Bound States in the Hot Electroweak Phase

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Abstract

The high temperature phase of the electroweak standard theory is described by a strongly coupled $SU(2)$-Higgs-model in three dimensions. As in the Abbott-Farhi-model Higgs and $W$-boson are low lying bound states. Using a method by Simonov based on the Feynman-Schwinger representation of correlators we calculate the masses of these states. Our results are compared with lattice masses.

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The electroweak standard theory successfully predicted massive vector bosons making use of the Higgs mechanism. Much of its strength is based on the possibility to do rigorous perturbative calculations. Later on interest also turned towards the behavior of this theory at high temperature that has been realized in the early universe. The proposal that there exists a hot phase where the Higgs field obtains a positive thermal mass and the SU(2)_L \otimes U(1)-symmetry is restored \[1\] is commonly accepted. Investigations of the electroweak phase transition have to treat the hot and the cold phase at the transition temperature in an appropriate way.

Perturbation theory in the broken phase has been studied extensively \[2\]-\[9\]. For \( \lambda \varphi^4 \)-couplings corresponding to zero temperature Higgs masses below 70 GeV it works very well. A perturbative treatment of the hot phase would predict a vanishing Higgs vacuum expectation value and a massless gauge boson. However perturbation theory alone is not a good guide in this phase due to IR-problems.

The electroweak standard theory at the transition temperature allows the high temperature expansion. To good accuracy it can be described by an effective 3-dimensional SU(2)-Higgs-model whose parameters can be calculated from the fundamental 4-dimensional theory \[10\]. All IR-problems appearing in the unbroken phase are represented by this effective theory. It might be compared to 3-dimensional QCD but with a fundamental scalar field \( \Phi \) instead of quarks and a SU(2)-gauge group. \( \Phi \) corresponds to the perturbative Higgs particle. We avoid this name to preclude confusions. Its mass \( m_3 \) will be called Lagrangian mass.

As in QCD the most secure way to calculate the properties of such a system is a Monte-Carlo simulation on the lattice. This is done intensely and first results are appearing \[11\]-\[14\]. Of course being aware of the long way to reliable hadron lattice results one has to question if the present lattice sizes are appropriate for measuring certain quantities. An interesting alternative way to such 3-dimensional systems at high temperature is the exact renormalization group approach \[15\]. It requires, however, simplifying assumptions for practical use, and these are not easily controlled. Recently it was speculated from the investigation of 1-loop gap-equations in the two settings with small and large Higgs vacuum expectation value, that the hot phase is just another Higgs phase with different parameters \[16\]. But this does not agree well with the masses \( m_W \) and \( m_H \) observed on the 3-dimensional lattice \[11, 12, 14\]. The subject of this letter is to propose a model which explains these correlation masses; we compare with the lattice results at the end.

With rising 3-dimensional gauge coupling in the infrared \[17\] it is reasonable to propose confinement like in QCD. As in the Abbott-Farhi-model \[18\] physical states are bound states and singlets with respect to the gauge group. They may nevertheless be in a nontrivial isospin representation. Let us start with the SU(2)-Higgs-model Lagrangian

\[
\mathcal{L} = -\frac{1}{2} \text{Tr } F_{kl} F_{kl} + (D_k \phi)^\dagger (D^k \phi) - V(\phi^\dagger \phi) .
\]

\( D_k = \partial_k - ig A_k \) is the covariant derivative with the gauge field \( A_k = A_k^i \tau^i \) (\( \tau^i \) are the Pauli matrices), \( F_{kl} = [D_k, D_l] \) is the field strength tensor and \( \phi = (\varphi_1^i, \varphi_2^i) \) is the scalar isospin doublet. \( \mathcal{L} \) is invariant under the gauge transformation \( U(x) \in SU(2) \)

\[
D_k \rightarrow UD_k U^\dagger \quad \phi \rightarrow U \phi \, .
\]

The scalar doublet can unambiguously be written as \( \phi = \Phi (1) \), where the 2×2-matrix \( \Phi \) is a linear combination of \( 1_{2\times 2} \) and \( \tau^i \), and \( \frac{1}{2} \text{Tr } \Phi^\dagger \Phi = \phi^\dagger \phi \). Using this notation
uncovers the $SU(2)\otimes SU(2)_I$ invariance of the Lagrangian (eq. (1)), namely $\Phi \rightarrow U\Phi V$, of which the left $SU(2)$ is gauged while the isospin $SU(2)_I$ is global.

The confining theory has (among others) bound states corresponding to the local interpolating fields

$$\text{Tr} \Phi\Phi^\dagger$$

and

$$\text{Tr} \Phi\Phi^\dagger D_k \Phi \tau^i$$

The first one is a scalar isospin singlet and is identified with the “Higgs particle”. The second one is a vector isospin triplet, the “W-boson”. The scalar triplet operator $\text{Tr} \Phi\Phi^\dagger \tau^i$ vanishes. The operator $\text{Tr} \Phi\Phi^\dagger D_k \Phi$ does not correspond to a vector singlet bound state but is identical to $\partial_k \text{Tr} \Phi\Phi^\dagger$ and generates the scalar singlet at non-zero momentum. There are no bound states with these quantum numbers.

The nonlocal versions of the operators in eq. (3) are proportional to

$$\text{Tr} \Phi(x)\Phi(x)^\dagger T(x,\bar{x})$$

and

$$\text{Tr} \Phi(x)\Phi(x)^\dagger T(x,\bar{x}) \Phi(\bar{x}) \tau^i$$

The space index of the vector operator is fixed by the direction of the link operator $T$ (also known as gauge field transporter). The latter is defined as ($P$ denotes the path ordering)

$$T(x,\bar{x}) = P \exp \left( ig \int_x^{\bar{x}} A_k dz_k \right)$$

These are the direct continuum counterparts of the lattice observables interpreted as Higgs and $W$-boson.

It has been shown in the framework of sum rules, that this model reproduces the phenomenology of the standard theory in the Higgs phase without any use of naive spontaneous symmetry breaking but with an appropriate pattern of $\text{Tr} \Phi\Phi^\dagger$ vacuum structure [19]-[21]. (There is no fundamental difference between the 3 and 4-dimensional case.) Of course already at that time there was the other option (the genuine Abbott-Farhi-model) that gluon condensates like in QCD lead to a strongly bound phase [22]. Indeed this is the picture we propose for the hot phase of electroweak theory.

The emerging physical picture would thus be very similar to QCD in 3 space time dimensions, but with spinless constituents, the fundamental scalar fields. It is tempting to try to make connection to the confining quark model of QCD which in a very transparent way can explain many features of the bound state spectra. Since the fundamental scalar particles are light compared to the string tension $\sigma$ we cannot use the simple non-relativistic quark model.

In the case of QCD Simonov [23] has used the Feynman-Schwinger representation of hadronic correlators in order to treat light quarks kept together by a linear confining potential. The method is easily adopted to the case treated here, i.e. a theory in 3 space-time dimensions with scalar “quarks”. In this case we even avoid the problems of the large spin-interactions which lead presumably to zero modes and chiral symmetry breaking. Another difference is that the $W$-boson is an isospin-triplet, a state which has no analogy in QCD. Simonov’s method is nevertheless applicable since the isospin structure is not changed by the propagation. The propagator of the matrix field $\Phi(x)_{\alpha a}$ to $\Phi(y)_{\beta b}$ is given in the Feynman-Schwinger representation by ($\alpha, \beta$ are gauge and $a, b$ are isospin indices)

$$G(x,y)_{\alpha a \beta b} = \int_0^\infty ds \exp \left( -m_{3\delta}^2 s \right) \int Dz \exp \left( -\frac{1}{4} \int_0^s \dot{z}_k^2 d\tau \right) T(x,y)_{\alpha \beta} \delta_{ab}$$
The path integral $\int \mathcal{D}z$ runs over all curves $z(\tau)$ with tangent $\dot{z}(\tau)$ connecting $x$ and $y$.

Using a similar expression for the anti-scalar propagator from $\bar{x}$ to $\bar{y}$ one can put things together to get the Green function for a scalar–anti-scalar state at position $(x, \bar{x})$ propagating to $(y, \bar{y})$ in the quenched approximation

$$G(x, \bar{x}, y, \bar{y}) = \int_0^\infty ds \int_0^\infty d\bar{s} \int \mathcal{D}z \mathcal{D}\bar{z} \exp \left(- (m_3^2(s + \bar{s})) \right) \times \exp \left(-\frac{1}{4} \int_0^s \dot{z}^2_k(\tau) d\tau - \frac{1}{4} \int_0^{\bar{s}} \dot{\bar{z}}^2_k(\bar{\tau}) d\bar{\tau} \right) \mathcal{P} \exp \left(ig \oint_{x,\bar{x},y,\bar{y}} A_k dz_k \right).$$

The last factor of eq. (10) is just the Wegner Wilson integral over the closed loop formed by the gauge field transporters $T$ of the propagators (eq. (7)) and the bound state operators (eq. (5)). In a first approach it is approximated by the area law

$$\mathcal{P} \exp \left(ig \oint_{x,\bar{x},y,\bar{y}} A_k dz_k \right) \propto \exp (-\sigma F),$$

where $F$ is the area and $\sigma$ the string tension.

We are evaluating the Green function eq. (10) following Simonov [23]. Making some plausible assumptions on the (nearly) minimal surfaces the surface $F$ is parameterized by $z_k(\tau)$ and $\bar{z}_k(\bar{\tau})$. In a next step the “center of mass trajectory” $R_k$ and the relative coordinate $u_k$ are introduced

$$R_k = \frac{s\bar{s}}{s + \bar{s}} \left(\frac{1}{s} z_k + \frac{1}{\bar{s}} \bar{z}_k\right) \quad u_k = z_k - \bar{z}_k .$$

The path integrals over $z_k$ and $\bar{z}_k$ transform into path integrals over $R_k$ and $u_k$.

Making the assumption that the center of mass motion is dominated by the classical path the path integral over $R_k$ can be replaced by the parameterization $R_k(\gamma)$.

Without loss of generality we choose $R_1 = R_2 = 0$, $R_3 = \gamma \Theta = \vartheta$ with $0 \leq \gamma \leq 1$, where $\Theta$ is the distance between the centers of mass of the bounded scalar–anti-scalar system at position $(x, \bar{x})$ and at position $(y, \bar{y})$. This assumption is the most stringent one and should become better for heavier states (as compared to $\sqrt{\sigma}$).

One then can rewrite the Green function (eq. (10)) as (for detail see ref. [23])

$$G(x, \bar{x}, y, \bar{y}) \propto \int_0^\infty ds \int_0^\infty d\bar{s} \int \mathcal{D}u \exp(-B)$$

with the three dimensional action

$$B = \int_0^\Theta d\theta \left[ m_3^2 \left(\frac{s + \bar{s}}{\Theta} \right) + \frac{s + \bar{s}}{4s\bar{s}} \Theta + \frac{\Theta}{4(s + \bar{s})} \left(\frac{\partial u}{\partial \theta}\right)^2 + \sigma \sqrt{u_1^2 + u_2^2} \right].$$

It is convenient to introduce the new variables

$$\mu_1 = \frac{\Theta}{2s} \quad \mu_2 = \frac{\Theta}{2\bar{s}} \quad \bar{\mu} = \frac{\mu_1 \mu_2}{\mu_1 + \mu_2} .$$

The $u_1$ and $u_2$ path integral can be substituted by the solutions of the two dimensional Schrödinger equation

$$H \psi(u_1, u_2) = \epsilon \psi(u_1, u_2) \quad \text{with} \quad H = -\frac{1}{2\bar{\mu}} \left(\frac{\partial^2}{\partial u_1^2} + \frac{\partial^2}{\partial u_2^2}\right) + \sigma \sqrt{u_1^2 + u_2^2} .$$
This is most easily seen if one interprets $B$ (eq. (11)) as the Euclidean action of a point particle in two space and one time dimension. The last path integral over $u_3$ is now trivial and gives a negligible contribution. The remaining $s$ and $\bar{s}$ integrals are finally evaluated by the method of steepest descent. The correlation mass $M$ of the bound states, defined by $G \propto \exp(-M\Theta)$, is hence found by minimizing the integrand of $B$ evaluated at the solutions of the Schrödinger equation with respect to $\mu_1$ and $\mu_2$. Since the Lagrangian masses of the both fundamental scalars are the same ($m_3$) one gets only one condition

$$\frac{\partial M(\mu)}{\partial \mu} = 0 \quad M(\mu) = \frac{m_3^2}{\mu} + \mu + \epsilon(\mu)$$

(14)

with $\mu_1 = \mu_2 = \mu = 2\tilde{\mu}$. $\epsilon(\mu)$ is the eigenvalue introduced in eq. (13). One may easily express $\epsilon(\mu)$ by the dimensionless eigenvalues $a_{nl}$ of the equation

$$-\psi''(\rho) - \frac{1}{\rho}\psi'(\rho) + \left(\frac{l^2}{\rho^2} + \rho\right)\psi(\rho) = a_{nl}\psi(\rho)$$

(15)

through

$$\epsilon_{nl}(\mu) = \frac{\sigma^{2/3}}{\mu^{1/3}} a_{nl}.$$

(16)

Here $n$ is the principal quantum number and $l$ the orbital one. The numerical values of $a_{nl}$ for the lowest quantum numbers have been obtained numerically. They are listed in table 1. In order to compare with the potential and with the lattice size the argument of the wave functions has been scaled back in units $(g^2T)^{-1}$.

### Table 1: The dimensionless eigenvalues $a_{nl}$ of eq. (15)

| $n$  | $l=0$ | $l=1$ | $l=2$ |
|------|-------|-------|-------|
| 1    | 1.74  | 2.87  | 3.82  |
| 2    | 3.67  | 4.49  | 5.26  |

In order to compare with the potential and with the lattice size the argument of the wave functions has been scaled back in units $(g^2T)^{-1}$. $\psi_{n=1,l=0}$ and $\psi_{n=1,l=1}$, the (radial symmetric) solutions of the Schrödinger equation, are plotted together with the potential in Fig. 2 versus $r = \sqrt{u_1^2 + u_2^2}$.

From the steepest descend equation (14) we obtain

$$\mu = \sqrt{\sigma} \ z(a_{nl}, m_3^2/\sigma)^{3/2},$$

(17)

where $z(a, y)$ is a solution of the cubic equation

$$z^3 - \frac{1}{3} a \ z - y = 0$$

(18)

yielding

$$z(a, y) = \frac{2^{1/3}}{3} a \left(27y + \sqrt{729y^2 - 4a^3}\right)^{-1/3} + \frac{1}{3} \frac{21/3}{2^{1/3}} \left(27y + \sqrt{729y^2 - 4a^3}\right)^{1/3}.$$

(19)

The mass $M_{nl}$ of the bound state with the quantum numbers $n$ and $l$ is read off eq. (14) using eq. (17) and eq. (17)

$$\frac{M_{nl}(m_3/\sqrt{\sigma})}{\sqrt{\sigma}} = 4z^{3/2} - \frac{2m_3^2}{\sigma z^{3/2}}.$$

(20)
Figure 1: The masses of the bound states vs. the squared Lagrangian mass. The full lines are the masses of the s-states, the dashed lines are those of the p-states and the dot-dashed line corresponds to the 1d-state.

$M_{nl}$ of the lowest bound states is plotted in Figure 1 versus the mass of the fundamental scalar (which corresponds to the current quark mass in QCD).

Our results show the expected hierarchy of states, i.e. the lowest lying state is an isoscalar S-state (the “Higgs”), followed by an isovector P-state (the “W-boson”). They should now be compared with existing lattice data. These data are given for certain values of lattice parameters $\beta_G$ and $\beta_H$, where $\beta_H$ essentially fixes the temperature, and the lattice constant is $a = \frac{4}{\beta_G g^2 T}$. Masses in units of $g^2 T$ are obtained by multiplying the lattice ones with $\frac{4}{\beta_G g^2 T}$.

We consider lattice data of ref. [14] corresponding to a zero temperature Higgs mass of 35 GeV, and concentrate on the data point $\beta_G = 12$, $\beta_H = 0.3412$ where the static force has been measured as well. The mass values are $m_H = (0.65 \pm 0.1) g^2 T$ and $m_W = (1.75 \pm 0.08) g^2 T$, measured on a $30^3$ lattice. The static force is consistent with a confining potential with string tension $\sigma = (0.13 \pm 0.02) (g^2 T)^2$ up to a distance of $4(g^2 T)^{-1}$ where it has been measured. At some (large) distances screening should set in, which, however, is presumably not relevant for the lowest bound states.

For a comparison, we have to estimate the corresponding Lagrangian mass $m_3$ of the fundamental scalar. From $\beta_H$, it is obtained as

$$\left( \frac{m_3}{g^2 T} \right)^2 = \left( \frac{\beta_G}{4} \right)^2 \left( \frac{2(1 - 2\beta_R - 3\beta_H)}{\beta_H} \right)_{\text{subtr}}$$

with $\beta_R = \frac{\beta_H}{g^2 \beta_G}$, and the subscript indicates that the expression has to be subtracted at the critical $\beta_H$, which is equal to 0.34138 in the present case. This leads to $m_3 = 0.17 g^2 T$ at $\beta_H = 0.3412$. In eq. (21) we assume for simplicity that $m_3$ vanishes at the critical $\beta_H$. Strictly speaking, $m_3$ depends on an arbitrary normalization scale $\mu_3$; in deriving eq. (21) we adopt the natural choice $\mu_3 \propto g^2 T$. 
A comparison with lattice data can now be performed based on equations (20) (see also Fig. 1). Using the string tension \( \sigma = (0.13 \pm 0.02) (g^2 T)^2 \) and the Lagrangian mass \( m_3 = 0.17 g^2 T \) from the lattice data one sees that our model yields qualitatively the right results of a substantial splitting between these states. The predicted mass of the W-boson \( (m_W = (1.47 \pm 0.11) g^2 T) \) agrees quite well with the lattice mass, in view of the approximations made. Taking into account finite-\( a \) effects the lattice value may come down improving the agreement. The mass of the composite Higgs \( (m_H = (1.06 \pm 0.07) g^2 T) \) calculated by us is a factor 1.6 larger than the lattice value.

One may also take into account the attractive Coulomb force. The force measured on the lattice is excellently fitted by

\[
F = 0.123 (g^2 T)^2 K_1(1.01 g^2 T r) + 0.13 (g^2 T)^2
\]

(cf. Fig. 5 of ref. 14). The potential is obtained from \( F \) by integration fixing it to \( \sigma r \) at large distances. The first term corresponds to the exchange force of a particle with mass \( 1.01 g^2 T \), the constant part is the string tension. The influence of the Coulomb force is small: it lowers the mass of the composite Higgs by only 2%, that of the W by 0.2%. This is due to the fact, that the bound state solutions are much larger than the range of the Coulomb force (cf. Fig. 3). The \( r \)-dependent part of the force can therefore be neglected in good accuracy justifying the use of the area law in eq. (8).

As far as our model calculations are concerned various approximations are involved. Perhaps the most important one is neglecting the quantum fluctuations of the center-of-mass coordinates \( R_k \). This approximation should work better for heavier (recurrent) bound states. The deviation of the Higgs correlation mass may be explained by this.

The lattice results are influenced by finite volume as well as finite lattice spacing effects. Although finite size investigations did not show significant effects on the masses one should
keep in mind that the predicted size of the bound states is comparable to the lattice size, which in this case is $10(g^2T)^{-1}$ with periodic boundary conditions (see Fig. 2).

Recurrent states in the isoscalar S-channel would be difficult to identify because they are predicted in the mass range of the $W$-balls of pure $SU(2)$ gauge theory $(1.5g^2T)^{-1}$ with periodic boundary conditions (see Fig. 2). More interesting are the d-states ($l=2$). The corresponding $W$-ball states have masses of the order of $2.5g^2T$ which is significantly heavier than the predicted $l=2$ scalar–anti-scalar bound state mass $m_{1d}$ of about $1.8g^2T$. Note, nevertheless, that this is above the threshold of the two Higgs channel corresponding to the measured masses. Lattice measurements of this channel are under way.

The predicted masses are based on the evaluation of the Green function eq. (6), which could only be performed with some approximations. Even without these approximation one finds that the masses of the bound state model depend only on the string tension $\sigma$ and the Lagrangian mass $m_3$. Considering lattice data in the whole $\beta_H$ (i.e. $m_3$) range one finds a stronger $m_3$-dependence of $m_W/(g^2T)$ and $m_H/(g^2T)$ than we would predict if $\sqrt{\sigma}/(g^2T)$ was constant. This might indicate that $\sqrt{\sigma}/(g^2T)$ itself depends significantly on $m_3$, dropping as $m_3$ approaches zero. It would be interesting to verify or disprove this by lattice studies. $\sqrt{\sigma}/(g^2T)$ could also show a sizable $\lambda$-dependence (decreasing with increasing $\lambda$).

One should also keep in mind that the static potential will not continue to rise linearly at larger distances but charge screening will set in. The screening length has not been measured so far. As a consequence of charge screening, heavier bound states will disappear. The mass of a bound state close to the threshold would be affected as well.

In conclusion the hot electroweak phase appears as an intriguing case of a confining 3-dimensional gauge system. It is similar to 3-dimensional QCD but with a fundamental scalar $\Phi$ instead of quarks and hence without spin-interactions and without the problems of chiral symmetry breaking. The masses of the low lying bound states have been calculated from the Green function within some approximations. The comparison with lattice results is satisfactory. Our discussion is open for refinement both on the lattice side because of the limited lattice size and on the side of calculating the masses of relativistic bound states, a notorious problem also in QCD. For the latter question also other methods can be envisaged, including sum rules for the 3-dimensional Abbott-Farhi-model [13]–[21]. An enlarged version of this letter is in progress [25].

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