WEAK GAUGE COUPLING EXPANSION IN THE LATTICE-REGULARIZED STANDARD SU(2) HIGGS MODEL

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The weak gauge coupling expansion is derived in the lattice-regularized SU(2) Higgs model with a scalar doublet field at an arbitrary point of the parameter space boundary with vanishing gauge coupling. Consequences of the triviality of the ϕ^4 component on gauge invariant Green's functions are formulated.

Quantum field theory is dealing with an infinite number of degrees of freedom. This implies that a mathematically meaningful formulation can only be given by first introducing some regularization. There are, in principle, many different regularizations possible, but the most popular ones are, no doubt, dimensional regularization (DR) [1] and lattice regularization (LR) [2]. It is generally expected that the physical content of a theory is independent of the choice of the regularization procedure. In fact, for instance in quantum chromodynamics some part of the calculations (like jet-calculus etc.) are usually done by DR, some other part (like hadron mass calculations etc.) by LR. In the case of the standard SU(2) ⊗ U(1) electro-weak theory the situation is somewhat different, because almost exclusively DR is considered. There are, at least, two facts which seem to speak against LR in this case: the apparent difficulty to put chiral fermions on the lattice [3] and the almost rigorously proven triviality of the ϕ^4 model on the lattice [4] [‡1]. It is not known whether these difficulties are just consequences of the use of LR or they signal profound features independent of regularization. In view of this, the study of the lattice-regularized standard SU(2) ⊗ U(1) model (and/or some other models with elementary scalar fields) is very important and interesting.

The lattice regularization allows for a variety of different approaches: besides the powerful exact theorems it is also possible to perform approximate numerical calculations or different sorts of analytic expansions. The general strategy of the analytic expansions is to reduce the number of coupling parameters by sending some of them to the boundary of the coupling parameter space. The small parameter in the expansion is the distance from the boundary in some appropriately chosen metric. In the standard Higgs model (i.e. SU(2) gauge field coupled to a complex scalar doublet), which will be considered throughout this paper, there are three couplings: the scalar self-coupling λ, the gauge coupling g (or $\beta \equiv 4g^{-2}$) and the hopping parameter κ representing in LR the mass parameter for the scalar field. A possible expansion in this model is the strong self-coupling expansion (SSCE) investigated in refs. [6,7]. In SSCE expectation values at some point of the parameter space are expressed by a series containing expectation values at infinite self-coupling ($\lambda = \infty$). In the present paper a similar weak gauge coupling expansion (WGCE) will be derived at an arbitrary point of the ($\beta = \infty$) plane. Since in this case the terms of the series depend on the expectation values in the four-component ϕ^4 model, the consequences of the expected triviality of ϕ^4 can be easily imposed.

The notations and conventions for the lattice description of the standard Higgs model will be in general the same here as in refs. [6–8]. Therefore, the SU(2) gauge link-variable will be denoted by $U(x, \mu) \in$ SU(2), the

‡1 For a review see ref. [5].

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length of the Higgs field will be $\rho_x \geqslant 0$ and the angular Higgs variable $\alpha_x \in SU(2)$. x denotes lattice points, $\mu = \pm 1, \pm 2, \pm 3, \pm 4$ are link directions and (x, μ) is the link from the point x to the neighbouring point $(x + \hat{\mu})$ in direction μ. For the Higgs field we shall also use

$$\varphi_x \equiv \rho_x \alpha_x . \tag{1}$$

The lattice action in these variables can be written like

$$S_{\lambda,\beta,\kappa} = \beta \sum_{P} (1 - \tfrac{1}{2}\mathrm{Tr}\, U_P) + \sum_x [\rho_x^2 - 3 \log \rho_z + \lambda(\rho_x^2 - 1)^2] - \kappa \sum_{(x\mu)} \mathrm{Tr}(\varphi^+_{x+\hat{\mu}} U(x,\mu) \varphi_x). \tag{2}$$

Here Σ_P stands for a summation over positively oriented plaquettes, and $\Sigma_{(x\mu)} \equiv \Sigma_{x,\mu>0}$ is a sum over positive links. The integration measure in the path integral is $d\rho_x d^3\alpha_x d^3 U(x,\mu)$ (where $d^3 g$ denotes the Haar-measure in SU(2)). In the limit $\beta \to \infty$ the gauge part of the action vanishes (the link-variables become gauge equivalent to unity), therefore the $\beta \to \infty$ action is

$$S_{\lambda,\beta=\infty,\kappa} = \sum_x [\rho_x^2 - 3 \log \rho_x + \lambda(\rho_x^2 - 1)^2] - \kappa \sum_{(x\mu)} \mathrm{Tr}(\varphi^+_{x+\hat{\mu}} \varphi_x). \tag{3}$$

This defines a four-component ϕ^4-model with global SU(2) ⊗ SU(2)- (or O(4)-) symmetry.

The derivation of the WGCE at the point $(\lambda, \beta = \infty, \kappa = \kappa_0)$ starts from the relation

$$S_{\lambda,\beta,\kappa} = S_{\lambda,\beta=\infty,\kappa_0} + S_g - \sum_{(x\mu)} (\kappa\, \mathrm{Tr}\{\varphi^+_{x+\hat{\mu}} [U(x,\mu) - 1] \varphi_x\} + (\kappa - \kappa_0) \mathrm{Tr}(\varphi^+_{x+\hat{\mu}} \varphi_x)). \tag{4}$$

Here S_g denotes the Wilson action for the SU(2) gauge field:

$$S_g \equiv \beta \sum_P (1 - \tfrac{1}{2}\mathrm{Tr}\, U_P). \tag{5}$$

The relation in eq. (4) corresponds to eq. (7) of ref. [7] for the SSCE. We are interested in the generating function Z of the gauge-invariant connected correlation functions:

$$Z[r,k]_{\lambda,\beta,\kappa} \equiv \log \left\langle \exp\left(\sum_x r_x \rho_x + \sum_{r(x\mu)} k_{rx\mu} \mathrm{Tr}[\tau_r \varphi^+_{x+\hat{\mu}} U(x,\mu) \varphi_x] \right) \right\rangle_{\lambda\beta\kappa} . \tag{6}$$

Here τ_r, $(r = 1, 2, 3)$ is a weak-isospin Pauli matrix. The derivatives of Z with respect to r give the connected correlation functions of the gauge-invariant Higgs-boson variable ρ (weak isospin zero), whereas derivatives with respect to k produce the connected correlation functions of the isospin 1 gauge-invariant W-boson variable $\mathrm{Tr}(\tau \varphi^+ U \varphi)$. (Note that one could use, in principle, also other interpolating fields, for instance, in the Higgs-boson channel $\mathrm{Tr}(\varphi^+ U \varphi)$ or in the W-boson channel $\mathrm{Tr}(\tau \alpha^+ U \alpha)$, but this would not change anything essential.)

The path integral needed in eq. (6) can be written as

$$\int [d\rho\, d^3\alpha\, d^3 U] \exp(-S_{\lambda,\beta=\infty,\kappa_0} - S_g) \exp\left(\sum_x r_x \rho_x + \sum_{(x\mu)} [k_{rx\mu} v_{rx\mu} + (\kappa - \kappa_0) s_{x\mu} + a_{x\mu} j_{x\mu} + a_{rx\mu} j_{rx\mu}] \right), \tag{7}$$

where an automatic summation over repeated isospin indices $(r, s, t, ... = 1, 2, 3)$ is understood. The definitions of the quantities appearing in the exponent, and some similar ones needed later, are

$$s_{x\mu} \equiv \mathrm{Tr}(\varphi^+_{x+\hat{\mu}}\varphi_x), \quad v_{rx\mu} \equiv \mathrm{Tr}(\varphi^+_{x+\hat{\mu}}\varphi_x\tau_r), \quad u_{rx\mu} \equiv \mathrm{Tr}(\varphi^+_{x+\hat{\mu}}\tau_r\varphi_x), \quad w_{rsx\mu} \equiv \mathrm{Tr}(\varphi^+_{x+\hat{\mu}}\tau_r\varphi_x\tau_s),$$

$$U(x,\mu) - 1 \equiv -a_{x\mu} + \mathrm{i}\tau_r a_{rx\mu}, \quad a_{rx\mu} \equiv \tfrac{1}{2}agA_{rx\mu}, \quad j_{x\mu} \equiv -\kappa s_{x\mu} - k_{rx\mu}v_{rx\mu}, \quad j_{rx\mu} \equiv \mathrm{i}\kappa u_{rx\mu} + \mathrm{i}k_{sx\mu}w_{rsx\mu}. \tag{8}$$

The integration over the gauge variables cannot be performed explicitly (unlike the integration over the Higgs field length in the case of SSCE), but the integral can be expanded in powers of g^2 in the same way as in ordinary lattice gauge perturbation theory [9,10]. Instead of the usual SU(N) gauge variables [9,10] we shall use here $a_{rx\mu}$ (or $A_{rx\mu}$) as defined in eq. (8). This gives simpler four-point (and more-point) vertices. Writing the SU(2) Haar-measure in terms of $a_{rx\mu}$, the necessary gauge integral is

$$\int [\mathrm{d}a_{rx\mu}] \exp\left(-S_\mathrm{g} + \frac{1}{2}\sum_{(x\mu)}\sum_{n=1}^{\infty}\frac{(a_{rx\mu}a_{rx\mu})^n}{n} + \sum_{(x\mu)}(a_{x\mu}j_{x\mu} + a_{rx\mu}j_{rx\mu})\right). \tag{9}$$

The gaussian part of the integration can be performed, for instance, by imposing the lattice version of the covariant gauge condition [9,10], but we shall pursue here, for simplicity, a short-cut by splitting-up a piece $-\frac{1}{2}b\, a_{rx\mu}a_{rx\mu}$ from $a_{x\mu}j_{x\mu}$ and adding it to the gauge action S_g. The parameter b is, for the moment, arbitrary (it will be fixed later by convenience). This makes the quadratic part of the action non-degenerate and the gauge fixing unnecessary. The lattice gauge propagator has in this case a "unitary gauge" form:

$$\Delta_{r_1x_1\mu_1,r_2x_2\mu_2} = \frac{\delta_{r_1r_2}}{Na^2}\sum_k \frac{\exp[-\mathrm{i}(k,x_1-x_2)]}{(aM_\mathrm{W})^2 + (\hat{k}^*,\hat{k})}\,[\delta_{\mu_1\mu_2} + \hat{k}^*_{\mu_1}\hat{k}_{\mu_2}/(aM_\mathrm{W})^2], \tag{10}$$

where N is the number of lattice points, a the lattice spacing, k denotes discrete lattice momenta and

$$\hat{k}_\mu \equiv 1 - \exp(-\mathrm{i}k_\mu), \quad (aM_\mathrm{W})^2 \equiv \tfrac{1}{4}g^2(b-1). \tag{11}$$

Besides the usual three-, four- and more-point vertices there are now also two-point insertions (and also additional four-, six- and more-point vertices) coming from the rest of $a_{x\mu}j_{x\mu}$. Otherwise the perturbative expansion is straightforward. There are, of course, no Faddeev–Popov ghosts, since the gauge is not fixed. At the end we shall briefly comment on the differences which occur if a latticized covariant gauge condition is imposed.

By using a relation like eq. (25) or eq. (33) in ref. [7], a formula for the gauge-invariant generating function Z can be derived. Before writing it down, let us introduce a shorthand notation reducing excessive repetitions:

$$\{f.\}^n_\nu \equiv f_{\nu_1}f_{\nu_2}\ldots f_{\nu_n}, \quad \sum_{n(\nu)} \equiv \sum_{\nu_1\ldots\nu_2}. \tag{12}$$

In terms of this let us define

$$\{A\}^{lmn}_{l(rx\mu)m(y\nu)n(z\lambda)} \equiv \{A.\}^l_{rx\mu}\{A_s.A_s.\}^m_{y\nu}\{A_{s_1}.A_{s_1}.A_{s_2}.A_{s_2}.\}^n_{z\lambda}, \tag{13}$$

and

$$\{C\}^{lmn}_{l(rx\mu)m(y\nu)n(z\lambda)} \equiv [\mathrm{i}^l(-1)^{m+n}/2^{m+3n}l!m!n!]\,\{\kappa u. + k_s.w_s.\}^l_{rx\mu}\{\kappa s. + k_s.v_s. - b\}^m_{y\nu}\{\kappa s. + k_s.v_s.\}^n_{z\lambda}. \tag{14}$$

Note that, for simplicity, unimportant higher order terms in the last factor of $\{A\}$ were neglected. Using the trick in eq. (12) twice, the master formula for the generating function Z can be written as

$$Z[r,k]_{\lambda,\beta,\kappa} = \sum_{L(Z)M(RX\xi)N(Y\eta)} \sum_{K(l(rx\mu)m(yv)n(z\lambda))} \frac{\{r.\}_Z^L \{k.\}_{RX\xi}^M (\kappa-\kappa_0)^N}{L!M!N!K!} \{(\tfrac{1}{2}ag)^{l+2m+4n} \langle\{A\}_{\dots}^{lmn}\rangle_g^{\mathrm{c}}\}_{l(rx\mu)m(yv)n(z\lambda)}^K$$

$$\times \langle \{\rho.\}_Z^L \{v.\}_{RX\xi}^M \{s.\}_{Y\eta}^N \{\{C\}_{\dots}^{lmn}\}_{l(rx\mu)m(yv)n(z\lambda)}^K \rangle_{\lambda\kappa_0}^{\mathrm{c}} . \tag{15}$$

In the g-dependent connection expectation value of gauge variables $\langle\dots\rangle_l^{\mathrm{c}}$ the quadratic (respectively quartic) terms in the second (respectively third) factor of $\{A\}^{lmn}$ have to be considered as single entities for connectedness, whereas in the connected expectation value $\langle\dots\rangle_{\lambda\kappa_0}^{\mathrm{c}}$ calculated at $(\lambda, \beta=\infty, \kappa=\kappa_0)$ every term $\{C\}^{lmn}$ counts as a single entity. Taking derivatives of eq. (15) it is also possible to obtain expressions for connected gauge invariant expectation values. The g-dependent connected expectation value $\langle\dots\rangle_g^{\mathrm{c}}$ is, of course, given by a g^2 perturbation series in terms of the known gauge propagators and vertices. The connected expectation value $\langle\dots\rangle_{\lambda\kappa_0}^{\mathrm{c}}$ is defined in the ϕ^4 model at $\beta=\infty$ by the action in eq. (3).

Up to now the expansion point $(\lambda, \beta=\infty, \kappa=\kappa_0)$ was kept general. Forgetting about questions of convergence, the master formula eq. (15) gives, in principle, the expectation values at any point (λ, β, κ) in terms of the expectation values at a $\beta=\infty$ point (λ, κ_0) with arbitrary κ_0. Since in the $\beta=\infty$ plane there is a critical line $(\lambda, \kappa_{\mathrm{cr}}(\lambda))$ with diverging correlation lengths, the limit $\kappa_0 \to \kappa_{\mathrm{cr}}(\lambda)$ defines a continuum limit of the perturbation series for any $0 \leqslant \lambda \leqslant \infty$. Moreover, the limit of the correlation functions at $\kappa_{\mathrm{cr}}(\lambda)$ is expected to coincide with the correlation functions of a free theory. Therefore, in the case of $\kappa_0 \to \kappa_{\mathrm{cr}}(\lambda)$ only the two-point correlation functions survive. In this way it is possible to impose the consequences of the triviality of ϕ^4 on the WGCE in eq. (15). The limits of the non-vanishing scalar propagators appearing in eq. (15) are:

$$\langle \rho_{x_1}\rho_{x_2}\rangle_{\lambda\kappa_0}^{\mathrm{c}} \to a^2 Z_{\rho\rho}(\lambda)\Delta_{x_1x_2}^{(M_{\mathrm{H}})}, \quad \langle \rho_{x_1} s_{x_2\mu_2}\rangle_{\lambda\kappa_0}^{\mathrm{c}} \to a^3 M_{\mathrm{H}} Z_{\rho s}(\lambda)\Delta_{x_1x_2}^{(M_{\mathrm{H}})},$$

$$\langle s_{x_1\mu_1} s_{x_2\mu_2}\rangle_{\lambda\kappa_0}^{\mathrm{c}} \to a^4 M_{\mathrm{H}}^2 Z_{ss}(\lambda)\Delta_{x_1x_2}^{(M_{\mathrm{H}})} . \tag{16}$$

Here $\Delta_{x_1x_2}^{(M)}$ denotes the scalar propagator with mass M, which we define as

$$\Delta_{x_1x_2}^{(M)} = \frac{1}{Na^2} \sum_k \frac{\exp[-\mathrm{i}(k, x_1-x_2)]}{(aM)^2 + (\hat{k}^*, \hat{k})} . \tag{17}$$

This is the usual lattice form, which in the continuum limit is equal to the continuum propagator, apart from negligible a^2 corrections. Note that in eq. (16) the λ-dependent normalization factors $Z \dots (\lambda)$ are dimensionless and finite, corresponding to the assumed naive (canonical) dimensions of the fields ρ and s.

In the following only the spontaneously broken symmetry phase $\kappa_0 \geqslant \kappa_{\mathrm{cr}}(\lambda)$ will be considered. In the two-point functions of the currents u, v, w the contributions of the zero-mass Goldstone-particles are:

$$\langle u_{r_1x_1\mu_1} u_{r_2x_2\mu_2}\rangle_{\lambda\kappa_0}^{\mathrm{c}} \to \langle v_{r_1x_1\mu_1} v_{r_2x_2\mu_2}\rangle_{\lambda\kappa_0}^{\mathrm{c}}$$

$$\to \langle u_{r_1x_1\mu_1} v_{r_2x_2\mu_2}\rangle_{\lambda\kappa_0}^{\mathrm{c}} \to a^4 (af_\lambda/\kappa_0)^2\, \delta_{r_1r_2} (\partial/\partial(ax_1)_{\mu_1})(\partial/\partial(ax_2)_{\mu_2})\Delta_{x_1x_2}^{(0)},$$

$$\langle w_{r_1s_1x_1\mu_1} w_{r_2s_2x_2\mu_2}\rangle_{\lambda\kappa_0}^{\mathrm{c}} \to a^2 Z_{ww}(\lambda)(af_\lambda)^2 \epsilon_{r_1s_1t}\epsilon_{r_2s_2t}\Delta_{x_1x_2}^{(0)},$$

$$\langle w_{r_1s_1x_1\mu_1} u_{r_2x_2\mu_2}\rangle_{\lambda\kappa_0}^{\mathrm{c}} \to \langle w_{r_1s_1x_1\mu_1} v_{r_2x_2\mu_2}\rangle_{\lambda\kappa_0}^{\mathrm{c}} \to a^3 Z_w(\lambda)(af_\lambda)^2 \epsilon_{r_1s_1r_2} (\mathrm{i}\partial/\partial(ax_2)_{\mu_2})\Delta_{x_1x_2}^{(0)} . \tag{18}$$

Here again canonical dimensions were assumed for $\kappa_0 \to \kappa_{\mathrm{cr}}(\lambda)$. The quantity f_λ is the mass parameter characterizing spontaneous symmetry breaking, which is usually denoted in QCD by f_π. Note that, apart from a normalization factor, the vector current is $(u-v)$ and the axial-vector current is $(u+v)$.

In the continuum limit $\kappa_0 \to \kappa_{\mathrm{cr}}(\lambda)$ the usual divergencies appear in the individual terms of WGCE. It is important first to localize the most dangerous quadratic divergences and to try to define a resummation of the series in

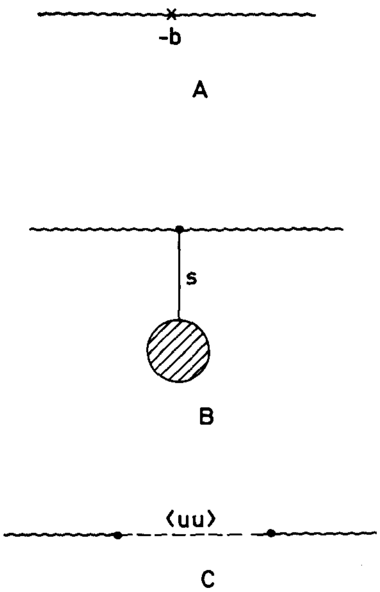

Fig. 1. Insertions in the gauge boson propagator. The wavy lines in the graphs represent the gauge propagator in eq. (10), solid lines are scalars (like ρ and s) and dashed lines are the currents (u, v and w). The blob in 1B stands for the sum of all graphs beginning with a scalar s-line.

such a way that only logarithmic divergences are left over. In the gauge propagator quadratically divergent insertions come from (κs.) and ($-b$) in the second factor of $\{C\}^{lmn}$ (see figs. 1a, 1b). Summing up all multiple insertions of the type in figs. 1a, 1b gives a shift of $(aM_{\mathrm{W}})^2$ in the gauge propagator. Taking now $b = 1$, which is the usual choice [9], and omitting $o(a^2g^2)$ pieces, instead of eq. (11) we obtain

$$(aM_{\mathrm{W}})^2 = \tfrac{1}{4}g^2\kappa\langle\mathrm{Tr}[\varphi^+_{x+\hat{\mu}}U(x,\mu)\varphi_x]\rangle_{\lambda\beta\kappa} \equiv \tfrac{1}{4}g^2\kappa\langle\mathcal{J}_{x\mu}\rangle_{\lambda\beta\kappa}. \tag{19}$$

This is a tree-level relation which corresponds in this gauge invariant formalism to the usual tree-level relation between vacuum expectation value of the scalar field and the W-mass. According to recent numerical Monte Carlo data [11] eq. (19) is well satisfied. For instance, the measured value $\langle\mathcal{J}_{x\mu}\rangle = 1.0589(3)$ at ($\lambda = 1.0, \beta = 8.0, \kappa = 0.28$) gives $aM_{\mathrm{W}} = 0.19252(3)$, to be compared to the Monte Carlo result $am_{\mathrm{W}} = 0.19(1)$. This shows that the loop corrections are in this point at most about 5%, in the same way as in conventional perturbation theory.

The multiple insertions of the current–current correlation function $\langle uu\rangle$ in fig. 1c do not shift the mass, but alter the spin structure of the gauge propagator. Summing up multiple insertions with the limiting form of $\langle uu\rangle$ in eq. (18) gives the gauge propagator

$$\overline{\Delta}_{r_1x_1\mu_1,r_2x_2\mu_2} = \frac{\delta_{r_1r_2}}{Na^2}\sum_k \frac{\exp[-\mathrm{i}(k, x_1 - x_2)]}{(aM_{\mathrm{W}})^2 + (\hat{k}^*, \hat{k})}[\delta_{\mu_1\mu_2} - (1-\xi_\lambda)^{-1}\hat{k}^*_{\mu_1}\hat{k}_{\mu_2}/(\hat{k}^*,\hat{k}) - \xi_\lambda(1-\xi_\lambda)^{-1}\hat{k}^*_{\mu_1}\hat{k}_{\mu_2}/(aM_{\mathrm{W}})^2], \tag{20}$$

The parameter ξ_λ is defined by

$$\xi_\lambda \equiv (2M_{\mathrm{W}}\kappa/gf_\lambda\kappa_0)^2. \tag{21}$$

It seems tempting to assume that in the continuum limit $\xi_\lambda = 0$, leaving us with a nice Landau-like gauge propagator, but the limiting value of f_λ (or f_λ/M_{H}) in the $\beta = \infty$ ϕ^4 model is unknown at present. This point clearly deserves further investigation.

Let us now briefly comment on the differences occurring if in the gauge integral the lattice version of the covariant gauge condition is imposed. In this case there are additional graphs containing closed loops of the

Faddeev–Popov ghost. The gauge propagator is first the same as in pure gauge theory [9,10], but summing up the insertions in fig. 1b one obtains the usual massive 't Hooft gauge propagator [12], which depends on a gauge parameter α. The simplest choice is $\alpha = 0$ ('t Hooft–Landau gauge), because then the insertion in fig. 1c vanishes. The good high-momentum behaviour of these gauge propagators is certainly advantageous, nevertheless the problem of quadratically divergent multi-loop contributions to the gauge propagator still remains.

A remarkable feature of the tree-level W-mass relation in eq. (19) is that, obviously, the right-hand side does not decrease fast enough for $g^2 \to 0$, $\kappa \to \kappa_{cr}(\lambda)$. This speaks against the existence of a non-trivial continuum limit of the standard Higgs model at $(\lambda, \beta = \infty, \kappa_{cr}(\lambda))$, in accordance with recent numerical investigations suggesting a first-order confinement–Higgs phase transition for finite β. In this case the lines of constant physics in the λ = const. planes would look like fig. 1b of ref. [13]. Another possibility is that due to remaining (multi-loop) quadratic divergences in the perturbation series the relation (19) is no longer valid in the vicinity of the phase transition, where the scale is much different from M_W. This could perhaps produce an exponential decrease of (aM_W) allowing a large cut-off compared to the W-mass.

In any case, the WGCE provides an analytic handle which can be used for a better understanding of the properties of the $\beta = \infty$ critical line in the standard Higgs model (and probably also in other models with elementary scalar fields). In a combination with SSCE it could also be useful for the study of λ-dependence.

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