SUSY SU(6) GUT
without Gauge Hierarchy Problem

Zurab Tavartkiladze

INFN Sezione di Ferrara, I-44100 Ferrara, Italy
and Institute of Physics of Georgian Academy of Sciences, 380077 Tbilisi, Georgia

Abstract

A solution of the doublet-triplet splitting problem in the supersymmetric $SU(6)$
gauge theory is suggested. The ‘missing doublet’ multiplet – 175-plet of the $SU(6)$
group as well as the custodial $SU(2)_{\text{cus}}$ global symmetry play crucial role for achieving
the doublet-triplet hierarchy. Two examples in which the doublet-triplet splitting occurs naturally are presented.

\textsuperscript{1}E-mail address: tavzur@axpfe1.fe.infn.it
The supersymmetric (SUSY) Grand Unified Theories (GUT) provide an attractive possibility to understand the stability of the electroweak symmetry breaking scale and the unification of the gauge couplings. It is well known [1] that in the minimal supersymmetric standard model (MSSM) the constants $g_{3,2,1}$ of the gauge group $G_{321}$ join at energies $M_X \sim 10^{16}$ GeV, at which scale $G_{321}$ can be consistently embedded in $SU(5)$ or some larger group $G$. This suggests the following paradigm: below the Planck scale $M_{Pl}$ the hypothetical “theory of everything” reduces to a SUSY GUT with gauge group $G$, which first breaks down to $SU(5)$ at scale $M_G \geq M_X$, and then at the scale $M_X \ SU(5)$ reduces to MSSM:

\[ G \xrightarrow{M_G} SU(5) \xrightarrow{M_X} G_{321} \]

(1)

Obviously, it is also possible that the $G$ breaks to $G_{321}$ at once, directly at the scale $M_G \sim 10^{16}$ GeV.

The main problem which emerges in SUSY GUTs is a problem of the doublet-triplet (DT) splitting. The MSSM Higgs doublets $(h_1, h_2)$ which induce the electroweak symmetry breaking and fermion masses should be light (with mass $\sim M_W$), while their colour-triplet partners in GUT supermultiplets should have masses of order of $M_X$ in order to avoid too fast decay of nucleon. Several mechanisms are known for solving the DT splitting problem without fine tuning:

(i) The missing partner mechanism [2], which is operative directly in $SU(5)$ theory. Besides the standard $\tilde{\mathbf{5}} + \mathbf{5}$ Higgses it requires the ‘missing doublet’ multiplets $\mathbf{50} + \mathbf{50}$ (which however contain the colour-triplets) and the Higgs 75-plet for the $SU(5)$ breaking.

(ii) The missing VEV mechanism [3] can be realized in $SO(10)$ model. Among other relevant Higgs multiplets it employs two 10-plets and a 45-plet with the specific direction of VEV towards the $B - L$ generator of $SO(10)$.

(iii) The Goldstone boson mechanism [4, 5]. In these scenarios the light Higgs doublets $h_{1,2}$ emerge as pseudo-goldstone modes, as a result of the spontaneous breaking of the larger global symmetry of the superpotential. In particular, the models [5] based on $SU(6)$ gauge symmetry includes Higgses 35 and $\tilde{\mathbf{6}} + \mathbf{6}$, and the Higgs superpotential possess an accidental global symmetry $SU(6) \times SU(6)$ independently transforming these two sets.

(iv) The custodial symmetry mechanism [6] is also based on the $SU(6)$ gauge group. The Higgs sector includes the 35-plet and two pairs of $\tilde{\mathbf{6}} + \mathbf{6}$ related by the ’custodial’ global symmetry $SU(2)_{cus}$. If the mass term of 35 is suppressed in the superpotential (or it is the SUSY breaking scale $\sim M_W$), then after $SU(6)$ breaking to $G_{321}$ the pair of doublet-antidoublet from $\tilde{\mathbf{6}} + \mathbf{6}$ which can serve as MSSM Higgses $h_{1,2}$ remain light.

In the case (i) the $SU(5)$ unification of the gauge constants is straightforward, with the possible uncertainties related to the GUT threshold corrections. In the cases (ii) and (iii) with $G = SO(10)$ and $SU(6)$ respectively the hierarchy $M_G \geq M_X$ in breaking [11] is consistent and even can have interesting understanding of fermion mass hierarchies [7, 8].
However, in the custodial symmetry mechanism the picture can not be achieved and the SU(6) gauge group breaking proceeds as

$$ SU(6) \xrightarrow{M_G} SU(3)_C \times SU(3)_W \times U(1)_I \xrightarrow{M_I} G_{321} $$  \hspace{1cm} (2)

where due to specifics of the model the intermediate scale $M_I$ emerges as a geometrical $M_I \sim \sqrt{M_G M_W}$. Consequently, unification of the gauge couplings is spoiled.

A possibility of improving this drawback was suggested in [9], where the 35-plet of the model was replaced by the 175-plet of $SU(6)$ — the traceless tensor $\Phi^{ABC}_{A'B'C'}$ antisymmetric in the up and down indices. It is crucial that 175, in contrast to 35, does not contain the $G_{321}$ doublet fragments. This feature allows to have $M_I \sim M_G$ and in principle the gauge couplings could be directly unified at the scale $M_G \sim 10^{16}$ GeV. However, there emerges the following problem: $\Phi$ has no renormalizable coupling to the Higgses $\tilde{H}, H$ in representations $6 + \bar{6}$, so that the renormalizable Higgs superpotential possesses an extra global symmetry, related to the independent $SU(6)$ transformations of $\Phi$ and $\tilde{H}, H$. In order to avoid the extra Goldstone degrees of freedom which in fact are the colour triplets, the nonrenormalizable couplings like $\frac{1}{M_P} \Phi^3 \tilde{H} H$ cutoff by the (reduced or genuine) Planck scale, $M_P \sim 10^{18} - 10^{19}$ GeV, should be introduced in the theory. However, since $M_G \ll M_P$, these colour triplets will get the masses no more than $\sim \frac{M_G^3}{M_P^2} = 10^{10} - 10^{12}$ GeV. This would affect the unification of the gauge couplings and, which is more dramatic, would lead to the unacceptably fast proton decay.

In this paper we present the possibility which do not suffer from these problems. It naturally provides the DT splitting while the extra global symmetries can be avoided already at level of the renormalizable superpotential.

Consider the supersymmetric gauge $SU(6)$ model with the global custodial symmetry $SU(2)_{\text{cus}}$. The fermion sector consists of the anomaly free chiral set of supermultiplets with the following content under the $SU(6) \times SU(2)_{\text{cus}}$ group per one generation: $\bar{6}^m \sim (6, 2)$ and $15 \sim (15, 1)$, where $m$ is the $SU(2)_{\text{cus}}$ index.

The Higgs sector contains the superfields $\Phi \sim (175, 1)$ needed for breaking of $SU(6)$ to $G_{331} = SU(3)_C \times SU(3)_W \times U(1)_I$, and $\Psi_m \sim (R, 2)$ and $\bar{\Psi}^m \sim (\bar{R}, 2)$, $m = 1, 2$, for the further breaking of $G_{331}$ down to $G_{321}$. Here $R$ can be 84 or 210 of $SU(6)$, which representations are uniquely selected by the following requirements:

a) It should acquire a VEV inducing the $G_{331}$ symmetry breaking down to $G_{321}$. Therefore, it should contain the $G_{321}$ doublet fragment which will be absorbed by corresponding vector superfields of $G_{331}$ through the Higgs mechanism.

b) The tensor product $\bar{R} \times R$ should contain 175 in order to avoid the accidental global symmetries in the Higgs superpotential due to the renormalizable coupling $\Phi \Psi \bar{\Psi}$. 


Therefore, the most general renormalizable Higgs superpotential has the form:

\[
W(\Phi, \Psi) = M_\Phi \Phi^2 + \lambda \Phi^3 + M_\Psi \Phi \bar{\Psi}_m + h \Phi \bar{\Psi}_m \Psi_m
\]  

(3)

We assume that the mass parameters \( M_\Phi, M_\Psi \) are of order of the GUT scale. Hence, the model contains only one mass scale \( \sim M_G \) and no intermediate scale will arise in the model, unlike the model of ref. [3].

The VEV of \( \Phi \) can induce the \( SU(6) \) breaking only to the \( G_{331} \) subgroup, among all maximal subgroups. In other words, \( SU(6) \) breaking to \( SU(5) \times U(1) \) and \( SU(4) \times SU(2) \times U(1) \) is not possible. Indeed, decomposition of 175 in terms of the \( G_{331} \) multiplets is:

\[
175 = (1, 1)_0 + (1, 1)_6 + (8, 8)_0 + (6, \bar{6})_2 + (\bar{6}, 6)_2 +
+(3, 3)_2 + (\bar{3}, \bar{3})_{-2} + (3, 3)_{-4} + (3, 3)_4
\]  

(4)

where subscripts denote the \( U(1)_I \) charges defined by the \( Y_I \sim \text{Diag}(1, 1, 1, -1, -1, -1) \) generator of \( SU(6) \). It is crucial that 175 does not contain \( SU(2)_W \) doublet, and so non of its fragments can participate in the breaking \( G_{331} \rightarrow G_{321} \). The component \((1, 1)_0\) from 175 has the following VEV structure:

\[
\langle \Phi^{abc}_{a'b'c'} \rangle = 3V \epsilon^{abc} \epsilon_{a'b'c'} , \quad \langle \Phi^{ijk}_{i'j'k'} \rangle = -3V \epsilon^{ijk} \epsilon_{i'j'k'}
\]

\[
\langle \Phi^{ia}_{i'a'} \rangle = -V \delta^i_a \epsilon^{abc} \epsilon_{ca'b'} , \quad \langle \Phi^{ja}_{j'a'} \rangle = V \delta^a_j \epsilon^{ijk} \epsilon_{k'i'j'}
\]  

(5)

where \( \epsilon \) is the \( SU(3)_C \) invariant antisymmetric tensor, \( a, b, ... \) and \( i, j, ... \) denote \( SU(3)_C \) and \( SU(3)_W \) indices respectively.

\( SU(3)_W \times U(1)_I \) is farther broken to \( SU(2)_W \times U(1)_Y \) by the VEVs of \( \bar{\Psi}_m + \Psi_m \) in \( G_{321} \) singlet components. Due to the \( SU(2)_{\text{cus}} \) symmetry their VEVs can be placed only on the first pair. Then the doublet-antidoublet which come from \( \bar{\Psi}_1 + \Psi_1 \) will be goldstones. As far the doublets from the \( \bar{\Psi}_2 + \Psi_2 \), they remain massless because they are related to the genuine Goldstone doublets by custodial \( SU(2)_{\text{cus}} \) symmetry. Below we present two examples of such \( SU(6) \times SU(2)_{\text{cus}} \) models, with \( R \) chosen as 84 or 210.

**R = 84:** In this case \( \bar{\Psi}, \Psi \) are three index tensors \( \Psi^A_{BC} \), where \( \Psi^A_{BC} = -\Psi^A_{CB} \) and \( \Sigma \Psi^A_{AB} = 0 \). The content of 84 with respect for \( SU(5) \) subgroup is:

\[
84 = 24_{-5} + 45_1 + 5_1 + 10_7
\]  

(6)

where subscripts denote the \( U(1) \) charge of generator \( Y' = \text{Diag}(1, 1, 1, 1, 1, -5) \) of \( SU(6) \). Therefore if \( 24_{-5} \) has a nonzero VEV in \( G_{321} \) direction the \( Y' \) generator is broken while ordinary hypercharge \( Y = \text{Diag}(2, 2, 2, -3, -3, 0) \) remains unbroken.
Analyzing the superpotential (3), one can see that there is an unique non-trivial supersymmetry conserving minima (with vanishing F and D terms) with the VEV configurations \( \langle \Phi \rangle \) and \( \langle \Psi \rangle \) that imply the \( SU(6) \) symmetry breaking to \( G_{321} \). More explicitly, \( \langle \Phi \rangle \) has a form (5) while \( \langle \Psi \rangle \) is the following:

\[
\langle \Psi_{BC,m}^A \rangle = U \left[ \delta^6_C [2\delta^4_B - 5(\delta^4_A \delta^4_B + \delta^5_A \delta^5_B)] - \delta^6_B [B \to C] \right] \delta_{1m}
\]  

(The VEV of \( \overline{\Psi} \) is the same). The magnitudes of these VEVs are the following:

\[
V = \frac{M_\Psi}{10h}, \quad U = \frac{3}{5h} \left( \frac{6}{5} M_\Phi M_\Psi \right)^{1/2}
\]

Since \( \text{Tr}(\Phi^3) = 0 \), \( V \) and \( U \) do not depend on the constant \( \lambda \).

We see from (3) that 84 contains two doublets which are compressed in the fragments 5 and 45. After the \( SU(6) \) symmetry breaking to \( G_{321} \) one combination of these doublets, which in terms of the \( \Psi \) fragments reads as

\[
h^m_w = \frac{1}{\sqrt{21}} (-2\Psi^c_{cw} + 3\Psi^w_{w'}^m)
\]

is massless, while another combination

\[
H^m_w = \frac{1}{\sqrt{7}} (\Psi^c_{cw} + 2\Psi^w_{w'}^m)
\]

has the mass of order \( M_X \) (here \( c \) is the \( SU(3)_C \) index and \( w, w' \) are the \( SU(2)_W \) indices). The same applies to the conjugated states \( \overline{h} \) and \( \overline{H} \). As far as the pairs \( (\overline{h}, h)_m \), \( m = 1, 2 \) are related by the \( SU(2)_{\text{cusp}} \) symmetry they are both massless. First pair \( (\overline{h} + h)^1 \) is a genuine Goldstone mode which is eaten up by the corresponding gauge superfield of \( SU(6) \), while the second one survives after \( G_X \) breaking as a pseudogoldstone mode which can get \( \sim M_W \) mass only after the supersymmetry breaking.

Note, that the coupling \( \Phi \Psi \overline{\Psi} \) does not affect the structure for the VEV \( \langle \Phi \rangle \) and it maintains the pattern (4). Although the term \( h\Phi \Psi^m \Psi_m \) in eq. (3) violates the extra global symmetries, from the ”view” of doublet-antidoublet fragments in \( \overline{\Psi} + \Psi \) the VEV \( \langle \Phi \rangle \) is a singlet. After substituting the VEVs of \( \Phi \) and \( \Psi, \overline{\Psi} \) in the superpotential, due to the \( SU(2)_{\text{cusp}} \) symmetry no mixing terms emerge between the doublet (antidoublet) modes of \( \Psi_1(\overline{\Psi}^1) \) and \( \Psi_2(\overline{\Psi}^2) \). Since \( \Phi \) itself does not include the doublet modes. So one doublet-antidoublet is massless pseudogoldstone until SUSY is unbroken. In this context this situation resembles the pseudogoldstone picture [4, 5, 8] but the difference is that we

\[\text{Here the indices 4 and 5 stand for } SU(2)_W, \text{ and } 1, 2, 3 \text{ for the } SU(3)_C. \text{ Index 6 corresponds to the broken sixth degree of freedom of the } SU(6) \text{ gauge group.} \]
do not have the $SU(6) \times SU(6)$ global symmetry in the superpotential, but due to the structure of $Φ$ and $SU(2)_{\text{cns}}$ symmetry the doublets can be rotated away from the Higgs superpotential. This can not be done for triplets from $Ψ$ because $Φ$ itself contains the triplet fragments and there occurs the mixing between triplets (antitriplets) from 175 and 84. Without losing of generality one can choose the basis in which the mass matrix has the form:

$$
\begin{pmatrix}
3_{175} & 3_{841} & 3'_{841} & 3_{842} & 3'_{842} \\
3_{175} & M_Φ & αU & βU & 0 & 0 \\
3_{841} & αU & M_Φ & 0 & 0 & 0 \\
3'_{841} & βU & 0 & M_Ψ & 0 & 0 \\
3_{842} & 0 & 0 & 0 & M_Ψ & 0 \\
3'_{842} & 0 & 0 & 0 & 0 & M_Ψ
\end{pmatrix}
$$

(11)

where

$$
\text{Det}
\begin{pmatrix}
M_Φ & αU & βU \\
αU & M_Ψ & 0 \\
βU & 0 & M_Ψ
\end{pmatrix}
= 0
$$

(12)

($α$ and $β$ are some Clebsch factors which are not important for us). Therefore, one eigenstate is massless and is identified as a Goldstone mode. Exact calculations show that all other states have masses of order $M_X$.

Fermion sector can be arranged in the same manner as in ref. [6]: the quark-lepton masses are generated from the following Yukawa superpotential:

$$
W_Y = g_d \bar{d}^m 15 m_1 15 Ψ_m Ψ_n \epsilon_{mn}
$$

(13)

Second coupling can be obtained by heavy particle exchange mechanism [10]; After introducing two 20-plets ($≡ 20^m$) the renormalizable couplings which are responsible for generation of $g_u$ term in (13) are

$$
W_Y^0 = g_1 15 20^m Ψ_m + M_0 20^m 20^m \epsilon_{mn}
$$

(14)

After integrating out the heavy $20^m$ states below the $M_0$ scale we are left with the $g_u$ effective coupling.

$\mathbf{R = 210}$: In this case $210 \equiv Ψ_A^{A'B'C'}$ is antisymmetric with respect to the up and down indices and $ΣΨ_A^{A'B'C'} = 0$. In terms of $SU(5)$ 210 reads as:

$$
210 = 75 + 50 + 45 + 40
$$

(15)

where the $G_{321}$ singlet is contained in 75 and the MSSM doublet fragment in 45.
The Higgs superpotential still has the same form (3), however now the coupling $h\Phi\bar{\Psi}\Psi$ implies three invariants:

$$h\Phi\bar{\Psi}\Psi \equiv \sum_{i=1}^{3} h_i I_i$$  \hspace{1cm} (16)

where

$$I_1 = \bar{\Psi}^{ABC}_1 \Phi^{A_1B_1C_1}_1 \Psi^{A'B'}_{A_1B_1C_1}
I_2 = \bar{\Psi}^{ABC}_1 \Phi^{A_1B_1C_1}_1 \Psi^{A'B'}_{A_1B_1C_1}
I_3 = \bar{\Psi}^{ABC}_1 \Phi^{A_1B_1C_1}_1 \Psi^{A'B'}_{A_1B_1C_1}$$  \hspace{1cm} (17)

The $G_{321}$ invariant VEV of $\Psi(\bar{\Psi})$ has the form:

$$\Psi^{12}_{126} = \Psi^{13}_{136} = \Psi^{23}_{236} = U
\Psi^{14}_{146} = \Psi^{24}_{246} = \Psi^{34}_{346} = \Psi^{15}_{156} = \Psi^{25}_{256} = \Psi^{35}_{356} = -U
\Psi^{45}_{456} = 3U$$  \hspace{1cm} (18)

Supersymmetric minima allows to have nonvanishing $V$ and $U$ with the magnitudes:

$$V = \frac{9M_\Psi}{8(9h_1-h_2+h_3)}, \quad U = \frac{9}{4(9h_1-h_2+h_3)} (\frac{M_\Psi M_\Phi}{2})^{1/2}$$  \hspace{1cm} (19)

and the "philosophy" is the same as in the 84-plets case: after the $SU(6)$ gauge symmetry breaking two pairs of the Higgs doublets $\bar{h}^m + h_m$ remain massless. One pair is absorbed by the appropriate gauge fields which became superheavy and the second one survives. Therefore in the effective low-energy theory we will have one pair of massless Higgs doublet-antidoublet.

In this case the couplings relevant for the quark and lepton masses are the following:

$$W_Y = \frac{g_d}{M_1} \bar{\Phi} \bar{\Psi}^m \epsilon_{mn} + \frac{g_u}{M_0} 15 \bar{\Psi}_{m} \Psi_{n} \epsilon_{mn}$$  \hspace{1cm} (20)

$g_u$ term can be generated in the same manner as was discussed in the 84-plet’s case, while the renormalizable Yukawa superpotential which is responsible for generation of the $g_d$ term has the form:

$$W'_Y = g' \bar{\Phi} \bar{\Psi}^m_{105} \epsilon_{mn} + g'' 15 \bar{\Psi} \Psi_{105} + M_1 \bar{105} 105$$  \hspace{1cm} (21)

Integrating out the heavy $105 + 105$ states, the first term of eq. (20) is obtained with $g_u \sim g'g''$.

As we see in this case the fermion sector requires more complicated multiplets because it is impossible to write renormalizable Yukawa couplings for down quarks and leptons.
However if the mass of 105-plets $M_1$ is $10^{18}$ GeV order, then after their decoupling the effective Yukawa constants for third generation of down quark and lepton will have just needed magnitude $-\frac{M_{X}}{M_1} \sim 10^{-2}$. More detailed study of fermion masses in our model will be presented elsewhere.

Concluding, we have suggested supersymmetric $SU(6)$ theory in which the DT splitting occurs naturally. Although this mechanism is based on the custodial symmetry, the lightness of Higgs doublets has the different origin then in the model [3]. Crucial feature is that the 175 not contain the Higgs doublets, and consequently there emerges no mixing between doublet components of the $SU(6)$ symmetry breaking scalars. This feature allows to achieve the one point unification of the gauge couplings at the scale $M_G \sim 10^{16}$ GeV. (Recall, that in the model [3] this mixing was rendering the scale $M_f$ the middle geometrical, $M_f \sim \sqrt{M_G M_W}$). Therefore, the ‘missing doublet’ multiplet 175 of $SU(6)$ is very attractive for the model building. Its properties admit Higgs doublets to be massless till SUSY is unbroken. Besides 175, the higher dimensional selfconjugate representations of $SU(6)$ which do not contain doublet fragments are 3963 and 4116 [11], which is clearly too much. As far the larger unitary groups $SU(6+N)$, one can make sure that they have no selfconjugate ”missing doublet” multiplets which could be used for the symmetry braking. Therefore, the $SU(6)$ group appears to be the single group in which the presented mechanism can be realized.

Acknowledgement
I am grateful to Z. Berezhiani, J. Chkareuli, G. Dvali, I. Gogoladze and A. Kobakhidze for useful discussions and important comments.

References

[1] P. Langacker and M. Luo, Phys. Rev. D44 (1991) 817;
U. Amaldi, W. de Boer and H. Furstenau, Phys. Lett. B260 (1991) 447;
J. Ellis, S. Kelley and D. Nanopoulos, Phys. Lett. B260 (1991) 131.

[2] S. Dimopoulos and F. Wilczek, in Erice Summer Lectures, Plenum, New York, 1981;
H. Georgi, Phys. Lett. B108 (1982) 283;
B. Grinstein, Nucl. Phys. B206 (1982) 387;
A. Masiero, D.V. Nanopoulos, K. Tamvakis and T. Yanagida, Phys. Lett. B115 (1982) 380;
J.L. Lopez and D.V. Nanopoulos, [hep-ph/9508253].
[3] S. Dimopoulos and F. Wilczek, NSF-ITP-82-07 (unpublished);
    M. Srednicki, Nucl. Phys. B202 (1982) 327;
    K.S. Babu and S.M. Barr, Phys. Rev. D48 (1993) 5354; Phys. Rev. D50 (1994) 3529.

[4] K. Inoue, A. Kakuto and T. Takano, Progr. Theor. Phys. 75 (1986) 664;
    A. Anselm and A. Johansen, Phys. Lett. B200 (1988) 331;
    R. Barbieri, G. Dvali and A. Strumia, Nucl. Phys. B391 (1993) 487.

[5] Z. Berezhiani and G. Dvali, Sov. Phys. Lebedev Institute Reports 5 (1989) 55;
    R. Barbieri, G. Dvali and M. Moretti, Phys. Lett. B312 (1993) 137;
    Z. Berezhiani, C. Csaki and L. Randall, Nucl. Phys. B444 (1995) 61.

[6] G. Dvali, Phys. Lett. B324 (1994) 59.

[7] G.V. Anderson et al., Phys. Rev. D49 (1994) 3660;  
    K.S. Babu and R.N. Mohapatra, Phys. Rev. Lett. 74 (1995) 2418;  
    K.S. Babu and S.M. Barr, ibid. 75 (1995) 2088;  
    Z. Berezhiani, Phys. Lett. B355 (1995) 178.

[8] R. Barbieri, G. Dvali, A. Strumia, Z. Berezhiani, L.Hall, Nucl.Phys. B432 (1994) 49;  
    Z. Berezhiani, Phys. Lett. B355 (1995) 481.

[9] I. Gogoladze, A. Kobakhidze and Z. Tavartkiladze, Phys. Lett. B372 (1996) 246.

[10] C.D. Froggatt and H.B. Nielsen, Nucl. Phys. B147 (1979) 277;  
    Z.G. Berezhiani, Phys. Lett. B129 (1983) 99; B150 (1985) 177;  
    S. Dimopoulos, Phys. Lett. B129 (1983) 417.

[11] R. Slansky, Phys. Rep. 79 (1981) 3.