sible values of the momentum of particle 1.[1,2,3] Ordinarily the value $\overline{m} = m_2 + m_3 + \ldots + m_n$ is taken for m_{eff}, a result which considers the velocities of particles 2, 3, ..., n to be equal in magnitude and direction.[2]

We shall show that the solution for the bounds $p_{1\,max}$, $p_{1\,min}$ of the momentum p_1 of particle 1 can be restricted if one takes into account the angles θ_{ij} between the other charged particles i and j (i, j = 2, 3, ..., n′) and if the lower limit $\tilde{p}_i$ of their momenta p_i is estimated.

An attempt has been made earlier[4] to take into account information about the angles and momenta of the particles for purposes of identification. Unlike reference 4, the present work includes this information directly in m_{eff}. In this way only knowledge of the lower bounds of the momenta is required (in reference 4, knowledge of the values of p_i themselves is required, which is difficult for large p_i and leads to an indeterminacy in the limiting values for p_1).

To deduce a necessary formula, we note that the equation for the momentum p_1 (or energy E_1) of one of the secondary particles having the total energy E and momentum P coincides with the equation for p_1 in the decay of a particle with energy E and momentum P into two particles with masses m_1 and m_{eff}, if we take

$$m_{eff}^2 = (E_2 + \ldots + E_n)^2 - (\mathbf{p}_2 + \ldots + \mathbf{p}_n)^2. \quad (1)$$

It is easy to show that the roots of the equation above for p_1 have the characteristic

$$dp_{1max}/dm_{eff} < 0, \ dp_{1min}/dm_{eff} \geqslant 0.$$

This means that increasing the estimate for m_{eff} shrinks the region of solutions for the value p_1.

To increase this estimate, we write (1) as three positive terms

$$m_{eff}^2 \equiv \sum_{i=2}^{n} m_i^2 + 2\sum_{i<j}^{n} (E_iE_j - p_ip_j) + 2\sum_{i<j}^{n} p_ip_j(1-\cos\vartheta_{ij}). \quad (2)$$

Taking into account $E_iE_j - p_ip_j \geq m_im_j$ and $p_i > \tilde{p}_i$ (where $\tilde{p}_i$ is the lower bound of p_i),† we immediately get the following estimate:

$$m_{eff}^2 \geqslant \tilde{m}^2 \equiv \overline{m}^2 + \Delta^2 \equiv \overline{m}^2 + 2\sum_{2\leqslant i<j}^{n'} \tilde{p}_i\tilde{p}_j(1-\cos\vartheta_{ij}). \quad (3)$$

Here the sum is carried out over all pairs of charged particles, except particle 1. The masses of the neutral particles are included in $\overline{m}$.

Thus, if we take $\tilde{m}$ instead of $\overline{m}$ for the value of m_{eff}, p_{1min} and p_{1max} come closer and closer together as $\tilde{p}_i$ increases and θ_{ij} becomes larger. For narrow beams of secondary particles the use of formula (3) gives no effect.

If for some particles i and j not only $\tilde{p}$ but also p is known, $\tilde{p}$ can be changed to p in the equations and the term $E_iE_j - p_ip_j - m_im_j$ can be added. This makes $p_{1\,max}$ and $p_{1\,min}$ converge even more.

The most probable contribution from neutral particles to m_{eff} can be taken into account by adding to Δ^2 the term $\frac{1}{2}n'(\frac{1}{2}n'-1)\overline{p^2}(1-\cos\overline{\theta})$, where $\overline{p}$ and $\overline{\theta}$ are the average values of p_i and θ_{ij} in the given interaction.

The results of this work are given in more detail in reference 5.

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*We consider high energy stars to be those in which there are tracks of relativistic particles.

†For gray tracks, for example, one can take $\tilde{p}_i = m_i$; for neutral particles, $\tilde{p}_i = 0$ is taken.

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NOTE ON A BARYON SCHEME

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LET us assume that each of the eight known baryons is described by a four-component wave function. The general equation for all the baryons is in this case an equation for a 32-component spinor. A 32-dimensional spinor space can be treated as a representation of a 10-dimensional vector space.[1] The basis vectors of this space are related to ten 32-row matrices Γ^a, where

$$\Gamma^a\Gamma^b + \Gamma^b\Gamma^a = 2\delta_{ab}, \quad a, b = 1, 2, \ldots, 10.$$

There corresponds to a reflection of the basis vectors e^b in the vector space the following transformation of the matrix Γ^a:

$$\Gamma^{a'} = -\Gamma^b\Gamma^a\Gamma^b, \quad a = 1, 2, \ldots, 10,$$

(no summation over b).

We divide the 10-dimensional vector space into a 4-dimensional Minkowski space (the matrices corresponding to the basis vectors are Γ^ν, $\nu = 1, 2, 3, 4$) and a 6-dimensional isotopic space (the matrices here are Γ^a, $a = 5, 6, \ldots, 10$). With the aid of the isotopic space transformations we can determine the following 3-component isotopic vectors, whose components satisfy the same relations as the components of ordinary spin:

$$2\mathbf{J}^1 = (\Gamma^6, \Gamma^5, i\Gamma^5\Gamma^6), \quad 2\mathbf{J}^2 = (\Gamma^8, \Gamma^7, i\Gamma^7\Gamma^8),$$
$$2\mathbf{J}^3 = (\Gamma^{10}, \Gamma^9, i\Gamma^9\Gamma^{10}), \quad 2\mathbf{J}^4 = (i\Gamma^6\Gamma^7, i\Gamma^7\Gamma^5, i\Gamma^5\Gamma^6).$$
$$2\mathbf{J}^5 = (i\Gamma^{10}\Gamma^8, i\Gamma^8\Gamma^9, i\Gamma^9\Gamma^{10}).$$

We assume that the general equation for all baryons is of the type[2]

$$\left\{\Gamma^\nu \partial / \partial x_\nu - k_0 I \exp\left[\Gamma^0 \sum_{k=1}^{5} a_k \mathbf{J}^k\right]\right\}\psi = 0, \qquad (1)$$

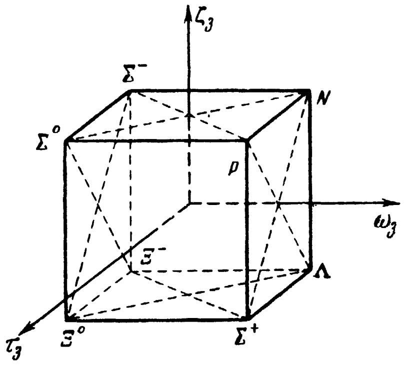
where $\Gamma^0 = \Gamma^1\Gamma^2\Gamma^3\Gamma^4$, $I = i\Gamma^0\Gamma^6\Gamma^7\Gamma^{10}$, and $\mathbf{a}_k$ are some small isotopic vectors whose components depend on the potentials of the meson and lepton fields.

Equation (1) is solved by perturbation theory, where we express the Hamiltonian of the unperturbed problem as

$$H = \hbar c\left\{\Gamma^4\Gamma^k \partial / \partial x_k - k_0 \Gamma^4 I \exp\left[\Gamma^0 \sum_{k=1}^{5} a_k^3 J_3^k\right]\right\}. \qquad (2)$$

The form of the baryon is characterized here by three quantum numbers τ_3, ω_3, and ζ_3, which appear as eigenvalues of the operators $(i/2)\Gamma^0\Gamma^7\Gamma^8$, $(i/2)\Gamma^0\Gamma^9\Gamma^{10}$, and $(i/2)\Gamma^0\Gamma^5\Gamma^6$ (see references 3 and 4). The baryon scheme can be described by a unit cube whose center lies at the origin of the coordinates τ_3, ω_3, and ζ_3 (see diagram).

From (1) and (2) we can determine the perturbation Hamiltonian which describes the weak interaction transitions between baryons. Every baryon can be displaced along the edges of the cube and along the diagonals marked in the diagram by dotted lines. All the indicated processes are possible in principle, although by the energy conservation law some of these, for example $\Xi^- \to \Sigma^- + \pi^0$, cannot be realized. Other proc-



esses, for example, $\Sigma^+ \to \Lambda^0 + \beta^+ + \nu$, are obviously realizable but are seldom found, since competing processes are more probable. The scheme described gives all known baryon transformations and does not give any forbidden transitions. All the weak interaction transitions noted can also proceed from positive energy states to negative energy states and vice versa. This means there can exist processes of the type $N + \overline{P} \to \beta^- + \nu$, whose probability is small compared with that of the strong interaction processes.

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ON THE DECAY SCHEME OF THE Bi^{210} *ISOMER*

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LEVY and Perlman[1] assumed that the α decay of the Bi^{210} isomer with $T_{1/2} = 2.6 \times 10^6$ yr and $E_\alpha = (4935 \pm 20)$ kev goes to the ground state of Tl^{206}. From the energy balance of the α and β transitions $Bi^{210} \to Pb^{206}$ the authors concluded that the long-lived state of Bi^{210} is the ground state and that RaE ($T_{1/2} = 5.01$ days) is metastable with an excitation energy of about 25 kev. We detected[2] the fine structure of the α spec-