Properties of the vacuum expectation values in $R_\xi$ and general gauge

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We have investigated the renormalization group evolution and the perturbative expansion of the vacuum expectation value(VEV) of the Abelian Higgs model both in the $R_\xi$ and the general gauge which extrapolates between the Fermi and the $R_\xi$ gauge. In $R_\xi$ gauge, the gamma function of the VEV was different from that of the scalar field. On the contrary, the gamma function of the VEV was the same as that of the scalar field in general gauge. Both in $R_\xi$ and $R_\xi$ gauge, the two-loop VEVs have an IR divergence in Landau gauge ($\xi = 0$) and these IR divergences do not occur when $\xi \neq 0$.

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Recently, the renormalization properties of the VEV in spontaneous broken gauge symmetry with $R_\xi$ gauge[1] were investigated[2][3]. It was shown that the gamma function of the VEV ($\gamma_\nu$) have additional contributions originating from the scalar background field and as a result, $\gamma_\nu$ was found to be different from that of the scalar field ($\gamma_\phi$). Moreover, it was found that perturbative expansion of the effective action has an IR divergence in the Landau($\xi = 0$) gauge[4][5]. The advantage of the $R_\xi$ gauge is that the mixing term between the gauge and the Goldstone fields in the quadratic part of the Lagrangian of the spontaneously broken symmetry phase cancels out with the gauge fixing term and hence the propagators between different fields do not exist. The $R_\xi$ gauge introduced by Kastening for some other reason[6] also has this property. In this paper, by using the Abelian-Higgs model, we will investigate the renormalization group(RG) function of the VEV from the RG invariance of the effective potential both in $R_\xi$ and the general gauge which extrapolates between the Fermi and the $R_\xi$ gauge and then we will calculate the perturbative expansion of the VEV from the no-tadpole condition in order to test the IR divergence in this gauge.

We start with the Euclidean Lagrangian density of the Abelian Higgs model given by

$$L_{AH}(\Phi_1, \Phi_2, A_\mu) = \frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{1}{2}(\partial_\mu \Phi_1 + gA_\mu \Phi_2)^2 + \frac{1}{2}(\partial_\mu \Phi_2 - gA_\mu \Phi_1)^2 + \frac{1}{2}m^2(\Phi_1^2 + \Phi_2^2) + \frac{\lambda}{24}(\Phi_1^2 + \Phi_2^2)^2 + \text{counter terms},$$

with

$$F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu. \quad (2)$$

where $f(A_\mu, \Phi)$ is the gauge-fixing function. When $m^2 > 0$, $L_{AH}(\Phi_1, \Phi_2, A_\mu)$ have a local $O(2)$ symmetry given by

$$\delta_\theta A_\mu = \partial_\mu \theta, \quad \delta_\theta \Phi_1 = \theta \Phi_2, \quad \delta_\theta \Phi_2 = -\theta \Phi_1,$$

and when $m^2 < 0$, $\Phi_1$ develops a VEV and the $O(2)$ symmetry is spontaneously broken. The Lagrangian density in the broken phase can be obtained by substituting $\Phi_1$ with $H + v$ where the Higgs field $H$ have a vanishing VEV.

A. $R_\xi$ gauge

The gauge-fixing function of the $R_\xi$ gauge is given in the broken symmetry phase of the theory as

$$f(A_\mu, \Phi_2) = \partial_\mu A_\mu - \xi g_0 \Phi_2, \quad (4)$$

and the total Lagrangian in the broken symmetry phase is given by

$$L_{BS}(H, \Phi_2, A_\mu) = L_{AH}(H + v, \Phi_2, A_\mu) + \frac{1}{2\xi}f(A_\mu, \Phi_2)^2 + \frac{\tau f(A_\mu, \Phi_2)^2}{\delta \theta}c + \text{counter - terms}, \quad (5)$$

where $H$ is the Higgs field which have vanishing VEV and $c$ and $\tau$ are the ghosts. As noted above, the quadratic term $gu_0 \Phi_2 A_\mu$ in the $L_{AH}(H + v, \Phi_2, A_\mu)$ cancels out with that in $\frac{f(A_\mu, \Phi_2)^2}{\delta \theta}c$ and the propagator between $\Phi_2$ and $A_\mu$ field does not exists in the total Lagrangian in the broken symmetry phase $L_{BS}(H, \Phi_2, A_\mu)$. In order to obtain the one-loop effective potential, we shift $H$ with $H + \phi$ where $\phi$ is the classical field to obtain $\mathcal{T}$ defined as[7]

$$\mathcal{T} = L_{BS}(H + \phi, \Phi_2, A_\mu) - L_{BS}(\phi, \Phi_2, A_\mu) - \left[ \frac{\delta L_{BS}(H, \Phi_2, A_\mu)}{\delta H} \right]_{H=\phi} H.$$

(6)
Then the quadratic parts of $\overline{L}$ in momentum space are given as

$$
\frac{1}{2} H D_H^{-1} H + \frac{1}{2} \left( \Phi_2 A_\mu \right) \left( D_G^{-1} B_\mu \right) \left( D^{-1}_\mu \right) \left( \Phi_2 \right) - \pi D^{-1}_c,
$$

where

$$
D_H^{-1} = p^2 + m_H^2,
$$

$$
D_G^{-1} = p^2 + m_G^2 + \xi g^2 v^2,
$$

$$
D^{-1}_{\mu\nu} = (p^2 + m_A^2)(\delta_{\mu\nu} - \frac{p_\mu p_\nu}{p^2}) + (\frac{p^2}{\xi} + m^2_A) \frac{p_\mu p_\nu}{p^2},
$$

$$
B_\mu = g \phi p_\mu,
$$

and

$$
D_g^{-1} = p^2 + m_g^2,
$$

with

$$
m_H^2 = m^2 + \frac{\lambda}{2}(\phi + v)^2, m_G^2 = m^2 + \frac{\lambda}{6}(\phi + v)^2, m_A = g(\phi + v), m_g^2 = \xi g v m_A.
$$

Let us define $X^{-1}_{\mu\nu}$ as

$$
X^{-1}_{\mu\nu} \equiv D^{-1}_{\mu\nu} + B_\mu D_G B_\nu.
$$

By using Eqs. (9), (10) and (11), we obtain

$$
X^{-1}_{\mu\nu} = (p^2 + m_A^2)(\delta_{\mu\nu} - \frac{p_\mu p_\nu}{p^2}) + \frac{D(p^2)}{m_G^2 + \xi m_A^2} \frac{p_\mu p_\nu}{\xi p^2},
$$

where

$$
D(p^2) = p^4 + (m_G^2 + 2 \xi g v m_A) p^2 + \xi m_A^2 (m_G^2 + \xi g^2 v^2) \equiv (p^2 + m_+^2)(p^2 + m_-^2).
$$

Then the one-loop effective potential $V_1$ is given by [8]

$$
V_1 = -\frac{h}{2} Tr \ln D_H^{-1} - \frac{h}{2} Tr \ln D_G^{-1} - \frac{h}{2} Tr \ln X^{-1} + h Tr \ln D_g^{-1}.
$$

By performing the one-loop momentum integral in $D \equiv 4 - 2\varepsilon$ dimension[9], we can obtain the renormalized one-loop effective action in the MS scheme as

$$
V = V_0 + V_1 = \frac{1}{2} m^2 (\phi + v)^2 + \frac{1}{24} \lambda (\phi + v)^4 + \frac{h}{16 \pi^2} \left\{ \frac{1}{4} m_H^4 (\ln m_H^2 - \frac{3}{2}) + \frac{1}{4} m_A^4 (\ln m_A^2 - \frac{3}{2}) + \frac{1}{4} m_g^4 (\ln m_g^2 - \frac{3}{2}) \right\} + \frac{3}{4} m_A^2 (\ln m_A^2 - \frac{5}{6}) - \frac{1}{2} m_g^2 (\ln m_g^2 - \frac{3}{2}) \right\}.
$$

where $\ln X \equiv \frac{X}{4 \pi^2} + \gamma$. The effective potential is independent of the renormalization mass scale $\mu$ and should satisfy the RG equation[10]

$$
\mu \frac{dV}{d\mu} = \left( \mu \frac{\partial}{\partial \mu} + \beta \frac{\partial}{\partial \lambda} + \beta_g \frac{\partial}{\partial g} + \beta_{m^2} \frac{\partial}{\partial m^2} + \gamma \frac{\partial}{\partial \gamma} + \gamma \frac{\partial}{\partial v} \right) V = 0.
$$
By using the RG functions in the MS scheme \[11\] \[12\]

\[
\beta_\lambda = \mu \frac{d\lambda}{d\mu} = \frac{\hbar}{(4\pi)^2} \left( \frac{10}{3} \lambda^2 - 12 \lambda g^2 + 36 g^4 \right) + \cdots,
\]

\[
\beta_g = \mu \frac{dg}{d\mu} = \frac{\hbar}{(4\pi)^2} \frac{g^3}{3} + \cdots,
\]

\[
\beta_{m^2} = \mu \frac{dm^2}{d\mu} = \frac{\hbar}{(4\pi)^2} \left( \frac{4}{3} \lambda - 6 g^2 \right) m^2 + \cdots,
\]

\[
\gamma_\phi = \frac{\mu}{\phi} \frac{d\phi}{d\mu} = \frac{\hbar}{(4\pi)^2} \left( 3 - \xi \right) g^2 + \cdots. \quad (20)
\]

Then, in order to satisfy the RG equation given in Eq. (19), we should have

\[
\gamma_v = \frac{\mu}{v} \frac{dv}{d\mu} = \frac{\hbar}{(4\pi)^2} (3 + \xi) g^2 + \cdots, \quad (21)
\]

which agrees with \[2\] \[3\]. The perturbative expansion of the VEV \( v = v_0 + \hbar v_1 + \hbar^2 v_2 + \cdots \) can be obtained from the no-tadpole condition

\[
\frac{\delta V}{\delta \phi} |_{\phi = 0, v = v_0} = 0, \quad (22)
\]

as

\[
\frac{\delta V_0}{\delta \phi} |_{\phi = 0, v = v_0} = 0, \quad (23)
\]

and

\[
v_1 \left[ \frac{\delta^2 V_0}{\delta \phi \delta v} \right] |_{\phi = 0, v = v_0} + \left[ \frac{\delta V_1}{\delta \phi} \right] |_{\phi = 0, v = v_0} = 0. \quad (24)
\]

By substituting the effective potential given in Eq. (18) into Eqs. (22) and (23), we obtain the perturbative expansion of the VEV up to one-loop as \( v_0 = \sqrt{-m^2_A} \) and

\[
v_1 = -\frac{3}{\lambda v_0^2} m_2^0 \left[ \frac{\delta V_1}{\delta \phi} \right] |_{\phi = 0, v = v_0} = \frac{1}{32\pi^2 m_2^0} \left[ \frac{1}{2} m_2^A \frac{\partial m_2^A}{\partial \phi} (\ln (m_2^A) - 1) + \frac{1}{2} m_2^+ \frac{\partial m_2^+}{\partial \phi} (\ln (m_2^+)) - 1 \right]
\]

\[
+ \frac{1}{2} m_2^- \frac{\partial m_2^-}{\partial \phi} \ln (m_2^-) - 1) + \frac{3}{2} m_2^A \frac{\partial m_2^A}{\partial \phi} (\ln m_2^A - 1) - m_2^X \frac{\partial m_2^X}{\partial \phi} (\ln m_2^X - 1) \right] |_{\phi = 0, v = v_0}. \quad (25)
\]

From Eq. (16) one can see that by writing \( m_2^X = p \pm \sqrt{q} \), both \( q \) and \( \frac{\partial q}{\partial \phi} \) goes to zero in the limit \( \phi \to 0 \) and \( v \to v_0 \) and hence \( \frac{\partial m_2^X}{\partial \phi} \) as well as \( v_1 \) have finite limit in \( R_\xi \) gauge. Now, let us consider IR divergence of the VEV in the case of the two-loop effective potential \[13\] \[14\] where the term which can give the IR divergence in effective potential \( V_2 \) and the VEV in case of the Landau gauge \( (\xi = 0) \) is \[4\] \[5\]

\[
m_2^X (m_2^2 + \xi g^2 v^2) \ln m_2^X \ln (m_2^2 + \xi g^2 v^2), \quad (26)
\]

coming from the two-loop Feynman diagram

\[
G \quad \circ \quad \circ \quad X, \quad (27)
\]

where \( G \) is the Goldstone boson and \( X \) can be Higgs, gauge boson or ghost. Then, by noting that the mass of the Goldstone boson given as \( m_2^X + \xi g^2 v^2 \) in the \( R_\xi \) gauge does not vanish in the limit \( v \to v_0 \) and \( \phi \to 0 \) when \( \xi \neq 0 \), the corresponding two-loop VEV obtained \[ \frac{\partial m_2^X}{\partial \phi} |_{\phi = 0, v = v_0} \] does not have an IR divergence in \( R_\xi \) gauge as long as \( \xi \neq 0 \).
B. The general gauge

The general gauge is defined as

\[
f(\Phi_1, \Phi_2, A_\mu) = \partial_\mu A_\mu - u\xi g \Phi_1 \Phi_2,
\]

which becomes the Fermi gauge when \( u = 0 \) and \( R_\xi \) gauge when \( u = 1 \). The resulting Lagrangian in the symmetric phase is given by \([6]\)

\[
L_{SYM}(\Phi_1, \Phi_2, A_\mu) = L_{AH}(\Phi_1, \Phi_2, A_\mu) + \frac{1}{2\xi} f(\Phi_1, \Phi_2, A_\mu)^2 + \tau (-\partial^2 + u\xi g(\Phi_1^2 - \Phi_2^2))c + \text{counter - terms}.
\]

When \( m^2 < 0 \), the \( O(2) \) symmetry breaks down spontaneously and we substitute \( \Phi_1 \) with \( H + v \) where \( H \) have vanishing VEV to obtain the Lagrangian density in the broken phase \( L_{BS}(H, \Phi_2, A_\mu, v) \) as

\[
L_{BS}(H, \Phi_2, A_\mu, v) = [L_{SYM}(\Phi_1, \Phi_2, A_\mu)]_{\Phi_1 \rightarrow H + v}.
\]

Recently, we have shown that if the Lagrangian in the symmetric phase \( L_{SYM}(\Phi_1, \Phi_2, A_\mu) \) and that in broken symmetry phase \( L_{BS}(H, \Phi_2, A_\mu, v) \) is related as in Eq.(30), we can prove that \( \gamma_v = \gamma_\phi \) and the RG functions of the broken symmetry phase is same as that of the RG functions of the symmetric phase\([15] [16]\). Note that in case of the \( R_\xi \) gauge, the gauge fixing and the ghost terms given in Eq.(5) does not have this relation. In case of the \( R_\xi \) gauge, let us define \( \xi \) as in case of the \( BS \)

\[
\xi = \frac{1}{2\xi} f(\Phi_1, \Phi_2, A_\mu)^2 + \text{counter terms}.
\]

Then the quadratic parts of \( \mathcal{L} \) in momentum space is given by

\[
\frac{1}{2} H D^{-1}_H H + \frac{1}{2} (\Phi_2 A_\mu) \left( \begin{array}{cc} D^{-1}_{uG} & B_{u\mu} \\ B_{u\mu} & D^{-1}_{uG} \end{array} \right) \left( \begin{array}{c} \Phi_2 \\ A_\mu \end{array} \right) - \tau D_{uG}^{-1} c,
\]

where the inverse propagators \( D^{-1}_H \) and \( D^{-1}_{uG} \) are given in Eqs.\((8)\) and \((10)\) and \( D^{-1}_{uG}, B_{u\mu} \) and \( D_{uG}^{-1} \) are given by

\[
D^{-1}_{uG} = p^2 + m_G^2 + \xi u^2 m_A^2,
\]

\[
B_{u\mu} = (1 - u)m_{AP\mu},
\]

and

\[
D_{uG}^{-1} = p^2 + \xi u m_A^2.
\]

As in case of the \( R_\xi \) gauge, let us define \( X^{-1}_{u\mu\nu} \) as

\[
X^{-1}_{u\mu\nu} \equiv D^{-1}_{u\mu\nu} + B_{u\mu} D_{uG} B_{u\nu} = (p^2 + m_A^2)(\delta_{\mu\nu} - \frac{p_\mu p_\nu}{p^2}) + \frac{D_u(p^2)}{m_G^2 + \xi u^2 m_A^2} \frac{p_\mu p_\nu}{\xi p^2},
\]

where

\[
D_u(p^2) = p^4 + (m_G^2 + 2 u\xi m_A^2)p^2 + \xi m_A^2 (m_G^2 + \xi u^2 m_A^2) \equiv (p^2 + m^2_{+}) (p^2 + m^2_{-}),
\]

to obtain the one-loop effective potential \( V_1 \) as

\[
V_1 = -\frac{1}{2} Tr \ln D^{-1}_H - \frac{1}{2} Tr \ln D^{-1}_{uG} - \frac{1}{2} Tr \ln X^{-1}_u + h Tr \ln D_{uG}^{-1}.
\]
Again, by performing the one-loop momentum integral in $D = 4 - 2\varepsilon$ dimension, we can obtain the renormalized one-loop effective action in the MS scheme as

$$V = V_0 + V_1 = \frac{1}{2}m^2(\phi + v)^2 + \frac{1}{24}\lambda(\phi + v)^4 + \frac{h}{16\pi^2}\left[\frac{1}{4}m_4^4(\ln m_H^2 - \frac{3}{2}) + \frac{1}{4}m_4^m(\ln m_{u_+}^2 - \frac{3}{2}) + \frac{1}{4}m_4^m(\ln m_{u_-}^2 - \frac{3}{2}) + \frac{3}{4}m_4^2(\ln m_{u_A}^2 - \frac{5}{6}) - \frac{1}{2}\xi^2 u^2 m_A^2(\ln \xi u m_{u_A}^2 - \frac{3}{2})\right].$$

(39)

The renormalization of the wave function for the scalar field can be done by extracting the $\frac{1}{4}$ pole term from the Feynman diagram $\Phi_1 \Phi_2 \Phi_3$ by using the Feynman rules for the vertex $\Phi_1 \Phi_2 A_\mu$ given as

$$\Phi_1 \Phi_2 A_\mu = (1 + u) g k_{1\mu} + (1 - u) g k_{2\mu},$$

(40)

and the propagators $D_{\mu\nu}$ and $D_{uG}$ given in Eqs.(10) and (33) for the Goldstone field $\Phi_2$ and the photon field $A_\mu$. As a result, we obtain

$$\Phi_{1B} = \left\{1 + \frac{h}{2\varepsilon(4\pi)^2}(3 - \xi + 2u \xi)g^2 + \cdots\right\} \Phi_1,$$

(41)

and obtain the one-loop gamma function as

$$\gamma_{u\phi} = \frac{h}{(4\pi)^2}(3 - \xi + 2u \xi)g^2 + \cdots,$$

(42)

which agrees with that of the Fermi gauge when $u = 0$ [11] and that of the $\overline{R}_\xi$ gauge[6] when $u = 1$. Then, one can check that the RG equation given in Eq.(19) is satisfied with $\beta_\lambda$ and $\beta_{u^2}$ given in Eq.(20) and with $\gamma_v = \gamma_{u\phi}$. As in case of $R_\xi$ gauge, we can obtain $v_0 = \sqrt{-\frac{6m_\lambda^2}{\tau}}$ and

$$v_1 = \frac{1}{32\pi^2 m_G^4}\left[\frac{1}{2}m_H^2 \frac{\partial m_H^2}{\partial \phi}(\ln(m_H^2) - 1) + \frac{1}{2}m_{u_+}^2 \frac{\partial m_{u_+}^2}{\partial \phi}(\ln(m_{u_+}^2) - 1) + \frac{1}{2}m_{u_-}^2 \frac{\partial m_{u_-}^2}{\partial \phi}(\ln(m_{u_-}^2) - 1)ight] + \frac{3}{2}m_A^2 \frac{\partial m_A^2}{\partial \phi}(\ln m_{u_A}^2 - 1) \big|_{\phi=0, v=v_0}.$$  

(43)

From Eq.(37), by writing $m_{u_\pm}^2 = p_u \pm \sqrt{q_u}$ we have

$$q_u = m_G^4 + 4\xi(u - 1)m_A^2m_G^2.$$  

(44)

Then one can see that $[\sqrt{q_u}]_{\phi=0, v=v_0} = 0$ and $[\frac{\partial m_{u_\pm}^2}{\partial \phi}]_{\phi=0, v=v_0} \neq 0$ when $u \neq 1$ or $\xi \neq 0$ and as a result, $[\frac{\partial m_{u_\pm}^2}{\partial \phi}]_{\phi=0, v=v_0}$ diverges. Although $[\frac{\partial m_A^2}{\partial \phi}]_{\phi=0, v=v_0}$ diverges, since $[m_{u_+}^2]_{\phi=0, v=v_0} = [m_{u_-}^2]_{\phi=0, v=v_0} = u\xi g^2 v_0^2$, these divergence contained in $\frac{1}{2}m_{u_+}^2 \frac{\partial m_{u_+}^2}{\partial \phi}(\ln(m_{u_+}^2) - 1)$ terms of Eq.(43) cancels out and the one-loop VEV $v_1$ converges for all values of $v$. In case of $\overline{R}_\xi$ gauge we obtain

$$v_1 = -\frac{v_0}{32\pi^2}\left\{\lambda(\ln(-2m^2) - 1) + \xi g^2(\ln(-\frac{6\xi^2 m^2}{\lambda}) - 1) + \frac{18g^4}{\lambda}(\ln(-\frac{6g^4 m^2}{\lambda}) - 1)\right\},$$

(45)

so that

$$\mu \frac{\partial v_1}{\partial \mu} = \frac{1}{16\pi^2}(\lambda + \xi g^2 + \frac{18g^4}{\lambda})v_0.$$  

(46)

By using Eq.(20) we have

$$\mu \frac{dv_0}{d\mu} = \frac{h}{16\pi^2}(\lambda + 3g^2 - \frac{18g^4}{\lambda})v_0,$$

(47)
and hence up to $O(h)$, we obtain

$$
\mu \frac{dv}{d\mu} = \frac{h}{16\pi^2}(3 + \xi)g^2 v.
$$

which is consistent with Eq.(42) when $u = 1$. Finally, consider the IR divergence of the VEV in case of the two-loop effective potential. Since the mass of the Goldstone boson given by $m_G^2 + \xi m_A^2$ in $R_\xi$ gauge does not vanish in the limit $v \to v_0$ and $\phi \to 0$ when $\xi \neq 0$, the term

$$
m_X^2 (m_G^2 + \xi m_A^2) \ln(m_X^2 + \ln(m_G^2 + \xi m_A^2))
$$

coming from the Feynman diagram given in Eq.(27) in two-loop effective potential and the resulting VEV coming from $\left[ \frac{dv}{d\phi} \right]_{\phi=0, v=v_0}$ which can give the IR divergence in case of the Landau gauge ($\xi = 0$) where Goldstone boson is massless[4][5] do not have an IR divergence as long as $\xi \neq 0$.

In this paper, we have investigated the RG function and the perturbative expansion of the VEV of the Abelian Higgs model th in $R_\xi$ and the general gauge which extrapolate between the Fermi and the $R_\xi$ gauge by requiring the RG invariance of one-loop effective potential. In case of the $R_\xi$ gauge, the gamma function of the VEV that satisfy the RG equation for the effective potential was different from the gamma function of the scalar field. In case of the general gauge, the gamma function of the VEV obtained from the RG invariance of the effective potential was same as that of the scalar field. When $u = 1$ which corresponds to the $R_\xi$ gauge and $\xi = 0$ which corresponds to Landau gauge in general gauge, the one-loop VEV obtained from the no-tadpole condition do not have the IR divergence and give correct RG behavior. Both in $R_\xi$ and $\overline{R_\xi}$ gauge, the two-loop VEV have an IR divergence in Landau gauge($\xi = 0$) and these IR divergence do not occur when $\xi \neq 0$.

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