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

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By Lochak (theory) and Urutskoev (experiment) the hypothesis has been suggested that during electric discharges in water (fluids) light magnetic monopoles can be created which according to Lochak should be considered as a kind of excited neutrinos. Based on a quantum field theoretic development of de Broglie's and Heisenberg's fusion ideas and the results of preceding papers a transparent proof is given that such magnetic monopoles can occur during discharges. In the theoretical description these circumstances are formulated within the scope of an extended (effective) Standard Model and the monopoles with vanishing electric charge arise from neutrinos whose states are modified by the symmetry breaking caused by the discharge. In the introduction some technical implications are referred to. The article is divided into two parts.

*Key words:* Parafermionic Boson and Lepton States; Leptonic Magnetic Monopoles.

## 1. Parafermionic Boson and Lepton States

To ensure transparency of the formalism a symbolic notation for the spinor field operators is used with $\psi_I := \psi_Z(x)$, where Z contains all algebraic indices of the field. Then in this representation a state $|a\rangle$ is characterized by the set of matrix elements

$$\tau_n(a) := \langle 0|\mathcal{A}(\psi_{I_1}\dots\psi_{I_n})|a\rangle, \quad n = 1\dots\infty, \quad (1)$$

where $\mathcal{A}$ means antisymmetrization in $I_1 \dots I_n$ (*only for conserved symmetries!*)

Within this formalism bound state equations can be defined for the set of matrix elements (1) for *fixed n*. These equations were introduced as generalized de Broglie-Bargmann-Wigner equations (GBBW equations) in previous papers. For their solutions as well as for corresponding test functions the symbols $C_{I_1\dots I_n}$ were defined by which the set of basis elements for the weak mapping method can be described in a simple way [Part I, 20].

Before going into details of the definition of wave functions, attention must be paid to two facts:

### *1.1. Representation with G-conjugated Fields*

The algebraic indices of the spinor field operator $\psi_I$ are defined by the four-dimensional indices $\kappa$ and $\alpha$ and the auxiliary indices $i$ for regularization. The four-dimensional index $\kappa$ can be splitted into the double index $\kappa = (B,b)$ and the superspinors are defined by

$$\psi_{\kappa\alpha i}(x) \equiv \psi_{Bb\alpha i}(x) = \begin{pmatrix} \psi_{b\alpha i}(x); B=1 \\ \psi^c_{b\alpha i}(x); B=2 \end{pmatrix}, \quad (2)$$

where $b$ denotes the isospin index.

Considering isospin transformations only, the spinor fields transform in isospace as

$$\psi' = \exp(-\mathrm{i}\varepsilon_k\sigma^k)\psi, \quad (3)$$

while the charge conjugated spinor fields transform in isospace according to

$$\psi^{c\prime} = \exp(\mathrm{i}\varepsilon_k\sigma^{*k})\psi^c. \quad (4)$$

Owing to the equivalence of the $\sigma$-algebra to the $\sigma^*$-algebra one can surmise that concerning the corresponding transformations the system contains a hidden symmetry. Indeed this symmetry can be realized by replacing the charge conjugated spinor fields by $G$-conjugated spinor fields. The latter are defined by [1]

$$\psi^G_{b\alpha i}(x) = c^{-1}_{bb'}\psi^c_{b'\alpha i}(x) \quad (5)$$

with $c := -\mathrm{i}\sigma_2$. One easily verifies that $\psi^G$ transforms according to (3) under isospin transformations.

Therefore for isospin transformations the fields $\psi$ and $\psi^G$ cannot be distinguished. In addition it can be shown that for Lorentz transformations the fields $\psi$ and $\psi^c$ as well as $\psi^G$ transform with the same transformation law [2, 3].

This fact allows to describe bound states by mixtures of these fields without destroying homogeneous transformation properties which are required for an appropriate physical interpretation of these states. In order to avoid a confusion of the index $G$ with magnetic fields, this kind of isospinors can be called a $D$-representation as in this representation the decomposition of wave functions into products between superspinors and spinors is possible.

### *1.2. Definition of Macro-Observables*

If by weak mapping macroscopic (effective) observables have to be derived, attention must be paid to the fact that exact solutions of the GBBW equations cannot be used as test functions themselves.

*If a composite particle is inserted into an assemblage of other particles, its internal structure must be adapted to the influence of this surrounding and the state of this particle can no longer be described by an eigensolution of the GBBW equations.*

Therefore in deriving effective theories, one is forced to consider test functions with freely variable parameters which can react to external forces.

In the case of CP-symmetry breaking for vector bosons exact state solutions of the GBBW equations were derived [I, 21]. Hence, if one uses approximations to simplify the calculation for corresponding test functions all group theoretical properties can be adopted from the exact solutions.

For the evaluation of the effective theory the single time wave functions are needed. To perform the transition to equal times we refer to [I, 17]. Owing to the translational invariance of the system we use the limit $t_1 = t_2 = 0$ without loss of generality. With $Z := (i, \alpha, \kappa)$ in this limit the wave functions of the vector bosons read with $(S^l + T^l)^D \equiv (S^l + T^l)^S$ (in this special case) where the $S$-label denotes ordinary superspin-isospin representation. Then one can define test functions by splitting the exact solutions into parts:

$$C^A_{Z_1Z_2}(\mathbf{r}_1, \mathbf{r}_2|\mathbf{k}, l, \mu) := (S^l + T^l)^S_{\kappa_1\kappa_2} e^{[-i\mathbf{k}\frac{1}{2}(\mathbf{r}_1+\mathbf{r}_2)]} \cdot (\gamma^\mu C)_{\alpha_1\alpha_2} f^A(\mathbf{r}_1 - \mathbf{r}_2|\mathbf{k}, l, \mu)_{i_1i_2},$$

$$C^G_{Z_1Z_2}(\mathbf{r}_1, \mathbf{r}_2|\mathbf{k}, l, \mu) := (S^l + T^l)^S_{\kappa_1\kappa_2} e^{[-i\mathbf{k}\frac{1}{2}(\mathbf{r}_1+\mathbf{r}_2)]} \cdot (\gamma^5\gamma^\mu C)_{\alpha_1\alpha_2} f^G(\mathbf{r}_1 - \mathbf{r}_2|\mathbf{k}, l, \mu)_{i_1i_2}$$

$$C^F_{Z_1Z_2}(\mathbf{r}_1, \mathbf{r}_2|\mathbf{k}, l, \mu, \nu) := (S^l + T^l)^S_{\kappa_1\kappa_2} e^{[-i\mathbf{k}\frac{1}{2}(\mathbf{r}_1+\mathbf{r}_2)]} \cdot (\Sigma^{\mu\nu} C)_{\alpha_1\alpha_2} f^F(\mathbf{r}_1 - \mathbf{r}_2|\mathbf{k}, l, \mu, \nu)_{i_1i_2}, \quad (6)$$

where $A$ means the electric vector potential, $G$ the magnetic axial vector, and $F$ the common field tensor. The influence of the surrounding can be expressed by variable coefficients which later on are to be identified with the effective field variables.

### *1.3. Lepton States for Broken Symmetry*

Concerning the fermion states, their group theoretical analysis has been performed in several papers [4 – 8].

For conserved symmetries the permutation group representations play an essential role in the construction of appropriate wave functions. Being based on the theory of representations of the permutation group elaborated by Kramer et al. [9], they guarantee the complete antisymmetrization in the basic spinor field quantities.

This can only be achieved by using mixed representations of the permutation group. Such mixed representations are generated by the application of Young operators. For two-dimensional representations these operators can be found in [6, 8, I, 32].

The quantum numbers of these states coincide with the phenomenological quantum numbers where the last column in [I, 32], eq. (37), corresponds to the phenomenological spinor fields $\psi^G$ afterwards.

Based on these sets the effective coupling of leptons to electroweak bosons was calculated in previous papers [I, 25; I, 32; I, 33].

*But attention must be paid to the fact that in these calculations the boson states (6) are referred to broken CP-symmetry, while according to their construction the lepton states are referred to conserved symmetries.*

*Therefore for being free from contradiction in analogy to the parafermionic boson states (6) also the lepton states must contain parafermionic elements in order to be adapted to the broken CP-symmetry.*

*For the constructions of such parafermionic lepton states holds: Any superspin-isospin symmetry breaking states must still allow to identify neutrinos.*

While for conserved symmetries in [6 – 8] GBBW equations were analysed which were invariant under the permutation group, it is obvious that for symmetry breaking this type of equations cannot be used for the construction of parafermionic states. In [4] such asymmetric GBBW equations were discussed.

The corresponding lepton states are then products between superspin-isospin states and spin-orbit states. According to [I, 17] such states read

$$C^{\kappa_1\kappa_2\kappa_3}_{\alpha_1\alpha_2\alpha_3}(\mathbf{r}_1,\mathbf{r}_2,\mathbf{r}_3|\mathbf{k},j,n) := \exp\left[-\mathrm{i}\mathbf{k}\frac{1}{3}(\mathbf{r}_1+\mathbf{r}_2+\mathbf{r}_3)\right] \cdot [\Theta^j_{\kappa_1\kappa_2\kappa_3}\Omega^n_{\alpha_1\alpha_2\alpha_3}\psi(\mathbf{r}_2-\mathbf{r}_1,\mathbf{r}_3-\mathbf{r}_2|\mathbf{k})]. \quad (7)$$

For the spin tensor $\Omega^n$ we apply lepton fields $l^j_\alpha(x)$ which are *not* eigenstates of the Dirac operator for a definite **k**-vector. Furthermore, as the leptons are assumed to occupy the ground states of the three-parton system, the spin tensor as well as the orbit functions must show the highest possible invariance under symmetry operations, which for these parts of the wave functions are the little group operations with all discrete transformations. This leads to the spin tensor and its charge conjugated counterpart

$$\Omega^n_{\alpha_1\alpha_2\alpha_3} = C_{\alpha_1\alpha_2}\xi^n_{\alpha_3}, \quad \bar{\Omega}^n_{\alpha_1\alpha_2\alpha_3} = C_{\alpha_1\alpha_2}C_{\alpha_3\alpha}\xi^n_\alpha, \quad (8)$$

where $\xi^n_\alpha$ are the four unit spinors $\delta_{\alpha n}$, $n=1,2,3,4$, while $C$ is invariant under rotations and the discrete operation PC [10]. The orbit part is assumed to have s-wave character which automatically is invariant under parity transformations.

In [I, 17] a set of superspin-isospin states was given which respects the identification of leptons by quantum numbers, but violates the complete antisymmetrization in (7). Hence, it has to be analysed whether these states are suitable canditates for parafermionic representations.

The corresponding superspin-isospin tensor for neutrinos is explicitly given in [I, 17] by

$$\Theta^2_{\kappa_1,\kappa_2,\kappa_3} = 3^{-1/2}\Big[\delta_{4,\kappa_1}\delta_{4,\kappa_2}\delta_{3,\kappa_3} + \delta_{4,\kappa_1}\delta_{3,\kappa_2}\delta_{4,\kappa_3} + \delta_{3,\kappa_1}\delta_{4,\kappa_2}\delta_{4,\kappa_3}\Big], \quad (9)$$

where the index 2 serves for the identification of the neutrino state in [I, 17]. The complete list of superspin-isospin states in [I, 17] shows that abandoning the Young construction leads to higher isospin states and higher charge states which at present have not been observed so far. It was shown in [6, 7] that for superspin-isospin states with permutation symmetry the ansatz (7) allows an exact solution of the symmetrical as well as the asymmetrical GBBW equations.

For instance with $\hat{\varphi}^{\kappa_1\kappa_2\kappa_3}_{\alpha_1\alpha_2\alpha_3} = \Theta^l_{\kappa_1\kappa_2\kappa_3}\hat{\varphi}_{\alpha_1\alpha_2\alpha_3}(x_1x_2x_3)$ from the asymmetric equations in [4] the following equation results:

$$\begin{aligned}
\Theta^l_{\kappa_1\kappa_2\kappa_3}\hat{\varphi}_{\alpha_1\alpha_2\alpha_3}(x_1,x_2,x_3) = \int \mathrm{d}x \sum_i \lambda_i G_{\alpha_3\alpha'_1}(x_3-x,m_i)\\
\cdot \sum_h 6\Big\{ -v^h_{\alpha'_1\beta'}(v^hC)_{\beta\beta''}\Theta^l_{\kappa_1\kappa_3\kappa_2}\\
\cdot \Big[-\sum_j \lambda_j(-i)F_{\beta\alpha_2}(x-x_2,m_j)\hat{\varphi}_{\alpha_1\beta'\beta''}(x_1,x,x)\Big]\\
-v^h_{\alpha'_1\beta'}(v^hC)_{\beta\beta''}\Theta^l_{\kappa_2\kappa_3\kappa_1}\\
\cdot \Big[-\sum_j \lambda_j(-i)F_{\beta\alpha_1}(x-x_1,m_j)\hat{\varphi}_{\alpha_2\beta'\beta''}(x_2,x,x)\Big]\\
-v^h_{\alpha'_1\beta''}(v^hC)_{\beta'\beta}\Theta^l_{\kappa_1\kappa_2\kappa_3}\\
\cdot \Big[-\sum_j \lambda_j(-i)F_{\beta\alpha_2}(x-x_2,m_j)\hat{\varphi}_{\alpha_1\beta'\beta''}(x_1,x,x)\Big]\\
-v^h_{\alpha'_1\beta''}(v^hC)_{\beta'\beta}\Theta^l_{\kappa_2\kappa_1\kappa_3}\\
\cdot \Big[-\sum_j \lambda_j(-i)F_{\beta\alpha_1}(x-x_1,m_j)\hat{\varphi}_{\alpha_2\beta'\beta''}(x_2,x,x)\Big]\Big\},
\end{aligned} \quad (10)$$

where for short the symmetry breaking part of the propagator has been omitted. One easily verifies that with (9) owing to its permutation invariance, the superspin-isospin part of (7) can be eliminated from (10).

*But although the Young-construction is avoided by (7), the states (9) cannot be the correct description of the CP-symmetry breaking because in this case the exact superspin-isospin boson states are neither symmetric nor antisymmetric. Therefore, in analogy to the boson states, for CP-symmetry breaking the superspin-isospin lepton states ought not have a permutation symmetry like (9).*

A detailed information about the consequences of this insight will be given in Section 3.

## 2. How Magnetic Monopoles are Linked to Discharges

The crucial formula which decides whether magnetic monopoles do exist follows from the effective lepton-boson coupling term given in [I, 32] and also [I, 20]:

$$\mathcal{H}^1_{bf} = 3W_{I_1I_2I_3I_4}R^q_{II'I_1}C^p_{II'I_4}C^l_{I_2I_3}f_q\partial^b_l\partial^f_p. \quad (11)$$

This term can be evaluated under the assumption that all wave functions are referred to broken CP-symmetry.

In the first step the expression $W_{I_1I_2I_3I_4}C^k_{I_2I_3}\partial^b_k$ has to be calculated. If $C^k_{I_2I_3}$ is projected on $W$ the third term of the algebraic part of the vertex drops out and the same holds for the terms connected with $\partial^E$ and $\partial^B$ (for $W$ cf. [I, 17]). Then one obtains

$$
\begin{aligned}
&W_{I_1I_2I_3I_4}C^k_{I_2I_3}\partial^b_k = \sum_{\mathbf{k}}\sum_{I_2I_3}\lambda_{i_1}B_{i_2i_3i_4}\iint \mathrm{d}^3r_2\mathrm{d}^3r_3\\
&\cdot\delta(\mathbf{r}_1-\mathbf{r}_2)\delta(\mathbf{r}_1-\mathbf{r}_3)\delta(\mathbf{r}_1-\mathbf{r}_4)\sum_h[(\gamma^0v^h)_{\beta_1\beta_2}\\
&\cdot(v^hC)_{\beta_3\beta_4}\delta_{\rho_1\rho_2}\gamma^5_{\rho_3\rho_4}-(\gamma^0v^h)_{\beta_1\beta_3}(v^hC)_{\beta_2\beta_4}\delta_{\rho_1\rho_3}\gamma^5_{\rho_2\rho_4}]\\
&\cdot(T^b+S^b)_{\rho_2\rho_3}[f^A_{i_2i_3}(\mathbf{r}_2-\mathbf{r}_3|\mathbf{k})(\gamma^kC)_{\beta_2\beta_3}\partial^A_{k,b}(\mathbf{k})\\
&+f^G_{i_2i_3}(\mathbf{r}_2-\mathbf{r}_3|\mathbf{k})(\gamma^k\gamma^5C)_{\beta_2\beta_3}\partial^G_{k,b}(\mathbf{k})]\\
&\cdot\exp\left[-\mathrm{i}\mathbf{k}\frac{1}{2}(\mathbf{r}_2+\mathbf{r}_3)\right] \qquad (12)
\end{aligned}
$$

which leads to

$$
\begin{aligned}
&W_{I_1I_2I_3I_4}C^k_{I_2I_3}\partial^b_k = -4[(\gamma^0\gamma^k)_{\beta_1\beta_4}(T^b\gamma^5)_{\rho_1\rho_4}\hat{f}^A(\mathbf{k})\\
&\cdot\partial^A_{k,b}(\mathbf{k})+(\gamma^0\gamma^k\gamma^5)_{\beta_1\beta_4}(S^b\gamma^5)_{\rho_1\rho_4}\hat{f}^G(\mathbf{k})\partial^G_{k,b}(\mathbf{k})]\lambda_{i_1}\\
&\cdot B_{i_4}\delta(\mathbf{r}_1-\mathbf{r}_4)\exp[-\mathrm{i}\mathbf{k}\mathbf{r}_1], \qquad (13)
\end{aligned}
$$

where $\hat{f}^A$ and $\hat{f}^G$ are the values of the corresponding boson functions (6) at the origin. Furthermore, the variable coefficients which are assumed to be the effective field variables are hidden in the state functional for composite particles $|\mathcal{P}(b,f)\rangle$ [I, 32].

If one substitutes the wave functions $C_{I_1I_2I_3}$ together with the dual fermion functions $R_{I_1I_2I_3}$ [5] into (11), integrates over $\mathbf{r}_4$ and renames $\mathbf{r}_1$ to $\mathbf{r}''$, then one gets with (13)

$$
\begin{aligned}
&\mathcal{H}^1_{bf} = -12\int\cdots\int \mathrm{d}^3k\mathrm{d}^3k'\mathrm{d}^3k''\mathrm{d}^3r\mathrm{d}^3r'\mathrm{d}^3r''\\
&\cdot[(\gamma^0\gamma^k)_{\beta_1\beta_4}(T^b\gamma^5)_{\rho_1\rho_4}\hat{f}^A(\mathbf{k})\partial^A_{k,b}(\mathbf{k})+(\gamma^0\gamma^k\gamma^5)_{\beta_1\beta_4}\\
&\cdot(S^b\gamma^5)_{\rho_1\rho_4}\hat{f}^G(\mathbf{k})\partial^G_{k,b}(\mathbf{k})]\sum_{i_1i_4i}\lambda_{i_1}R^{\rho\rho'\rho_1}_{\beta\beta'\beta_1}(\mathbf{r},\mathbf{r}',\mathbf{r}''|\mathbf{k}',l,n)_{ii'i_1}\\
&\cdot\exp\left[\mathrm{i}\mathbf{k}'\frac{1}{3}(\mathbf{r}+\mathbf{r}'+\mathbf{r}'')\right]C^{\rho_4\rho\rho'}_{\beta_4\beta\beta'}(\mathbf{r}'',\mathbf{r},\mathbf{r}'|\mathbf{k}'',j,m)_{ii'i_4}\\
&\cdot\exp\left[-\mathrm{i}\mathbf{k}''\frac{1}{3}(\mathbf{r}+\mathbf{r}'+\mathbf{r}'')\right]\\
&\cdot\exp\left[-\mathrm{i}\mathbf{k}\mathbf{r}''\right]f_{ln}(\mathbf{k}')\partial^f_{jm}(\mathbf{k}'') \qquad (14)
\end{aligned}
$$

with $l$, $j$ as superspin-isospin state numbers, and $n$, $m$ as spin state numbers. Note that in (11) the summation convention has been used which in (12) and (14) is explicitly expressed by integrations!

Introduction of center of mass coordinates

$$
z=\frac{1}{3}(\mathbf{r}+\mathbf{r}'+\mathbf{r}''); \quad \mathbf{u}=\mathbf{r}'-\mathbf{r}; \quad \mathbf{v}=\mathbf{r}''-\mathbf{r}' \quad (15)
$$

and

$$
\begin{aligned}
&\mathbf{r}=\mathbf{z}-\frac{2}{3}\mathbf{u}-\frac{1}{3}\mathbf{v}; \quad \mathbf{r}'=\mathbf{z}+\frac{1}{3}\mathbf{u}-\frac{1}{3}\mathbf{v};\\
&\mathbf{r}''=\mathbf{z}+\frac{1}{3}\mathbf{u}+\frac{2}{3}\mathbf{v}
\end{aligned} \qquad (16)
$$

yields

$$
\begin{aligned}
&\mathcal{H}^1_{bf} = -12\int\cdots\int \mathrm{d}^3k\mathrm{d}^3k'\mathrm{d}^3k''\mathrm{d}^3z\mathrm{d}^3u\mathrm{d}^3v\\
&\cdot[(\gamma^0\gamma^k)_{\beta_1\beta_4}(T^b\gamma^5)_{\rho_1\rho_4}\hat{f}^A(\mathbf{k})\partial^A_{k,b}(\mathbf{k})+(\gamma^0\gamma^k\gamma^5)_{\beta_1\beta_4}\\
&\cdot(S^b\gamma^5)_{\rho_1\rho_4}\hat{f}^G(\mathbf{k})\partial^G_{k,b}(\mathbf{k})]\sum_{i_1i_4i}\lambda_{i_1}R^{\rho\rho'\rho_1}_{\beta\beta'\beta_1}(\mathbf{u},\mathbf{v}|\mathbf{k}',l,n)_{ii'i_1}\\
&\cdot C^{\rho\rho'\rho_4}_{\beta\beta'\beta_4}(\mathbf{u},\mathbf{v}|\mathbf{k}'',j,m)_{ii'i_4}f_{ln}(\mathbf{k}')\partial^f_{jm}(\mathbf{k}'')\exp(\mathrm{i}\mathbf{k}'\mathbf{z})\\
&\cdot\exp(-\mathrm{i}\mathbf{k}''\mathbf{z})\exp\left[-\mathrm{i}\mathbf{k}\left(\mathbf{z}+\frac{1}{3}\mathbf{u}+\frac{2}{3}\mathbf{v}\right)\right]. \qquad (17)
\end{aligned}
$$

The further evaluation depends upon the fermionic wave functions (7) and (9). In particular with (7) and with

$$
\begin{aligned}
&Y^s\in\{(T^a\gamma^5),(S^a\gamma^5), \quad a=0,1,2,3\},\\
&X^t\in\{(\gamma^0\gamma^k),(\gamma^0\gamma^k\gamma^5), \quad k=1,2,3\}
\end{aligned} \qquad (18)
$$

the parts containing the wave functions in (17) can be written as follows:

$$
\begin{aligned}
&R^{\rho\rho'\rho_1}_{\beta\beta'\beta_1}(\mathbf{u},\mathbf{v}|ln)Y^s_{\rho_1\rho_4}X^t_{\beta_1\beta_4}C^{\rho\rho'\rho_4}_{\beta\beta'\beta_4}(\mathbf{u},\mathbf{v}|jm)\equiv\\
&\sum_{i_1i_2i_3i_4}\{(\Theta^l)_{\rho\rho'\rho_1}[\Omega^n_{\beta\beta'\beta_1}\psi^*(\mathbf{u},\mathbf{v})]\}_{i_1i_2i_3}Y^s_{\rho_1\rho_4}X^t_{\beta_1\beta_4} \quad (19)\\
&\qquad\cdot\{(\Theta^j)_{\rho\rho'\rho_4}[\Omega^m_{\beta\beta'\beta_4}\psi(\mathbf{u},\mathbf{v})]\}_{i_1i_2i_4},
\end{aligned}
$$

where the summation over the auxiliary indices is explicitly indicated. The special form of the wave functions in (19) allows the formal definition

$$
\langle\Theta^l|Y^s_{(3)}|\Theta^j\rangle =: Y^s_{lj} \quad \forall s, \qquad (20)
$$

and as the broken symmetry manifests itself mainly in the superspin-isospin part (9), we adopt the exact spin formulas for conserved symmetry [I, 32] as

$$
\begin{aligned}
&\langle C_{12}\bar{\Omega}^n\psi|X^t_{(3)}|C_{12}\bar{\Omega}^m\psi\rangle =\\
&\xi^n_\alpha X^t_{\alpha\beta}\xi^m_\beta\Upsilon(\mathbf{u},\mathbf{v},\mathbf{k}',\mathbf{k}'',l,n,j,m).
\end{aligned} \qquad (21)
$$

With $\xi^n_\alpha X^t_{\alpha\beta}\xi^m_\beta = X^t_{nm}$ this gives for (17)

$$\mathcal{H}^1_{bf} = -12\int\cdots\int \mathrm{d}^3k\mathrm{d}^3k'\mathrm{d}^3k''\mathrm{d}^3z\mathrm{d}^3u\mathrm{d}^3v \cdot[(\gamma^0\gamma^k)_{nm}(T^b\gamma^5)_{lj}\hat{f}^A(\mathbf{k})\partial^A_{kb}(\mathbf{k}) + (\gamma^0\gamma^k\gamma^5)_{nm}(S^b\gamma^5)_{lj} \cdot \hat{f}^G(\mathbf{k})\partial^G_{kb}(\mathbf{k})]\chi(\mathbf{u},\mathbf{v}|\mathbf{k}',\mathbf{k}'',n,m,l,j)f_{ln}(\mathbf{k}')\partial^f_{jm}(\mathbf{k}'') \cdot \exp(\mathrm{i}\mathbf{k}'\mathbf{z})\exp(-\mathrm{i}\mathbf{k}''\mathbf{z})\exp\left[-\mathrm{i}\mathbf{k}\left(\mathbf{z}+\frac{1}{3}\mathbf{u}+\frac{2}{3}\mathbf{v}\right)\right]. \quad (22)$$

Fourier transformation of the functional operators $\partial^A(\mathbf{k})$, $\partial^G(\mathbf{k})$, $f(\mathbf{k}')$, $\partial^f(\mathbf{k}'')$ yields for (22)

$$\mathcal{H}^1_{bf} = -12\int\cdots\int \mathrm{d}^3k\mathrm{d}^3k'\mathrm{d}^3k''\mathrm{d}^3z\mathrm{d}^3u\mathrm{d}^3v\mathrm{d}^3x\mathrm{d}^3p\mathrm{d}^3y \cdot[(\gamma^0\gamma^k)_{nm}(T^b\gamma^5)_{lj}\hat{f}^A(\mathbf{k})\tilde{\partial}^A_{kb}(\mathbf{x}) + (\gamma^0\gamma^k\gamma^5)_{nm}(S^b\gamma^5)_{lj} \cdot \hat{f}^G(\mathbf{k})\tilde{\partial}^G_{kb}(\mathbf{x})]\chi(\mathbf{u},\mathbf{v}|\mathbf{k}',\mathbf{k}'',n,m,l,j)\tilde{f}_{ln}(\mathbf{y})\tilde{\partial}^f_{jm}(\mathbf{p}) \cdot \exp\left[\mathrm{i}\mathbf{k}'(\mathbf{z}+\mathbf{y})\exp\left[-\mathrm{i}\mathbf{k}''(\mathbf{z}+\mathbf{p})\right]\right] \cdot \exp\left[-\mathrm{i}\mathbf{k}\left(\mathbf{z}+\frac{1}{3}\mathbf{u}+\frac{2}{3}\mathbf{v}+\mathbf{x}\right)\right], \quad (23)$$

where the transformed functional operators are denoted by tilde.

A reduction of this expression can be achieved if assumptions about the form factors are made.

By careful calculations it was demonstrated in [11] that for conserved symmetries the dependence of $\hat{f}^A(\mathbf{k})$ on $\mathbf{k}$ drops out. If one transfers this to the symmetry breaking boson functions and extends this to the magnetic boson value at the origin, a further simplification of (23) can be achieved.

*But* before proceeding further the formalism must be adapted to the concept of excited neutrinos.

*Following the idea of excited neutrinos as magnetic monopoles, from the point of view of weak mapping theorems, the excited neutrino states must be introduced as new additional basis states, i. e. particles, into the theory*.

We introduce functional sources $f^*(\mathbf{k})$ and their duals $\partial^*(\mathbf{k})$ for the excited neutrinos. In analogy to ordinary leptons the latter must possess a functional Dirac Hamiltonian. This gives

$$\mathcal{H}^*_f = \int \mathrm{d}^3z f^*_\alpha(\mathbf{z})[-\mathrm{i}(\gamma^0\gamma^k)\partial^z_k + m^*\gamma^0]_{\alpha\beta}\partial^{f*}_\beta(\mathbf{z}). \quad (24)$$

Furthermore, we need the coupling of the excited neutrinos to the boson fields. The general formula for this lepton-boson coupling is given by (11). The symbolic state numbers and corresponding quantum numbers $q,p,k$ run over all relevant particle states of the system including the hypothetical excited neutrino states.

A characteristic property of (11) is that this equation in principle any $q$ can interact with any $p$ and any $k$. Really, this enormous assemplage of interaction terms can be reduced by appropriate evaluation, for instance, to the interactions terms in (23). In addition, searching for monopoles one can suppress the 1,2-charged vector fields in (23).

To separate the excited neutrino states from the other leptonic states we postulate:

**Postulate**: *The excited neutrino states are orthogonal to the lepton ground states*

By this obvious postulate no mixture between ground states and excited states can occur and if one simply replaces the functional sources for ordinary leptons by the star functionals, then from (23) one obtains

$$\mathcal{H}^{1*}_{bf} = -12c^b\int\cdots\int \mathrm{d}^3k\mathrm{d}^3k'\mathrm{d}^3k''\mathrm{d}^3z\mathrm{d}^3u\mathrm{d}^3v\mathrm{d}^3x\mathrm{d}^3y\mathrm{d}^3p \cdot[\xi^n_\alpha(\gamma^0\gamma^k)_{\alpha\beta}\xi^m_\beta(T^b\gamma^5)_{lj}\hat{f}^A(0)\tilde{\partial}^A_{kb}(\mathbf{x}) + \xi^n_\alpha(\gamma^0\gamma^k\gamma^5)_{\alpha\beta} \cdot \xi^m_\beta(S^b\gamma^5)_{lj}\hat{f}^G(0)\tilde{\partial}^G_{kb}(\mathbf{x})]\Upsilon(\mathbf{u},\mathbf{v}|\mathbf{k}',\mathbf{k}'',n,m,l,j)\tilde{f}^*_{ln}(\mathbf{y}) \cdot \tilde{\partial}^{f*}_{jm}(\mathbf{p})\exp[\mathrm{i}\mathbf{k}'(\mathbf{z}+\mathbf{y})]\exp[-\mathrm{i}\mathbf{k}''(\mathbf{z}+\mathbf{p})] \cdot \exp\left[-\mathrm{i}\mathbf{k}\left(\mathbf{z}+\tfrac{1}{3}\mathbf{u}+\tfrac{2}{3}\mathbf{v}+\mathbf{x}\right)\right]. \quad (25)$$

*In this context note that according to (20), (21) the indices n, l, j, m represent the quantum numbers of the states involved which should not be confused with the ordinary algebraic indices of spin, superspin, and isospin degrees, denoted by Greek letters*.

To emphasize this difference we introduce, in addition to the spin unit spinors $\xi^n_\alpha$ in (25), unit spinors $\zeta^l_\rho$ in superspin-isospin space. The latter are a consequence of the wave functions for broken symmetry and result from an evaluation of formula (19) and definition (20). With

$$\zeta^l_\rho \otimes \zeta^j_{\rho'} := \Theta^l_{\kappa,\kappa',\rho}\Theta^j_{\kappa,\kappa',\rho'} \quad (26)$$

one gets for (20)

$$(T^b\gamma^5)_{lj} \equiv \zeta^l_\rho(T^b\gamma^5)_{\rho\rho'}\zeta^j_{\rho'}, \qquad (S^b\gamma^5)_{lj} \equiv \zeta^l_\rho(S^b\gamma^5)_{\rho\rho'}\zeta^j_{\rho'}. \quad (27)$$

Now attention must be paid to the fact that with (25) the coupling of one and the same particle to the boson fields has to be examined. Hence in (25) the 'ingoing' particles characterized by the functional operator $\tilde{f}^*_{ln}$ and the 'outgoing' particles characterized by the operator $\tilde{\partial}^{f*}_{jm}$ must be identical. This means that the superspin-isospin quantum numbers as well as the spin quantum numbers of these particles must be the same.

In particular for the superspin-isospin states of the excited neutrino, one obtains with (9) the expression

$$\zeta_\rho(\nu^*) \otimes \zeta_{\rho'}(\nu^*) := \Theta^2_{\kappa,\kappa',\rho}\Theta^2_{\kappa,\kappa',\rho'}. \quad (28)$$

Suppressing a possible dependence of $\Upsilon$ on the quantum numbers, one can eliminate the unit spinors in (25) leading to the replacement $n \to \alpha$ and $m \to \beta$ and $f^*_{ln} \to f^*_{2\alpha}$, $\partial^{f*}_{jm} \to \partial^{f*}_{2\beta}$.

With the transformation to the new variable $\mathbf{s} = \mathbf{z} + 1/3\mathbf{u} + 2/3\mathbf{v}$ and the above replacements (25) reads

$$\begin{aligned}\mathcal{H}^{1*}_{bf} =& -12c^b \int \cdots \int \mathrm{d}^3k\mathrm{d}^3k'\mathrm{d}^3k''\mathrm{d}^3s\mathrm{d}^3u\mathrm{d}^3v\mathrm{d}^3x\mathrm{d}^3y\mathrm{d}^3p \\ &\cdot [(\gamma^0\gamma^k)_{\alpha\beta}\zeta(\nu^*)_\rho(T^b\gamma^5)_{\rho\rho'}\zeta(\nu^*)_{\rho'}\hat{f}^A(0)\tilde{\partial}^A_{kb}(\mathbf{x}) \\ &+ (\gamma^0\gamma^k\gamma^5)_{\alpha\beta}\zeta(\nu^*)_\rho(S^b\gamma^5)_{\rho\rho'}\zeta(\nu^*)_{\rho'}\hat{f}^G(0)\tilde{\partial}^G_{kb}(\mathbf{x})] \\ &\cdot \Upsilon(\mathbf{u},\mathbf{v})\tilde{f}^*_{2\alpha}(\mathbf{y})\tilde{\partial}^{f*}_{2\beta}(\mathbf{p})\exp\Big[\mathrm{i}\mathbf{k}'\Big(\mathbf{s} - \frac{1}{3}\mathbf{u} - \frac{2}{3}\mathbf{v} + \mathbf{y}\Big)\Big] \\ &\cdot \exp\Big[-\mathrm{i}\mathbf{k}''\Big(\mathbf{s} - \frac{1}{3}\mathbf{u} - \frac{2}{3}\mathbf{v} + \mathbf{p}\Big)\Big]\exp[-\mathrm{i}\mathbf{k}(\mathbf{s}+\mathbf{x})]. \quad (29)\end{aligned}$$

The further evaluation depends on an information about the neutrino wave functions. The three-body GBBW-bound state equations of the (composite) neutrino were discussed in [4, 5, 8]. As the exact neutrino wave functions are unknown, we use test functions to describe a possible neutrino structure which should be in accordance with the properties of the exact solutions of the GBBW equations. The following results were obtained:

(i) The product wave function (7) is a compatible solution of the GBBW equations if the superspin-isospin part of (9) satisfies

$$\gamma^5_{\kappa_2\kappa_3}\Theta^2_{\kappa_1\kappa_2\kappa_3} = 0. \quad (30)$$

(ii) The condition (30) is satisfied by construction and holds for conserved symmetries as well as for CP-violation. In the latter case (9) and (7) can still be applied, but for CP-violation the GBBW equations themselves are modified. Obviously, if use is made of such an ansatz in any case a special physical interpretation of the resulting set of states is required.

(iii) For conserved symmetries the ground state wave functions can be constructed by products of mixed representations of superspin-isospin states as well as of orbital and spin states for the little group and the permutation group. This leads to an orbital state which is a completely symmetric s-state under permutations of the Cartesian coordinates [8].

In view of these conditions the following states for orbital test functions can be defined where it has to be noted that the function $\Upsilon(\mathbf{u},\mathbf{v})$ contains densities of wave functions with respect to $\mathbf{u}$ and $\mathbf{v}$.

$$\Upsilon(\mathbf{u},\mathbf{v}) = \mathrm{e}^{-a\mathbf{u}^2}\Big(\frac{a}{\pi}\Big)^{3/2}\mathrm{e}^{-a\mathbf{v}^2}\Big(\frac{a}{\pi}\Big)^{3/2} \quad (31)$$

with $a \gg 1$ but leave it open to find a meaningful value later on. Then the integrals

$$\begin{aligned}&\int \mathrm{d}^3u \exp\Big[-\mathrm{i}(\mathbf{k}'-\mathbf{k}'')\frac{1}{3}\mathbf{u} - a\mathbf{u}^2\Big]\int \mathrm{d}^3v \\ &\cdot \exp\Big[-\mathrm{i}(\mathbf{k}'-\mathbf{k}'')\frac{2}{3}\mathbf{v} - a\mathbf{v}^2\Big]\Big(\frac{a}{\pi}\Big)^3 \quad (32) \\ &= \exp\Big[-\frac{5}{9a}(\mathbf{k}'-\mathbf{k}'')^2\Big]\end{aligned}$$

can be substituted in (29). After integration over $\mathbf{k}$ and $\mathbf{s}$ formula (29) reads

$$\begin{aligned}\mathcal{H}^{1*}_{bf} =& -12c^b \int \cdots \int \mathrm{d}^3k'\mathrm{d}^3k''\mathrm{d}^3x\mathrm{d}^3y\mathrm{d}^3p \\ &\cdot [(\gamma^0\gamma^k)_{\alpha\beta}\zeta(\nu^*)_\rho(T^b\gamma^5)_{\rho\rho'})\zeta(\nu^*)_{\rho'}\hat{f}^A(0)\tilde{\partial}^A_{kb}(\mathbf{x}) \\ &+ (\gamma^0\gamma^k\gamma^5)_{\alpha\beta}\zeta(\nu^*)_\rho(S^b\gamma^5)_{\rho\rho'}\zeta(\nu^*)_{\rho'}\hat{f}^G(0)\tilde{\partial}^G_{kb}(\mathbf{x})] \\ &\cdot \tilde{f}^*_{2\alpha}(\mathbf{y})\tilde{\partial}^{f*}_{2\beta}(\mathbf{p})\exp[\mathrm{i}\mathbf{k}'(\mathbf{x}+\mathbf{y})]\exp[-\mathrm{i}\mathbf{k}''(\mathbf{x}+\mathbf{p})] \\ &\cdot \exp\Big[-\frac{5}{9a}(\mathbf{k}'-\mathbf{k}'')^2\Big]. \quad (33)\end{aligned}$$

By means of the transformation $\mathbf{k}' = \mathbf{v} + \mathbf{h}$ and $\mathbf{k}'' = \mathbf{v}$ (33) can be changed into a form which allows exact integrations. Furthermore, expression (33) is invariant under the replacement $\mathbf{x}$ by $-\mathbf{x}$. After these operations (33) passes into

$$\begin{aligned}\mathcal{H}^{1*}_{bf} =& -12c^b \iint \mathrm{d}^3x\mathrm{d}^3y[(\gamma^0\gamma^k)_{\alpha\beta}\zeta(\nu^*)_\rho(T^b\gamma^5)_{\rho\rho'} \\ &\cdot \zeta(\nu^*)_{\rho'}\hat{f}^A(0)\tilde{\partial}^A_{kb}(\mathbf{x}) + (\gamma^0\gamma^k\gamma^5)_{\alpha\beta} \\ &\cdot \zeta(\nu^*)_\rho(S^b\gamma^5)_{\rho\rho'}\zeta(\nu^*)_{\rho'}\hat{f}^G(0)\tilde{\partial}^G_{kb}(\mathbf{x})]\tilde{f}^*_{2\alpha}(\mathbf{y})\tilde{\partial}^{f*}_{2\beta}(\mathbf{y}) \\ &\cdot \exp\Big[-\frac{1}{2}(\mathbf{y}-\mathbf{x})^2\Big]\Big(\frac{\pi 9a}{5}\Big)^{1/2} \quad (34)\end{aligned}$$

which approximately yields after integration over $\mathbf{y}$

$$\mathcal{H}^{1*}_{bf} = -12c^b \int \mathrm{d}^3x \Big[(\gamma^0\gamma^k)_{\alpha\beta}\zeta(v^*)_\rho (T^b\gamma^5)_{\rho\rho'} \cdot \zeta(v^*)_{\rho'} \hat{f}^A(0)\tilde{\partial}^A_{kb}(\mathbf{x}) + (\gamma^0\gamma^k\gamma^5)_{\alpha\beta}\zeta(v^*)_\rho (S^b\gamma^5)_{\rho\rho'} \cdot \zeta(v^*)_{\rho'} \hat{f}^G(0)\tilde{\partial}^G_{kb}(\mathbf{x})\Big] \tilde{f}^*_{2\alpha}(\mathbf{x})\tilde{\partial}^{f*}_{2\beta}(\mathbf{x})\left(\frac{2\pi^2 9a}{5}\right)^{1/2}. \quad (35)$$

In analogy to the derivation of classical equations from an effective functional equation [I, 17] from (24) and (35) an effective Dirac equation can be derived. In this equation the coupling to the charged vector bosons $b = 1,2$ is not relevant to the monopol problem. Thus these parts will be omitted. Then one gets for the effective spinor amplitude of the excited neutrino $\psi(x)$ the reduced equation

$$\begin{aligned} \mathrm{i}\partial_t \psi_{2\alpha}(x) &= [-\mathrm{i}(\gamma^0\gamma^k)_{\alpha\beta}\partial_k + \gamma^0_{\alpha\beta} m^*]\psi_{2\beta}(x) \\ &+ \{g^*_A(\gamma^0\gamma^k)_{\alpha\beta}\zeta_\rho(v^*)(T^0\gamma^5)^D_{\rho\rho'}\zeta_{\rho'}(v^*)A^0_k(x) \\ &- g'^*_A(\gamma^0\gamma^k)_{\alpha\beta}\zeta_\rho(v^*)(T^3\gamma^5)^D_{\rho\rho'}\zeta_{\rho'}(v^*)A^3_k(x)\}\psi_{2\beta}(x) \\ &+ \{\mathrm{i}g^*_G(\gamma^0\gamma^k\gamma^5)_{\alpha\beta}\zeta_\rho(v^*)(S^0\gamma^5)^D_{\rho\rho'}\zeta_{\rho'}(v^*)G^0_k(x) \\ &- \mathrm{i}g'^*_G(\gamma^0\gamma^k\gamma^5)_{\alpha\beta}\zeta_\rho(v^*)(S^3\gamma^5)^D_{\rho\rho'} \\ &\cdot \zeta_{\rho'}(v^*)G^3_k(x)\}\psi_{2\beta}(x) \end{aligned} \quad (36)$$

with the effective coupling constants

$$g^*_Z := \hat{f}^Z(0)\left(\frac{2\pi^2 9a}{5}\right)^{1/2}; \qquad g'^*_Z := \frac{1}{3}\hat{f}^Z(0)\left(\frac{2\pi^2 9a}{5}\right)^{1/2}; \quad Z = A, G. \quad (37)$$

First we evaluate the superspin-isospin parts in (36).

### 2.1. Coupling to Symmetric Superspin-Isospin States

For broken CP-symmetry the lepton as well as the boson states must violate the antisymmetry condition and become parafermionic states. A possible candidate for such a parafermionic state was proposed by (9). But at the end of Section 1 objections were raised to the use of this state. *Nevertheless it is instructive to treat the coupling terms of this state to bosons first.*

For this state one gets from (28) and (9)

$$\zeta^S_\rho(v^*) \otimes \zeta^S_{\rho'}(v^*) = \frac{1}{3}[\delta_{3\rho}\delta_{3\rho'} + 2\delta_{4\rho}\delta_{4\rho'}], \quad (38)$$

where the index $S$ means standard representation, i. e. a representation of the superspinors by $\psi$ and $\psi^c$ which is the formulation originally used for the spinor field [I, 17] and which was used for the construction of (9).

In the meantime $G$-conjugated spinors have been introduced and preferred as the latter permit a completely homogenous transformation of the superspinors for the Lorentz group as well as for the isospin group. If the correspondence

$$\begin{aligned} \rho = 1 &\to \Lambda = 1, \quad A = 1, \\ 2 &\to \Lambda = 1, \quad A = 2, \\ 3 &\to \Lambda = 2, \quad A = 1, \\ 4 &\to \Lambda = 2, \quad A = 2 \end{aligned} \quad (39)$$

is used where $A$ is the isospin index, while $\Lambda$ is the superspin index, defined by ordinary spinors $\psi$ and $\psi^c$ fields, the transformation to $G$-conjugated spinors reads

$$\psi^D_{\rho\alpha} = G_{\rho\rho'}\psi^S_{\rho'\alpha}, \quad G = \begin{pmatrix} \mathbf{1} & 0 \\ o & c^{-1} \end{pmatrix} \quad (40)$$

with $c = -\mathrm{i}\sigma_2$. Application of this transformation to (38) yields

$$\zeta^D_\rho(v^*) \otimes \zeta^D_{\rho'}(v^*) = \frac{1}{3}[\delta_{4\rho}\delta_{4\rho'} + 2\delta_{3\rho}\delta_{3\rho'}]. \quad (41)$$

Substitution of such states into (36), then leads to the superspin-isospin matrix elements

$$\begin{aligned} \zeta^{D+}(v^*)(T^0\gamma^5)^D\zeta^D(v^*) &= -1, \\ \zeta^{D+}(v^*)(T^3\gamma^5)^D\zeta^D(v^*) &= \frac{1}{3} \end{aligned} \quad (42)$$

and

$$\begin{aligned} \zeta^{D+}(v^*)(S^0\gamma^5)^D\zeta^D(v^*) &= 1, \\ \zeta^{D+}(v^*)(S^3\gamma^5)^D\zeta^D(v^*) &= -\frac{1}{3}, \end{aligned} \quad (43)$$

where the explicit expressions of the $S\gamma^5$- and $T\gamma^5$-elements in $G$-conjugated form, i. e. $D$-representation can be found in [I, 33].

If these values are substituted into (36) this results in

$$\begin{aligned} \mathrm{i}\partial_t\psi_{2\alpha}(x) &= [-\mathrm{i}(\gamma^0\gamma^k)_{\alpha\beta}\partial_k + \gamma^0_{\alpha\beta}m^*]\psi_{2\beta}(x) \\ &- (\gamma^0\gamma^k)_{\alpha\beta}\left[g^*_A A^0_k(x) + \frac{1}{3}g'^*_A A^3_k(x)\right]\psi_{2\beta}(x) \\ &+ \mathrm{i}(\gamma^0\gamma^k\gamma^5)_{\alpha\beta}\left[g^*_G G^0_k(x) + \frac{1}{3}g'^*_G G^3_k(x)\right]\psi_{2\beta}(x). \end{aligned} \quad (44)$$

In order to study the coupling of the excited neutrino to the 'physical' bosons, the Weinberg transformation must be applied to obtain the boson mixing of the Standard Model which is imperative for the experimental identification of theoretically predicted effects.

Applying this transformation to (44) one has to take into account three conditions:

(i) Although in general there are various lepton-boson couplings, for reasons of consistency there can be applied only one universal Weinberg transformation for all couplings.

(ii) For reasons of consistency the universal Weinberg transformation for magnetic bosons must be the same as that for the electric bosons.

(iii) The universal Weinberg transformation should be chosen as the phenomenological one. The latter transformation slightly differs from that 'Weinberg' transformation which is required for the diagonalization of the boson mass matrix [I, 25].

If the extended Standard Model is to be formulated by means of a Lagrangian, a renormalization of the boson fields and the fermion fields must be performed in order to secure the consistency of the currents in the various field equations. But this will not be carried out in detail.

Due to the regularization all renormalization constants are finite. Hence, their introduction poses no essential problem and is of minor interest with respect to the monopol problem. So one can omit this step and can directly apply the universal Weinberg transformation to the monopol equations (44).

If for conserved symmetries the electric vector potential $A^0$ is coupled to $g$ and $A^3$ to $g'$ then the universal Weinberg angle is introduced by the definitions [12], where in contrast to definition (45) the coupling constants in [12] (6.23) are interchanged

$$\sin\Theta_W = \frac{g}{N}, \quad \cos\Theta_W = \frac{g'}{N} \tag{45}$$

with $N := (g^2 + g'^2)^{1/2}$. The Weinberg transformation can then be written as

$$A_k^3 = Z_k \frac{g'}{N} + A_k \frac{g}{N}, \quad A_k^0 = -Z_k \frac{g}{N} + A_k \frac{g'}{N} \tag{46}$$

and

$$G_k^3 = X_k \frac{g'}{N} + G_k \frac{g}{N}, \quad G_k^0 = -X_k \frac{g}{N} + G_k \frac{g'}{N}, \tag{47}$$

where $A_k$ and $G_k$ are the genuine electromagnetic vector potentials.

Substitution of these transformations into (44) then yields for the coupling terms in (44)

$$\begin{aligned} &-(\gamma^0\gamma^k)_{\alpha\beta}\left[A_k(x)\left(\frac{g_A^* g'}{N} + \frac{1}{3}\frac{g_A'^* g}{N}\right)\right. \\ &\left.-Z_k(x)\left(-\frac{g_A^* g}{N} + \frac{1}{3}\frac{g_A'^* g'}{N}\right)\right]\psi_{2\beta}(x), \\ &\mathrm{i}(\gamma^0\gamma^k\gamma^5)\left[G_k(x)\left(\frac{g_G^* g'}{N} + \frac{1}{3}\frac{g_G'^* g}{N}\right)\right. \\ &\left.+X_k(x)\left(-\frac{g_G^* g}{N} + \frac{1}{3}\frac{g_G'^* g'}{N}\right)\right]\psi_{2\beta}(x). \end{aligned} \tag{48}$$

To simplify matters I assume that the coupling constants to the electric boson fields correspond to those values which are used for the definition of the universal Weinberg transformation, i. e. $g_A^* = g$ and $g_A'^* = g'$. Then the $A$-term in (48) has the nonvanishing coupling constant

$$e = \frac{4}{3}\frac{gg'}{N}, \tag{49}$$

*which clearly shows that for a nonvanishing electric charge the symmetric superspin-isospin state (9) does not allow the identification of the hypothetical 'excited neutrino' with a neutrino in the effective theory. Therefore the objections of Section 1 are justified, i. e. the state (9) can not describe the broken CP-symmetry correctly.*

## 2.2. Superspin-Isospin States with Broken Permutation Symmetry

In [6–8] GBBW equations were analysed which were invariant under the permutation group, while asymmetric GBBW equations were treated in [4]. It is therefore reasonable to look whether for broken permutation symmetry only one term of the sum (9) can be a solution of the latter equation in [4].

**Proposition:** *If for a single term $\theta^2_{\kappa_1\kappa_2\kappa_3}$ of (9) the condition (30) is satisfied, then this term is a solution of the asymmetric GBBW equation in [4] with the propagator in [21], provided the subsidiary conditions for the spin-orbit states are satisfied which arise from projections on this equation with $\theta^2_{\kappa_x\kappa_y\kappa_z}$ where $\kappa_x, \kappa_y, \kappa_z$ are permutations of $\kappa_1, \kappa_2, \kappa_3$.*

**Proof**: The term $\theta^2$ satisfies condition (30) and from [4] it follows that the symmetry breaking part of the propagator does not require the fulfillment of an additional condition. The remaining part of the statement can be verified by means of (10). □

The single term $\theta^2_{\kappa_1\kappa_2\kappa_3} = \delta_{4\kappa_1}\delta_{4\kappa_2}\delta_{3\kappa_3}$ leads to

$$\zeta^D_\rho(\nu^*)\otimes\zeta^D_{\rho'}(\nu^*) = \delta_{3\rho}\delta_{3\rho'}. \quad (50)$$

Application of the Weinberg transformation then yields for the corresponding coupling terms

$$(\gamma^0\gamma^k)A_k\frac{gg'}{N}[(T^0\gamma^5)_{33}+(T^3\gamma^5)_{33}] = 0,$$
$$\mathrm{i}(\gamma^0\gamma^k\gamma^5)G_k[g'^*_G g(S^0\gamma^5)_{33}+g^*_G g'(S^3\gamma^5)_{33}] \neq 0, \quad (51)$$

i. e. the *electric charge vanishes*, which agrees with experimental findings [13], while for the coupling constant $e^*_\nu$ to the magnetic vector potential $G_k$ one gets the *nonvanishing magnetic charge*

$$e^*_\nu = \frac{8}{9}\hat{f}^G(0)\left(\frac{2\pi^2 9a}{5}\right)^{1/2}. \quad (52)$$

From (31) it follows that the constant $a^{-1/2}$ is proportional to the neutrino radius $r^*_\nu$.

Some additional comments are given in the next section.

## 3. Summary of Part Two and Conclusions

According to Lochak (theory) and Urutskoev (experiment) magnetic monopoles should be considered as a new species of light leptonic particles with a magnetic charge but without an electric charge. This hypothesis includes that for these particles no duality transformations should exist which could lead to a transmutation of a magnetic charge into an electric charge and vice versa. Lochak showed that the postulated independence of both types of charges must be theoretically expressed by an additional magnetic vector potential besides the ordinary Maxwellian electro potential and that also for these quantities no mutual duality transformation should be possible.

A hint about the existence of such quantities was found by Lochak in de Broglie's photon theory. In this theory both types of potentials can be obtained but with the reservation that either electric vector potentials or magnetic vector potentials can be derived, i. e. in de Broglie's theory a coexistence of both kinds of solutions is impossible. Nevertheless, if magnetic monopoles do really exist, they must coexist with electrically charged particles. A solution of this problem was found in a quantum field theoretic generalization of de Broglie's photon theory where a coexistence of both species can be achieved. The corresponding states are defined by solutions of generalized de Broglie-Bargmann-Wigner (GBBW) equations. These equations explicitly depend on the groundstate of the system and it can be shown that the simultaneous existence of electric and magnetic boson states is only possible under the condition that a CP-symmetry breaking of their common ground state is present.

Based on this result it is imperative to apply de Broglie's fusion idea also to the states of the leptonic monopoles as according to Lochak the latter should be considered as excited neutrinos which necessarily implies a substructure of these particles.

The properties of such hypothetical monopoles can be studied by calculating their coupling to the electromagnetic vector potentials which include electric as well as the new magnetic parts. Such coupling terms stem from the effective dynamical equations of the extended Standard Model which result by use of a mapping formalism applied to the basic spinor field model. The mapping is based on many (composite) particle states which are algebraically generated from the set of one (composite) particle states of the GBBW equations. The fact that this procedure depends on the structure of the set of composite one-particle states implies that in this way modifications and extensions of the ordinary Standard Model can be generated, if 'new' one particle states appear in this set, as for instance, magnetic bosons and excited neutrinos.

A further complication must be taken into account: According to the present knowledge photons are mixtures of U(1)-states and SU(2)-states. Hence independently of the interest in only the electromagnetic parts of the theory which classically are connected with the U(1) group, in quantum field theory the full electroweak formalism must be applied as otherwise electromagnetic properties cannot be studied. The necessity to start for physical reasons with a $\mathrm{U}(1)\otimes\mathrm{SU}(2)$ spinor field theory from the beginning, prevents the application of duality transformations for the resulting effective Standard Model. Hence, the electric and magnetic charges are really independent quantities.

Furthermore, the CP-symmetry breaking (and the additional isospin symmetry breaking) cannot be restricted to bosons only. Due to the common vacuum the fermions must be also included. This leads to the formation of parafermionic boson as well as of parafermionic fermion states. The exact boson calculation shows that the CP-symmetry breaking manifests itself primarily in the group theoretical structure of the superspin-isospin part of the boson wave functions. If this by symmetry breaking induced modification of the boson states is analogously applied to the superspin-isospin group structure of the leptons, then the existence of magnetic monopoles can be proven.

Finally it should be noted: The CP-symmetry breaking depends on the presence of a suitable medium, for instance water. Sparks in water tanks are bounded space-time phenomena. Hence such a discharge generates only a local CP-symmetry breaking. If excited neutrinos leave the region of broken symmetry and enter into the domain of ordinary conserved symmetry, the reason for being excited, i. e. magnetic, is dropped. Thus there will be a tendency that the excited neutrinos rearrange themselves into ordinary neutrinos in order to adapt themselves to the new symmetry conserving medium, for instance air. Thus the magnetic monopoles become unstable particles in dependence of boundary conditions according to comments about experiments by Ivoilov [13] and Lochak [14]. In addition in the effective Dirac equation no conservation law can be derived for magnetic charges.

In summary it holds:

*(i) Without symmetry breaking no magnetic monopoles can be created.*
*(ii) The excited neutrinos are neutrino states for broken CP-isospin symmetry.*
*(iii) The magnetic charges are new independent quantities.*
*(iv) Magnetic charges have no conservation law.*
*(v) The decay of magnetic monopoles depends on boundary conditions.*

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