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ЕРЕВАНСКИЙ ФИЗИЧЕСКИЙ ИНСТИТУТ
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FIVE-DIMENSIONAL SUPERSYMMETRICAL EXPANSION OF DIRAC AND CHIRAL ELECTRODYNAMICS

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Էղ. ԵԳՈՐՅԱՆ, Ռ.Պ. ՄԱՆՎԵԼՅԱՆ

ԴԻՐԱԿՅԱՆ ԵՎ ԿԻՐԱԼ ԷԼԵԿՏՐԱԴԻՆԱՄԻԿԱՆԵՐԻ ՀՆԳԱՉԱՓ ԳԵՐՀԱՄԱՉԱՓ ԸՆԴԱՐՁԱԿՈՒՄԸ

Դիրակյան և կիրալ էլեկտրադինամիկաների համար աշխատանքում առաջարկված է գերինվարիանտ հնգաչափ քվանտային գործողություն: Ապացուցվում է, որ դեպի քառաչափ տարածություն որոշակի ռեդուկցիայի դեպքում, դիրակյան և կիրալ էլեկտրադինամիկաների ընդունված, ոչ կանոնավորված քառաչափ Ֆեյնմանյան դիագրամները վերականգնվում են: Սակայն, կիրալ /դիրակյան/ Ֆերմիոնների հնգաչափ տեսության փոխազդեցության մեջ մտնող գերհոսանքը ոչ անոմալ է և հետևաբար, հնարավոր է, որ հնգաչափ գործողությունը ծառայի որպես լավ քվանտային գործողություն կիրալ էլեկտրադինամիկայի քառաչափ տեսության համար:

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Эд.Ш.ЕГОРЯН, Р.П.МАНВЕЛЯН

ПЯТИМЕРНОЕ СУПЕРСИММЕТРИЧНОЕ РАСШИРЕНИЕ ДИРАКОВОЙ И КИРАЛЬНОЙ ЭЛЕКТРОДИНАМИК

В работе предлагается суперинвариантное пятимерное квантовое действие для дираковской и киральной электродинамик. Доказывается, что при определенной редукции в четырехмерное пространство стандартные нерегуляризованные четырехмерные фейнмановские диаграммы дираковской и киральной электродинамик восстанавливаются. Однако супертох, входящий во взаимодействие пятимерной теории киральных (дираковских) фермионов не аномален и поэтому пятимерное действие возможно может служить хорошим квантовым действием для четырехмерной теории киральной электродинамики.

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E.Sh. EGORIAN, R.P. MANVELYAN

FIVE-DIMENSIONAL SUPERSYMMETRICAL EXPANSION OF DIRAC AND CHIRAL ELECTRODYNAMICS

A superinvariant five-dimensional quantum action for Dirac and chiral electrodynamics is proposed in this paper. It is proved that certain reduction to a four-dimensional space will restore the standard unregularized four-dimensional Feynman diagrams of Dirac and chiral electrodynamics. But the supercurrent included in the interaction of five-dimensional theory of chiral (Dirac) fermions is not anomalous, and hence, the five-dimensional action can possibly be a good quantum action for the four-dimensional theory of chiral electrodynamics.

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Yerevan 1987

1. Introduction

In ref.[1] an alternative to the Neveu-West [2] second quantized theory of a spinless relativistic particle is proposed and a five-dimensional field theory with linear supersymmetry is obtained. The quantum field depends not only on the four-dimensional x_μ , but on additional bosonic t (dimension - 2) and Grassmanian $\theta, \bar\theta$ (dimension - 1) coordinates too. The Grassmanian $\theta, \bar\theta$ and bosonic t coordinates correspond to ghost, antighost fields and Lagrange multiplier t in the first quantized thèory, respectively. This linear supersymmetry corresponds to the reparametrization quantum BRST symmetry.

In the spin 1/2 case this linear supersymmetry arises too [3].

One can suppose that an arbitrary quantum field theory results from the second quantization of a relativistic invariant system of particles. Then the above mentioned linear supersymmetry can be realized for all quantum field-theories.

But the particle systems corresponding to all field-theories (in particular, to Dirac and chiral electrodynamics) is not yet found, and we make a hypothesis that a five-dimensional supersymmetric quantum Lagrangian corresponding to an arbitrary four-dimensional field theory does exist. Then, the four-dimensional field theory results from the five-dimensional one after reduction which concludes in an integration over non-physical fields.

The hypothesis mentioned above is closely connected also with the stochastic quantization scheme of Parisi-Wu. There is hidden one-dimensional linear supersymmetry in this scheme [4,5]. This supersymmetry becomes explicit at transition from Langevin equation to a superfield formulation [4]. A five-dimensional supersymmetric theory arises in result, which reduces to a standard four-dimensional theory. Analogous transition from Langevin equations to a five-dimensional supersymmetric theory takes place also in case of the fermion theory in external classical boson field [6].

However, there are some problems in case of the fermion theory in a quantum electromagnetic field. For united stochastic quantization of the fermion and electromagnetic fields, one must introduce some operator $\hat{K}$ of dimension 1 into the Langevin equation for fermions [7,8]. If $\hat{K} = i\not{\partial}$ (noncovariant) is chosen, then the corresponding Langevin equations bring to a supersymmetric formulation, but when $\hat{K} = i\not{D}$ (covariant) is chosen, there is no supersymmetric transition.

The last example shows that the supersymmetry of a quantum action seems not always to be the result of stochastic quantization. The supersymmetry principle seems to us to be more fundamental than the Langevin equation, because just the supersymmetry provides proper reduction to a four-dimensional space [4,9] as well as renormalization of expanded theory [10,11].

After the aforesaid, it is interesting to study the ability of our hypothesis in a nontrivial (i.e. when there is good quantum action) case like the anomalous theory of chiral electrodynamics. Here we shall put down some supersymmetrical five-dimensional Lagrangian which on a naive level (i.e. on the level of unregularized Feynman diagrams) is reduced to the four-dimensional chiral electrodynamics. The interaction supercurrent of the five-dimensional supersymmetrical Lagrangian (it does not coincide with the four-dimensional anomalous current) on the quantum level is conserved, i.e. is not anomalous and hence, the five-dimensional theory is renormalizable. Since a formal, naive reduction of five-dimensional supertheory to four-dimensional chiral electrodynamics took place, it is natural to consider the five-dimensional renormalized theory as definition of the non-renormalizable four--dimensional one. But the problem of the reduction of regularized and renormalized theory is still open.

2. Five-Dimensional Superexpansion

We assume that Dirac (or chiral) fermions and the Abelian vector field additionally depend on the bosonic (t) and Grassmanian (θ, $\bar{\theta}$) coordinates, i.e. we introduce superfields $\Psi(x_\mu, t, \theta, \bar{\theta})$, $\bar{\Psi}(x_\mu, t, \theta, \bar{\theta})$, $\Phi_\mu(x_\mu, t, \theta, \bar{\theta})$. We do not introduce the index of the fermion field chirality, for our formulae are correct for both Dirac and chiral fermions. The superfields have the following components:

$$\Phi_\mu(x,t,\theta,\bar\theta) = A_\mu(x,t) + \bar\theta\psi_\mu(x,t) + \bar\Psi(x,t)\theta + \bar\theta\theta C_\mu(x,t)$$
$$\Psi(x,t,\theta,\bar\theta) = \Psi(x,t) + \bar\theta\varphi_1(x,t) + \theta\varphi_2(x,t) + \bar\theta\theta\omega(x,t)$$
$$\bar\Psi(x,t,\theta,\bar\theta) = \bar\Psi(x,t) + \theta\bar\varphi_1(x,t) + \theta\bar\varphi_2(x,t) + \bar\theta\theta\bar\omega(x,t) \qquad (1)$$

The first components of these superfields at fixed t_0 will be physical (our theory will be translation-invariant over t, so the reduction will be independent of t_0 and one can take $t_0 = 0$).

Let us postulate the following superinvariance of the quantum theory:

$$t \to t - \bar\theta\varepsilon$$
$$\theta \to \theta + \varepsilon \qquad (2)$$
$$\bar\theta \to \bar\theta + \bar\varepsilon$$

The corresponding field transitions read:

$$\delta\Phi_\mu = (\varepsilon Q + \bar\varepsilon\bar Q)\Phi_\mu \ , \quad \delta\Psi = (\varepsilon Q + \bar\varepsilon\bar Q)\Psi$$
$$\delta\bar\Psi = (\varepsilon Q + \bar\varepsilon\bar Q)\bar\Psi \qquad (3)$$

where

$$Q = \partial_\theta + \bar\theta\partial_t$$
$$\bar Q = \partial_{\bar\theta} \qquad (4)$$
$$\{Q, \bar Q\} = \partial_t$$

This supersymmetry is the chiral representation of a usual one-dimensional supersymmetry and is characteristic of the stochastic quantization [4].

Let us write down the following supersymmetric action:

$$S^{(\hat{K})}=\int \mathcal{L}^{(K)}dt\,d\theta\, d\bar{\theta}\, d^4x$$

$$\mathcal{L}^{(\hat{K})}=-\frac{1}{4}(\partial_\mu \varphi_\nu-\partial_\nu \varphi_\mu)^2+\frac{1}{2\alpha}(\partial_\mu \varphi_\mu)^2+\frac{1}{2}\varphi_\mu B \varphi_\mu \tag{5}$$
$$+\bar{\Psi} i \not{\mathcal{D}}\Psi + D\bar{\Psi}(i\hat{K})^{-1}\bar{D}\Psi - \bar{D}\bar{\Psi}(i\hat{K})^{-1}D\Psi$$

where $\mathcal{D}_\mu=\partial_\mu - ie\varphi_\mu$, $\hat{K}$ is an operator of dimension 1 (e.g., $\hat{K}=\not{\partial},\not{\mathcal{D}}$), and $D, \bar{D}$ are supercovariant derivatives:

$$\bar{D}=\partial_{\bar{\theta}}-\theta\partial_t\ , \qquad D=\partial_\theta\ , \qquad \{D,\bar{D}\}=-\partial_t\ ,$$
$$B\equiv[D,\bar{D}]$$

The action (5) is assumed to be Euclidean over x_μ . $\hat{K}$ must be chosen proportional to one γ -matrix in order the last two terms in (5) do not go to zero in case of chiral fermions. We shall discuss two possible Lorentz-covariant variants for the operator $\hat{K}$: $\hat{K}=\not{\partial}$ and $\hat{K}=\not{\mathcal{D}}$.

Note, that without the gauge-fixing second term in $\mathcal{L}^{(\hat{K})}$ the action (5) at $\hat{K}=\not{\mathcal{D}}$ is invariant under the gauge transformations independent of t , θ , $\bar{\theta}$:

$$\varphi'_\mu(x,t,\theta,\bar{\theta})=\varphi_\mu(x,t,\theta,\bar{\theta})+\partial_\mu\alpha(x)$$
$$\Psi'(x,t,\theta,\bar{\theta})=\exp\{i\alpha(x)\}\,\Psi(x,t,\theta,\bar{\theta}) \tag{6}$$
$$\bar{\Psi}'(x,t,\theta,\bar{\theta})=\bar{\Psi}(x,t,\theta,\bar{\theta})\exp\{-i\alpha(x)\}$$

In result, the theory with the action $S^{\not{\mathcal{D}}}$, as we shall see in section 4, has conserved current, that is why the action $S^{\not{\mathcal{D}}}$ we shall call covariant.

The actions $S^{\not{\partial}}$ and $S^{\not{\mathcal{D}}}$ in a five-dimensional space describe different theories, but when reduced to a four-dimensional space they bring to the same theory (Dirac or chiral electrodynamics).

The fact that the theory with $S^{\not{\partial}}$ action is reduced to electrodynamics can be proved in two ways: first - the theory with $S^{\not{\partial}}$ action can be obtained from the Langevin equations by transformation to the superfield formalizm as done in refs. [4,6], second - the proof from ref.[4] can be used here.

In this paper we shall prove that on the naive level the theory with $S^{\not{\mathcal{D}}}$ action is reduced to one obtained after the reduction of the theory with $S^{\not{\partial}}$.

3. Reduction

Our purpose is to prove that the producing functional

$$Z(H_\mu,J,\bar{J})=\int \exp\left\{-\int[\mathcal{L}^{\not{\mathcal{D}}}-\Phi_\mu H_\mu-\bar{J}\Psi-\bar{\Psi}J]d^4x\,dt\,d\theta\, d\bar{\theta}\right\}\mathcal{D}A_\mu\mathcal{D}\Psi\mathcal{D}\bar{\Psi} \tag{7}$$

at the following choice of sources

$$H^0_\mu=h_\mu(x)\bar{\theta}\theta\delta(t),\quad J^0=J(x)\bar{\theta}\theta\delta(t),\quad \bar{J}^0=\bar{J}(x)\bar{\theta}\theta\delta(t) \tag{8}$$

is reduced to the producing functional of the four-dimensional electrodynamics with the sources of the vector field $h_\mu(x)$ and fermion fields $J(x)$, $\bar{J}$) :

$$Z(H^0_\mu,J^0,\bar{J}^0)=\int \exp\left\{-\int d^4x\left[-\frac{1}{4}F^2_{\mu\nu}+\frac{1}{2\alpha}(\partial_\mu A_\mu)^2+\right.\right.$$
$$\left.\left.+\bar{\Psi}(x)\not{\mathcal{D}}\Psi(x)+h_\mu(x)A_\mu(x)+J(x)\Psi(x)+\bar{\Psi}(x)J(x)\right]\right\}\mathcal{D}A_\mu\,\mathcal{D}\Psi\,\mathcal{D}\bar{\Psi} \tag{9}$$

The choice of sources like (8) in the five-dimensional producing functional (7) is called reduction. Thus, reduction means to calculate averages from the first components of superfields at $t = 0$.

To prove the statement (9), let us remove nonlocality in the eq.(7) by introducing fields χ , $\hat{\chi}$, η , $\bar{\eta}$. It will result in

$$Z=\int \exp\left\{-\mathcal{L}_Q + H_\mu \varphi_\mu + \bar{J}\Psi + \bar{\Psi}J\right\} \mathcal{D}\varphi_\mu \mathcal{D}\Psi \mathcal{D}\bar{\Psi} \mathcal{D}\chi \mathcal{D}\bar{\chi} \mathcal{D}\eta \mathcal{D}\bar{\eta} \qquad (10)$$

where integration over θ , $\bar{\theta}$, x , t is supposed in the exponent, and $\mathcal{L}_Q$ is determined by

$$\mathcal{L}_Q = -\frac{1}{4} F_{\mu\nu} F_{\mu\nu} + \frac{1}{2\alpha} (\partial_\mu \varphi_\mu)^2 + \frac{1}{2} \varphi_\mu B \varphi_\mu + \bar{\Psi} i \mathcal{D} \Psi +$$

$$+ \bar{\chi} i \not{\mathcal{D}} \chi + \bar{\eta} i \not{\mathcal{D}} \eta + \bar{\chi} D \Psi + \bar{D} \bar{\Psi} \chi - \bar{\eta} \bar{D} \Psi + D \bar{\Psi} \eta \qquad (11)$$

Let us prove that the vertices $\bar{\chi} \not{\mathcal{D}} \chi$ and $\bar{\eta} \not{\mathcal{D}} \eta$ do not contribute into the reduced producing functional Z(H^0, J^0, J^0). It will mean that effective contribution of $[i\not{\mathcal{D}}]^{-1}$ in the initial Lagrangian coincides with $[i\not{\partial}]^{-1}$. Let us write out all the fermion propagators

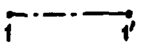
$$\underset{1}{\bullet}\!\!\longleftarrow\!\!\underset{1'}{\bullet} \equiv \langle \Psi(p,\omega,\theta)\bar{\Psi}(p',\omega',\theta')\rangle = \not{p}(p^4+\omega^2)^{-1}[2+(p^2+i\omega)\overline{\theta-\theta}'\theta +$$
$$+(p^2-i\omega)\overline{\theta'-\theta}\theta']\delta(p+p')\delta(\omega+\omega')$$

$$\underset{1}{\bullet}\!\!\longleftarrow\!\!\underset{1'}{\bullet} \equiv \langle \chi(1)\bar{\chi}(1')\rangle = \not{p}^{-1}\{\overline{\theta-\theta}'(\theta-\theta')-(p^2-i\omega)^{-1}[1-i\omega\theta\overline{\theta-\theta}']\}$$
$$\delta(p+p')\delta(\omega+\omega')$$

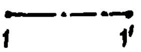
$$\underset{1}{\bullet}\!\!\longleftarrow\!\!\underset{1'}{\bullet} \equiv \langle \eta(1)\eta(1')\rangle = \not{p}^{-1}\{\overline{\theta-\theta}'(\theta-\theta')+(p^2+i\omega)^{-1}[-1+i\omega\theta\overline{\theta-\theta}']$$
$$\delta(p+p')\delta(\omega+\omega')$$

$$\underset{1}{\bullet}\!\!\longleftarrow\!\!\underset{1'}{\bullet} \equiv \langle \chi(1)\bar{\Psi}(1')\rangle = (p^2-i\omega)^{-1}\overline{\theta-\theta}'\delta(p+p')\delta(\omega+\omega')$$

$$\underset{1}{\bullet}\!\!\longleftarrow\!\!\underset{1'}{\bullet} \equiv \langle \Psi(1)\bar{\chi}(1')\rangle = (p^2-i\omega)^{-1}(\theta-\theta'+i\omega\theta(\overline{\theta-\theta'})\theta')\delta(p+p')\delta(\omega+\omega')$$



$$\equiv \langle \eta(1)\bar{\Psi}(1')\rangle = (p^2+i\omega)^{-1}(\theta-\theta'+i\omega\theta\,(\overline{\theta-\theta'})\,\theta')\,\delta(p+p')\,\delta(\omega+\omega')$$



$$\equiv \langle \Psi(1)\bar{\eta}(1')\rangle = (p^2+i\omega)^{-1}\,\overline{\theta'-\theta}\,\delta(p+p')\,\delta(\omega+\omega')$$

$$\langle \chi(1)\bar{\eta}(1')\rangle = \langle \eta(1)\bar{\chi}(1')\rangle = 0$$

Fig.1

where Fourier transformation over x and t is done.

Now consider an arbitrary diagram in the reduced theory:

$$Z(H^\circ_\mu, J^\circ, \bar{J}^\circ) = \int \exp\left\{-\mathcal{L}_Q + H^\circ_\mu \Phi_\mu + \bar{J}^\circ\Psi + \bar{\Psi}J^\circ\right\}\mathcal{D}\Phi_\mu$$

$$\mathcal{D}\Psi\,\mathcal{D}\bar{\Psi}\,\mathcal{D}\chi\,\mathcal{D}\bar{\chi}\,\mathcal{D}\eta\,\mathcal{D}\bar{\eta}\,. \tag{12}$$

where integration over $d^4x\,dt\,d\theta\,d\bar{\theta}$ is supposed in the exponent. Let us extract in such diagrams all the possible vertices which include χ , $\bar{\chi}$, η , $\bar{\eta}$:

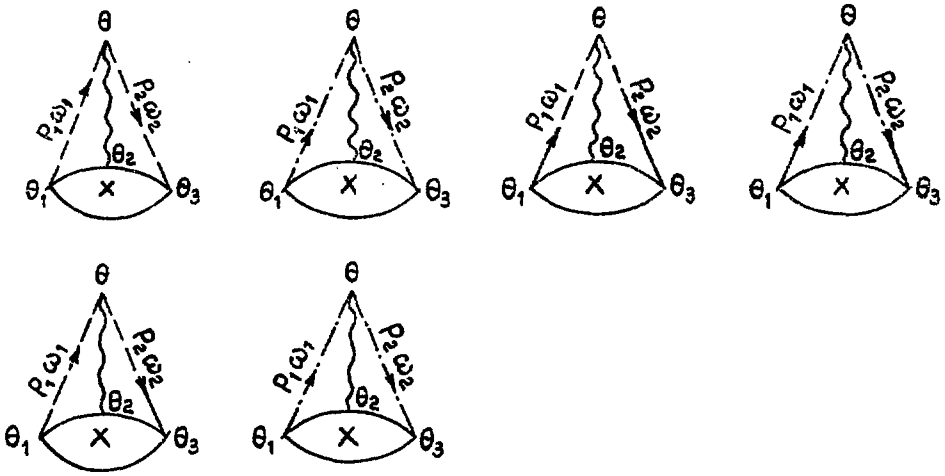


Fig.2

The diagrams of the theories (12) depend only on external four-momenta, and integration over all the additional momenta and Grassmanian coordinates $\theta, \bar{\theta}$ is made. Presenting δ-functions of conservation of the additional ω momenta by means of an exponent, one obtains the following general expression for the diagrams in Fig.2 :

$$\int \frac{d\omega_1\, d\omega_2\, d\omega_3}{(p_1^2 \pm i\omega_1)^{\alpha_1} (p_2^2 \pm i\omega_2)^{\alpha_2}} X(\omega_i, p_i) \exp\{it(\omega_1 - \omega_2 + \omega_3)\} dt \qquad (13)$$

where the function $X(\omega_i, p_i)$ corresponds to the three-vertex subdiagram (block) and is analytical over ω_1 , ω_2 ; α_i in the denominator of the eq.(13) take the values 0 and 1 depending on particular diagram. Substituting $\omega_2 \to -\omega_2$ in (13), one obtains

$$\int \frac{d\omega_1\, d\omega_2\, d\omega_3}{(p_1^2 \pm i\omega_1)^{\alpha_1} (p_2^2 \mp i\omega_2)^{\alpha_2}} X(\omega_i, p_i) \exp\{it(\omega_1 + \omega_2 + \omega_3)\} dt \qquad (14)$$

Now it is easy to understand that the integral (14) turns to zero. At $t > 0$ the integration contour must be closed in the upper half-plane (due to the exponent). Then, due to oppcsite signs of ω_1 and ω_2 in the numerator of (14), the integral over one of them turns to zero. At $t < 0$ the contour must be closed in the lower half-plane and the integral over the other ω momentum turns to zero. So, only the vertex $\bar{\Psi} \Phi_\mu \gamma_\mu \Psi$ contributes to the producing functional (12), i.e. the S^{∂} part of the covariant action $S^{\mathcal{D}}$ works only in the reduction and hence, we shall have the four-dimensional electrodynamics.

Thus, our five-dimensional "covariant supersymmetrical theory after naive (unregularized) reduction gives the usual four-dimensional quantum averages.

4. The Conserved Current

Let us make the following substitutions of variables

$$\delta\Psi(x,t,\theta,\bar\theta)=i\alpha(x)\Psi(x,t,\theta,\bar\theta),\ \delta\bar\chi(x,t,\theta,\bar\theta)=-i\alpha(x)\bar\chi(x,t,\theta,\bar\theta)$$
$$\delta\bar\Psi(x,t,\theta,\bar\theta)=-i\alpha(x)\bar\Psi(x,t,\theta,\bar\theta),\ \delta\eta(x,t,\theta,\bar\theta)=i\alpha(x)\eta(x,t,\theta,\bar\theta) \tag{15}$$
$$\delta\chi(x,t,\theta,\bar\theta)=i\alpha(x)\chi(x,t,\theta,\bar\theta),\ \delta\bar\eta(x,t,\theta,\bar\theta)=-i\alpha(x)\bar\eta(x,t,\theta,\bar\theta)$$

in the producing functional (10) at zero sources. The following current conservation is obtained in result:

$$\partial_\mu\int dt\,d^2\theta\,J_\mu(x,t,\theta,\bar\theta) \tag{16}$$

where the supercurrent J_μ is determined by

$$J_\mu=\bar\Psi\gamma_\mu\Psi+\bar\chi\gamma_\mu\chi+\bar\eta\gamma_\mu\eta \tag{17}$$

The eq.(16) is satisfied on the quantum level too (for both Dirac and chiral fermions), as, due to the equality of boson and fermion components in the superfield, the transformations (15) are not anomalous, i.e. no Fujikawa determinant appears. Due to existence of conserved current, the theory with the producing functional (10) is renormalizable (the details will be published). Thus, the five-dimensional supersymmetrical theory with Lagrangian (11) is renormalizable and in case of chiral fermions can be a good quantum Langevian for the four-dimensional chiral electrodynamics. It should be pointed out, that the conservation of the supercurrent J_μ

in the quantum level does not contradict to the anomaly of the four-dimensional chiral current j_μ^5 . This anomaly does not disappear [12-15] . Just j_μ^5 is no more the current of interaction. Most likely it is an external current (like the axial current in the Dirac electrodynamics) and that is why does not interfere in the renormalizability.

In conclusion we want to explain how the five-dimensional renormalizable theory (11) of chiral electrodynamics can bring after reduction to a non-renormalizable and non-unitary four--dimensional one. For this purpose, it is necessary to investigate the reduction of the regularized and renormalized theory. We think, that just the anomaly of the four-dimensional current j_μ^5 plays the crucial role in the regularized reduction, and non-polynomial counterterms arise after the integration over the non-physical fields. This brings to the non-renormalizability of the four-dimensional theory. The non-unitarity of the four-dimensional theory is possibly connected with the non-zero transition $\Psi \to \chi, \eta$ on the renormalized level. The investigation of the regularized and renormalized reduction of the theory (11) is under way.

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