Renormalization of the Electroweak Theory in the Background Field Method

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Abstract

The applicability of the background field method in spontaneously broken gauge theories is examined with new features emphasized. An explicit one loop analysis in the electroweak theory shows that the method can be consistently implemented in the on-shell renormalization scheme, and that the choice of the background gauge cannot be arbitrary and must be fixed in the Landau gauge if one calculates scattering amplitudes involving unphysical Goldstone bosons. Some possible applications are also briefly indicated.

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background field method, electroweak theory, renormalization, Ward identities.
Although the S matrix keeps gauge invariant in gauge theories, this property is lost at intermediate stages in the conventional procedure of quantization. The background field method (BFM) is a technique which can preserve explicit gauge invariance at intermediate stages, so that Green functions satisfy the naive Ward identities which are more restrictive than the Slavnov-Taylor identities from BRS invariance and therefore lead to a simpler renormalization structure.

The method was developed some time ago for unbroken gauge theories\(^1\), in a manner analogous to the conventional one and The applicability of the background field method in spontaneously broken gauge theories is examined with new features emphasized. An explicit one loop analysis in the electroweak theory shows that the method can be consistently implemented in the on-shell renormalization scheme, and that the choice of the background gauge cannot be arbitrary and must be fixed in the Landau gauge if one calculates scattering amplitudes involving unphysical Goldstone bosons. Some possible applications are also briefly indicated.

applicable to all orders in perturbation theory. However, a parallel work, in the case of spontaneously broken gauge theories in general and the electroweak theory in particular, has not been available to our knowledge\(^1\). Although the formal construction in the two cases is similar, some subtle points associated with renormalization may be easily missed since bare quantities are usually involved in formal manipulations. An explicit study is therefore necessary for the resolution of these points. Indeed, spontaneously broken gauge theories differ from unbroken ones mainly in their patterns of symmetry realization. As will be pointed out below, the constraints on renormalization constants imposed by the BFM are not apparently understand-\(^1\)

\(^1\)The method was used to study symmetry restoration in spontaneously broken gauge theories in a constant electromagnetic background field in Ref. \(^2\), which however does not apply to the construction of Green functions. A formally similar construction also appeared in Ref. \(^3\) in the context of the gauge coupling running of the QCD and electroweak theory in grand unified theories. Recently the method was compared to the pinch technique \(^4\) in Ref. \(^5\) which however only treats bare quantities.
able in spontaneously broken gauge theories as in unbroken theories. Furthermore, renormalization constants also depend on subtraction schemes used. While in QCD mass-independent subtractions (e.g. the minimal subtraction and its modifications) are usually used, in the electroweak theory many physical scales are involved and mass-dependent subtractions (e.g. the on-shell renormalization scheme) are usually favoured. It is not clear at all whether such subtractions are consistent with the constraints imposed by the BFM in spontaneously broken gauge theories. It is probable that the consistent implementation of the BFM inversely picks out appropriate subtraction schemes. Another feature not encountered in the QCD case is that the gauge choice for background gauge fields cannot be arbitrary and must be fixed in the Landau gauge if one calculates scattering amplitudes involving unphysical Goldstone bosons, while in QCD the gauge choices for quantum and background gauge fields are independent and can be both arbitrary.

In this letter we study the renormalization of the electroweak theory in the BFM by an explicit one loop analysis. Certainly, the full power of the method can be best enjoyed in higher order calculations, yet the feasibility of the method must begin at the one loop level and some features can also be glimpsed at this level. We will show that the on-shell renormalization scheme is consistent with the BFM and that this consistency can also be set up by examining the renormalized Ward identities. Finally we summarize our results and briefly mention some possible applications.

We begin with the formal construction of the bare Lagrangian. The generating functional for connected Green functions of background fields, $W[J, ˆ{f}]$, is defined by

$$\exp(iW[J, ˆ{f}]) = \int (Df\ D\omega\ D\bar{\omega}) \exp i \int d^4x [L_{\text{classical}}(f + ˆ{f}) + L_{\text{g.f.}} + L_{\text{FP}} + Jf], \quad (1)$$

where $f$ collectively stands for the generic background field and $f$ the generic quantum field to be integrated over, $\omega$ and $\bar{\omega}$ are Faddeev-Popov ghosts. Note that the external source $J$ is introduced only for the $f$ field. The gauge fixing term $L_{\text{g.f.}}$ is
so chosen as to preserve the background gauge invariance of $W[0, \hat{f}]$ while breaking the quantum gauge invariance. The background field effective action is defined by the Legendre transform,

$$\Gamma[\hat{f}, \tilde{f}] = W[J, \hat{f}] - \int d^4x J \tilde{f},$$

$$\tilde{f} = \frac{\delta W}{\delta J}.$$ (2)

$\Gamma[0, \hat{f}]$ is the gauge invariant effective action that one computes in the BFM. When augmented by a gauge-fixing term for background gauge fields, it can be used to generate the $S$ matrix by constructing trees using its vertices and propagators.

Now we restrict ourselves to the electroweak theory of $SU(2) \otimes U(1)$. For simplicity only the bosonic sector is included. The inclusion of fermions is straightforward and does not pose any new problems since fermions constitute a gauge invariant subset by themselves and there is no need to discriminate quantum from background fields. We write,

$$L_{\text{classical}} = L_Y + L_{\text{scalars}},$$

$$L_Y = -\frac{1}{4} W_a^{\mu \nu} W^a_{\mu \nu} - \frac{1}{4} B_{\mu \nu} B^{\mu \nu},$$

$$W^a_{\mu \nu} = \partial_\mu W^a_\nu - \partial_\nu W^a_\mu + g_2 \epsilon^{abc} W^b_\mu W^c_\nu,$$

$$B_{\mu \nu} = \partial_\mu B_\nu - \partial_\nu B_\mu,$$

$$L_{\text{scalars}} = (D_\mu \Phi)^\dagger (D^\mu \Phi) + \mu^2 \Phi^\dagger \Phi + \lambda (\Phi^\dagger \Phi)^2,$$

$$D_\mu \Phi = (\partial_\mu - ig_2 \frac{\tau^a}{2} W^a_\mu - ig_1 \frac{1}{2} B_\mu) \Phi,$$

$$\Phi = \left( \frac{1}{\sqrt{2}} (\phi_1 + v + i \phi_2) \right),$$

where the usual notations for fields and parameters have been used. $L_{\text{classical}}(f + \hat{f})$ is obtained from the above by the replacements,

$$W^a_\mu \rightarrow W^a_\mu + \hat{W}^a_\mu, \quad B_\mu \rightarrow B_\mu + \hat{B}_\mu,$$

$$\Phi \rightarrow \Phi + \hat{\Phi}, \quad \Phi = \left( \frac{1}{\sqrt{2}} (\phi_1 + i \phi_2) \right), \quad \hat{\Phi} = \left( \frac{1}{\sqrt{2}} (\phi_1 + v + i \phi_2) \right).$$ (4)

Note that the background scalar $\hat{\Phi}$ develops a non-zero VEV while the quantum scalar $\Phi$ does not. The invariance of $W[0, \hat{f}]$ (and thus $\Gamma[0, \hat{f}]$) under gauge
transformations of the $\hat{f}$ requires that $\mathcal{L}_{g.f.}$ (and thus $\mathcal{L}_{\text{FP}}$) be invariant under the union of gauge transformations of the $\hat{f}$ and appropriate changes of the integration variables $f$. A straightforward generalization of the so-called “background field $R$ gauge” in QCD [1] and the conventional $R_\xi$ gauge in the electroweak theory [7] leads to the following construction [2][3][8][5],

$$
\mathcal{L}_{g.f.} = -\frac{1}{2\xi_1}[(\partial_\mu \delta^{ac} + g_2\epsilon^{abc}\hat{W}^b_\mu)W^c,\mu - ig_2\xi_2\frac{1}{2}(\Phi^i,\tau^a\Phi - \Phi^i,\tau^a\Phi)]^2
$$

Note that the two terms in the first pair of square brackets transform separately in the adjoint representation under the aforementioned union. $\mathcal{L}_{\text{FP}}$ is then determined by using quantum gauge transformations that keep $\mathcal{L}_{\text{classical}}(f + \hat{f})$ invariant.

To go beyond tree level, renormalizations have to be carried out. The only renormalizations required in the BFM are those of physical parameters, background fields and gauge parameters $\xi_i$ [1]. However, for the purpose of calculating 1PI functions of background fields to one loop, even the renormalization of $\xi_i$ is not required since $\xi_i$ appears only in vertices that are at least quadratic in quantum fields. For the other renormalizations we take as usual [9],

$$
\dot{B}_\mu \rightarrow (Z^B_2)^{1/2}\dot{B}_\mu, \quad \dot{W}^a_\mu \rightarrow (Z^W_2)^{1/2}\dot{W}^a_\mu,
$$

$$
(\dot{\phi}^\pm, \dot{\phi}_1, \dot{\phi}_2) \rightarrow (Z_\phi)^{1/2}(\dot{\phi}^\pm, \dot{\phi}_1, \dot{\phi}_2),
$$

$$
g_1 \rightarrow Z^B_1(Z^B_2)^{-3/2}g_1, \quad g_2 \rightarrow Z^W_1(Z^W_2)^{-3/2}g_2,
$$

$$
\mu^2 \rightarrow (Z_\phi)^{-1}(\mu^2 - \delta\mu^2), \quad \lambda \rightarrow (Z_\phi)^{-2}Z_\lambda, \quad \nu \rightarrow (Z_\phi)^{1/2}(v - \delta v).
$$

Explicit gauge invariance of $\Gamma[0, \hat{f}]$ requires that covariant derivatives be renormalized in the following way, e.g.,

$$
\dot{D}_\mu \hat{\Phi} = (\partial_\mu - ig_2\frac{1}{2}\hat{B}_\mu - ig_2\frac{\tau^a_2}{2}\hat{W}^a_\mu)\hat{\Phi} \rightarrow (Z_\phi)^{1/2}\dot{D}_\mu \hat{\Phi},
$$

where, on the rhs, $v$ has been replaced by $v - \delta v$ and all quantities in $\dot{D}_\mu \hat{\Phi}$ are renormalized or finite. But this is possible only if

$$
Z^W_1 = Z^W_2, \quad Z^B_1 = Z^B_2,
$$

5
\[ \delta v = \text{finite.} \quad (9) \]

Eqn. (8) is not self-evident, at least for the finite parts since \( Z_2 \)'s are generally related to wavefunction renormalizations while \( Z_1 \)'s go into mass renormalizations.

( This would be clearer if \( Z_1 \) were replaced by \( Z_1(Z_2)^{3/2} \). ) The similar situation does not occur in QCD where the symmetry is unbroken. Nevertheless, we will show explicitly that the constraints of Eqns. (8) and (9) are satisfied at least in the on-shell renormalization scheme so that naive Ward identities are indeed saturated by 1PI functions of background fields.

The Feynman rules and counterterms can now be written down. Due to the nonlinearity of gauge condition functions and the discrimination between quantum and background fields, the Feynman rules are different from those in the conventional approach. It is worth mentioning that the background Goldstone bosons are massless since the background gauge has not been fixed. To avoid tree level \( Z_\mu - A_\mu \) mixing we assume \( \xi_1 = \xi_2 = \xi \). For simplicity we work in the quantum \( 't \) Hooft – Feynman gauge, \( \xi = 1 \). We found that calculations in the \( \xi = 1 \) gauge in the BFM are much simpler than in the conventional approach. The renormalization constants are determined by the following conditions:

1. vanishing tadpole
   \[ \hat{T} = 0 \]

2. on-shell definition of masses
   \[ \text{Re} \hat{\Sigma}^W_T(M_W^2) = \text{Re} \hat{\Sigma}^Z_T(M_Z^2) = \text{Re} \hat{\Sigma}^\phi_1(M_H^2) = 0 \]

\[ ^2 \text{The following conventions are assumed: (1) } \Sigma \text{'s or } \Gamma \text{'s with a hat, a tilde or nothing are respectively renormalized, counterterm or unrenormalized one loop contributions; (2) factor } i \text{ has been separated out from } \Sigma \text{'s or } \Gamma \text{'s; (3) for gauge bosons } \Sigma_T \text{ refers to the coefficient of } g_{\mu\nu}; (4) \text{ momenta are taken to be incoming.} \]
3. explicit $U(1)_{\text{e.m.}}$ symmetry

$$\left[ \frac{1}{p^2} \hat{\Sigma}_T^A(p^2) \right]_{p^2=0} = 0, \quad \hat{\Sigma}_T^A(0) = 0$$

4. unity residue of the $\hat{\phi}^\pm$ propagator

$$\left[ \frac{1}{p^2} \hat{\Sigma}_T^{\hat{\phi}^\pm}(p^2) \right]_{p^2=0} = 0$$

5. electromagnetic coupling defined in the Thomson limit

$$\left[ \hat{\Gamma}_\mu^A \hat{\phi}^+ + \hat{\Gamma}_\mu^A \hat{\phi}^- (p, p_+, p_-) \right]_{p^2, p_+^2 \to 0} = e(p_+ - p_-)_{\mu} \cdot \delta Z^A$$

The electromagnetic coupling definition requires unity residues of the photon and the reference particle propagators. Since there is no extra freedom to require unity residue for the $\hat{\phi}^\pm$ propagator when this has been done for the $\hat{A}$ propagator, we have used $\hat{\phi}^\pm$ as the reference particle and required unity residue of its propagator instead of the $\hat{\phi}_1$ propagator. If we include fermions we may use, e.g. the electron, as the reference particle and obtain the same result for the charge renormalization.

The counterterm contributions to the above quantities are,

$$\tilde{T} = -v[\delta \mu^2 + \lambda v^2(\delta Z_\lambda - 2 \frac{\delta v}{v})],$$

$$\tilde{\Sigma}_W^\mu = -\delta Z_2^W(p^2 g_{\mu\nu} - p_{\mu}p_{\nu}) + g_{\mu\nu}(\delta M_H^2 + M_W^2 \delta Z_2^W),$$

$$\tilde{\Sigma}_Z^\mu = -\delta Z_2^Z(p^2 g_{\mu\nu} - p_{\mu}p_{\nu}) + g_{\mu\nu}(\delta M_Z^2 + M_Z^2 \delta Z_2^Z),$$

$$\tilde{\Sigma}_A^\mu = -\delta Z_2^A(p^2 g_{\mu\nu} - p_{\mu}p_{\nu}),$$

$$\tilde{\Sigma}_A^{\tilde{Z}} = -\delta Z_2^{\tilde{Z}}(p^2 g_{\mu\nu} - p_{\mu}p_{\nu}) + g_{\mu\nu}(\delta M_Z^2 + M_Z^2 \delta Z_2^{\tilde{Z}}),$$

$$\tilde{\Sigma}_Z^{\tilde{Z}} = \delta Z_2^\phi p^2 - (\delta M_H^2 + M_H^2 \delta Z_\phi),$$

$$\tilde{\Sigma}_Z^{\tilde{Z}} = \delta Z_2^\phi p^2 - [\delta \mu^2 + \lambda v^2(\delta Z_\lambda - 2 \frac{\delta v}{v})],$$

$$\tilde{\Gamma}_\mu^A \hat{\phi}^+ \hat{\phi}^- (p, p_+, p_-) = e(p_+ - p_-)_{\mu}[ \delta Z_\phi + \delta Z_1^A - \delta Z_2^A + \frac{e^2 - 2}{2c_s}(\delta Z_1^A - \delta Z_2^A)].$$

\(^3\text{Only the charge interaction part is relevant.}\)
We will not list the explicit expressions of the unrenormalized one loop contributions.

In determining renormalization constants we note the following crucial properties:

\[
\begin{align*}
\Sigma_{\mu
u}^A &\propto (p^2 g_{\mu
u} - p_\mu p_\nu), \\
\Sigma_{\mu
u}^{\bar{A}} &\propto (p^2 g_{\mu
u} - p_\mu p_\nu), \\
\Sigma^\phi_\mu(0) & = \Sigma^\phi_{\mu} + \delta\mu^2 + \lambda v^2(\delta Z_\lambda - 2\delta v), \\
\Gamma_{\mu}^{\phi\phi^+}(p, p_+, p_-) &\equiv -e(p_+ - p_-)\mu\delta Z_\phi.
\end{align*}
\]  

It is clear that Eqn. (8) is satisfied and the unphysical Goldstone bosons \(\phi_2\) and \(\phi^\pm\) keep massless to one loop. \(\delta v/v\) is determined from \(\delta M_\phi^2\), \(\delta M_2^2\), \(\delta Z_\phi^2\) and \(\delta Z_\phi\) and the result is finite.

The self-consistency of the BFM in the on-shell scheme may also be checked by examining Ward identities. We derive Ward identities directly for renormalized 1PI functions. In this procedure, Eqns. (8) and (9) are necessary to obtain genuine renormalized gauge transformations from the bare ones. The gauge invariance of \(\Gamma[0, \hat{f}] \equiv \hat{\Gamma}\) (Here \(\hat{\cdot}\) means both “background” and “renormalized”) gives,

\[
\begin{align*}
\frac{1}{g_2^2} \partial_\mu \frac{\delta \hat{\Gamma}}{\delta \phi_\mu} + ic(\hat{\phi}_2 + s \hat{A}_\mu) \frac{\delta \hat{\Gamma}}{\delta \phi_\mu} &= 0, \\
\frac{1}{\sqrt{g_1^2 + g_2^2}} \partial_\mu \frac{\delta \hat{\Gamma}}{\delta \phi_\mu} + i(c^2 - s^2)(\hat{\phi} \frac{\delta \hat{\Gamma}}{\delta \phi} - \hat{\phi}^+ \frac{\delta \hat{\Gamma}}{\delta \phi^+}) &= 0, \\
\frac{1}{e} \partial_\mu \frac{\delta \hat{\Gamma}}{\delta A_\mu} + i(\hat{W}_\mu \frac{\delta \hat{\Gamma}}{\delta W_\mu} - \hat{W}_\mu^+ \frac{\delta \hat{\Gamma}}{\delta W_\mu^+}) &= 0,
\end{align*}
\]

and the Hermitian conjugate of Eqn. (12). The above equations are the starting point of all identities. For example, we have in momentum space,

\[
\begin{align*}
p^\mu \hat{\Gamma}_{\mu\rho\sigma}^{\bar{A}W^+\bar{W}^-}(p, p_+, p_-) &= e[\hat{\Gamma}_{\rho\sigma}^{\bar{W}^+\bar{W}^-}(p_+) - \hat{\Gamma}_{\rho\sigma}^{\bar{W}^+\bar{W}^-}(p_-)], \\
p^\rho \hat{\Gamma}_{\mu\rho\sigma}^{\bar{A}W^+\bar{W}^-}(p, p_+, p_-) - M_W(1 - \frac{\delta v}{v})\hat{\Gamma}_{\rho\sigma}^{\phi\phi^+\phi^-(p, p_+, p_-)} &= e[\hat{\Gamma}_{\mu\rho\sigma}^{\bar{A}\bar{A}}(p) + \frac{c}{s} \hat{\Gamma}_{\mu\rho\sigma}^{\bar{A}\bar{Z}}(p) - \hat{\Gamma}_{\rho\sigma}^{\bar{W}^+\bar{W}^-}(p_-)].
\end{align*}
\]
At tree level the above are easily checked. At one loop level they are separately satisfied by counterterms and unrenormalized one loop contributions. For the former the presence of $\delta v/v$ and Eqn. (8) are crucial. Typical examples for the latter are shown in Figs. 1 and 2.

We have discussed the feasibility of the BFM in spontaneously broken gauge theories. In comparison to unbroken gauge theories new features appear that are peculiar to spontaneously broken gauge theories. These are associated with the presence of unphysical Goldstone bosons and the renormalizations of masses generated by spontaneous symmetry breaking, making the applicability of the BFM less evident. We have shown by explicit one loop calculations that the method can be consistently carried through in the on-shell renormalization scheme of the electroweak theory. Actually the method is consistent with any scheme which treats the QED subpart normally, i.e. ,with the electromagnetic coupling defined in the Thomson limit and the $U(1)_{e.m.}$ symmetry explicitly preserved all the way. ( By the way we may mention that the $\hat{A}_\mu - \hat{\phi}_{1,2}$ mixing is identically zero in the BFM. ) Due to the masslessness of unphysical Goldstone bosons ( we have actually verified this in general quantum $\xi$ gauges ) the choice of the background gauge ( parameterized by $\hat{\xi}$ and independent of the quantum gauge parameters $\xi$ ) in S matrix calculations cannot be arbitrary and must be fixed in the Landau gauge so that external propagators may be appropriately amputated, if the S matrix involves external unphysical Goldstone bosons. ( Indeed this is not the S matrix in the exact sense though the usage prevails in the literature. )

The advantage of the BFM over the conventional approach is clear. Since the Slavnov-Taylor identities are replaced by the naive Ward identities, the renormalization structure is simplified and less independent renormalization constants are needed. Technically, calculations in the 't Hooft-Feynman gauge of the BFM are
easier than in the conventional approach. It is reasonable to expect that this simplicity can be best enjoyed in even higher order calculations. Since the relations between 1PI functions are much simplified in the BFM, we expect that some processes, which involve formidably complicated gauge cancellations in the conventional approach, e.g., the longitudinal gauge boson scattering in the electroweak theory, can be more easily computed in the BFM. The BFM may also find its applications in the electroweak chiral Lagrangians where gauge non-invariant terms are generally involved when calculations are carried out to higher orders in general $R_\xi$ gauges of the conventional method. Work on these aspects is now in progress.

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Appendix

Some notations are listed in this appendix.

\begin{equation}
Z_\mu = c W^3_\mu - s B_\mu, \quad A_\mu = s W^3_\mu + c B_\mu, \\
\hat{Z}_\mu = c \hat{W}^3_\mu - s \hat{B}_\mu, \quad \hat{A}_\mu = s \hat{W}^3_\mu + c \hat{B}_\mu,
\end{equation}

where $c$ and $s$ are defined by renormalized couplings,

\begin{equation}
c = \frac{g_2}{\sqrt{g_1^2 + g_2^2}}, \quad s = \frac{g_1}{\sqrt{g_1^2 + g_2^2}}.
\end{equation}

Some renormalization constants appearing in the text are \((i = 1, 2)\)

\begin{equation}
\begin{align*}
Z_i^Z &= c^2 Z_i^W + s^2 Z_i^B, \quad Z_i^A = s^2 Z_i^W + c^2 Z_i^B, \\
\delta Z_i^{AZ} &= cs(\delta Z_i^W - \delta Z_i^B), \\
\delta M_{W_i}^2 &= M_{W_i}^2(-2\frac{\delta v}{v} + 2\delta Z_1^W - 3\delta Z_2^W + \delta Z_\phi), \\
\delta M_{Z_i}^2 &= M_{Z_i}^2(-2\frac{\delta v}{v} + 2\delta Z_1^Z - 3\delta Z_2^Z + \delta Z_\phi), \\
\delta M_{H_i}^2 &= M_{H_i}^2(-3\frac{\delta v}{v} + \frac{3}{2}\delta Z_\lambda - \delta Z_\phi + \frac{\delta \mu^2}{M_{H_i}^2}).
\end{align*}
\end{equation}

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**Figure Captions**

1. A typical example of unrenormalized one loop contributions to Eqn. (15). The virtual particles in loops are the quantum fields $W^\pm$, $Z$ or $A$. The solid (dashed) external lines are the background gauge (unphysical Goldstone) bosons.

2. A typical example of unrenormalized one loop contributions to Eqn. (16).
This figure "fig1-1.png" is available in "png" format from:

http://arxiv.org/ps/hep-ph/9409401v1