Symmetry Breaking by Electric Discharges in Water and Formation of Light Magnetic Monopoles in an Extended Standard Model (Part III)

Harald Stumpf

Institute of Theoretical Physics, University Tuebingen, Germany

Reprint requests to H. S., Ahornweg 6, D-72072 Tuebingen, Germany

Z. Naturforsch. **67a,** 173 – 179 (2012) / DOI: 10.5560/ZNA.2012-0009
Received October 11, 2011

By Lochak (theory) and Urutskoev (experiment) the hypothesis has been suggested that during electric discharges in water light magnetic monopoles can be created and according to Lochak these monopoles should be looked upon as a kind of excited neutrinos. Based on a quantum field theoretic derivation of an extended (effective) Standard Model, which describes the effects of symmetry breaking caused by the discharge, in Part I and Part II (H. Stumpf, Z. Naturforsch. **66a**, 205 and 329 (2011)), the possibility was investigated whether under these circumstances magnetic effects can occur. While in the above two parts the focus of the theory was directed to show the evidence of magnetically active neutrinos, in this part the behaviour of the leptonic doublets (ν, e^-) and $(\bar{\nu}, \mathrm{e}^+)$ is studied under the (symmetry breaking) conditions of such an electric discharge. The results obtained by the quantum field theoretic treatment of this problem are in agreement with symmetry considerations earlier performed by Lochak.

Key words: Magnetic Monopoles; Discharges in Water, Extended Standard Model.

1. Introduction

By Lochak (theory) and Urutskoev (experiment) the hypothesis has been suggested that during discharges in water (fluids) light magnetic monopoles can be created which, according to Lochak, should be considered as a kind of excited neutrinos.

This hypothesis was motivated by the idea that weak nuclear reactions could be catalysed by strong magnetic fields which magnetic monopoles can exert on their surroundings [1 – 3].

In Part I [4] it was explained that a theoretical evaluation of this hypothesis must be based on the ideas of de Broglie and of Heisenberg about a fermionic substructure of elementary particles. In a modern theoretical description this finally leads to the derivation of an extended Standard Model which includes composite particles and new effects, in particular in connection with magnetic monopoles.

Within this formalism in Part II [5] a proof for the existence of such magnetic monopoles was given. An *essential* prerequisite for the success of these calculations is the symmetry breaking generated by the discharge in water.

Furthermore, it was demonstrated in Part I that the discharges not only generate symmetry breaking, but in addition to the electronic current also induce proton captures by the nucleons of the anode which in this way are transmuted into unstable elements. Afterwards this eventually ends up in weak decays of the unstable elements due to their proton surplus.

By the decay $p = n + e^+ + \nu$ the weak reactions are responsible for the appearence of magnetic monopoles which result from neutrinos emitted under symmetry breaking conditions during this process.

The investigations in Part II are only concerned with these under symmetry breaking modified or deformed (excited) neutrino states and their interactions with electroweak bosons which are likewise modified by the symmetry breaking. But the above reaction equation at once raises the question: How react the charged leptons to this symmetry breaking. The present Part III is devoted to the answer of this *urgent* question.

Like in Part I and Part II it is also unavoidable to refer to preceding papers in this part without giving renewed deduction of results obtained there. In the same way this holds for the work of Lochak and Urutskoev for which extensive references were given in Part I and Part II. The deduction of the extended Standard Model itself is based on a nonlinear spinor field theory with local interaction, canonical quantization, self-regularization, and probability interpretation. The corresponding proofs were given in [6, 7] and in [8] (decomposition theorem).

These proofs are basic for the whole formalism because the real problem of a nonlinear spinor theory is the *compatibility* with various demands like relativistic invariance, locality and causality, finiteness, and probability interpretation. In the past in such spinor theory approaches these demands could not be satisfied in one or the other way, and it is just the progress with respect to the compatibility of these demands which was achieved by [6 – 8].

2. Symmetry and Symmetry Breaking

The algebraic representation is the basic formulation of the spinor field model. In order to avoid lengthy deductions, we give only the basic formula of this model and refer for details to [9, 10]. The corresponding Lagrangian density reads, see [9] (2.52),

$$\begin{aligned}\mathcal{L}(x) := &\sum_{i=1}^{3} \lambda_i^{-1} \bar{\psi}_{A\alpha i}(x)(\mathrm{i}\gamma^\mu \partial_\mu - m_i)_{\alpha\beta} \delta_{AB} \psi_{B\beta i}(x) \\ &- \frac{1}{2} g \sum_{h=1}^{2} \delta_{AB} \delta_{CD} v^h_{\alpha\beta} v^h_{\gamma\delta} \\ &\cdot \sum_{i,j,k,l=1}^{3} \bar{\psi}_{A\alpha i}(x) \psi_{B\beta j}(x) \bar{\psi}_{C\gamma k}(x) \psi_{D\delta l}(x)\end{aligned} \tag{1}$$

with $v^1 := \mathbf{1}$ and $v^2 := \mathrm{i}\gamma^5$. The field operators are assumed to be Dirac spinors with index $\alpha = 1,2,3,4$ and additional isospin with index $A = 1,2$ as well as auxiliary fields with index $i = 1,2,3$ for nonperturbative regularization [9] (2.51) and [6]. The algebra of the field operators is defined by anticommutators originating from canonical quantization [9].

2.1. Conserved Symmetries

With respect to the Poincare group this topic was treated in great detail in [9] Sects. 2.7 and 3.3. Concerning internal symmetry groups for the $SU(2) \otimes U(1)$ invariance of the initial spinor field equation, or Lagrangian (1), respectively, the good quantum numbers result as eigenvalues of the symmetry generators in [9] (6.80) – (6.83) which in functional space can be formulated in the form of (6.84) – (6.89) in [9].

Provided that the propagator is invariant under the action of these groups, the corresponding eigenvalue conditions of [9] (6.84) – (6.89) hold for normal transformed states too.

2.2. Discrete Transformations

In literature, corresponding operations and operators are explicitly constructed in Fock-space with Dirac-vacuum, cf. for instance [11, 12]. In [13] these definitions were generalized to be valid in the algebraic formalism. In particular it can be proved that the Lagragian (1) is invariant under parity operation, cf. [14] p. 62, 426, and [15] p. 49. I.e. the Lagrangian (1) defines a *left-right symmetric model*. Furthermore, by construction of the model the conditions for the application of the PCT- theorem, cf. [11], are fulfilled. I.e. *the model is invariant under the PCT-operation* and the fermion propagator for conserved symmetries is PCT- and CP-invariant too, cf. [16].

2.3. Antisymmetrization

For Fermi-fields time-ordered antisymmetric products of field operators are defined in the covariant perturbation theory. The definition of such products was applied to the spinor fields of (1), and dynamical field equations were derived in functional space [9] Chapt. 3: Against the use of this formalism beyond perturbation theory, i.e. for non-perturbative, theoretical calculations serious mathematical objections can be raised [9] Sect. 3.7.

An appropriate alternative treatment was developed by means of the *algebraic Schrödinger representation* which is based on the GNS-construction (Gelfand–Naimark–Segal) [9] Sect. 4.4, [10] Sect. 3.5, and which is *exclusively* referred to in the following.

The GNS-construction provides genuine representations of the basic field operator algebra, and the set of basis vectors is generated by the application of antisymmetric products of field operators on a spacelike hyperplane to the groundstate (vacuum) $|0\rangle$ of the system. Furthermore, in order to obtain a simplified formulation of the states expressed by the GNS-construction, the basis states can be integrated into generating functional states in functional space [17] (7). Such functional states can then be chosen as eigenstates of energy [17] (10).

But as an infinite number of inequivalent vacuum states exists, it is impossible to draw any conclusions from the functional eigenvalue equation [17] (10) *without specifying the vacuum.*

2.4. Symmetry Breaking and Parafermionic States

The selection of a definite set of basis states can be done by application of the normal ordering transformation [17] (16). A symmetry breaking can then be implemented by the assumption of an appropriately modified propagator in the normal ordering transformation which replaces the free propagator for conserved symmetries. In [16] it was demonstrated that a symmetry breaking is caused by a discharge in water. In water, one can discriminate the positive charges from the negative charges by their different damping. The latter leads to an irreversible time evolution, i.e., the time reversal invariance is violated. Owing to the conserved PCT-invariance of the whole theory the violation of time reversal has to be compensated by a CP-symmetry breaking in order to maintain the PCT-invariance of the theory. (An alternative argument for this CP-violation is given in [16].)

CP-violation is characteristic for weak processes, cf. [18], Sect. 4.1 and if in this case left-right symmetric models are applied the following comment refers to this approach:

"*The basic premise of the left–right symmetric models (in electroweak theory) is ... that the parity asymmetry observed in nature, such as β decay and μ decay arise from vacuum (states) being noninvariant under parity symmetry*" and "*we wish to emphasize as a motivation for left–right symmetric theories that there is now convincing evidence for neutrino masses from observation*" Mohapatra [18] Sect. 6.1 and also [19].

The left-right symmetric model (1) differs from Mohapatra's theory, but it shares the basic premises with it. Thus it has to be expected that for CP-symmetry breaking discharges in water there is an opportunity to find those magnetic monopoles which are correlated with weak interactions and which were assumed and looked for by Urutskoev and Lochak.

The fact that by symmetry breaking a discrimination between various particle species is possible prevents the use of antisymmetric states and leads to the application of *parafermionic* configurations studied in Part II.

3. Simplified Mapping with Parafermionic States

A description of the effective motion of composite particles can be obtained by weak mapping theorems. A prerequisite of these theorems is the complete antisymmetrization of the states for conserved symmetries. But due to symmetry breaking the resulting parafermionic states imply the loss of antisymmetry. To compensate the loss of antisymmetry the exact mapping theorems have to be replaced by a simplified version of the weak mapping procedure in order to avoid unnecessary complications.

To this end, one has to note that the results of weak mappings are used for the derivation of effective theories. In effective theories exchange terms are neglected. With the neglection of these terms the necessity for the application of faithful mapping theorems is dropped and one can apply the socalled chain rule which by construction refrains from exchange terms, cf. [9] Sects. 5.6 – 5.8 and also [10].

With respect to the monopol problem, the coupling of the composite fermions (leptons) to the composite vector bosons is of interest. The corresponding term is of lowest order and the chain rule expression differs from the expression of the exact mapping term only by an additional overlap term.

For both approaches the common term reads

$$\begin{aligned}\mathcal{H}^2_{bf} &:= W_2^{qlp}\partial_l^b f_q \partial_p^f \\ &= 3W_{I_1I_2I_3I_4}R^q_{II'I_1}C^l_{I_2I_3}C^p_{I_4II'}\partial_l^b f_q\partial_p^f\end{aligned} \tag{2}$$

and can be taken from [9] (5.74) where the overlap term here and in the following is ignored.

In the above term the symmetry breaking manifests itself exclusively in the internal structures of the wave functions and (II.11), the starting point of our monopol investigation in Part II, is identical with (2).

It should be noted that the parafermionic wave functions of Part II are all factorized into products of superspin–isospin configurations times the spin–orbit configurations in (II.6), (II.7) which *inevitably involves the application of G-conjugated spinors*, defined in Sect. 1 of Part II.

Owing to the identity (II.11) ≡ (2), one can adopt the calculations of Part II, provided these calculations can be extended to other lepton states apart from the neutrino states of Part II. A convenient starting point for this adoption is the Dirac equation (II.36) where the neutrino indices can be replaced by the general lepton index l.

Three remarks should be added about the meaning of this equation:

i) In this equation, the coupling terms of the leptons to the charged vector bosons have been omitted because this coupling is not relevant for our problem as will be shown in the following section.

ii) Originally, this equation was meant for the proof of excited neutrino states as monopoles. But one can easily see that this equation can likewise be applied to any other lepton state because, apart from the group theoretical modifications, the necessary adaption concerns only the orbital wave function Υ in (II.25). By definition (II.31), this function is assumed to be independent of the state quantum numbers. In this context, one should recall that in the conventional theory leptons are considered as structureless point particles, so that (II.31) seems to be an acceptable assumption and as long as no information about the solutions of the BBW-equation for leptons is available only some global deformation parameters can be introduced, see below.

iii) In analogy to the derivation of classical effective equations from a functional equation, cf. [9] (8.71) – (8.73), in Part II in (II.36) an effective Dirac equation for neutrinos was obtained. The generalization of (II.36) to lepton states then reads

$$
\begin{aligned}
&\mathrm{i}\partial_t \psi_{l\alpha}(x) = \\
&\left[-\mathrm{i}(\gamma^0\gamma^k)_{\alpha\beta}\partial_k + \gamma^0_{\alpha\beta} m_l^*\right]\psi_{l\beta}(x) \\
&\left\{+g_A^*(\gamma^0\gamma^k)_{\alpha\beta}\zeta_\rho(l)^D (T^0\gamma^5)^D_{\rho\rho'}\zeta_{\rho'}(l)^D A_k^0(x)\right. \\
&\left.-g_A'^*(\gamma^0\gamma^k)_{\alpha\beta}\zeta_\rho(l)^D (T^3\gamma^5)^D_{\rho\rho'}\zeta_{\rho'}(l)^D A_k^3(x)\right\}\psi_{l\beta}(x) \\
&\left\{+\mathrm{i}g_G^*(\gamma^0\gamma^k\gamma^5)_{\alpha\beta}\zeta_\rho(l)^D (S^0\gamma^5)^D_{\rho\rho'}\zeta_{\rho'}(l)^D G_k^0(x)\right. \\
&\left.-\mathrm{i}g_G'^*(\gamma^0\gamma^k\gamma^5)_{\alpha\beta}\zeta_\rho(l)^D (S^3\gamma^5)^D_{\rho\rho'}\zeta_{\rho'}(l)^D G_k^3(x)\right\}\psi_{l\beta}(x)
\end{aligned}
\tag{3}
$$

with the effective coupling constants

$$
\begin{aligned}
g_Z^* &:= \hat{f}^Z(0)\left(\frac{2\pi^2 9a}{5}\right)^{1/2}, \\
g_Z'^* &:= c^3\hat{f}^Z(0)\left(\frac{2\pi^2 9a}{5}\right)^{1/2}, \quad Z = A, G,
\end{aligned}
\tag{4}
$$

where c^3 is an element of the set of deformation parameters c^1, c^2, c^3 which were introduced in (II.25) – (II.35) (and in preceding papers) but not further pursued in (II.36) and (II.37).

From (4) it follows $g_A'^* = c^3 g_A^*$ and without further calculation one can fit c^3 to that value which is used to define the universal (eletric) Weinberg transformation. Let $\Theta_\mathrm{W}^\mathrm{e}$ be the universal (electric) Weinberg angle. Then with [20] (6.29), c^3 can be defined by

$$
\tan\Theta_\mathrm{W}^\mathrm{e} = \frac{g_A^*}{g_A'^*} = \frac{1}{c^3}
\tag{5}
$$

if one observes that in [20] (6.23) the roles of g and g' are reversed in comparison with those in (3).

The (electric) Weinberg transformation can then be defined by use of the definitions

$$
\sin\Theta_\mathrm{W}^\mathrm{e} = \frac{g_A^*}{N_A} = \beta, \quad \cos\Theta_\mathrm{W}^\mathrm{e} = \frac{g_A'^*}{N_A} = \beta'
\tag{6}
$$

with $N_A = (g_A^{*2} + g_A'^{*2})^{1/2}$ which leads to

$$
\begin{aligned}
\mathbf{A^3} &= \beta'\mathbf{Z} + \beta\mathbf{A}, \\
\mathbf{A^0} &= -\beta\mathbf{Z} + \beta'\mathbf{A}
\end{aligned}
\tag{7}
$$

and corresponds to [20] (6.24).

With respect to the coupling of the leptons to the magnetic bosons, a similar transformation has to be applied, and one can conclude that the magnetic Weinberg angle $\Theta_\mathrm{W}^\mathrm{m}$ must be equal to the eletric angle $\Theta_\mathrm{W}^\mathrm{e}$. Putting $f^G(0) = \lambda f^A(0)$, one obtains

$$
\begin{aligned}
g_G^* &= \lambda g_A^*, \\
g_G'^* &= \lambda g_A'^*,
\end{aligned}
\tag{8}
$$

and consequently it must hold

$$
\frac{1}{c^3} = \frac{g_G^*}{g_G'^*} = \tan\Theta_\mathrm{W}^\mathrm{m},
\tag{9}
$$

where due to (8) the definitions (6) also hold if the electric coupling constants are replaced by the the magnetic ones.

The appearance of magnetic potentials is new in the theory and although the electric and magnetic Weinberg angles are equal, there is *no compelling evidence* to apply the same Weinberg transformation to the magnetic potentials as to the electric ones.

On the contrary, as there exists a duality among electric and magnetic quantities, it is a heuristic working hypothesis to apply the *dual* Weinberg transformation to the magnetic quantities which gives

$$
\begin{aligned}
\mathbf{G^3} &= \beta'\mathbf{X} - \beta\mathbf{G}, \\
\mathbf{G^0} &= \beta\mathbf{X} + \beta'\mathbf{G},
\end{aligned}
\tag{10}
$$

and owing to $\Theta_{\mathrm{W}}^{\mathrm{e}} = \Theta_{\mathrm{W}}^{m}$, one gets

$$\beta = \beta_A = \frac{g_A^*}{N_A} = \frac{g_G^*}{N_G} = \beta_G\,, \qquad \beta' = \beta_A' = \frac{g_A'^*}{N_A} = \frac{g_G'^*}{N_G} = \beta_G'\,. \tag{11}$$

These equalities define the A- and the G-representation, and one can formulate (7) in the A-representation just as (10) can be formulated in the G-representation.

If one substitutes (7) and (10) into (3), only the **A** and the **G** terms are of interest because they define the coupling of the electric and the magnetic bosons to the leptons. One gets from (3)

$$\begin{aligned} H_{bf}^{\mathrm{em}}\psi_l &= \frac{g_A^* g_A'^*}{N_A}\Big[\zeta_\rho^D(l)(T^0\gamma^5)_{\rho\rho'}^D\zeta_{\rho'}^D(l) \\ &- \zeta_\rho^D(l)(T^3\gamma^5)_{\rho\rho'}^D\zeta_{\rho'}^D(l)\Big](\gamma^0\gamma)_{\alpha\beta}\mathbf{A}(x)\psi_{l\beta}(x) \\ &+ \imath\frac{g_G^* g_G'^*}{N_G}\Big[\zeta_\rho^D(l)(S^0\gamma^5)_{\rho\rho'}^D\zeta_{\rho'}^D(l) \\ &+ \zeta_\rho^D(l)(S^3\gamma^5)_{\rho\rho'}^D\zeta_{\rho'}^D(l)\Big](\gamma^0\gamma\gamma^5)_{\alpha\beta}\mathbf{G}(x)\psi_{l\beta}(x)\,. \end{aligned} \tag{12}$$

It should be emphasized that the superspin–isospin parts in the brackets of (12) are real and that the signs and factors in front of the brackets are in agreement with the result of Lochak's quantum mechanical monopole theory.

4. Discharge Effects on Leptonic Doublets

The superspin–isospin matrices of (12) were given in [21] (33) and (34) in D- (i.e. G-conjugated) representation. For convenience, they are explicitly reproduced:

$$(S^0\gamma^5)^D = \begin{pmatrix} \mathbf{1} & 0 \\ 0 & \mathbf{1} \end{pmatrix}, \quad (T^0\gamma^5)^D = \begin{pmatrix} \mathbf{1} & 0 \\ 0 & -\mathbf{1} \end{pmatrix} \tag{13}$$

and

$$(S^3\gamma^5)^D = \begin{pmatrix} \sigma^3 & 0 \\ 0 & -\sigma^3 \end{pmatrix}, \quad (T^3\gamma^5)^D = \begin{pmatrix} \sigma^3 & 0 \\ 0 & \sigma^3 \end{pmatrix}. \tag{14}$$

If these matrices are inserted into (12), one gets

$$\zeta_\rho^D(l)\begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & 2 & 0 & 0 \\ 0 & 0 & -2 & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix}_{\rho\rho'} \cdot\,\zeta_{\rho'}^D(l) \cdot \frac{g_A^* g_A'^*}{N_A}(\gamma^0\gamma)_{\alpha\beta}\mathbf{A}(x)\psi_{l\beta}(x) \tag{15}$$

and

$$\zeta_\rho^D(l)\begin{pmatrix} 2 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 2 \end{pmatrix}_{\rho\rho'} \cdot\,\zeta_{\rho'}^D(l) \cdot \frac{g_G^* g_G'^*}{N_G}(\gamma^0\gamma\gamma^5)_{\alpha\beta}\mathbf{G}(x)\psi_{l\beta}(x)\,. \tag{16}$$

The superspin–isospin states $\zeta(l)$ are defined in (II.26). For broken symmetry, the superspin–isospin three-particle states Θ^l are reduced to a single product of these one-particle states, cf. Part II Sect. 2.2. But in addition to (II.26), the one-particle states are further restricted in the present case. For diagonal matrices like in (15) and (16), (II.26) changes into

$$\zeta_\rho(l)\otimes\zeta_{\rho'}(l) := \Theta_{\kappa\kappa'\rho}^l\Theta_{\kappa\kappa'\rho'}^l\,, \tag{17}$$

i.e. there remain only four superspin–isospin unit vectors for the state classification, as Θ^l itself is a direct product of such unit vectors.

However, before any further conclusion can be drawn, the question must be answered: Does CP-symmetry breaking change the conventional electroweak quantum numbers of the leptons? These numbers are defined by isospin t and isospin component t_3, fermion number f, hypercharge Y, and electric charge q. These quantities are linked with the numbering of representations of a $SU(2)\otimes U(1)$ group in the superspin–isospin space by the eigenvalues of the group generators, cf. [9] Chapt. 6.

An answer to the above question is given by proposition 6.5 in [9] p. 165:

The normal ordered state functionals (i.e. the states referred to the modified groundstate) satisfy the same group theoretical eigenvalue equations as the states for conserved symmetry, provided the propagator is invariant under the $SU(2)\otimes U(1)$ group.

For a proof of the invariance of F, the relevant terms are the operators $\gamma_{\kappa\kappa'}^5$ and $(\gamma^5\gamma^0)_{\kappa\kappa'}$ in superspin–isospin space where the latter is responsible for CP-symmetry breaking. Both these terms are singlets of the superspin–isospin group representations and leave F invariant under such transformations.

Thus for CP-symmetry breaking the conventional quantum numbers of leptons can be adopted to classify the vectors $\zeta_\rho(l)$.

On account of this result, it is no longer necessary to make use of the table in [9] (6.108). Rather from

(15), one can guess that the identities $\delta_{\rho\,2} \equiv \mathrm{e}^+$ and $\delta_{\rho\,3} \equiv \mathrm{e}^-$ must hold, and that from (16) the identities $\delta_{\rho\,1} \equiv \bar{\nu}$ and $\delta_{\rho\,4} \equiv \nu$ can be concluded.

So in D-representation, one gets

$$\begin{pmatrix} \psi_{1\beta} \\ \psi_{2\beta} \end{pmatrix}^D \equiv \begin{pmatrix} \psi_\beta(\bar{\nu}) \\ \psi_\beta(\mathrm{e}^+) \end{pmatrix}^D, \quad \begin{pmatrix} \psi_{3\beta} \\ \psi_{4\beta} \end{pmatrix}^D \equiv \begin{pmatrix} \psi_\beta(\mathrm{e}^-) \\ \psi_\beta(\nu) \end{pmatrix}^D. \tag{18}$$

For a physical classification of the state vectors, the D-representation is not very well suited because in the conventional Standard Model the S-representation is preferred. By application of the transformation (II.5), one can return to the S-representation which leads to the relations

$$\begin{pmatrix} \psi_{1\beta} \\ \psi_{2\beta} \end{pmatrix}^D \to \begin{pmatrix} \psi_\beta(\bar{\nu}) \\ \psi_\beta(\mathrm{e}^+) \end{pmatrix}^S, \quad \begin{pmatrix} \psi_{3\beta} \\ \psi_{4\beta} \end{pmatrix}^D \to \begin{pmatrix} \psi_\beta(\nu) \\ \psi_\beta(\mathrm{e}^-) \end{pmatrix}^S, \tag{19}$$

i.e. one gets the *doublets of the first lepton generation* in the conventional Standard form. From (15), (16), and (18) it follows that *the states* $\psi^D_{1,\beta} \equiv \bar{\nu}$, $\psi^D_{4,\beta} \equiv \nu$ *have equal magnetic charges and zero electric charge* and *the states* $\psi^D_{2,\beta} \equiv \mathrm{e}^+$ $\psi^D_{3,\beta} \equiv \mathrm{e}^-$ *have opposite electric charges and zero magnetic charge*.

The strange behaviour of the magnetic charge g under electric, i.e. ordinary spinorial, charge conjugation is in accordance with the quantum mechanical investigations of Lochak: "*The most important feature appears in the formula (26): the charge conjugation does not change the sign of the magnetic charge g*" [23], cf. also [3] (7.2).

The above statements must be completed by some additional remarks:

i) In literature, it is reported that in the Standard Model with (nearly) vanishing neutrino masses the lepton states are not influenced by CP-symmetry breaking, cf. [20] p. 116 and [22] p. 47. Thus by the application of this statement in [17], the calculations of the effects of CP-symmetry breaking in an extended Standard Model have led to the transformation of charged lepton states into dyon states.

But this result cannot be maintained in view of the new facts presented in this part and of the fact that the conventional Standard Model and the derivation of the extended Standard Model are based on different presuppositions and, hence, are not equivalent.

ii) Another point is the problem of the generations. In the spinor field model, the set of states of the first lepton generation with fermionic substructure is only characterized by the set of different superspin–isospin states while the orbit parts of the complete wave functions are (assumed) to be equal and are given (approximately) by formula (II.31). In contrast to this construction the wave functions for the higher lepton generations are characterized by the same superspin–isospin configurations as those of the first generation, but the orbit parts of the wave functions must differ from the wave function of the first generation by orbital excitations.

As was shown in previous papers, the same holds for all quark generations (with fermionic substructure) but with the difference that already the orbital parts of the wave functions for the first generation must have nontrivial excitation states which allow the distinction between leptons and quarks and which are the origin of the color degrees of quarks [7, 24].

iii) A further point which must be emphasized is the difference between the electric Weinberg transformation and the magnetic one. In (II.46) and (II.47) both transformations were assumed to be identical, while in the present part the magnetic transformation was assumed to be just the dual of the electric one. The duality between electric fields and magnetic fields applies to all electroweak states and thus is a strong argument in favour of the definition in this part.

iv) In the Dirac theory of a magnetic monopol, the famous Dirac relation between the magnetic pole strength and the electric pole strength can be deduced, cf. [25] (9.34). Does this relation apply to the pole strengths of the lepton–boson couplings in (15) and (16) too? The answer is clearly:

The Dirac relation is not applicable because this relation depends on the presupposition of stable (conserved) electric as well as magnetic charges, cf. [25] Sect. 9.4, which is *not* satisfied in our model.

On the other hand, it should be possible to derive modified relations between electric and magnetic charges which, however, in contrast to the 'universal' Dirac relation will depend on the decay time of the magnetic monopoles and *are not universal*, but are connected with the experimental arrangement as was explained in Part II, Sect. 3.

[1] L. I. Urutskoev, Ann. Fond. L. de Broglie **29**, 1149 (2004).
[2] G. Lochak, L. Urutskoev, Low-Energy Nuclear Reactions and the Leptonic Monopole Condensed Matter Nuclear Science (Ed. J. P. Biberian), World Scientific, Singapore 2006, p. 421 – 437.
[3] G. Lochak, Z. Naturforsch. **62a**, 231 (2007).
[4] H. Stumpf, Z. Naturforsch. **66a**, 205 (2011).
[5] H. Stumpf, Z. Naturforsch. **66a**, 329 (2011).
[6] H. Stumpf, Z. Naturforsch. **55a**, 415 (2000).
[7] H. Stumpf, Z. Naturforsch. **59a**, 750 (2004).
[8] H. Stumpf, Z. Naturforsch. **37a**, 1295 (1982).
[9] T. Borne, G. Lochak, and H. Stumpf, Nonperturbative Quantum Field Theory and the Structure of Matter, Kluver Inc., Dordrecht 2001.
[10] H. Stumpf and T. Borne, Composite Particle Dynamics in Quantum Field Theory, Vieweg, Wiesbaden 1994.
[11] F. Gross, Relativistic Quantum Mechanics and Field Theory, Wiley Inc., New York 1999.
[12] L. Fonda and G. C. Ghirardi, Symmetry Principles in Quantum Physics, Dekker Inc., New York 1970.
[13] H. Stumpf, Z. Naturforsch. **58a**, 481 (2003).
[14] O. Nachtmann, Elementarteilchenphysik, Phänomene und Konzepte, Vieweg, Wiesbaden 1996.
[15] G. Grimm, Elektroschwache Standardtheorie als Effektive Dynamik von Bindungszuständen in Rahmen einer nichtlinearen Spinortheorie, Thesis, Tuebingen 1994.
[16] H. Stumpf, Z. Naturforsch. **63a**, 301 (2008).
[17] H. Stumpf, Z. Naturforsch. **60a**, 696 (2005).
[18] R. N. Mohapatra, Unification and Supersymmetry, 3rd Edn., Springer, Heidelberg 2003.
[19] M. Lindauer and C. Weinheimer, Phys. J. **10**, 31 (2011).
[20] K. Huang, Quarks, Leptons and Gauge Fields, 2nd Edn., World Scientific, Singapore 1992.
[21] H. Stumpf, Z. Naturforsch. **61a**, 439 (2006).
[22] H. V. Klapdor-Kleingrothaus and K. Zuber, Teilchenastrophysik, Teubner, Stuttgart 1997.
[23] G. Lochak, Int. J. Theor. Phys. **24**, 1019 (1985).
[24] H. Stumpf, Z. Naturforsch. **41a**, 1399 (1986).
[25] B. Felsager, Geometry, Particles and Fields, Springer, New York 1998.