-nucleus, the



1 and *2* are the angular distributions derived theoretically for finite mass of Be_4^9 with $R_0 = 6.1 \times 10^{-13}$ cm and $R_0 = 4.5 \times 10^{-13}$ cm, respectively; *3* is the proton distribution according to Butler's theory with $r_0 = 4.5 \times 10^{-13}$ cm [2].

emitted nucleon and the captured nucleon. Obviously, $W_1 = E_d M_a / M$; $W_2 = W_1 + Q$. Here E_d is the deuteron energy in the laboratory system and Q is the reaction energy. When $M_a \to \infty$ the formula is reduced to that of Butler.

Allowance for the finiteness of the nuclear mass M_a leads to a change in magnitude of the reaction radius R_0 (which is always an increase). The angular distribution of the reaction products is also somewhat deformed (especially at angles $\theta > 50°$).

We give as an illustration the angular distribution of the group of protons resulting from the Be_4^9 $(d, p)Be_4^{10}$ reaction at $E_d = 3.6$ mev with neutron capture in the ground state of Be_4^{10} ; $l_n = 1$, $Q = 4.59$ mev (see the Figure); the experimental data are taken from Ref. 2. Allowance for the finiteness of M_a leads to larger values of R_0. Moreover, when R_0 is determined from agreement of the first maximum (or first minimum) the positions of the secondary maxima and minima of the distribution curve also agree with experiment.

This is not a unique example. Thus, in the He_2^3 $(d, p)He_2^4$ reaction at $E_d = 10.2$ mev[3] [or $H_1^3(d,n)$ He_2^4] an analogous calculation gives $R_0 = 5.3 \times 10^{-13}$ cm instead of $r_0 = 4.3 \times 10^{-13}$ cm. As in the preceding instance, agreement of the theoretical and experimental first minima leads directly to agreement of the succeeding maxima and minima. It should be noted that with increasing deuteron energy, the correction for finiteness of the mass M_a becomes more important. Therefore, the decrease of the radius r_0 which was noted by Gordon[4] probably resulted in part from the fact that he did not make allowance for the finiteness of M_a.

In conclusion I wish to thank A. Z. Dolginov for suggesting the problem and for his continued interest.

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[1] S. T. Butler, Proc. Roy. Soc. (London) **A208**, 559 (1951).

[2] Fulbright, Bruner, Bromley and Goldman, Phys. Rev. **88**, 700 (1952).

[3] S. T. Butler and J. L. Symonds, Phys. Rev. **83**, 858

(1951).

[4] M. M. Gordon, Phys. Rev. 99, 1625A (1955).

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On the Mechanism of Initiation of Boiling of Liquid Metastable Systems by Ionizing Radiation

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(Submitted to JETP editor July 3, 1956)
J. Exptl. Theoret. Phys. (U.S.S.R.) 31, 897-899
(November, 1956)

A CERTAIN concentration of like-charged ions undoubtedly gives rise to an embryonic micro-burst of a liquid in which an essential role of the electropoles of the ions appears not only in the repulsion of the ions but also in the realization of binding of ions to molecules. In the calcuation of the number of bubbles which are formed as a result of the aggregation of ions in a liquid supersaturated with vapor[1] or gas[2,3], we shall assume as sufficient for the induced creation of a bubble that there originate in a region whose dimensions do not exceed the "interval of localizations" $\lambda(\xi)$ (ξ are parameters which characterize the degree of supersaturation) $\nu(\xi)$ positively charged ions, the electrons of which acquired an energy $\epsilon > \epsilon'_{\rm imp}(\xi)$ upon ionization. (We shall call such ions noncompensated). The parameter $\epsilon'_{\rm imp}$ characterizes the conditions of separation (averaged over all possible configurational dispositions of ν ions) for which the effect of oppositely charged centers becomes unimportant, at least in the time of formation of the embryonic cavity.

The probability $w(\epsilon_0)$ of creation of conditions adequate for the formation of an embryonic cavity somewhere along the track of a δ-electron, which has an initial energy ϵ_0, depends on the complex spatial distribution on noncompensated ions. We can, however, with sufficient confidence, assume that the most numerous δ-electrons with low energies give the chief contribution to the creation of the fluctuations which interest us. This circumstance facilitates the modeling of a distribution of noncompensated ions. Thus, for example, we can hypothesize that a large part of the noncompensated ions is formed near the tracks of low energy δ-electrons and the distributions of such ions is approximated closely by a Poisson distribution, in which the mean specific number of these noncompensated ions is related to the specific energy loss of an electron $\eta \approx (1/\epsilon_{\rm imp})(d\epsilon/dx) \approx a/\epsilon_{\rm imp}\epsilon$, where by $\epsilon_{\rm imp}$ should be understood the energy loss which occurs (on the average) per noncompensated ion formed near the track ($\epsilon_{\rm imp} \gtrsim \epsilon'_{\rm imp}$). Then the probability of creation of the required groups per element of length dx averaged over the track of a δ-electron is

$$dw \approx \eta \frac{(\lambda\eta)^{\nu-1}}{(\nu-1)!} dx \quad \text{for} \quad \lambda\eta \ll 1;$$

$$\text{and} \; w(\varepsilon_0) \approx \int_{\nu\varepsilon_{\rm imp}}^{\varepsilon_0} dw(\varepsilon) + w_1(\varepsilon_0);$$

where $w_1(\epsilon_0)$ is the additional probability of a sought fluctuation at the origin of the track of a δ-electron. Integrating over the spectrum of the δ-electrons, we obtain the specific number of bubbles:

$$n_{\rm elect} \approx \int_{\nu\varepsilon_{\rm imp}}^{\varepsilon_{\max}} w(\varepsilon_0) \frac{KZ^2}{\beta^2} \frac{d\varepsilon_0}{\varepsilon_0^2}$$

$$\approx \frac{KZ^2}{\beta^2} \frac{(a\lambda)^{\nu-1}}{(\nu-1)!\,(\nu\varepsilon_{\rm imp})^{\nu-1}\,\varepsilon_{\rm imp}^{\nu}} \left(\frac{1}{\nu-1} + \frac{1}{\nu^2}\right) \approx \frac{Z^2}{\beta^2} \psi_\nu(\xi),$$

on disregarding the relativistic corrections and the weak dependence on the upper limit (Z and β are the charge and speed of an ionizing particle).

A characteristic of the process of ionic initiation in certain liquids should be noted. Thus, for example, in pure liquids of inert elements in which formation of negative molecular ions does not occur, the energy necessary for sufficient removal of electrons from ionized molecules, which take part in the formation of cavities, must be anomalously large to avoid recombination of ions with high mobility electrons before completion of formation of a cavity. Therefore, the intensity of a track in bubble chambers with inert liquids must increase with addition of certain substances which form negative ions.

In a series of events, a certain apparently relatively small number of bubbles can be formed as a result of thermal spikes generated by nuclear collisions. In calculation of the number of such bubbles