Contemporary status and prospects of high-energy physics[1)]

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A concise review of the most recent major achievements of elementary-particle physics is given. The successes and problems of gauge theories of the strong and electroweak interactions are discussed. A comparison is made of the possible alternatives in the development of physics in the transition to laboratory energies of the order of a tera-electron-volt. Models of grand unification and superunification of the various types of fundamental interactions are considered. A number of examples are used to demonstrate the connection between the properties of elementary particles and the properties of astronomical objects and of the Universe as a whole.

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1. INTRODUCTION

The main difficulty in describing the contemporary status of high-energy physics is perhaps the fact that this science represents a combination of a very large number of very dissimilar fields. The very designation "high-energy physics" is in a sense conventional, since it concerns not only physical experiments using large accelerators or cosmic rays, but also reactor neutrinos, laser beams, mass spectroscopy of the matter around us, tests of its stability, astrophysical observations, and much else.

The theoretical approaches are no less diverse: from the phenomenological parametrization of experimental data to the quest for new mathematical methods for constructing a quantum theory of gravitation.

All these very different forms of experimental and theoretical activity relate to the single problem of establishing the basic physical laws of Nature.

The main difficulty in describing the prospects of fundamental physics is that the development of physics seems logically consistent only in retrospect. But if we consider not "postdictions" but predictions, the next important step is almost always unexpected and very frequently is not taken seriously not only by those with a detached view, but also by those who actually take this step.

If we nevertheless attempt to distinguish the main peculiarity of the contemporary picture of high-energy physics and its main trend, we can say the following. For the first time in the history of physics, a quantum field theory has now been formulated for three of the four basic forces of nature: electromagnetic, weak, and strong. The efforts are directed at the creation of a unified theory of all forces, including a quantum theory of gravitation, and at the construction of a complete picture of the world on the basis of this theory.

2. BASIC PRINCIPLES

The most important achievements of recent decades in elementary-particle physics are connected with symmetries. Symmetries—even broken symmetries—form the heat of contemporary physics.

One of the most profound results of 20th-century physics is the formulation of the principle according to which a symmetry determines not only the kinematics of the fundamental processes, but also their dynamics. More precisely, the dynamics is determined by a local

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[1)]This paper contains an expanded text of a similar report prepared at the invitation of the organizing committee of the Twentieth International Conference of High Energy Physics (Madison, USA, July 1980). Usually, the concluding report at a conference does not contain a list of references. But for a paper the absence of a bibiliography is undoubtedly a defect. To remedy this defect, at least we have listed at the end of the paper some reviews of the topics considered here, which have been published in Uspekhi Fizicheskikh Nauk in recent years.

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symmetry. This principle is not a new one: it is the basis of the classical theory of gravitation which was formulated at the beginning of the century—the general theory of relativity, whose equations are invariant with respect to local coordinate transformations.

In quantum field theory, the principle of local symmetry led to the creation of theories of the electroweak and strong interactions. The equations of these theories are invariant with respect to so-called gauge transformations—local transformations of the phases of the fields. In the case of the electromagnetic interaction, these are Abelian (mutually commuting) transformations of the phases of electrically charged fields, corresponding to the group U(1). In the case of the strong interaction, they are non-Abelian (mutually noncommuting) transformations of the phases of the fields carrying so-called color changes, corresponding to the group SU(3). The unified theory of the electroweak interaction is described by the broken local symmetry U(1) × SU(2). In contrast to the Poincaré group P of coordinate transformations, which is said to be a geometrical group, groups of transformations of the types U(1), SU(2), SU(3), etc., are called internal groups, since they refer to so-called internal (i.e., not space-time) variables such as isotopic spin.

Thus, theories of the electroweak and strong interactions are based on the local internal symmetry $H = \mathrm{U}(1) \times \mathrm{SU}(2) \times \mathrm{SU}(3)$.

It must be emphasized that, despite all the specific features of gauge symmetries and of the mechanisms of their breakdown, the theories that have been created on the basis of these symmetries have not had to go beyond the framework of standard quantum field theory. The theory of the strong interaction and the theory of the electroweak interaction are both described by Lagrangians. Niels Bohr would hardly have counted them among the "crazy" theories about whose necessity he spoke at the end of the fifties. As a legacy of the theory of relativity our theories have acquired the velocity of light c, and as a legacy of quantum mechanics they have acquired the quantum of action $\hbar$. In spite of numerous forecasts, we have so far not detected any discreteness of space and time, and we have so far not supplemented the fundamental constants c and $\hbar$ with a third fundamental constant—an elementary length (or maximal elementary momentum transfer). Moreover, it appears that the conjecture is taken ever more seriously that the natural dimensional unit of high-energy physics is the Planck mass $m_P = (G_N/\hbar c)^{-1/2} \approx 10^{19}\ \mathrm{GeV}/c^2$, a quantity which Planck introduced into physics in the first year of the present century by combining Newton's gravitational constant G_N with $\hbar$ and c. At distances $1/m_P$ (we henceforth employ units in which $\hbar = c = 1$), the gravitational interaction becomes strong and must be appreciably influenced by quantum effects.

The fact that m_P is a natural dimensional unit is suggested by the following argument involving symmetry. The strict (unbroken) local symmetry of nature has the form

$$P \times \mathrm{U}(1)_{\mathrm{em}} \times \mathrm{SU}(3)_{\mathrm{c}},$$

where the group P is related to the gravitational fields, $\mathrm{U}(1)_{\mathrm{em}}$ is related to the electromagnetic field, and SU(3) is related to the gluon field. Thus, the interaction constants of photons $\alpha_{\mathrm{em}} (= e^2/\hbar c)$ and of gluons α_s are dimensionless, and at the fundamental level the only dimensional constant is Newton's constant G_N and, consequently, m_P.

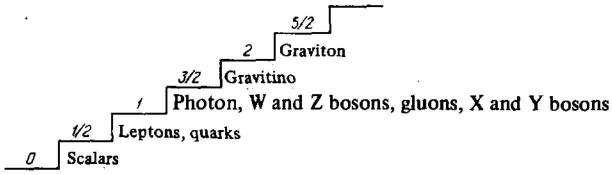
Now that we have characterized the theories of the individual interactions in their general features and have introduced the units, we can attempt to formulate what we expect from the future.

First of all, we hope to find a symmetry group G which includes as subgroups all the currently known internal symmetries, $G \supset H$, and describes all types of gauge interactions by means of a single coupling constant α_G. This is the program of so-called grand unification. The fundamental possibility of grand unification is based on the fact that, because of the phenomenon of polarization of the vacuum, the "charges" of the various gauge interactions—the electromagnetic α_{em}, weak α_W, and strong α_s—depend on distance. For example, the effective electric charge of a particle grows as we penetrate it more deeply, while its strong and weak charges fall off. As a result, charges which differ at large distances may become identical at small distances. Undoubtedly, new types of forces and corresponding gauge fields will be discovered as grand unification is developed. Some of these fields are predicted by the existing grand models; others may appear quite unexpectedly and lead to the construction of more complex schemes.

The second stage, which is much more fundamental, would be the unification of the internal symmetry G and the geometrical symmetry P into a single supersymmetry $S \supset P \times G$, which combines all forces, including gravitation. This is the program of superunification. We expect that superunification would make it possible not only to determine the constant α_G, but also to calculate relations between the various scales which exist in nature. It would establish a connection between the Planck mass m_P and other masses: the proton mass m_p, the electron mass m_e, etc., i.e., it would account for the hierarchy of masses $m_P, \ldots, m_p, \ldots, m_e, \ldots$. In a sense, the explanation of this hierarchy is the final problem of high-energy physics. At the same time, we would explain the whole world. (Of course, this statement should not be taken too literally, but, so to speak, with a grain of salt.)

3. FUNDAMENTAL PARTICLES

Before we turn to a detailed discussion of the individual interactions and problems connected with their unification, we shall give a brief survey of the fundamental particles. We shall discuss what is known about them from experiment, what we would like to know about them, and why they are necessary, i.e., what role they play in nature. The particles can be naturally arranged on a ladder of spins:



These particles have different experimental status. There are no reservations about the existence of the photon γ and the leptons (e, μ, τ, ν_e, ν_μ; the existence of ν_τ has almost been proved). The existence of colored quarks (q) and gluons (g) seems certain for the physicists working with them, but for the general public these particles will still arouse distrust for a long time because they cannot be extracted from the hadrons which are constructed from them. We know five types (flavors) of quarks (q = u, d, s, c, b), each of which exists in three color varieties (yellow, blue, and red). The u and c quarks have electric charge $+2/3$, and the d, s, and b quarks have charge $-1/3$. The gluons—the carriers of the strong interaction—are flavor singlets (they are electrically neutral) and are components of a color octet.

The properties of the intermediate bosons W^+, W^-, and Z^0—the carriers of the weak interaction—are uniquely predicted by the theory of the electroweak interaction. In particular, $m_W \approx 80$ GeV and $m_Z \approx 90$ GeV. We expect that these particles will be discovered during the next two years in the colliding proton–antiproton beams whose construction is nearing completion at CERN.

The situation regarding elementary scalar particles (so-called Higgs bosons) is much more uncertain. These particles play an important role in the theory, being the basis of the mechanism which leads to the generation of the masses of the leptons, quarks, and intermediate bosons. Unfortunately, however, we do not know what the masses of the scalar bosons themselves must be. They may be comparatively light (~10 GeV), but they may also be very heavy, much heavier than the intermediate bosons. In the latter case, we encounter a new short-range strong interaction at an energy of order 1 TeV. One of the central problems of elementary-particle physics concerns the possible existence and properties of scalar bosons.

The vector X and Y bosons, whose existence is predicted by the majority of grand-unification schemes, are very heavy. Their expected masses are of order 10^{15} GeV, so that their production would require accelerators about a light year in length. But we might discover that these particles exist without resorting to accelerators at all. The point is that the interactions of these particles with leptons and quarks do not conserve the baryon and lepton quantum numbers and leads to instability of the proton. If the searches which are now being carried out actually reveal decay of the proton, this will be a triumph for the idea of grand unification.

It is well known that the experimental detection of the quanta of the gravitational field—gravitons—is a fantastically difficult problem because of the weakness of the gravitational interaction. So far, physicists have not succeeded in observing even classical gravitational waves under laboratory conditions. As to particles with $J=3/2$, the so-called gravitinos predicted by superunification schemes, at the present time it is difficult even to say whether their detection is a comparatively simple experimental problem or whether this problem is even more complex than the detection of gravitons. This is due, in particular, to the fact that we do not know what the gravitino mass must be.

Elementary particles with $J>2$ are usually rejected by theoreticians, since attempts to give consistent theoretical descriptions of them encounter serious difficulties. Nevertheless, it seems to me that it is too early to discount them: they will play an important role in the future theory of all particles and all interactions.

The short excursion up the ladder of spins leads to a paradoxical conclusion: of the known types of fundamental particles of contemporary physics, the majority (five out of nine) have so far not been detected experimentally. It is as if we are living in a house in which there are two walls and part of a roof, while all the rest exists only on paper. This impression becomes even stronger if we take into account the fact that at each rung of the ladder, in addition to the particles which we have already discussed, there may (or must, if we believe certain concrete theoretical models) exist many other particles. Among them are superheavy Higgs bosons, which are required in grand models, additional particles with $J=\frac{1}{2}$—the photino and gluino, which are predicted by supergravity, additional vector bosons, and so forth.

Another paradoxical conclusion is that we understand much better why we require those particles which have not yet been discovered than many of those of whose existence we have been convinced experimentally. To see this, let us consider the leptons and quarks more carefully. These are usually subdivided into three generations with similar quantum numbers, which display a specific symmetry between the leptons and quarks:

u	c	t (?)
d	s	b
ν_e	ν_μ	ν_τ
e	μ	τ .

(In this scheme, there is one missing particle—the t quark. A search for the reaction of its production in the colliding electron–positron beams of PETRA led to the conclusion that the mass of the t quark is greater than 18 GeV. There is nothing unnatural in the possibility that the t quark is even heavier, and our subsequent discussion is based on the assumption that the t quark exists. An alternative theoretical scheme not containing the t quark will be considered later in the section devoted to grand models.)

The role of the leptons and quarks in nature has two aspects, which by convention we shall call the macroscopic and microscopic roles. The macroscopic role of the particles of the first generation is clear: these constitute the material from which the world around us is constructed and without which it cannot function. The nuclei of atoms are constructed from u and d quarks, the atomic shells are constructed from electrons, and the electron neutrino is an essential participant in thermonuclear reactions, without which the Sun and the stars would be extinguished. As to the macroscopic role of the particles of the succeeding generations, at first sight it appears to be insignificant. These particles resemble the rough drafts which the Creator discarded as unsuccessful and which with our refined techniques we have unearthed in his wastepaper basket. We are now beginning to suspect that these particles played an important role in the first moments of the big bang. It is evidently precisely because of them that there arose the small excess of baryons over antibaryons, as a result of which there exists in the Universe not only relic photons and neutrinos, but also ordinary matter.

The microscopic role of the fermions is even less clear to us. By what principle is their presence in the Lagrangian necessary at all? Vector particles are required for gauge symmetry, and scalar particles are needed to break this symmetry. But why are fermions necessary? An answer to this question should be given by supersymmetry, which combines bosons and fermions into common multiplets. But supersymmetry will evidently also be able to tell us what the total number of fermion generations should be. In the appropriate sections, we shall discuss certain restrictions on this number imposed, on the one hand, by grand models, and, on the other, by astrophysics.

The truly fairground abundance and diversity of elementary particles is in sharp contrast to that trend towards unity about which we spoke earlier in the discussion of the fundamental principles of physics. The esthetic criterion of simplicity suggests to us that not all these particles are really elementary and, what is more, perhaps that they are all nonelementary and consist of a small number of truly fundamental particles. Very many different names for these fundamental particles exist in the literature: prequarks, metaquarks, preons, rishons, glicks, etc., but, unfortunately, there is as yet no simple and elegant model of prematter.

4. ELECTROMAGNETIC INTERACTION

Most of the processes—physical, chemical, and biological—with which we deal in everyday life are electromagnetic processes. It is therefore not surprising that the electromagnetic forces have been studied far better than the other fundamental forces. The theory of the interaction of photons with electrons and with other charged leptons—quantum electrodynamics—is the most highly elaborated of all physical theories. Its predictions have unprecedented accuracy. The experiments in which these predictions have been tested and confirmed (in particular, the measurement of the magnetic moment of the muon) also have unprecedented accuracy.

Electrons and muons are widely used as electromagnetic tools for investigating the particles that have been more poorly studied. Examples are provided by processes of hadron production in colliding electron–positron beams or reactions of deep inelastic scattering of leptons by nucleons.

Quantum electrodynamics with its gauge invariance and renormalizable perturbation theory has served as a prototype for the construction of more complex theories—quantum chromodynamics, the theory of the electroweak interaction, and models of grand unification. It is instructive that in the creation of these theories constructive use has been made not only of the merits of quantum electrodynamics, but also of its defects. The point is that quantum electrodynamics by itself is not a completely self-consistent theory. Owing to the polarization of the vacuum, the effective charge of the electron, which is partially screened at large distances (i.e., at small momentum transfers), grows with increasing momentum transfer and becomes infinite at sufficiently large momentum. Although in pure electrodynamics this effect sets in at such large momenta that they can never be obtained under laboratory conditions, from a purely fundamental point of view it has nevertheless aroused dissatisfaction. Further theoretical study of the polarization of the vacuum led to the discovery of the remarkable fact that in non-Abelian gauge theories, instead of vacuum screening of a charge, there may occur antiscreening, so that the effective charge does not grow, but falls off with increasing momentum transfer. This effect, known as asymptotic freedom, plays a principal role in quantum chromodynamics and in models of grand unification.

Quantum electrodynamics has now become a constituent part of the unified theory of the electromagnetic and weak interactions. Therefore it seems natural at first sight to turn directly to the discussion of weak-interaction physics. However, since the electroweak symmetry is spontaneously broken, its treatment entails a discussion of the properties of scalar particles and other related questions. Therefore, before turning to the weak interaction, we shall consider the strong interaction, whose theory is self-contained to a great extent.

5. STRONG INTERACTION

During the seventies, striking progress has been achieved in understanding the properties of the strong interaction. Physicists have formulated and developed quantum chromodynamics—a theory of the strong interaction which now has no rivals and whose conclusions are being confirmed in a large number of different physical phenomena.

Quantum chromodynamics has provided a natural qualitative explanation for a large number of previously established empirical or, more accurately, phenomenological regularities, such as isotopic invariance of the strong interaction, "flavor" SU(3) symmetry, the chiral symmetry $SU(2)_L \times SU(2)_R$, and its generalization $SU(3)_L \times SU(3)_R$. According to quantum chromodynamics, these symmetries follow from the universality of the gluon–quark coupling constant, from the vector character of the gluons, and, as will be clear from what follows, from the lightness of the u and d quarks and (to a lesser degree) of the s quark.

Quantum chromodynamics has provided an explanation and, at the same time, the natural limits of the nonrelativistic quark model, which describes the systematics of hadrons, and the parton model, which describes deep inelastic processes.

Finally, the validity of quantum chromodynamics has been confirmed by the discovery in recent years of a number of new objects and phenomena. A short enumeration of some of them follows.

1. Charmonium, the $c\bar{c}$ system, the first level of which—the J/ψ meson ($m = 3.097$ GeV)—was discovered in 1974. Since then, other $^3S_1$ levels of this system have been found: $\psi'(3.685)$, $\psi''(3.77)$, and $\psi'''(4.03)$, as well as a number of ^3P_J levels: $\chi_0(3.45)$, $\chi_1(3.51)$, and $\chi_2(3.55)$. Finally, the $\eta_c(2.98)$ meson, which is the ground state of paracharmonium pairs ($^1S_0$), was found at Stanford in 1979. (The numbers in parentheses indicate the masses in giga-electron-volts.)

2. Particles with explicit charm— the mesons $D^0[c\bar{u}]$ (1.863), $D^+[c\bar{d}]$ (1.868), $F^+[c\bar{s}]$ (2.04), and the baryon $\Lambda_c^+[cdu]$ (2.27).

3. Upsilonium, the $b\bar{b}$ system, the first level of which —the $\Upsilon(9.46)$—was found in 1976. This was followed by other S_1 levels: $\Upsilon(10.02)$, $\Upsilon''(10.40)$, and $\Upsilon''(10.55)$. The last of these, which was discovered recently using the intersecting storage rings at Cornell, has a width of order 10 MeV, whereas the widths of the preceding levels are measured in kilo-electron-volts. The natural interpretation is that the Υ''' decays strongly into meson pairs B^+B^- and/or $B^0\bar{B}^0$, where $B^- = b\bar{u}$ and $B^0 = b\bar{d}$.

4. Quark jets in e^+e^- annihilation, which were discovered in 1975 and which represent narrow beams of hadrons arising, according to the currently accepted interpretation, from the "fragmentation" of the quarks in the process $e^+e^- \to q\bar{q}$.

5. Gluon jets in e^+e^- annihilation, which were discovered using the electron–positron storage rings at DESY in 1979. These jets arise as a result of hadronic fragmentation of the gluons in the processes $e^+e^- \to q\bar{q}g$ and $\Upsilon \to 3g$.

Thus, a multitude of facts relating to hadron physics agrees qualitatively, and in a number of cases also quantitatively, with quantum chromodynamics. There is not a single fact which might contradict this theory. There is a not a single theoretical model which might compete with it. And nevertheless a theory of the strong interactions has by no means been constructed. It is true that at the present time we have no doubt that we have found the correct Lagrangian on which the strong interactions are based. However, we are able to solve the equations which follow from it only in the limit of small distances.

As we have already indicated above, the effective gluon–quark coupling constant falls off with increasing momentum transfer Q. As $Q \to \infty$, asymptotic freedom reigns:

$$\alpha_s(Q) \approx \frac{2\pi}{b(Q)\ln(Q/\Lambda)} \to 0;$$

here $b(Q) = 11 - (2/3)n_f(Q)$, where $n_f(Q)$ is the number of types of quarks whose masses satisfy the condition $m \ll Q$. (By virtue of the uncertainty relation, the screening of color charge due to heavier quarks with masses m does not change at distances $r \gtrsim 1/2m$.) If we move towards small momenta, $\alpha_s(Q)$ grows. The dimensional constant Λ (whose value will be discussed below) indicates the value of the momentum at which the expression for α_s becomes completely meaningless, tending to infinity. It is sometimes said in this connection that α_s has an infrared pole at $Q = \Lambda$. For $Q \gg \Lambda$, perturbation theory works, and we are then able to perform calculations. For $Q \approx \Lambda$, we fall in the region of the really strong interaction and are not able to perform calculations. Here conjecture and intuition are at work.

It is usually assumed that at distances of order $1/\Lambda$ the interaction between the mutually complementary color charges in a colorless hadron is so strong that it is not possible to separate them to distances greater than $1/\Lambda$. This hypothesis that the quarks and gluons are confined inside the white hadrons is in agreement with everything that we know from experiment about the strong interactions. But we do not yet know how to prove the correctness of this hypothesis on the basis of quantum chromodynamics and, what is important, to make a theoretical calculation of the confinement mechanism. This means, for example, that we do not know how to calculate the process of hadron fragmentation in quark or gluon jets. We are not able to go all the way from the Lagrangian to the nonrelativistic potential model or to the so-called bag model.

The unusual features of chromodynamics in comparison with electrodynamics are due to the fact that the color gauge group SU(3) is non-Abelian, so that the gluons carry color charges and themselves emit gluons. This "luminescence of light" leads to antiscreening of the color charges. As a result of the gluon self-interaction, the vacuum in quantum chromodynamics has a complex structure. Thus, in the classical limit it represents a degenerate system of infinitely many states differing from one another only by the topological properties of purely gauge potentials, each of which corresponds to zero field intensities. In the quasiclassical approximation, there are tunneling transitions between these degenerate states of the vacuum, known as instantons. Instantons are quasiclassical vacuum fluctuations whose amplitudes are proportional to $e^{-2\pi/\alpha_s}$. By virtue of the property of asymptotic freedom, instantons of small dimensions correspond to $\alpha_s \ll 1$, and the probability of their appearance is small. In this case, we speak of a rarefied gas of instantons. As the dimensions of the instantons approach $1/\Lambda$, the probability of their appearance rises and they "stick together" or "fuse": the vacuum fluctuations become very complex. It is possible that it is complex infrared vacuum fluctuations such as "melted instantons" that lead to the confinement of gluons and quarks.

In spite of the fact that the problem of confinement has not been solved, in a number of cases it is possible to obtain infrared-stable results. Thus, for example, the total cross section for the annihilation process

$$e^+e^- \to \gamma^* \to \text{hadrons}$$

at asymptotically large energies is

$$4\pi \frac{\alpha^2}{s} \sum_{q} e_q^2 \left[1 + \frac{\alpha_s(Q)}{\pi} + (1.98 - 0.115 n_f(Q)) \left(\frac{\alpha_s(Q)}{\pi} \right)^2 + \ldots \right];$$

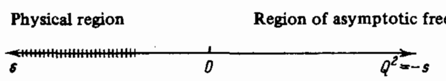
here e_q is the quark charge and the sum is taken over all types of quarks, s is the square of the total energy of the colliding electron and positron in the center-of-mass system, and $Q^2 = -s$. Thus, the cross section at large time-like momenta is expressed in terms of the value of α_s at large space-like momenta. Unfortunately, the accuracy of the corresponding experiments is as yet insufficient for a reliable determination of $\alpha_s(Q)$.

A bridge between the region of asymptotic freedom and the physical region is provided by dispersion relations, with whose use it is possible to obtain quantum-chromodynamical sum rules having the form

$$\int \sigma_{\exp}(s)\, \rho(s, Q^2)\, ds = F_{(\rho)}(Q^2)_{\text{theor}}.$$

The left-hand side of this equation is an integral over the physical region of the experimentally measured cross section, taken with some specially chosen weight function ρ.

The amplitude F on the right-hand side is calculated theoretically in the region in which perturbation theory still works:



The idea is to make an appropriate choice of $\rho(s, Q^2)$ so as to separate on the left-hand side the contribution to $\sigma_{\exp}$ from a resonance in which we are interested, and to have on the right-hand side controllably small corrections to perturbation theory. In the language of quark–gluon diagrams, the theoretical amplitude can be represented in the form of a sum of four terms:

$$F_{\text{theor}} = F_{\text{AF}} + F_{\text{pert}} + F_{\text{semipert}} + F_{\text{nonpert}}.$$

The diagrams which contribute to these terms are shown in Fig. 1, where the solid lines represent quark propagators, the dashed lines represent gluon propagators given by perturbation theory, and the wavy lines represent external currents. The breaks in the lines marked by crosses indicate that the momenta corresponding to the appropriate propagators are small and that these propagators are not described by perturbation theory.

The first term corresponds to the graph associated with asymptotically free quarks (Fig. 1a).

The second term takes into account the interaction of the quarks and gluons according to perturbation theory. The simplest graph of this kind is shown in Fig. 1b.

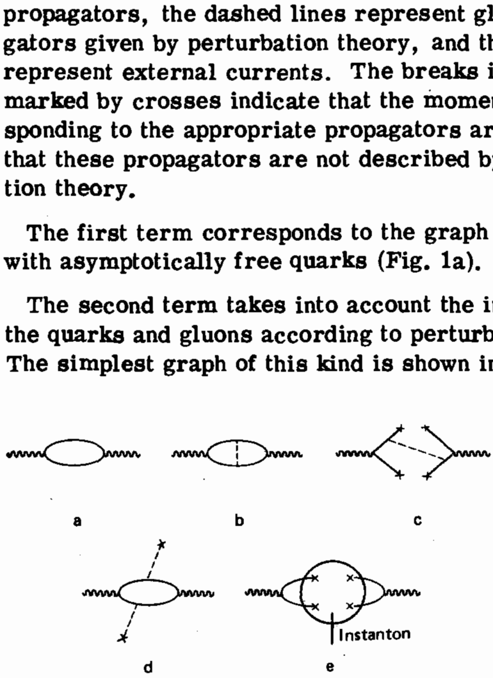


FIG. 1.

Here all the virtual particles carry large momenta, and their propagators are therefore described by perturbation theory.

The third term corresponds to diagrams in which large momenta flow only along certain propagators, while the momenta flowing along the other propagators are close to zero (Figs. 1c and 1d). The latter are formed under the action of large-scale vacuum fluctuations, and their contribution is proportional to the nonzero vacuum expectation values of the quark fields $\langle q\bar{q} \rangle$ and of the gluon fields $\langle GG \rangle$. In Figs. 1c and 1d they correspond to the broken quark lines and gluon lines, respectively.

Finally, the fourth term corresponds to those graphs (Fig. 1e) in which none of the virtual particles carries large momentum; all the momentum flows through a small-scale vacuum fluctuation (an instanton).

Quantum-chromodynamical sum rules have been applied to a number of quarkonia—mesonic systems of the type $q\bar{q}$: $J/\psi, \chi, \eta_c, \Upsilon, \rho, \omega, \pi, K$. Using a small number of parameters, they have made it possible to find relations between (and in certain cases, even to predict) a large number of physical observables: the masses and decay widths of the various mesons. (In particular, mention should be made of the prediction of the mass of the η_c meson, which has been confirmed by experiment, and the prediction of the mass of the P level of upsilonium, $m_{1P} - m_{\Upsilon} = 370 \pm 30$ MeV, which is awaiting confirmation.) These parameters include the quark and gluon vacuum expectation values ($\langle \bar{q}q \rangle$ and $\langle GG \rangle$) and, finally, the parameter Λ, which determines the value of α_s for a given Q.

In quantum chromodynamics, the quark masses decrease with increasing momentum transfer to the quark. In the case of light quarks, the theory gives the following expression for the masses of the so-called current quarks (i.e., for $Q \gg \Lambda$):

$$m_q(Q) = m_q^0 [\alpha_s(Q)]^{4/b(Q)}, \quad q = u, d, s,$$

where the parameters m_q^0 found from experiment have the values

$$m_u^0 = 3.5 - 5 \text{ MeV},$$
$$m_d^0 = 6.5 - 10 \text{ MeV},$$
$$m_s^0 = 100 - 250 \text{ MeV}.$$

Formally, these values of m_q^0 correspond to values $Q \approx e\Lambda$ for which $\alpha_s(Q) \approx 1$. However, at such small values of Q the foregoing expression for m_q is inapplicable, and we are dealing not with current quarks, but with so-called constituent quarks. For heavy quarks ($q = c, b$) when $Q \gg m_q$, the parametrization takes the form

$$m_c(Q) \approx 1.37 \left[\frac{\alpha_s(Q)}{\alpha_s(m_c)} \right]^{4/b} \text{GeV},$$
$$m_b(Q) \approx 4.8 \left[\frac{\alpha_s(Q)}{\alpha_s(m_b)} \right]^{4/b} \text{GeV}.$$

(We note that isotopic invariance of the strong interactions is broken not only because of the different charges of the u and d quarks, but also because of their different masses. The neutron is heavier than the proton because the d quark is heavier than the u quark.)

The vacuum expectation values of the fields of the light quarks are

$$\langle \bar{u}u\rangle = \langle \bar{d}d\rangle \approx \langle \bar{s}s\rangle \approx -\left(\frac{1\text{ GeV}}{4}\right)^3.$$

The energy density ε of the physical vacuum can be expressed in terms of the vacuum expectation value of the gluon fields:

$$\varepsilon = -\frac{9}{32\pi}\langle \alpha_s G^a_{\mu\nu} G^a_{\mu\nu}\rangle \approx -\left(\frac{1\text{ GeV}}{4}\right)^4$$

(here $G^a_{\mu\nu}$ are the gluon field intensities, for $a=1, 2, \ldots, 8$). The sign of ε is in agreement with what might have been expected on the basis of the so-called bag model: the breakdown of the physical vacuum inside the bag creates a positive volume density B of the bag. However, the value of B used in describing hadrons consisting of light quarks is about an order of magnitude smaller than ε. This means that there is only a partial breakdown of the physical vacuum in such hadrons.

The parameter Λ obtained from chromodynamical sum rules has a value of about 100 MeV or even somewhat less. This means that α_s is small for momenta $\lesssim 1$ GeV and that asymptotic freedom is destroyed not by the higher orders of perturbation theory in α_s, but by the vacuum fluctuations (graphs of the type shown in Figs. 1c–1e) giving power corrections of the form $\langle \bar{q}q\rangle^2/Q^6$, $\langle G^2\rangle/Q^4$, Λ^{11}/Q^{11}, etc.

Until recently, much larger values of Λ ($\gtrsim 500$ MeV) were given by the analysis of experimental data on deep inelastic interactions of neutrinos and muons with nucleons. However, recent data on the scattering of muons seem to be in agreement with $\Lambda \sim 100$ MeV. It would be highly desirable to return once again to neutrino experiments. The interpretation of the data on deep inelastic processes is more complex than the analysis of quarkonia by means of sum rules, since in deep inelastic processes the amplitude depends on a large number of kinematic invariants and contains external hadron lines. As a result, there is a danger of mistaking power corrections in Q (in particular, so-called terms of twist 4) for logarithmic corrections in Q proportional to α_s. The situation becomes unambiguous for $Q^2 \gg 10$ GeV^2, where the power corrections should be negligibly small.

A reliable knowledge of Λ (and hence of α_s) is essential both for an understanding of the strong interaction proper and for the analysis of models of grand synthesis, about which we shall say more in the appropriate section. If the quark masses are neglected, which is a completely legitimate approximation for hadrons consisting of light quarks, the Lagrangian of quantum chromodynamics has no dimensional parameter, since the constant α_s is dimensionless. The so-called dimensional transmutation—the appearance of a scale Λ in a scale-invariant theory—is related to the fact that because of quantum corrections α_s is a running coupling constant depending on Q. We cannot specify its value without stipulating at what value of Q this quantity is given. If the Planck mass m_P is assumed to be the natural scale in physics, we can conclude that Λ is determined by the value of $\alpha_s(m_P)$.

Ideally, having mastered quantum chromodynamics, we must learn to express in terms of Λ (and the quark masses) all quantities characterizing the hadrons. We shall mention here only some of the most fundamental of these quantities.

1. The constant f_π. Experimentally, this constant, which determines the decay of the pion, has a value of about 130 MeV. Theoretically, f_π characterizes the spontaneous breakdown of chiral $SU(2)_L \times SU(2)_R$ symmetry for which, from the massless quarks, there appear massive nucleons and massless Goldstone excitations—the pions.

2. The pion mass m_π. In essence, the calculation of m_π reduces to the problem of expressing the vacuum expectation value of the quark operators $\langle \bar{u}u\rangle = \langle \bar{d}d\rangle$ in terms of Λ, since there exists the well-known relation

$$m_\pi^2 f_\pi^2 = -(m_d + m_u)\langle 0|\bar{u}u + \bar{d}d|0\rangle.$$

3. The mass of the η' meson. At first sight, in the limit of massless quarks, the η' mass, like the masses of the π and η mesons, must be zero. Taking into account the perturbations due to the quark masses, but not taking instantons into account, it can be shown that the η' meson must be lighter than the η meson. The fact that in reality it is much heavier is the essence of the so-called U(1) problem. In this case, by U(1) we mean the symmetry which corresponds to conservation of the SU(3)-singlet axial current. The solution of the U(1) problem is related to the fact that conservation of the singlet axial current is violated in quantum chromodynamics by the axial gluon anomaly. A calculation of the mass of the η' meson would mean a calculation of the contribution from the gluon fluctuations of the vacuum.

The discovery and study of particles containing "live" valence gluons, whose existence is predicted by quantum chromodynamics, would be of special interest. We have in mind here mainly states of gluonium, or, as they are also called, glueballs—mesons containing no valence quarks at all. The simplest glueballs should consist of two or three gluons: $G = 2g$ or $G = 3g$. Quantum-chromodynamical sum rules permit only very rough estimates of the glueball mass M and width Γ:

	$G=2g$				$G=3g$	
J^P:	0^+	0^-	2^+	2^-	1^+	1^-
M, GeV:	1.2	2.5	1.2	1.2	2.5	2.5
Γ, MeV:	300	300	30	30	100	100

The uncertainties here are about ± 0.3 GeV for the masses and a factor $2^{\pm 1}$ for the widths. The best place in which to search for gluonium is the decays of the $b\bar{b}$ and $c\bar{c}$ systems, in particular, the radiative decays $\Upsilon \to G\gamma$ and $J/\psi \to G\gamma$, which proceed through intermediate states $gg\gamma$.

There have been recent observations of decays $J/\psi \to E(1420) + \gamma$, where the E meson decays via the channel $E \to \pi\delta \to \pi K\bar{K}$, in contrast to decays of the type $E \to \bar{K}K^* \to \bar{K}K\pi$, which are characteristic of the ordinary E mesons with $J^P = 1^+$. If the new E meson is a narrow gluonium state with $J = 2$, it has $P = -1$. Preliminary estimates have not predicted copious production of this particle in J/ψ decay.

The family of gluonia also includes the η' meson, which gives the dominant contribution to the sum rule for the pseudoscalar gluon current $G^a_{\mu\nu}G^a_{\rho\sigma}\varepsilon^{\mu\nu\rho\sigma}$. However, taking into account the fact that the dominant structure component in the η' corresponds to pairs of valence quarks (this is confirmed by the decay $\eta' \to 2\gamma$), it would be more correct to refer to the η' as glukonium—a peculiar hybrid of gluonium and quarkonium. We note that these glukonia also include certain other states contained in the table, in particular, the 0^{++} state with $M = 1.2$ GeV. This state is described by two close broad levels. Experimentally, it corresponds to the $\varepsilon(1300)$ meson.

It is of interest to search for mesons containing single valence gluons, for example, $c\bar{c}g$ or $b\bar{b}g$. An analysis which has recently been made in the bag model indicates that the state $b\bar{b}g$ with an exotic (i.e., absent in the quark–antiquark system) set of quantum numbers $\Upsilon^{PC} = 1^{-+}$ may be lighter than the Υ''.

Further development of strong-interaction theory, as hitherto, will take place on a broad front, on one flank of which we have the phenomenological interpretation of experimental data, and on the other particularly abstract mathematical constructions—what is sometimes called theoretical theory.

It is remarkable that it is of great value for the theory to have experiments performed not only at high energies, but also at low energies. The latter make it possible to bring order into the spectroscopy of hadrons, in particular, exotic (not of the type $q\bar{q}$ and qqq) mesons and baryons, baryonia, and two-baryon resonances. Here the old accelerators may still play a role, but particularly valuable information will be given by the low-energy antiproton ring under construction at CERN. It would also be highly desirable to construct colliding electron–positron beams with energies of the order of several GeV and with high luminosity ($L \gtrsim 10^{33}$ cm^{-2} sec^{-1}).

At high energies, the main efforts will be directed towards the search for new hadrons, both those containing new heavy quarks (t ?) and those consisting of light quarks and gluons, as well as the study of exclusive and inclusive reactions, and especially hadronic jets, whose detailed investigation may help us to understand the relationship between interactions at small and large distances. We must also bear in mind the unusual phenomena observed in cosmic rays (which we shall discuss later). It is possible that some of these phenomena are due to the existence of new heavy hadrons (superheavy quanta with masses 5–10 GeV, ultraheavy quanta with masses 30–80 GeV, and minicentaurs) consisting of light quarks and gluons. This last possibility is suggested by the large cross sections for production of these hadronic states.

As to "theoretical theory," we see here very great and promising activity. I shall mention only several approaches:

Computer experiments with gluodynamics (chromodynamics without quarks) on a lattice.

The study of gluodynamics in the limit of infinitely many colors ($N_c \to \infty$). Analysis of the relations between gluon loops and dual models.

Minidimensional exactly solvable models (the sine-Gordon model, the σ model, and the CP^{n-1} model, all in two-dimensional space, $d = 2$, the three-dimensional analog of the Ising model, etc.). The purpose of studying these models is to investigate the relations between instantons, asymptotic freedom, and confinement and to find methods of discovering hidden symmetries, whose existence is suspected by some theoreticians in gluodynamics.

In all these approaches, there is a conspicuous connection with statistical physics, particularly with the theory of phase transitions. Many of the ideas which have arisen here have already penetrated into the real world with continuous four-dimensional space and three colored quarks. This process will undoubtedly continue.

6. WEAK INTERACTION

Two characteristic features of the gauge theory of the electroweak interaction which differ sharply from pure electrodynamics and from chromodynamics are as follows: 1) the electroweak symmetry $SU(2)_L \times U(1)$ is spontaneously broken, as a consequence of which the intermediate bosons $W^{\pm}$ and Z^0 are massive; 2) the theory is mirror nonsymmetric. This asymmetry is a basic principle of the theory; the left-handed components of the fermions $\psi_L = \frac{1}{2}(1+\gamma_5)\psi$ form isotopic doublets with respect to the group $SU(2)_L$,

$$\binom{u}{d'}_L,\ \binom{c}{s'}_L,\ \binom{t?}{b'}_L,\ \binom{\nu_e}{e}_L,\ \binom{\nu_\mu}{\mu}_L,\ \binom{\nu_\tau}{\tau}_L,$$

whereas the right-handed components $\psi_R = \frac{1}{2}(1-\gamma_5)\psi$ form isotopic singlets.

The lower components of the doublets d′, s′, and b′ are linear superpositions of d, s, and b—eigenstates of the strong interaction. The 3×3 matrix which relates them is determined by four physical parameters: three angles ϑ_1, ϑ_2, ϑ_3 and one CP-noninvariant phase δ. The coefficients of all nine charged quark currents $\bar{u}d$, $\bar{u}s, \ldots, \bar{t}b$ can be expressed in terms of these parameters. If the neutrino masses are nonzero, the analogous matrix of charged lepton currents may be even more complex (see below).

The only parameter which is well known experimentally is ϑ_1, which characterizes mainly the mixing of the d and s quarks: $\vartheta_1 \approx 0.23$. In essence, it is this parameter which manifests itself in the decays of the τ lepton and in the decays of charmed particles; the other parameters manifest themselves weakly here.

Recent measurements of τ-lepton decays are in perfect agreement with the theoretical expectations. As to the charmed particles, the first experiments here yielded an unexpected result. We are referring to the nonleptonic decays of the D mesons. It was expected that the c quarks in these decays would decay "without paying attention" to which quarks are in their vicinity. If this naive picture held, the lifetime of the D^+ and D^0 mesons would be the same. Experiments showed, however, that $\tau_{D^+} \approx (3-5)\tau_{D^0}$. This may mean that an intensive annihilation process $c\bar{u} \to s\bar{d}$ takes place in the D^0 meson as a result of the interaction of the charged currents ($\bar{s}c$) and ($\bar{d}u$) and the emission of bremsstrahlung gluons which evidently accompanies them. The complication of the picture of semileptonic decays is associated only with the virtual strong interactions and does not affect our ideas about the weak interaction. This is also indicated, in particular, by the fact that, judging by the partial widths, the semileptonic decays of the c quark take place in complete agreement with expectations.

We shall now briefly discuss three other parameters which determine the charged quark currents. Unfortunately, they are rather poorly known: $0.1 \lesssim |\vartheta_2| \lesssim 0.50$, $|\vartheta_3| \lesssim 0.5$, and $|\delta| = 0.02-0.04$. A more recent observation of the decay products of b quarks (in the reaction $\Upsilon \to B\bar{B} \to \ldots$) confirms that the current $\bar{c}b$ is "stronger" than the current $\bar{u}b$. Measurement of the decays of charmed particles, and also of the relative probabilities of the various decays of the b quark and its lifetime, would make it possible to find all the angles ϑ_i. The parameter δ might manifest itself in the charge asymmetries in the decays of B^0 mesons. The discovery of the t quark and a measurement of its decays would serve as a check of the correctness of the entire picture.

In the second order in the weak interaction of the charged quark currents, they contribute to the mass difference of the K_L and K_S mesons and to the parameters characterizing the violation of CP invariance in the decays of these mesons. In particular, it turns out that, in contrast to the model of superweak CP violation, it is predicted here that there is a nonzero difference $\delta\eta$ between the amplitudes η_{+-} and η_{00} characterizing the decays $K_L \to \pi^+\pi^-$ and $K_L \to \pi^0\pi^0$, respectively. The predicted difference $\delta\eta \approx 2\times 10^{-5}$ is very small, but its measurement is of very great interest, since it would provide information about the mechanism of CP violation. We note that the expected contribution of the parameter δ to the other low-energy CP-odd effect, namely, the dipole moment of the neutron, is at least eight orders of magnitude smaller than the existing upper limit (1.6×10^{-24} e·cm) and is practically inaccessible experimentally. The possible contributions of Higgs bosons to $\delta\eta$ and d_n will be discussed elsewhere.

We turn now to the neutral currents. The standard theory of the electroweak interaction contains only diagonal currents of the type $\bar{e}e$, $\bar{d}d$, etc., in accordance with the fact that no manifestations of nondiagonal neutral currents (of the type $\bar{e}\mu$ or $\bar{d}s$) have been observed in nature, despite the high accuracy of the search experiments (this refers to the decays $\mu \to e\gamma$, $K \to \mu\bar{\mu}$, $K \to \mu\bar{e}$, etc.).

The ratio of the constants of the axial currents—the neutral and charged currents—is characterized by a constant ρ, which in the standard theory should be equal to unity. Experimentally, $\rho = 0.985 \pm 0.023 \pm 0.013$, where the first error is statistical and the second is systematic. As to the vector currents, their ratio is characterized by a free parameter of the theory—the weak angle θ_W. The three most accurate experiments give the following values for $\sin^2\theta_W$:

0.228 ± 0.018 — CDHS, 1979 (ν),
0.218 ± 0.014 — CHARM, 1979 (ν),
$0.224 \pm 0.02 \pm 0.01$ — SLAC, 1979 (e).

The first two values were obtained in two different neutrino experiments at CERN. The third value was obtained in an experiment on deep inelastic scattering of electrons by deuterons and protons; we indicate here the statistical and systematic errors, as well as the error due to the theoretical uncertainties associated with quantum chromodynamics. Thus, the most probable value of $\sin^2\theta_W$ lies in the interval 0.21–0.24. This value of θ_W is also in agreement with the results of an experiment at Novosibirsk, in which it was observed that there is a rotation of the plane of polarization of laser light when it passes through atomic bismuth vapor. This phenomenon is due to the weak nonconservation of parity in the interaction of the atomic electrons with the nucleus. Groups at Seattle and at Oxford have so far observed a somewhat smaller effect. Most recently, a group at the P. N. Lebedev Physics Institute (Moscow), which did not observe rotation of the plane of polarization, claimed that this result is in conflict with the theory.

In principle, it is possible to construct theoretical models which, while reproducing the prediction of the standard model of the electroweak interaction for deep inelastic scattering of both neutrinos and electrons, nevertheless admit a wide range of possibilities for the interaction of atomic electrons with nuclei, including the the absence of a P-odd effect. These models correspond to more complex gauge symmetries:

$$SU(2)_L \times SU(2)_R \times U(1) \quad \text{or} \quad SU(2)_L \times U(1) \times U(1);$$

they require the existence of additional Z^0 bosons. I do not yet see any reason to discard the standard theory in favor of these models.

Whereas the experiments on nonconservation of parity in atoms admit a unique theoretical interpretation, the effects of nonconservation of parity in nuclei are much more complex and are more poorly understood. Here both neutral and charged currents should contribute. In recent years, a large ($\sim 10^{-4}$) P-odd asymmetry has been discovered and measured in the angular distribution of the fragments of fission of the nuclei ^{234}U, ^{236}U, and ^{240}Pu induced by polarized neutrons. These results are important for an understanding of the fission mechanism. In general, experiments on nonconservation of parity in nuclear forces tend to reveal P-odd effects which exceed the theoretical predictions by an order of magnitude. This happened several years ago in the case of radiative neutron capture $np \to d\gamma$. There have now appeared data indicating that anomalous large P-odd effects exist in pp scattering and in the scattering of neutrons by tin. It is very important to achieve complete experimental clarity here.

The year 1980 was marked by a "neutrino boom." Neutrino masses and neutrino oscillations suddenly attracted the attention not only of physicists, but also of the mass media. The problem of neutrino oscillations is not new: it has been discussed continuously since the mid-fifties. There are several reasons for the splash of interest in this problem. The main reason is apparently the fact that the idea of a massless neutrino no longer seems elegant. With the advent of the latest gauge theories, the concept of the elegant and the natural has changed. Whereas previously a small but nonzero ratio m_ν/m_e was regarded as unnatural and "inelegant," we now understand that the vanishing of any mass is unnatural if this is not required by gauge symmetry, as is the case for the photon. But we do not know any gauge symmetries which would require the vanishing of the neutrino mass. Moreover, a number of models of grand unification predict extremely small but nonzero neutrino masses. Changes in the public consciousness have led to the appearance of an avalanche-like process of new interpretations of old experimental (laboratory) and observational (astrophysical) data and to the performance of new experiments. Astrophysical manifestations of a neutrino mass will be considered in a special section. Here we shall give a brief review of the laboratory experiments.

We begin with reactor experiments. A group working at the Savannah River reactor reported the observation of oscillations of a beam of reactor antineutrinos $\bar{\nu}_e$. We assume that $\bar{\nu}_e$ is a superposition of two states ν_1 and ν_2 having definite masses m_1 and m_2:

$$\bar{\nu}_e = \nu_1 \cos\alpha + \nu_2 \sin\alpha,$$

where α is a mixing angle. Then the probability that a neutrino $\bar{\nu}_e$ remains unchanged after traversing a certain distance has the form

$$P(\bar{\nu}_e \to \bar{\nu}_e) = 1 - \sin^2 2\alpha \cdot \sin^2\left(1.27 \frac{\delta m^2 L}{E}\right);$$

here E is the neutrino energy (in MeV), L is the distance from the source to the detector (in meters), and $\delta m^2 \doteq m_1^2 - m_2^2$ (in eV^2). As a consequence of such oscillations, the observed $\bar{\nu}_e$ spectrum should depend on the distance L. The $\bar{\nu}_e$ spectrum was measured at the Savannah River reactor in 1978 at two distances: 6 and 11.2 m. A comparison of these results with the calculated spectrum in 1980 led the group of authors to the conclusion that $\delta m^2 = 1$ eV^2.

After this, the results of an analysis of another experiment at the Savannah River reactor carried out in 1979 at a distance $L = 11.2$ m were published. In this experiment, a comparison was made of the yields of two reactions: $\bar{\nu}_e d \to nne^+$, induced by charged currents, and $\bar{\nu} d \to np\bar{\nu}$, induced by neutral currents. This comparison is less dependent on the possible uncertainties in the calculated $\bar{\nu}_e$ spectrum of the reactor. The second reaction is the standard one: its yield does not oscillate, since the neutral currents are the same for all types of neutrinos. The authors concluded that $P(\bar{\nu}_e \to \bar{\nu}_e) \approx 0.4 \pm 0.2$, and also $\delta m^2 \approx 1$ eV^2 and $\alpha \approx 0.5$.

Unfortunately, this conclusion is in direct conflict with the preliminary result of an experiment performed at Grenoble in 1980. This experiment, to measure the reaction $\bar{\nu}_e p \to e^+ n$ at $L = 8.7$ m, did not reveal any oscillations $[P(\bar{\nu}_e \to \bar{\nu}_e) \approx 0.87 \pm 0.14]$ and, in the opinion of the authors, rules out the values $\delta m^2 \approx 1$ eV^2 and $\alpha \approx 0.5$.

Another group of experiments was carried out using accelerators with beams of high-energy neutrinos with characteristic values $E/L \approx 10^2$. Experiments with ordinary beams of neutrinos from the decays of π^- and K mesons did not detect any oscillation effects; in particular, they did not detect a "loss" of muon neutrinos. In addition, no oscillations were revealed by measurements of the yields of the reactions $\nu_e d \to ppe^-$, $\bar{\nu}_e d \to nne^+$, and $\bar{\nu}_e p \to ne^+$ at the meson factory at Los Alamos. The neutrino sources for these reactions were stopped muons; the average distance from the source to the detector was $L = 9$ m.

It is possible that experiments with so-called direct neutrinos (from the decays of charmed particles) indicate a "loss" of ν_e. In these experiments, a beam of protons from an accelerator is incident on a massive metallic target, in which long-lived particles (pions, kaons, and hyperons) are absorbed without having a chance to decay, and the neutrino beam is enriched in neutrinos resulting from the decays of charmed particles. Such a beam of direct neutrinos should contain equal proportions of muon and electron neutrinos, provided that we have a good understanding of the decays of charmed particles. Experiments with direct neutrinos have been carried out at CERN, where three detectors were placed in the path of the beam. These detectors gave the following values for the ratio of the number of ν_e to the number of ν_μ:

BEBC:	$0.49^{+0.25}_{-0.15}$,
CHARM:	0.49 ± 0.21,
CDHS:	$0.77 \pm 0.18 \pm 0.24$.

The first two results indicate a possible loss of about half of the electron neutrinos, while the third is compatible with the absence of any loss. Evidently, the situation here will not be clarified without new neutrino experiments using accelerators.

In the spring of 1980 a preliminary summary was made of a long-term experiment by a group at the Institute of Theoretical and Experimental Physics at Moscow to measure the upper end of the electron spectrum in the decay of tritium, $H^3 \to He^3 + e^- + \bar{\nu}$. By analyzing the shape of the spectrum, the authors concluded that 14 eV $\lesssim m_{\bar{\nu}_e} \lesssim$ 46 eV at the 99% confidence level. This was a very skillful experiment; the spectrometer with which it was performed has record accuracy. Nevertheless (or perhaps because of this), a definitive conclusion that the mass of the electron neutrino is really close to 30 eV can be made only after performing independent measurements of the spectrum of tritium in other experiments. In particular, it would be good to adopt another source of tritium instead of the tritium-enriched valine ($NH_2C_4H_8COOH$) used in the Moscow experiment.

Let us now consider the phenomenology of neutrino oscillations. If the lepton quantum number is conserved, the neutrinos can have only Dirac masses, which transform into one another the states with given lepton number L and opposite helicities. In this case, after diagonalization, as in the case of quarks, the three generations of leptons are described by three neutrino masses, and the nine left-handed charged lepton currents are described by three angles and one phase.

Another simple case corresponds to nonconservation of the lepton quantum number in the presence of neutrino states with only left-handed helicity and antineutrino states with only right-handed helicity. In this case, the neutrinos have left-handed Majorana masses, which transform left-handed neutrinos and right-handed antineutrinos into one another. After diagonalization, we again have three masses, and the nine left-handed charged lepton currents are characterized by three angles and three phases.

On the basis of models of grand unification, it is to be expected that it is the case of left-helicity Majorana neutrinos that is realized in nature. Here, in contrast to the Dirac case, there must exist oscillations involving a change of the lepton number and, for example, an initial beam of antineutrinos can in principle produce not only positrons but also electrons. However, owing to the conservation of helicity, the yield of electrons here is suppressed by a factor $(m/E)^2$, where m is the mass of the neutrino and E is its energy.

In the general case, when there exist both left-handed and right-handed neutrinos and antineutrinos and the lepton number is not conserved, the Lagrangian may contain three types of mass terms: Dirac, left-handed Majorana, and right-handed Majorana masses. After diagonalization, there are six characteristic diagonal neutrino states. They can be called Majorana states, since each of these states, like a Majorana particle, occurs by itself: it does not have a mass-degenerate twin state—an antiparticle. On the other hand, in contrast to the usual case, the interactions of these Majorana neutrinos may in general be C- and CP-noninvariant.

So far, we have discussed the phenomenology of weak current–current interactions. Let us make a few remarks about intermediate bosons. It is expected that these particles will be detected in 1982 in the colliding beams of protons and antiprotons (2×270 GeV) at CERN. In the mid-eighties, mass production of Z^0 bosons should begin using the colliding e^+e^- beams of the accelerator LEP, the first run of which will have energy 2×50 GeV. We shall not discuss here what will happen if the $W^\pm$ and Z^0 bosons are not discovered. The possibility that these particles do not exist does not seem to me at all probable.

The expected properties of the intermediate bosons are predicted with high accuracy by the standard model. We note, in particular, that inclusion of radiative corrections increases the expected masses of these particles by about 3 GeV. These large radiative corrections occur because the boson masses are very large in comparison with the masses of the light leptons and quarks, where the electromagnetic constant is usually normalized.

Models containing additional W and Z bosons, some of which may be even lighter than the standard W and Z, have been widely discussed in the literature. Strict limits on the existence of such particles are imposed by experimental data to test quantum electrodynamics in the reaction $e^+e^-\to\mu^+\mu^-$, obtained recently at the maximum accessible energies of the PETRA storage rings. It will be much more difficult to make experimental tests of models in which all the additional bosons are heavier than the standard ones. In particular, experimental data on the longitudinal polarizations of the particles in β decay or in the decays of pions and kaons do not exclude the existence of right-handed currents if the masses of the corresponding right-handed bosons are greater than or of the order of 300 GeV. In models containing such bosons, for example, $SU(2)_L\times SU(2)_R$ or $U(2)_L\times U(2)_R$, mirror symmetry is restored at large momentum transfers.

7. SCALARS

In the contemporary theoretical literature scalar Higgs bosons (H bosons) play the role of secret benefactors: they give masses to all particles and determine all the mixing angles in the weak charged currents, but themselves remain imperceptible. So far, not a single direct experimental manifestation of these particles has been detected.

The mechanism of spontaneous breakdown of the gauge symmetry $SU(2)_L\times U(1)$ comes about through the scalars. What is responsible for this breakdown in the standard model is a specific self-interaction potential $\lambda(|\varphi|^2-\eta^2)^2$, where λ is a dimensionless constant, φ is an isodoublet scalar field, and $\eta=(\sqrt{2}\,G)^{-1/2}\approx 1/4$ TeV. The minimum of this potential corresponds to the existence of a scalar condensate and a nonzero vacuum expectation value of the field φ: $|\varphi|=\eta$. Owing to the interaction of the scalar field with the gauge fields, the appearance of the condensate leads to masses of the intermediate bosons: $m_W=\eta\sqrt{\pi\alpha}/\sin\theta_W$ and $m_Z=m_W/\cos\theta_W$; Yukawa interactions of the scalars with the leptons and quarks of the type $f\bar{\psi}_L\psi_L\varphi+f^*\bar{\psi}_R\psi_L\bar{\varphi}$ give masses to these fermions: $m=f\eta$. Because of the small ratio m/η, the Yukawa constants f must be small, and this is the reason for the imperceptibility of the scalars: their production cross sections contain the small factor $(m/\eta)^2$.

Each mass (diagonal or nondiagonal) has its own Yukawa constant. After diagonalization of the mass matrix, mixing angles appear in the weak currents. Even without allowance for the neutrino masses (see above), the number of Yukawa constants is very large: six quark masses, three charged-lepton masses, three angles, and one phase. This abundance of arbitrary constants which must be fixed "by hand" indicates that the theory is manifestly incomplete. The scalar potential with a nonzero vacuum condensate itself also seems unnatural to many physicists.

In the simplest variant of the standard theory of the electroweak interaction, there is one physical neutral scalar boson H^0. In more complex variants, one introduces additional bosons, both neutral and charged. However, their introduction requires precautions, since they may lead to transitions which are forbidden experimentally, for example, intense transitions of the type $K^0 \longleftrightarrow \bar{K}^0$. These transitions, like the nonconservation of flavor in the neutral currents in general, are not easy to avoid if allowance is made not only for the so-called tree approximation, but also for loops of radiative corrections. This difficulty is encountered, in particular, by theoretical attempts to obtain relations between the mixing angles and quark masses (for example, $\vartheta_1 \approx \sqrt{m_d/m_s}$) on the basis of some discrete symmetries.

Additional scalar bosons provide further possibilities for the violation of CP invariance. In principle, CP parity might be violated in the nonlinear potential of their self-interaction. The expected effects in this case might be close to the experimental upper limits for both the difference $\eta_{+-} - \eta_{00}$ and the dipole moment of the neutron. We recall that if the only source of CP violation is the phase δ in the charged-current matrix (i.e., in the Yukawa vertices), the expected CP-odd effects are small. This is especially true for d_n. Thus, the observation of a dipole moment of the neutron would serve as indirect evidence that elementary scalars exist.

One of the reasons why scalar bosons are difficult to find experimentally is that, in contrast to the W bosons, the masses of the H bosons are not predicted by the theory. In the standard minimal model, there exists the relationship $m_H = 2\sqrt{\lambda}\eta$. But the quantity λ is a free parameter.

The value taken on by λ has a significant effect on the entire physics of scalar and vector bosons at high energies in the region of a tera-electron-volt. If λ is small (the H boson in this case is light), perturbation theory is valid up to arbitrarily high energies. (We note that in the minimal model the minimum mass of the H boson is of the order of 10 GeV, which corresponds to $\lambda \sim \alpha^2$.) But if λ is large, so that $m_H \gg m_W$, then beyond about a tera-electron-volt there must be a strong interaction involving scalar and vector bosons. This new short-range ($r \sim 10^{-17}$ cm) interaction may lead to the appearance of a whole family of specific resonances with masses of the order of a TeV. It is very important to stress that the presence of this strong interaction between bosons at high energies should have practically no effect on the interactions between fermions at low energies. All the predictions of the standard model for the current-current amplitudes remain unchanged. In particular, the equality $\rho = 1$ holds as before, provided that the scalar particles are isospinors. Everything is calm outside, even if a storm is raging inside.

Owing to the radiative corrections, the effective self-interaction of the scalar fields contains contributions proportional to g^4 and f^4, where g are gauge constants and f are Yukawa constants. Taking into account these contributions and assuming the validity of perturbation theory (light H bosons), we can obtain bounds on the values of the Yukawa constants and hence on the fermion masses. These bounds have the form

$$m\ (\text{quark}) \leqslant 76\ \text{GeV},$$
$$m\ (\text{lepton}) \leqslant 100\ \text{GeV}.$$

Working within the framework of perturbation theory, it can be shown that larger fermion masses would lead to instability in the scalar sector. If leptons or quarks with larger masses were discovered, this would mean that the strong interaction discussed above occurs at small distances.

The experimental search for scalar particles is a primary task. Unfortunately, the expected cross sections for production of H^0 bosons are very small (as a rule, $\leq 10^{-35}$ cm^2). What seems very promising is the associated production of Z and H bosons: $e^+e^- \to Z^0H^0$. The search for this reaction is one of the most important tasks of the accelerator LEP. It will be easier to discover the $H^\pm$ bosons if they exist, since H^+H^- pairs should be produced electromagnetically, for example, $e^+e^- \to H^+H^-$.

8. TECHNICOLOR. NEW PHYSICS IN THE TeV REGION

Technicolor is a not completely appropriate name for a hypothetical strong interaction with confinement radius of order 10^{-17} cm and characteristic energy scale of order 1 TeV. This and similar interactions are also known in the literature as "metacolor" and "hypercolor." But it seems reasonable to use these other names in discussing hypothetical strong interactions at much higher energies up to m_P, reserving the prefix "techni" for an interaction whose experimental investigation is within the scope of the technical possibilities of the present century.

The hypothesis of technicolor was originally put forward as a means of getting rid of the elementary scalar particles with their numerous disadvantages and replacing them by composite particles constructed from fermions, so-called techniquarks. According to this hypothesis, there exist special techniquarks which interact via the exchange of technigluons and form technihadrons. The masses of the majority of technihadrons should be of the order of a tera-electron-volt. However, some of them should be massless (these are the Goldstone bosons, corresponding to violation of the exact global symmetries obeyed by techniquarks) or very light (these are the pseudo-Goldstone bosons, corresponding to approximate symmetries). Assuming that the technicolor Lagrangian possesses strict chiral SU(2) symmetry, the triplet of Goldstone bosons corresponding to violation of this symmetry is "consumed" by a triplet of massless vector bosons, which as a result become massive. Then, by virtue of the SU(2) invariance of technicolor, we have the relations $M_W = M_Z \times \cos\theta_W$ and $\rho = 1$, which are characteristic of the case of spontaneous breakdown of the symmetry by elementary isospinor scalars.

Unfortunately, a single technicolor is insufficient to obtain nonzero masses of the leptons and quarks. To solve this problem, it is assumed that technicolor (TC) is part of an extended technicolor (ETC). With the emission of an ETC boson, an ordinary quark q is converted into a techniquark Q. As a result, $m_q \sim \langle\bar{Q}Q\rangle / m^2_{ETC}$, where $\langle\bar{Q}Q\rangle \approx \text{TeV}^3$ is the vacuum expectation value of the techniquark condensate, and m_{ETC} are the masses of the vector ETC bosons. It can be seen that these masses must be of the order of 100 TeV. They can be obtained as a result of one further interaction—"technicolor prime" (TC′). We see that instead of a single—even if possessing certain defects—scalar, it is necessary to introduce a whole hierarchy of interactions, which seems much less attractive.

Like the ordinary Higgs scheme, the model of technicolor predicts the existence of particles with zero spin. Some of these must be pseudoscalars, since they arise as a result of the violation of chiral symmetry. We note, in particular, the technianalog of the ordinary η (or η') meson, whose properties have recently been the subject of many papers. As we have already mentioned, there may also exist numerous light Goldstone bosons, in particular, bosons corresponding to so-called horizontal symmetry, i.e., symmetry between right-handed quarks of the same charge: $u_R \leftrightarrow c_R \leftrightarrow t_R$ and $d_R \leftrightarrow s_R \leftrightarrow b_R$. To eliminate the undesirable Goldstone bosons, it is necessary to assume that the horizontal symmetry is local and that the corresponding gauge constant is sufficiently small.

I do not share the prejudice of many theoreticians against elementary scalar particles. It seems to me that $J=0$ is no worse than $J=1/2$ or $J=1$. Thus, if the latter are assumed to be elementary, it is reasonable to expect that the former are also elementary. Conversely, if we construct, in spite of everything, a model in which the scalar bosons are composite, it is natural for the leptons, the quarks, and even the "sacred" vector bosons to be composite.

Recently, the possibility of constructing composite leptons and quarks in models such as technicolor has been widely discussed. A complication here is that if we follow the analogy with ordinary composite baryons, the expected masses of these composite leptons and quarks are enormous, of the order of the inverse confinement radius. This estimate follows directly from the uncertainty relation. The only way of circumventing this difficulty which has been discussed in the literature involves the use of chiral symmetry. If the elementary fermions are massless and possess chiral-invariant interactions, there are in principle two possibilities for composite particles.

1. Chiral symmetry is realized linearly: the axial current is conserved because the composite fermions are massless. Then all the composite bosons are in general massive.

2. Chiral symmetry is realized nonlinearly: the axial current is conserved because there exist massless pseudoscalar Goldstone bosons. Then the composite baryons are massive.

In the case of quantum chromodynamics, the second variant is realized. It is not yet clear whether the first variant can be realized at all. A necessary condition for this is the equality of the axial anomalies for composite and elementary fermions. This condition is not satisfied in the case of SU(3) symmetry, since the multiplets of composite and elementary fermions have different trialities. Thus, quantum chromodynamics has no choice. However, even in those cases when the equality of the anomalies is satisfied [for example, for the groups $O(2n+1)$], there is serious doubt that the composite fermions remain massless. It looks as if this does not happen, at any rate, in the limit of large n. It appears that to obtain light composite fermions it is necessary to go beyond the customary models and ideas and to pursue the search for completely new mechanisms, and possibly principles, of interaction of particles at small distances. In this connection, it seems very interesting to study conformally invariant theories: if bound states existed in such theories, they would have to be massless.

Is there experimental evidence that some new physics is operative at distances of order 10^{-17} cm $\approx$ TeV^{-1}? So far, energies of the order of or greater than a TeV in the center-of-mass system have been obtained only in cosmic rays. A number of unusual phenomena have been discovered at these energies.

First, there are six events of the centaur type (of which one is particularly convincing). In these events, at laboratory energies of order 10^3 TeV, about 100 charged particles were produced with total mass of order 200–300 GeV, but no π^0 mesons were produced. No satisfactory explanation of these events has yet been found. Some authors think that we are dealing here with the decay of massive particles of a new type—clusters of superdense quark matter. Second, there are 13 events of the minicentaur type, corresponding to production of several tens of particles with total mass 30–40 GeV.

Third, there are data indicating that at laboratory energies of order 100 TeV in showers there are secondary particles with anomalously high penetrating power (the so-called Tien-Shan effect).

Fourth, there are data indicating a rapid growth with energy of the multiplicity of secondary particles and their transverse momenta. Thus, at laboratory energy of $\sim 10^3$ TeV the multiplicity is much greater than that obtained with a logarithmic extrapolation of the accelerator data. It is even said that the growth obeys an $E^{1/4}$ law. The measured transverse momenta are as high as 10 GeV. The data also reveal so-called binocular events: two narrow jets with large transverse momenta. It is not excluded that this whole group of phenomena can be interpreted, at any rate partially, in terms of quantum chromodynamics.

It does not seem to me that the phenomena enumerated above are related to technicolor. The point is that, by virtue of the small dimensions of the region of techniconfinement, the cross sections for production of techniparticles should be very small: of order 10^{-4}, or probably even 10^{-6}, of the total cross section. At the same time, the effects discussed by the cosmic-ray physicists have, according to them, cross sections of order 10^{-2}–10^{-1} of the total cross section.

From the foregoing discussion, it should be obvious that there is no problem more important in its prospects than that of entering the TeV region of energies by means of accelerators. The first step in this direction would be the creation of an accelerator-storage-ring complex with colliding pp beams of energy 2×3 TeV. At the end of the seventies, there were many discussions about the possible construction at the end of this century of a very big international accelerator with colliding $p\bar{p}$ beams of energy 2×30 TeV. Finally, if the idea of colliding linear electron–positron beams is realized, we may also have electron accelerators in the TeV region.

9. GRAND UNIFICATION

The simplest variants which realize the idea of grand unification are based on the premise that there are no new fundamental forces not only at a tera-electron-volt, but also at much higher energies, up to colossal energies of order 10^{15} GeV. Here we shall consider the successes and difficulties of this approach. We begin with SU(5), the simplest of the currently active models of grand unification, after which we shall consider more complex models: SO(10), E_6, etc.

The group SU(5) has minimum rank among the semisimple Lie groups containing as a subgroup the product $H=SU(3)\times SU(2)\times U(1)$.

Each generation of fermions is described in the SU(5) model by a reducible representation. Thus, for example, the 16 left-helicity spinors of the first generation can be decomposed as follows into irreducible multiplets:

$$10\,(u_i,\ \tilde{u}_i,\ d_i,\ e^+)_L+\bar{5}\,(\tilde{d}_i,\ e^-,\ \nu)_L+1\,(\tilde{\nu})_L;$$

here $i=1,2,3$ is a color index. Usually, the singlets—the left-handed neutrino and the corresponding antiparticle, the right-handed antineutrino—are eliminated from the discussion from the outset, since they do not interact with the gauge fields.

Since electric charge is one of the generators of the group SU(5), the total charge of an SU(5) multiplet must be equal to zero. This establishes a relation between the charges of the leptons and quarks. In particular, it follows from the equality $Q_\nu=0$ that $Q_d=+(1/3)Q_e$ and $Q_u=-(2/3)Q_e$. Thus, we obtain an explanation of the fractional charges of the quarks.

The theory contains 24 gauge fields: γ, $W^{\pm}$, Z, eight gluons, six X bosons, and six Y bosons. The properties of the X and Y bosons are unusual. They are colored and have fractional electric charges: $X_i^{\pm 4/3}$ and $Y_i^{\pm 1/3}$. The masses of these bosons can be estimated on the basis of the fact that at momenta larger than these masses the symmetry $SU(3)\times SU(2)\times U(1)$ is reduced to SU(5).

Let us look at the "map" of the fundamental forces (Fig. 2). Here we have plotted along the horizontal axis the logarithm of the momentum transfer, and along the vertical axis the inverse squares of the "running" gauge constants of the strong $(1/\alpha_s)$, weak $(1/\alpha_W)$, and electromagnetic $(1/\alpha_{em})$ interactions. (The factor 3/8 gives the electromagnetic constant the same normalization as the constants α_W and α_s. Here the sum of the squares of the charges of the particles in a multiplet is equal to the sum of the squares of their isospin projections.)

The dependence of the constants α_i on the momentum

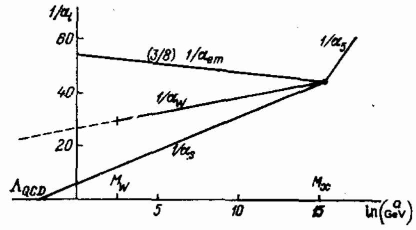


FIG. 2.

Q is due to the polarization of the vacuum and is given by the relation

$$\alpha_i^{-1}(Q_1)-\alpha_i^{-1}(Q_2)=\frac{b_i}{2\pi}\ln\frac{Q_1}{Q_2};$$

here

$$b_{em}=\frac{22}{3}\sum q_1^2-\frac{4}{3}\sum q_{1/2}^2-\frac{1}{3}\sum q_0^2,$$

$$b_s=11-\frac{2}{3}n_f(Q),$$

$$b_W=\frac{22}{3}-\frac{2}{3}n_f-\frac{1}{6}n_H,$$

where q_1, $q_{1/2}$, and q_0 are the electric charges of the particles with spin 1, $\frac{1}{2}$, and 0, respectively, the sum being taken over all these particles, $n_f(Q)$ is the number of different flavors of quarks with masses $m\ll Q$, and n_H is the number of scalar doublets with $m\ll Q$. The dashed line shows the continuation of the trajectory $1/\alpha_W$ in the case when the W bosons and the fermions are massless.

We see that the running constants meet at a single point. The coordinates of this remarkable point are $\alpha_{GU}\approx 0.02$ and $Q_{GU}\approx 5\times 10^{14}$ GeV. (The subscript GU stands for "grand unification.")

The masses of the X and Y bosons should be expected to be of order Q_{GU}. By considering the interaction of the X and Y bosons with fermions, it is easy to see that the X and Y should transform into both antilepton–antiquark pairs and quark pairs:

$$e^+\bar{d}\leftarrow X\rightarrow uu,$$

$$\bar{\nu}_e\bar{d}\leftarrow Y\rightarrow ud.$$

As a result of these transitions, which take place with nonconservation of baryon charge, nucleons may be converted into leptons:

$$p=uud\rightarrow e^+,$$

$$n=udd\rightarrow\bar{\nu}.$$

(Of course, by virtue of the conservation of energy and momentum, a single lepton cannot be emitted, and we are concerned with decays of the type $p\rightarrow e^+\pi^0$, $p\rightarrow e^+\pi^+\pi^-$, etc.). According to theoretical estimates, the lifetime of the proton is $\tau_p^{theor}=10^{30\pm3}$ yr. Experiment gives $\tau_p^{exp}>10^{30}$ yr. The main uncertainty in τ_p^{theor} is associated with the X-boson mass M_X. We note that $\tau_p\propto M_X$, while $M_X\propto\Lambda_{QCD}$ (see the "map"). The theoretical value 10^{30} yr corresponds to a value $\Lambda=400$ MeV. If we take $\Lambda=100$ MeV, in accordance with the arguments given above in the discussion of quantum chromodynamics, a barrel of water would be sufficient to detect the decay of the proton. Compare this with the scale of underground experiments, both planned and in progress, which lead to masses 10^3–10^4 tons and are capable of measuring $\tau_p \approx 10^{33}$ yr.

The foregoing theoretical estimate of the lifetime of the proton refers to a variant of the so-called minimal SU(5) model with the minimal set of multiplets of scalar particles: 24, 5, and $\overline{5}$. If we add to them a 45-plet of scalars, we can raise the expected lifetime of the proton above the possibilities of the most sensitive of the experiments under preparation. We note that the uncertainties in the prediction of τ_p are just as large in other grand models based on groups of higher rank, which we shall discuss later. Thus, if the decay of the proton is detected experimentally, it will be necessary to regard this as a special gift of Nature to physicists.

The observation of proton decay would be extraordinarily important. This would be the experiment of the century. Like a tuning fork, it would tune all high-energy physics to Planck frequencies. By proving the validity of extrapolations by 14 orders of magnitude in energy, this experiment would determine the further development of high-energy physics for many years.

By measuring the partial widths of the individual channels, by measuring the angular distributions and polarizations of the decay products, and by establishing the selection rules obeyed by the individual decay channels, it would be possible to learn much about the physics of very small distances.

Among the most interesting problems, we mention the question of whether the difference between the baryon and lepton quantum numbers, $B-L$, is conserved. [It is conserved in the minimal SU(5) theory.] Another interesting question is the mixing of fermions of different generations when they interact with X and Y bosons. If this mixing is small, as in the minimal SU(5) model, then $\Gamma(p\to e^+\pi^0) \gg \Gamma(p\to\mu^+\pi^0)$. Otherwise, this is not so.

If the decay of the proton were detected, we would be able to peek into the "hot laboratory" of grand unification, as through a keyhole. Unfortunately, we cannot count on more at the present time.

Nonconservation of baryon charge makes it possible in principle to explain the observed baryon asymmetry of the Universe. This question will be considered in more detail in the section devoted to astrophysics.

Let us consider briefly the other predictions of the SU(5) model. The model predicts the value $\sin^2\theta_W = 0.20-0.21$. This value, on the one hand, is so close to the experimental value (0.22–0.24) that a fortuitous coincidence seems unlikely. On the other hand, however, there is some discrepancy between the two values, which with further refinement of both theory and experiment may prove to be fatal for the model.

In unbroken SU(5) above the energy of grand unification, the masses of the charged lepton and "lower" quark (with $Q=-1/3$) in each generation must be equal. If we take into account the growth of the quark mass in going from $\sim 10^{15}$ GeV to ~giga-electron-volt (when the lepton mass changes insignificantly), we can obtain

$$\frac{m_b}{m_\tau} \approx \frac{m_s}{m_\mu} \approx \frac{m_d}{m_e} \approx 3.$$

For the b quark and τ lepton, this relation is satisfied reasonably well. For the other two pairs, it is not satisfied. The discrepancy can be eliminated by means of additional scalar multiplets or by introducing an additional nonrenormalizable interaction of the type $(1/M)\overline{\varphi}\varphi\overline{\psi}\psi$, where φ and ψ are scalar and spinor fields, respectively, and the constant M is of the order of the Planck mass.

From the formula for the mass of the b quark, $m_b(Q) \approx 4.8[\alpha_s(Q)/\alpha_s(m_b)]^{4/b(Q)}$ GeV, it is clear that the extrapolation depends on the number of quark flavors n_f, since $b = 11-(2/3)n_f$. It is easy to see that the relation $m_b/m_\tau \approx 3$ is strongly violated if the number of fermion generations is not 3 but, say, 6. It is important to stress that this limit on the number of fermion generations also refers to the contributions of very heavy fermions.

The SU(5) model does not give strict predictions for the neutrino mass. Although massless neutrinos are considered most natural for this model, in principle the model also admits other possibilities. For example, by introducing a scalar 15-plet, we can assign a Majorana mass m_L to the left-handed neutrino. However, if this mass is to be acceptably small, the corresponding Yukawa constant must be extremely small, and this seems unnatural. By invoking a singlet ν_R, we can in principle also obtain a Dirac mass m_D. We note that a Majorana mass m_L can also be given by a nonrenormalizable interaction $\sim(1/M)\varphi\varphi\psi\psi$; in this case,

$$m_L \sim \frac{M_p^2}{M} \sim 10^{-6}-10^{-10}\ \text{eV}.$$

The most complex unsolved problem of the SU(5) model is the problem of the hierarchy of masses. The masses of all the particles in this model are generated by the mechanism of spontaneous symmetry breakdown in the scalar sector of the theory. In the simplest variant, the breaking of SU(5) to SU(3)× SU(2)× U(1) gives the vacuum expectation value $\langle H_{24}\rangle \approx 10^{15}$ GeV, and the breaking of SU(3)× SU(2)× U(1) to SU(3)× U(1) gives the vacuum expectation value $\langle H_5\rangle \approx 10^2$ GeV. The scalar bosons H_{24} and H_5 are related. In any case, such a relation must exist because of the exchanges of gauge fields. So far, it is not understood how, despite the radiative corrections, a difference between the vacuum expectation values by 13 orders of magnitude can occur and not be broken. The problem of the hierarchy of scales exists not only in SU(5), but also in other grand models, to whose discussion we now turn.

The fifth-rank group SO(10) has certain advantages in comparison with SU(5). First of all, all the fermions of a single generation are contained in an irreducible representation of SO(10): $16 = 1+\overline{5}+10$. The SO(10) model has 45 gauge bosons. So-called triangle anomalies are automatically absent in the model. The minimal SO(10) model contains four scalar multiplets: 10, 16, $\overline{16}$, and 45. Variants including scalar multiplets 120 and 126 are also considered. [We give the SU(5) composition of some of the SO(10) representations: $10 = 5+\overline{5}$, $45 = 24+10+\overline{10}+1$, $120 = 45+\overline{45}+10+\overline{10}+5+\overline{5}$, and $126 = 50+45+\overline{15}+10+5+1$.]

The breakdown of SO(10) symmetry can occur in various ways:

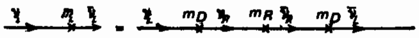
1) SO(10) → SU(5) → SU(3) × SU(2) × U(1) → SU(3) × U(1),

2) $SO(10) \xrightarrow{M_A} SU(4) \times SU(2) \times U(1) \xrightarrow{M_B} SU(3) \times SU(2) \times U(1) \xrightarrow{M_C} SU(3) \times U(1)$,

3) $SO(10) \to \ldots \to SU(3) \times SU(2)_L \times SU(2)_R \times U(1) \xrightarrow{M_R} SU(3) \times SU(2) \times U(1) \to SU(3) \times U(1)$.

In the first chain, the vacuum expectation value $\langle H_{16} \rangle$ breaks SO(10) down to SU(5); $\langle H_{45} \rangle$ breaks SU(5) down to SU(3)×SU(2)×U(1); $\langle H_{10} \rangle$ gives the last breakdown. We note that the other two chains make it possible to obtain values of $\sin^2\theta_W$ greater than in the SU(5) model. Thus, for example, in the second chain, $\sin^2\theta_W = 0.23$ for $M_A \approx 10^{15}$ GeV and $M_B \approx 10^{12}$ GeV. In the third chain, $0.25 \geq \sin^2\theta_W \geq 0.21$ for $10^6 \leq M_R \leq 10^{11}$ GeV.

A small mass m_L of the left-handed neutrino can occur in SO (10) in a natural way as a result of the mixing $\nu_L \to \nu_R$:



Here we have taken into account the fact that m_L gives the transitions $\nu_L \leftrightarrow \bar{\nu}_L$, m_R gives the transitions $\nu_R \leftrightarrow \bar{\nu}_R$, and m_D gives the transitions $\nu_L \leftrightarrow \nu_R$ and $\bar{\nu}_L \leftrightarrow \bar{\nu}_R$. The natural value of m_R in SO(10) must be very large. For the variant containing a scalar 126-plet, $m_R \sim 10^{15}$ GeV, since the 126-plet includes an SU(5) singlet. For the minimal variant, a value $m_R \neq 0$ arises only when allowance is made for radiative corrections and is therefore smaller: $m_R \sim 10^{10}$ GeV. If we take $m_D \sim 1$ GeV, as for the c quark, then by looking at the figure it is easy to obtain the following value for the effective Majorana mass of the left-handed neutrino:

$$m_L^{\text{eff}} \sim \frac{m_D^2}{m_R} \sim 10^{-1} - 10^{-6}\,\text{eV}.$$

This estimate refers to ν_μ. It would be natural in this case to expect smaller values for ν_e and larger values for ν_τ.

In models of grand unification, and in particular in the SO(10) model, in addition to the ordinary decay of the proton, there may be processes in which the baryon quantum number changes by two units. This applies to decays in which two nucleons are converted into mesons, and also oscillation transitions neutron ↔ antineutron in the vacuum. A source of such processes is the interaction of quarks with scalar fields of the type shown in the diagram of Fig. 3. Here the wavy lines describe scalar particles, and the cross indicates the vacuum expectation value $\langle H_{126} \rangle$.

Dimensional considerations enable us to relate the frequency of the oscillations to the decay probability on the basis of the fact that the first quantity is linear, while the second is quadratic, in the effective interaction constant:

$$\tau_{\text{dec}} \sim \tau_{\text{osc}}^2 m_{\text{proton}}.$$

It follows from this relation that a lifetime of order 10^{30} yr corresponds to an oscillation period of the order of a year. The experimental serach for such oscillations in intense beams of slow neutrons seems feasible and is of very great interest. Perhaps it would make sense to search experimentally for oscillation transitions of ordinary atoms into antiatoms (for example, $e^-p \to e^+\bar{p}$), although at the present time we can see no theoretical basis for such transitions; even if such transitions existed, they would be very slow, owing to the large dimensions of the atoms.

A model based on the exceptional sixth-rank group E_6 has a number of unusual properties.

At the present time, the E_6 model looks somewhat too big for the description of the known particles. The group E_6 contains SO(10) as its subgroup. Thus, for example, the 16 left-helicity fermions of a single generation, which form the fundamental representation of SO(10), belong in the group E_6 to a 27-plet: $27 = 16 + 10 + 1$. Of the remaining 11 states, six belong to a singlet quark with $Q = -1/3$, one belongs to a singlet Majorana lepton N^0, and four belong to a doublet of leptons (L^+, L^0). By assumption, they are all superheavy. The E_6 model contains no triangle anomalies. It includes 78 gauge bosons $(78 = 45 + 16 + \overline{16} + 1)$.

The simplest Higgs multiplets in E_6 are the 27 and 351 $(27 \times 27 = \overline{27} + 351'_S + 351_A$, where $351 = 144 + 126 + 54 + 16 + 10 + 1$ and $351' = 144 + 120 + 45 + 16 + \overline{16} + 10)$. It can be seen that the Higgs multiplets (27 and 351′) have the property that they can be represented in the form of bilinear combinations of fermions. Therefore the scalar bosons can in principle be regarded as composite. This is not so in the case of the group SO(10), since the Higgs multiplets in this group include the 16 and 45, which cannot be represented as bilinear combinations of fermions $(16 \times 16 = 126 + 120 + 10)$. Besides the chain of symmetry breakings $E_6 \to SO(10) \to SU(5)$, the group E_6 can also be broken in many other ways. In particular, it can be broken directly down to SU(3) × SU(2) × U(1).

All three models considered above have the property that the three fermion generations are contained in them independently of each other, without being combined into a single multiplet. The second and third generations are, as it were, xerox copies of the first. There are many other models of this "xerox" type, for example, $SU(8)_L \times SU(8)_R$ or SU(16). The higher the rank of the group, the more particles there are in its multiplets and the more diverse are the chains of symmetry breaking. In the SU(5) model, there is no new physics from 10^2 GeV to 10^{14} GeV, but only a slow variation of the gauge constants. This energy region is frequently called a gauge desert. In models of higher rank, oases of new interactions appear in the desert. There are more and more of them with increasing rank of the group. They gradually merge into an enormous exotic garden.

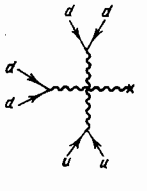


FIG. 3.

As a rule, models in which all the generations are combined into a single big family are even richer in new phenomena. Several types of such "truly grand" models have been considered:

1) Orthogonal models: SO(18), SO(22),..., SO($4n+2$) with 2^{2n} particles in a spinor multiplet.

2) Unitary models: SU(8), SU(11), SU(14),....

3) Products of simple groups related by a discrete symmetry, so that there is only one gauge constant:

$[SU(5)]^2$, $[SO(10)]^2$, $[SU(6)]^4$, ...

4) Exceptional groups: E_6, E_7, E_8.

These models contain gauge bosons which give so-called horizontal transitions between the generations. If the masses of these bosons are sufficiently small, there must exist rare decays of the type $\mu \to e\gamma$.

Some of the "truly grand" models are so "tight" and "rigid" that it will evidently be possible to test them in the very near future. This is particularly true of the E_6 model (which should not be confused with the "xerox" E_6 considered above). In this model, all the known fermions belong to a single 27-plet, which contains two quarks with $Q=+2/3$, four quarks with $Q=-1/3$, five leptons with $Q=0$, and four leptons with $Q=-1$:

$$\begin{array}{lllll} u & c & & & \\ d & s & b & h & \\ \nu_e & \nu_\mu & \nu_\tau & \nu_\lambda & \nu_\rho \\ e & \mu & \tau & \lambda & \end{array}$$

(If color is taken into account, we have 18 quarks and 9 leptons.) Thus, the model predicts the existence of four more hitherto unobserved particles. Note that there is not a quark among them! There must exist in the model neutral currents which change the quark flavor. According to variants of the model considered in the literature in which these neutral currents contain no transitions between light quarks, charged currents which transform b quarks into the c quark are excluded. This prediction seems to be in conflict with experimental data obtained very recently in the study of the decay products of the Υ''' meson.

A special point in the ocean of models is the most senior one of the exceptional groups, E_8. Its interesting property is the fact that the dimensions of the fundamental and adjoint representations are the same: 248 fermions and 248 gauge bosons. Unfortunately, there are thousands of scalar particles in the model. This circumstance has frightened theoreticians, and there is no detailed study of the E_8 model in the literature.

To summarize our survey of grand models, we must conclude that the idea of grand unification is undoubtedly very attractive and promising. There exist a number of indications that grand unification is actually realized in nature. On the other hand, there are many difficulties and unsolved problems. At the present time, we are not yet in a position to discriminate between different competing models. It is as if we were attempting to establish the appearance of a dinosaur from several bones. It is difficult to imagine that it will be possible to cope with this task if the decay of the proton is not detected. Of course, precise data on the neutral currents or on the neutrino masses are very important. But only the experimental discovery of nonconservation of baryon charge will transform grand unification from the hypothetical to the actual. To paraphrase the statement by M. Goldhaber, "If the proton is condemned to die, let it die in our hands, and quickly."

Among the most serious defects of the models of grand unification, the following two stand out: 1) these models do not include gravitation; 2) their Higgs sector, unlike the gauge sector, looks arbitrary and unnatural. There is hope that we will recover from both of these defects when we proceed to superunification, in the framework of extended supergravity.

10. SUPERUNIFICATION

The algebra of supersymmetry contains not only the ordinary generators of the Poincaré group—space-time displacements P_μ and rotations $M_{\mu\nu}$—but also spinor generators Q_α. In addition to the ordinary commutators, this algebra contains the anticommutator of the spinor generators (such algebras are called graded algebras): $\{Q_\alpha \bar{Q}_\beta\} = -2P_\mu(\gamma_\mu)_{\alpha\beta}$, where γ_μ are the Dirac matrices.

Supersymmetric multiplets contain both bosons and fermions. The simplest example is a multiplet containing the photon and a massless Majorana neutrino. The boson and fermion states in a multiplet are related by the spinor generators.

Supersymmetry opens up the unique possibility of combining internal symmetries with geometrical symmetries. In practice, this is accomplished by attracting an internal index i to the spinor generator: Q^i_α. What is obtained in this way is called extended supersymmetry.

Supersymmetry can be either global or local. An example of a model with extended global supersymmetry is the model containing 11 massless particles: one with $J=1$, four with $J=\frac{1}{2}$, and six with $J=0$ (a total of eight boson and eight fermion states with definite helicity). This model possesses global SU(4) symmetry.

Since invariance with respect to local coordinate transformations entails inclusion of the gravitational field, locally a supersymmetry theory must contain gravitons—massless particles with $J=2$. In addition, it must contain gravitinos—massless particles with spin 3/2. The simplest example of a locally supersymmetric theory (supergravity) is the theory containing one graviton and one gravitino. The simplest extended supergravity contains one graviton, two gravitinos, and one photon. This is so-called $N=2$ supergravity. By acting successively on the graviton and on the following components of the multiplet with the spinor generator Q^i_α, it is easy to see that the maximally extended supergravity not containing particles with $J>2$ corresponds to $N=8$. The multiplet of $N=8$ supergravity contains the following massless particles: one graviton, eight gravitinos, 28 bosons with $J=1$, 56 fermions with $J=\frac{1}{2}$, and 70 scalars (i.e., 128 boson and 128 fermion states with given helicity). This theory possesses SO(8) symmetry. There are very great expectations ("superexpectations") for extended supergravity: there is hope that the development of supergravity will lead to a unified theory of all interactions. The incorporation of all particles into a single supermultiplet fixes the dimensions of the various grand multiplets, and of the scalar multiplets in particular, and it also fixes the Yukawa constants.

It is well known that the loops of virtual fermions and bosons give contributions of opposite sign. In supersymmetry, these contributions cancel. Thus, for example, it has recently been shown that in the above-mentioned SU(4) model with an arbitrary additional gauge invariance the Gell-Mann–Low function is equal to zero not only in the one-loop and two-loop approximations, but also in the three-loop approximation. Such a cancellation might also occur for the masses of the scalar particles, ensuring the hierarchy of scales. If the cancellation occurs in all orders, the theory is conformally invariant.

There are also manifestations of fermion–boson cancellations in supergravity. Thus, for ordinary quantum gravity (without matter) it has been proved that there is no divergence in the single-loop approximation. At the same time, for $N=1$ supergravity there are also no divergences in the two-loop approximation.

A very interesting physical quantity for which these cancellations may be vitally important is the so-called cosmological term λ, which describes the gravitational effect of the vacuum. The λ term must occur because of the quantum fluctuations in the vacuum. On the basis of naive dimensional arguments, one might expect that $\lambda \sim m_P^4 \sim 10^{78}$ GeV^4. But this energy density is inconceivably great: as if all the nucleons of the Universe were contained in each cell of space having the nucleon Compton volume. Even if such a volume contained the mass of only a single nucleon, we would still obtain an unacceptably large value $\lambda \sim m_p^4 \sim 1$ GeV^4. But observational astronomical data indicate that $\lambda < 10^{-47}$ GeV^4 (which corresponds to about one proton mass per cubic meter of the vacuum). This means that there is a superfine cancellation between the vacuum fluctuations of different fields, and it is only because of this that the vacuum looks like empty space. As I. Ya. Pomeranchuk said, "The vacuum is full of deep physical content."

Supergravity is still in its infancy. It was only recently that $N=2$ supergravity with all the auxiliary fields was constructed. Larger values of N still await their turn. The problem of symmetry breaking in supergravity is also very far from its solution.

When the group SO(N) with charge e is gauged in N-supergravity, there is a gravitino mass of order em_P and a cosmological term of order $e^2m_P^4$. Symmetry breaking must occur in such a way as to completely compensate this enormous term.

In the construction of supergravitational models, an important role is played by spaces of higher dimensions, up to $d=11$, which include our four-dimensional space as a subspace. It is possible that the additional seven dimensions will one day become just as physical as the quarks are today. We also mention here an interesting attempt to combine into a single space not only the ordinary coordinates x_μ and the anticommuting spinor variables θ_α^i (which form so-called superspace), but also the physical fields belonging to the supermultiplet.

The edifice of contemporary physics is being constructed simultaneously at different levels, and we turn to a discussion of phenomenological applications of the as yet nonexistent broken extended supergravity. One of the serious difficulties here is that even the largest of the groups—SO(8)—is too tight to accommodate the group SU(3) × SU(2) × U(1) of the known interactions, with all the necessary gauge and fermion fields. Two possible ways of overcoming this difficulty are being discussed:

1. To turn to $N>8$ by including fields with spins $J>2$.

2. To find and use some hidden symmetry of the theory with $N=8$.

The first approach has so far had no real success. But this approach cannot be easy. In a sense, the progress of the theory of elementary particles can be regarded as an ascent (and at times a descent, for example, from $J=2$ to $J=3/2$) on the ladder of spins discussed at the beginning of this paper. Those steps which we have already mastered do not seem difficult to us. But it required truly colossal techniques to master each of them.

In essence, all the dizzy successes of non-Abelian gauge theories reduce to the fact that we have learned (or think we have learned?) to work not only with vector long-range interactions (photons), but also with vector short-range interactions (gluons, W, Z, and H bosons). Clearly, a working quantum field theory with $J>2$ would be a more serious achievement than all the previously created theories.

In the second approach, an interesting development has occurred in the past two or three years. It was discovered that $N=8$ supergravity with global SO(8) symmetry contains a hidden local nonlinear SU(8) symmetry. The next step was the assumption that owing to quantum corrections the auxiliary vector fields acquire kinetic terms and begin to propagate. (This effect was discovered in the two-dimensional nonlinear CP^{n-1} model.) In essence, it was assumed that the 63 gauge fields of SU(8) are bound states of the original fundamental fields belonging to the multiplet of $N=8$ supergravity. The group SU(8) is sufficiently large to accommodate color, weak isospin, and the generations of fermions, so that by assumption the SU(8) supermultiplet of bound states contains not only our gauge fields, but also quarks, leptons, and scalar bosons. All these are now composite. Moreover, it also contains a large number of other bound states, in particular, with $J>1$.

The concept of a renormalizable subset of particles is very important for what follows. The SU(8) theory described above is nonrenormalizable, and the characteristic scale in it is the Planck mass m_P. Therefore the masses of the particles in this theory must in general be of order m_P. An exception is the subset of particles with $J \leqslant 1$, the interaction between which is renormalizable and which as a result may remain massless. In essence, we are dealing here with dimensional arguments. Nonrenormalizable interactions have a power growth up to momenta of order m_P, and with a renormalizable logarithmic dependence on the momentum the contribution of small distances to the quantum corrections is small. If a particle is to "survive" and "escape" from the region m_P and remain massless and dynamical at low energies, it must, together with certain other "selected" particles, confine its interactions to those described by an effective renormalizable gauge Lagrangian. All other particles become superheavy (of order m_P) because of their nonrenormalizable strong interactions.

The concept of a "group of escaped particles" having a renormalizable gauge interaction is more general than the SU(8) model discussed above. It may be the explanation of the important role played by renormalizable gauge interactions in Nature. Using this concept, one can attempt, in particular, to go beyond $N=8$ and consider supergravitational schemes with $N>8$ containing particles with $J>2$. Nevertheless, these particles will have masses of order m_P and will not manifest themselves at large distances.

The attempt to distinguish a set of "escaped particles" in the case of the SU(8) theory led to a very interesting and instructive result. With certain additional assumptions, it turned out that these particles and the interactions between them are described by the gauge SU(5) symmetry with its 24 vector bosons. Then—and this is very important—there exist only three fermion generations with multiplets $\bar{5}+10$ in each, and a minimal set of scalar bosons: $5+\bar{5}+24$. Thus, the mountain of extended supergravity gave birth to the mouse of the minimal SU(5) model. If this picture is correct, there should not be higher grand symmetries of the type SO(10), E_6, etc.; there should not be technicolor; there should not be supersymmetric particles accessible to observation: the gravitino gluino, and photino. But there should exist proton decay with a lifetime not exceeding 10^{33} yr. There should exist a "gauge desert," and a new strong interaction should manifest itself only near m_P.

Most physicists do not believe in the prospect of a gauge desert. We are accustomed to the fact that hitherto each new advance along the scale of energies has revealed new physical phenomena. But if the clouds around the Earth were so dense that only our generation, having broken through them, saw the stars for the first time, would we be prepared to believe that there is nothing but emptiness between us and the nearest star?

To conclude the discussion of supergravity, it is appropriate to make a few remarks about the gravitational interaction proper and the constant G_N. There have been a number of attempts to construct a renormalizable theory of gravitation. Here G_N is not regarded as a fundamental constant, but arises as a result of spontaneous symmetry breaking. If we succeeded in constructing such a theory, the gravitational interaction in it would not be strong not only below m_P, but also in the post-Planck region.

On the other hand, we cannot take it to be ruled out that the Planck mass is simply a mirage, just as the classical electron radius αm_e^{-1} was in its time. It is possible that the gravitational interaction is modified at not very small distances. Indeed, Newton's law has been verified experimentally only at centimeter distances, and then not with very high accuracy. (In general, Newton's constant G_N is known only with an accuracy of four significant figures.) The above-mentioned modification does not seem to me a plausible possibility, but it cannot be excluded.

11. ASTROPHYSICS AND COSMOLOGY

Ten or twenty years ago, it seemed that of the three basic elements of physics—observations, experiment, and theory—the first belonged entirely to history. The role of astronomy in establishing new physical laws, which was so decisive in Newton's time, seemed to come to nothing. Recently, however, the position has changed sharply. Astrophysical and cosmological data are playing an ever increasing role in fundamental physics, especially in estimating the viability of particular physical models of the structure of particles. On the other hand, physicists are claiming more and more to calculate, on the basis of the edifice of the properties of elementary particles, those properties of the Universe which have traditionally been considered as "sacred initial conditions." The claims of physics are growing.

Contemporary cosmology is the basis of the theory of the hot Universe, which achieved almost universal recognition after the discovery of the relic heat radiation. We know that the younger the Universe was, the hotter it was. From the fact that the theory of the primeval nucleosynthesis is in good agreement with observational data on the abundance of the elements, it follows that we can confidently describe the behavior of the Universe in the first seconds of its existence, when the temperature was measured in MeV. At the present time, we see no obstacles to the extrapolation of our description to higher temperatures, up to m_P. The behavior of the Universe at these temperatures is determined by the behavior of elementary particles at those small distances at which the dynamical properties of theories of superunification and grand unification should manifest themselves. Thus, the early Universe is, as it were, a natural laboratory for testing these theories.

Among the basic parameters characterizing the Universe, there are several which have recently been attracting particular attention. We have already spoken about the surprisingly small value (or vanishing) of the cosmological constant λ. Another no less surprising fact is that the observed mean density ρ of matter in the Universe is close to the critical density ρ_c. To explain the meaning of ρ_c, let us mentally cut out a sphere in the Universe and consider the nonrelativistic motion of a test mass on the surface of the sphere. For $\rho=\rho_c$, we have the so-called flat case, when the kinetic energy

$T_{\rm kin}$ and the potential energy U of the test body are equal. For $\rho > \rho_c$, the expansion of the Universe must be replaced by a contraction, since $|U| > T_{\rm kin}$; for $\rho < \rho_c$ $|U| < T_{\rm kin}$ and the expansion is unlimited. (These three cases are similar to the behavior of a thrown stone, when its velocity is equal to, less than, or greater than the second cosmic velocity.) Considering the fact that at the initial instant both $T_{\rm kin}$ and U were very large, while the sum $T_{\rm kin} + U$ was the same as it is today, we see that in this case there is a cancellation between two large quantities which is just as surprising as in the case of the cosmological constant.

We turn now to the problem of the baryon asymmetry of the Universe. One of the most profound ideas put forward in the mid-sixties is that this asymmetry arose as a result of nonconservation of baryon number and the violation of CP and C invariance. The baryon asymmetry of the Universe is characterized by the ratio of the mean density of baryons n_B to the mean density of photons n_γ. Observations give $n_B/n_\gamma \approx 10^{-9}$. In grand models, the main contribution to the baryon asymmetry should come from decays of superheavy Higgs bosons. These bosons live longer than the gauge bosons X and Y and decay in an epoch when the temperature of the Universe is lower than its mass ($T \ll M$), so that the excess of baryons arising in the decay of these bosons does not vanish. In the SU(5) model, the observed value of n_B/n_γ can be obtained only with two sets of scalar quintets: $5+\bar{5}$. To obtain the asymmetry in the SO(10) model, the left-right symmetry inherent in the model must be broken very early: the masses of the right-handed bosons (related to right-handed currents) must be very large: $m_{W_R} \gtrsim 10^{11}$ GeV. Otherwise, the C-invariant interactions "burn up" the excess of baryons.

Even in the absence of interactions which violate the conservation of baryon charge, a baryon excess might arise as a result of CP- and C-nonsymmetric evaporation of black holes. However, the quantitative estimates here are less reliable.

Recently, there has appeared a series of papers devoted to the cosmological fate of magnetic monopoles. On the one hand, as is well known, monopoles have not been observed in nature. On the other hand, grand models such as SU(5), SO(10), and E_6 contain classical monopole solutions, the masses of these monopoles being about two orders of magnitude greater than the grand-unification mass, i.e., of order 10^{16} GeV. Naive estimates of the production and extinction of monpoles made several years ago gave an unacceptably large concentration of relic monopoles and thereby cast doubt on the validity of the above-mentioned grand models. Very recently, a way out of this difficulty may have been found. To explain its essence, we recall that according to the theory of the electroweak interaction the photon is a linear superposition of two fields—an isovector field W^0 and an isoscalar field B^0: $A = B^0 \cos\theta_W + W^0 \sin\theta_W$, where θ_W is the weak angle. It is known that a pair of B monopoles of opposite sign cannot break up because of the phenomenon of magnetic confinement for $T < \eta \approx \frac{1}{4}$ TeV, since the flux of magnetic lines of force, which is conserved for $T < \eta$, is compressed into a tube by the scalar condensate. A new observation consists in the fact that W^0 monopoles must also be subject to confinement, but in a hot plasma rather than in a cold phase, for $T > \eta$. If this is correct, then both components of the electromagnetic monopoles must annihilate.

Astrophysics gives very rigid constraints on the possible properties of various hypothetical particles, and in particular on the properties of heavy neutral leptons. An analysis of the extinction of hypothetical free quarks gives such unacceptably large values for the present concentration of these particles that we must conclude that free fractionally charged quarks do not exist.

A soft spontaneous breakdown of CP invariance, if it were realized in nature, would lead, in the process of cooling of the Universe, to the production of a domain structure of the vacuum. The CP-odd condensate would in general have different signs in causally unconnected regions. The moving domain walls would strongly perturb the isotropy of the relic radiation. Observations do not reveal such a perturbation, and this is an argument that CP invariance is strictly violated in nature.

The existence of galactic magnetic fields with dimensions of order 10^{22} cm enables us to conclude that the Compton wavelength of the photon is not smaller than these dimensions and hence that its mass does not exceed 10^{-27} eV. It can be shown that the fact that the photon is practically massless, in its turn, rules out the possibility of processes involving nonconservation of electric charge, such as $e^- \to \nu\gamma$ or $\mathrm{Ga} \to \mathrm{Ge} + \gamma$.

But perhaps astrophysical observations give us more unique information about the neutrino than about any other particle. The observed abundance of ^{4}He limits the number N_ν of different types of massless or light ($m_\nu \ll 1$ MeV) neutrinos. This sensitivity of the abundance of helium to the value of N_ν is due to the fact that at the time when the present ratio of abundances of neutrons and protons was formed it was the neutrinos (together with photons) that determined the energy density in the Universe and, consequently, the rate of its expansion. One can often find the estimate $N_\nu \leqslant 4$ in the literature, but some authors are more careful and prefer $N_\nu \leqslant 4-6$.

The data on the Hubble recession of the galaxies, in conjunction with independent data on the age of the Earth, Sun, and stars, restrict the density of matter in the Universe: $\rho \lesssim \rho_c$. On the other hand, the theory of the hot Universe enables us to express the number of relic neutrinos, whose direct observation is practically impossible, in terms of the known number of relic photons. As a result, it is possible to find an upper limit for the sum of the masses of the various types of neutrinos:

$$\sum_{i=1}^{N_\nu} m_{\nu_i} \leqslant \begin{cases} 30-40\ \text{eV}, & \text{if}\ \lambda = 0, \\ 100-200\ \text{eV}, & \text{if}\ \lambda < 0. \end{cases}$$

We recall that the cosmological term λ characterizes the gravitational effect of the vacuum. If it turned out, for example, that $\Sigma m_{\nu_i} \sim 100$ eV, this would mean that $\lambda < 0$ and hence that the vacuum antigravitates.

Astrophysical observations provide various evidence not only for an upper limit on the neutrino masses but also for a lower limit, although this evidence is less certain.

An argument in favor of a nonzero neutrino mass is the existence in the Universe of the so-called hidden mass in the invisible halos surrounding the galaxies and clusters of galaxies. (The existence of these halos shows up in the analysis of the distribution of stellar velocities.) According to recent estimates, the hidden mass is approximately 30 times as large as the visible mass of these objects. So far, it is completely unproved that this hidden mass actually belongs to neutrinos. For example, it might exist in the form of cold stars of the type of Jupiter. However, the observed abundance of deuterium indicates that the hidden particles are not nucleons.

In general, cosmologists would welcome neutrino masses of the order of 10–20 eV, although they emphasize that their arguments do not claim to be proofs. In particular, neutrinos with such masses provide a natural explanation of the manner of formation of the largest structural cells of the Universe—the superclusters (the number of which in the observable part of the Universe is of the order of 10^6). For $T \lesssim m \approx 10$ eV, the cold neutrinos become gravitationally unstable and begin to form neutrino clouds which later, for $T < 1$ eV, begin to collect atomic dust. The density of neutrino clouds is $\rho \sim T^4 \sim m_\nu^4$, and their characteristic dimensions are of the order of the time t for their formation (recall that $c = 1$), which in the theory of the hot Universe is determined by the relation $t \approx m_p T^{-2} \sim m_p m_\nu^{-2}$. As a result, the volume of a supercluster is $V \sim t^3 \sim m_p^3 m_\nu^{-6}$, and we obtain a surprisingly simple and elegant relation which expresses the masses of the superclusters—the largest objects in the Universe—in terms of the masses of the neutrinos—the lightest elementary particles:

$$M \sim V\rho \sim m_p^3 m_\nu^{-2},$$

which for $m_\nu \sim 10$ eV gives a value close to the observed value.

It is possible that the existence of neutrino oscillations is indicated by the results of the long hunt for solar neutrinos. To detect these neutrinos, use is made of the reaction $\nu_e + {}^{37}\mathrm{Cl} \to e^- + {}^{37}\mathrm{Ar}$, whose yield is measured in solar neutrino units (1 SNU = 1 capture per second per 10^{36} nuclei). The expected effect should have been 7–8 SNU, whereas the measurements gave 2.2 ± 0.4. One of the possible explanations of this discrepancy is that because of oscillations the solar electron neutrinos near the Earth are partially transformed into an inert state, for example, ν_τ. It is not excluded, however, that the reason for the discrepancy is the inadequacy of the model for the internal structure of the Sun. This problem might be solved by measurements of the flux of soft solar neutrinos which are to be made by means of a gallium detector. We shall not consider many other phenomena discussed in the literature: supernova explosions, black holes, and the possible existence of an unstable (false) vacuum. Today's physics is indeed headed for the sky.

12. NEW STABLE PARTICLES? NEW LONG-RANGE FORCES?

So far, in speaking of future experiments, we have had in mind the search for those particles and phenomena whose existence is to be expected to some extent on the basis of known facts and theoretical constructions. In contrast, in the present section we shall discuss those hypothetical particles and interactions for whose possible existence there is no, even indirect, evidence. Nobody needs them at the present time. It would be incorrect, however, to confine the role of experiment to tests of existing theories. Indeed, nobody needed the muon in 1937 or the τ lepton in 1975. The discovery of a CP-noninvariant interaction was also completely unexpected in 1964. These examples can easily be multiplied. We can have no doubt that experimentalists will continue to make "unplanned" discoveries, following Galileo's motto: "Measure everything that can be measured, and make phenomena that are inaccessible to measurement accessible to it."

We begin with a discussion of the search for new stable particles. A lower limit $M \lesssim 17$ GeV for the masses of new charged particles comes from experiments using the colliding beams of PETRA, which did not reveal production of new particles up to $\sqrt{s} \approx 35$ GeV. Mass-spectroscopic analysis of stable matter (water) has been made up to $M \lesssim 300$ GeV. The object of the search was anomalously heavy "exotic hydrogen," for the enrichment with which water was subjected to electrolysis. A lower limit was obtained for the ratio of the concentrations of "exotic" and ordinary hydrogen in water: $n_x/n_p < 10^{-21}$. It seems to me very important to extend the range of masses in every possible way and to raise the sensitivity of these searches. It would also be very interesting to use other substances instead of water, in particular, to search for anomalously heavy atoms in meteorites, in heavy minerals, and in the products of active volcanos. A simple estimate shows that if the mass of the particles under consideration is less than about 10^{15}–10^{16} GeV, atoms containing these particles are still not attracted strongly enough by the Earth for this attraction to break their chemical bonds. Heavier particles, in particular, so-called maximons with $m \sim m_P$, must fall to the center of the Earth.

The discovery of a "deposit" of negatively charged heavy particles might present more than just scientific interest. It might completely revolutionize energy production. The point is that such particles might be used for bold catalysis of fusion reactions such as $\mathrm{d} + \mathrm{t} + \mathrm{X}^- \to \mathrm{He}^4 + \mathrm{n} + \mathrm{X}^- + 17$ MeV, where X^- is a heavy stable particle. Unfortunately, the efficiency of such catalysis is limited by the adherence of X^- to He^4. The breakdown of $\mathrm{X}^-\mathrm{He}^4$ "ions" and the regeneration of X^- require time, energy, and complex technology, so that with the small quantities of X^- particles which might be obtained using the most intense accelerators of the future "the porridge won't cook": the energy obtained would be insufficient for an electric stove. If, however, we succeeded in finding a deposit of at least several kilograms of X^- particles, this might solve all the energy problems of mankind.

Heavy stable particles with mass ≈ 1 TeV are predicted by the technicolor model. But it seems to me that new particles should be sought over the entire accessible range of masses, without adhering to particular theoretical schemes. The heavier a particle, the easier it is to detect it in a spectrometer, both from the anomalously large time of flight and from the anomalously small deflection in a magnetic field.

The search for fractionally charged particles is also very interesting. The observation of fractional charges in niobium balls has been reported. Recently, a possible reason for the appearance of a false effect in such experiments has been suggested. The search for fractional charges is of great interest. It should be noted that fractional charges may be a property not only of colored free quarks (in whose existence I do not believe), but also of colorless particles—fractons. (One of the varieties of fractons is the so-called hydrons, which are hadrons in which one of the quarks is replaced by a neutral particle belonging to a color triplet of hypothetical Higgs bosons.)

Finally, we cannot exclude the existence of particles which, while having only the gravitational interaction with ordinary matter, can nevertheless interact quite strongly with one another, forming specific y-matter. Recently, there have been discussions of the hypothesis that the oscillations of the Sun with a period of 160 min, whose existence was reported several years ago by astronomers from the Crimean Observatory and subsequently by astronomers from Stanford, may be due to a y-planet with mass of the order of the mass of the Moon, moving at a depth of order 20 000 km from the surface of the Sun.

In connection with the observed perturbations in the motion of Neptune, Pluto, and Halley's comet, at the beginning of the seventies there were discussions of the hypothesis that these perturbations are due to a hypothetical tenth planet with the mass of Jupiter, situated at a distance of 60 astronomical units from the Sun. A search for this planet gave a negative result. Moreover, subsequent refined calculations showed that such a planet, even if it existed, would not lead to agreement between the observations and the calculations. This example shows, however, that massive, hitherto unobserved bodies may exist in the solar system. It would be interesting to undertake a search for invisible planets and asteroids in the solar system by means of cosmic probes.

Gravimetric experiments are also of interest in connection with other questions, namely, the search for new long-range forces. Here we are concerned with laboratory experiments such as the Eötvös experiment. Precise experiments to verify the equality of the inertial and gravitational masses rule out Coulomb-like or Newton-like forces between leptons and/or nucleons, even if these forces are ten orders of magnitude weaker than the gravitational attraction between these particles. But nowadays it does not seem improbable that with further improvement in the accuracy of experiments new long-range forces may be discovered. One of the important lessons of recent years was that we should not be afraid of very large (or very small) numbers. Until recently, large numbers were encountered only in cosmology (the number of nucleons, the age of the Universe). In the physics of weak, strong, and electromagnetic interactions, the largest parameter was perhaps m_p/m_e and, for example, a neutrino mass of order 10 eV was considered "esthetically unacceptable." Now, with the advent of grand models (and particularly supermodels with their scale m_P), our psychology has changed. It no longer seems strange to us to have a light (but not massless) neutrino or a long-lived (but decaying) proton. Small coupling constants should also not seem improbable to us.

In speaking of long-range forces, we have had in mind potentials of the type $1/r$, which correspond to massless particles. There are much weaker restrictions on hypothetical interactions of the type $e^{-\mu r}/r$ with a finite but large range, say, of the order of a kilometer or a centimeter. To search for such forces, experiments such as the Cavendish experiment are preferred over experiments such as the Eötvös experiment. It would be particularly interesting to seek long-range forces which grow with increasing energy of the colliding particles (or of the colliding systems of particles, for example, nuclei or lumps of matter) more rapidly than the gravitational forces. Such forces would correspond to the exchange of very light particles with $J > 2$.

It must be borne in mind that the particles which are known to us may appear "neutral" with respect to new long-range forces, possessing not "charges" but only corresponding dipole "magnetic" moments, if their spin is $\frac{1}{2}$. (For $J > \frac{1}{2}$, they might have higher "electric" and "magnetic" multipole moments. If a new interaction is CP-noninvariant, the particles must also have "electric" dipole moments.)

Among the possible forces with respect to which the particles known to us are neutral, forces corresponding to non-Abelian gauge fields with a large (even macroscopic) confinement radius would have particularly interesting properties. If we look at the map of the fundamental forces given earlier (see Fig. 2), the idea of a possible macroscopic confinement radius does not seem strange. Indeed, it can be seen on this map (see the dashed line) that if the W and Z bosons were massless, the confinement radius for the ordinary weak forces would be measured in millimeters.

Suppose for a moment that, in addition to the known gauge symmetries, there is one further symmetry $SU(2)_\theta$, whose gauge fields are three θ gluons. The particles that are known to us do not have θ charges—they are θ-neutral; but there may exist heavier particles which carry θ charges. In principle, these particles may also have the ordinary electroweak and strong interactions, and we would be able to create them using accelerators when we exceed the threshold for their production. If the generally accepted ideas about confinement are valid, a pair of θ particles produced in an accelerator and having opposite θ charges must be connected by a θ-gluon string. The thickness of this θ string must be of the order of the confinement radius R_θ, the specific linear density is of order R_θ^{-2}, and its length L is limited only by the available energy: $L = ER_\theta^2$, where E is the kinetic energy of the pair of θ particles. If we assume that $R_\theta \approx 10^{-6}$ cm and take $E \approx 10$ GeV, then L is of the order of a meter. The particles at the ends of the string might, as a result of the electromagnetic interaction (if they have negative electric charge) or the strong interaction (if their charge is positive or zero), be bound to the nuclei of ordinary atoms. Moreover, if the tension of the string is less than the force of chemical binding between the atoms (of order eV/10^{-8} cm), then the atoms at the ends of the string might adhere chemically to ordinary matter (for this, it is necessary that $R_\theta \gg 10^{-6}$ cm).

There cannot be a very thick θ string, since for $R_\theta \gtrsim 1$ cm it must dissolve in a gas of relic θ gluons whose temperature is of order 3° K.

The properties of θ strings are striking: they must be absolutely permanent, they can be stretched indefinitely, and they can completely freely cut through walls, mountains, and even the Earth!

13. CONCLUDING REMARKS

Elementary-particle physics has in its ever-expanding arsenal the most varied tools: nuclear reactors, lasers, mass spectrometers, optical, radio, and neutrino telescopes, and even gravimeters. Nevertheless, accelerators have been and will no doubt remain the basis of this science. It is difficult to imagine a future growth without a growth of the energy of the accelerators. What are the limits to this growth?

About 30 years ago, Fermi spoke of a future accelerator encircling the Earth. Embarking on a science-fiction journey, let us imagine a gigantic ring around the Earth with radius $\rho \approx 7000$ km. At the height of order 1000 km at which it is situated, there is a good vacuum ($\lesssim 10^5$ particles/cm^3) and an abundance of solar energy (≈ 1.4 kV/m^2). If we fill this entire ring with magnets having $H \approx 50$ kG, the energy of the protons in the ring will be of order 10^7 GeV. We recall that

$$\left(\frac{pc}{\text{GeV}}\right) = 30\left(\frac{H}{\text{kG}}\right)\left(\frac{\rho}{\text{km}}\right).$$

An energy of 10^7 GeV in the center-of-mass system of colliding protons corresponds in the case of a stationary target to an energy of the order of 10^{23} eV in the laboratory system. Such energies have hitherto not been observed even in cosmic rays. Is it possible to go even higher by increasing the magnetic field or the radius of the accelerator, bringing it, say, to 40 000 km—the radius of a geostationary orbit? Unfortunately, owing to synchrotron radiation, it is not possible to go much above 10^8 GeV. The energy radiated in one revolution is

$$\left(\frac{E_{\text{rad}}}{\text{GeV}}\right) = 8\left(\frac{pc}{10^7\,\text{GeV}}\right)^4\left(\frac{1000\text{ km}}{\rho}\right)$$

or

$$\frac{E_{\text{rad}}}{pc} = \frac{4\pi}{3}\frac{\alpha\gamma^3}{m\rho},$$

where m is the mass of the proton and $\gamma \approx pc/m$ is its Lorentz factor. For $pc \approx 10^9$ GeV, the entire energy of the proton is radiated in several revolutions. For electrons, the energy ceiling is at about 10^4 GeV.

There is another, more realistic approach than the creation of gigantic cosmotrons, namely, the creation of comparatively small accelerators with high field intensities. Thus, the design for colliding linear electron–positron beams which is now being discussed is based on a rate of acceleration MeV/cm. This gives an energy of order 10^3 GeV for an accelerator length of order 10 km. There have been published proposals to use beams of large proton accelerators to accelerate electrons and unstable particles with an acceleration rate of order 30 MeV/cm. Some physicists pin their hopes on laser methods of acceleration. Progress is required not only to raise the energy of accelerators, but also to raise their luminosity. This last circumstance is particularly important if we take into account the fact that the cross sections for the most interesting processes fall off quadratically with increasing energy.

As an objection to the construction of more and more powerful accelerators, one may say that even 10^7 GeV is still 12 orders of magnitude smaller than the Planck mass, which seems to be the natural energy scale of physics and which will remain just as inaccessible after hundreds of years. However, it must be borne in mind that the frontier of knowledge does not expand uniformly as we go up the scale of energies, and in ascending a new hill we may see completely new lands and vistas.

Let us return from the distant future to the next decades. Our immediate goal is energies of order 0.1 TeV in the center-of-mass system of the colliding leptons and/or quarks. It is to be hoped that this region will be studied in sufficient detail towards the end of the eighties. It is technologically feasible to conquer the next region of energies even before the end of this century by constructing colliding proton–antiproton beams with energy equal to several tens of TeV. If the development of linear electron accelerators with a high acceleration rate proves to be successful, such installations might be supplemented by colliding electron–positron beams with energies equal to several TeV. Everything that we now know on the basis of both theoretical extrapolations and exploratory data obtained in cosmic rays guarantees that there are many interesting physical phenomena awaiting investigation in this region. It is this region that contains the answers to many of the questions discussed above. It undoubtedly also conceals a multitude of even more profound questions.

BIBLIOGRAPHY

Basic principles

C. N. Yang, "Einstein's impact on theoretical physics," Phys. Today **33**, No. 6, 42 (1980) [Russ. transl., Usp. Fiz. Nauk **132**, 169 (1980)].

General Relativity: An Einstein Centenary Survey (ed. S. W. Hawking and W. Israel), Cambridge, 1979, Chap. 1 [Russ. transl., Usp. Fiz. Nauk **133**, 139 (1981)].

P. G. Bergmann, "Unitary field theories," Phys. Today **32**, No. 3, 44 (1979) [Russ. transl., Usp. Fiz. Nauk **132**, 177 (1980)].

Nobel lectures in physics, 1979: S. Weinberg, "Conceptual foundations of the unified theory of weak and electromagnetic interactions," Rev. Mod. Phys. **52**, 515 (1980) [Russ. transl., Usp. Fiz. Nauk **132**, 201 (1980)]; S. L. Glashow, "Towards a unified theory—threads in a tapestry," Rev. Mod. Phys. **52**, 529 (1980) [Russ. transl., Usp. Fiz. Nauk **132**, 219 (1980)]; A. Salam, "Gauge unification of fundamental forces," Rev. Mod. Phys. **52**, 525 (1980) [Russ. transl., Usp. Fiz. Nauk **142**, 229 (1980)].

J. Ilipoulos, "An introduction to gauge theories," Preprint 76-11, CERN, Geneva, 1976 [Russ. transl., Usp. Fiz. Nauk **123**, 565 (1977)].

B. A. Arbuzov and A. A. Logunov, "Structure of elementary particles and relationships between the different forces of nature," Usp. Fiz. Nauk **123**, 505 (1977) [Sov. Phys. Usp. **20**, 956 (1977)].

V. B. Berestetskiĭ, "Zero-change and asymptotic freedom," Usp. Fiz. Nauk **120**, 439 (1976) [Sov. Phys. Usp. **19**, 934 (1976)].

A. A. Vladimirov and D. V. Shirkov, "Renormalization group and ultraviolet asymptotics," Usp. Fiz. Nauk **129**, 407 (1979) [Sov. Phys. Usp. **22**, 860 (1979)].

Fundamental particles

Nobel lectures in physics, 1976: B. Richter, "From the psi to charm—the experiments of 1975 and 1976," Czech. J. Phys. **A28**, 323 (1978) [Russ. transl., Usp. Fiz. Nauk **125**, 201 (1978)]; S. C. C. Ting, "Discovery of the J particle: A personal recollection," Czech. J. Phys. **A28**, 124 (1978) [Russ. transl., Usp. Fiz. Nauk **125**, 227 (1978)].

M. L. Perl, "The τ heavy lepton: a recently discovered elementary particle," Preprint SLAC-PUB-2153, Stanford University, 1978 [Russ. transl., Usp. Fiz. Nauk **129**, 671 (1979)].

Ya. I. Azimov, L. L. Frankfurt, and V. A. Khoze, "New Particle in e^+e^- annihilation: the heavy lepton $\tau^{\pm}$," Usp. Fiz. Nauk **124**, 459 (1978) [Sov. Phys. Usp. **21**, 225 (1978)].

J. Kirkby, "Review of e^+e^- reactions in the energy range 3 to 9 GeV," Invited talk at the 9th Intern. Symposium on Lepton and Photon Interactions at High Energies, Batavia, 1979 SLAC-PUB-2419, October 1979 [Russ. transl., Usp. Fiz. Nauk **122**, 309 (1981)].

Ya. I. Azimov and V. A. Khoze, "Contemporary status of the τ lepton," Usp. Fiz. Nauk **132**, 379 (1980) [Sov. Phys. Usp. **23**, 699 (1980)].

L. M. Lederman, "The upsilon particle," Sci. Am. **239**, No. 4, 60 (1978) [Russ. transl., Usp. Fiz. Nauk **128**, 693 (1979)].

S. L. Glashow, "Quarks with color and flavor," Sci. Am. **233**, No. 4, 38 (1975) [Russ. transl., Usp. Fiz. Nauk **119**, 715 (1976)].

D. B. Bline, A. K. Mann, and C. Rubbia, "The search for new families of elementary particles," Sci. Am. **234**, No. 1, 44 (1976) [Russ. transl., Usp. Fiz. Nauk **120**, 113 (1976)].

Electromagnetic interaction

J. H. Field, E. Picasso, and F. Combley, "Tests of fundamental physical theories from measurements of free charged leptons," Preprint, CERN, Geneva, 1978 [Russ. transl., Usp. Fiz. Nauk **127**, 553 (1979)].

S. D. Drell, "Experimental status of quantum electrodynamics," Physica **96A**, 3 (1979) [Russ. transl., Usp. Fiz. Nauk **130**, 507 (1980)].

Strong interaction

A. I. Vaĭnshteĭn, M. B. Voloshin, V. I. Zakharov, V. A. Novikov, L. B. Okun', and M. A. Shifman, "Charmonium and quantum electrodynamics," Usp. Fiz. Nauk **123**, 217 (1977) [Sov. Phys. Usp. **20**, 796 (1977)].

I. M. Dremin, "Gluon jets," Usp. Fiz. Nauk **131**, 715 (1980) [Sov. Phys. Usp. **23**, 515 (1980)].

Ya. I. Azimov, Yu. L. Dokshitser, and V. A. Khoze, "Gluons," Usp. Fiz. Nauk **132**, 443 (1980) [Sov. Phys. Usp. **23**, 732 (1980)].

M. Jacob and P. V. Landshoff, "The inner structure of the proton," Sci. Am. **242**, No. 3, 46 (1980) [Russ. transl., Usp. Fiz. Nauk **133**, 505 (1981)].

J. D. Bjorken and B. J. Ioffe, "e^+e^- annihilation into hadrons," Usp. Fiz. Nauk **116**, 115 (1975) [Sov. Phys. Usp. **18**, 361 (1975)].

V. G. Grishin, "Inclusive processes in high-energy hadron interactions," Usp. Fiz. Nauk **127**, 51 (1979) [Sov. Phys. Usp. **22**, 1 (1979)].

I. M. Dremin and C. Quigg, "Clusters in hadron multiple production processes," Usp. Fiz. Nauk **124**, 535 (1978) [Sov. Phys. Usp. **21**, 265 (1978)].

Y. Nambu, "The confinement of quarks," Sci. Am. **235**, No. 5, 48 (1976) [Russ. transl., Usp. Fiz. Nauk **124**, 147 (1978)].

Weak interaction

S. M. Bilen'kiĭ, and B. M. Pontecorvo, "Lepton mixing and neutrino oscillations," Usp. Fiz. Nauk **123**, 181 (1977) [Sov. Phys. Usp. **20**, 776 (1977)].

V. M. Shekhter, "Weak interactions involving neutral currents," Usp. Fiz. Nauk **119**, 593 (1976) [Sov. Phys. Usp. **19**, 645 (1976)].

P. F. Ermolov and A. I. Mukhin, "Neutrino experiments at high energies," Usp. Fiz. Nauk **124**, 385 (1978) [Sov. Phys. Usp. **21**, 185 (1978)].

V. A. Alekseev, Ya. B. Zel'dovich, and I. I. Sobel'man, "Parity nonconservation effects in atoms," Usp. Fiz. Nauk **118**, 385 (1976) [Sov. Phys. Usp. **19**, 207 (1976)].

A. N. Moskalev, R. M. Ryndin, and I. B. Khriplovich, "Possible lines of research into weak-interaction effects in atomic physics," Usp. Fiz. Nauk **118**, 409 (1976) [Sov. Phys. Usp. **19**, 220 (1976)].

G. V. Danilyan, "Parity violation in nuclear fission," Usp. Fiz. Nauk **131**, 329 (1980) [Sov. Phys. Usp. **23**, 323 (1980)].

L. M. Barkov, M. S. Zolotarev, and I. B. Khriplovich, "The observation of parity nonconservation in atoms," Usp. Fiz. Nauk **132**, 409 (1980) [Sov. Phys. Usp. **23**, 713 (1980)].

Scalars

A. I. Vaĭnshteĭn, V. I. Zakharov, and M. A. Shifman, "Higgs particles," Usp. Fiz. Nauk **131**, 537 (1980) [Sov. Phys. Usp. **23**, 429 (1980)].

Grand unification

M. Gell-Mann, P. Ramond, and R. Slansky, "Color embeddings, charge assignments, and proton stability in unified gauge theories," Rev. Mod. Phys. **50**, 721 (1978) [Russ. transl., Usp. Fiz. Nauk **130**, 459 (1980)].

S. G. Matinyan, "Toward the unification of weak, electromagnetic, and strong interactions: SU(5)," Usp. Fiz. Nauk **130**, 3 (1980) [Sov. Phys. Usp. **23**, 1 (1980)].

Superunification

V. I. Ogievetskiĭ and L. Mezincescu, "Boson-fermion symmetries and superfields," Usp. Fiz. Nauk **117**, 637 (1975) [Sov. Phys. Usp. **18**, 960 (1975)].

A. A. Slavnov, "Supersymmetric gauge theories and theories possible applications to the weak and electromagnetic interactions," Usp. Fiz. Nauk **124**, 487 (1978) [Sov. Phys. Usp. **21**, 240 (1978)].

D. Z. Freedman and P. van Nieuwenhuizen, "Supergravity and the unification of the laws of physics," Sci. Am. **238**, No. 2, 126 (1978) [Russ. transl., Usp. Fiz. Nauk **128**, 135 (1979)].

Astrophysics and cosmology

A. D. Dolgov and Ya. B. Zel'dovich, "Cosmology and elementary particles," Rev. Mod. Phys. **53**, 1 (1981) [Russ. transl. of earlier version, Usp. Fiz. Nauk **130**, 559 (1980)].

Ya. B. Zel'dovich, "Gravitation, charges, cosmology, and

coherence," Usp. Fiz. Nauk **123**, 487 (1977) [Sov. Phys. Usp. **20**, 945 (1977)].
L. P. Grishchuk, "Gravitational waves in the cosmos and the laboratory," Usp. Fiz. Nauk **121**, 629 (1977) [Sov. Phys. Usp. **20**, 319 (1977)].
N. P. Konopleva, "Gravitational experiments in space," Usp. Fiz. Nauk **123**, 537 (1977) [Sov. Phys. Usp. **20**, 973 (1977)].
V. N. Rudenko, "Relativistic experiments in gravitational fields," Usp. Fiz. Nauk **126**, 361 (1978) [Sov. Phys. Usp. **21**, 893 (1978)].
Ya. É. Éinasto, "The structure of galactic systems," Usp. Fiz. Nauk **120**, 497 (1976) [Sov. Phys. Usp. **19**, 955 (1976)].

Accelerators

G. I. Budker and A. N. Skrinskiĭ, "Electron cooling and new possibilities in elementary particle physics," Usp. Fiz. Nauk **124**, 561 (1978) [Sov. Phys. Usp. **21**, 277 (1978)].
B. Richter, "The next generation of electron–positron colliding beam machines," Preprint SLAC-PUB-2274, Stanford University, 1979 [Russ. transl., Usp. Fiz. Nauk **130**, 707 (1980].
B. Richter, "The next generation of accelerators," IEEE Trans. Nucl. Sci. **NS-26**, 4261 (1979) [Russ. transl., Usp. Fiz. Nauk **130**, 717 (1980)].
S. P. Kapitsa, "Seminar on big European projects," Usp. Fiz. Nauk **129**, 549 (1979) [Sov. Phys. Usp. **22**, 939 (1979)].
V. A. Yarba, "Present and future accelerators of particles at superhigh energies," Usp. Fiz. Nauk **129**, 347 (1979) [Sov. Phys. Usp. **22**, 841 (1979)].

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