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ON THE VACUUM STABILITY IN THE SUPERRENORMALIZED YUKAWA-TYPE THEORY

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I. Introduction

The investigation of the vacuum stability in the framework of different field theoretical approaches presents a considerable interest [1-4]. This problem is closely tied with many questions arising in quantum field theory and concerning the structure of the vacuum. For instance, there is an indication [5] that in the case of the massless Yang-Mills theory the one-loop radiative corrections to the classical action lead to instability of the vacuum. Moreover, it was pointed out [6] that the magnetic instability of the perturbative vacuum of the non-Abelian Yang-Mills field enables us probably to solve the problem of the colour confinement in QCD. The nontrivial expectation value of the vacuum field may arise in two different ways: either on the classical level when the Lagrangian has a term like a negative squared mass [7,8] which leads to the broken symmetry of the ground state or on the quantum level when the quantum corrections can give rise to the vacuum instability [9]. The quantum corrections lead to divergences which should be removed by the renormalization procedure. As a result, the contribution of quantum fluctuations may seriously change the classical picture. This situation can be seen, for example, in the $(h\varphi^4)_4$ scalar field model [9-11]. The theory is described by the Lagrangian

$$\mathcal{L}_B(x)=\frac{1}{2}\varphi(x)[\Box_x-m^2]\varphi(x)+h\varphi^4(x), \qquad (I)$$

where h is the bare self-interacting coupling constant and m is the boson mass. In the real R^4 space-time this theory is renormalizable. The classical potential is bounded below and gives rise to a ground state with a zero expectation value of the field φ. The quantum corrections after eliminating the corresponding divergences lead [9,10] to the effective potential unbounded below, i.e. to instability of the model.

The theory under consideration may be made finite by introducing [11] into it a nonlocal interaction. Then, the boson propagator in the Euclidean space-time is given by

$$\widetilde{D}(k^2) = \frac{V(k^2\ell^2)}{m^2+k^2}. \tag{2}$$

Here $V(k^2\ell^2)$ is a nonlocal formfactor and ℓ is a parameter of nonlocality. Due to the fastly decreasing (2) as $k^2 \to \infty$ the theory has no ultraviolet divergences. The quantum contribution has a finite value and the stability of the theory and its ground state is preserved.

But the situation may be changed by introduction of an additional local Yukawa-type interaction with fermions. The Lagrangian of the fermion sector is

$$\mathcal{L}_F(x) = \overline{\Psi}(x)[i\hat{\partial} - M]\Psi(x) + g\,\overline{\Psi}(x)\Gamma\Psi(x)\varphi(x)\,, \tag{3}$$

where M and g are the fermion mass and the Yukawa coupling constant, respectively.

The fermion propagator has the standard form

$$\widetilde{S}(i\hat{k}) = \frac{1}{M - i\hat{k}}. \tag{4}$$

Here $\hat{k} = \gamma_\mu k_\mu$ and γ_μ are the Dirac matrices. The Yukawa-type model with the Lagrangian

$$\mathcal{L}(x) = \mathcal{L}'_B(x) + \mathcal{L}'_F(x) \tag{5}$$

is superrenormalizable because of the finite boson part (1). Here $\mathcal{L}'_B$ is the Lagrangian of the nonlocal boson sector.

When the quantum corrections are taken into account there arise divergences connected with fermion loops. After eliminating these divergences by renormalization one obtains a finite contribution from a boson-fermion interaction which can give rise to instability in the theory.

The aim of the present paper is to investigate the role of renormalization procedure in the vacuum instability in this model.

For a qualitative investigation of the first quantum correlations one usually uses perturbative methods, such as loop expansion [1,9] of the effective potential. In the interesting for us region of strong coupling nonperturbative methods should be used. On the way, the variational estimations [11,13,15] are more attractive, first of all, from the point of view of universality and simplicity of calculations. In this paper we will use the variational estimation method suggested in Ref. 11 .

2. The superrenormalizable Yukawa-type model

The Yukawa-type model with the boson-fermion interaction is widely used [11-15] for testing different methods and new ideas in the field theory. In the chiral-invariant Yukawa theory with scalar coupling, the fermion part of the effective potential has been investigated [12] in the one-loop approximation. It has been found that at certain ratio of coupling constants g and h the instability of the initial ground state arises. The presence of fermions encourages spontaneous symmetry breaking and destabilizes the theory. Stability requires [13] the bare g being infinitesimal. It should be noted that this contrasts with large-N studies [14] in which g remains finite. The last paper contains an indication of the existence of the phase with a nontrivial vacuum expectation value of the bosonic field.

The investigation of this problem within the framework of the modified variational Gaussian approximation [15] gives the effective potential unbounded below. The Yukawa-type theory turns out to be unstable.

All the above-mentioned papers deal with the local Yukawa-type model. In the superrenormalizable theory considered here we have another picture due to the nonlocality of the bosonic sector. This model has the boson and fermion propagators defined by (2) and (4), respectively, and the total Lagrangian is written as

$$\mathcal{L}(x)=\frac{1}{2}\varphi(x)[\Box-m^2]\varphi(x)+\bar{\Psi}(x)[i\hat{\partial}-M]\Psi(x)-h\,[K(\ell^2\Box)\varphi(x)]^4-ig\,\bar{\Psi}(x)\gamma_5\Psi(x)[K(\ell^2\Box)\varphi(x)] \quad (6)$$

where the nonlocal boson field $\phi(x)\equiv K(\ell^2\Box)\varphi(x)$ is introduced. The nonlocal operator $V(-\ell^2\Box)\equiv[K(\ell^2\Box)]^2$ obeys some special conditions (see Ref. 11). To concretize eq. (6) we choose the pseudoscalar interaction $\Gamma=i\gamma_5$ there.

The Lagrangian (6) is invariant under the transformations

$$\varphi(x)\longrightarrow-\varphi(-x)\;,$$
$$\Psi(x)\longrightarrow\gamma_0\Psi(-x)\,. \quad (7)$$

But the ground state may have no this symmetry due to the contribution of the fermionic sector.

As is well known [9,16,17] , the vacuum state of a theory is defined by the absolute minimum of the effective potential.

We investigate now the behaviour of the effective potential which is defined by [16]

$$V_{eff}[\phi_c]=-\lim_{\Omega\to\infty}\frac{1}{\Omega}\ln\langle S_\Omega[\phi_c]\rangle \quad (8)$$

where

$$\langle S_\Omega[\phi_c]\rangle=\frac{1}{N}\int\delta\varphi\,\delta\bar{\Psi}\,\delta\Psi\;\delta\Big(\phi_c-\frac{1}{\Omega}\int_\Omega dx\,\varphi(x)\Big)\exp\Big\{\int_\Omega dx\,\mathcal{L}(\varphi)\Big\}. \quad (9)$$

Here Ω is the space-time total volume and $\mathcal{L}$ is the Lagrangian given by (6). The unique point $\phi_c=\phi_0$ which corresponds to the minimum of $V_{eff}[\phi_c]$ is treated as vacuum expectation value of the field.

In the superrenormalizable Yukawa-type model under consideration, in order to remove all the divergences induced by fermionic loops it is sufficient to introduce counterterms only for the renormalization of the boson mass and the self-interaction coupling constant h .

The total Lagrangian (6) can be rewritten as

$$\mathcal{L}_R=\frac{1-Z_2}{2}\varphi[\Box-m^2]\varphi+\bar{\Psi}[i\hat{\partial}-M]\Psi-ig\bar{\Psi}\gamma_5\Psi\phi+ h(1-Z_4)\phi^4-\frac{\delta m^2}{2}\phi^2\;, \quad (10)$$

where

$$Z_2 = -g^2 \cdot \tilde{\Pi}'(-m^2) \tag{11}$$

$$Z_4 = \frac{g^4}{4} \tilde{\Gamma}(0,0,0,0)$$

$$\delta m^2 = g^2 \cdot \tilde{\Pi}(-m^2) \qquad .$$

Here

$$\tilde{\Pi}(k^2) = \int \frac{d^4p}{(2\pi)^4}\, tr\{\tilde{S}(i\hat{p})\,\tilde{S}(i\hat{p}+i\hat{k})\}$$

$$\tilde{\Gamma}(p_1,p_2,p_3,p_4) = \int \frac{d^4k}{(2\pi)^4}\, tr\,\{\tilde{S}(i\hat{p}_1+i\hat{k})\dots\tilde{S}(i\hat{p}_4+i\hat{k})\} \tag{12}$$

We make now the following regularization

$$\tilde{\Pi}_{REG}(k^2) = \tilde{\Pi}(k^2) - \tilde{\Pi}(-m^2) - (m^2+k^2)\,\tilde{\Pi}'(-m^2) \tag{13}$$

$$\tilde{\Gamma}_{REG}(p_1,\dots,p_4) = \tilde{\Gamma}(p_1,\dots,p_4) - \tilde{\Gamma}(0,0,0,0)\,.$$

Substitution of (IO)-(13) into (9) removes all the divergences in the theory.

We will calculate the integrals in (9) in the following way. The Gaussian integration over fermionic variables can be carried out exactly. Further, we employ the variational method to the remaining functional integral over the bosonic field. Finally, one obtains the estimation for the right-hand side of equation (9).

It is convenient to make the following change of the integration variables [11] in functional integral (9)

$$\psi(x) = \int dz \cdot \sigma(x-z)\cdot \eta(z) \tag{14}$$

$$\phi(x) = \int dz \cdot \Delta(x-z)\cdot u(z)\,,$$

where the functions $\sigma(x-z)$ and $\Delta(x-z)$ are defined by

$$S(x-y) = \int dz \cdot \sigma(x-z)\cdot \bar{\sigma}(z-y) \tag{15}$$

$$D(x-y) = \int dz \cdot \Delta(x-z)\cdot \Delta(z-y)\,.$$

The new integration variables $u(z)$ and $\eta(z)$ can be expanded into the basis orthonormalized functions [11] $\{g_n(x)\}$

$$u(z) = \sum_n u_n\, g_n(z) \tag{16}$$

$$\eta(z) = \sum_n \eta_n\, g_n(z) \qquad .$$

In (7) we pick out the Gaussian integration measures

$$d\sigma_\varphi = \prod_n \frac{du_n}{\sqrt{2\pi}}\, exp[-\frac{u_n^2}{2}] \tag{17}$$

$$d\sigma_\psi = \prod_n d\bar{\eta}_n\, d\eta_n\, exp[-\bar{\eta}_n \eta_n]$$

which obey the usual conditions

$$\int d\sigma_\varphi = 1 \quad , \quad \int d\sigma_\psi = 1\,. \tag{18}$$

The functional integral over fermionic fields in (9) is carried out exactly

$$\int d\sigma_\psi \cdot exp\{-ig\int_\Omega dx\, \bar{\psi}(x)\gamma_5\psi(x)\cdot[\phi(x)+\phi_c]\} = det(I+igH)\,, \tag{19}$$

where the matrix H_{mn} has the form

$$H_{mn} = \int dx\, \bar{\sigma}_m(x)\,\gamma_5\,\sigma_n(x)\cdot[\phi(x)+\phi_c]\,. \tag{20}$$

Note that in the Euclidean metric where the fermion propagator has the form (4) the matrix (20) is nonhermitian. But it is possible to substitute the fermionic propagator $\tilde{S}(i\hat{k})$ by

$$\tilde{G}(k^2) = |\tilde{S}(i\hat{k})| \qquad (21)$$

As a result, the matrix H_{mn} becomes hermitian provided that

$$G(x-y) = \int dz\, \sigma(x-z)\cdot\sigma(z-y) = \sum_n \sigma_n(x)\,\sigma_n(y)\,. \qquad (22)$$

The replacement (21) does not alter the ultraviolet behaviour of the fermionic propagator which plays the main role in our consideration.

Further, we may use the following inequality ;

$$\det(I+igH) \geqslant \exp\left\{\frac{g^2}{2}\,tr\,HH^+ - \frac{g^4}{4}\,tr\,H^4\right\} \qquad (23)$$

valid for arbitrary hermitian matrices H . Substituting (23) into the integral (9) we have

$$\langle S_\Omega[\phi_c]\rangle \geqslant \exp\left[-\Omega\left(\frac{m^2}{2}\phi_c^2 + h\phi_c^4\right)\right]\cdot\int d\sigma_\varphi \exp\Big\{\frac{g^2}{2}\,tr\,HH^+ - \frac{g^4}{4}\,tr\,H^4 + \frac{Z_2}{2}\iint_\Omega dx\,dy\,\phi(x)D^{-1}(xy)\phi(y) + \frac{Z_2}{2}m^2\phi_c^2\cdot\Omega - (h-Z_4)\int_\Omega dx\left(\phi^4(x) + 6\phi^2(x)\phi_c^2 + \phi_c^4\right) - \frac{\delta m^2}{2}\int_\Omega dx\left(\phi^2(x)+\phi_c^2\right)\Big\}\,. \qquad (24)$$

Let us introduce new integration variables $u_n \to u_n(1+q_n)^{-1/2}$, where $\{q_n\}$ are positive numbers obeying the condition $\sum_n q_n < \infty$. Using the inequality

$$\int d\sigma_\varphi \exp W[\varphi] \geqslant \exp\int d\sigma_\varphi\, W[\varphi] \qquad (25)$$

valid for real functionals $W[\varphi]$ we obtain the following estimation for (9):

$$\langle S_\Omega[\phi_c]\rangle \geqslant \exp\left\{-\frac{1}{2}\sum_n\left[\ln(1+q_n) - \frac{q_n}{1+q_n}\right]\right\}\cdot\exp\left\{-\Omega\left[\frac{m^2}{2}\phi_c^2 + h\phi_c^4\right]\right\}\times$$

$$\times\exp\int d\sigma_\varphi\Big\{\frac{g^2}{2}\,tr\,H_qH_q^+ - \frac{g^4}{4}\,tr\,H_q^4 + \frac{Z_2}{2}\iint_\Omega dx\,dy\,\phi_q(x)D^{-1}(xy)\phi_q(y) + \frac{Z_2}{2}m^2\phi_c^2\,\Omega - (h-Z_4)\int_\Omega dx\left[\phi_q^4(x) + 6\phi_q^2(x)\phi_c^2 + \phi_c^4\right] - \frac{\delta m^2}{2}\int_\Omega dx\left[\phi_q^2(x)+\phi_c^2\right]\Big\} \qquad (26)$$

where

$$\phi_q(x) = \sum_n \frac{\Delta_n(x)}{\sqrt{1+q_n}}\,u_n$$
$$H_q = (1+q)^{-1/2}\,H\,(1+q)^{-1/2}\,. \qquad (27)$$

From (26) we obtain the final estimation for the effective potential

$$V_{eff}[\phi_c] \leqslant V_+[\phi_c] = \min_{\{q_n\}}\lim_{\Omega\to\infty}\Big\{\frac{1}{2\Omega}\sum_n\left[\ln(1+q_n) - \frac{q_n}{1+q_n}\right] + \frac{m^2}{2}\phi_c^2 + h\phi_c^4 - \frac{g^2}{2\Omega}\int d\sigma_\varphi\, tr\,H_qH_q^+ + \frac{g^4}{4\Omega}\int d\sigma_\varphi\, tr\,H_q^4 - \frac{Z_2}{2\Omega}\iint_\Omega dx\,dy\,\phi_q(x)D^{-1}(xy)\phi_q(x) - \frac{Z_2}{2}m^2\phi_c^2 + (h-Z_4)\frac{1}{\Omega}\int_\Omega dx\left[\phi_q^4(x) + 6\phi_q^2(x)\phi_c^2 + \phi_c^4\right] + \frac{\delta m^2}{2\Omega}\int_\Omega dx\left[\phi_q^2(x)+\phi_c^2\right]\Big\}\,. \qquad (28)$$

The variational problem (28) with arbitrary positive numbers $\{q_n\}$ is very complicated and for simplicity we can choose q_n such that the following asymptotic inequalities are satisfied

$$\lim_{\Omega\to\infty}\frac{1}{\Omega}\sum_n\left[\ln(1+q_n) - \frac{q_n}{1+q_n}\right] = \int\frac{d^4k}{(2\pi)^4}\left[\ln(1+q(k^2)) - \frac{q(k^2)}{1+q(k^2)}\right]$$
$$\lim_{\Omega\to\infty}\sum_n\frac{\Delta_n^2(x)}{1+q_n} = \int\frac{d^4k}{(2\pi)^4}\,\frac{\tilde{D}(k^2)}{1+q(k^2)}\,, \qquad (29)$$

where $q(K^2)$ is a positive function obeying the condition $\int dq\, q(K^2) < \infty$.

Let us direct the volume to infinity: $\Omega \to \infty$. Substituting (29) into (26) and picking out the dimensional factor m^4 we obtain a variational estimation for the upper bound of the effective potential

$$V_+[\phi_c] = m^4 \cdot \min_{q>0} \Big\{ \frac{1}{2}\int\frac{d^4K}{(2\pi)^4}\Big[\ln\big(1+q(K^2)\big) - \frac{q(K^2)}{1+q(K^2)}\Big] -$$

$$\frac{g^2}{2}\int\frac{d^4K}{(2\pi)^4}\widetilde{D}_q(K^2)\widetilde{\Pi}_{REG}(K^2) + \frac{g^4}{4}\iint\frac{d^4p\,d^4K}{(2\pi)^8}\widetilde{D}_q(p^2)\widetilde{D}_q(K^2)\cdot\Big[2\cdot\widetilde{\Gamma}_{REG}(0,p,K,0) +$$

$$\widetilde{\Gamma}_{REG}(0,p,K,p+K) + \frac{12\cdot h}{g^4}\Big] + \phi_c^2\cdot\Big[\frac{1}{2} - \frac{g^2}{2}\widetilde{\Pi}_{REG}(0) + \frac{g^4}{4}\int\frac{d^4p}{(2\pi)^4}\widetilde{D}_q(p^2)\cdot$$

$$\Big[4\cdot\widetilde{\Gamma}_{REG}(0,0,0,p) + 2\cdot\widetilde{\Gamma}_{REG}(0,0,p,p) + \frac{24\cdot h}{g^4}\Big]\Big] + h\phi_c^4\Big\}, \qquad (30)$$

where

$$\widetilde{D}_q(K^2) = \frac{\widetilde{D}(K^2)}{1+q(K^2)} \qquad (31)$$

$$\widetilde{\Gamma}_{REG}(0,0,0,p) = -\frac{1}{(2\pi)^3}\int_0^1 d\alpha\sqrt{\frac{1}{\alpha}-1}\cdot\ln\Big(1+\frac{\alpha(1-\alpha)p^2}{M^2}\Big)$$

$$\widetilde{\Gamma}_{REG}(0,0,p,p) = -\frac{1}{4\pi^2}\int_0^1 d\alpha\,\ln\Big(1+\frac{\alpha(1-\alpha)p^2}{M^2}\Big)$$

$$\widetilde{\Pi}_{REG}(K^2) = \frac{1}{32\pi^2}\Big[1-K^2+\frac{8}{\pi}\int_0^1 d\alpha[\alpha(1-\alpha)]^{-1/2}\big(M^2+K^2\alpha(1-\alpha)\big)\ln\frac{M^2+K^2\alpha(1-\alpha)}{M^2+\alpha(1-\alpha)}\Big].$$

It is convenient to represent formula (30) by diagrams. If we introduce conventional notation shown in Fig.1, the right-hand side of inequality (30) can be described by the set of graphs shown in Fig.2.

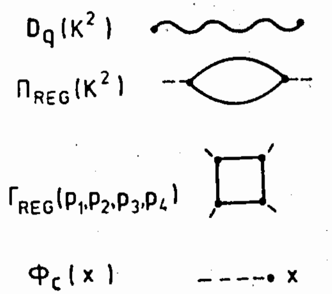

Fig.1. Notation for the Feynman diagrams.

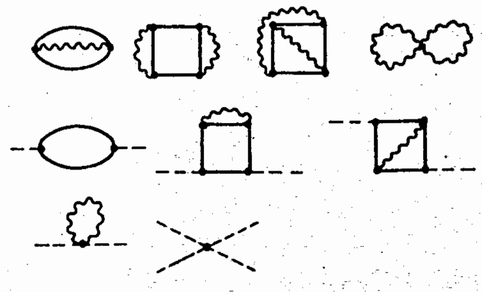

Fig.2. Graphical representation of the effective potential.

Minimization procedure applied to the right-hand side of (30) results in the complicated integral equation. But we can overcome this obstacle in the following way. Let us take the function

$$q(K^2) = f\cdot m^2\cdot\widetilde{D}(K^2), \qquad (32)$$

where $f = f(g)$ is a variational parameter. Now we look for the solution of the equation

$$\frac{dV_+}{df}\Big|_{f=f_c(g;h)} = 0 \qquad (33)$$

which obeys the condition $d^2V_+/df^2 > 0$. Substituting the solution of eq. (33) into (30) and taking into account (29) we finally get

$$V_+[\phi_c] = m^4\cdot\{C_0 + a(g;h;f_0)\cdot\phi_c^2 + h\phi_c^4\}. \qquad (34)$$

The field configuration $\phi_c = \phi_0$ giving the absolute minimum of the potential (34) corresponds to the vacuum state of the system. We see in (34) that the potential is bound below at an arbitrary constant $h > 0$, i.e., the theory remains stable. The factor $a(g;h;f_0)$ is given by

$$a(g;h;f_0) = \frac{1}{2} - \frac{g^2}{2}\widetilde{\Pi}_{REG}(0) + \frac{g^4}{4}\int\frac{d^4k}{(2\pi)^4}\widetilde{D}_{f_0}(p^2)\cdot$$
$$\left[4\cdot\widetilde{\Gamma}_{REG}(0,0,0,p) + 2\cdot\widetilde{\Gamma}_{REG}(0,0,p,p) + \frac{24h}{g^4}\right]. \qquad (35)$$

From (28) it is easy to see that $a(g;h;f_0)$ may have different signs at fixed h, i.e.

$$a(g;h;f_0) \geqslant 0 \quad \text{at} \quad g \leqslant g_{crit}(h)$$

and

$$a(g;h;f_0) < 0 \quad \text{at} \quad g > g_{crit}(h) .$$

So, in this theory the initial vacuum stability may be broken and there exist two phases with different vacuum expectation values of the field. It should be noted that here the second order phase transition takes place.

The qualitative behaviour of the effective potential $V_+[\phi_c]$ at different values of g is shown in Fig.3. At small ϕ_c the potential V_+ behaves as $\sim\phi_c^2$, $\sim\phi_c^4$ and $\sim-\phi_c^2$ for curves 1,2 and 3, respectively (see Fig.3).

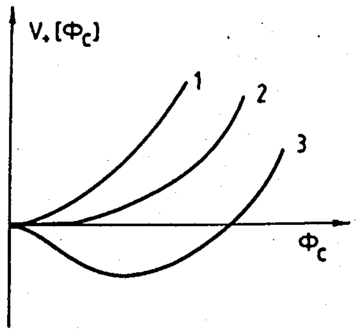

Fig.3. The behaviour of $V_+[\phi_c]$ for fixed h at different g. The curves 1,2 and 3 correspond to the coupling constants g', g'' and g''', respectively, where $g' < g'' = g_{crit} < g'''$

3. Conclusion

We have investigated the appearance of the instability in the superrenormalizable Yukawa-type field theory which is an extension of the $h\varphi^4$ scalar field model with a nonlocal self-interaction. At the infinitesimal coupling constant g we have a finite theory. In this case, the vacuum state is stable and has a trivial expectation value of the field. An additional interaction with fermions breaks the vacuum stability. As a result, a finite critical value for the boson-fermion coupling constant arises. For $g<g_{crit}$ the interaction is small and the system is found in the phase with "normal" vacuum. When the constant g increases the fermionic destabilizing role becomes more important and for $g>g_{crit}$ the boson field φ acquires nontrivial vacuum expectation value and the system goes into a new phase. Here, the second order phase transition takes place. Thus, in the renormalized theory under consideration the symmetry $\varphi\to-\varphi$ turns out to be spontaneously broken.

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References

1. S.G.Matinyan and G.K.Savvidy, Soviet Journ.Yad.Fiz. 25 (1977) 218.
2. H.Leutwyler, Nucl. Phys. B179 (1981) 129.

3. A.A.Grib, S.G.Mamayev and V.M.Mostepanenko, Vacuum quantum effects in strong fields (Energoatomizdat, Moscow, 1988).

4. P.M.Stevenson and R.Tarrach, Phys.Lett. B176 (1986) 436.

5. G.K.Savvidy, Phys.Lett. B71 (1977) 133.

6. G.Preparata, Phys.Lett. B201 (1988) 139, Nuovo Cim. 96A (1986) 366 .

7. J.Goldstone, A.Salam and S.Weinberg, Phys.Rev. 127 (1962) 965.

8. P.Higgs, Phys.Lett. 12 (1964) 132.

9. S.Coleman and E.Weinberg, Phys.Rev. D7 (1973) 1888.

10. P.M.Stevenson, Phys.Rev. D32 (1985) 1389.

11. G.V.Efimov, Problems of quantum theory of nonlocal interactions (Nauka, Moscow, 1985).

12. P.Ramond, Field Theory: A Modern Primer (Benjamin/Cummings, London, 1981).

13. P.M.Stevenson and G.A.Hajj, Phys.Rev. D34 (1986) 3117.

14. B.Basseti, F.Riva and D.Zanon, Phys.Lett. B100 (1981)407.

15. I.Yotsuyanagi, Kanazawa University report No. DPKU-8701, 1987.

16. R.Fukuda and E.Kyriakopoulas, Nucl.Phys. B85 (1975) 354.

17. L.O'Raifeartaigh, A.Wipe and H.Yoneyama, Nucl.Phys. B271 (1986) 653.

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Нестабильность вакуума в суперперенормируемой теории Юкавы

Исследуется устойчивость основного состояния и возможность появления фазового перехода в суперперенормируемой нелокальной модели Юкавы. Получена вариационная оценка сверху для эффективного потенциала. Показано существование конечного критического значения для константы бозон-фермионного взаимодействия. При более сильной связи исходный вакуум теряет устойчивость и система переходит к новой фазе с ненулевым вакуумным средним значением поля.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

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On the Vacuun Stability in the Superrenormalized Yukava-Type Theory

The stability of the ground state and the possibility of appearance of the phase transition in the superrenormalizable nonlocal Yukawa-type field theory are investigated. A variational estimation of the upper bound for the effective potential is obtained. It is shown that there exists a finite critical value for the boson-fermion coupling constant. The initial vacuum becomes unstable when this coupling constant exceeds the critical value. As a result, the system under consideration goes into the phase with nonvanishing expectation value of the field.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

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