Discharges in fluids as a possible source of Electric and Magnetic Electroweak Bosons

Harald Stumpf

Institute of Theoretical Physics, University Tübingen, Germany

ABSTRACT . Based on the assumption that electroweak bosons, leptons and quarks possess a substructure of fermionic constituents, in [1] it was demonstrated that under CP-symmetry breaking "electric" and "magnetic" electroweak bosons coexist as solutions of corresponding bound state equations. The derivation of these states as two fermion composites is treated in an improved way by analyzing the influence of the vacuum on the formation of these states. The vacuum is represented by the propagator of the fermionic constituents and for the case of CP-symmetry breaking it is shown that the symmetry breaking term in the propagator can be correlated to the experimental arrangement. In the associated effective theory for composite bosons and fermions, leptons are coupled to electric as well as to magnetic bosons. A discharge can be theoretically modelled by injecting high energetic electrons in the system the bremsstrahlung of which leads to the generation of both kinds of electroweak bosons.

1 Introduction

In the past decade electric discharges in fluids were reported which induce low energy nuclear reactions. The latter cannot be explained in conventional theory, [2]. It was speculated that these effects might be triggered by the formation or the presence of low mass magnetic monopoles, [2],[3]. As monopoles are to be identified by their fields, the analysis of fields is an inseparable part of monopole hunting. This problem will be treated in this paper.

Until now modern theories in monopole physics have been dominated by the idea that magnetic monopoles manifest themselves as topological field configurations. That idea was first introduced by Dirac in electrodynamics, [4]. In this approach electric as well as magnetic monopoles are associated to one and the same vector potential.

In selfdual electrodynamics electric monopoles and topological magnetic monopoles can be mutually interchanged by duality transformations, leaving the theory invariant and thus making the two types of monopoles indistinguishable, [5]. Nonabelian theories, in general, do not admit selfduality. So in that case electric charges and topological magnetic charges should be truly independent physical quantities. In particular, electric charges are characterized by the conservation of $U(1)$ gauge invariance while the magnetic charges are conserved by topological quantum numbers, [6], p.242.

Although in Dirac's topological approach the quantization of electric charge is demonstrated, the nonabelian topological monopoles seem to be not very suitable for physical applications. They are too heavy for playing any role in the present day accessible energy ranges, their constructions are based only on approximate solutions of the associated field equations and are too sophisticated for being conclusive and their topological stability might be rather inconvenient in considering them as possibly transient phenomena.

A more economical way of treating this matter was started by de Broglie, [7], and later on by Cabibbo and Ferrari, [8]. To formulate electrodynamics (and quantum electrodynamics) in the presence of electric as well as of magnetic charge densities, they introduced two vector potentials instead of one vector potential in the conventional theory and in Dirac's topological approach. This idea leads to the definition of electric and magnetic bosons which are characterized by their different vector potentials contributing in a different way to the common field strength tensor.

In the latter case no topological construction is necessary and electric and magnetic charges are treated on equal footing. In electrodynamics the fields of these bosons are linked by duality transformations which forstall the independence of electric and magnetic bosons and associated charges, [9]. But as was indicated above for nonabelian fields the independence of these bosons can be taken for granted. The extension of the de Broglie-Cabibbo-Ferrari formalism to the full electroweak theory is imperative. In the Standard Model photons are represented by mixtures of $U(1)$ and $SU(2)$ fields and thus the inclusion of the weak isospin is unavoidable.

One characteristic feature of this formalism consists in the fact that in the definition of magnetic charges the fields are not involved, in contrast to the topological case where the magnetic charges are constituted by field configurations themselves. This fact opens the search for magnetic monopoles in analogy to electric monopoles.

As a counterconcept to the topological magnetic monopoles in grand unified theories Lochak showed that massless Dirac fermions can couple to the magnetic bosons and can in this way reveal their magnetic monopole property, [10],[11],[12]. As neutrinos have small masses and as antineutrinos and neutrinos are emitted in large numbers from reactors and the sun without showing electromagnetic reactions, Lochak's fermionic magnetic monopoles cannot be identified with neutrinos directly. Softening the zero mass condition Lochak thus proposed his fermionic monopoles to be identical with excited neutrinos, [13],[14].

Such an assumption has an immediate consequence: As point particles cannot be internally excited, the magnetically excited neutrinos must be composites. This is in accordance with a further discovery by Lochak. He showed that in de Broglies photon theory apart from electric photons also magnetic photons can be derived, both of which are composed of two fermions, [15]. Therefore in the discussion of monopole phenomena the application of de Broglie's fusion idea should or even must be included.

However: in de Broglie's photon theory only magnetic *or* electric photon states can be calculated, i.e. theoretically these states cannot exist simultaneaously and the whole theory is referred to single composite particle states,[16],chap.1. Thus a field theoretic version of de Broglie's and Lochak's discoveries is required which should lead to an extended electroweak Standard Model as an effective theory for electric and magnetic electroweak bosons as well as for fermions. This was advocated by Lochak,[3].

To treat these problems we use a model which is based on a relativistically invariant nonlinear spinor field theory with local interaction, canonical quantization, selfregularization and probability interpretation. This model implies that in the sense of de Broglie and of Heisenberg the present "elementary" particles are assumed to possess a fermionic substructure. The model is expounded in detail in [16].

By purely theoretical reasoning it was demonstrated in [1] that in the above spinor field model electric and magnetic electroweak boson states can coexist if the CP-symmetry of the vacuum is violated. This finding is in accordance with the phenomenological observation that the existence of magnetic monopoles and dyons implies CP-symmetry breaking,[8],[17],[18]. Therefore theoretically one can assume that CP-symmetry breaking is a crucial condition for the discovery of magnetic monopole effects and it is the task to propose scenarios to realize this symmetry breaking in practice. It will be shown in the following that the formal way of theoretical symmetry breaking in [1] reflects mathematically the experimental situation.

With respect to the treatment of this case it has to be emphasized that in our model the method of introducing CP-violation is completely different from the corresponding method in the conventional theory. While in the Standard Model the CP-symmetry breaking is formally introduced by quark mass matrices with complex parameters,cf. [19], chap.26, in our approach this symmetry breaking is effected by an appropriate change of the vacuum. Mathematically this indicates the transition to a new inequivalent field representation which is a common method in algebraic field theory successfully applied in solid state physics, cf. [20],[21].

In consequence of this difference of the methods, the results differ considerably too. While the formal phenomenological method of the Standard Model is intended to explain the decay of K-mesons, the algebraic method of the model under consideration leads to a completely new formulation and structure of the whole theory due to the new inequivalent vacuum. The physical consequences of this algebraic approach are remarkable and can be phenomenologically summarized by an effective Lagrangian.

As our exposition is based on the results of the preceding papers [1],[22],[23], it is unavoidable that for brevity we have to refer to these results without giving renewed deductions. In particular we skip the formulation of the algebraic representation of the basic spinor field itself, a clear exposition of which was given in [24], sect. 2. Furthermore in our calculations no use was made of the decomposition into left-handed and right-handed fermions for simplicity. Insofar the model under consideration is a simplified version of the mathematical structure of the Standard Model. This is justified as already in this version the crucial effects of CP- symmetry breaking can be demonstrated.

2 GBBW-boson equations with CP violation

We start from the two body GBBW (generalized de Broglie-Bargmann-Wigner) equations for bosons without symmetry breaking. With $x \in M^4$ and $Z = (i, \kappa, \alpha)$ where κ means superspin-isospin index, $\alpha =$ Dirac spinor index, $i =$ auxiliary field index, these equations read, cf. [16]

$$[D^{\mu}_{Z_1 X_1} \partial_{\mu}(x_1) - m_{Z_1 X_1}] \varphi_{X_1 Z_2}(x_1, x_2) = \\ 3 U_{Z_1 X_2 X_3 X_4} F_{X_4 Z_2}(x_1 - x_2) \varphi_{X_2 X_3}(x_1, x_1) \tag{1}$$

By interchange of all arguments in (1) one obtains a second set of equations which, however, imposes no further conditions on the wave function φ, if φ is antisymmetric.

In this representation the following definitions are used:

$$D^{\mu}_{Z_1 Z_2} := i \gamma^{\mu}_{\alpha_1 \alpha_2} \delta_{\kappa_1 \kappa_2} \delta_{i_1 i_2}; \qquad m_{Z_1 Z_2} := m_{i_1} \delta_{\alpha_1 \alpha_2} \delta_{\kappa_1 \kappa_2} \delta_{i_1 i_2} \tag{2}$$

with $F_{Z_1 Z_2}(x_1 - x_2)$ as the free spinor field Feynman propagator. This propagator reflects the properties of the vacuum, i.e., the experimental situation. For high energy processes the vacuum can be considered as the background of an ideal gas which leads to the free spinor field propagator.

The vertex term in (1) is fixed by the definition

$$U_{Z_1 Z_2 Z_3 Z_4} := \lambda_{i_1} B_{i_2 i_3 i_4} V^{\kappa_1 \kappa_2 \kappa_3 \kappa_4}_{\alpha_1 \alpha_2 \alpha_3 \alpha_4} \tag{3}$$

where $B_{i_2 i_3 i_4}$ indicates the summation over the auxiliary field indices and where the vertex is given by

$$V^{\kappa_1 \kappa_2 \kappa_3 \kappa_4}_{\alpha_1 \alpha_2 \alpha_3 \alpha_4} := \frac{g}{2} \{ [\delta_{\alpha_1 \alpha_2} C_{\alpha_3 \alpha_4} - \gamma^5_{\alpha_1 \alpha_2} (\gamma^5 C)_{\alpha_3 \alpha_4}] \delta_{\kappa_1 \kappa_2} [\gamma^5 (1 - \gamma^0)_{\kappa_3 \kappa_4}] \}_{as(2,3,4)} \tag{4}$$

The parameters λ_i originate from the regularization procedure.

Equations (1) are relativistically invariant quantum mechanical two body equations with nontrivial interaction, selfregularization and probability interpretation. They possess an associated exact single time energy (relativistic Schroedinger) equation and admit exact solutions. More details can be found in [25],[26],[27],[28].

For vector bosons the exact solutions read in a general form

$$\varphi_{Z_1 Z_2}(x_1, x_2) = T^a_{\kappa_1 \kappa_2} exp[-i \frac{k}{2}(x_1 + x_2)] A^{\mu} \chi^{i_1 i_2}_{\mu, \alpha_1 \alpha_2}(x_1 - x_2 | k) \tag{5}$$

If one excludes states with fermion number 2, which decouple from the effective dynamics, for electroweak bosons the superspin-isospin part is given by a singlet and a triplet matrix representation and defined by the following sets of symmetric and antisymmetric matrices:

$$S^l := \begin{pmatrix} 0 & \sigma^l \\ (-1)^{l+1}\sigma^l & 0 \end{pmatrix}; \qquad T^l := \begin{pmatrix} 0 & \sigma^l \\ (-1)^l \sigma^l & 0 \end{pmatrix}, \qquad l = 1,2,3 \tag{6}$$

for the triplet and

$$S^0 := \begin{pmatrix} 0 & \sigma^0 \\ \sigma^0 & 0 \end{pmatrix}; \qquad T^0 := \begin{pmatrix} 0 & \sigma^0 \\ -\sigma^0 & 0 \end{pmatrix} \tag{7}$$

for the singlet.

The quantum numbers of the bosons can be expressed by the eigenvalues of Q (charge operator) and F (fermion number operator), where Q and F are exclusively defined in superspin-isospin space in accordance with the field theoretic formalism. One obtains, see [16],[26]

$$Q = (G_3 + \frac{1}{2}F) \tag{8}$$

with the isospin generator

$$G_3 := \frac{1}{2}\begin{pmatrix} \sigma^3 & 0 \\ 0 & -\sigma^3 \end{pmatrix}_{\kappa_1\kappa_2} \tag{9}$$

and the fermion number generator

$$F := \frac{1}{3}\begin{pmatrix} \mathbf{1} & 0 \\ 0 & -\mathbf{1} \end{pmatrix}_{\kappa_1\kappa_2} \tag{10}$$

In the two body system an eigenstate of (8) and (10) is represented by a superspin-isospin tensor Θ which satisfies the conditions

$$Q_{\kappa_1\kappa}\Theta_{\kappa\kappa_2} + \Theta_{\kappa_1\kappa}Q_{\kappa\kappa_2} = q\Theta_{\kappa_1\kappa_2} \tag{11}$$

and

$$F_{\kappa_1\kappa}\Theta_{\kappa\kappa_2} + \Theta_{\kappa_1\kappa}F_{\kappa\kappa_2} = f\Theta_{\kappa_1\kappa_2} \tag{12}$$

Corresponding eigenstates can be formed by linear combinations of elements of the sets (6) or (7), respectively. This gives the following result with states and quantum numbers:

$$
\begin{aligned}
\frac{1}{2}(S^0 + S^3) &= \Theta_1^s(q = 0, f = 0) \\
\frac{1}{2}(S^0 - S^3) &= \Theta_2^s(q = 0, f = 0) \\
\frac{1}{2}(S^1 + S^2) &= \Theta_3^s(q = 1, f = 0) \\
\frac{1}{2}(S^1 - S^2) &= \Theta_4^s(q = -1, f = 0)
\end{aligned}
\tag{13}
$$

in the symmetric case and if the S-states are replaced by T-states one obtains the antisymmetric eigenstates:

$$
\begin{aligned}
\frac{1}{2}(T^0 + T^3) &= \Theta_1^a(q = 0, f = 0) \\
\frac{1}{2}(T^0 - T^3) &= \Theta_2^a(q = 0, f = 0) \\
\frac{1}{2}(T^1 + T^2) &= \Theta_3^a(q = 1, f = 0) \\
\frac{1}{2}(T^1 - T^2) &= \Theta_4^a(q = -1, f = 0)
\end{aligned}
\tag{14}
$$

Without additional symmetry breaking the states (14) are associated to the quantum numbers of electroweak vector bosons, where the photon and the Z-boson correspond to linear combinations of Θ_1^a and Θ_2^a which make via isospin symmetry breaking the Z-boson massive, while the photon remains massless. By the same mechanism the charged bosons aquire masses. For a discussion of this mechanism we refer to [16].

Based on these considerations one can proceed to the discussion of CP-violation. The possible modifications of the GBBW-equations are limited by the condition that in any case CPT-invariance must be conserved. For our discussion we adopt the results of a preceding paper on discrete transformations, see [29].

Symmetry breaking takes place via the modification of the vacuum. For GBBW-equations the vacuum is represented by the propagator F in (1). Written in full the free spinor field propagator reads

$$
F := -i\lambda_{i_1}\delta_{i_1 i_2}\gamma^5_{\kappa_1\kappa_2}\int d^4p[(\gamma^\mu p_\mu + m_{i_1})C]_{\alpha_1\alpha_2} e^{-ip(x_1-x_2)}(p^2 - m_{i_1}^2)^{-1}(2\pi)^{-4}
\tag{15}
$$

In [29] the PCT-transformation of time ordered matrix elements was derived which applied to (15) yields the following relation for $x' = -x$

$$F^{\kappa_1\kappa_2}_{\alpha_1\alpha_2}(x_1,x_2)_{i_1i_2} = -\gamma^0_{\kappa_1\kappa_1'}\gamma^5_{\alpha_1\alpha_1'}\gamma^0_{\kappa_2\kappa_2'}\gamma^5_{\alpha_2\alpha_2'}F^{\kappa_1'\kappa_2'}_{\alpha_1'\alpha_2'}(x_1',x_2')'_{i_1i_2} \tag{16}$$

By direct calculation one can prove the invariance of (15), i.e. $F = F'$.

In the same manner by means of the formulas in [29] the CP-transformation of (15) is given by

$$F^{\kappa_1\kappa_2}_{\alpha_1\alpha_2}(x_1-x_2)_{i_1i_2} = (\gamma^0\gamma^5)_{\kappa_1\kappa_1'}\gamma^0_{\alpha_1\alpha_1'}(\gamma^0\gamma^5)_{\kappa_2\kappa_2'}\gamma^0_{\alpha_1\alpha_1'}F^{\kappa_1'\kappa_2'}_{\alpha_1'\alpha_2'}(x_1'-x_2')'_{i_1i_2} \tag{17}$$

with $x' = (-\mathbf{r}, t)$. Again by direct calculation one can prove the invariance of ( ), i.e., F under this transformation.

From this it follows: if in F a CP-violating term X has to be established, owing to the required CPT-invariance this term must satisfy the condition

$$-\gamma^0_{\kappa_1\kappa_1'}\gamma^5_{\alpha_1\alpha_1'}\gamma^0_{\kappa_2\kappa_2'}\gamma^5_{\alpha_2\alpha_2'}X^{\kappa_1'\kappa_2'}_{\alpha_1'\alpha_2'} = X^{\kappa_1\kappa_2}_{\alpha_1\alpha_2} \tag{18}$$

while the CP-violation must lead to

$$(\gamma^0\gamma^5)_{\kappa_1\kappa_1'}\gamma^0_{\alpha_1\alpha_1'}(\gamma^0\gamma^5)_{\kappa_2\kappa_2'}\gamma^0_{\alpha_2\alpha_2'}X^{\kappa_1'\kappa_2'}_{\alpha_1'\alpha_2'} \neq X^{\kappa_1\kappa_2}_{\alpha_1\alpha_2} \tag{19}$$

One easily verifies that such a term can be defined by

$$X^{\kappa_1\kappa_2}_{\alpha_1\alpha_2} = (\gamma^5\gamma^0)_{\kappa_1\kappa_2}C_{\alpha_1\alpha_2} \tag{20}$$

The CP-symmetry breaking term (20) can only be incorporated into F by means of a mass correction term. Thus in accordance with (20) the PC-violating propagator $F(\delta m)$ can be written in the following form

$$\begin{aligned} F(\delta m) =& -i\lambda_{i_1}\delta_{i_1i_2}\gamma^5_{\kappa_1\kappa}(2\pi)^{-4} \\ & \int d^4p\frac{[(\gamma^\mu p_\mu + m_{i_1})\delta_{\kappa\kappa_2} + \delta m\gamma^0_{\kappa\kappa_2}]_{\alpha_1\alpha}C_{\alpha\alpha_2}}{(p^2-m^2_{i_1})}e^{-ip(x_1-x_2)} \\ =& F_0 + F_1 \end{aligned} \tag{21}$$

where F_0 is to be identified with (15) and where δm contains such a phase parameter. From the representation (21) of this modified vacuum expectation value it follows that owing to the decomposition of γ^0

by $\gamma^0 = \sigma^3_{\Lambda_1\Lambda_2} \otimes \delta_{A_1A_2}$, the mass correction term acts in a different way on spinors and charge conjugated spinors causing thus C-violation. As a consequence of this symmetry breaking in the phenomenological treatment matter can be absolutely discriminated from antimatter and an unambiguous definition of positive charge is possible, [30]. Thus C- or CP-violation leads to the discrimination of positive from negative charges, even at the parton level, and thus enforces the introduction of parafermions or parapartons which manifests itself in the violation of the antisymmetry of states.

For the following we accept (21) as a postulate and study the consequences for the formation of electric and magnetic photons. The relation to experiments and the justification of this postulate will be given in section 4.

As the starting point for a detailed calculation of the corresponding boson states we use the integral form of the GBBW-equations (1). Owing to the antisymmetry of the vertex (3), this equation can be written in the form

$$\varphi_{Z_1Z_2}(x_1,x_2) = 3\int d^4x G_{Z_1X_1}(x_1-x)U_{X_1X_2X_3X_4}F_{X_2Z_2}(x-x_2)\varphi_{X_3X_4}(x,x) \tag{22}$$

If in (22) the modified propagator (21) is substituted, the CP-invariance violation by (21) induces a violation of the permutation antisymmetry of the wave functions in (22).

It should be noted that the symmetry properties of the vertex U in (22) are not affected by CP-symmetry breaking because U is related to the dynamical law of the underlying quantum field which does not depend on the special representation induced by associated special vacuum expectation values.

3 Parafermionic electric and magnetic boson states

Symmetry breaking means that the invariance of the Hamiltonian under a certain symmetry group is violated. In consequence the degeneracy of the energy eigenvalues of the Hamiltonian is removed with respect to the members of group representations, i.e. the members of a multiplet aquire different energy eigenvalues and can thus individually identified, i.e. they are no longer indistinguishable.

It is obvious that this fact must find its expression in the properties of the field operator algebra. As the field operators themselves are members of a multiplet of the corresponding group their energetic non-degeneracy makes them distinguishable and thus the anticommutation relations for the case of the unbroken group must be partially replaced by commutation relations for the case of the broken group. I.e., each member of the multiplet has its own fermion state space which does not interact with the fermion state spaces of the other members via the exclusion principle. This necessity is simply illustrated by the fact that in the hydrogen atom the wave function is not antisymmetrized between proton and electron.

In this case according to Green, [31], such a construction with mixed commutators and anticommutators of fermion field operators should therefore be called a parafermi statistics. Without giving further details of this parafermi algebra for field operators, see [32], we directly discuss the effect of this modified algebra on representations. In particular we consider the effect of CP-symmetry breaking on boson wave functions.

For conserved symmetries the exact boson eigenstates can be written in the general form (5). If we decompose the superspin-isospin index κ into Λ = superspinor index and A = isospinor index, then the set of antisymmetric T-matrices in (6) and (7) can be expressed in the form

$$\begin{aligned} &T^0_{\kappa_1\kappa_2} \equiv i\sigma^2_{\Lambda_1\Lambda_2} \otimes \sigma^0_{A_1A_2} \qquad &&T^1_{\kappa_1\kappa_2} \equiv i\sigma^2_{\Lambda_1\Lambda_2} \otimes \sigma^1_{A_1A_2}, \\ &T^2_{\kappa_1\kappa_2} \equiv \sigma^1_{\Lambda_1\Lambda_2} \otimes \sigma^2_{A_1A_2}, \qquad &&T^3_{\kappa_1\kappa_2} \equiv i\sigma^2_{\Lambda_1\Lambda_2} \otimes \sigma^3_{A_1A_2} \end{aligned} \tag{23}$$

The whole wave function is antisymmetric. To establish its antisymmetry the matrices T^l must exactly have the form (23). The matrices $\sigma^a_{\Lambda\Lambda'}$ characterize the content of spinor fields and charge conjugated spinor fields in the corresponding boson state. As in σ^2 as well as in σ^1 only the elements σ^2_{12}, σ^2_{21} and σ^1_{12}, σ^1_{21} are unequal zero, this means that in all T-matrices either symmetrized or antisymmetrized products of one spinor field and one charge conjugated spinor field occur.

That these superspin quantum numbers can be treated separately from the other quantum numbers, depends on the fact, that the GBBW-equations do not comprise the full field dynamics of the underlying spinor field. This property prevents the rediscovery of the field operators $\psi^i_{\Lambda A\alpha}(x)$ by a decomposition of the eigenstates (5).

If CP-symmetry is broken and accordingly spinors and charge conjugated spinors can be discriminated, the permutation representations for these fields are destroyed. In this case it follows from the above considerations that in the eigenstates of the GBBW-equations the only part of the wave functions must be the matrices $\sigma_{\Lambda\Lambda'}$ where this discrimination can be introduced. For comparison we discuss the treatment of conserved and of broken CP-symmetry in parallel terms.

In doing so, we base our argumentation not on the phenomenological superspin-isospin combinations (13),(14), but only on the elementary superspin-isospin matrices (6),(7). The latter are basic for the derivation of the effective gauge theories by means of the application of weak mapping theorems, while the former arise in the course of the further evaluation of the derived effective theories.

Furthermore as all deductions for the basic superspin-isospin matrices run along the same lines, it is sufficient to treat only one example, say T^0 or S^0.

i) symmetry conserving case

In a general form the exact wave functions are given by (5) for $a = 0, 1, 2, 3$.

On the right hand side of the integral equation (22) appears the reduced wave function which for $a = 0$ reads

$$\hat{\varphi}^{\kappa_1\kappa_2}_{\alpha_1\alpha_2}(x,x) = exp(-ikx)\hat{\chi}_{\alpha_1\alpha_2}(0|k)T^0_{\kappa_1\kappa_2} \tag{24}$$

where $\hat{\chi}$ can be expanded on a spinor basis in the form

$$\hat{\chi}_{\alpha_1\alpha_2}(0|k) := [A^0_\mu(k)(\gamma^\mu C)_{\alpha_1\alpha_2} + F^0_{\mu\nu}(k)(\Sigma^{\mu\nu}C)_{\alpha_1\alpha_2}] \tag{25}$$

We substitute in (22) the symmetry conserving propagator (15) and express the Green function in (22) by its Fourier transform

$$G_{Z_1Z_2}(x_1-x_2) = \delta_{i_1i_2}\delta_{\kappa_1\kappa_2}(2\pi)^{-4}\int d^4p(\gamma^\mu p_\mu+m_{i_1})_{\alpha_1\alpha_2}(p^2-m^2_{i_1})^{-1}e^{-ip(x_1-x_2)} \tag{26}$$

Then after insertion of (21), (24) and (26) in (22), the center of mass motion can be eliminated and an equation for χ only results. With $k' := \frac{1}{2}k$ this equation reads

$$\chi^{\kappa_1\kappa_2}_{\alpha_1\alpha_2}(x_1 - x_2|k)_{i_1 i_2} = g' \int d^4p e^{-ip(x_1-x_2)} \lambda_{i_1} f^+_{i_1}[(p_\mu + k'_\mu)\gamma^\mu + m_{i_1}]_{\alpha_1\beta_1}\delta_{\kappa_1\nu_1} \times$$
$$\sum_h [v^h_{\beta_1\beta_2}\delta_{\nu_1\nu_2}(v^h C)_{\beta_3\beta_4}\gamma^5_{\nu_3\nu_4} - v^h_{\beta_1\beta_3}\delta_{\nu_1\nu_3}(v^h C)_{\beta_2\beta_4}\gamma^5_{\nu_2\nu_4}$$
$$-v^h_{\beta_1\beta_4}\delta_{\nu_1\nu_4}(v^h C)_{\beta_3\beta_2}\gamma^5_{\nu_3\nu_2}] \times \tag{27}$$
$$(-i)\lambda_{i_2}\{f^-_{i_2}[(p_\lambda - k'_\lambda)\gamma^\lambda C + m_{i_2}C]_{\beta_2\alpha_2}\gamma^5_{\nu_2\kappa_2}\}\hat{\chi}^{\nu_3\nu_4}_{\beta_3\beta_4}(0|k)$$

where the definitions

$$f^+_i := [(p+k')^2 - m_i^2]^{-1}, \qquad f^-_i := [(p-k')^2 - m_i^2]^{-1} \tag{28}$$

have been introduced. In addition in g' all numerical constants are enclosed being of no relevance for our intended proof.

Then after some algebra one obtains from (27)

$$\chi^{i_1 i_2}_{\alpha_1\alpha_2}(x_1 - x_2|k) i\sigma^2_{\Lambda_1\Lambda_2}\sigma^0_{A_1A_2} =$$
$$g' \int d^4p e^{-ip(x_1-x_2)} \lambda_{i_1} f^+_{i_1}[(p_\mu + k'_\mu)\gamma^\mu + m_{i_1}C]_{\alpha_1\alpha_2}$$
$$A^0_\eta(k)\gamma^\eta_{\beta_1\beta_2} i\sigma^2_{\Lambda_1\Lambda_2}\sigma^0_{A_1A_2} i\lambda_{i_2} f^-_{i_2}[(p_\lambda - k'_\lambda)\gamma^\lambda C + m_{i_2}C]_{\beta_2\alpha_2} \tag{29}$$

Multiplication of (29) with $(i\sigma^2_{\Lambda_1\Lambda_2})^+$ and formation of the trace yields

$$\chi^{i_1 i_2}_{\alpha_1\alpha_2}(x_1 - x_2|k)\sigma^0_{A_1A_2} = g' \int d^4p e^{-ip(x_1-x_2)} \lambda_{i_1} f^+_{i_1}[(p_\mu + k'_\mu)\gamma^\mu + m_{i_1}]_{\alpha_1\beta_1}$$
$$A^0_\eta(k)\gamma^\eta_{\beta_1\beta_2}\sigma^0_{A_1A_2} i\lambda_{i_2} f^-_{i_2}[(p_\lambda - k'_\lambda)\gamma^\lambda C + m_{i_2}C]_{\beta_2\alpha_2} \tag{30}$$

Of course $\sigma^0_{A_1A_2}$ can be eliminated. Then if on the left hand side of (30) the limit $x_1 \to x_2$ and the sum over i_1, i_2 are performed, with substitution of (25) one gets a set of equations for the calculation of $A^0_\mu(k)$ and $F^0_{\mu\nu}(k)$ together with the mass eigenvalue $k^2 = m_B^2$.

ii) case of broken CP-symmetry

For the treatment of this case some additional relations are needed.

Owing to symmetry breaking also the set of symmetric matrices S^a comes into play. For this set the following representation holds

$$\begin{array}{ll} S^0_{\kappa_1\kappa_2} \equiv \sigma^1_{\Lambda_1\Lambda_2} \otimes \sigma^0_{A_1A_2} & S^1_{\kappa_1\kappa_2} \equiv \sigma^1_{\Lambda_1\Lambda_2} \otimes \sigma^1_{A_1A_2}, \\ S^2_{\kappa_1\kappa_2} \equiv i\sigma^2_{\Lambda_1\Lambda_2} \otimes \sigma^2_{A_1A_2}, & S^3_{\kappa_1\kappa_2} \equiv \sigma^1_{\Lambda_1\Lambda_2} \otimes \sigma^3_{A_1A_2} \end{array} \tag{31}$$

Then for the sets (23) and (31) the relations

$$(T^a\gamma^0) = S^a; \qquad (S^a\gamma^0) = T^a \qquad a = 0,1,2,3 \tag{32}$$

hold which are essential for the solution procedure. From (32) it follows

$$(T^a + S^a)\gamma^0 = (S^a + T^a) \qquad a = 0,1,2,3 \tag{33}$$

and for these combinations one obtains the representation

$$(T^a + S^a)_{\kappa_1\kappa_2} = (\sigma^1 + i\sigma^2)_{\Lambda_1\Lambda_2} \otimes \sigma^a_{A_1A_2} \qquad a = 0,1,2,3 \tag{34}$$

In comparison with the superspin parts of the conserved symmetry this matrix is neither symmetric nor antisymmetric. But if one adds equations (13) and (14) one sees that the elements of the set $\{(\Theta^a_b + \Theta^s_b), b = 0,1,2,3\}$ are eigenstates of the operators (8) and (10) too, i.e. the set (34) allows the same state classification as in the CP-symmetry conserving case. In particular with respect to the phenomenological quantum numbers the set (34) or its corresponding combinations are complete.

In addition to the set (34) one can form the set

$$\begin{array}{ll} (T^0 - S^0)_{\kappa_1\kappa_2} = -(\sigma^1 - i\sigma^2)_{\Lambda_1\Lambda_2} \otimes \sigma^0_{A_1A_2} & \\ (T^a - S^a)_{\kappa_1\kappa_2} = (-1)^a(\sigma^1 - i\sigma^2)_{\Lambda_1\Lambda_2} \otimes \sigma^a_{A_1A_2} & a = 1,2,3 \end{array} \tag{35}$$

Hence both sets should be used to describe CP-symmetry breaking within the scope of solutions of GBBW-equations. Therefore we apply the ansatz

$$\begin{aligned} \varphi_{Z_1Z_2}(x_1,x_2|k,a) =& (T^a + S^a)_{\kappa_1\kappa_2} exp[-i\frac{k}{2}(x_1+x_2)]\chi^{i_1i_2}_{\alpha_1\alpha_2}(x_1-x_2|k,a) \\ &+ (T^a - S^a)_{\kappa_1\kappa_2} exp[-i\frac{k}{2}(x_1+x_2)]\tilde{\chi}^{i_1i_2}_{\alpha_1\alpha_2}(x_1-x_2|k,a) \end{aligned} \tag{36}$$

for the left hand side which for the right hand side yields

$$\begin{aligned}
&\hat{\varphi}^{\kappa_1\kappa_2}_{\alpha_1\alpha_2}(x,x|k,a) \qquad (37)\\
=&e^{-ikx}\{[A^a_\mu(k)(\gamma^\mu C)_{\alpha_1\alpha_2}+G^a_\mu(k)(\gamma^5\gamma^\mu C)_{\alpha_1\alpha_2}+F^a_{\mu\nu}(k)(\Sigma^{\mu\nu}C)_{\alpha_1\alpha_2}](T^a+S^a)_{\kappa_1\kappa_2}\\
+&[\tilde{A}^a_\mu(k)(\gamma^\mu C)_{\alpha_1\alpha_2}+\tilde{G}^a_\mu(k)(\gamma^5\gamma^\mu C)_{\alpha_1\alpha_2}+F^a_{\mu\nu}(k)(\Sigma^{\mu\nu}C)_{\alpha_1\alpha_2}](T^a-S^a)_{\kappa_1\kappa_2}\}
\end{aligned}$$

and which replaces the expansion (24),(25) .

The set $\{(T^0-S^0),(-1)^{a+1}(T^a-S^a),a=1,2,3\}$ is the CP-transform of the set (34). Furthermore forming the combinations $\{(\Theta^a_0-\Theta^s_0),(-1)^{a+1}(\Theta^a_a-\Theta^s_a),a=1,2,3\}$ one obtains likewise eigenstates of (8) and (10) with corresponding eigenvalues (13) or (14).

Thus physically the set (35) is redundant. Nevertheless for broken CP-symmetry the states (34) and their CP-transforms must be slightly different. If one substitutes (36) and (37) into (22) one gets an exact system of equations for the calculation of the expansion coefficients in (37).

As the principal effect of CP-symmetry breaking is taken into account by the parafermionic ansatz of the wave function and the set (34) is phenomenological complete, for the sake of brevity we omit the contribution of the CP-transform and apply only the ansatz

$$\begin{aligned}
\hat{\chi}^{\kappa_1\kappa_2}_{\alpha_1\alpha_2}(0|k,a)=[&A^a_\mu(k)(\gamma^\mu C)_{\alpha_1\alpha_2}+G^a_\mu(k)(\gamma^5\gamma^\mu C)_{\alpha_1\alpha_2}\\
&+F^a_{\mu\nu}(k)(\Sigma^{\mu\nu}C)_{\alpha_1\alpha_2}](T^a+S^a)_{\kappa_1\kappa_2}
\end{aligned} \qquad (38)$$

If in (22) the symmetry breaking propagator (21) is substituted then for χ the following integral equation results

$$\begin{aligned}
&\chi^{\kappa_1\kappa_2}_{\alpha_1\alpha_2}(x_1-x_2|k)_{i_1i_2}= \qquad (39)\\
&g'\int d^4pe^{-ip(x_1-x_2)}\lambda_{i_1}f^+_{i_1}[(p_\mu+k'_\mu)\gamma^\mu+m_{i_1}]_{\alpha_1\beta_1}\delta_{\kappa_1\nu_1}\times\\
&\sum_h[v^h_{\beta_1\beta_2}\delta_{\nu_1\nu_2}(v^hC)_{\beta_3\beta_4}\gamma^5_{\nu_3\nu_4}-v^h_{\beta_1\beta_3}\delta_{\nu_1\nu_3}(v^hC)_{\beta_2\beta_4}\gamma^5_{\nu_2\rho_4}\\
&\qquad -v^h_{\beta_1\beta_4}\delta_{\nu_1\nu_4}(v^hC)_{\beta_3\beta_2}\gamma^5_{\nu_3\nu_2}]\times\\
&(-i)\lambda_{i_2}\{f^-_{i_2}[(p_\lambda-k'_\lambda)\gamma^\lambda C+m_{i_2}C]_{\beta_2\alpha_2}\gamma^5_{\nu_2\kappa_2}\\
&\qquad +f^-_{i_2}(\delta m)C_{\beta_2\alpha_2}(\gamma^5\gamma^0)_{\nu_2\kappa_2}\}\hat{\chi}^{\nu_3\nu_4}_{\beta_3\beta_4}(0|k)
\end{aligned}$$

and after some algebra one obtains from (39) the following equation

$$\chi^{i_1 i_2}_{\alpha_1\alpha_2}(x_1 - x_2|k)(\sigma^1 + i\sigma^2)_{\Lambda_1\Lambda_2}\sigma^0_{A_1A_2} = \tag{40}$$
$$g'' \int d^4p e^{-ip(x_1-x_2)} \lambda_{i_1} f^+_{i_1}[(p_\mu + k'_\mu)\gamma^\mu + m_{i_1}]_{\alpha_1\beta_1} \times$$
$$\{[A^0_\eta(k)\gamma^\eta_{\beta_1\beta_2}T^0_{\kappa_1\kappa_2} + G^0_\eta(k)(\gamma^5\gamma^\eta)_{\beta_1\beta_2}S^0_{\kappa_1\kappa_2}]i\lambda_{i_2}f^-_{i_2}[(p_\lambda - k'_\lambda)(\gamma^\lambda C) + m_{i_2}C]_{\beta_2\alpha_2}$$
$$+[A^0_\eta(k)\gamma^\eta_{\beta_1\beta_2}S^0_{\kappa_1\kappa_2} + G^0_\eta(k)(\gamma^5\gamma^\eta)_{\beta_1\beta_2}T^0_{\kappa_1\kappa_2}]i\lambda_{i_2}f^-_{i_2}(\delta m)C_{\beta_2\alpha_2}\}$$

Substitution of (23) and (31) for $a = 0$ into (40), multiplication with $[(\sigma^1 + i\sigma^2)_{\Lambda_1\Lambda_2}]^+$ and formation of the trace yields

$$\chi^{i_1 i_2}_{\alpha_1\alpha_2}(x_1 - x_2|k)\sigma^0_{A_1A_2} \tag{41}$$
$$g'' \int d^4p e^{-ip(x_1-x_2)} \lambda_{i_1} f^+_{i_1}[(p_\mu + k'_\mu)\gamma^\mu + m_{i_1}]_{\alpha_1\beta_1} \times$$
$$\{[A^0_\eta(k)\gamma^\eta_{\beta_1\beta_2}\sigma^0_{A_1A_2} + G^0_\eta(k)(\gamma^5\gamma^\eta)_{\beta_1\beta_2}\sigma^0_{A_1A_2}]i\lambda_{i_2}f^-_{i_2}[(p_\lambda - k'_\lambda)(\gamma^\lambda C) + m_{i_2}C]_{\beta_2\alpha_2}$$
$$+[A^0_\eta(k)\gamma^\eta_{\beta_1\beta_2}\sigma^0_{A_1A_2} + G^0_\eta(k)(\gamma^5\gamma^\eta)_{\beta_1\beta_2}\sigma^0_{A_1A_2}]i\lambda_{i_2}f^-_{i_2}(\delta m)C_{\beta_2\alpha_2}\}$$

Analogous equations can be derived for $a = 1, 2, 3$. Of course $\sigma^0_{A_1A_2}$ can be eliminated from (41). If on the left hand side the limit $x_1 \to x_2$ and the sum over i_1, i_2 are performed, with substitution of (38) into the left hand side of the resulting equation, one gets the set of equations for the calculation of $A^0_\mu(k)$, $G^0_\mu(k)$ and $F^0_{\mu\nu}(k)$ as well as of the mass eigenvalue $k^2 = m^2_B$. For this evaluation we refer to [1].

It is not necessary to discuss details of the eigenvalue problem as the mass eigenvalues of the single (composite) boson states undergo a mass renormalization if the corresponding theory for the quantum field theoretic effective boson-fermion dynamics is derived.

But if one assumes that the secular problem for the corresponding set of equations is solved, then with the definite values of the above field variables the expansion (38) is completely fixed and owing to the Lorentz-invariance of equation (22) this holds for any four-vector k on the hyperboloid $k^2 = m^2_B$. This information is sufficient to fix via equation (22) or (41), respectively, the boson wave functions completely. Due to the selfregularization in all calculations no divergencies occur, i.e. for CP-invariant as well as for broken CP-symmetry exact finite boson eigenstates can be derived.

So the question remains:

Where are the electric and magnetic electroweak bosons?

This question is really answerered by the weak mapping theorems. The latter proceeding allows to study the consequences of boson (and fermion) bound state calculations on the phenomenological level. An extensive discussion and corresponding calculations with respect to this topic have been given in [23]. It would exceed the scope of this paper to repeat the deductions given there. So we refer to this paper and argue in the following on the basis of the results obtained there.

4 Correspondence of the propagator to experiments

In theory the propagator represents the vacuum. To show its correspondence to experiments the following information about the structure of the propagator is required:

Proposition 1: For free spinor fields χ_Z the associated superspin-isospin propagator $F_{Z_1Z_2}(x_1 - x_2)$ can be decomposed into the sum of the conventional fermion and the conventional antifermion propagators at any time $\tau := t_1 - t_2$. The fermion and the antifermion propagators are identical.

Proof: The index Z is defined by $Z :=$ (A; Λ, α, i). For the proof the indices A and i are spectator indices and will be suppressed for the sake of brevity.

$$F^{\alpha_1\alpha_2}_{\Lambda_1\Lambda_2}(x_1 - x_2) := \Theta(t_1 - t_2)\langle 0|\chi_{\Lambda_1\alpha_1}(x_1)\chi_{\Lambda_2\alpha_2}(x_2)|0\rangle + \Theta(t_2 - t_1)\langle 0|\chi_{\Lambda_2\alpha_2}(x_2)\chi_{\Lambda_1\alpha_1}(x_1)|0\rangle \quad (42)$$

i) Without loss of generality we assume $t_1 - t_2 > 0$. Then (42) reads

$$\begin{aligned} F^{\alpha_1\alpha_2}_{\Lambda_1\Lambda_2}(x_1 - x_2) &= \langle 0|\chi_{\Lambda_1\alpha_1}(x_1)\chi_{\Lambda_2\alpha_2}(x_2)|0\rangle \\ &= \delta_{\Lambda_1 1}\delta_{\Lambda_2 2}\langle 0|\chi_{1\alpha_1}(x_1)\chi_{2\alpha_2}(x_2)|0\rangle + \delta_{\Lambda_1 2}\delta_{\Lambda_2 1}\langle 0|\chi_{2\alpha_1}(x_1)\chi_{1\alpha_2}(x_2)|0\rangle \\ &= \delta_{\Lambda_1 1}\delta_{\Lambda_2 2}\langle 0|\chi_{\alpha_1}(x_1)\chi^c_{\alpha_2}(x_2)|0\rangle + \delta_{\Lambda_1 2}\delta_{\Lambda_2 1}\langle 0|\chi^c_{\alpha_1}(x_1)\chi_{\alpha_2}(x_2)|0\rangle \end{aligned} \quad (43)$$

For free fields the vacuum is invariant under discrete transformations P, C, T. This means $\mathcal{C}|0\rangle \equiv |0\rangle$ and

$$\begin{aligned} \langle 0|\chi^c_{\alpha_1}(x_1)\chi_{\alpha_2}(x_2)|0\rangle &= \langle 0|\mathcal{C}^{-1}\mathcal{C}\chi^c_{\alpha_1}(x_1)\mathcal{C}^{-1}\mathcal{C}\chi_{\alpha_2}(x_2)\mathcal{C}^{-1}\mathcal{C}|0\rangle \\ &= \langle 0|\chi_{\alpha_1}(x_1)\chi^c_{\alpha_2}(x_2)|0\rangle \end{aligned} \quad (44)$$

Hence the both parts in (43) are identical.

Furthermore the two parts have to be shown to be particle and antiparticle propagators.

ii) According to [33], sect 15c it is $\chi = \chi^+ + \chi^-$ and one gets

$$\begin{aligned}\langle 0|\chi_{\alpha_1}(x_1)\chi^c_{\alpha_2}(x_2)|0\rangle &= C_{\alpha_2\alpha'_2}\langle 0|\chi_{\alpha_1}(x_1)\bar{\chi}_{\alpha'_2}(x_2)|0\rangle \\ &= -iC_{\alpha_2\alpha'_2}S^+_{\alpha_1\alpha'_2}(x_1 - x_2)\end{aligned} \tag{45}$$

In the same way it follows

$$\langle 0|\chi^c_{\alpha_1}(x_1)\chi_{\alpha_2}(x_2)|0\rangle = -iC_{\alpha_1\alpha'_1}S^-_{\alpha_2\alpha'_1}(x_2 - x_1) \tag{46}$$

where S^+ is the conventional particle propagator while S^- is the conventional antiparticle propagator. $\diamondsuit$

The corresponding experiments of Urutskoev et al.,[2], are rather intricate. Discharges between metallic foils in vessels filled with various fluids lead to the evidence of numerous elements being not present in the system before the explosion and depending on the special foils and fluids. Inspite of a lot of interesting semiempirical studies,see for instance [34], the physical mechanism underlying these low energy nuclear reactions is unknown. Only the action of strong forces is definitely excluded.

It is the aim of this paper to exclusively study the basic effect of such discharges in matter which manifests itself in the production of electric and magnetic electroweak bosons. As a secondary effect such bosons might induce nuclear reactions which is not the topic to be discussed here. In accordance with the statements in the introduction we assume that discharges in matter should be interpreted as the cause of a symmetry breaking of the vacuum, leading to a new inequivalent vacuum state, a view which was already put forward in three preceding papers, [1],[22],[23].

As a suitable canditate for this symmetry breaking an experimental violation of the commonly assumed CP-invariance was considered. Hence in the following we try to show how this symmetry breaking comes about and how the theoretical counterpart of the experimental arrangement has to be formulated.

Based on the above theorem one can analyze the modifications which are necessary to adapt the theoretical description of the vacuum to the experimental arrangements. Two properties are characteristic for these experiments:

i) the reactions are confined to the interior of a closed vessel

ii) the reactions by discharges proceed within a fluid medium

We concentrate on the theoretical description of ii) because in contrast to i) an essential change of the properties of the vacuum has to be expected. This is due to the fact that for discharges in fluid media the motion of the charge carriers is damped, i.e. accompanied by energy losses. One of the standard fluids used in these experiments is water. To be definite we refer our arguments to water.

The stopping power of matter for fast particles has been extensively discussed in [35], Sect. 23. For electrons in water the following ranges have been calculated

Primary energy	0.1	1	10	100	1000	mc^2
H_2O	0.47×10^{-2}	0.19	2.6	19	78	cm

For low energies the magnitudes of these ranges fit into the dimensions of the vessels in the above experiments.

The fact that positrons can be annihilated somewhere in their paths diminishes the average ranges of positrons in comparison with those of electrons. Corresponding formulas describing these differences are given in [35], Sect. 23, eqs.(21),(22). Apart from the importance of numerical values, the principal effect of these differences consists in the signal of a symmetry breaking. Charged particles and their antiparticles behave differently in a medium which leads to C- or CP-symmetry breaking, respectively.

Discharges are triggered by electrons which on their way in the fluid ionize molecules, generate secondary electrons, etc.. But to avoid theoretical difficulties in the description of the rather complicated processes of a discharge, we simplify the theoretical treatment by considering only an average damping effect of electrons.

In such discharges no positrons occur and their real presence is not necessary, because the inclusion of positrons is a theoretical concept to study the behavior of the system under charge conjugation. The absence of positrons in experiments is therefore no argument against positrons in the theoretical treatment.

In our model electrons and positrons are assumed to have a fermionic substructure. Therefore the question is: do the above considerations apply to their fermionic constituents as well?

The group theoretical representation of lepton states with respect to superspin-isospin combinations is exact. In particular for the superspin part one obtains, see [22],

$$e^{+} \quad \rightarrow \quad \delta_{1\Lambda_1}\delta_{1\Lambda_2}\delta_{1\Lambda_3}; \qquad e^{-} \quad \rightarrow \quad \delta_{2\Lambda_1}\delta_{2\Lambda_2}\delta_{2\Lambda_3} \tag{47}$$

where Λ is the superspinor index. With respect to this index the fermion number f is defined. In terms of the fermion numbers one gets for e^+ the configuration (1/3,1/3,1/3) while e^- leads to the set (- 1/3,- 1/3, - 1/3). Fermion numbers are used to discriminate particles from antiparticles by changing f into -f by convention.

Thus from the above sets of fermion numbers it follows: electrons consist only of partons, positrons only of antipartons. In the complete representation of the wave functions this property is not changed. Therefore the partons of the electrons and the antipartons of the positrons share their behavior with that of electrons or positrons, respectively, and we can base our arguments on the spinor field propagator F instead of the phenomenological electron-positron Feynman propagator.

In the next step we consider the influence of damping on the motion of partons and antipartons. Their motion is described by the propagator and if damping is effective their motion ceases in a finite time interval. This fact can be expressed by a damping factor in the integral representation (15) of the propagator. According to Feynman this integral can be evaluated by giving the mass an infinitesimal negative imaginary part, i.e., $m \rightarrow m - i\delta$, $\delta > 0$. If δ is allowed to have a finite value this leads for $t_1 - t_2 > 0$ to a damping factor $exp[-\delta'(t_1 - t_2)]$ with $\delta' := m\delta$, while for $t_1 - t_2 < 0$ one obtains the damping factor $exp[\delta'(t_1 - t_2)]$ in the space-time representation of the propagator.

The damping factor introduced in this way is independent of superspin-isospin states and it does not allow a different behavior of partons and antipartons. But according to Prop. 1 the propagator (15) can be decomposed into a pure particle and a pure antiparticle propagator. Thus these propagators can be treated separately with different damping factors which leads to the experimentally observed different behavior of particles and antiparticles in the fluid. Formally this can be described by giving in the corresponding propagator equation the mass a superspin-isospin dependent damping factor which respects the above decomposition.

For the free superspin-isospin propagator (15) this equation reads

$$[i\gamma^{\mu}_{\alpha\alpha_1}\partial_\mu(x_1)-m_i\delta_{\alpha\alpha_1}]\delta_{\kappa\kappa_1}\delta_{ii_1}F^{\kappa_1\kappa_2}_{\alpha_1\alpha_2}(x_1,x_2)_{i_1i_2}=-iC_{\alpha\alpha_2}\gamma^5_{\kappa\kappa_2}\lambda_i\delta_{ii_2}\delta(x_1-x_2) \tag{48}$$

According to [35] the damping factors of particles and antiparticles differ only weakly compared with their average absolute values. We therefore use the following formulation of this fact

$$\delta_{ii'}\delta_{\kappa\kappa'}\delta_{\alpha\alpha'}m \to (m^*\delta_{\kappa\kappa'}-i\delta\gamma^0_{\kappa\kappa'})\delta_{\alpha\alpha'}\delta_{ii'} \tag{49}$$

with $m^* := m - i\delta^*$ where δ^* is the average damping factor, while δ means the small difference between the damping factors of particles and antiparticles.

Proposition 2: In the propagator equation (48) the mass term (49) violates the CP forminvariance of the propagator and of the equation.

Proof: The propagator equation (48) for mass m and the propagator itself are forminvariant under CP-transformation. If in (48) the mass m is replaced by eq. (49) a violation of CP-invariance can only be caused by this modified mass term.

The CP-transformation of the propagator equation (48) with mass term (49) is carried out in accordance with the general transformation law (19). This gives for the modified mass term the following expression:

$$\begin{aligned}&(\gamma^0\gamma^5)^+_{\lambda\kappa}(\gamma^0\gamma^5)^+_{\lambda_2\kappa_2}\gamma^0_{\beta\alpha}\gamma^0_{\beta_2\alpha_2}(m^*\delta_{\kappa\kappa_1}-i\delta\gamma^0_{\kappa\kappa_1})\delta_{\alpha\alpha_1}\delta_{ii_1}\times\\&(\gamma^0\gamma^5)_{\kappa_1\kappa_1'}(\gamma^0\gamma^5)_{\kappa_2\kappa_2'}\gamma^0_{\alpha_1\alpha_1'}\gamma^0_{\alpha_2\alpha_2'}F^{\kappa_1'\kappa_2'}_{\alpha_1'\alpha_2'}(x_1',x_2')'_{i_1i_2}\\\equiv&(m^*\delta_{\lambda\kappa_1'}+i\delta\gamma^0_{\lambda\kappa_1'})\delta_{\beta\alpha_1'}\delta_{ii_1}F^{\kappa_1'\lambda_2}_{\alpha_1'\beta_2}(x_1',x_2')'_{ii_2}\end{aligned} \tag{50}$$

If in this way the whole equation is transformed one obtains

$$\begin{aligned}&[i\gamma^{\mu}_{\beta\alpha_1'}\partial_\mu(x_1')\delta_{\lambda\kappa_1'}-(m^*\delta_{\lambda\kappa_1'}+i\delta\gamma^0_{\lambda\kappa_1'})\delta_{\beta\alpha_1'}]\delta_{ii_1}F^{\kappa_1'\lambda_2}_{\alpha_1'\beta_2}(x_1',x_2')'_{i_1i_2}\\&=-iC_{\beta\beta_2}\gamma^5_{\lambda\lambda_2}\lambda_i\delta_{ii_2}\delta(x_1'-x_2')\end{aligned} \tag{51}$$

By an appropriate change of indexing one can reestablish the original denotation of equation (48). Then a comparison between eq. (48) with mass (49) and eq. (51) shows: The damping term has changed its sign under CP-transformation, i.e. CP-invariance is violated. This then holds for the solution of eq. (51), i.e. the propagator too.$\diamondsuit$

Apart from discrete symmetry operations the propagator equation (48) admits the application of the continuous $SU(2)$-isospin group and the abelian $U(1)$-fermion number group. We are particularly interested in the isospin group.

Proposition3: The propagator equation (48) is forminvariant under global isospin transformations.

In the same manner one can show

Proposition 4: The mass term (49) with additional CP-symmetry breaking damping term is forminvariant under global isospin transformations.

However, invariance under local and global electroweak isospin transformations is a theoretical concept which is not met in physical reality. Although isospin symmetry breaking manifests itself on the phenomenological level only, the proper treatment of the phenomenological theory as an effective theory requires isospin symmetry breaking already on the parton level, i.e., in the parton propagator F.

This symmetry violating parton propagator was extensively treated in [16],sect.8.3. We adopt this symmetry breaking mass correction term from [16], and introduce it directly in the propagator equation. With inclusion of the CP-symmetry violating damping term and referred to the representation with charge conjugated spinors, this equation reads

$$\begin{aligned}&[i\gamma^{\mu}_{\alpha\alpha_1}\delta_{\kappa\kappa_1}\partial_\mu(x_1) - m_i^*\delta_{\alpha\alpha_1}\delta_{\kappa\kappa_1} - i\theta\delta_{\alpha\alpha_1}\gamma^0_{\kappa\kappa_1} + \theta'\delta_{\alpha\alpha_1}(\gamma^0\gamma^3)_{\kappa\kappa_1}] \times \\ &\delta_{ii_1}F^{\kappa_1\kappa_2}_{\alpha_1\alpha_2}(x_1,x_2)_{i_1i_2} = -iC_{\alpha\alpha_2}\gamma^5_{\kappa\kappa_2}\lambda_i\delta_{ii_2}\delta(x_1-x_2)\end{aligned} \quad (52)$$

One easily verifies

Proposition 5 Equation (52) breaks CP- as well as global isospin invariance.

If the corresponding propagator is used for the derivation of an effective electroweak theory it causes also the violation of the local isospin invariance.

5 Boson production by Bremsstrahlung

The consequences of changes of the vacuum can be studied by the derivation of a corresponding effectice theory. This was done in preceding papers and for the sake of brevity we refer to the results obtained there without giving renewed deductions. It is possible to express the resulting effective theory by an effective Lagrangian derived in [23]. For CP-symmetry breaking this Lagrangian is given by

$$
\begin{aligned}
\mathcal{L}{:=}&-\frac{1}{4}F^a_{\mu\nu}\eta^{\mu\varrho}\eta^{\nu\kappa}F^a_{\varrho\kappa}+\frac{1}{2}\mu_A^2 A^a_\mu\eta^{\mu\varrho}A^a_\varrho+\frac{1}{2}\mu_G^2 G^a_\mu\eta^{\mu\varrho}G^a_\varrho\times \\
&\frac{i}{2}[\bar\psi\gamma^\mu\partial_\mu\psi+(\partial_\mu\bar\psi)\gamma^\mu\psi]-m\bar\psi\psi \\
&-g_e\sum_{a=1}^{3}A^a_\mu j^\mu_a+g'_eA^0_\mu j^\mu_0-ig_m\sum_{a=1}^{3}G^a_\mu J^\mu_a+ig'_mG^0_\mu J^\mu_0
\end{aligned}
\tag{53}
$$

with $j^\mu_a := \frac{1}{2}\bar\psi\sigma_a\gamma^\mu\psi$ and $J^\mu_a := \frac{1}{2}\bar\psi\sigma_a(\gamma^5\gamma^\mu)\psi$, a = 1,2,3,0.

The effective Lagrangian density (53) is limited to a finite range of energy. Above a certain energy threshold it must be modified by formfactors. In this way one does not encounter the divergence difficulties of conventional field theories for this type of Lagrangian.

Furthermore it is $\psi_1 := \psi_\nu$, $\psi_2 := \psi_e$ in isospace and the field strength tensor is defined by, see [23]

$$
\begin{aligned}
F^a_{\mu\nu}{:=}&\partial_\mu A^a_\nu-\partial_\nu A^a_\mu-\varepsilon_{\mu\nu\varrho\sigma}\eta^{\varrho\varrho'}\eta^{\sigma\sigma'}\partial_{\varrho'}G^a_{\sigma'} \\
&+\eta^{abc}(g_1A^b_\mu A^c_\nu+g_2G^b_\mu G^c_\nu+g_3\varepsilon_{\mu\nu\varrho\sigma}\eta^{\varrho\varrho'}\eta^{\sigma\sigma'}A^b_{\varrho'}G^c_{\sigma'})
\end{aligned}
\tag{54}
$$

In comparison with the Lagrangian of a $SU(2)\otimes U(1)$ gauge theory, the striking properties of the Lagrangian (53) with fields (54) are the appearance of an additional magnetic vector potential G_μ and the coupling of leptons (and quarks) to the magnetic currents as a consequence of the CP-symmetry breaking of the vacuum. As this CP-symmetry breaking violates local isospin symmetry but conserves global isospin symmetry, see [25], it is obvious that an additional global isospin symmetry breaking must be introduced in order to obtain the physical relevant electroweak boson and fermion fields.

Then in the weak mapping procedure an important additional effect arises by the appearance of a nontrivial boson mass matrix for the electric and magnetic electroweak boson fields, if CP-symmetry and isospin invariance are simultaneously broken. The consequences of this combined symmetry breaking can be summarized in the following way:

i) The charged sector of the vector bosons undergoes a synchronous mass splitting with the neutral sector. This mass splitting leads to very light negatively charged bosons and very heavy positively ones. As the bare coupling constants of the electromagnetic and of the weak processes are of the same magnitude, cf. [30],eqs. (10.41),(10.42), and the same holds for the masses of the W^--boson and the photon, the transition rates for these bosons must be of the same magnitude too. In consequence this produces a considerable change of the transition rates of processes where the W^--bosons are involved in comparison with the symmetry conserving case. On the other hand due to the increase in the mass of the W^+-boson corresponding processes are suppressed.

ii) The charged fermions of the CP-symmetric theory are transmuted into dyons if the CP-symmetry is broken, a fact which is independent of isospin symmetry breaking. This can be read off from the corresponding effective Lagrangian for the fermions. This Lagrangian is derived from the fermion part of the Lagrangian density (53) and can be expressed by the “electric” and “magnetic” currents in the form

$$\mathcal{L}_f(x) = \bar{\psi}(-i\gamma^\mu\partial_\mu + m)\psi + \frac{1}{2}g\tilde{A}^a_\mu\tilde{j}^\mu_a + i\frac{1}{2}g\tilde{G}^a_\mu\tilde{J}^\mu_a \tag{55}$$

where the physical fields and currents are denoted by the tilde sign.

Based on these results the production of electric and magnetic electroweak bosons can be understood if in this system the effect of a discharge is included. Such a discharge can be simulated by the sudden injection of high energetic electrons which can be theoretically described by an initial condition at a certain time, say $t = 0$. The main effects of such an injection of electrons are inelastic collisions with atoms of the surrounding matter and bremsstrahlung by deflection processes. If according to (55) the electrons are coupled to electric as well as to magnetic vector potentials this gives rise to the emission of electric as well as to magnetic electroweak bosons, i.e. to their production. i.e. the conventional theory of the stopping power of matter for fast particles must be extended to include the processes with magnetic bosons.

A further effect of such an explosion might be the excitation of neutrinos to magnetically charged particles which at present has not yet been theoretically analyzed. In any case neither the dyons nor the excited neutrinos are carriers of conserved magnetic charges. The latter are simulated if magnetic electroweak bosons are generated by the discharge or the explosion, respectively and fade away if the magnetic bosons are absorbed or dispersed. In any case the energy to create such bosons comes from the original high energetic electrons and due to the modified mass values of the electroweak bosons not only photons but also W-bosons might be generated in interaction with the medium and can in this way induce nuclear reactions.

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