Preprint YERPHI-1143(20)-89

ԵՐԵՎԱՆԻ ՖԻԶԻԿԱՅԻ ԻՆՍՏԻՏՈՒՏ
ЕРЕВАНСКИЙ ФИЗИЧЕСКИЙ ИНСТИТУТ
YEREVAN PHYSICS INSTITUTE

R.P.GRIGORIAN, I.V.TYUTIN

MULTIPLICATIVE RENORMALIZATION OF N=1 SUPERSYMMETRIC YANG-MILLS THEORIES: THE SU(2) GROUP

ЦНИИатоминформ
ЕРЕВАН - 1989

Նախնատիպ ԵՓԻ-1143(20)-89

Ռ.Պ. ԳՐԻԳՈՐՅԱՆ, Ի.Վ. ՏՅՈՒՏԻՆ*

N=1 ԳԵՐՀԱՄԱՉԱՓ ՅԱՆԳ-ՄԻԼՍՅԱՆ ՏԵՍՈՒԹՅՈՒՆՆԵՐԻ ԲԱԶՄԱՊԱՏԿԱԿԱՆ ՎԵՐԱՆՈՐՄԱՎՈՐՈՒՄԸ. SU(2) ԽՈՒՄԲԸ

Ապացուցված է, որ գերտրամաչափական մուլտիպլետի հատուկ ընտրված պարամետրիզացիայում, գերհամաչափ տրամաչափությունում գտնվող, չընդարձակված (N=1) գերհամաչափություն ունեցող ընդհանուր կարգի գերտրամաչափական տեսությունները կարելի է բազմիցս վերանորմավորել։

Երևանի Ֆիզիկայի ինստիտուտ
Երևան 1989

---

*Մոսկվայի էլեկտրոնային մեքենաշինության ինստիտուտ

Preprint YERPHI-1143(20)-89

R.P. GRIGORYAN, I.V. TYUTIN*

MULTIPLICATIVE RENORMALIZATION OF N=1 SUPERSYMMETRIC YANG-MILLS THEORIES: THE SU(2) GROUP

In supersymmetric gauge in a specially chosen parametrization of a supergauge multiplet the multiplicativity of renormalization of general-class supergauge theories with nonexpanded (N=1) supersymmetry is proved.

Yerevan Physics Institute

Yerevan 1989

* Moscow Institute of Electronic Engineering

Препринт ЕФИ-1143(20)-89

Р.П.ГРИГОРЯН, И.В.ТЮТИН*

МУЛЬТИПЛИКАТИВНАЯ ПЕРЕНОРМИРОВКА N = 1 СУПЕРСИММЕТРИЧНЫХ ЯНГ-МИЛЛСОВСКИХ ТЕОРИЙ: ГРУППА SU(2)

Доказана мультипликативность переномировки суперкалибровочных теорий общего класса с нерасширенной (N = 1) суперсимметрией в суперсимметричной калибровке в специально выбранной параметризации суперкалибровочного мультиплета.

Ереванский физический институт

Ереван 1989

* Московский институт электронного машиностроения

1. This paper is devoted to prove multiplicativity of re-normalization of general-class supergauge theories with nonexpanded (N=1) supersymmetry [1,2] in supersymmetric gauge in a specially chosen parametrization of a supergauge multiplet. Many works are devoted to the problem of the renormalizability of supersymmetric theories. Limiting ourselves to theories with N=1, let us present the results of refs [3,4], where the renormalization structure was investigated for the general-class supergauge theories (including matter fields) in supersymmetric gauge as well as in the Wess-Zumino gauge [1]. It was in particular shown that in both gauges renormalization conserves the supergauge and supersymmetries. In the Wess-Zumino gauge the renormalization is a multiplicative one, but in the supersymmetric gauge the renormalization, which is multiplicative for all the parameters and fields except for the gauge supermultiplet and which for the gauge supermultiplet has the form of a nonlinear substitution of variables, does not contradict to the Ward's identities. Such situation is due to dimensionless gauge supermultiplet and is typical for all the theories with nonlinearly realized symmetry. In particular, renormalization of general - class two - dimensional chiral theories is also accompanied by nonlinear reparametrization [5]. Nevertheless, in some two-dimensional symmetric chiral models (e.g., $O(n)/O(n-1)$ [6], $SU(n) \times SU(n)$ [7])

the renormalization multiplicativity has been proved. The proof was performed for special parametrization of theory using the Zinn-Justin method [8]. By modifying the scheme [7] it was later proved that a wide class of two-dimensional chiral theories could be multiplicatively renormalized [9].

The main problem arising is that of the choice of a new parametrization.

The above-mentioned considerations as well as similarity of the structures of renormalization of two-dimensional chiral and supergauge theories allow one to assume that in supergauge N=1 theories too, by suitable choice of parametrization one can achieve an overall multiplicativity of renormalization in the supersymmetric gauge. The content of this paper just consists in the solution of this task. The form of the supergauge multiplet parametrization will be described below. The paper is planned as follows. In section 2 the modified action S of the theory is constructed, an equation for the producing Γ-functional resulting from the BRST-symmetry is derived and its consequences are studied to find the structure of theory divergences. In section 3 the Ward's identities resulting from global gauge symmetry are analyzed for the same purpose. In section 4, on the basis of the results from the sections 2 and 3, the multiplicativity of renormalization for the case with the gauge symmetry SU(2) group and a concrete parametrization is proved. We shall as far as possible use condensed notations. Derivatives with respect to the fields are the right hand ones, and those with respect to the sources are the left hand ones.

2. Thus, we shall consider a theory containing chiral matter superfields $\Psi_{\pm}^{i}(x,\theta)$ as well as Hermitian fields of the supergauge multiplet $V \equiv V^{ij}(x,\theta) = \Gamma_a^{ij} V^a(x,\theta)$. The anti-Hermitian matrices Γ_a are generators of super gauge transformations of matter fields the parameters of which are the chiral superfields $\Lambda_+ = \Gamma_a \Lambda_+^a (x,\theta)$:

$$\Psi_+ \to \Psi_+' = e^{\Lambda_+}\Psi_+ \, ; \quad \Psi_- \to \Psi_-' = e^{-\Lambda_+^+}\Psi_- \tag{1}$$

The gauge supermultiplet is transformed as follows:

$$e^{iV} \to e^{iV'} = e^{-\Lambda_+^+} e^{iV} e^{-\Lambda_+} \tag{2}$$

The supergauge-and supersymmetry-invariant part of the action has the form

$$S_{SS} = \frac{1}{g^2} S_V(V) + S_{V,\Psi} + S_\Psi ,$$

where S_Ψ is expressed as

$$S_\Psi \sim m_+ \Psi_+^2 + f_+ \Psi_+^3 + m_- \Psi_-^2 + f_- \Psi_-^3 + h.c.$$

and S_V and $S_{V,\Psi}$ represent the self-interaction of the super - gauge multiplet and its interaction with matter fields super-multiplets, respectively (explicit form of S_{SS} will not be hence-forth necessary (see, e.g. [2] ; the set of all the matter fields $\Psi_\pm, \Psi_\pm^+$ will be hereinafter denoted by Φ).

Not yet defining concretely the form of the supergauge multiplet new parametrization, let us generally represent transition to it by

$$\left(e^{i\tilde{\Gamma}_a V^a}\right)^{ij} \equiv M^{ij}(x,\theta) = i\tilde{\Gamma}_a^{ij} \pi^a(V(x,\theta)) + T_n^{ij} \sigma^n(\pi(V)) \tag{3}$$

where $\tilde{\Gamma}_a$ does not coincide with Γ_a in general; $\tilde{\Gamma}_a$ and T_n represent the full set of matrices. The gauge transformation (2) now will be written as $(\Lambda_+ = \Lambda_+^a \tilde{\Gamma}_a)$

$$M^{ij}(\pi(x,\theta)) \longrightarrow M'^{ij}(\pi(x,\theta)) = \left(e^{-\Lambda_+^+} M(\pi(x,\theta)) e^{-\Lambda_+}\right)^{ij}.$$

For infinitesimal $\Lambda_+(x,\theta)$ we have

$$\delta M(\pi) = -\Lambda_+^+ M - M\Lambda_+ = \tilde{\Gamma}_a \delta\pi^a + T_n \delta\sigma^n(\pi),$$

whence the law(linear) of transformation of π and $\sigma(\pi)$ one via the other is obtained at group transformations. Further we shall consider a theory where instead of V^a the fields π^a are chosen as the main supergauge fields, and we shall introduce the set $(\pi^a, \sigma^n(\pi)) \equiv \omega^r(\pi)$. In terms of the fields ω^r the infinitesimal supergauge transformations have the form

$$\delta\omega^r = R^r_{+a}(\omega)\Lambda_+^a + R^r_{-a}(\omega)\Lambda_+^{+a} = R^r_A(\omega)\Lambda^A = R^r_{AS}\omega^S\Lambda^A, \qquad (4)$$

where the following notations are introduced

$$R^r_{+a} + R^r_{-a} = R^r_{(\alpha)a}, \quad i(R^r_{+a} - R^r_{-a}) = R^r_{(\beta)a}, \quad (R^r_{(\alpha)a}, R^r_{(\beta)a}) \equiv R^r_A$$

$$\Lambda_+^a + \Lambda_+^{+a} = 2\Lambda^a_{(\alpha)}, \quad \Lambda_+^a - \Lambda_+^{+a} = 2i\Lambda^a_{(\beta)}, \quad (\Lambda^a_{(\alpha)}, \Lambda^a_{(\beta)}) \equiv \Lambda^A, \qquad (5)$$

and R^r_{AS} are already numerical matrices.

Let us now describe the modified actions. As a gauge function choose the function $T^A = t^A_r \bar{D}D\omega^r$, where t^A_r are some (gauge) parameters and the gauge-fixing term S_{gf} in the total action represent as

$$S_{gf} = \int d^4x\, d^4\theta\, [B_{+a}B_{-a} + B_A t^A_r \bar{D}D\omega^r] =$$

$$= \int d^4x\, d^4\theta\, [B_{+a}B_{-a} + (B_{+a}t^{-a}_r + B_{-a}t^{+a}_r)\bar{D}D\omega^r], \qquad (6)$$

$$B_A = \left(\frac{1}{2}(B_{+a} + B_{-a}), \frac{1}{2i}(B_{+a} - B_{-a})\right), \quad t^A_r = \left(t^{+a}_r + t^{-a}_r, \; i(t^{-a}_r - t^{+a}_r)\right).$$

Let us by general rules introduce the following action for the ghost fields

$$S_{gh}=\int d^4x d^4\theta \mathcal{D}_A \bar{D}D t_r^A R_B^r C^B = \int d^4x d^4\theta \tilde{\mathcal{D}}_A\, t_r^A\, R_B^r C^B,$$

where notation $\bar{D}D\mathcal{D}_A=\tilde{\mathcal{D}}_A$ is used, and the fields C^B and $\mathcal{D}_A$ consist of chiral fields and fields Hermitianly-conjugated with them in analogy with the formulae (5).

Now we can write the modified action

$$S = S_{SS}(\pi,\Phi)+S_{gf}+S_{gh}+\int[K_{\Phi,I}R^I_{\Phi,A}(\Phi)C^A+K_r R^r_A(\omega)C^A+ \tag{7}$$

$$+\frac{1}{2}L_D f^D_{AB}C^A C^B+\lambda_r\omega^r+\chi^A_r\omega^r\tilde{\mathcal{D}}_A]\equiv\bar{S}+\int[\lambda_r\omega^r+\chi^A_r\omega^r\tilde{\mathcal{D}}_A],$$

where λ_r is an external source, a general real superfield; χ^A_r are constant Grassmanian numbers; $K_{\Phi,I}$, K_r and L_D are external sources, the chiral properties of which are presented in ref.[3]; f^D_{AB} are structure constants:

$$R^b_B\frac{\delta R^a_A}{\delta\pi^b}-R^b_A\frac{\delta R^a_B}{\delta\pi^b}=f^D_{BA}R^a_D. \tag{8}$$

The modified action S, as a consequence of the BRST-symmetry, satisfies the following equation (this kind of equation for the pure Young-Mills theory was for the first time written by Zinn-Justin [8]).

$$\frac{\delta S}{\delta\Phi^I}\frac{\delta S}{\delta K_{\Phi,I}}+\frac{\delta S}{\delta\pi^a}\frac{\delta S}{\delta K_{\pi,a}}+\frac{\delta S}{\delta C^A}\frac{\delta S}{\delta L_A}+\frac{\delta S}{\delta\tilde{\mathcal{D}}_A}\bar{D}DB_A=$$

$$=\lambda_r\frac{\delta S}{\delta K_r}+\chi^A_r\frac{\delta S}{\delta t^A_r}.$$

The last term in the right hand side of this equation clarifies the meaning of χ^A_r : χ^A_r is the BRST-variation of the parameter t^A_r [10]. Introducing in a standard way the producing functional Γ of vertex functions, we obtain that it satisfies a similar, as S does, equation

$$\frac{\delta\Gamma}{\delta\Phi^{I}}\frac{\delta\Gamma}{\delta K_{\Phi,I}}+\frac{\delta\Gamma}{\delta\pi^{a}}\frac{\delta\Gamma}{\delta K_{\pi,a}}+\frac{\delta\Gamma}{\delta C^{A}}\frac{\delta\Gamma}{\delta L_{A}}+\frac{\delta\Gamma}{\delta\tilde{\mathcal{D}}_{A}}\bar{D}DB_{A}=$$

$$=\lambda_{r}\frac{\delta\Gamma}{\delta K_{r}}+\chi_{r}^{A}\frac{\delta\Gamma}{\delta t_{r}^{A}}.$$

For the divergent part of the one-loop approximation of Γ (denote it by Γ_{1d}) we obtain:

$$\frac{\delta S}{\delta\Phi^{I}}\frac{\delta\Gamma_{1d}}{\delta K_{\Phi,I}}+\frac{\delta\Gamma_{1d}}{\delta\Phi^{I}}\frac{\delta S}{\delta K_{\Phi,I}}+\frac{\delta S}{\delta\pi^{a}}\frac{\delta\Gamma_{1d}}{\delta K_{\pi,a}}+\frac{\delta\Gamma_{1d}}{\delta\pi^{a}}\frac{\delta S}{\delta K_{\pi,a}}+$$

$$+\frac{\delta S}{\delta C^{A}}\frac{\delta\Gamma_{1d}}{\delta L_{A}}+\frac{\delta\Gamma_{1d}}{\delta C^{A}}\frac{\delta S}{\delta L_{A}}+\frac{\delta\Gamma_{1d}}{\delta\tilde{\mathcal{D}}_{A}}\bar{D}DB_{A}=\lambda_{r}\frac{\delta\Gamma_{1d}}{\delta K_{r}}+\chi_{r}^{A}\frac{\delta\Gamma_{1d}}{\delta t_{r}^{A}}. \tag{9}$$

Note that since the supersymmetry transformations are linear over the fields, the supersymmetry of S results in that of Γ. Consequently, Γ_{1d} is a supersymmetric functional. Let us now construct the general expression for Γ_{1d} issuing from the Feynman diagram divergence index. The latter (with account of all additional vertice) is equal to

$$\omega_{G}=2-N-2N_{\lambda} \tag{10}$$

where N is the total number of the external to the diagram superfields: C^{A}, $\tilde{\mathcal{D}}_{A}$, Φ^{I}, $K_{\Phi,I}$, L_{A} and χ_{r}^{A} (χ_{r}^{A} has a mass dimension and its ghost number is -1), N_{λ} denotes the number of vertices $\lambda\omega$. It is seen from (10) that with account of the ghost number conservation Γ_{1d} does not depend on L_{A} at all, the source K_{r} can be included in Γ_{1d} only linearly, in the form of production of KC and $K\chi$; χ , in its turn, can also be included only linearly, and besides $K\chi$ - in the form of the production $\chi\tilde{\mathcal{D}}$. The matter fields Φ can be included in Γ_{1d}

no more than quadratically. There are no divergences having the structure of the production of the fields and sources with the same chirality [11,12]. Thus, the general structure of Γ_{1d} is

$$\Gamma_{1d} = S_1(\Phi,\pi,t) + \int [B_{+a} f^{aB}(\pi,t) B_{-B} + B_{+a} g_-^a(\pi,t) + B_{-a} g_+^a(\pi,t) + \widetilde{\mathcal{D}}_D R_{1B}^D(\pi,t) C^B + K_r Q_{1B}^r(\pi,t) C^B + \lambda_r \omega_1^r(\pi,t) + \chi_r^A \varkappa_{1A}^r(\pi,t) \widetilde{\mathcal{D}}_B + \chi_r^A \xi_{1A}^{rs}(\pi,t) K_s], \tag{11}$$

where S_1 contains matter fields in the combination $\Phi\Phi^+$. The functional S_1 as well as the coefficient functions f^{aB}, $g_\pm^a$, R_{1B}^D, Q_{1B}^r, ω_1^r, $\varkappa_{1A}^{rB}$, ξ_{1A}^{rs} depend in general on π and t. Besides, it follows from the supersymmetry of Γ_{1d} that these coefficient functions are superfields and S_1 is a supersymmetric functional. It particularly means, that all the coefficient functions are unambiguously restored from their values at $\theta = 0$.

Substituting the expression (11) for Γ_{1d} into (9) and setting to zero the expressions before different field structures, we obtain the following set of equations:

$$\frac{\delta R_B^r}{\delta \pi^a} Q_{1A}^a + \frac{\delta Q_{1B}^r}{\delta \pi^a} R_A^a - (A \leftrightarrow B) = f_{AB}^D Q_{1D}^r , \tag{12}$$

$$\frac{\delta \omega^r}{\delta \pi^a} Q_{1A}^a + \frac{\delta \omega_1^r}{\delta \pi^a} R_A^a = Q_{1A}^r , \tag{13}$$

$$-\frac{\delta R_B^s}{\delta \pi^a} \xi_{1A}^{ra} + \frac{\delta \xi_{1A}^{rs}}{\delta \pi^a} R_B^a = \frac{\delta Q_{1B}^s}{\delta t_r^A} , \tag{14}$$

$$\frac{\delta \omega^s}{\delta \pi^a} \xi_{1A}^{ra} = \xi_{1A}^{rs} - \frac{\delta \omega_1^s}{\delta t_r^A} , \tag{15}$$

$$\frac{\delta \xi_{1B}^{st}}{\delta t_r^A} - \frac{\delta \xi_{1A}^{rt}}{\delta t_s^B} = 0, \tag{16}$$

$$\varphi_a(\pi,t)\Big|_{t_0=0} = T^j_a \pi^j f(\pi^2). \tag{42}$$

Let us come back to the general case when $t^a_0 \neq 0$. Let us take the opportunity that the functions $\varphi_0(\pi,t)$ and $\varphi_a(\pi,t)$ are invariants of the Lorentz group, i.e.

$$\varphi_0(\pi,t) = \varphi_0(\pi_\Lambda, t_\Lambda), \qquad \varphi_a(\pi,t) = \varphi_a(\pi_\Lambda, t_\Lambda),$$

where π_Λ and t_Λ are the values of π and t obtained as a result of Lorentz-transformation Λ. Transformation Λ (see (36)) choose from the condition when $t^a_{\Lambda,0} = 0$. Obviously this transformation is always in the vicinity of $t^a_0 = 0$. Then, according to (41), (42) we have

$$\varphi_0(\pi_\Lambda, t_\Lambda) = f_0(\pi^2_\Lambda), \qquad \varphi_a(\pi_\Lambda, t_\Lambda) = T^j_{\Lambda,a} \pi^j_\Lambda f(\pi^2_\Lambda).$$

Now consider the following Lorentz-invariants:

$$\rho^{ab} = t^a_r \bar{t}^{b,r} = \rho^{ab}_\Lambda = t^a_{\Lambda,i} t^b_{\Lambda,i},$$

$$L_{ab} = L_{\Lambda,ab} = T^i_{\Lambda,a} T^i_{\Lambda,b}, \qquad \rho^{ab} L_{bc} = \delta^a_c,$$

$$\omega^r t^a_r = \omega^r_\Lambda t^a_{\Lambda,r} = \pi^i_\Lambda t^a_{\Lambda,i} = \pi^i_\Lambda \bar{t}^{a,i}_\Lambda.$$

Then we have

$$\pi^2_\Lambda \equiv X = \pi^i_\Lambda \bar{t}^{a,i}_\Lambda T^j_{\Lambda,a} T^j_{\Lambda,b} \bar{t}^{b,k}_\Lambda \pi^k_\Lambda =$$

$$= \omega^r_\Lambda t^a_{\Lambda,r} L_{\Lambda,ab} t^b_{\Lambda,s} \omega^s_\Lambda = \omega^r t^a_r L_{ab} t^b_s \omega^s, \tag{43}$$

$$T^j_{\Lambda,a} \pi^j_\Lambda = T^k_{\Lambda,a} T^k_{\Lambda,b} \bar{t}^{b,j}_\Lambda \pi^j_\Lambda = L_{\Lambda,ab} t^b_{\Lambda,r} \omega^r_\Lambda = L_{ab} t^b_r \omega^r,$$

and from the formula (38) we obt...

Let us, at last, solve the set (12)-(18) taking (19) into consideration. From (16) we find the general solution for ξ_{1A}^{rt} :

$$\xi_{1A}^{rt}(\pi,t) = \frac{\delta \tilde{u}^{t}(\pi,t)}{\delta t_r^A} \tag{20}$$

where $\tilde{u}^t(\pi,t)$ is some function. With regard to (20) one obtains from (15)

$$\omega_1^s(\pi,t) = a^s(\pi) + \tilde{u}^s(\pi,t) - \frac{\delta\omega^s}{\delta\pi^a}\tilde{u}^a(\pi,t). \tag{21}$$

Taking into consideration (19), from which it follows that $a^a(\pi)=0$ and introducing the new functions $\tilde{\tilde{u}}^s(\pi,t)$ by

$$\tilde{\tilde{u}}^a(\pi,t) = \tilde{u}^a(\pi,t), \qquad \tilde{\tilde{u}}^n(\pi,t) = \tilde{u}^n(\pi,t) + a^n(\pi),$$

rewrite (21) as

$$\omega_1^s(\pi,t) = \tilde{\tilde{u}}^s(\pi,t) - \frac{\delta\omega^s}{\delta\pi^a}\tilde{\tilde{u}}^a(\pi,t). \tag{22}$$

At the same time, rewrite (20) via $\tilde{\tilde{u}}$:

$$\xi_{1A}^{rt} = \frac{\delta\tilde{\tilde{u}}^t(\pi,t)}{\delta t_r^A}. \tag{23}$$

The equation (12) for $r=a$ was solved in ref.[3].

$$Q_{1A}^a = \text{æ}^b(\pi,t)\frac{\delta R_A^a}{\delta\pi^b} - R_A^b\frac{\delta \text{æ}^a(\pi,t)}{\delta\pi^b}, \tag{24}$$

where $\text{æ}^b(\pi,t)$ are some functions. With account of eq.(13) eq.(1 is identically satisfied for $r=n$.

Now let us find the solution of the equation (18) (see ref. [3]):

$$S_1 = \varkappa^{b}(\pi,t)\frac{\delta S_{cc}}{\delta\pi^{b}} + \frac{2\delta Z_1}{g^2}S_{\pi}(\pi) - 2\delta Z_2 S_{\pi,\Phi} .$$

And the eq. (17) leads to

$$\frac{\delta}{\delta t_r^{A}}\delta Z_i = 0 , \quad i=1,2 ; \qquad \frac{\delta}{\delta t_r^{A}}(\varkappa^{a}(\pi,t)+\tilde{\tilde{u}}^{a}(\pi,t))=0.$$

Now consider the functions $u_0^{a}(\pi) = \varkappa^{a}(\pi,t)+\tilde{\tilde{u}}^{a}(\pi,t)$, then the functions

$$u^{r} = \tilde{\tilde{u}}^{r} - \frac{\delta\omega^{r}}{\delta\pi^{b}}u_0^{b}$$

Then one can write (see eqs. (22), (23))

$$\omega_1^{r} = u^{r} - \frac{\delta\omega^{r}}{\delta\pi^{a}}u^{a} , \qquad \xi_{1A}^{rt} = \frac{\delta u^{t}(\pi,t)}{\delta t_r^{A}} , \qquad \varkappa^{a} = -u^{a}. \quad (25)$$

From eq.(13) one finds

$$Q_{1A}^{r} = -u^{b}\frac{\delta R^{r}}{\delta\pi^{b}} + R_A^{b}\frac{\delta u^{r}}{\delta\pi^{b}} , \quad (26)$$

(compare with (24)); the eq.(14) is satisfied identically. This finishes the solution of the set (12)-(18). The obtained solutions together with the relations (19) allow one all the coefficient functions in (11) to express via the functions $u^{r}(\pi,t)$. Our task is to define the form of all these functions concretely.

3. In this section the Ward's identities as a consequence of global gauge symmetry will be derived and analyzed. Infinite-

simal transformations of this symmetry can be obtained from (4) if independent on χ and θ Λ_+^a parameters are chosen: $\Lambda_+^a = (\alpha + i\beta)^a$ where α and β are real numbers. So,

$$\delta_g \omega^r = R_\alpha^r(\omega)\xi^A = R_{AS}^r \omega^S \xi^A, \quad \xi_{(\alpha)}^a \equiv \alpha^a, \quad \xi_{(\beta)}^a \equiv \beta^a, \tag{27}$$

where δ_g denotes global variation. The fields C^A are transformed as

$$\delta_g C^D = f_{AB}^D C^A \xi^B,$$

and the matter fields ϕ are transformed according to (1) with constant Λ_+^a. Making global transformations of the fields in (7), we come to an identity for the modified action S. Then, by usual procedure we obtain the same identity for the Γ-functional:

$$\frac{\delta\Gamma}{\delta\pi^a} R_{AS}^a \frac{\delta\Gamma}{\delta\lambda_S} + \frac{\delta\Gamma}{\delta\phi^I} R_{\phi,AJ}^I \phi^J + \frac{\delta\Gamma}{\delta C^D} f_{BA}^D C^B = t_r^D R_{AS}^r \frac{\delta\Gamma}{\delta t_S^D} +$$

$$+ K_{\phi,I} R_{\phi,AJ}^I \frac{\delta\Gamma}{\delta K_{\phi,J}} + K_r R_{AS}^r \frac{\delta\Gamma}{\delta K_S} + L_C f_{BA}^C \frac{\delta\Gamma}{\delta L_B} + \tag{28}$$

$$+ \lambda_r R_{AS}^r \frac{\delta\Gamma}{\delta\lambda_S} + \chi_r^B R_{AS}^r \frac{\delta\Gamma}{\delta\chi_S^B}.$$

In future it is sound practice to exclude from consideration the value $\delta\Gamma/\delta t_S^D$ present in (28). It can be done by means of eq.(9), differentiating it with respect to χ beforehand. As a result we obtain the following identity for Γ_{1d}:

$$\frac{\delta\Gamma_{1d}}{\delta\pi^a} R_{AS}^a \frac{\delta S}{\delta\lambda_S} + \frac{\delta S}{\delta\pi^a}\left(R_{AS}^a \frac{\delta\Gamma_{1d}}{\delta\lambda_S} + t_r^F R_{At}^r \xi_{1F}^{ta}\right) + \frac{\delta\Gamma_{1d}}{\delta\phi^I} R_{\phi,AJ}^I \phi^J +$$

$$\frac{\delta\Gamma_{1d}}{\delta C^D} f^D_{BA} C^B = K_{\Phi,I} R^I_{\Phi,AJ} \frac{\delta\Gamma_{1d}}{\delta K_{\Phi,J}} + K_r R^r_{AS} \frac{\delta\Gamma_{1d}}{\delta K_S} + L_c f^c_{BA} \frac{\delta\Gamma_{1d}}{\delta L_B} +$$

$$+ \lambda_r R^r_{AS} \frac{\delta\Gamma_{1d}}{\delta\lambda_S} + \chi^B_r R^r_{AS} \frac{\delta\Gamma_{1d}}{\delta\chi^B_S} + t^F_r R^r_{At} \Big[\frac{\delta\omega^t}{\delta\pi^a} \tilde{\mathcal{D}}_F Q^a_{1B} C^B -$$

$$- \frac{\delta\omega^t}{\delta\pi^a} \tilde{\mathcal{D}}_F \chi^B_r \xi^{ra}_{1B} + \frac{\delta æ^{tB}_{1F}}{\delta\pi^a} \tilde{\mathcal{D}}_B R^a_D C^D + \frac{\delta\xi^{ts}_{1F}}{\delta\pi^a} K_S R^a_B C^B +$$

$$+ æ^{tB}_{1F} \bar{D}D\theta_B + \lambda_r \xi^{tr}_{1F} + \chi^D_r \frac{\delta æ^{tB}_{1F}}{\delta t^D_r} \tilde{\mathcal{D}}_B + \chi^B_r \frac{\delta\xi^{ts}_{1F}}{\delta t^B_r} K_S \Big].$$

Substituting expression (11) for Γ_{1d} into this, one obtains the following equations:

$$\frac{\delta\omega^r}{\delta\pi^a}\left(R^a_{AS}\omega^S_1 + t^F_S R^S_{At} \xi^{ta}_{1F}\right) + \frac{\delta\omega^r_1}{\delta\pi^a} R^a_{AS}\omega^S = t^F_S R^S_{At}\xi^{tr}_{1F} + R^r_{AS}\omega^S_1 \quad (29)$$

$$\frac{\delta S_{cc}}{\delta\pi^a}\left(R^a_{AS}\omega^S_1 + t^F_S R^S_{At}\xi^{ta}_{1F}\right) + \frac{\delta S_1}{\delta\pi^a} R^a_{AS}\omega^S + \frac{\delta S_1}{\delta\Phi^I} R^I_{\Phi,AJ}\Phi^J = 0 \quad (30)$$

(the remaining equations are satisfied identically that is why they are not written).

Comparing (30) with (18) one obtains

$$Q^a_{1A} = R^a_{AS}\omega^S_1 + t^F_S R^S_{At}\xi^{ta}_{1F}$$

and, with account of this relation, the comparison of (29) with (13) yields

$$G^r_{1A} = R^r_{AS}\omega^S_1 + t^F_S R^S_{At}\xi^{tr}_{1F} \,, \tag{31}$$

which is the generalization of the former expression.

Substituting instead of Q^r_{1A}, ω^s_1 and ξ^{tr}_{1F} their expressions via $u(\pi,t)$ into (31) (see (25), (26)) one obtains

$$\hat{Q}_A u^r = R^r_{As} u^s, \quad \hat{Q}_A \equiv R^i_A \frac{\delta}{\delta \pi^i} - t^F_s R^s_{Au} \frac{\delta}{\delta t^F_u} \qquad (32)$$

The relation (32) testifies, that under the action of the operators $\hat{Q}^A$ the values u^r are transformed as vectors. As to the operators themselves, it is not difficult to check that they form an algebra which is isomorphous to the algebra of the operators $\hat{R}_A = R^i_A \delta/\delta\pi^i$:

$$[\hat{Q}_A, \hat{Q}_B] = f^D_{BA} \hat{Q}_D$$

(cf. (8)). Also note, that under the action of the operator $\hat{Q}^A$, ω^s and t^B_s are also transformed as vectors:

$$\hat{Q}_A \omega^s = R^s_{At} \omega^t, \qquad \hat{Q}_A t^B_s = -R^r_{As} t^B_r. \qquad (33)$$

4. All our further considerations will apply to the theory with the SU(2) group and a concrete parametrization which will be now discussed. As $i\tilde{\Gamma}_a$ (see (3)) consider the doublet-representation generators τ_i, where τ_i are the Pauli's matrices:

$$e^{\tau_i v^i} = \operatorname{ch} v + \tau_i \frac{v^i}{v} \operatorname{sh} v \equiv M(x,\theta) = \tau_i \pi^i + \sigma(\pi),$$

whence

$$\pi^i(v) = \frac{v^i}{v} \operatorname{sh} v, \quad \sigma(\pi) = \operatorname{ch} v, \quad \sigma^2 - \pi^2 = 1, \quad v = \sqrt{(v^i)^2} \qquad (34)$$

Thus, the field ω^r is a four-component one: $\omega^r = (\pi^i, \omega^0) \equiv (\pi^i, \sigma)$ (in this section in the indices like r, s, t, ... , instead of a, b, c, ... we shall put i, j, k, ...: $s=(i,o)$). From (34) it is seen, that the value of $\sigma^2-\pi^2$ is an invariant, and the global gauge transformations

$$M(x,\theta) \rightarrow M'(x,\theta) = e^{i\tau_\kappa(\alpha-i\beta)^\kappa} M(x,\theta)\, e^{-i\tau_\ell(\alpha+i\beta)^\ell} \tag{35}$$

are transformations from the group SL(2,C) which is isomorphous to the Lorentz group. The four-vector q^r transformation law at Lorentz transformations is ($\alpha=0$)

$$q^0_\Lambda = q^0 \operatorname{ch} 2\beta + q_i \frac{\beta^i}{\beta} \operatorname{sh} 2\beta\,,$$

$$q^i_\Lambda = q^i + q^0 \frac{\beta^i}{\beta} \operatorname{sh} 2\beta + 2q_\kappa \frac{\beta^i\beta^\kappa}{\beta^2} \operatorname{sh}^2\beta\,, \quad \beta = \sqrt{(\beta^i)^2} \tag{36}$$

The transformation laws of the fields π^i and σ at the transformations (35) for infinitesimal α and β have the form of (27), where

$$R^i_{(\alpha)j\kappa} = 2\varepsilon_{i\kappa j}\,, \quad R^i_{(\alpha)j0} = 0\,, \quad R^0_{(\alpha)j\kappa} = 0\,, \quad R^0_{(\alpha)j0} = 0,$$

$$R^i_{(\beta)j\kappa} = 0\,, \quad R^i_{(\beta)j0} = 2\delta_{ij}\,, \quad R^0_{(\beta)j\kappa} = 2\delta_{j\kappa}\,, \quad R^0_{(\beta)j0} = 0. \tag{37}$$

In accordance with the noticing made in the preceding section (see (33)), the vectors ω^s and $t^A_s = (t^{+a}_s + t^{-a}_s,\ i(t^{-a}_s - t^{+a}_s))$ are in the theory. We shall consider the case of $t^{+a}_s = t^{-a}_s = t^a_s$. Introduce the vectors $\tilde{t}^{a,s}$ via

$$\tilde{t}^{a,s} = \eta^{sr} t^a_r\,, \qquad \eta^{sr} = \begin{pmatrix} +1 & & & 0 \\ & +1 & & \\ & & +1 & \\ 0 & & & -1 \end{pmatrix}$$

and present $U^s(\pi,t)$ as an expansion

$$U^s(\pi,t) = \omega^s(\pi)\,\varphi_0(\pi,t) + \tilde{t}^{a,s}\,\varphi_a(\pi,t)\,. \tag{38}$$

Substituting this expansion into eq.(32), one obtains

$$\hat{Q}_A \varphi_o(\pi,t) = \hat{Q}_A \varphi_a(\pi,t) = 0\,, \tag{39}$$

i.e. the coefficient functions φ_o and φ_a are the Lorentz group invariants.

Next will be considered the so-called singular gauge, when $S_{gf} = \int B_A T^A = \int (B_{+a} t_r^{-a} + B_a t_r^{+a}) \bar{D} D \omega^r$ (see(6)). Formally one can pass to this gauge by the substitutions $t_r^A \to \alpha t_r^A$, $B_A \to (1/\alpha) B_A$ with following tending of α to ∞. This gauge is noteworthy due to the following observation. In the modified action S there are two summands which are proportional to the parameters t_r^A: S_{gf} and S_{gh}. At $t_i^{+a} = t_i^{-a} = t_i^a$ one can change the variables in S

$$B'_{\pm i} = B_{\pm a} t_i^a\,,\quad \tilde{\mathcal{D}}'_{\pm i} = \tilde{\mathcal{D}}_{\pm a} t_i^a\,,\quad \chi_r^{\prime \pm i} = \chi_r^{\pm a} T_a^i\,,$$

$$\tilde{\mathcal{D}}_{-a} = \tilde{\mathcal{D}}_+^{+a}\,,\quad t_i^a T_b^i = \bar{t}^{a,i} T_b^i = \delta_b^a\,,\quad \bar{t}^{a,i} T_a^j = \delta^{ij}\,, \tag{40}$$

passing in the functional integral to integration over the primed fields, after which, at $t_o^a = 0$, the dependence on t_r^A vanishes in the theory. This observation is equivalent to the statement that in the singular gauge, at $t_o^a = 0$, the theory is independent of the parameters t_r^A. Let us first consider the case when $t_o^a = 0$. In this case U^S does not depend on t. Considering the expansion (38) for $S = 0$, we obtain, that at $t_o^a = 0$, φ_o is independent of t. And from (39) it follows that φ_o is a function of π^2 only:

$$\varphi_o(\pi,t)\Big|_{t_o=0} = f_o(\pi^2)\,. \tag{41}$$

Likewise, for $S = i$ from (38) we obtain:

$$-\frac{\delta S_{cc}}{\delta \pi^{a}} \xi^{ra}_{1A} = \frac{\delta S_1}{\delta t^{A}_{r}} , \tag{17}$$

$$\frac{\delta S_1}{\delta \Phi^{I}} R^{I}_{\Phi,A} + \frac{\delta S_{cc}}{\delta \pi^{a}} Q^{a}_{1A} + \frac{\delta S_1}{\delta \pi^{a}} R^{a}_{A} = 0 . \tag{18}$$

The equations being identically satisfied with account of the eqs (12-18) and a series of additional relations are not written here. These relations will be derived below. One can come to these relations issuing from the equalities for the modified action S which follow from its expression (7):

$$\frac{\delta S}{\delta \bar{\mathcal{D}}_{A}} = -E_K t^{A}_{r} \frac{\delta S}{\delta K_r} + E_K \chi^{A}_{r} \frac{\delta S}{\delta \lambda_r}$$

$$\frac{\delta S}{\delta B_{\pm a}} = B_{\mp a} + E_{\mp} \bar{D} D t^{\mp a}_{r} \frac{\delta S}{\delta \lambda_r} ,$$

$$\frac{\delta S}{\delta \lambda^{a}} = \pi^{a}$$

where $E_K = E_+ + E_-$, E_+ , E_- are projectors on the chiral superfields. The functional Γ also satisfies these equations. Directly write equations for Γ_{1d} :

$$\frac{\delta \Gamma_{1d}}{\delta \bar{\mathcal{D}}_{A}} = -E_K t^{A}_{r} \frac{\delta \Gamma_{1d}}{\delta K_r} + E_K \chi^{A}_{r} \frac{\delta \Gamma_{1d}}{\delta \lambda_r} ,$$

$$\frac{\delta \Gamma_{1d}}{\delta B_{\pm a}} = E_{\mp} \bar{D} D t^{\mp a}_{r} \frac{\delta \Gamma_{1d}}{\delta \lambda_r} ,$$

$$\frac{\delta \Gamma_{1d}}{\delta \lambda^{a}} = 0 .$$

Substituting the expression for Γ_{1d} (11) into this, one obtains

$$\omega^{a}_{1}(\pi,t) = 0 , \quad f^{aB}(\pi,t) = 0 , \quad R^{D}_{1B} = t^{D}_{r} Q^{r}_{1B} , \tag{19}$$

$$g^{a}_{\pm} = E_{\pm} \bar{D} D t^{\pm a}_{r} \omega^{r}_{1} , \qquad \text{æ}^{rA}_{1B} = t^{A}_{s} \xi^{rs}_{1B} + \delta^{A}_{B} \omega^{r}_{1} .$$

$$u^r(\pi,t) = \omega^r f_o(x) + \bar{t}^{a,r} L_{a\beta} t^\beta_s \omega^s f(x). \tag{44}$$

Next, let us present $\ell^{a\beta}$, $L_{a\beta}$ in a more convenient form:

$$\ell^{a\beta} = t^a_i \hat{\ell}^{ij} t^\beta_j, \qquad \hat{\ell}^{ij} = \delta^{ij} - \gamma^i\gamma^j, \qquad \gamma^i = T^i_a t^a_o,$$
$$L_{a\beta} = T^i_a \hat{L}^{ij} T^j_\beta, \qquad \hat{L}^{ij}\hat{\ell}^{jk} = \delta^{ik}. \tag{45}$$

Rewrite the relations (43), (44) in terms of γ^i:

$$x = (\pi^k + \sigma\gamma^k)\hat{L}^{ke}(\pi^e + \sigma\gamma^e),$$
$$u^i(\pi,t) = \pi^i \tilde{f}_o(x) + [\hat{L}^{ij}(\pi^j + \sigma\gamma^j) - \pi^i] f(x),$$
$$u^o(\pi,t) = \sigma\tilde{f}_o(x) - [\gamma^i \hat{L}^{ij}(\pi^j + \sigma\gamma^j) + \sigma] f(x), \quad \tilde{f} = f_o + f.$$

The advantage of writing in terms of the values introduced in (45) is in the fact that there the whole dependence of the theory on the parameters t^a_r is expressed in terms of γ^i. Indeed, if in the functionals S_{gf} and S_{gh} the substitution (40) is accomplished, then S_{gf} and S_{gh} will have the form

$$S_{gf} = (B'_{+i} + B'_{-i})\bar{D}D(\pi^i + \gamma^i\sigma), \quad S_{gh} = (\tilde{\mathcal{D}}'_{+i} + \tilde{\mathcal{D}}'_{-i})[R^i_B C^B + \gamma^i R^o_B C^B].$$

Hence it is seen that γ^i is included in the interaction vertices only. It means, that any quantum field value, u^r among them, in any order of the perturbation theory must be a finite γ^i polynomial. Consider the consequences of this fact for the functions $\tilde{f}_o(x)$ and $f(x)$.

Choose γ^i without restriction to generality in the form $\gamma^i = \gamma\delta^{i1}$. Then we have:

$$\hat{\ell}^{ij} = \delta^{ij} - \gamma^2\delta^{i1}\delta^{j1} = \begin{pmatrix} 1-\gamma^2 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 1 \end{pmatrix}, \qquad \hat{L}^{ij} = \begin{pmatrix} \frac{1}{1-\gamma^2} & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 1 \end{pmatrix}$$

$$X = \frac{(\pi^1+\gamma\sigma)^2}{1-\gamma^2} + (\pi^2)^2 + (\pi^3)^2, \qquad u^0(\pi,t) = \sigma\tilde{f}_0(X) - \frac{\pi^1\gamma+\sigma}{1-\gamma^2} f(X),$$

$$u^1(\pi,t) = \pi^1\tilde{f}_0(X) + \frac{\gamma(\pi^1\gamma+\sigma)}{1-\gamma^2} f(X), \quad u^2(\pi,t) = \pi^2\tilde{f}_0(X), \quad u^3(\pi,t) = \pi^3\tilde{f}_0(X)$$

For simplicity take $\pi^2 = \pi^3 = 0$:

$$X = \frac{(\pi^1+\gamma\sigma)^2}{1-\gamma^2}; \qquad \sigma = \sqrt{1+(\pi^1)^2}. \tag{46}$$

According to the aforesaid $\tilde{f}_0(X)$ and $f_1(\pi,\gamma) \equiv f(X)(\pi^1\gamma+\sigma)/1-\gamma^2$ must be polynomials over γ. Consider X at $\pi^1 = 0$: $X|_{\pi^1=0} = \gamma^2/1-\gamma^2$. Then $(\gamma^2/1-\gamma^2)\tilde{f}_0 = P_0(\gamma^2)$, ($P_0$ – polynomial) or $\tilde{f}_0(X) = P_0(X/1+X)$. Substituting the value of X from (46) into the right hand side of the last relation, we obtain

$$\tilde{f}_0(X) = P_0\left[\frac{(z+\gamma)^2}{(1+\gamma z)^2}\right], \qquad z = \frac{\pi^1}{\sigma}, \tag{47}$$

The right hand side must have no singularities. The expression (47) is symmetric over γ and z, hence it is a final-order polinomial both over γ and z and, consequently, has no peculiarities over γ and z. Hence it follows that the polinomial P_0 in (47) is a constant, i.e. $\tilde{f}_0(X) = C$. And for $f_1(\pi_0\gamma)$ we have:

$$f_1\Big|_{\pi^1=0} = \frac{1}{1-\gamma^2} f\left(\frac{\gamma^2}{1-\gamma^2}\right) = P_1(\gamma^2)$$

whence

$$f(X) = \frac{1}{1+X} P_1\left(\frac{X}{1+X}\right)$$

Substituting the value of X from (46) we obtain

$$f_1(\pi,\gamma) = \frac{1}{\sigma}\,\frac{1}{1+z\gamma}\, P_1\left[\frac{(z+\gamma)^2}{(1+z\gamma)^2}\right],$$

where, as apposed to $\tilde{f}_0(X)$, there is always a pole at $z = 1/\gamma$. That is why $P_1(\gamma) \equiv 0$.

Thus, we finally obtain the expansion of $u^r(\pi,t)$:

$u^r(\pi,t)=\delta z_\pi \omega^r(\pi)$,where $\delta z_\pi \equiv C$ does not depend on π and t.

Now it is seen, that the one-loop renormalized action $S-\eta\Gamma_{1d}$ is obtained by expanding S_{1R} with an accuracy to the terms having the first order over η, which is constructed by the initial modified action S by the substitution:

$$\pi \to z_\pi \pi,\quad \phi \to z_\phi^{1/2}\phi,\quad c \to c,\quad \tilde{\mathcal{D}} \to \tilde{\mathcal{D}},$$

$$K \to z_\pi^{-1}K,\quad K_\phi \to z_\phi^{-1}K_\phi,\quad L \to L,\quad \lambda \to z_\pi^{-1}\lambda,\quad \chi \to z_\pi^{-1}\chi,$$

$$t \to z_\pi^{-1}t,\quad g \to z_g g,\quad f \to z_\phi^{-3}f,\quad m \to z_\phi^{-2}m,$$

$$z_\pi = 1+\eta\delta z_\pi,\quad z_g = 1+\eta\delta z_1,\quad z_\phi = 1+\eta\delta z_2.$$

Then, considering the theory based on S_{1R} we find that S_{1R} satisfies all the equations satisfied by the action S. Consequently, the producing Γ-functional of the theory developed over S_{1R} will satisfy all the identities satisfied by the Γ-functional of the theory with the action S, etc. Therefore, in the highest approximation one can use the method of mathematical induction.

Thus, finally we obtain, that in the singular gauge the theory is renormalized multiplicatively.

References

1. Wess J., Zumino B., Supergauge invariant extension of quantum electrodynamics. Nucl.Phys., 1974, vol.B78, N.1, p.1.
2. Salam A., Strathdee J., Supersymmetry and non-Abelian gauge. Phys.Lett., 1974, vol.B51, p.353-355; Delbourgo R., Salam A., Strathdee J., Supersymmetric V-A gauges and fermion number. Phys.Lett., 1974, vol.B51, p.475-478.
3. Тютин И.В. Перенормировка суперкалибровочных теорий с нерасширенной суперсимметрией. ЯФ, I983, т.37, № 3, с.76I-77I.
4. Piguet O., Sibold K., Renormalization of N=1 Supersymmetric Yang-Mills Theories (I). The Classical Theory. Nucl.Phys., 1982, vol.B197, p.257-271; Piguet O., Sibold K., Gauge Independence in N=1 Supersymmetric Yang-Mills Theories. Nucl.Phys. 1984, vol.B248, p.301-335.
5. Воронов Б.Л., Тютин И.В. Перенормировка двухмерных киральных теорий. ЯФ, I98I, т.33, № 4, с.II37-II47.
6. Brezin E., Zinn-Justin J., Le Gouillou J.C. Renormalization of the nonlinear σ-model in $2+\varepsilon$ dimensions. Phys.Rev. 1976, vol.D14, 2615-2621.
7. Bardeen W.A., Shizuya K., Structure and renormalizability of massive Yang-Mills field theories. Phys.Rev., 1978, vol.D18 N.6, p.1969-1982.
8. Zinn-Justin J., Renormalization of gauge theories. In: Lecture Notes in Physics. Springer, 1975, vol.37, p.2-39.
9. Братчиков А.В., Тютин И.В. Мультипликативная перенормировка двумерных киральных теорий. [illegible], т.66, № 3, с.[illegible];

Братчиков А.В., Тютин И.В. Мультипликативная перенормировка двухмерных киральных теорий II. ТМФ, 1987, т.70, № 3, с.405-411.

10. Kluberg-Stern H., Zuber J.B., Ward identities and some clues to the renormalization of gauge invariant operators. Phys.Rev., 1975, vol.D12, p.467-481; Kluberg-Stern H., Zuber J.B., Renormalization of non-Abelian gauge theories in a background-field gauge. I,II. Phys.Rev., 1975, vol.12, p.482-488; 3159-3180.

11. Ferrara S., Piguet O., Perturbation theory and renormalization of supersymmetric Yang-Mills theories. Nucl.Phys., 1975, vol.B93, p.261-302.

12. Кривощёков В.К., Славнов А., Файзуллаев Б.А. Супердиаграммная техника для калибровочных теорий. ТМФ, 1976, т.26, с.147-161.

The manuscript was received 2 December 1988

Р.П.ГРИГОРЯН, И.В.ТЮТИН

МУЛЬТИПЛИКАТИВНАЯ ПЕРЕНОРМИРОВКА N = I СУПЕРСИММЕТРИЧНЫХ ЯНГ-МИЛЛСОВСКИХ ТЕОРИЙ: ГРУППА SU(2)

(на английском языке, перевод Папяна Г.А.)

Редактор Л.П.Мукаян

Технический редактор А.С.Абрамян

Подписано в печать 24/УII-89г.	ВФ-02057 Формат 60x84/16
Офсетная печать.Уч.изд.л.I,0	Тираж 299 экз. Ц. I5 к.
Зак.тип.№ I399	Индекс 3649

Отпечатано в Ереванском физическом институте

Ереван 36, ул.Братьев Алиханян,2.

The address for requests:
Information Department
Yerevan Physics Institute
Alikhanian Brothers 2,
Yrevan, 375036
Armenia, USSR

ИНДЕКС 3649

ЕРЕВАНСКИЙ ФИЗИЧЕСКИЙ ИНСТИТУТ