where $\bar{g}(t)$ is the effective coupling constant

$$\bar{g}^2(t) = g^2\left[1 - \frac{11}{24\pi^2} C_{YM} g^2 t + \cdots\right] \qquad (3)$$

for the pure Yang-Mills theory.

If the integral $\int_{-\infty} \bar{g}^2(t)\,dt$ is divergent, (2) may have a singularity at $g = 0$, in contrast to QED. This would be a crucial difference between QCD and QED.

Note added after the Conference

Our result (at least for the virtual diagrams) was given prior to us, see J.M.Cornwall' Confinement and Infra-red Properties of Yang-Mills Theory' (UCLA/76/TEP/10) and also J.M.Cornwall and G.Tiktopoulos "infra-red properties of Non-Abelian Gauge Theories" (UCLA/76/TEP/2).

P L E N A R Y R E P O R T

GAUGE AND SUPERGAUGE FIELD THEORIES

A. Slavnov
Steklov Mathematical Institute of the Academy of Sciences, Moscow, USSR

1. Introduction

The present section was initially supposed to include all the reports concerning gauge fields. It happened, however, that most of the theoretical contributions to this conference satisfy this requirement. This fact can be interpreted as follows:

Gauge Field Theory = Modern Quantum Field Theory

I want to use this comment to apologize for the fact that I shall not be able to discuss here many interesting investigations. A great part of my report will be devoted to the problems which I think to be most actual. These are, first of all, the problems concerning the explanation of the observed mass spectrum and possible generalizations of gauge theories (supersymmetric gauge theories).

As the theory of gauge fields was discussed in detail at two previous conference I shall give only a short historical review.

The theory of gauge fields began with the classical paper by Yang and Mills /1/ where the principle of gauge invariance, taken from electrodynamics, was generalized to the case of nonabelian group. The role of gauge invariance as a dynamical principle defining the interaction of elementary particles was emphasized by Sakurai /2/. A short time later Weinberg /3/ and Salam /4/ proposed gauge invariant unified models of weak and electromagnetic interactions and the possibility of using spontaneous symmetry breaking for the creation of the gauge field masses was conjectured /5,6/.

However, all these ideas could not find their practical application until the problem of quantization and renormalization of gauge field theory was solved. The solution of this problem encountered serious difficulties and possible ways to overcome them were for the first time outlined by Feynman /7/.

The pioneer papers by Popov and Faddeev/8/ and De Witt/9/ opened a new period in the development of gauge fields. In these papers the consistent procedure of Yang-Mills field quantization was constructed and the relativistically invariant diagram technique was formulated. In papers by Slavnov/10/ and Taylor/11/ the invariant renormalization procedure for gauge theories was developed and the gauge invariance and unitarity of renormalized theory was proved. At last 'tHooft extended the formalism of quantization, developed in papers/8,9/, to gauge theories with spontaneous symmetry breaking and B.Lee and Zinn-Justin/13/, and 't Hooft and Veltman/14/ generalized the renormalization procedure to this case.

As a result of all these and many other papers, which I have no opportunity to list here (the detailed bibliography can be found in rapporteur talks by B.Lee/15/ and J.Illiopoulos/16/), the quantum theory of gauge field was essentially completed and a way for constructing realistic gauge models of elementary particle interaction was opened.

Before discussing recent results I want to remind the main problems of gauge field theory before the London Conference in 1974.

1. Up to now the only observed gauge vector particle is photon. The existence of heavy intermediate vector mesons is not established.

2. The renormalizable models of weak and electromagnetic interactions include, as a rule, besides e, μ, ν_μ and ν_e the hypothetical heavy leptons.

3. The unified models of weak and electromagnetic interactions of hadrons are in a poor agreement with the quark structure of hadrons. The new quarks are necessary.

4. The V-A structure of weak interactions is brought into the unified models from outside.

5. Most of the unified models are described by non-simple Lie groups and therefore are not really universal.

6. The gauge models of strong interactions have no reliable theoretical foundation yet. The popular hypothesis of quark confinement is verified only in nonrealistic models.

7. The problem of infrared divergences in nonabelian theories is not investigated.

8. The Higgs mechanism on which the unified models of weak and electromagnetic interactions are based, is not natural from the viewpoint of standard gauge theories. The incorporation of the Higgs scalar leads to the appearance of new free parameters decreasing the predictive power of the theory.

I have not listed here a number of problems having more particular character as, for example, the rule $\Delta T=1/2$, etc. These questions are discussed in parallel sections and I shall not consider them.

II. Modern Status of Gauge Theories

1. Intermediate Vector Meson.

I can not say anything essentially new on this topic. The intermediate W-bosons, predicted by gauge models must be very heavy ($\gtrsim 50$ GeV) and, hence, their direct observation lies beyond modern experimental possibilities. I would like only to note that the observation of an intermediate meson should not be considered as the only reliable test of gauge theories. The hypothesis about the existence of very heavy objects carrying the interactions, is equivalent in some sense to the conjecture about the structure of space-time at small distances, which is hard to check directly. One should look for indirect consequences.

2. Heavy Leptons

Most of renormalizable models of weak and electromagnetic interactions predict the existence of heavy leptons which are necessary either from group requirements, as in models of the Georgi-Glashow type/17/, or for the compensation of Adler-Bell-Jackiw/18/ anomalies, as in Weinberg-Salam model. This fact considered by most of the people as the drawback of gauge models, now after the appearance of the data indicating the existence of a lepton with the mass $\sim 1,8$ GeV may transform into an argument in favour of the gauge theories.

3. Weak Interactions of Hadrons. New Quarks.

The Gauge models of weak interactions of hadrons are in poor agreement with the hadron SU(3)-structure. The hypothesis about the Cabbibo structure of weak current $J_\mu = \bar{p}\Gamma_\mu n \cos\theta + \bar{p}\Gamma_\mu \lambda \sin\theta$ leads in the gauge theory to the appearance of the strangeness changing neutral currents of the type $\bar{n}\Gamma\lambda$. The difficulty could be avoided if one introduces a new "charmed" quark/19/. So the gauge theories of weak interactions predict the existence of charmed states. At the last conference Illiopoulos betted that the main event of the next conference will be the discovery of charmed states. It looks like that he has won.

4. V-A Structure

In the Salam-Weinberg type models the distinction between the weak and electromagnetic interaction is imposed practically from outside. Besides pure esthetical objections this leads to difficulties in the renormalization due to the existence of axial vector anomalies. More natural from the viewpoint of the unified theories are the models with pure vector interaction where the parity violation arises either spontaneously or due to mass terms.

This idea is realized in so called vector-like models/20/. Up to now mainly the vector-like models based on SU(2)xU(1) group were considered (see, however/21/). The "minimal" model containing 6 doublets of two-component leptons

$$\begin{pmatrix}\nu_e\\ e\end{pmatrix}_L, \begin{pmatrix}N_\mu\\ M\end{pmatrix}_L, \begin{pmatrix}\nu_\mu\\ \mu\end{pmatrix}_L, \begin{pmatrix}N_\mu\\ e\end{pmatrix}_R, \begin{pmatrix}\nu_e\\ M\end{pmatrix}_R, \begin{pmatrix}N_\mu\\ \mu\end{pmatrix}_R \quad (4.1)$$

and 6 doublets of two-component quarks, was widely discussed.

The characteristic feature of the "minimal" model is the nonconservation of lepton number, resulting in the appearance of small but nonzero neutrino mass. The nonconservation of lepton charge leads to a characteristic phenomenon of neutrino oscillations which was previously discussed by Pontecorvo et al./22/ as a possible explanation of the "deficite" of solar neutrinos.

I shall not consider the consequences of the model (4.1) since recent experiments have shown that it does not correspond to reality. In this model as well as in other vector-like models where all the fermions transform according to the same representation of SU(2) the neutral weak current is a pure vectorial one. The experiment indicates the presence of the axial part of neutral current (see the rapporteur talk by S.S.Gerstein).These experiments however, do not exclude the vector-like models at all. The vector character of the neutral current is not necessary consequence of this hypothesis. In particular, the axial part of neutral current automatically arises in super-symmetric vector-like models which will be discussed later.

5. Universal Interaction

The experimental situation is well described in the framework of the gauge group $(SU(2) \times U(1))_{\text{weak+electromagnetic int.}} \times SU(3)_{\text{color}}$. According to the hypothesis of complete universality this group arises as a result of spontaneous breaking of some simple group G. There are many models based on this hypothesis. As a group of the symmetry of the universe the group SU(4) x x SU(4)/23/, SU(5)/24,25/, exceptional groups E_6, E_7/26,27/ were considered. The detailed analysis can be found in paper/28/. For lack of experimental information all these groups can be considered as proper candidates. It is less evident however, that in the nearest future, we shall have this experimental information which would allow us to make a unique choice. Besides, the true inification occurs only for the energies

$\gtrsim 10^{18}$ GeV when the gravitational effects become essential. Therefore the "superunified" models should include the gravitation also. Such attempts have been undertaken already. However, perhaps we should care about the physisists of the future millennium so as not to leave them unemployed.

6. Strong Interactions. "Confinement"

Three main ingredients of the gauge quark theories of hadrons are:

1) The quarks possess the additional degrees of freedom (color). This hypothesis was put forward by Bogolubov, Struminsky and Tavkhelidze/29/ and by Han and Nambu/30/ for the explanation of the hadron state classification and it happens to be very useful for the description of strong interaction dynamics.

2) The strong interactions are mediated by the exchange of color Yang-Mills mesons ("gluons") interacting with quark color degrees of freedom /31,32,23/.

3) There exists a mechanism providing the unobservability of "color" objects.

I shall discuss mainly the third point, the most questionable one.

There are two principally different hypotheses about the nature of this mechanism. According to the first one, the color symmetry is spontaneously broken, the color gluons obtain big masses by the Higgs mechanism. There exist the real quarks with integer charge and belonging to the same multiplet with leptons/23/. The quarks are unstable and decay on leptons+pions with the life-time $\sim 10^{-6}-10^{-13}$ sec., which explains the unsuccessful attempts to find them. It is necessary, however, that after the spontaneous symmetry breaking the global color symmetry remained valid at least in a good approximation. There is the following alternative/33/. If the quarks possess the integer charge, then in spontaneously broken theory it is possible to preserve the approximate color SU(3) symmetry which is necessary for the classification of hadron states. If the quarks possess the fractional charges the color symmetry can be preserved only in the models containing an abelian subgroup. In the nonabelian models of quarks with fractional charges local color symmetry should remain unbroken and the color gluons should be massless. In this case only the second way of explanation of the unobservability of quarks remains, namely, the "confinement" hypothesis/31,32/.

The hypothesis of quark confinement is based on the proposal that the local color symmetry is an exact symmetry of the universe. Then from the invariance of the vacuum under the gauge transformations

$$\Omega|0\rangle = |0\rangle$$

it follows that

$$\langle T\bar{\psi}(x)\psi(y)\rangle_0 = \langle T\Omega\bar{\psi}(x)\psi(y)\Omega^{-1}\rangle_0 \sim \delta(x-y), \quad (6.1)$$

i.e. the Green functions of all color objects vanish identically. The observables are only the gauge invariant states, for example, the singlet bound states which can be identified with hadrons. The intuitive image of this picture is the system of particles with the interaction increasing with distances. In such a system the quarks are likely to be bound with some elastic strings which do not allow them to spread on a macroscopic distance.

I would like to illustrate how such a situation can arise in the field theory by an example which has no direct connection with the gauge fields but possesses many common features with them. I mean the nonlinear σ-model/34,35,36/. I shall follow the paper by Bresin and Zinn-Justin/34/. In d-dimensional nonlinear σ-model the β-function from the renormalization group equation is

$$\beta(g) = (d-2)g - (n-2)\frac{g^2}{2\pi} + O(g^3), \quad (6.2)$$

where n is an order of symmetry group of the Lagrangian. For $d>2$ there exists the ultraviolet fixed point

$$g^* = \frac{(d-2)2\pi}{n-2}. \quad (6.3)$$

For $g<g^*$ the perturbation theory is sensible, the symmetry is spontaneously broken and the massless field π is a Goldstone boson. The region $g>g^*$ can not be described in the framework of the perturbation theory. The transition through the point g^* corresponds to the phase transition. For $g>g^*$ the spectrum consists of (n-1) massive π-particle and the bound state σ with the same mass. The spontaneous symmetry breaking disappears. For d=2, as it is seen from (6.2) the model is asymptotically free and the perturbation theory is senseless for any g. (Formally it manifests in the appearance of infrared divergences).

The four dimensional Yang-Mills theory in this aspect resembles the two dimensional σ-model. In particular Olesen[37] gave plausible arguments that in the Yang-Mills theory the β-function have no zeros at finite g. It indicates that in this case we have the symmetric phase as well. The effective quark interaction increases infinitely in the infrared region, i.e. the confinement takes place.

As far as the confinement is related to the nonvalidity of perturbation theory, up to now this effect was established either in solvable models (mostly two-dimensional) or in the models where the method of strong coupling is applied. The confinement in the two-dimensional model is almost trivial for in this case the Coulomb potential increases linearly with distance. Nevertheless the two-dimensional models are a good laboratory as they allow one to investigate different aspects of this phenomenon. The simplest model of this kind is the Shwinger model, the two-dimensional electrondynamics[38]. In this model with massive fermions the interaction energy of two external charges separated by the distance L is[39].

$$E=\left[\varepsilon\left(\theta-\frac{2\pi Q}{e}\right)-\varepsilon(\theta)\right]L+\cdots, \qquad (6.4)$$

where ε is a periodic function. Hence, for $Q=ne$ the interaction disappears. As in the massless Shwinger model there arise the screening of charge and quark confinement. At the same time for arbitrary the long-distance interaction is present and, hence, contrary to the massless case, the symmetry is unbroken.

Interesting results were obtained while investigating the two-dimensional SU(N) Yang-Mills theory[40,41]. In this model there is a discrete set of singlet bound states with finite masses. The model is unitary in a gauge invariant (singlet) sector, i.e. the transitions between the singlet and color states are absent. The long-distance interaction between the singlet states is also absent. At the same time the behaviour of the theory at small distances is determined by the noninteracting quarks. The authors[41] made an important observation. It is a common viewpoint that the confinement is equivalent to the infinite mass of color states. It happens, however, that the properties of color gauge noninvariant states depend essentially on the procedure of infrared regularization. In particular such a regularization can be indicated[42] that the quarks obtain finite masses. The crucial point is that for any regularization the transitions between the singlet and color states are absent.

The possible way of realization of "infrared prison" in four dimensional space-time was proposed by Wilson[43] who assumed that to construct a prison in the real world the lattice is necessary. Wilson generalized the principle of gauge invariance to the case of lattice Euclidean space-time. With every net of a lattice the field of matter $\varphi(x)$ is associated which is transformed as follows:

$$\varphi(x)\rightarrow R(x)\varphi(x) \qquad (6.5)$$

and with every link the bilocal gauge field $A(x,x')$ transformed as

$$A(x,x')\rightarrow R(x)A(x,x')R^{-1}(x'). \qquad (6.6)$$

The introduction of a lattice leads to the ultraviolet finite theory and at the same time removes the divergences connected with the degeneracy of gauge invariant action. The main advantage which gives a lattice is the possibility of application of approximate methods different from perturbation theory, for instanoe, the approximation of strong coupling or mean field method. A number of papers[44,45,46,47,48] was devoted to the investigation of gauge theories on a lattice. Their conclusions confirm the qualitative picture of confinement discussed above.

However, nowadays, it is difficult to say to what extent all these results are applicable to the real nature as the introduction of a lattice violates the relativistic (and even Euclidian) invariance. The results obtained do not admit the limit $a\to 0$ and the spectra of "observable" states strongly depends on a. (The model could be considered as satisfactory if the masses were much smaller than inverse lattice parameter. In fact the situation is reverse, i.e. $M \gtrsim a^{-1}$). It is possible that this situation can be avoided using more effective calculation methods. Thus, for example, the replacement of the series in $(ag)^{-1}$ with the Pade approximation in the massive Shwinger model enables to obtain a good agreement with the results for the continuous space[49].

Essentially another possibility which is not connected with the limit $a\to 0$ is proposed by Wilson[50]. In general the idea is as follows. The propagator of gauge field can be devided into a "low-energy" and "high-energy" part in such a way

$$\frac{1}{p^2} \sim \frac{1}{p^2+\Lambda^2} + \frac{\Lambda^2}{p^2(p^2+\Lambda^2)} , \qquad (6.7)$$

where Λ - is a cut-off momenta (the inverse lattice parameter). First it is proposed to sum up the contributions of all high energy parts. For Λ large enough in asymptotically free theories, we can use the usual perturbation theory. As a result of such summation there arise an effective Lagrangian containing the infinite number of vertices with different number of "low-energy" legs. At the second stage of calculations it is proposed to use the strong coupling method where a low-energy propagator should be taken as a Green function of a lattice theory. In other words the calculation of relativistically invariant S-matrix is reduced to the calculation of the S-matrix in a lattice theory with some effective action.

In principle, the proposed program does not meet objections, but it seems to me that when we really try to put it into practice we again would face the violation of relativistic invariance because the two stage calculations can be carried out only approximately and the conditions of applicability of perturbation theory at the first stage (large Λ) and of the strong coupling method at the second one (small Λ) are opposed to each other. But of course everything depends on the skill and, possibly, professor Wilson or somebody else will demonstrate us the vitality of this idea.

It is not excluded that correct formulation of lattice theories requires changing the geometry of space-time.

Donkov, Kadyshevsky, Mateev, Mir-Kasimov[51] formulated quantum field theory incorporating a fundamental length ℓ and rigorously satisfying Poincare invariance. The mathematical realization of such a theory is possible in curved momentum space with radius $1/\ell$. This changes drastically the spacetime geometry at short distances $\sim \ell$; in particular time becomes discrete. The usual difficulty which one meets changing the geometry is to define properly the invariant time ordered product. In the approach[51] this is achieved automatically. A gauge field theory in this framework is in a certain sence similar to gauge theories on a lattice: the zero component i.e. the scalar potential of the electromagnetic field is necessarily an angular variable $|\varphi| \leqslant \frac{\hbar c}{\ell e}\pi$

and the charges of the elementary particles are integer multiples of e: thus charges are conmensurate.

The attempts are made to obtain some quantitative consequences from the confinement hypothesis. A.A.Migdal/52/ assumed, that the nonanalyticity of hadron masses in coupling constant has a universal character. Owing to this the ratio of the masses can be determined on the basis of perturbation theory. The Green functions of singlet operators are looked for in the form of the matrix Pade approximation

$$\langle T O_i O_j \rangle = \frac{P_M}{Q_N} \qquad (6.8)$$

For the determination of Q_N and P_M the condition of sewing at some remote Euclidian point $p^2=-\Lambda^2$ is used. If the renormalized coupling constant g_p satisfies some consistency condition then in the limit $M, N, \Lambda\to\infty$ the discrete spectrum survives and the mass ratios can be found in the form of series in g_p. Unfortunately, it is very difficult to test the consistency condition and in real calculations g_p is used as a fitting parameter. The physical sense of the used limiting procedure is also rather obscure.

7. Infrared divergences in perturbation theory.

Though the statement that the perturbation theory for the massless Yang-Mills fields "suffers very strong infrared singularities" became a "folklore" the real nature of these singularities is still unclear. That is why we should welcome the appearance of papers/53-59/ where serious attempts have been made to clear up this problem. The calculations in the lowest orders of pertubation theory for the Yang-Mills field/53,54/ confirm the theorem by Kinoshita and Lee-Nauenberg/60,61/ about the absence of mass singularity in inclusive cross section. The scattering cross section being averaged over initial and summed over final degenerate (color) states tends to a definite limit while removing infrared cutt-off. It should be noted that the theorem by Kinoshita-Lee-Nauenberg does not concern the question of ultraviolet renormalizations. Unfortunately this question is practically not considered also in papers/53,54/, because the calculations are carried out with the subtraction at nonphysical Euclidian point. The problem of determination of the physical charge is open.

The result of/53,54/ does not exclude the possibility that beyond the framework of perturbation theory or with the proper summation of the series in g the infrared singularities can lead to the suppression of the processes with the radiation of color objects. In particular it is stated that if we sum the contribution of the leading (double logarithmic) singularities over all orders of perturbation theory the amplitude of the process with the radiation of a color particle vanishes (exponentially in the regularization mass)/55,56/. An analogous effect was observed in massless electrodynamics by Fomin et al./57/. The difference between the results of papers/53,54/ and /55,56/ can be interpreted in terms of Bogolubov's quasiaverages/62/. The objects calculated in/55,56/ are essentially the quasiaverages while the results of/53,54/ refer to the ordinary expectation values. One should also to keep in mind that the results/55,56/ can have no relation to the confinement because the contribution of nonleading logarithms in principle can lead to finite cross sections. And besides it is not evident that the account of infinite number of soft colour quanta would not lead to the finite result. To make these questions clear it is very important to perform a consistent investigation of infrared singularities of arbitrary diagrams. Such an investigation was carried out in the paper by Kinoshita and Ukawa/58/ contributed to this conference, in which the algorithm for the determination of infrared singularities in the Yang-Mills theory is given for an arbitrary order of perturbation theory. Interesting possibilities are also opened due to the application of the renormalization group methods to this problem[59].

8. Particle-like solutions in gauge theories.

The characteristic feature of gauge theories is the existence of classical particle-like solutions with finite energy, the solitons. The well-known example of such solutions is the magnetic monopole in Georgi-Glashow model[63]. The different aspects of solitons were considered in the talk by Faddeev[64]. So I restrict myself to the discussion demonstrating their possible role in the structure of gauge theories.

In the formalism of continual integral the Euclidian Green functions of the fields A_μ are determined by the products of $A_\mu(x_i)$ integrated with the weight

$$\exp\left\{-S(A)\right\} = \exp\left\{-\frac{1}{4g^2}\int \mathcal{L}(A)dx\right\} \quad (8.1)$$

For small g^2 the main contribution to the integral is given by the fields A_μ in the neighbourhood of classical solutions

$$\frac{\delta S}{\delta A_\mu} = 0 \ ; \ S(A) < \infty \ . \quad (8.2)$$

The ordinary perturbation theory corresponds to the consideration of the trivial minimum $A_\mu = 0$ only. In the presence of long-distance interaction it became senseless due to the infrared divergences. A.M.Polyakov conjectured[65] that nontrivial solutions of (8.2) can lead to the existence of finite correlation length, i.e. the finite mass of vector fields. (The considered system can be represented as a plasma consisting of magnetic monopoles). He succeeded in showing that this effect really takes place in electrodynamics in three-dimensional lattice space and in three-dimensional Euclidian Georgi-Glashow model.

However, it is not clear whether these considerations have something to do with the real four-dimensional case because the structure of a gauge group crucially depends on the number of dimensions.

In the natural gauge $A_0=0$ the remaining gauge transformations $g(\vec{x})$ depend only on space coordinates. For $d=4$, under the condition

$$g(\vec{x}) \xrightarrow[\vec{x}\to\infty]{} 1 \quad (8.3)$$

the gauge group is divided into classes according to the values of topological charge

$$Q = \int d^3x\ \varepsilon^{ikj}\, \mathrm{Tr}\, T_i T_k T_j \ ; \ T_i = \partial_i g(\vec{x}) g^{-1}(\vec{x}) \quad (8.4)$$

In other words the group transformations include, apart from continious infinitesimal transformations, shifts T changing the topological charge by 1.

Every class of this type nulifies the action and therefore generates its own "vacuum" $|n\rangle$. Euclidian solutions of four-dimensional Yang-Mills equations (instantons), found by Belavin et al[66] describe a tunnel effect between two adjacent "vacua".

There are two possibilities of interpretation of this fact. According to the first one [67,68] the gauge group consists of only connected transformations, all states $|n\rangle$ having the same energy and real gauge invariant vacuum being a coherent superposition

$$|\theta\rangle = \sum_n e^{in\theta}|n\rangle + \ldots \qquad T|\theta\rangle = e^{i\theta}|\theta\rangle \quad (8.5)$$

In the other interpretation, proposed by Faddeev[69], the gauge group is a group of all transformations, including shifts T. Then only periodic states $(\theta=\vartheta)$ are permitted. Note that an analogy with a solid state case suggests that the states with different should have different energies and the minimal one corresponds to $\theta=0$.

Such a structure of the vacuum can be related to the $U(1)$-problem. Indeed, an axial current anomaly

$$B(A) = \frac{1}{4\pi^2}\varepsilon^{ikj}\,\mathrm{Tr}\left(A_i A_k A_j - \frac{3}{2}F_{ik}A_j\right) \quad (8.6)$$

is changed by 2 under the action of the operator T. Therefore, in the space of states with the fixed θ only the operator $e^{i\pi Q_5}$ makes sense, but not the operator $e^{i\alpha Q_5}$ for arbitrary α, i.e. continious γ^5-transformations are forbidden and there is no reason for the existence of isoscalar Goldstone bosons (η-problem)[70,67,68], Faddeev (private communication)).

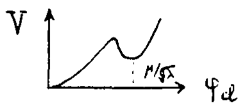
All these considerations assume the asymptotic condition (8.3). This condition must follow from the physics. In my opinion the only possible justification of the condition (8.3) is a fixation of a neutral direction at infinity. Such a fixation certainly contradicts the "confinement" ideology. Thus, confinement in a real world still remains an appealing but rather controversial hypothesis.

9. Spontaneous symmetry breaking.

At present the only reliable method of gauge field mass generation is the Higgs mechanism. The classical Higgs effect arises in the potential

$$V = \lambda\varphi^4 - \frac{\mu^2\varphi^2}{2} \qquad (9.1)$$

For $\mu=0$ in the tree approximation the effect is absent but nevertheless the radiative corrections can result in the symmetry breaking /71/. The opposite effect is also possible. In the system, consisting of the field φ, interacting with gauge field, under a certain relation between the constants g and λ, the effect of spontaneously symmetry breaking being present in the tree approximation, can vanish while considering the loop corrections in g /72/. The effective potential then has the form



Hence we have the limitation on the masses of Higgs mesons. In the Salam-Weinberg type models the symmetry breaking takes place if

$$m > 5 \text{ GeV} \qquad (9.2)$$

Thus, the possibility of existence of light Higgs mesons is excluded. The qualitative picture of Higgs effect can change also under a macroscopic influence: heating, density changing, external fields. These effects were discussed in detail in the talk by Kirzhnits /73/.

Ivanov and Ogievetsky /74/ investigated a connection between the Higgs effect and gauge fields. They have shown that gauge theory can be considered as a result of spontaneous symmetry breaking with respect to some infinite-parametric group where Yang-Mills field is a Goldstone boson realizing a nonlinear representation of this group. This observation could have interesting applications if one succeeded in constructing the corresponding linear representation (infinite-dimensional analog of σ-model). Such a theory could be referred to dual models.

The last comment conserning the Higgs effect is related to asymptotically free theories. Despite the common opinion /75/ there is an opportunity of spontaneous symmetry breaking resulting in the fact that all physical particles acquire nonzero masses and at the same time the theory remains asymptotically free. This possibility is based on the use of unstable solutions of renormalization group equations, for the system of Yang-Mills, Ukawa and φ^4 interaction. For certain relations between the coupling constants g, h and λ there exist unstable solutions for which all the invariant charges tend to zero when the arguments tend to infinity, i.e. corresponding to the asymptotically free theory /76/. Using this idea Fradkin and Kalashnikov have shown that $SU(5)$ symmetry, unifying weak, electromagnetic and strong interactions can be broken to $U(2)\times SU_{3c}$ without loss of asymptotical freedom. It seems to me, however, that this mechanism can have a relation to reality only in the case, when the unstable solution corresponds to some symmetry group (in this case the notion of unstability itself becomes senseless as there is only one coupling constant in the theory). Such a situation takes place, for example, in supersymmetric theories. If the unstable solution does not correspond to any symmetry group, an arbitrary weak external influence can change the situation drastically.

Summarizing we can say that the Higgs mechanism seems to be quite a reasonable method of intermediate meson mass generation in weak interactions. However, from the point of view of gauge fields it is pure "external" as the appearance of Higgs scalars does not follow from the original principles. As the parameters of Higgs mesons are not fixed by the gauge invariance their introduction decreases considerably the predictive power of the theory and makes it difficult to choose the concrete model. There arises a natural desire to formulate the theory in such a way that the Higgs mesons arised with necessity and did not bring extra arbitrariness. This can be achieved in two ways. The first way is the dynamical symmetry breaking where the Higgs mesons arise as bound states of basic fermions. This idea originates from the classical papers by Bardeen, Cooper and Shriffer and N.N.Bogolubov on the theory of superconductivity and was used in the field theory in the papers /77,78/. In the framework of gauge theories this idea was recently developed in papers/79/. Unfortunately in this approach it is difficult to obtain any reliable quantitative results as one has to solve the equation for the bound states of a Bethe-Salpeter type.

The second way, where the Higgs mesons are described by elementary scalar fields and their presence is dictated by symmetry requirements seems to me more promising. The corresponding transformations should mix the basic particles of gauge theory (fermions and vector fields) with the Higgs scalars. It is the property the supersymmetry transformations possess, and the remaining part of my talk will be devoted to this subject.

III. Supersymmetry

1. The main concepts.

The algebra of symmetry connecting Fermi and Bose fields should contain anticommuting elements, i.e. should be not an ordinary but graduated Lie algebra. The minimal algebra incorporating anticommuting elements and containing as a subalgebra the algebra of Poincare group is

$$[P_\mu, P_\nu] = 0, \quad [P_\mu, S_\alpha] = 0, \quad [S_\alpha, S_\beta]_+ = -(\gamma_\mu C)^{\alpha\beta} P_\mu \qquad (1.1)$$

where P_μ are the generators of four-dimensional translations, S_α are those of supersymmetry transformations being the Majorana spinors, C is a charge conjugation matrix. The algebra (1.1) was firstly considered in papers/80)/ in connection with the parity nonconservation problem. Nonlinear realizations of this algebra were investigated in paper/81/, where the hypothesis that neutrino is a Goldstone fermion was put forward. The papers/80,81/ did not get a wide resonance as there was not proposed any example of renormalizable supersymmetric model. Such a model was built by Wess and Zumino/82/ who came to the supersymmetry idea independently generalizing two-dimensional dual models. The papers /82,83/ initiated a number of investigations and I shall try to present their main results below.

Salam and Strathdee introduced a concept of "superfield" which allows one to formulate supersymmetric theories in an elegant and suitable way/84/. It is convenient/84/ to realize the representation of algebra (1.1) in space of functions of 8 variables (x_μ, θ_α)/81/, where x_μ are the commuting real parameters and θ_α are anticommuting Majorana spinors.

The algebra (1.1) corresponds to the transformation of superspace

$$x'_\mu = x_\mu + \frac{i}{2}\bar{\varepsilon}\gamma^\mu\theta, \quad \theta'_\alpha = \theta_\alpha + \varepsilon_\alpha \qquad (1.2)$$

The scalar with respect to the transformation (1.2) is the function $\Psi(x,\theta)$ transformed as follows

$$\Psi(x,\theta) = \Psi'(x',\theta') \qquad (1.3)$$

It follows from the anticommutativity of parameters θ that $\theta_\alpha^2 = 0$ and, hence, any function of θ is a finite polynomial. In particular

$$\Psi(x,\theta) = C(x) + \bar\theta \chi(x) + \frac{1}{4}\Big\{\bar\theta\theta M(x) + \bar\theta\gamma^5\theta N(x) + \bar\theta i\gamma_\nu\gamma_5\theta A_\nu(x) + (\bar\theta\theta)(\bar\theta\lambda(x))\Big\} + \frac{1}{32}(\bar\theta\theta)^2 D(x) \quad (1.4)$$

The superfield of general form $\Psi(x,\theta)$ is equivalent to the supermultiplet of ordinary fields containing (pseudo) scalars C,M,N,D, Majorana spinors χ, λ and vector field A_μ. The superfield $\Psi(x,\theta)$ provides invariant expansion into the sum of three superfields with smaller number of components

$$\Psi(x,\theta) = \Phi_+(x,\theta) + \Phi_-(x,\theta) + \Psi_\perp(x,\theta) \quad (1.5)$$

The "chiral" superfields $\Phi_\pm$ are equivalent to the multiplets $\{A_\pm, \psi_\pm, F_\pm\}$ consisting of (pseudo) scalars $A_\pm$, $F_\pm$ and two-component spinor $\psi_\pm$. As the fermions and bosons enter into the same multiplet some scalar particles acquire nonzero fermion number. It imposes severe constraints on the possible fermion-number conserving models[85].

The action principle can be formulated in terms of superfields[86,87,88] if one uses the integration over the Grassman algebra introduced in[89].

2. Supersymmetric gauge theories.

The most interesting for applications are the supersymmetric theories invariant under gauge transformations. For the description of matter fields one can use the chiral superfields $\Phi_\pm$ and the gauge fields can be incorporated into the superfield of general form.

The generalized gauge transformations depend on 8 arbitrary functions. Hence, 8 components of the superfield of a general form are nonphysical and can be excluded. The invariant kinetic term for the fields of matter has the form[90,91,92]

$$\frac{1}{4}\int d^4x\, d^4\theta \left\{\Phi_+^+ e^{g\Psi}\Phi_+ + \Phi_-^+ e^{-g\Psi}\Phi_-\right\} \quad (2.1)$$

As always in gauge invariant theory the action is degenerate and when quantizing it is necessary to put an subsidiary condition fixing the gauge. Zumino and Wess[90] proposed to use the gauge freedom to put the components C,X,M,N of the field Ψ equal to zero. This gauge has a remarkable property: the series (2.1) become a finite polynomial and after excluding of the auxiliary fields F_+ and D the action becomes (for simplicity it is written for the Abelian case):

$$\begin{aligned} S = \int d^4x \Big\{ &-\frac{1}{4}(\partial_\mu A_\nu - \partial_\nu A_\mu)^2 + \frac{i}{2}\bar\lambda\hat\partial\lambda + \frac{1}{2}D^2 + F_+^+F_+ + \\ &+ (\partial_\mu A_+^+ + \frac{ig}{2}A_+^+A_\mu)^2 + i\bar\psi\gamma_\mu(\partial_\mu - \frac{ig}{2}A_\mu)\psi + \\ &+ \frac{g}{2}(A_+^+A_+ - A_-^+A_-)D + M(A_-^+F_+ - \bar\psi\psi) + \frac{ig}{\sqrt2}A_+^+\bar\lambda\psi_+ - \\ &- M^2(A_+^+A_+ + A_-^+A_-) - \frac{g^2}{2}(A_+^+A_+ - A_-^+A_-)^2 + h.c.\Big\} \end{aligned} \quad (2.2)$$

The supersymmetric Abelian gauge theory describes the minimal electromagnetic interaction of spinor and scalar fields (candidates for Higgs mesons) plus contact interactions $\bar\Psi\Gamma\psi\varphi$ and φ^4 characterized by the same constant g.

In the Zumino-Wess gauge all the "unphysical" components are excluded and the interpretation is evident. However, the condition

$$C = X = M = N = 0$$

violates the explicit supersymmetry of the theory that makes difficult to carry out the renormalization procedure (analogous difficulties arise in QED when using the Coulomb gauge). The explicitly invariant perturbation theory for a supersymmetric QED was constructed in paper[93]. Instead of noninvariant condition (2.3) the supersymmetric condition of the type:

$$\Psi_+ = 0 \quad (2.4)$$

was used.

In this gauge the action remains nonpolynomial and the perturbation series contains the infinite number of types of primitively divergent diagrams. It could be shown, however, that the index of divergence is not larger than 2 and because of the existence of infinite system of generalized Ward identities all the counterterms are expressed in terms of two independent constants being the wave function renormalization of matter and gauge fields. An analogous procedure can be applied to nonabelian supersymmetric theories[94].

In this case like the ordinary Yang-Mills theory there arises some additional charge renormalization. An independent matter fields mass renormalization is absent. Besides in supersymmetric theories the sum of all vacuum diagrams cancels[95].

3. Spontaneous Breaking of Supersymmetry

As the mass degeneration of scalar and spinor particles is absent in the nature, the supersymmetry should be broken. This problem is nontrivial one as the supersymmetry puts down severe constraints on the form of an effective potential. The possibility of spontaneous supersymmetry breaking was demonstrated for the first time by Fayet and Illiopoulos[96]. The proposed mechanism can be applied to Abelian supersymmetric theories. Later it was shown that the spontaneous supersymmetry breaking is also possible in a system of n interacting chiral fields for $n \geqslant 3$ [97,98].

The mentioned mechanism can be applied to a rather restricted class of theories and the attempts to use them for a construction of a realistic model[99] were not a success. Besides they lead to the difficulties with the physical interpretation of Goldstone fermion accompanying spontaneous symmetry breaking[81]. For the processes with the Goldstone fermion there exist the low-energy theorems[81,100,101] analogous to the Adler theorem for the Goldstone bosons. It follows from them that if the electronic neutrino is a Goldstone fermion then the amplitude of β-decay should tend to zero for zero momenta of the neutrino[100,101]. Such a behaviour disagrees with experiment. There is of course the opportunity to identify the Goldstone fermion with muonic or some new "neutrino"[102].

I want to discuss one more mechanism of supersymmetry breaking which enlarges essentially the admissible class of theories and is free of difficulty with the Goldstone neutrino[103,104].

With the help of this mechanism arbitrary mass terms for scalar component of superfields can be generated in any supersymmetric theory.

Let me exemplify this statement. To any supersymmetric action we can add an invariant term of the form

$$\{\Phi_+^+\Phi_- R_- + \Phi_-^+\Phi_+ R_+\}_F + \frac{1}{2}\{R_+\tilde{R}_- + \tilde{R}_+ R_-\}_D + c(\tilde{R}_+ + \tilde{R}_-)_F , \quad (3.1)$$

where $R_\pm$ and $\tilde{R}_\pm$ are auxiliary chiral fields, which are the singlets with respect to the gauge group.

As far as $\langle R_F \rangle_0 \neq 0$ the canonical transformation $R_F \to R_F + C$ produces the mass term

$$C\,(A_-^+ A_+ + A_+^+ A_-) \quad (3.2)$$

which removes the degeneracy.

The variation over $\tilde{R}$ leads to the free equations for R. So the auxiliary fields R and $\tilde{R}$ are completely separated from physical fields and their only observable effect is the appearance of the mass term (3.2). At the same time the explicit supersymmetry of (3.1) enables us to apply the renormalization procedure developed for the symmetrical theories. Analogously one can obtain the mass terms of the form

$$a A_+^+ A_+ + b A_-^+ A_- . \quad (3.3)$$

The Goldstone fermion enters into the auxiliary supermultiplets ($R_\pm$) and does not take part in observable processes[103,104,105].

The described mechanism can be applied to a wide class of theories, in particular, to nonabelian supersymmetric ones. The nonabelian supersymmetric models are an example of the asymptotically free theory with scalar particles. The spontaneous supersymmetry breaking enables us to make all observable vector and scalar particles massive and to construct the models which are both asymptotically free and infrared convergent[103]. At the same time this mechanism opens the way for the construction of realistic models of weak and electromagnetic interactions.

4. On the Way to a Realistic Model

The supersymmetry strongly decreases the number of possible candidates for realistic models. The minimal supersymmetric model of leptons/104/ can be constructed on the basis of gauge group SU(2)xSU(1). The fields of matter are described by two chiral isodoublets $\Phi_{\pm 1,2}$ and one chiral singlet $S_+ = S_-^+$. The gauge fields compose an isotriplet Ψ_a and singlet Ψ. The interaction has the form

$$\mathcal{L} = \frac{1}{2}\left\{ \Phi_{+i}^{+} e^{g\Psi_a \tau_a + g_1 \Psi} \Phi_{+i} + \Phi_{-i}^{+} e^{-g\Psi_a\tau_a - g_1\Psi} \Phi_{-i} \right\}_D - \frac{1}{2}\left(a_{ij} S_- \Phi_{+i}^{+} \Phi_{-j} \right)_F + h.c \quad (4.1)$$

The spectra of observed leptons looks as foolws:

The charged sector consists of electron, muon and heavy lepton E. The neutral sector includes electronic, muonic neutrinos, one new "neutrino" ν_x and two heavy Dirac fermions with masses $\gg m_K$ (K-meson). The model is a vector-like one but the effective neutral current contains both the vector and axial parts.

In spite of the fact that the gauge group is not simple one, the mass of the charged intermediate meson is uniquely fixed and is equal to m_W==37.3 GeV. The mass of heavy charged lepton is $\sim m_W$. The masses of heavy neutral lepton can vary in a wide range.

The model correctly reproduces the spectrum of "light" leptons and includes the standard V-A theory. Another model was proposed recently by Fayet/106/. In this model, however, electron and muon are massless, at least, in the tree approximation.

The construction of realistic supersymmetric models makes its first steps and I hope that in the nearest future we shall be the witnesses of further progress in this direction.

5. "Super"Supersymmetry

There exist more pretentious hopes to construct the "superunified" models in the framework of supersymmetry. I mean, on the one hand, the possibility of nontrivial unification of supersymmetry with internal symmetries and on the other hand, the supersymmetrical generalizations of gravity.

Concerning the first possibility, as far as the supersymmetry algebra is not a usual Lie algebra, the standard "no-go" theorems are not applied to it and the nontrivial synthesis of supersymmetry and internal symmetry is possible/107/. If we supply the generators of supersymmetry by "inner" indices i, j, then in a local theory with massive particles the admissible algebra has the form/108/

$$[S_{\alpha i}, S_{\beta j}]_+ = \delta^{ij}(C\gamma_\mu)_{\alpha\beta} P_\mu, \quad [S_{\alpha i}, P_\mu] = 0. \quad (5.1)$$

However, the representation of this algebra even in the case of the simplest groups of internal symmetry contains a huge number of fields which are difficult to connect with experiment and the corresponding Lagrangians correspond to nonrenormalizable theory/109/. Some possibility to avoid these difficulties was proposed in/110/, but at present it is not clear whether it is possible to construct along these lines a physically interesting model.

Supersymmetric generalizations of gravity were also considered/111-118/. The most direct way is to replace a flat superspace by the Riemannian superspace $Z_\mu(x_\mu, \theta_\mu)$ with a metric tensor $g_{MN}(z)$/112/. Apart from gravitational field this tensor includes fields with spins 3, 5/2 and so on. The supersymmetric "Einstein equations" are invariant with respect to general coordinate transformations. This invariance can be broken spontaneously up to the global supersymmetry. It appears that the possibility of symmetry breaking imposes strong constraints on a possible internal symmetry group/113/. Although this scheme seems appealing nobody succeeded in demostrating its selfconsistency and in constructing an acceptable model. The origin of these difficulties and possible way out were discussed in/114,115/. There exist alternative approaches to supergravity. The gravitation field can be included into a vector superfield $h_\mu(x,\theta)$ the source of which is a conserved supercurrent /116/, which is a supersymmetric generalization of a momentum-energy tensor/119/. Another possibility which seems to be the most promising is based on the observation that the minimal interaction of massless Rarita-Schwinger field with gravitation is invariant under the local supersymmetry transformations/117/. This model provides a consistent way of describing the interacting spin 3/2 field. Das and Freedman/118/ quantized this model and demonstrated the absence of acausal effects in the tree approximation.

The renormalization problem for supergravity has not been yet seriously considered, but there are hopes that such a theory will be less divergent than the usual Eistein gravity.

Conclusion

It is impossible to cover in one-hour talk all interesting results obtained in gauge theories during last two years. A considerable work was performed, but still much more should be done. Gauge theories already provide self consistent apparatus for the theory of weak and electromagnetic interactions. Further progress in this field demands an accumulation of experimental data and, on the other hand, improvement of the theoretical criteria for a realistic model. I think, the supersymmetry can serve for this purpose. As to the strong interactions here we have a lot of data and a number of ideas but a consistent gauge models still wait to be discovered. However, if beauty is indeed a criterion of truth, then gauge theories here also have good chances. Finally, now we are doing the first real steps to the final goal - unification of all interactions. Let us hope that we have chosen a right way.

Acknowlegements

I would like to thank all the participants of the seminar directed by N.N.Bogolubov and A.A.Logunov for discussion of a preliminary version of this talk.

I am indebted to L.Faddeev for a critical reading of the manuscript and helpful suggestions. I thank V.P.Akulov, D.I.Kazakov and P.P.Kulish who helped me in the preparation of the final version of this report.

References

1. C.N.Yang, R.Mills.Phys.Rev.,96,191,1954.
2. J.J.Sakurai. Ann.Phys.,11,1,1960.
3. S.Weinberg. Phys.Rev.Lett.,19,1264,1967.
4. A.Salam."Elementary Particle Physics", Stockholm,1968.
5. P.W.Higgs. Phys.Rev.Lett.,13,508,1964; Phys. Rev.,1156,1966.
6. T.W.B.Kibble.Phys.Rev.,155,1554,1967.
7. R.P.Feynman. Acta Phys.Pol.,24,697,1963.
8. L.D.Faddeev, V.N.Popov.Phys.Lett.,B25,29,1967; Preprint ITP-67-36,Kiev,1967.
9. B.S.De Witt.Phys.Rev.,162,1195,1239,1976.
10. A.A.Slavnov.Preprint ITP-71,83E, Kiev 1971, TMF, 10,153,1972; 13,174,1972.
11. J.C.Taylor.Nucl.Phys.,B33,435,1971.
12. G't.Hooft.Nucl.Phys.,B35,173,1971.
13. B.W.Lee, J.Zinn-Justin.Phys.Rev.,D5,3121,3137, 3155,1972.
14. G't.Hooft,M.Veltman.Nucl.Phys.,B44,189,1972; B50,318,1972.
15. B.W.Lee."Proceedings of the XVI International Conference on High Energy Physics" (NAL Batavia, III,1972) v.IV.
16. J.I.Illiopoulos. "Proceedings of the XVII International Conference on High Energy Physics". (London,1974).
17. H.Georgy, S.Glashow.Phys.Rev.,D6,429,1972.
18. S.Adler.Phys.Rev.,177,2426,1969; J.C.Bell, R.Jackiw. Nuovo Cimento,60,47,1969.
19. S.Glashow, J.Illiopoulos, L.Maiani.Phys.Rev., D2,1285,1970.
20. H.Fritzsch, M.Gell-Mann,P.Minkowski.Phys. Lett.,59B, 256,1975.
A.De Rujula, H.Georgi, S.Glashow.Phys.Rev.,D12 2768,1975.
F.A.Wilczek, A.Zee,R.L.Kingsley, Treiman.Phys. Rev., D12,2768,1975; S.Pakwara,W.Simmons,S.P.Tuom.Phys.Rev.Lett.,35,702,1972.
21. H.Fritzsch,P.Minkowski.Preprint CERN TH,2165, 1975.
22. B.Pontecorvo.JETP,34,247,1958; V.N.Gribov, B.Pontecorvo.Phys.Lett.,24B,493,1969.
23. J.Pati,A.Salam.Phys.Rev.,D8,1240,1973; D10,275, 1974; D11,703,1975.
24. H.Georgi,S.Glashow.Phys.Rev.Lett.,32,438,1974.

25. O.K.Kalashnikov, E.S.Fradkin,533.Contributed papers.
26. P.Ramond, Cal.Tech.Preprint. Calt-68-540,1976.
27. F.Gursey,P.Sikivie.Yale Preprint COO-3075-133, 1976.
28. H.Fritzsoh, P.Minkowski.Ann. of Phys.,93, 193,1975.
29. N.N.Bogolubov, B.V.Struminsky, A.N.Tavkhelidze. JINR Preprint D-1968,Dubna,1965.
30. M.Y.Han, Y.Nambu.Phys.Rev.,139,1006,1965.
31. S.Weinberg.Phys.Rev.Lett.,31,494,1973.
32. H.Fritzsoh, M.Gell-Mann,H.Leutwyler.Phys. Lett.,47B,365,1973.
33. R.N.Mohapatra, J.C.Pati,A.Salam.Phys.Rev., D13,1733,1976.
34. A.M.Polyakov.Phys.Lett.,59B,79,1975.
35. A.A.Migdal.JETP,69,1457,1975.
36. E.Brezin, J.Zinn-Justin.Phys.Rev.Lett.,36, 691,1976; Phys.Rev.B.(in print).
37. P.Olesen.Phys.Rev.,D12,3181,1975.
38. J.Shwinger.Phys.Rev.,125,397,1962;128,2425; 1962; J.H.Lowenstein,J.A.Swieoa.Ann.Phys., 68,172,1971; A.Casher, J.Kogut, L.Susskind. Phys.Rev.,D10,732,1974.
39. S.Coleman, R.Jaokiw,L.Susskind.Ann.Phys.,93, 267,1975.
40. G.'t Hooft. Nuol.Phys.,B72,461,1974; B75,461, 1974.
41. C.B.Callan Jr., N.Coote, D.J.Gross.Phys.Rev., D13,1649,1976.
42. C.R.Hagen.Nucl.Phys.,B95,477,1975.
43. K.Wilson.Phys.Rev.,D10,2445,1974; Erioe Lecture Notes 1975, CLNS-321.
44. R.Balian, M.Drouffe, C.Itzykson.Phys.Rev., D10,3376,1974.
45. J.Kogut, L.Susskind.Phys.Rev.,D11,395,1975.
46. W.Bardeen. Contributed papers.
47. S.Drell, M.Weinstein, Sh.Yankielowioz.SLAC-PUB-1752,1976.
48. H.G.Dosch, V.F.Muller, Heidelberg Preprint HD-THEP-75-21.
49. T.Banks, J.Kogut, L.Susskind.Phys.Rev., D13, 1043,1976.
50. K.Wilson.Report at Coral Conference 1976, CLNS-327.
51. A.D.Donkov, V.G.Kadyshevsky, M.D.Mateev, R.M.Mir-Kasimov, JINR Preprint E2-6992, Dubna, 1973; V.G.Kadyshevsky, M.D.Mateev, R.M.Mir-Kasimov, JINR Preprint E2-8892,1975.
52. A.A.Migdal. Chernogolovka Preprint,1976.
53. T.Appelquist, J.Carazzone, H.Kluberg-Stern, Roth M. Phys.Rev.Lett.,36,768,1976.
54. Y.P.Yao. Phys.Rev.Lett.,36,653,1976.
55. J.M.Cornwall. Tiktopoulos.Phys.Rev.Lett., 35,338,1975.
56. J.Frenkel, J.C.Taylor. Oxford Preprint 24/76.
57. P.I.Fomin, V.A.Miransky, Yu.A.Sitenko, 535, Contributed papers; P.I.Fomin, V.P.Gusynin,V.A.Miransky, Yu.A.Sitenko.Preprint ITP-75-932, Kiev.
58. T.Kinoshita, A.Ukawa.Phys.Rev.,D13,1573,1976; 940, Contributed papers.
59. C.P.Korthales Altes, E. de Rafael.CNRS,Marseille preprint,1976.
60. T.J.Kinoshita. Math.Phys.,3,650,1962.
61. T.D.Lee, M.Nauenberg.Phys.Rev.,133,1549,1964.
62. N.N.Bogolubov, Quasi-Averages in Statistioal Mechanics. JINR Preprint Dubna,1961.
63. G.'t Hooft. Nucl.Phys.,B79,276,1974; A.M.Polyakov. Pizma JETP,20,430,1974.
64. L.D.Faddeev,425. Contributed papers.
65. A.M.Polyakov.Phys.Lett.,59B,82,1975. Report on the International Conference on Nonlooal Field Theory, Alushta,1976.
66. A.A.Belavin,A.M.Polyakov,A.S.Schwartz,Yu.S.Tyupkin.Phys.Lett.,59B,85,1975.
67. R.Jaokiw,C.Rebbi.MIT-Preprint,1976.
68. C.B.Callan Jr., R.Dashen, D.J.Gross. Prinoeton Preprint COO-2220-75.
69. L.D.Faddeev.Report on the International Conferenoe on Nonlocal Field Theory, Alushta, 1976.
70. G.'t Hooft. Harvard Preprint,1976.
71. S.Coleman,E.Weinberg.Phys.Rev.,D7,1888,1973.
72. A.D.Linde. Preprint Lebedev Phys.Inst.,123,1975. S.Weinberg.Phys.Rev.Lett.,36,294,1976.
73. D.A.Kirzhnits. Contributed papers.
74. E.A.Ivanov,V.I.Ogievetsky.Pizma JETP,23,661, 1976.
75. D.J.Groes,F.Wilczek. Phys.Rev.Lett.,30,1343, 1973.
76. M.Suzuki.Nucl.Phys., B83,269,1974; N.P.Chang, Phys.Rev.,D10,2706,1974; E.Ma, Phys.Rev.,D11, 322,1975; B.L.Voronov,I.V.Tyutin.Pizma JETP; 21,396,1975; D.I.Kazakov,D.V.Shirkov.JINR Preprint E2-8974, Dubna,1975; E.S.Fradkin, O.K.Kalashnikov.Phys.Lett.,59B,159,1975.
77. Y.Nambu, G.Iona-Lasinio.Phys.Rev.,122,345,1961.
78. B.A.Arbusov,A.N.Tavkhelidze,R.N.Faustov. Doklady AN USSR,139,345,1961.
79. R.Jaokiw,K.Johnson.Phys.Rev.,D8,2386,1973; J.M.Cornwall,R.E.Norton.Phys.Rev.,D8,3338, 1973; E.J.Eichtein,F.L.Feinberg.Phys.Rev., D10,3524,1974; J.M.Cornwall,R.Jackiw,E.Tomboulis.Phys.Rev.,D10,2438,1975; E.C.Poggio, E.Tomboulis,S-H.H.Tye.Phys.Rev.,D11,2839,1975.
80. Yu.A.Golfand,E.P.Likhtman.Pizma JETP,13,452, 1971.
81. D.V.Volkov,V.P.Akulov.Pizma JETP,16, 621,1972; Phys.Lett.,B46,109,1973.
82. J.Wess,B.Zumino.Nucl.Phys.,B70,39,1974; Phys.Lett.,49B,52,1974.
83. J.Illiopoulos,B.Zumino.Nucl.Phys.,B76,310, 1974.
84. A.Salam,J.Strathdee.Phys.Rev.,D11,1521,1975; Nucl.Phys.,B76,477,1974.

85. A.Salam,J.Strathdee.Preprint IC/75/49.
P.Fayet.Nuovo Cim.,31A4,626(1976).
86. L.Mezinchesku, V.I.Ogievetsky.JINR Preprint E2-8277,Dubna,1974; Uspekhi Fiz.Nauk,117,637 1975.
87. K.Fujikawa,W.Lang.Nucl.Phys.,B88,61 ,1975.
88. V.K.Krivoshekov,A.A.Slavnov,B.Faysullaev. TMF,26,147,1976.
89. F.A.Beresin.Method of Second Quantization.
9o. J.Wess, B.Zumino.Nucl.Phys.,B78,1,1974.
91. A.Salam, J.Strathdee. Nuol.Phys.,B51,353,1974.
92. S.Ferrara, B.Zumino.Nucl.Phys.,B79,413,1974.
93. A.A.Slavnov.JINR Preprint E2-8308,1974; TMF,23,3,1975.
94. A.A.Slavnov.JINR Preprint E2-8449,Dubna,1974;
Nucl.Phys.,B97,155,1975;
B. de Wit, Stony Brook preprint ITP-8B-75-14;
S.Ferrara,O.Piguet.Nucl.Phys.,B93,261,1975;
O.Piguet,A.Rouet.Nuol.Phys.B99,458,1975;
Munich Preprint MPI-PA-36/75.
95. B.Zumino.Nucl.Phys.,B89,535,1975.
96. P.Fayet,J.Illiopoulos.Phys.Lett.,51B,461,1974.
97. P.Fayet.Phys.Lett.,58B,67,1975.
98. L.O'Raifertaigh.Nucl.Phys.,B96,331,1975;
S.Brown,L.O'Raifertaigh,T.Sherry.Preprint DINS-TP-75-16;.
99. P.Fayet.Nucl.Phys.,B90,104,1975; Mainland G.B. K.Tanaka.Phys.Rev.,D12,2394,1975.
100. W.Bardeen, in press.
101. B. de Witt, D.Z.Freedman.Phys.Rev.Lett.,35, 827,1975.
102. D.Capper, A.Salam, J.Strathdee.Preprint IC/76/16.
103. A.A.Slavnov.JINR Preprint E2-9389,Dubna,1975.
104. A.A.Slavnov.Report on the International Conference on Nonlocal Field Theory, Alushta, 1976; TMF (in press).
105. P.Fayet.Preprint PTENS76/10.
106. P.Fayet.Preprint PTENS 76/17.
107. D.V.Volkov, V.P.Akulov,TMF,18,39,1974;
P.H.Dondi,M.Sohnins.Nucl.Phys.,B81,317,1974;
A.Salam,J.Strathdee.Nucl.Phys.,B80,499,1974.
108. R.Haag,J.T.Lopuszansky,M.Sohnins.Nucl.Phys., B88,257,1975.
109. D.M.Capper,G.Leibrandt.Nucl.Phys.,B85,503, 1975; J.Wess.Acta Physica 0/3462.
110. P.Fayet.Preprint PTENS 75/18.
111. D.V.Volkov, Y.A.Soroka.Pizma JETP; 18,529, 1973.
112. R.Arnowitt,P.Nath.Phys.Lett.,56B,177,1975;
R.Arnowitt,P.Nath,B.Zumino.Phys.Lett.,56B, 81,1975.
113. R.Arnowitt, P.Nath.Harward Preprint,1976.
114. B.Zumino.Preprint TH-2120 CERN,1975.
115. D.V.Volkov,V.P.Akulov,V.A.Soroka,545, Contributed papers.
116. V.I.Ogievetsky,E.Sokachev,517,Contributed papers.
117. D.Z.Freedman, van Niewenhuizen, S.Ferrara; Preprint SUSB-ITP-76-23;
S.Deser,B.Zumino.Preprint TH-2164-CERN,1976.
118. A.Das,D.Z.Freedman.Stony Brook Preprint ITP-SB-76-35.
119. S.Ferrara, B.Zumino.Nucl.Phys.,B93,207,1975.