Boson condensations in the Glashow-Salam-Weinberg electroweak theory

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Condensations of the boson fields in the Glashow-Salam-Weinberg model of the electroweak interactions are considered. Apart from the well-known phase with the Higgs boson condensation the new phases appear. The electroweak magnetic phase with a very deep energy density minimum of $\mathcal{E}\approx 44.382$ GeV^4 for charge density equal approximately to 0.5539 GeV^3 emerges. In this phase the $W_\mu^\pm$ bosons are massless and the photon A_μ acquire the nonzero mass. Droplets of this phase could be experimentally observed by their very small ratio $Q/M_Q \leq 1/80.13$ GeV^{-1} (where Q is the electric charge of the droplet with the mass M_Q). Another phase with Z_μ condensation and its stability are also examined. The experimental knowledge of a droplet of this phase with the upper possible mass M_{I^3} (where I^3 is the weak isotopic charge of the droplet) could give us the value of the Higgs mass.

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I. INTRODUCTION

The Glashow-Salam-Weinberg (GSW) theory of weak interactions is quite a well-established and verified model [1], especially in its perturbative regime. The aim of this paper is to examine the nonperturbative phenomenon of boson condensations in this theory. The gauge field condensations [2] are interesting from the theoretical point of view where they may play the same role as the Higgs fields in the process of breaking the symmetry [3]. Boson condensations are also the subject of growing interest in the field of astrophysics where the presence of superdense matter is taken into account (for example, neutron starts or even more exotic cases [4]).

The boson condensation induced by the external charge may change the physical system. The new vacuum state may be interpreted as the coherent state [5] (see also Appendix A). It may drastically change the physical system. It is suggested that the electromagnetic vacuum in the presence of the external charge is unstable [6]. As a result the new charge vacuum which is accompanied with particle-antiparticle pair production may appear. Such phenomenon could happen in heavy-ion collisions [7] or inside astrophysical compact objects (neutron stars, strange stars, or boson stars [4]).

In this paper we emphasize boson condensation in the Glashow-Salam-Weinberg (GSW) model in the presence of external sources. The GSW model gives the rich structure of possible phases. Apart from the well known phase with the Higgs boson condensation the new phases with the gauge boson condensations may appear. In this paper it will be shown that the energy density of the new phase with the "electroweak magnetic field" has a very deep local minimum of $\mathcal{E}\approx(2.5811\ \mathrm{GeV})^4$ for charge density equal approximately to 0.5539 GeV^3. In this phase the $W_\mu^\pm$ bosons are massless and the photons A_μ acquire a nonzero mass. This means that the electromagnetic interactions will be suppressed in a similar fashion to the superconductivity case. We may expect such boson condensations in very dense objects where they produce locally the lowering of the energy.

Also another phase [8] appears with Z_μ boson condensation. There exists an upper limit for the Higgs nonlinear λ parameter in a certain range of the Z_μ boson condensation.

II. THE GENERAL THEORY

The Lagrangian density of the electroweak $\mathrm{SU}_L(2)\times \mathrm{U}_Y(1)$ model is summarized as

$$\mathcal{L}=-\tfrac{1}{4}F^a_{\mu\nu}F^{a\mu\nu}-\tfrac{1}{4}B_{\mu\nu}B^{\mu\nu}+(D_\mu H)^+D^\mu H-U(H)+\mathcal{L}_f \tag{1}$$

with the fermionic part $\mathcal{L}_f$ given by

$$\mathcal{L}_f=i\bar{L}\gamma^\mu D_\mu L+i\bar{R}\gamma^\mu D_\mu R-\sqrt{2}\frac{m}{v}(\bar{L}HR+\mathrm{H.c.})\ , \tag{2}$$

where m is the physical mass of the electron and v is the constant parameter. Here the $U_Y(1)$ field tensor is defined as

$$B_{\mu\nu}=\partial_\mu B_\nu-\partial_\nu B_\mu \tag{3}$$

and the $\mathrm{SU}_L(2)$ Yang-Mills field tensor as

$$F^a_{\mu\nu}=\partial_\mu W^a_\nu-\partial_\nu W^a_\mu-g\varepsilon_{abc}W^b_\mu W^c_\nu\ , \tag{4}$$

where the ε_{abc} are the structure constants for $\mathrm{SU}_L(2)$ (ε_{abc} is antisymmetric under the interchange of two neighboring indexes and $\varepsilon_{123}=+1$).

The covariant differentiation D_μ is given by

$$D_\mu H=\partial_\mu H+igW_\mu H+\tfrac{1}{2}ig'YB_\mu H\ , \tag{5}$$

$$D_\mu L=\partial_\mu L+igW_\mu L+\tfrac{1}{2}ig'YB_\mu L\ , \tag{6}$$

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$$D_\mu R=\partial_\mu R+\tfrac{1}{2}ig'YB_\mu R\ , \tag{7}$$

where

$$W_\mu=W^a_\mu\frac{\sigma^a}{2} \tag{8}$$

is the gauge field decomposition with respect to the su(2) algebra generators.

The potential of the scalar fields is

$$U(H)=\lambda(H^+H-\tfrac{1}{2}v^2)^2 \tag{9}$$

with the Higgs doublet

$$H=\begin{bmatrix}H^+\\ H^0\end{bmatrix}\ , \tag{10}$$

which after making a local symmetry transformation can be written as

$$H=\frac{1}{\sqrt{2}}\begin{bmatrix}0\\ \varphi\end{bmatrix}\ . \tag{11}$$

Here φ is the Higgs field.

In our notation we specify only the electron and its neutrino. The contributions from quarks and other leptons can be treated in a similar way. Here we adopt the notation

$$L=\begin{bmatrix}\nu_L\\ e_L\end{bmatrix}\quad\text{and } R=(e_R)\ . \tag{12}$$

The coupling constant for $SU_L(2)$ is called g, and by convention the $U_Y(1)$ coupling is $g'/2$. The weak hypercharge operator for the $U_Y(1)$ group is called Y. Quantum numbers in the electroweak $SU_L(2)\times U_Y(1)$ model are given in Table I.

The relations among the Weinberg angle Θ_W, g, and g' are as follows:

$$\cos\Theta_W=\frac{g}{\sqrt{g^2+g'^2}}\quad\text{and}\quad\sin\Theta_W=\frac{g'}{\sqrt{g^2+g'^2}}\ . \tag{13}$$

The field equations for the Yang-Mills fields are ($\Box=-\partial_\nu\partial^\nu$) for B^μ

$$\Box B^\mu+\partial^\mu\partial_\nu B^\nu=-\tfrac{1}{4}gg'\varphi^2W^{3\mu}+\tfrac{1}{4}g'^2\varphi^2B^\mu-\frac{g'}{2}j^\mu_Y\ , \tag{14}$$

for $W^{a\mu}(a=1,2)$

$$\Box W^{a\mu}+g\varepsilon_{abc}W^{b\nu}\partial_\nu W^{c\mu}$$
$$=g^2(\tfrac{1}{4}\varphi^2W^{a\mu}-W^b_\nu W^{b\nu}W^{a\mu}+W^{a\nu}W^b_\nu W^{b\mu})-gj^{a\mu}\ , \tag{15}$$

and for $W^{3\mu}$

$$\Box W^{3\mu}+g\varepsilon_{3bc}W^{b\nu}\partial_\nu W^{c\mu}$$
$$=\tfrac{1}{4}g^2\varphi^2W^{3\mu}-\tfrac{1}{4}gg'\varphi^2B^\mu-g^2W^b_\nu W^{b\nu}W^{3\mu}$$
$$+g^2W^{3\nu}W^b_\nu W^{b\mu}-gj^{3\mu}\ . \tag{16}$$

Here the matter current densities are given by the equations

$$j^\mu_Y=\bar{L}\gamma^\mu YL+\bar{R}\gamma^\mu YR\ , \tag{17}$$

$$j^{a\mu}=\bar{L}\gamma^\mu\frac{\sigma^a}{2}L,\quad\text{where } a=1,2,3\ . \tag{18}$$

Accordingly, the Higgs field satisfies

$$\Box\varphi=(-\tfrac{1}{4}g^2W^a_\nu W^{a\nu}-\tfrac{1}{4}g'^2B_\nu B^\nu+gg'W^3_\nu B^\nu)\varphi$$
$$-\lambda v^2\varphi+\lambda\varphi^3+m\frac{\varphi}{v}(\bar{e}_Le_R+\text{H.c.})\ . \tag{19}$$

III. THE BOSON CONDENSATIONS

The effective potential of our model is given as the vacuum expectation value of the Lagrangian density

$$\mathcal{U}_{\text{ef}}=-\langle\mathcal{L}\rangle_{\tilde{0}}\ , \tag{20}$$

where the new vacuum state $|\tilde{0}\rangle$ (the ground state of the system configuration) is the Glauber coherent state (see the Appendix).

We decompose the initial fields into the quantum fluctuating fields and the classical condensates (see the Appendix):

$$\begin{aligned}W^a_\mu&=\tilde{W}^a_\mu+\omega^a_\mu\ ,\\ B_\mu&=\tilde{B}_\mu+b_\mu\ ,\\ \varphi&=\tilde{\varphi}+\delta\ .\end{aligned} \tag{21}$$

Here $\tilde{W}^a_\mu$, $\tilde{B}_\mu$, and $\tilde{\varphi}$ are quantum fields with a vanishing vacuum expectation value and ω^a_μ, b_μ, and δ are classical constant fields related to them. The appearance of these classical fields can be interpreted as a consequence of the condensation. We choose the coordinate system in which the condensations ω^a_μ and b_μ are to be of the following form:

$$\omega^a_\mu=\begin{cases}\omega^a_0=\sigma n^a\ ,\\ \omega^a_i=\vartheta\varepsilon_{aib}n^b\quad\text{and}\quad n^an^a=1\ ,\end{cases} \tag{22}$$

$$b_\mu=\begin{cases}b_0=\beta\ ,\\ b_i=0\ .\end{cases} \tag{23}$$

In Eq. (22) (n^a) plays the role of the unit vector in the adjoint representation of the Lie algebra su(2). It chooses a direction for the condensation. It is easy to see that

$$\omega^a_\mu\omega^{a\mu}=\sigma^2-2\vartheta^2\quad\text{and}\quad b_\mu b^\mu=\beta^2\ . \tag{24}$$

When we define the "electroweak magnetic field" as $\mathcal{B}^a_i=1/2\varepsilon_{ijk}F^a_{jk}$ and the "electroweak electric field" as $\mathcal{E}^a_i=F^a_{i0}$ then in the homogeneous case [$\vartheta=$const, $\sigma=$const$(n^a)=$const] we receive for $\vartheta\neq0$ the "electroweak magnetic field condensation $\langle\mathcal{B}^a_i\rangle_{\tilde{0}}$" and the "electroweak electric field condensation $\langle\mathcal{E}^a_i\rangle_{\tilde{0}}$" in the form

$$\langle\mathcal{B}^a_i\rangle_{\tilde{0}}=-g\vartheta^2n^in^a$$

and (25)

$$\langle\mathcal{E}^a_i\rangle_{\tilde{0}}=g\sigma\vartheta(\delta_{ai}-n^an^i)\ .$$

From Eqs. (1)–(12) and (20)–(21) we obtain the classical part of the effective potential for the "boson condensates induced by external matter sources" configuration (hereafter, we will call it the BCMS configuration):

$$\mathcal{U}_{\rm ef}=\tfrac{1}{4}g^2\varepsilon_{abc}\varepsilon_{ade}\omega^b_\mu\omega^{d\mu}\omega^c_\nu\omega^{e\nu}-\tfrac{1}{8}g^2\delta^2\omega^a_\mu\omega^{a\mu} -\frac{1}{4}gg'Y\delta^2\omega^3_\mu b^\mu-\frac{1}{8}Y^2g'^2\delta^2 b_\mu b^\mu+gJ^{a\mu}\omega^a_\mu +\frac{g'}{2}J^\mu_Y b_\mu+\tfrac{1}{4}\lambda(\delta^2-v^2)^2+\cdots\ , \tag{26}$$

where the mean matter current densities are as follows:

$$J^{a\mu}=\left\langle \bar{L}\gamma^\mu\frac{\sigma^a}{2}L\right\rangle_{\bar{0}}$$

and (27)

$$J^\mu_Y=(\langle \bar{L}\gamma^\mu YL\rangle_{\bar{0}}+\langle \bar{R}\gamma^\mu YR\rangle_{\bar{0}})\ .$$

The dots in Eq. (26) and afterwards signify some quantum corrections. The Lagrangian density given by the Eqs. (1)–(12) leads together with Eq. (21) to the classical massive Lagrangian density for boson fields

$$\mathcal{L}_{\rm mass}=-\tfrac{1}{2}g^2\varepsilon_{abc}\varepsilon_{ade}\omega^b_\mu\omega^{d\mu}\tilde{W}^c_\nu\tilde{W}^{e\nu}+\tfrac{1}{8}g^2\delta^2\tilde{W}^a_\mu\tilde{W}^{a\mu}-\tfrac{1}{4}gg'Y\delta^2\tilde{W}^3_\mu\tilde{B}^\mu+\tfrac{1}{8}Y^2g'^2\delta^2\tilde{B}_\mu\tilde{B}^\mu+\tfrac{1}{8}g^2\omega^a_\mu\omega^{a\mu}\tilde{\varphi}^2-\tfrac{1}{4}gg'Y\omega^3_\mu b^\mu\tilde{\varphi}^2 +\tfrac{1}{8}g'^2Y^2b_\mu b^\mu\tilde{\varphi}^2-\lambda\delta^2\tilde{\varphi}^2-\tfrac{1}{2}\lambda(\delta^2-v^2)\tilde{\varphi}^2\ . \tag{28}$$

We now assume that we are in the local rest coordinate system in which

$$J^0_Y=\rho_Y,\quad J^i_Y=0,\quad J^{a0}=\rho^a\ \text{ and }\ J^{ai}=0\ , \tag{29}$$

where ρ_Y and ρ^a are the matter charge densities related to $U_Y(1)$ and $SU_L(2)$, respectively. Using Eqs. (22)–(24) we can rewrite Eq. (26) as

$$\mathcal{U}_{\rm ef}(\vartheta,\sigma,\beta,\delta)=-g^2\sigma^2\vartheta^2+\tfrac{1}{2}g^2\vartheta^4-\tfrac{1}{8}g^2\delta^2(\sigma^2-2\vartheta^2) +\tfrac{1}{4}gg'\delta^2\beta\sigma n^3-\tfrac{1}{8}g'^2\delta^2\beta^2+g\rho^a n^a\sigma +\frac{g'}{2}\rho_Y\beta+\tfrac{1}{4}\lambda(\delta^2-v^2)^2+\cdots\ . \tag{30}$$

Now from the field equations, Eqs. (14)–(19), we can obtain four equations:

$$\partial_\vartheta\mathcal{U}_{\rm ef}=\partial_\sigma\mathcal{U}_{\rm ef}=\partial_\beta\mathcal{U}_{\rm ef}=\partial_\delta\mathcal{U}_{\rm ef}=0\ . \tag{31}$$

These equations lead to four algebraic equations for the condensations ϑ, σ, β, and δ:

$$(\tfrac{1}{2}\delta^2-2\sigma^2+2\vartheta^2)\vartheta=0\ , \tag{32}$$

$$-g(2\vartheta^2+\tfrac{1}{4}\delta^2)\sigma+\tfrac{1}{4}g'\delta^2\beta n^3+\rho^a n^a=0\ , \tag{33}$$

$$\tfrac{1}{2}(g\sigma n^3-g'\beta)\delta^2+\rho_Y=0\ , \tag{34}$$

$$[-\tfrac{1}{4}g^2(\sigma^2-2\vartheta^2)+\tfrac{1}{2}gg'\sigma\beta n^3-\tfrac{1}{4}g'^2\beta^2 +\lambda(\delta^2-v^2)]\delta=0\ . \tag{35}$$

Now we choose

$$(n^a)=(0,0,1)\ . \tag{36}$$

In this case we have for $\vartheta\neq0$ an "electroweak magnetic field condensation" different from zero $\langle\mathcal{B}^3_3\rangle_{\bar{0}}=-g\vartheta^2$ pointed in the x^3 spatial direction and the "electroweak electric field condensations" $\langle\mathcal{E}^1_1\rangle_{\bar{0}}=\langle\mathcal{E}^2_2\rangle_{\bar{0}}=g\sigma\vartheta$ pointed in the x^1 and x^2 spatial directions, respectively.

Using Eqs. (28), (22)–(24), and Eq. (36) we receive in the classical regime the square masses of the boson fields as follows:

$$m^2_{\tilde{W}^{1,2}}=g^2(\tfrac{1}{4}\delta^2-\sigma^2+\vartheta^2)\ , \tag{37}$$

$$m^2_{\tilde{W}^3}=g^2(\tfrac{1}{4}\delta^2+2\vartheta^2)\ , \tag{38}$$

$$m^2_{\tilde{B}}=\tfrac{1}{4}g'^2\delta^2\ , \tag{39}$$

$$\tfrac{1}{2}m^2_{\tilde{\varphi}}=\lambda\delta^2+\tfrac{1}{2}\lambda(\delta^2-v^2)-\tfrac{1}{8}g^2(\sigma^2-2\vartheta^2) +\tfrac{1}{4}gg'\sigma\beta n^3-\tfrac{1}{8}g'^2\beta^2\ . \tag{40}$$

Let us perform for $\delta\neq0$ the "rotation" of the $\tilde{W}^3_\mu$ and $\tilde{B}_\mu$ fields to the physical fields $\tilde{Z}_\mu$ and $\tilde{A}_\mu$

$$\begin{bmatrix}\tilde{Z}_\mu\\ \tilde{A}_\mu\end{bmatrix}=\begin{bmatrix}\cos\Theta & -\sin\Theta\\ \sin\Theta & \cos\Theta\end{bmatrix}\begin{bmatrix}\tilde{W}^3_\mu\\ \tilde{B}_\mu\end{bmatrix} \tag{41}$$

and at the same time the "rotation" of the σ and β condensations to their counterparts ζ and α as well as the "rotation" of the charge densities ρ^3 and ρ_Y to the physical ρ_Z and ρ_Q

$$\begin{bmatrix}\zeta\\ \alpha\end{bmatrix}=\begin{bmatrix}\cos\Theta & -\sin\Theta\\ \sin\Theta & \cos\Theta\end{bmatrix}\begin{bmatrix}\sigma\\ \beta\end{bmatrix}, \tag{42}$$

$$\begin{bmatrix}(g/\cos\Theta)\rho_Z\\ (g\sin\Theta)\rho_Q\end{bmatrix}=\begin{bmatrix}\cos\Theta & -\sin\Theta\\ \sin\Theta & \cos\Theta\end{bmatrix}\begin{bmatrix}(g)\rho^a n^a\\ (g'/2)\rho_Y\end{bmatrix}. \tag{43}$$

Now using Eqs. (37)–(40) and defining the $\tilde{W}^\pm$ fields as $\tilde{W}^\pm=(\tilde{W}^1\mp i\tilde{W}^2)/\sqrt{2}$ we can rewrite the square masses of the physical boson fields as follows:

$$m^2_{\tilde{W}^\pm}=g^2[\tfrac{1}{4}\delta^2-(\zeta\cos\Theta+\alpha\sin\Theta)^2+\vartheta^2]\ , \tag{44}$$

$$m^2_{\tilde{Z}}=\tfrac{1}{2}[m^2_{Z\,\rm SM}+2g^2\vartheta^2 +\sqrt{(m^2_{Z\,\rm SM}+2g^2\vartheta^2)^2-2(gg'\delta\vartheta)^2}\,]\ , \tag{45}$$

$$m^2_{\tilde{A}}=\tfrac{1}{2}[m^2_{Z\,\rm SM}+2g^2\vartheta^2 -\sqrt{(m^2_{Z\,\rm SM}+2g^2\vartheta^2)^2-2(gg'\delta\vartheta)^2}\,]\ , \tag{46}$$

$$\tfrac{1}{2}m_{\tilde{\varphi}}^2=\lambda\delta^2+\frac{1}{2}\lambda(\delta^2-v^2)-\frac{1}{2\delta^2}(m_{\tilde{Z}}^2\zeta^2-m_{\tilde{A}}^2\alpha^2)+g^2\left[\frac{1}{\delta^2}(\zeta\cos\Theta+\alpha\sin\Theta)^2+\frac{1}{4}\right]\vartheta^2\ , \tag{47}$$

where $m_{Z\,\mathrm{SM}}^2$ is the standard counterpart for the boson $\tilde{Z}^\mu$ square mass which is equal to

$$m_{Z\,\mathrm{SM}}^2=\tfrac{1}{4}(g^2+g'^2)\delta^2\ . \tag{48}$$

It is useful to write the relations between the matter weak isotopic charge density ρ^3 [see Eqs. (29) and (27)], the matter weak hypercharge density ρ_Y [see Eqs. (29) and (27)], the standard electric charge density $\rho_{Q\,\mathrm{SM}}$, the standard weak charge density $\rho_{Z\,\mathrm{SM}}$ and their generalizations in our model; that is the electric charge density ρ_Q and weak charge density ρ_Z:

$$\rho_Q=\rho_{Q\,\mathrm{SM}}+\frac{1}{2}\left[\frac{g'}{g}\cot\Theta-1\right]\rho_Y\ , \tag{49}$$

$$\rho_Z=\rho^3-\rho_Q\sin^2\Theta\ , \tag{50}$$

$$\rho_{Q\,\mathrm{SM}}=\rho^3+\tfrac{1}{2}\rho_Y$$

and (51)

$$\rho_{Z\,\mathrm{SM}}=\rho^3-\rho_{Q\,\mathrm{SM}}\sin^2\Theta_W\ .$$

Here the Θ angle is the modified mixing angle which is given by the formula

$$\tan\Theta=\left\{\frac{-[1+8(\vartheta/\delta)^2]g^2+g'^2}{2gg'}+\left[\left[\frac{[1+8(\vartheta/\delta)^2]g^2-g'^2}{2gg'}\right]^2+1\right]^{1/2}\right\}. \tag{52}$$

When $\vartheta\to 0$ then it is not difficult to ascertain that out of Eqs. (49)–(52) the well-known GSW results emerge.

IV. DISCUSSION

The calculations below are done for the boson condensations in extrema of the effective potential $\mathcal{U}_{\mathrm{ef}}$ unless it is stated differently. From Eqs. (32) and (37) we can see that the solutions of Eqs. (32)–(35) for boson condensations in the extrema of the effective potential $\mathcal{U}_{\mathrm{ef}}$ split into two major cases. The first one for $\vartheta\neq 0$ gives us the phase with $m_{\tilde{W}^\pm}^2=0$. The second case for $\vartheta=0$ gives us the phase with $m_{\tilde{W}^\pm}^2\neq 0$ which depends on the values of condensations (with $m_{\tilde{W}^\pm}^2=0$ as the limit of the stability for this phase). Each of these two splits then into the $\delta\neq 0$ and $\delta=0$ cases. We chose in our numerical calculations the standard boson W^μ mass $m_{W\,\mathrm{SM}}=80.13$ GeV, the standard boson Z^μ mass $m_{Z\,\mathrm{SM}}=91.187$ GeV, and the fine structure constant = 1/137.

A. Condensations $\vartheta\neq 0$ and $\delta\neq 0$

Equations (32)–(35) can be now rewritten as follows:

$$\sigma=\frac{1}{2g\vartheta^2}\rho_{Q\,\mathrm{SM}}\ , \tag{53}$$

$$\beta=\frac{1}{g'}\left[g\sigma n^3+2\frac{\rho_Y}{\delta^2}\right], \tag{54}$$

$$\vartheta^6+\frac{1}{4}\delta^2\vartheta^4-\frac{1}{4g^2}\ \rho_{Q\,\mathrm{SM}}^2=0\ , \tag{55}$$

$$\delta^6+\left[\frac{g^2}{2\lambda}\vartheta^2-v^2\right]\delta^4-\frac{1}{\lambda}\rho_Y^2=0\ . \tag{56}$$

From Eq. (55) we see that the condensation $\vartheta\neq 0$ only when $\rho_{Q\,\mathrm{SM}}\neq 0$.

When we notice that the relation between the weak hypercharge quantum number Y and the electromagnetic charge quantum number Q can be written for matter fields in the form $Q=pY/2$ where suitable p ($p\neq 0$) are given in Table I, then the relation between the weak hypercharge density ρ_Y and the standard electromagnetic charge density $\rho_{Q\mathrm{SM}}$ can be written in the similar form

$$\rho_{Q\,\mathrm{SM}}=p\frac{\rho_Y}{2}\ . \tag{57}$$

After using Eq. (57) we solved numerically Eqs. (53)–(56) and we obtained the condensations squared ϑ^2 and δ^2 as functions of $\rho_{Q\,\mathrm{SM}}$ with p as a parameter. Different values for p (see Table I) represent different matter fields which could be the sources of charge densities.

The results of solving Eqs. (53)–(56) for the α and ζ condensations [see Eq. (42)] and the ϑ and δ condensations are shown in Figs. 1–4.

Now Eq. (21) has the form

$$\begin{aligned}&W_{0,3}^\pm=\tilde{W}_{0,3}^\pm,\quad W_1^\pm=\tilde{W}_1^\pm\pm i\vartheta/\sqrt{2},\quad W_2^\pm=\tilde{W}_2^\pm+\vartheta/\sqrt{2}\ ,\\&Z_i=\tilde{Z}_i,\quad Z_0=\tilde{Z}_0+\zeta,\quad \text{where }\ \zeta=\sigma\cos\Theta-\beta\sin\Theta\ ,\\&A_i=\tilde{A}_i,\quad A_0=\tilde{A}_0+\alpha,\quad \text{where }\ \alpha=\sigma\sin\Theta+\beta\cos\Theta\ ,\\&\varphi=\tilde{\varphi}+\delta\ .\end{aligned} \tag{58}$$

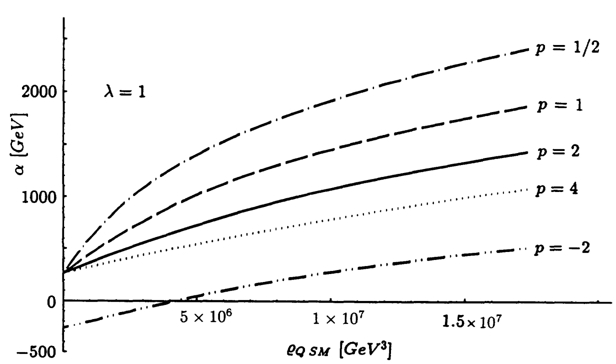

FIG. 1. The α condensation of the A_0 gauge boson fields as the function of the standard electric charge density $\rho_{Q\,\mathrm{SM}}$ ($\vartheta\neq 0,\delta\neq 0$).

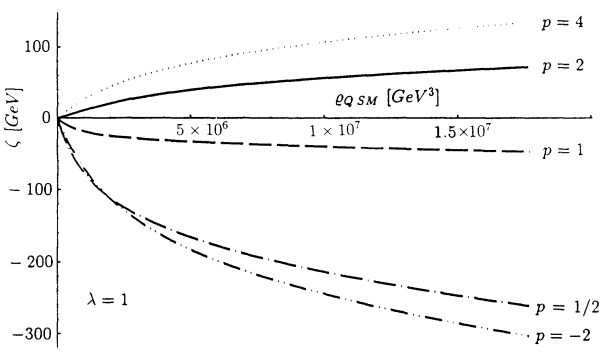

FIG. 2. The ζ condensation of the Z_0 gauge boson fields as the function of the standard electric charge density $\rho_{Q\,\mathrm{SM}}(\vartheta\neq0,\delta\neq0)$.

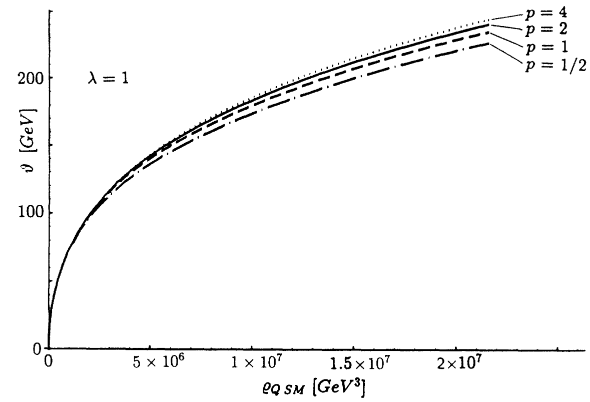

FIG. 3. The ϑ condensation of the $W^{a=1}_{i=2}$ and $W^{a=2}_{i=1}$ gauge boson fields as the function of the standard electric charge density $\rho_{Q\,\mathrm{SM}}(\vartheta\neq0,\delta\neq0)$.

The masses of the $\tilde{\varphi}$, $\tilde{Z}$, and $\tilde{A}$ were calculated according to Eqs. (45)–(46) and (40) and the appropriate results are shown in Figs. 4–6. The masses of the $\tilde{W}^{\pm}$ fields are, according to Eqs. (32) and (37) (for the $\vartheta\neq0$ phase), equal to 0.

The results for the ratio $\sin\Theta/\sin\Theta_W$ [see Eq. (52)] and the physical charge density ρ_Q [see Eq. (49)] for boson condensates given by Eqs. (53)–(56) as functions of $\rho_{Q\,\mathrm{SM}}$ are presented in Figs. 7 and 8, respectively.

In all the figures the curves for different values of p coverage for relatively small values of $\rho_{Q\,\mathrm{SM}}$ (i.e., for $\rho_{Q\,\mathrm{SM}}$ in the range up to values approximately 10^3 times bigger than these which correspond to matter densities in

TABLE I. Quantum numbers in the $SU_L(2)\times U_Y(1)$ electroweak theory.

	Weak isotopic charge I^3	Weak hypercharge Y	Electric charge Q $Q=I^3+Y/2$	$p=2Q/Y$
Quarks				
u_L	$\frac{1}{2}$	$\frac{1}{3}$	$\frac{2}{3}$	4
d_L	$-\frac{1}{2}$	$\frac{1}{3}$	$-\frac{1}{3}$	-2
u_R	0	$\frac{4}{3}$	$\frac{2}{3}$	1
d_R	0	$-\frac{2}{3}$	$-\frac{1}{3}$	1
Leptons				
ν_L	$\frac{1}{2}$	-1	0	0
e_L	$-\frac{1}{2}$	-1	-1	2
e_R	0	-2	-1	1
Gauge bosons				
W^+	1	0	1	
W^3	0	0	0	
W^-	-1	0	-1	
B	0	0	0	
Higgs boson				
H^+	$\frac{1}{2}$	1	1	2
H^0	$-\frac{1}{2}$	1	0	0
Quark configurations				
$(u_L d_L d_L)$	$-\frac{1}{2}$	1	0	0
$(u_L u_L d_L)$	$\frac{1}{2}$	1	1	2
$(3\times u_L-3\times d_L)$	0	2	1	1
$(5\times u_L-7\times d_L)$	-1	4	1	$\frac{1}{2}$

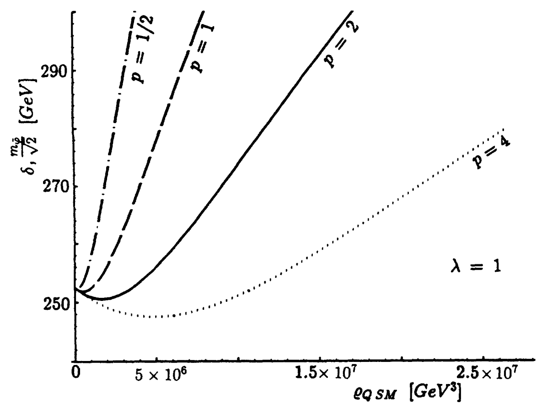

FIG. 4. The δ condensation of the φ Higgs boson field and the mass $m_{\tilde{\varphi}}/\sqrt{2}$ of the $\tilde{\varphi}$ Higgs boson field as the function of the standard electric charge density $\rho_{Q\,\mathrm{SM}}(\vartheta\neq 0,\delta\neq 0)$.

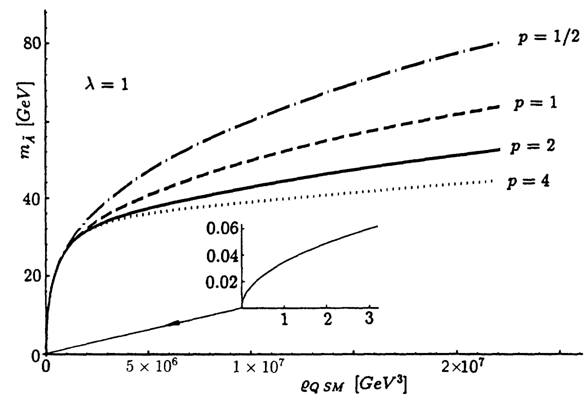

FIG. 6. The mass $m_{\tilde{A}}$ of the $\tilde{A}^{\mu}$ gauge boson fields as the function of the standard electric charge density $\rho_{Q\,\mathrm{SM}}$ $(\vartheta\neq 0,\delta\neq 0)$. The region for the standard electric charge density $\rho_{Q\,\mathrm{SM}}$ which is in the range up to values approximately 10^3 times bigger than those for nucleon matter is indicated by the arrow.

nucleon matter). In that range of values for $\rho_{Q\,\mathrm{SM}}$ we have also that $\rho_Q \approx \rho_{Q\,\mathrm{SM}}$ (see Fig. 8). For these reasons the following calculations in that regime were done for $p=1$ (for other $p\neq 0$ in Table I it would be the same).

The minimal energy density of the BCMS configuration

$$\mathcal{E}_{\min}(\rho_{Q\,\mathrm{SM}})=\mathcal{U}_{\mathrm{ef}}(\vartheta\neq 0,\delta\neq 0)$$

[see Eq. (30)] for boson condensates given by Eqs. (53)–(56) as functions of $\rho_{Q\,\mathrm{SM}}$ is presented in Fig. 9. For big charge density $\rho_{Q\,\mathrm{SM}}$ (i.e., for $\rho_{Q\,\mathrm{SM}}$ which corresponds to matter densities approximately 10^3 times bigger than those in nucleon matter) the minimal energy density $\mathcal{E}_{\min}(\rho_{Q\,\mathrm{SM}})$ is extremely big increasing rapidly with $\rho_{Q\,\mathrm{SM}}$ (for example, $\mathcal{E}_{\min}\approx 2.9\ 10^{178}\ \mathrm{GeV}^4$ for $\rho_{Q\,\mathrm{SM}}\approx 1.3\ 10^7\ \mathrm{GeV}^3$).

It is very interesting that there emerges a subtle structure when we investigate more carefully the function $\mathcal{E}_{\min}(\rho_{Q\,\mathrm{SM}})$. It appears a "stable" (BCMS) configuration of charge density with $\rho_{Q\,\mathrm{SM}}\neq 0$ (see Fig. 9) different from that for the standard model [with $\rho_{Q\,\mathrm{SM}}=0$ and $\mathcal{E}_{\min}(0)=0$]. The numerical calculations for the value of the local minimum of the function $\mathcal{E}_{\min}(\rho_{Q\,\mathrm{SM}})$ reveal little dependence on the λ parameter of the Higgs potential (see Fig. 9) and the results are as follows:

$$\mathcal{E}_{\min}(\rho_{Q\,\mathrm{SM}})\approx(2.5811\ \mathrm{GeV})^4$$

for (59)

$$\rho_{Q\,\mathrm{SM}}\approx 0.5539\ \mathrm{GeV}^3\ .$$

This charged BCMS configuration is separated from the uncharged standard model configuration by a high barrier $\Delta\mathcal{E}_{\min}$ which depends on the λ parameter (see Fig. 9). For example, when $\lambda=1$ then $\Delta\mathcal{E}_{\min}\approx(180\ \mathrm{GeV})^4$. It is not difficult to ascertain that $\mathcal{E}_{\min}\to 0$ as $\rho_{Q\,\mathrm{SM}}\to 0$ for all considered values of $\lambda>0$ and $p\neq 0$ (see Table I).

When we notice that the mass of an electric charged BCMS configuration is

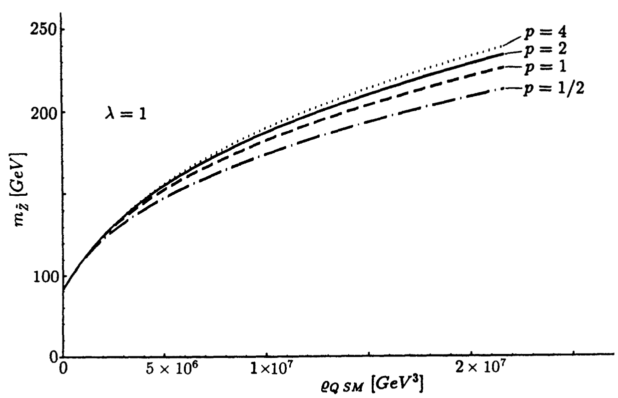

FIG. 5. The mass $m_{\tilde{Z}}$ of the $\tilde{Z}^{\mu}$ gauge boson fields as the function of the standard electric charge density $\rho_{Q\,\mathrm{SM}}$ $(\vartheta\neq 0,\delta\neq 0)$.

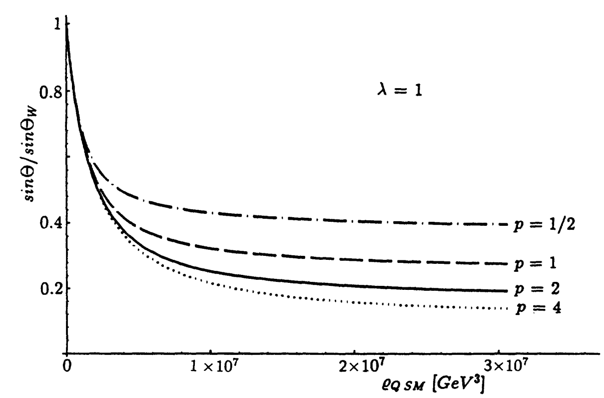

FIG. 7. The ratio $\sin\Theta/\sin\Theta_W$ (the Θ angle is the modified mixing angle [see Eq. (52)]) as the function of the standard electric charge density $\rho_{Q\,\mathrm{SM}}$ $(\vartheta\neq 0,\delta\neq 0)$.

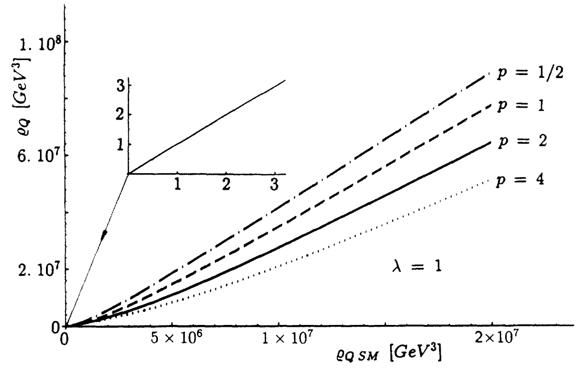

FIG. 8. The physical electric charge density ρ_Q [see Eq. (49)] as the function of the standard electric charge density $\rho_{Q\,\mathrm{SM}}$ ($\vartheta\neq 0,\delta\neq 0$). The region for the standard electric charge density $\rho_{Q\,\mathrm{SM}}$ which is in the range up to values approximately 10^3 times bigger than those for nucleon matter is indicated by the arrow.

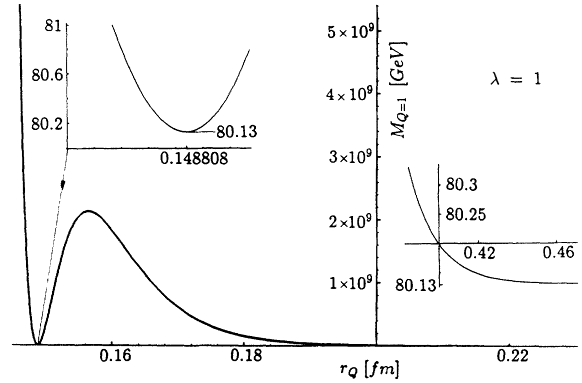

FIG. 10. The mass $M_{Q=1}$ of the BCMS configuration (with the electric charge $Q=1$) for boson condensates given by Eqs. (53)–(56) as the function of the mean electric charge radius r_Q ($\vartheta\neq 0,\delta\neq 0$ and for all values of $p\neq 0$ from Table I). The region in the vicinity of the "stable" BCMS configuration is indicated by the arrow.

$$M_Q=\tfrac{4}{3}\pi r_Q^3\mathcal{E}_{\min}(\rho_{Q\,\mathrm{SM}})\ ,$$

where r_Q is the "mean electric charge radius" of a BCMS configuration and that the electric charge $Q=\frac{4}{3}\pi r_Q^3\rho_{Q\,\mathrm{SM}}$. Then from Eqs. (30) and (53)–(56) we receive $M_Q\to Qgv/2=Q\times 80.13$ GeV as $\rho_{Q\,\mathrm{SM}}\to 0$ for all considered values of $\lambda>0$ and $p\neq 0$ and (see Table I). The function $M_{Q=1}(r_Q)$ is presented in Fig. 10. Configurations of this phase lie only on the M_Q-r_Q curve. For example, a droplet of the new phase with charge $Q=1$ and described by Eq. (59) will have the "mean charge radius" $r_Q=0.149$ fm (in comparison for a proton $r_Q\approx 0.805$ fm) and the mass $M_{Q=1}\approx 80.13$ GeV. (In Figs. 3–7 and Figs. 9 and 10 the curves for p and $-p$ cover.)

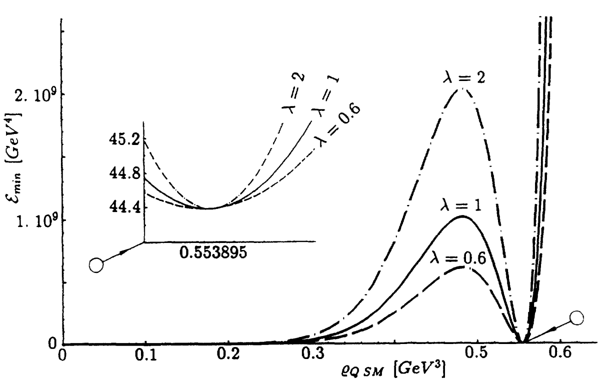

FIG. 9. The minimal energy density of the BCMS configuration $\mathcal{E}_{\min}(\rho_{Q\,\mathrm{SM}})=\mathcal{U}_{\mathrm{ef}}(\vartheta\neq 0,\delta\neq 0)$ [see Eq. (30)] for boson condensates given by Eqs. (53)–(56) (for all values of $p\neq 0$ from Table I) as the function of the standard electric charge density $\rho_{Q\,\mathrm{SM}}$ which is in the range up to values approximately 10^3 times bigger than those for nucleon matter. The region in the vicinity of the "stable" BCMS configuration [see Eq. (59)] is indicated by the arrow.

B. Condensations $\vartheta\neq 0$ and $\delta=0$

In that case Eqs. (32)–(35) lead to

$$\rho_Y=0\ , \tag{60}$$

$$\sigma=\sqrt[3]{\rho^3/2g}\quad\text{and}\quad\vartheta=\pm\sigma\ . \tag{61}$$

These equations together with Eqs. (37)–(40) give us the square masses of boson fields as follows:

$$m^2_{\tilde W^{1,2}}=m^2_{\tilde B}=0\ , \tag{62}$$

$$m^2_{\tilde W^3}=2^{1/3}g^{4/3}(\rho^3)^{2/3}\ , \tag{63}$$

$$\tfrac{1}{2}m^2_{\tilde\varphi}=-\frac{\lambda}{2}v^2+2^{-11/3}g^{4/3}(\rho^3)^{2/3}\ , \tag{64}$$

and the energy density $\mathcal{E}(\rho^3)$ which is equal to

$$\mathcal{E}(\rho^3)=\mathcal{U}_{\mathrm{ef}}(\vartheta\neq 0,\delta=0)$$

$$=32^{-7/3}g^{2/3}(\rho^3)^{4/3}+\frac{\lambda}{4}v^4\ . \tag{65}$$

The quantum numbers of the matter fields in the standard model (see Table I) give as a result of Eq. (60) the implication

$$\rho_Y=0\Longrightarrow\rho^3=0\ . \tag{66}$$

So we see that in that case Eqs. (60)–(65) reproduce the GSW model for $\delta=0$ [the unbroken $SU_L(2)\times U_Y(1)$ high-temperature phase].

C. Condensations $\vartheta=0$ and $\delta\neq 0$

Using Eqs. (42)–(43) we can rewrite the effective potential $\mathcal{U}_{\mathrm{ef}}$ given by Eq. (30) in a very simple form:

$$\mathcal{U}_{\mathrm{ef}}(\zeta,\alpha,\delta)=-\tfrac{1}{8}(g^2+g'^2)\delta^2\zeta^2+\rho_{Z\,\mathrm{SM}}\zeta+\rho_{Q\,\mathrm{SM}}\alpha$$

$$+\tfrac{1}{4}\lambda(\delta^2-v^2)^2\ . \tag{67}$$

This potential together with the equations $\partial_\alpha \mathcal{U}_{ef}=\partial_\zeta \mathcal{U}_{ef}=\partial_\delta \mathcal{U}_{ef}=0$ yield, respectively, the equations for the charge density $\rho_{Q\,SM}$ and the condensations ζ and δ [instead of these calculations we could use Eqs. (32)–(35) and (42)–(43)]:

$$\rho_{Q\,SM}=0\ , \tag{68}$$

$$\tfrac{1}{4}\sqrt{g^2+g'^2}\delta^2\zeta=\rho_{Z\,SM}\ , \tag{69}$$

and

$$\lambda(\delta^2-v^2)-\tfrac{1}{4}(g^2+g'^2)\zeta^2=0\ . \tag{70}$$

The nonzero value of the weak charge density $\rho_{Z\,SM}$ leads inevitably to the nonzero ζ condensation and to the forced symmetry breaking.

Combining Eqs. (68) and (69) with (67) we obtain the classical counterpart of the GSW effective potential for $\vartheta=0$ (see Fig. 11)

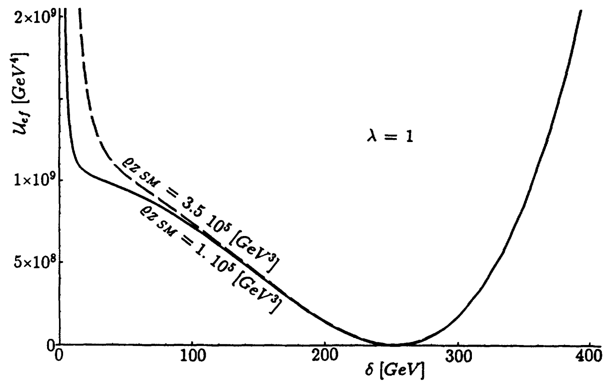

FIG. 11. The classical effective potential $\mathcal{U}_{ef}(\delta;\vartheta=0,\rho_{Q\,SM}=0)$ as the function of the δ Higgs field condensation ($\lambda=1$).

$$\mathcal{U}_{ef}(\delta;\vartheta=0,\rho_{Q\,SM}=0)=\frac{2\rho_Z^2}{\delta^2}+\frac{1}{4}\lambda(\delta^2-v^2)^2\ . \tag{71}$$

The solution of Eqs. (69) and (70) leads to

$$\zeta(\rho_{Z\,SM})=\frac{2\lambda^{1/3}}{\sqrt{g^2+g'^2}}\left\{\left[\rho_{ZSM}+\left[\rho_{ZSM}^2+\frac{\lambda v^6}{27}\right]^{1/2}\right]^{1/3}+\left[\rho_{ZSM}-\left[\rho_{ZSM}^2+\frac{\lambda v^6}{27}\right]^{1/2}\right]^{1/3}\right\}\geq 0 \tag{72}$$

and

$$\delta^2(\rho_{Z\,SM})=\frac{4\rho_{Z\,SM}}{\sqrt{g^2+g'^2}\zeta}\ , \tag{73}$$

where ζ and δ^2 are only the functions of $\rho_{Z\,SM}$. It is not difficult to ascertain that $\delta\to v$ and $\zeta\to 0$ as $\rho_{Z\,SM}\to 0$, and the well-known GSW broken low temperature phase ($\delta=v$) with $U_Q(1)$ symmetry emerges.

Using Eqs. (22)–(23) and (41)–(42) we can rewrite Eq. (21) for the physical field A_μ in the form

$$A_\mu=\tilde{A}_\mu+a_\mu,\quad \text{where}\quad a_\mu=(\alpha,0,0,0)\ . \tag{74}$$

Let us notice from Eqs. (68)–(71) that α is not a dynamical parameter so Eq. (74) gains the gauge transformation interpretation. The a_μ condensate corresponds to the nonphysical degree of freedom (unphysical photon) and it can be removed by the appropriate gauge transformation. So the requirement that the $U_Q(1)$ group has to survive untouched during the symmetry breaking gives us

$$\alpha=\sigma\sin\Theta_W+\beta\cos\Theta_W=0\ . \tag{75}$$

We have the result that when the $\tilde{W}_\mu^{1,2}$ and $\tilde{Z}_\mu$ fields acquire nonzero masses then the condensations in Eq. (21) can be rewritten as follows:

$$\begin{aligned}&W_\mu^{1,2}=\tilde{W}_\mu^{1,2},\quad W_i^3=\tilde{W}_i^3\ ,\\&W_0^3=\tilde{W}_0^3-\beta\cot\Theta_W\ ,\\&B_0=\tilde{B}_0+\beta\ ,\\&B_i=\tilde{B}_i\ ,\\&\varphi=\tilde{\varphi}+\delta\ ,\end{aligned} \tag{76}$$

or in the physical fields

$$\begin{aligned}&W_\mu^{\pm}=\tilde{W}_\mu^{\pm},\quad Z_i=\tilde{Z}_i\ ,\\&Z_0=\tilde{Z}_0+\zeta\quad \text{where}\quad \zeta=-\frac{1}{\sin\Theta_W}\beta\ ,\\&A_\mu=\tilde{A}_\mu\ ,\\&\varphi=\tilde{\varphi}+\delta\ .\end{aligned} \tag{77}$$

The appearance of the boson condensates strongly influences the masses of the fields in the model and from Eqs. (44)–(47) ($\vartheta=0$ and $\alpha=0$) we obtain (see Figs. 12–14)

$$m_{\tilde{W}^\pm}^2=\tfrac{1}{4}g^2\delta^2-g^2\zeta^2\cos^2\Theta_W\ , \tag{78}$$

$$m_{\tilde{Z}}^2=\tfrac{1}{4}(g^2+g'^2)\delta^2\ , \tag{79}$$

$$m_{\tilde{A}}^2=0\ , \tag{80}$$

$$\tfrac{1}{2}m_{\tilde{\varphi}}^2=\lambda\delta^2+\tfrac{1}{2}\lambda(\delta^2-v^2)-\tfrac{1}{8}(g^2+g'^2)\zeta^2\ . \tag{81}$$

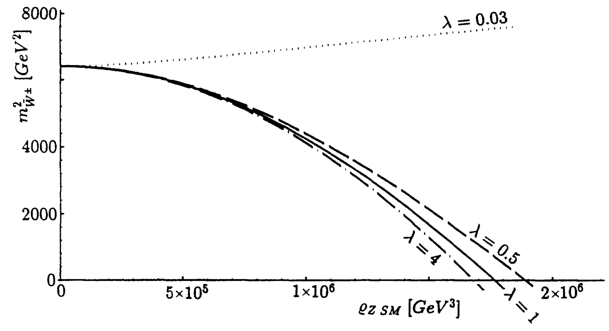

FIG. 12. The square mass $m^2_{\tilde{W}^\pm}$ of the $\tilde{W}^\pm_\mu$ gauge boson fields [see Eq. (78)] as the function of the standard weak charge density $\rho_{Z\,\mathrm{SM}}(\vartheta=0,\delta\neq 0,\rho_{Q\,\mathrm{SM}}=0)$.

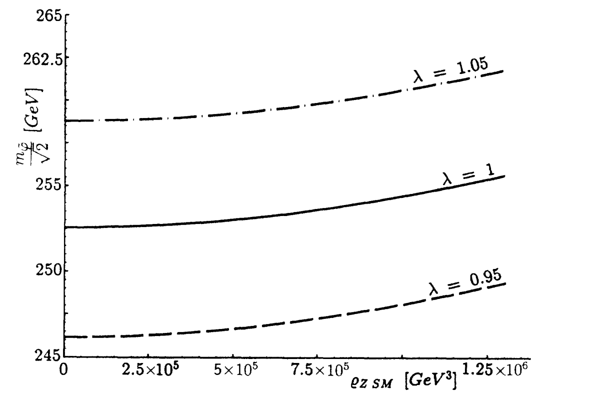

FIG. 14. The mass $m_{\tilde{\varphi}}/\sqrt{2}$ of the $\tilde{\varphi}$ Higgs boson field [see Eq. (81)] as the function of the standard weak charge density $\rho_{Z\,\mathrm{SM}}\,(\vartheta=0,\delta\neq 0,\rho_{Q\,\mathrm{SM}}=0)$.

The minimal energy density of the BCMS configuration $\mathscr{E}_{\min}(\rho_{Z\,\mathrm{SM}})=\mathscr{U}_{\mathrm{ef}}(\vartheta=0,\delta\neq 0)$ is as follows (see Fig. 15):

$$\mathscr{E}_{\min}(\rho_{Z\,\mathrm{SM}})=\tfrac{1}{2}\zeta\sqrt{g^2+g'^2}\rho_{Z\,\mathrm{SM}}+\frac{1}{64\lambda}(g^2+g'^2)^2\zeta^4\ . \tag{82}$$

From Eq. (78) it is clear that the boson condensation $\zeta>0$ leads to the instability in the $\tilde{W}^\pm_\mu$ sector if

$$\zeta^3>\frac{\sqrt{g^2+g'^2}}{g^2}\rho_{Z\,\mathrm{SM}}\ . \tag{83}$$

When the equality

$$\zeta^3=\rho_{Z\,\mathrm{SM}}\sqrt{g^2+g'^2}/g^2$$

is taken into account we obtain the relationship between $\lambda_{\max}$ and $\rho_{Z\max}$ where $\lambda_{\max}$ is the value of λ and $\rho_{Z\max}$ is the value of $\rho_{Z\,\mathrm{SM}}$ for which we have $m^2_{\tilde{W}^\pm}=0$ (see Fig. 16). The region of possible configurations of this phase is on and below the $\lambda_{\max}$-$\rho_{Z\max}$ curve.

For weak charge densities

$$\rho_{Z\,\mathrm{SM}}\leq gv^3/(8\cos^2\Theta_W)\approx 1.655\times 10^6\ \mathrm{GeV}^3$$

this phase is stable for an arbitrary λ (see Fig. 16). For values bigger than 1.655×10^6 GeV^3 this phase for given λ will be destabilized at certain value of $\rho_{Z\,\mathrm{SM}}=\rho_{Z\max}$ and the system could reach the charged ($\rho_{Q\,\mathrm{SM}}\neq 0$) stable phase with $\vartheta\neq 0$. For $\lambda<g^2/(16\cos^4\Theta_W)\approx 0.0422$ the phase is stable for all values of weak charge density $\rho_{Z\,\mathrm{SM}}$ (see Fig. 16).

We can also examine the mass

$$M_{I^3}=\tfrac{4}{3}\pi r_Z^3\mathscr{E}_{\min}(\rho_{Z\,\mathrm{SM}})$$

of a BCMS configuration with nonzero weak charge density. Here r_Z is the "mean weak charge radius" of this configuration which has the weak isotopic charge $I^3=\frac{4}{3}\pi r_Z^3\rho_{Z\,\mathrm{SM}}$. We receive the upper (according to the stability of this phase within the $\tilde{W}^\pm$ sector) limit $M_{I^3,\max}$ for the value of the mass M_{I^3} with the region of possible configurations of this phase which lie on and below the $M_{I^3,\max}$-$\lambda_{\max}$ curve (see Fig. 17).

The experimental knowledge of a BCMS configuration with the upper possible mass $M_{I^3,\max}$ could give us from this curve the value of λ and eventually the value of the

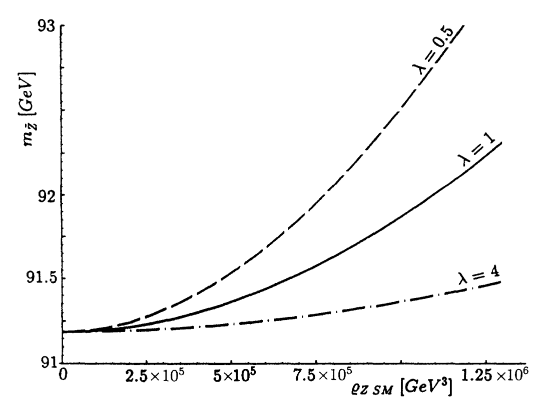

FIG. 13. The mass $m_{\tilde{Z}}$ of the $\tilde{Z}^\mu$ gauge boson fields [see Eq. (79)] as the function of the standard weak charge density $\rho_{Z\,\mathrm{SM}}$ $(\vartheta=0,\delta\neq 0,\rho_{Q\,\mathrm{SM}}=0)$.

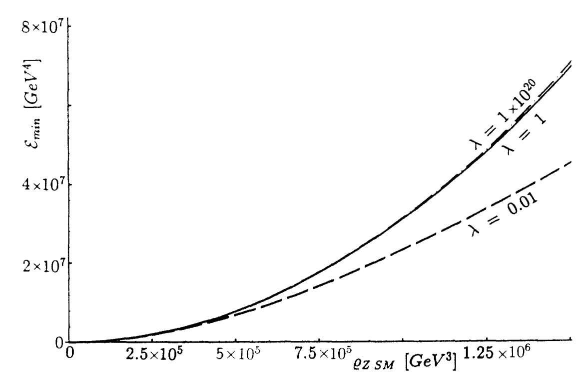

FIG. 15. The minimal energy density of the BCMS configuration [see Eq. (82)] $\mathscr{E}_{\min}(\rho_{Z\,\mathrm{SM}})=\mathscr{U}_{\mathrm{ef}}(\vartheta=0,\delta\neq 0,\rho_{Q\,\mathrm{SM}}=0)$ [see Eq. (82)].

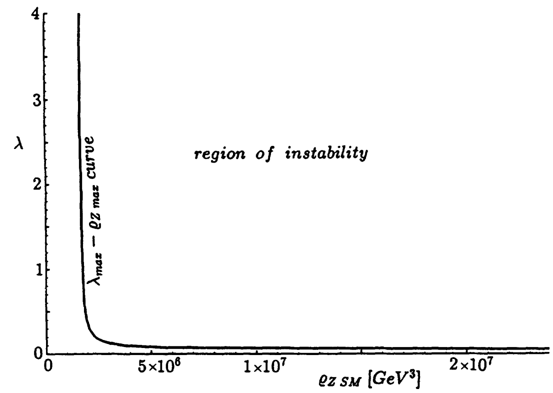

FIG. 16. The partition of the ($\lambda\rho_{Z\,\mathrm{SM}}$) plane into two regions of stability and instability of the phase with $\vartheta=0$ and $\delta\neq0$. The region of possible configurations of this phase is on and below the $\lambda_{\max}$-$\rho_{Z\,\max}$ curve where $\lambda_{\max}$ is the value of λ and $\rho_{Z\,\max}$ is the value of $\rho_{Z\,\mathrm{SM}}$ for which we have $m^2_{\tilde{W}^\pm}=0$.

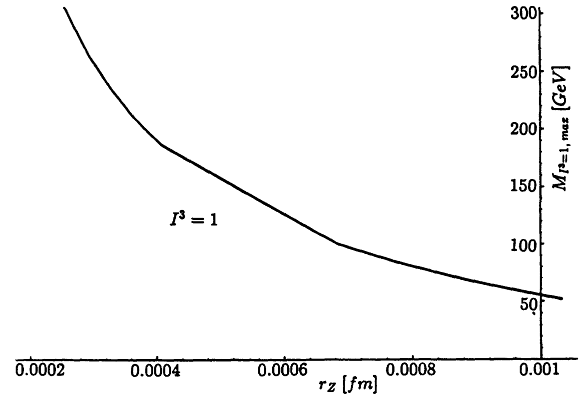

FIG. 18. The upper (according to the stability of the phase with $\vartheta=0$ and $\delta\neq0$ in the $\tilde{W}^\pm$ sector) mass $M_{I^3=1,\max}(r_Z)$ of a BCMS configuration (with the weak isotopic charge $I^3=1$) as the function of the "mean weak charge radius" r_Z of this configuration.

Higgs boson mass. The function $M_{I^3=1}(r_Z)$ is also presented on Fig. 18.

D. Condensations $\vartheta=0$ and $\delta=0$

In that case the GSW model for $\delta=0$ is reproduced [the unbroken $\mathrm{SU}_L(2)\times\mathrm{U}_Y(1)$ high-temperature phase].

V. CONCLUSIONS

In this paper the boson Higgs field and gauge field condensations in the presence of the external matter sources (BCMS configurations) were examined. Apart from the Higgs boson condensation ($\delta\neq0$) the rich phases structure is observed. In general, we notice two physically different phases. The first with $\vartheta\neq0$ and $\delta\neq0$ appears only when the charge density $\rho_{Q\,\mathrm{SM}}\neq0$. In this phase the $\tilde{W}^\pm_\mu$ bosons are massless while the photons $\tilde{A}_\mu$ and bosons $\tilde{Z}_\mu$ are massive. We observe very deep energy density minimum (see Fig. 9) with $\mathcal{E}_{\min}\approx44.382$ GeV^4 and the charge density $\rho_{Q\,\mathrm{SM}}\approx0.5539$ GeV^3 for which we obtained (for $Q=1$) a droplet of this phase with the mean electric charge radius $r_Q\approx0.149$ fm and the mass $M_{Q=1}=80.13$ GeV. The mass of a droplet of this phase

$$M_Q\to Qgv/2=Q\times80.13\ \mathrm{GeV}$$

as $\rho_{Q\,\mathrm{SM}}\to0$ for all values of $\lambda>0$. Configurations of this phase lie only on the M_Q-r_Q curve (see Fig. 10) or equivalently on the $E_{\min}$-$\rho_{Q\,\mathrm{SM}}$ curve (see Fig. 9).

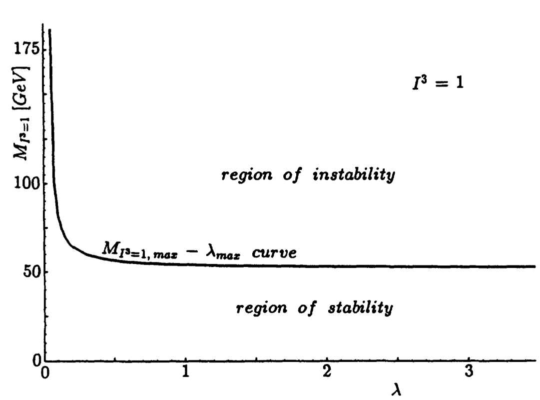

FIG. 17. The upper mass $M_{I^3=1,\max}$ of a BCMS configuration of the phase with $\vartheta=0$ and $\delta\neq0$ as the function of the $\lambda_{\max}$ nonlinear Higgs parameter where $\lambda_{\max}$ and $M_{I^3=1,\max}$ are the values of λ and $M_{I^3=1}$, respectively, for which $m^2_{\tilde{W}^\pm}=0$ (on the curve). The region of possible configurations of this phase is on and below the $\lambda_{\max}$-$M_{I^3,\max}$ curve.

The second phase ($\vartheta=0$ and $\zeta\neq0$) with the Z_μ gauge field condensation appears only when the weak charge density $\rho_{Z\,\mathrm{SM}}\neq0$. This phase is very similar (especially when $\rho_{Z\,\mathrm{SM}}\to0$) to the GSW low-temperature phase. The region of possible configurations of this phase is on and below the $\lambda_{\max}$-$\rho_{Z\,\max}$ curve (see Fig. 16). For

$$\lambda<g^2/(16\cos^4\Theta_W)\approx0.042\,2$$

the phase is stable for all values of weak charge density $\rho_{Z\,\mathrm{SM}}$. However, it is stable for an arbitrary λ only if

$$\rho_{Z\,\mathrm{SM}}\leq gv^3/(8\cos^2\Theta_W)\approx1.655\times10^6\ \mathrm{GeV}^3\ .$$

For $\rho_{Z\,\mathrm{SM}}>\rho_{Z\,\max}\approx1.655\times10^6$ GeV^3 this phase is unstable in the $W^\pm$ sector. When $\rho_{Z\,\mathrm{SM}}>\rho_{Z\,\max}$ the system may reach the charged stable phase ($\vartheta\neq0$).

The further evolution of this system seems to be very interesting. During its evolutions the system will emit charged particles lowering its energy and the charge. This emission may come from particle-antiparticle pair production which will neutralize the charged vacuum (the Greiner state [6]). Depending on the initial mass of the BCMS configuration ($M_{I^3}>M_Q$ or $M_{I^3}<M_Q$) the system in this process may reach the metastable point (with $\vartheta\neq0$) or the other ordinary standard model vacuum (with $\vartheta=0$). Such a process may happen in heavy ion collisions or in astrophysical objects.

ACKNOWLEDGMENTS

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APPENDIX A

Let us consider a system of quantum boson fields ϕ_A and a new system of quantum boson fields $\tilde{\phi}_A$ related to ϕ_A by

$$\phi_A=\tilde{\phi}_A+\xi_A \ , \tag{A1}$$

where the shifts ξ_A are the classical fields. These shift transformations can be expressed as

$$\tilde{\phi}_A=\mathcal{D}(\xi_A)\phi_A\mathcal{D}^{\dagger}(\xi_A) \ ,$$

where

$$\mathcal{D}(\xi_A)=\exp\sum_A\sum_\eta\int d^3k\,(\xi^{\eta}_{A\mathbf{k}}a^{\dagger}_{A\mathbf{k}\eta}-\xi^{\eta *}_{A\mathbf{k}}a_{A\mathbf{k}\eta}) \ .$$

Here $\sum_A$ is the sum over all shifted fields and $\sum_\eta$ means the sum over all degrees of freedom for these fields. The $a_{A\mathbf{k}\eta}$, and $a^{\dagger}_{A\mathbf{k}\eta}$ are the annihilation and creation operators for the ϕ_A field. The coefficients $\xi^{\eta}_{A\mathbf{k}}$ are the Fourier transformations of the ξ_A fields.

Now we assume that in the Hilbert space $\mathcal{H}$ there exists a normalized vacuum vector $|0\rangle$ which is annihilated by the operators $a_{A\mathbf{k}\eta}$

$$a_{A\mathbf{k}\eta}|0\rangle=0 \quad\text{and}\quad \langle 0|0\rangle=1 \ .$$

The shifts cause the changing of the ground state of a system according to the relation:

$$|0\rangle\rightarrow|\tilde{0}\rangle=\mathcal{D}(\xi_A)|0\rangle \ .$$

The new vacuum state $|\tilde{0}\rangle$ is simply the Glauber coherent state. This state includes the infinite number of excited states of ϕ_A fields. The state $|\tilde{0}\rangle$ is also normalized, i.e., $\langle\tilde{0}|\tilde{0}\rangle=1$.

As the state $|0\rangle$ is the vacuum state for the ϕ_A fields also the state $|\tilde{0}\rangle$ may be considered as the vacuum state for the $\tilde{\phi}_A$ fields. Hence, when we have $\langle 0|\phi_A|0\rangle=0$ we also have $\langle\tilde{0}|\tilde{\phi}_A|\tilde{0}\rangle=0$ and

$$\langle\tilde{0}|\phi_A|\tilde{0}\rangle=\langle\tilde{0}|\tilde{\phi}_A|\tilde{0}\rangle+\xi_A=\xi_A \ .$$

The point is that when the ground state $|\tilde{0}\rangle$ is attained as the result of the transformation as in Eq. (A1) which is not the gauge symmetry transformation or as the result of the appearance of some new external charges in the system it leads to the conclusion that the Fock spaces which are built on the ground states $|0\rangle$ or $|\tilde{0}\rangle$, respectively, are not unitary equivalent. This means that some classical boson fields ξ_A may attain physical interpretation.

We use the notation $\langle\mathcal{A}\rangle_{\tilde{0}}$ for the mean value $\langle\tilde{0}|\mathcal{A}|\tilde{0}\rangle$.

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