into a μ^+-meson with track length 630 μ, and finally decays into a positron. The whole chain of decays lies in the plane of a single emulsion layer.

Second Case (found by scanner K. A. Abashidze). The incident particle of unknown mass, emitted from a star which has 4 black and 3 relativistic tracks, travels for 5600 μ and decays into a π^--meson with range 353 μ, forms a σ-star consisting of three protons. The decay of the unknown mass particle and the σ-star are in a single emulsion layer.

Third Case (found by scanner L. N. Gabunia). The incident particle of unknown mass, having a range in emulsion of 6500 μ, stops and decays into a π^+-meson, which has a range 354 μ. The π^+-meson in its turn, decays into a μ^+-meson, which after 570 μ decays by the emission of a positron. The entire chain of decays lies in a single emulsion layer.

If we had only a single case to deal with, it would undoubtedly be interpreted as the decay of a τ'-meson, according to the scheme

$$\tau'^{\pm} \rightarrow \pi^{\pm} + 2\pi^0.$$

However, as can be seen from the description of these cases, the common characteristic for all is the occurrence of a π-meson track of length 357 μ $\pm 2\%$. Since these π-mesons are monochromatic, the decay of the $\pm$ particle of unknown mass is with high probability a two particle process.

A two particle decay of an incident meson into a π-meson of 3.4 mev energy (corresponding to 357 μ), is so far unknown.

The gradient of emulsion grains along the tracks of the unknown initial particles and also the nature of their multiple scattering does not allow us to differentiate between the possibilities of a fork caused by the decay of a neutral meson, or a two prong star, or a sudden change of direction of the initial particle in a single scattering event.

Since the exact measurement of initial particle mass by one of the indirect methods was made difficult by the inconvenient placement of the tracks relative to the emulsion, we are at present limited to an examination of various possible decay modes using particles of known mass.

Variation I. The decay scheme of the unknown particle is

$$?^{\pm} \rightarrow \pi^{\pm} + \pi^0 + Q.$$

Then its mass is

$$m_?^{\pm} = 560\ m_e,\ Q = 6.8 \text{ mev}.$$

Variation II. The decay scheme is

$$?^{\pm} \rightarrow \pi^{\pm} + \theta^0 + Q.$$

Then

$$m_?^{\pm} = 1260\ m_e,\ Q = 4.4 \text{ mev}.$$

Variation III. The decay scheme is:

$$?^{\pm} \rightarrow \pi^{\pm} + \nu + Q.$$

Then

$$m_?^{\pm} = 350\ m_e,\quad Q = 33.4 \text{ mev}.$$

Variation IV. A K$\pm$-meson of mass 970 m_e decays with the scheme

$$K^{\pm} \rightarrow \pi^{\pm} + ?^0 + Q.$$

Then $$m_?^0 = 680\ m_e,\ Q = 4.8 \text{ mev}.$$

It is interesting to note that three of the suggested variations give new masses for the initial incident particle, while the fourth gives a new value of mass for the neutral secondary meson. Attention should be directed to the fact that in one of the cases a negative initial particle stopped in the emulsion and was not captured by a nucleus, but decayed into a π^--meson which in turn formed a σ-star.

Indirect measurement of the incident particle mass continues.

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Connection between α-Decay and Nuclear Deformation

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IN this communication a connection between the deformation of the nuclear surface[1] and the relative intensities of α-groups in complex α-spectra of radioactive nuclei will be established; the results calculated apply to the α-spectrum of RdAc.[2] Among those factors which influence the intensity of α-groups should be listed the exponential factor in the well-known Gamow-Bethe formula giving the coefficient for the passage of the α-particle through the potential barrier (penetration factor) and the factor in front of the exponential, expressing the probability of "barrier-less" emission of the α-particle by the radioactive nucleus. We note that the application of the exact Gamow-Bethe formula[3] with a fixed value of the radius R of the nuclear potential well does not give essentially better results than the preceding work of the author[4] obtained with approximate formulas. The role of the factor in front of the exponential is not essential in the questions of interest here with either the one-body model[2] or in the many-body model.[3]

It is of interest here to consider the influence of nuclear deformation on the intensity of α-groups in connection with the work of Ref. 5. In that work a modified formula was obtained for the probability of penetrating the barrier:

$$w = w_0 \exp[-S(E, R)]; \qquad (1)$$

$$S(E, R) = (8Ze^2/\hbar v)\left[\arccos\sqrt{x} - \sqrt{x(1-x)} - {}^4/_5\sqrt{x(1-x)}\,\beta_{\max}\right]; \quad x = E/V.$$

The difference from the usual formula consists in the introduction of the term with the relative deformation β ($\beta_{\max} \approx \Delta R/R$ at the place of the largest compression or extension of the nuclear surface). The remaining notation is: Ze is the electric charge of the daughter nucleus, v and E are the velocity and energy of the α-particle, V is the maximum height of the potential barrier. It is sufficient to allow the perturbation of the daughter nucleus by the emitted α-particle to produce a change in β relative to the ground level of only 10% of the value of β in order to obtain a change in the intensity by a factor of two or three, and this is a very attractive feature. In the single-particle and many-particle models of the nucleus the orders of magnitude of β are completely different. In the former, the α-particle moves in a constant mean field of the remaining nucleons of the nucleus, which produce a potential well with very sharp sides; the changes in the width of the well for comparitively small changes of the α-particle energy ΔE cannot be at all important. In the collective model of the nucleus changes in the deformation β of the required order are entirely conceivable, because with a sufficiently strong coupling between the nucleons and the nuclear surface, the parameters characterizing the potential well will adiabatically follow the changes in state of the nucleus and alpha-particle. In the transition from the ground state to the excited one, a change in β ($\Delta\beta$) will be brought about by the centrifugal force which arises as a result of the rotation of the nucleus, the mechanical moment of which is determined by the resultant moment of all excited rotators I (see Ref. 4). Therefore it is permissable to use the dependence for $\Delta\beta$ obtained in Ref. 6.

$$\Delta\beta = (\eta K/C)\, I(I+1), \qquad (2)$$

where C is the surface coefficient equal to $M_2\omega_2^2 R_c^2$ (M_2 is the effective mass of the vibration of the second vibrational harmonic, ω_2 is the frequency of this vibration and R_c is the nuclear radius) and K is the coefficient of the first-order (in β) coupling of the nucleon with the nuclear surface

$$K = k\sqrt{5/4\pi}\,[3l^2 - j(j+1)]/4j(j+1), \qquad (3)$$

where k is a constant of the coupling with the surface, l is the quantized moment of the α-particle and j is the change in the moment of the nucleon which is outside the shell and which participates in the balance of energy and momentum of the system; η is a numerical coefficient equal —in order of magnitude— to several hundred units.

The size of K in Eq. (2) cannot be immediately identified with the coupling constant k. In fact, in the derivation of Eq. (2) it was necessary to equate the generalized force $C\beta$, which comes from the change of the nuclear surface as a result of deformation, to the generalized force of the centrifugal potential; in view of this the factor $d\Delta R/d\beta$ enters into the right-hand side of Eq. (2). This can be calculated, for example, by writing the expression for the quadrupole moment on one hand, through[7] ΔR and on the other hand, through[8] β , after which it is easy to convince oneself that the above argument is correct. These calculations show that the factor taking into account the influence of nuclear deformation in Eq. (1) does not change the statistical weights of the states with $l=0$ (main group), but considerably increases the statistical weights of the states with $l > 2$ and sharply decreases them in the states with $i < 2$ (under the condition $k > 0$). Choosing , in accordance with Ref. 9, $C=65$ mev, $k = 25$ mev and $R = 10^{-12}$ cm, we find that the calculation of nuclear deformations together with the decay scheme of Ref. 4 significantly improves the agreement with experiment,[10] which can be considered as an indirect confirmation of the collective (generalized) model of the nucleus.

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On the Probabilities of Σ-Particle Decay

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RECENTLY Gatto[1] indicated the interest in applying arguments developed first by Fermi[2] for the photoproduction of π-mesons to the problem of the decay of "strange" particles. As is well known, Fermi showed, on the basis of the requirements of unitarity and symmetry of the S-matrix, that the phases of strongly interacting particles which are produced in weak processes are defined by their mutual scattering. Gatto considered the decay of the Λ-particle from this point of view and came to the conclusion that the restrictions imposed by the unitarity and symmetry of the S-matrix on the ratio of probabilities of the two possible modes of decay of the Λ-hyperon ($\Lambda \to p + \pi^-$ and $\Lambda \to n + \pi^0$) are very weak. This conclusion, as will be evident from the following, depends to a large extent on the fact that the scattering of π-mesons by nucleons at energies ~40 mev is still small and the corresponding phases are small.

We consider in analogous fashion the decay of the Σ-hyperon. In this decay, ~ 115 mev is given off, corresponding to the scattering of π-mesons of ~ 140 mev in the laboratory system by nucleons. It is well known that the scattering of π-mesons by nucleons is already considerable (according to Orear[3] the phases $\alpha_{33} \sim 40°$, $\alpha_3 \sim -10°$, $\alpha_1 \sim 15°$, $\alpha_{31} = \alpha_{13} = \alpha_{11} = 0$; the d-phases do not exceed 5°).

The π-meson-nucleon system formed as a result of Σ-particle decay, has the following values J of total and L of orbital angular momentum, depending on the spin S and parity P of the hyperon:

Spin and parity of the Σ-hyperon		State of the $\pi + N$ system		
S	P	J	L	Phases
1/2	+	1/2	1	α_{31}, α_{11}
1/2	—	1/2	0	α_3, α_1
3/2	+	3/2	1	α_{33}, α_{13}
3/2	—	3/2	2	

(P designates the parity of the Σ-hyperon relative to the nucleon). Thus, L and J of the π-meson-nucleon system are determined unambiguously by the spin and parity of the Σ-hyperon.

We consider two possible decay modes for the Σ-hyperon:

$$\Sigma^+ \to p + \pi^0 \text{ and } \Sigma^+ \to n + \pi^+.$$

If weak interactions are excluded, then the S-matrix has only diagonal elements, transforming Σ^+ into Σ^+ and the π-meson-nucleon system with a given T into a state with the same T. These diagonal elements are equal to 1, $e^{i2\alpha_3}$ and $e^{i2\alpha_1}$ (the second index of the phases α_3 and α_1 is omitted).

If weak interactions are now included, nondiagonal elements arise which transform the Σ-particle into states of the π-meson-nucleon system with $T = 3/2$ and $1/2$. It follows from the unitarity and symmetry of the S-matrix[2] that these diagonal