also true of ρ and φ. If we apply a similar gauge transformation to all[8] fields ψ, we will have[9]

$$\psi = S\hat{\psi}, \quad D_\mu \psi = SD_\mu \hat{\psi}, \tag{11}$$

where

$$D_\mu \hat{\psi} = (\partial_\mu + ig\hat{X}_\mu)\hat{\psi}.$$

Now, if the radial fields are replaced by their vacuum expectation values, the first term in Eq. (1) becomes, as usual, a mass term for the $\hat{X}_\mu$ fields and the total Lagrangian $\mathcal{L} + \mathcal{L}_m$ can be written as

$$\mathcal{L}_{tot} = -\tfrac{1}{4}\hat{F}_{\mu\nu}{}^a \hat{F}_{\mu\nu}{}^a - \tfrac{1}{2}m_a{}^2 \hat{X}_\mu{}^a \hat{X}_\mu{}^a + \mathcal{L}'(\hat{\psi}, D_\mu \hat{\psi}), \tag{12}$$

with

$$m_a{}^2 = m_0{}^2 + \delta m_a{}^2, \tag{13}$$

where $\delta m_a{}^2$ is the contribution from the first term in (1). In the Lagrangian (12), the gauge fields $X_\mu{}^a$ have been replaced everywhere by the massive vector fields $\hat{X}_\mu{}^a$ and the Goldstone fields (the angular fields θ_a) have completely disappeared.

Before closing, I would like to point out that the polar decomposition of the scalar fields was carried out above only to establish contact with Kibble's work[3] and is by no means essential. One can directly write

$$\partial_\mu \varphi = (\partial_\mu \varphi)(\varphi^\dagger \varphi)^{-1} \varphi^\dagger \varphi \equiv W_\mu \varphi.$$

We again have Eq. (6), but now

$$Y_\mu = X_\mu + (1/ig)W_\mu.$$

It should also be clear that for writing down the gauge-invariant Lagrangian $\mathcal{L}_m$ it is by no means necessary to assume spontaneous breakdown of the symmetry; when the latter does take place, the Higgs-Kibble mechanism provides the mass correction $\delta m_a{}^2$.

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PHYSICAL BASIS OF THE TRANSPORT AND COMPOSITION OF COSMIC RAYS IN THE GALAXY*

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Consideration of cosmic-ray propagation leads naturally to an equilibrium model for cosmic-ray origin and storage in the galaxy. In this model the relative abundances of nuclei with substantially different charge values are nearly energy independent. We conclude that this simple, one-"component" model is consistent with observations of both the galactic primary (i.e., H^1, He^4, C, N, O, Fe) and secondary particles (He^3, Li, Be, B).

Because of ignorance of the parameters describing cosmic-ray transport in the galaxy, various simplifying models have been utilized to interpret "local interstellar" energy spectra and relative abundances of cosmic rays in terms of input, or "source" spectra. For example, during the past decade or so it has been demonstrated that the observed abundance of He^3 and light nuclei (Li, Be, B) at energies greater than a few hundred MeV/nucleon may be understood to be a consequence of propagation of a beam of "primary" particles through 3-4 g/cm² of material.[1] This simple "slab" model, in which the actual physics of particle transport is suppressed, is inconsistent with observations below ~200 MeV/nucleon where ionization energy loss becomes important.[2] Generalization to a distribution of path lengths tends to give better agreement at low energies.[2,3]

In this Letter, we consider this problem from the more direct and physical point of view of cosmic-ray propagation and show that (a) a simple equilibrium model is suggested by present knowledge of the galactic magnetic field and (b) the predictions of this model are consistent with the nearly energy-independent relative cosmic-ray abundances observed. Although the mathematical development is similar to that usually used for exponential distribution of path lengths, the physical content is more clearly specified and analytical solutions are obtained.

The motion of cosmic rays is dominated by the galactic magnetic field which is best regarded as a turbulent or random field.[4,5] As pointed out by Jokipii and Parker[5] the low observed anisotropies of cosmic rays imply that a typical line of force of the random magnetic field in the galactic disk reaches the "surface" of the disk in a distance of 300 pc. The cosmic-ray particles are not free to stream unimpeded out of the galaxy along the magnetic lines of force because the resulting anisotropies would be too large. Two mechanisms for particle confinement appear to be most reasonable, and the actual mechanism is probably some combination of them. Either particles are diffusively confined by resonant scattering due to small-scale fluctuations in the ambient magnetic field along the entire line of force,[6] or particles move freely along the lines of force, having only a small probability of escape from the galaxy each time the line of force reaches the surface of the disk (see discussion in Jokipii and Parker[5]). The latter case includes the possibility of a galactic halo, with particles returning to the disk after propagating in the magnetic field of the halo. In the former case, a diffusion operator with boundary conditions must be used. In general, the two points of view are different,[7] but they are similar enough that the simple "leakage lifetime" approach will be adequate to illustrate the basic phenomena.

Evidence from meteorites indicates that cosmic rays have been present at roughly their present intensity for the past 10^9 yr.[8] Therefore one may assume a steady state with thorough mixing and uniformly distributed sources.

We suppose that $N_i(T)$ represents the steady-state density of cosmic-ray particles of species i. The galaxy is represented by a "leaky box" in which the cosmic rays bounce around; each particle of species i has a probability $1/\tau_i$ of escape from galaxy per unit time, where τ_i is related to the reflection probability at the boundary. We neglect continuous acceleration, but include continuous deceleration due to ionization of interstellar neutral hydrogen. Finally, the source of primary particles with kinetic energy per nucleon T is represented by $Q_i(T)$ and the gain or loss of particles due to nuclear interactions with the interstellar gas is included. The continuity equation then reads[1,9]

$$\frac{\partial N_i}{\partial t} = 0 = -\frac{N_i}{\tau_i} + Q_i(T) + \frac{\partial}{\partial T}(b_i N_i) - n_{\mathrm{H}} c\beta\sigma_i N_i + \sum_{k\neq i}\int_0^\infty n_{\mathrm{H}} c\beta' N_k(T')\sigma_{ki}(T,T')dT', \tag{1}$$

where $b_i = -(dT/dt)_i$ is the rate of energy loss per nucleon due to ionization, β is the ratio of particle velocity to the velocity of light c, σ_i is the cross section for annihilation, and σ_{ki} is the cross section for formation of element i at energy T from element k at energy T' (and velocity $\beta' c$) due to interaction with the interstellar hydrogen at density n_{H}.

It is useful to express Eq. (1) in terms of the differential intensity $j_i(T) = N_i(T)\beta c/4\pi$, which is the actual measured quantity. We also express the lifetime τ_i in terms of grams of neutral hydrogen traversed per square centimeter. Thus, defining the mean path length $\Lambda_i = c\beta M_p n_{\mathrm{H}}\tau_i$, where M_p is the proton mass in grams, setting $q_i = Q_i/4\pi M_p n_{\mathrm{H}}$, and letting $(dT/d\xi)_i = -b_i/c\beta M_p n_{\mathrm{H}}$ be the kinetic energy lost per nucleon due to ionization per g/cm² of interstellar gas, we have from Eq. (1)

$$\frac{\partial}{\partial T}\left[\left(\frac{dT}{d\xi}\right)_i j_i(T)\right] + \left[\frac{1}{\Lambda_i} + \frac{\sigma_i}{M_p}\right] j_i(T) = q_i(T) + \sum_{k\neq i}\frac{1}{M_p}\int_0^\infty j_k(T')\sigma_{ki}(T,T')dT'. \tag{2}$$

For a particle of charge Z_i and mass A_iM_p, $(dT/d\xi)_i = -(Z_i^2/A_i)f(T)$, where $f(T)$ is a tabulated function of particle velocity or kinetic energy per nucleon.[1,10] The desired solution to Eq. (2) may be written immediately as

$$j_i(T) = \frac{A_i}{Z_i^2 f(T)} \int_T^\infty dT' \left[q_i(T') + \sum_{k\neq i} \frac{1}{M_p} \int_0^\infty dT'' j_k(T'') \sigma_{ki}(T',T'') \right] \times \exp\left\{ -\int_T^{T'} dT^* \left[\frac{A_i}{Z_i^2} \frac{1}{f(T^*)} \left(\frac{1}{\Lambda_i} + \frac{\sigma_i}{M_p} \right) \right] \right\}. \quad (3)$$

It is of interest that a similar expression may be obtained by assuming an exponential distribution of path lengths.[11] Clearly if the source spectra $q_i(T)$ are known, this becomes a series of coupled integral equations for the $j_i(T)$. Certain limits and special cases serve to illuminate the content of Eq. (3).

In the high-energy limit, energy loss by ionization is not important and the second term on the left in Eq. (2) dominates. The equilibrium flux then reduces to

$$j_i(T) \simeq \left[q_i(T) + \sum_{k\neq i} \frac{1}{M_p} \int_0^\infty dT' j_k(T') \sigma_{ki}(T,T') \right] \left[\frac{1}{\Lambda_i} + \frac{\sigma_i}{M_p} \right]^{-1} \quad (4)$$

for energies T such that $(dT/d\xi)/T \ll 1/\Lambda_i + \sigma_i/M_p$. At lower energies where energy loss is important the spectrum is more difficult to compute.

First, we consider the case of "primary" nuclei. We take $q_i(T)$ to be a power law in total energy, $q_i(T) = \alpha_i(1+T/T_0)^{-\gamma}$, where T_0 is the proton rest-energy, neglect the relatively small rate of production of these nuclei by nuclear interactions, and assume Λ_i and σ_i to be independent of energy. Setting $\Lambda_i =$ const implies that $\tau_i \propto 1/\beta$ which is reasonable since slower particles reach the boundary of the galaxy less often. Furthermore, σ_i/M_p is found to be generally smaller than $1/\Lambda_i$ and is nearly independent of T. Defining

$$R_i = (Z_i^2/A_i)/(1/\Lambda_i + \sigma_i/M_p), \quad R(T) = \int_0^T dT'/f(T'),$$

and considering nonrelativistic energies $T \ll T_0$, Eq. (3) becomes, correct to lowest order in T/T_0,

$$j_i(T) \simeq \frac{\alpha_i A_i}{Z_i^2 f(T)} \left(1+\frac{T}{T_0}\right)^{-\gamma} \int_T^\infty dT' \exp\left(-\frac{R(T')-R(T)}{R_i} \right). \quad (5)$$

Expressions similar to Eq. (5) are obtained for any source spectrum relatively energy independent or "flat" below ~500 MeV/nucleon. For T between ~1 and ~500 MeV/nucleon, $R(T)$ may be approximated by[10] $R(T) = 8.12\times10^{-4}T^{1.825}$ g/cm². Substituting into Eq. (5) we obtain

$$j_i(T) = 4.06\times10^{-2} \frac{\alpha_i A_i}{Z_i^2} R_i^{0.55} \left(1+\frac{T}{T_0}\right)^{-\gamma} T^{0.825} \exp(x_i)\Gamma(0.55, x_i), \quad (6)$$

where $\Gamma(\alpha,x) = \int_x^\infty t^{\alpha-1}e^{-t}dt$ is the incomplete Γ function[12] and $x_i = T^{1.825}/1230R_i$. From Eq. (6) it is clear that there is a characteristic energy $T_i = 50R_i^{0.55}$ MeV/nucleon, at which the argument of the Γ function, x_i, is unity. For $T \gg T_i$, Eq. (6) may be shown to reduce to the high-energy limit, Eq. (4), with the appropriate $q_i(T)$ and neglect of production by nuclear interactions. For $T \ll T_i$, Eq. (6) reduces to $j_i(T) \simeq 0.066\alpha_iA_iZ_i^{-2}R_i^{0.55}T^{0.825}$. Values of R_i and T_i for elements covering a broad range of A and Z are given Table I. We have used for Λ_i the value of 4 g/cm² and for σ_i have taken as an upper limit $\sigma_i = 50A_i^{2/3}$ mb.[1] Note that T_i varies only moderately.

A striking result is that the intensity of species i relative to j is not strongly dependent on the ener-

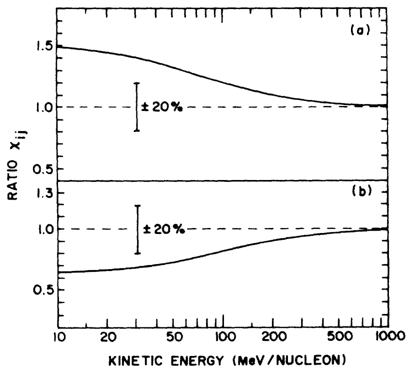

FIG. 1. Energy dependence of the normalized abundance ratio $\chi_{ij}(T)$ for (a) He^4 to CNO and (b) B^{11} to CNO.

gy per nucleon T. To illustrate this, we define a parameter

$$\chi_{ij}(T)=\frac{j_i/j_j}{(j_i/j_j)_{\text{high energy}}}.$$

Thus $\chi_{ij}(T)=1$ at energies $T \gg T_i$ and any departure from unity at lower energies shows the variation in the relative intensity with T. From Eq. (6) we have, for a source spectrum of the form of a power law in total energy,

$$\chi_{ij}(T)=\left(\frac{R_i}{R_j}\right)^{-0.45}\exp(x_i-x_j)\frac{\Gamma(0.55,x_i)}{\Gamma(0.55,x_j)}, \qquad (7)$$

where again $x_i = T^{1.825}/(1230R_i)$. Figure 1(a) shows the computed $\chi_{ij}(T)$ for the normalized ratio of He^4 to medium (CNO) nuclei. Clearly the ratio does not depend strongly on energy. For example, at 50 MeV/nucleon the increase is only ~30 %. Values of $\chi_{He^4 j}$ at 50 MeV/nucleon are given in the last column of Table I for other typical elements. Again note the small variation with energy even for elements with large Z and A values. This relatively small variation in the relative composition is in agreement with observations.[2,3,13]

Next, we consider nuclei such as He^3, Li, Be, and B which are believed to be nearly absent at the source and hence produced primarily by nuclear interactions on the ambient gas. Again we find the abundances of these "daughter" nuclei relative to their "parent" nuclei to be essentially constant as a function of kinetic energy per nucleon T. With the usual approximation[1] that $\sigma_{ki}(T,T')=\sigma_{ki}(T)\delta(T-T')$ for these nuclei, Eq. (3) becomes

$$j_i(T)=\frac{A_i}{Z_i^2 f(T)}\int_T^\infty dT'\sum_{k\neq i} j_k\frac{\sigma_{ki}}{M_p}\exp\left(-\frac{R(T')-R(T)}{R_i}\right). \qquad (8)$$

The high-energy limit is again most readily extracted from the differential equation (2) and one finds immediately[14]

$$j_i(T)=\left(\sum_{k\neq i} j_k(T)\frac{\sigma_{ki}}{M_p}\right)\left(\frac{1}{\Lambda_i}+\frac{\sigma_i}{M_p}\right)^{-1}. \qquad (9)$$

At lower energies the spectrum again depends on the equilibrium spectra of the "parent" nuclei and hence on their source spectra. If, as above, the $q_k(T)$ are power laws in total energy, the $j_k(T)$ are

Table I. Characteristic energies and normalized abundance ratios for typical cosmic-ray nuclei.

Chemical element	A	Z	R_i (g/cm^2)	T_i (MeV/nucleon)	$\chi_{He^4 j}$ at 50 MeV/nucleon
Proton	1	1	3.6	101	1.04
He^3	3	2	4.3	112	1.09
He^4	4	2	3.1	93	1.00
O	16	8	9.0	168	1.36
Fe	56	26	17.5	241	1.69
Pb	208	82	24.9	293	1.92

given in Eq. (6). For nonrelativistic energies, $T \ll T_0$, one finds, upon substitution into Eq. (8) and taking σ_{ki} independent of T,

$$j_i(T) = \frac{A_i}{Z_i^2} \sum_{k \neq i} \frac{\sigma_{ki} j_k(T) R_i R_k}{M_p (R_k - R_i)} \left[1 - \left(\frac{R_i}{R_k}\right)^{0.55} \exp(x_i - x_k) \frac{\Gamma(0.55, x_i)}{\Gamma(0.55, x_k)}\right]. \tag{10}$$

For $T \gg T_i = 50R_i^{0.55}$, Eq. (10) goes over to Eq. (9), as it should. Note that the observed values of the light-to-medium ratio and the He^3/He^4 ratio at high energies are obtained from Eq. (9) using $\Lambda_i \sim 4-5$ g/cm^2 and the appropriate values for σ_{ki} and σ_i.

We now consider the energy dependence of the "daughter-to-parent" ratio. For simplicity consider the case where only one parent nucleus k contributes to the production of $j_i(T)$ (for the production of light nuclei by medium nuclei, k and i may be defined as suitable averages). Then only one term appears on the right in Eq. (10) and, defining $\chi_{ik}(T)$ as before, we have

$$\chi_{ik}(T) = \left[1 - \left(\frac{R_i}{R_k}\right)^{0.55} \exp(x_i - x_k) \frac{\Gamma(0.55, x_i)}{\Gamma(0.55, x_k)}\right] \left[1 - \frac{R_i}{R_k}\right]^{-1}, \tag{11}$$

where $x_i = T^{1.825}/1230R_i$ and $x_k = T^{1.825}/1230R_k$. In Fig. 1(b) is plotted $\chi_{ik}(T)$ for the production of boron by CNO, using $\Lambda_i = \Lambda_k = 4$ g/cm^2, $\sigma_i = 50A_i^{2/3}$ mb, and $\sigma_k = 50A_k^{2/3}$ mb. The production cross section σ_{ki} is assumed nearly independent of T which is justified by experimental measurements.[15] Again the small dependence of this ratio on energy per nucleon is apparent, a result which is in agreement with observations.[3,13]

We note that the above results may be changed by assuming Λ to vary with energy. To assign a specific dependence of Λ on T would, however, require a precise knowledge of cosmic-ray propagation in the galaxy.

Note finally that the relative abundance of "daughter" to "parent" nuclei is completely independent of the lifetime of the "parent" nuclei in the galaxy (or the average amount of matter traversed by the "parent" nuclei), in contrast to the usual interpretation of the "slab" model.

We have shown that a simple equilibrium model of cosmic-ray transport is suggested by our knowledge of the galaxy. The results of this model are consistent with the observations which show relative abundances of nuclei with widely different Z and A values to vary only by small amounts down to energies of ~40 MeV/nucleon. A two-source model is not necessary to fit the observations. These simple calculations also show that elements with very high charge values[16] are expected to be present down to low energies. Future work will extend these results and use more accurate cross-section information.

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of $\chi_{ij}(T)$ on energy T.

[14]For "daughter" nuclei which are radioactive with half-life τ_d, Eq. (9) would read

$$j_i(T) = \frac{\sum j_k(T)\sigma_{ki}\Lambda_i/M_p}{1+\sigma_i\Lambda_i/M_p + \tau_i/(\eta\tau_d)},$$

where $\eta = (1+T/T_0)$.

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HYPERON-NUCLEON INTERACTION AT LOW ENERGIES

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The hyperon-nucleon scattering is described by a potential model. The Σ^+p scattering data require a ${}^3S_1$ resonance in the Λp scattering, for which resonance there seems to be experimental evidence. Its position and the Σ^+p scattering data determine the Σ^+p effective-range parameters, and combined with the Λp scattering data it determines the Λp effective-range parameters.

Recently experiments were done on low-energy hyperon-nucleon scattering, at about 10-MeV Σ laboratory energy for Σp scattering, and at 10- to 40-MeV Λ laboratory energy for Λp scattering.[1-7] They have yielded more precise data on $\Sigma^- p$ and Λp scattering, and they have also given some data on Σ^+p scattering. These data are in themselves insufficient to determine the effective-range parameters for Λp scattering ($T=\frac{1}{2}$) or Σ^+p scattering ($T=\frac{3}{2}$). Using a reasonable theoretical model[8] for the hyperon-nucleon interaction we find that a fit with the combined $\Sigma^\pm p$ scattering data implies the existence of a ${}^3S_1$ Λp resonance below the ΣN threshold.[9] Evidence for such a Λp resonance has been reported.[10] From the position of this resonance and the $\Sigma^\pm p$ scattering data we can determine the Σ^+p effective-range parameters.

We find[11]

$$a_s = -6 \pm 1 \text{ fm}, \quad r_s = 2.1 \pm 0.3 \text{ fm},$$

$$a_t = -0.2 \pm 0.05 \text{ fm}, \quad r_t = -40 \text{ fm}. \tag{1}$$

From the position of the resonance and the Λp scattering data we can determine the Λp effective-range parameters; then we find

$$a_s = -1.7 \pm 0.5 \text{ fm}, \quad r_s = 2.5^{+1.0}_{-0.5} \text{ fm},$$

$$a_t = -1.5 \pm 0.05 \text{ fm}, \quad r_t = 2.0 \pm 0.05 \text{ fm}. \tag{2}$$

The estimates (2) are less reliable than the estimates (1)—as will be explained in the discussion—because of the large energy difference between the resonance and the low-energy Λp scattering. This could explain the difference between the estimates (2) and the scattering length as obtained from hyperfragment calculations.

The model.—It is assumed that the hyperon-nucleon interaction can be described by a meson-exchange potential, which can be inserted in a coupled-channel Schrödinger equation. The potential consists of one-meson—exchange Born terms to which the contribution of two-pion exchange has been added by use of the prescription of Brueckner and Watson. We have taken into account the exchange of the pseudoscalar mesons π, K, η, and X^0, and the vector mesons ρ, K, ω, and φ.[12] Relations between the various coupling constants are imposed using symmetry arguments. It is assumed that the interaction Hamiltonian can be written as

$$H_I = g_p\sqrt{2}(\tfrac{2}{5}\{\bar{B}BP\}_F + \tfrac{3}{5}\{\bar{B}BP\}_D + \tfrac{1}{5}\{\bar{B}BP\}_S)$$
$$+ g_V\sqrt{2}(\{\bar{B}BV\}_S) + f_V\sqrt{2}(\tfrac{2}{5}\{\bar{B}BV\}_F$$
$$+ \tfrac{3}{5}\{\bar{B}BV\}_D + \tfrac{1}{5}\{\bar{B}BV\}_S),$$

where $\bar{B}$ and B describe the baryon octet, P is the nonet of pseudoscalar mesons, for which no η-X^0 mixing is assumed, and V is the nonet of vector mesons, assuming the Okubo Ansatz for the ω-φ mixing. The SU(3) invariants $\{\cdots\}_F$, etc., are defined as usual.[13] The vector-meson-baryon coupling constants g_V and f_V describe the electric and magnetic coupling, respectively. This interaction Hamiltonian is in accordance with relativistic generalizations of SU(6) in lowest order of q^2/m^2.[14] For those coupling constants which are not determined by symmetry