Regularization Dependence of
Quadratic Divergence Cancellations

Gary Kleppe

Department of Physics
Virginia Polytechnic Institute
and State University
Blacksburg VA 24060

ABSTRACT: Certain results related to the cancellation of quadratic divergences, which had been obtained using dimensional reduction, are reconsidered using a nonlocal regulator. The results obtained are shown to depend strongly on the regulator. Specifically, it is shown that a certain result of Al-sarhi, Jack, and Jones no longer holds, even if a nontrivial measure factor is used; also that there are no values of the top and Higgs mass for which the one-loop quadratic divergence in the standard model cancels independently of the renormalization scale, whether or not strong interaction effects are ignored.
1. Introduction

Over the last ten years, there has been some interest \[1,2,3\] in studying the quadratic
divergences of gauge theories, particularly those of the standard model. The smallness of
the Higgs mass is looked at as unnatural in the presence of the quadratic divergences. From
the point of view of standard renormalization theory it is hard to see what the difficulty is.
From this perspective, the quadratic divergences only exist to be subtracted away. There is
not even any good definition of “small”, because the cutoff scale is taken to infinity.

One may instead take the viewpoint of Wilson \[4\], in which the cutoff is thought of as
some large but finite scale. Field theories such as the standard model are thought of as
effective theories which are only valid at scales below the cutoff. Above the cutoff scale some
unknown new physics takes over. Thus although this new physics will determine the values
of renormalized parameters, we expect that for quadratically divergent parameters the values
should be of the order of the cutoff scale squared times the appropriate coupling constant(s);
values which are much smaller are unnatural, as they depend on precise fine-tuning at the
new physics scale. In the case of the Higgs mass, this would require the new physics scale
to be uncomfortably small, around the 1 TeV range.

One way to avoid this problem is through supersymmetry, which generally causes quadratic
divergences to be cancelled. However, no supersymmetric partners for known particles have
ever been detected experimentally. Therefore it is natural to look for other theories in which
the quadratic divergences cancel.

A curious result concerning quadratic divergences was found by Al-sarhi, Jack, and Jones
\[2,3\]. They considered the quadratic divergences of a theory of scalars and fermions using
dimensional reduction \[5\] and minimal subtraction. They found that by demanding that
the quadratic divergences cancel at one loop, and that this cancellation be unaffected by
the action of the renormalization group, cancellation automatically occurred at two loops;
and by demanding the two loop cancellation be invariant under the renormalization group, cancellation at three loops was automatic.

It was found that similar results did not occur at the next loop order [2], or for gauge theories [3]. However, the gauge theory result could be fixed if the peculiarities of the dimensional reduction technique were accounted for. Furthermore, it is known that dimensional reduction fails at four loops [6], and the renormalization group beta functions at this order depend on a choice of renormalization prescription, so it is possible that these failures are due to the choice of regularization and renormalization scheme. These authors [3] also consider the case of the standard model [7] using the same techniques. They find that there is no simultaneous solution of the one and two loop constraints, however the quadratic divergences at one loop cancel independently of the renormalization group scale for \( m_t = 115 \text{ GeV} \) and \( m_H = 180 \text{ GeV} \), but only if the strong interactions’ contribution to the beta functions are ignored.

There is another good reason to be suspicious of dimensional reduction. In this scheme, as in dimensional regularization, one does not see quadratic divergences directly. They are only inferred by looking at poles in lower dimensions. At one loop this cannot be argued with, but at two or more loops it is hard to justify why lower dimensional poles are related to quadratic divergences. Thus it is worth reexamining these results using a different regulator, to try to determine just how much they depend on what techniques are used.

In this paper the question of quadratic divergence cancellation is reexamined using the technique of nonlocal regularization [8,9] (a version of which was used at one loop in ref. 3). In section 2 the theory of scalars and fermions is considered, and it is shown that the result of Al-sarhi, Jack and Jones does not hold with this technique, even if a measure factor is added to the interactions. In section 3 the standard model is considered, and it is shown that with the minimal measure factor, there are no real values for \( m_t \) and \( m_H \) which make the quadratic divergence cancel independently of the scale, with or without strong interactions.
A nonminimal measure could alter this result, however.

2. Scalar-Fermi Theory

Consider the renormalizable theory of scalars and fermions whose Lagrangian is
\[
\mathcal{L} = \frac{1}{2} \phi_a (\partial^2 - m_a^2) \phi_a - \frac{1}{6} g_{abc} \phi_a \phi_b \phi_c - \frac{1}{24} \lambda_{abcd} \phi_a \phi_b \phi_c \phi_d
\]
\[-i \bar{\psi} \sigma \cdot \partial \psi - \frac{1}{2} \left[ \psi (M + Y_a \phi_a) \psi + \text{c.c.} \right] \]

where \( \psi \) is a two component spinor field, with arbitrary “flavor” indices which we will not show, and \( g \) and \( \lambda \) are totally symmetric. The metric used is \((-1,1,1,1)\). Our convention is that repeated scalar indices are summed over, even if repeated more than once as in \( m_a^2 \phi_a^2 \).

Nonlocal regularization is described in detail in ref. 9; detailed calculational techniques may be found in ref. 10. For our purposes, the technique may be summarized in two steps. The first step is to expand all propagators in the Schwinger parametrization, e.g.
\[
\frac{1}{p^2 + m^2} = \int_0^\infty \frac{d\tau}{\Lambda^2} \exp \left( -\tau \left( \frac{p^2 + m^2}{\Lambda^2} \right) \right)
\]
and to delete the regions of Schwinger parameter space for which all parameters around a closed loop in some diagram are simultaneously less than 1. The second step is to add a measure factor, a set of extra interactions which assure that the path integral measure is invariant under nonlocally distorted versions of any symmetries which were present in the original theory. In this first simple case we need not be concerned with this, since there are no symmetries in this theory which need to be preserved.

The quadratic divergences in (1) of course occur in the scalar two-point function. At one loop there are two contributions to this, shown in figures 1a and 1b.* The graph of figure 1a gives a contribution
\[
-\frac{\lambda_{abcc}}{2} \int \frac{d^4 q}{(2\pi)^4} \int_1^\infty \frac{d\tau}{\Lambda^2} \exp \left[ -\frac{\tau}{\Lambda^2} (q^2 + m_c^2) \right]
\]

* Solid lines denote scalars, dashed lines denote fermions.
\[-\frac{i\lambda_{abc}}{2(4\pi)^2} \int_1^\infty \frac{d\tau}{\tau^2} \exp \left[ -\frac{\tau m_c^2}{\Lambda^2} \right] \]

where \(\mu\) is an arbitrary scale. The graph of figure 1b gives a contribution

\[-\frac{1}{2} \text{Tr} \left\{ Y_a \int \frac{d^4 q}{(2\pi)^4} \int_{R_i^2} \frac{d\tau_1 d\tau_2}{\Lambda^4} \text{tr} \left[ (p + q) \cdot \bar{\sigma} q \cdot \sigma \right] \times \exp \left[ -\frac{1}{\Lambda^2} \left( \tau_1 q^2 + \tau_2 (p + q)^2 + \tau_1 \tau_2 M^2 \right) \right] Y_b^* + \text{c.c.} \right\} \]

Expanding (4) we get

\[\frac{i}{(4\pi)^2} \text{Tr} \left\{ Y_a \int_{R_i^2} \frac{d\tau_1 d\tau_2}{\tau_1^2 \tau_2^2} \left( \frac{2\Lambda^2}{\tau_1^2} - \frac{\tau_1 \tau_2}{\tau_1^2} \right) \exp \left[ -\frac{1}{\Lambda^2} \left( \tau_1 \tau_2 (p^2 + \tau_1 \tau_2 M^2) \right) \right] Y_b^* + \text{c.c.} \right\} + \text{finite} \]

Combining (3) with (9) we find that the one loop quadratic divergence cancels* if

\[\lambda_{abc} = 3 \text{Tr} \left[ Y_a^* Y_b + Y_b^* Y_a \right] \]

Note that in this and in the following, we neglect many graphs which have no quadratic divergences. In a similar manner to the above, we obtain the scalar field strength renormalization

\[\delta Z^\phi_{ab} = -\frac{1}{2(4\pi)^2} \text{Tr} \left[ Y_a Y_b^* + Y_b^* Y_a \right] \ln \frac{\Lambda^2}{\mu^2} \]

* The scalar one-point function also diverges quadratically. Cancellation of this would give another condition, this one involving \(g\) and \(M\) in addition to \(\lambda\) and \(Y\). For our purposes it is not necessary to consider this additional complication.
where we use a minimal subtraction scheme at scale \( \mu \).

Likewise, from the graph of figure 2 we obtain the fermion field strength renormalization

\[
\delta Z^\psi = -\frac{1}{2} \frac{\text{Tr}[Y_a Y_a^*]}{(4\pi)^2} \ln \frac{\Lambda^2}{\mu^2},
\]

from figure 3 the Yukawa coupling renormalization

\[
\delta Y_a = \frac{Y_b Y_a^* Y_b^*}{(4\pi)^2} \ln \frac{\Lambda^2}{\mu^2},
\]

and from figures 4a and 4b the scalar quartic coupling renormalization

\[
\delta \lambda_{abcd} = \frac{1}{(4\pi)^2} \left[ \frac{1}{2}(\lambda_{abef} \lambda_{cdfe} + \lambda_{acef} \lambda_{bdfe} + \lambda_{adef} \lambda_{bcfe}) 
- 2\text{Tr}[Y_a Y_b Y_c Y_d + Y_a^* Y_b^* Y_c^* Y_d + Y_a^* Y_b Y_c Y_d^* + c.c.] \right] \ln \frac{\Lambda^2}{\mu^2}
\]

There are of course mass and cubic scalar potential counterterms, but these will be unnecessary for our purposes. From the above we may calculate the beta functions for \( \lambda \) and \( Y \), which agree with the standard results (see e.g. ref. 3). Then demanding that \( \frac{\partial}{\partial \mu} \) on (10) gives zero, we get

\[
- \lambda_{abef} \lambda_{ceef} - 2\lambda_{acef} \lambda_{bceef} + \left( 16\text{Tr}[Y_a Y_c Y_b Y_d] + 14\text{Tr}[Y_a^* Y_b Y_c Y_d^*] - \frac{1}{2} \lambda_{bcda} \text{Tr}[Y_a^* Y_c] 
+ 3\text{Tr}[Y_a Y_c^*] (\text{Tr}[Y_b Y_c^*] + \text{Tr}[Y_b^* Y_c]) - \frac{1}{2} \lambda_{acda} \text{Tr}[Y_b Y_c^*] + c.c. \right) - 2\lambda_{abcd} \text{Tr}[Y_c Y_d^*] = 0
\]

Now let us calculate the two loop results for comparison. At two loops there will be both \( \Lambda^2 \ln \frac{\Lambda^2}{\mu^2} \) and \( \Lambda^2 \) divergences. Since these divergences are always independent of the momenta, we may set the external momenta to zero to simplify the computations. Let us first consider the most divergent terms, the \( \Lambda^2 \ln \frac{\Lambda^2}{\mu^2} \). Only graphs which are one loop graphs with some internal line corrected by either another one loop graph or a renormalization will contribute to the leading divergence. The first such graph is shown in figure 5a. It evaluates to

\[
- \frac{i \lambda_{abcd} \lambda_{cdce}}{4} \int \frac{d^4 p}{(2\pi)^8} \int \frac{d \tau_1 d \tau_2}{\Lambda^4} \int_1^\infty \frac{d \tau_3}{\Lambda^2} \exp \left[ -\frac{1}{\Lambda^2} \left( \tau_1 p^2 + \tau_3 q^2 + \tau_1 m_c^2 + \tau_2 m_d^2 + \tau_3 m_c^2 \right) \right]
\]
\[
\frac{i}{(4\pi)^4} \lambda_{abcd} \Lambda^2 \ln \frac{\Lambda^2}{\mu^2} + O(\Lambda^2)
\] (16)

The graph of figure 5b contributes
\[
\frac{i}{2} \lambda_{abcd} \text{Tr}[Y_c^* Y_d] \int \frac{d^4 p \, d^4 q}{(2\pi)^8} \int \frac{d\tau_1 \ldots d\tau_4}{\Lambda^8} \text{Tr}[q \cdot \sigma(p + q) \cdot \sigma] \\
\times \exp \left[ -\frac{1}{\Lambda^2} \left( \tau_1 p^2 + \tau_2 (p + q)^2 + \tau_3 q^2 + \tau_4 m^2_c + \tau_5 m^2_d \right) \right]
\] (17)

where there are two regions of parameter integrations:
1) \((\tau_1, \tau_2) = R^2_1, (\tau_3, \tau_4) = R^2_1\)

and
2) \(\tau_1 = (0, 1), \quad \tau_2 = (0, 1), \quad \tau_3 = (1, \infty), \quad \tau_4 = (1, \infty)\)

The second region does not give any divergence at this order; the first does give an amount
\[
\frac{i}{(4\pi)^4} \cdot \frac{3}{2} \lambda_{abcd} \text{Tr}[Y_c^* Y_d] \Lambda^2 \ln \frac{\Lambda^2}{\mu^2}
\] (18)

The graph of figure 5c appears to also give a \(\Lambda^2 \ln \frac{\Lambda^2}{\mu^2}\) contribution, but is in fact only quadratically divergent. All that remains are the graphs of figure 6 a-d, which are one loop graphs with renormalization insertions. Figure 6a comes from the quadratic vertex correction. It evaluates to
\[
\frac{i}{(4\pi)^4} \left\{ -\frac{1}{4} (\lambda_{abcd} \lambda_{cdee} + 2 \lambda_{acde} \lambda_{bcde}) + (2 \text{Tr}[Y_a^* Y_b Y_c^* Y_e] + \text{Tr}[Y_a Y_c Y_b Y_e] + \text{c.c.}) \right\} \Lambda^2 \ln \frac{\Lambda^2}{\mu^2}
\] (19)

Figure 6b, from the Yukawa renormalization, gives
\[
\frac{i}{(4\pi)^4} \cdot 3 \ \text{Tr}[Y_a Y_c^* Y_b Y_c^* + \text{c.c.}] \Lambda^2 \ln \frac{\Lambda^2}{\mu^2}
\] (20)

The scalar field strength renormalization insertion graph in figure 6c gives
\[
-\frac{i}{(4\pi)^4} \cdot \frac{1}{2} \lambda_{abcd} \text{Tr}[Y_c^* Y_d^*] \Lambda^2 \ln \frac{\Lambda^2}{\mu^2}
\] (21)
Finally, the fermion field strength renormalization insertion graph in figure 6d gives

\[ i \left( \frac{4\pi}{4}\right)^4 \frac{3}{2} \text{Tr}[Y_a Y_b Y_c Y_c^* + \text{c.c.}] \Lambda^2 \ln \frac{\Lambda^2}{\mu^2} \]  \tag{22} \]

Putting it all together gives the total

\[ i \left( \frac{4\pi}{4}\right)^4 \left( -\frac{1}{2} \lambda_{acde} \lambda_{bcde} + \lambda_{abcc} \text{Tr}[Y_c Y_c] + \left\{ \frac{7}{2} \text{Tr}[Y_a Y_b Y_c Y_c^*] + 4 \text{Tr}[Y_a Y_c^* Y_b Y_c^*] + \text{c.c.} \right\} \right) \Lambda^2 \ln \frac{\Lambda^2}{\mu^2} \]  \tag{23} \]

Comparing (23) with (10) and (15), we see that the results do not agree as they did in the dimensional reduction case. Thus the result in question is seen to fail when nonlocal regularization is used.

The above discussion assumed that no measure factor was used. As stated previously, for the scalar-fermi theory considered here there is no need for a measure factor, but on the other hand neither is there any reason not to include one. In particular, note that if we choose parameters in (1) such that the theory has supersymmetry, it would [11] be necessary to add a measure factor to preserve the supersymmetry in loop amplitudes*. So let us consider adding a minimal measure term

\[ S_m = \frac{1}{(4\pi)^2} \alpha_{ab} \Lambda^2 \int d^4 x \phi_a(x) \phi_b(x) \]  \tag{24} \]

where we shall take

\[ \alpha_{ab} = A \text{ Tr}[Y_a Y_b^* + Y_a^* Y_b] + B \lambda_{abcc} \]  \tag{25} \]

The one loop results (10) and (15) become

\[ (B - 1) \lambda_{abcc} + (A + 3) \text{ Tr}[Y_a Y_b^* + Y_a^* Y_b] = 0 \]  \tag{26} \]

* If supersymmetry auxiliary fields are used, as in the superfield formalism, then no measure factor is required as in this case supersymmetry is linearly realized. When such a theory is nonlocally regulated, the auxiliary fields no longer enter the action purely quadratically. So if one integrates out the auxiliary fields after nonlocalization, extra interactions are generated which are the same as the measure factor which would be needed if one just nonlocalized the theory without the auxiliary fields.
and (after using (26) to eliminate some terms)

\[(B - 1) (\lambda_{aef}\lambda_{cde} + 2\lambda_{acef}\lambda_{bce}) + (-8B - 2A + 14)Tr[Y_aY_b^*Y_c^*Y_c + Y_a^*Y_bY_c^*Y_c] \]

\[+ (-4B + 4A + 8)Tr[Y_a^*Y_c^*Y_c + Y_aY_c^*Y_b^*Y_c] - \frac{(B - 1)^2}{A + 3} \lambda_{abcd}\lambda_{cdee} \]

(27)

The presence of the measure term creates an extra graph at two loops, the one shown in figure 7. This graph gives an extra contribution to the two loop result of

\[
\frac{i}{(4\pi)^4} \left[ \frac{A(B - 1)}{A + 3} - B \right] \lambda_{abcd}\lambda_{cdee} \Lambda^2 \ln \frac{\Lambda^2}{\mu^2} \]

(28)

giving a new two loop total of

\[
\frac{i}{(4\pi)^4} \left[ \left( \frac{(A - \frac{1}{2})(B - 1)}{A + 3} - B \right) \lambda_{abcd}\lambda_{cdee} + \frac{7}{2} Tr[Y_a^*Y_c^*Y_c + Y_a^*Y_cY_c^*] \right] \Lambda^2 \ln \frac{\Lambda^2}{\mu^2} \]

(29)

where we have again used (26). We see that there are no values for A and B which will make (29) equivalent to (27).

There are of course other measure factors that may be added. However, it turns out that any other scalar measure interaction which does not spoil the renormalizability of the theory will affect neither (26) or (29). This is because the measure factor interaction is required to be analytic in $p^2$. Integrals of the form $\int d^4q f(q)$ may only be log divergent if $f(q)$ goes like $q^4$ for large $q$. One could consider adding additional measure terms involving fermions instead of scalars. Of course such terms would not affect (10). They could only affect (23) if they were of order $\Lambda^2$ or higher. But such terms would also ruin the power counting behavior of the theory.

So much for the leading order $\Lambda^2 \ln \frac{\Lambda^2}{\mu^2}$ divergence. The next order $\Lambda^2$ divergence will be totally dependent on $\phi^3$ and $\phi^4$ terms in the measure, so it is almost certain that the result in question could be made to work for these terms with the right choice of measure. Not only that, but one could take (26) to be true only to first order, and add higher order terms to this condition which would then change the condition at second order. None of
this, however, will affect the leading order $\Lambda^2 \ln \frac{\Lambda^2}{\mu^2}$ two loop divergences, which clearly do not satisfy the conjecture of Al-sarhi, Jack, and Jones when nonlocal regularization is used.

3. The Standard Model

In this section we follow the conventions used in ref. 3. In particular, we work in the unbroken phase of the standard model, in which the quadratic divergences occur in the Higgs self-energy, and neglect all Yukawa couplings except that of the top quark, which is denoted by $h$. The $SU(2)$ and $U(1)$ couplings are given by $g$ and $g'/2$ respectively. We will use the minimal measure factor, following the construction outlined in ref. 9; thus the measure factor will not affect the quadratic divergences at one loop. The calculational procedure then follows that shown in the last section, so we will just give the results:

The graph of figure 8 gives $-\frac{i}{(4\pi)^2} \cdot \lambda \Lambda^2$.

The graph of figure 9 gives $-\frac{i}{(4\pi)^2} \cdot \frac{3}{8} g'^2 \Lambda^2$.

The graph of figure 10 gives $-\frac{i}{(4\pi)^2} \cdot \frac{9}{8} g^2 \Lambda^2$.

The graph of figure 11 gives $\frac{i}{(4\pi)^2} \cdot 3h^2 \Lambda^2$.

The graph of figure 12 gives $\frac{i}{(4\pi)^2} \cdot g'^2 \Lambda^2$.

The graph of figure 13 gives $\frac{i}{(4\pi)^2} \cdot 2g^2 \Lambda^2$. This graph and the preceeding one would have been zero in dimensional regularization (or reduction).

Thus the quadratic divergences will cancel for

$$-\lambda + \frac{5}{8} g'^2 + \frac{7}{8} g^2 + 3h^2 = 0$$

(30)

Using the standard model beta functions [3], (30) is independent of $\mu$ only if

$$-12\lambda^2 + (3g'^2 + 9g^2)\lambda + \frac{187}{24}(g'^4 + g^4) - \frac{3}{2} g'^2 g^2 - 12\lambda h^2 + 39h^4 - \frac{17}{2} g'^2 h^2 - \frac{27}{2} g^2 h^2 - 48g^2 h^2$$

(31)
where $g_3$ is the QCD coupling constant. Using (30) to eliminate $\lambda$ in (31), we get

$$\frac{238}{48}g'^4 - \frac{51}{8}g'^2g^2 - \frac{437}{48}g^4 - 105h^4 - 52h^2g'^2 - 60g^2h^2 - 48g'^2h^2 = 0$$

(32)

Then using the relations

$$g^2 = 4m_W^2/v^2, \quad g'^2 = 4(m_Z^2 - m_W^2)/v^2,$$

(33)

and the values [12]

$$m_W = 80.6 \text{ GeV} \quad m_Z = 91.16 \text{ GeV},$$

(34)

we see that (32) has no solution for $h^2 > 0$, regardless of the value of $g_3$. Thus there are no real values for $m_t$ and $m_h$ which make the quadratic divergences cancel using this method. Of course, it may be possible to add a non-minimal measure factor, as in the last section, if one may do so while maintaining gauge invariance. If this can be done, then it is likely that (32) can be altered to get whatever answer one would like for the masses.

4. Conclusions

The purpose of this paper has been to determine how certain results related to the cancellation of quadratic divergences depend on how a theory is regulated. We have studied several results of Al-sarhi, Jack, and Jones from dimensional reduction, and shown that these results are not valid when one instead uses the technique of nonlocal regularization. With hindsight, we may remark that this regularization dependence should not have been unexpected. It is true that we should obtain the same answer independently of a regularization and renormalization program when we calculate any meaningful physical quantity. However, the quadratic divergences do not belong to the category of meaningful physical quantities, so there is no reason not to expect vastly different answers for quadratic divergences using different methods.*

* If cancellation of quadratic divergences depends on a regularization scheme, then why is such
If quadratic divergences depend strongly on a regularization procedure, then is there a “correct” procedure which gives a meaningfully correct answer, e.g. which correctly tells us when we need to rely on new physics to keep a parameter small? The answer seems to be that we should regulate in a way which approximates physics at scales just below the “new physics” cutoff scale. Unfortunately we do not know what this physics is like. Most would agree that this physics is probably not described by a dimensionally reduced theory**. It is more reasonable that a nonlocal field theory might be appropriate, but the construction used in this paper is only one of a probable large number of types of nonlocal theories, and even within this construction there is still ambiguity in the measure. So at this point there seems to be little reason to trust the results of either the dimensional reduction method or the nonlocal method.

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** Lower dimensional theories such as “matrix models” have been studied by researchers interested in fundamental physics, however these theories at most represent toy models and are not by themselves intended to describe our universe.
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FIGURE CAPTIONS

Figure 1: Contributions to the scalar two-point function.

Figure 2: Fermion two-point function graph.

Figure 3: Yukawa coupling renormalization graph.

Figure 4: Scalar quartic coupling renormalization graphs.

Figure 5: Contributions to the leading two loop quadratic divergence.

Figure 6: Two loop quadratic divergences from one loop renormalizations.

Figure 7: Two loop contribution due to a measure factor.

Figure 8: Quadratic divergence from Higgs self-interaction.

Figure 9: Quadratic divergence from hypercharge interaction.

Figure 10: Quadratic divergence from SU(2) interaction.

Figure 11: Quadratic divergence from top quark Yukawa coupling.

Figure 12: Quadratic divergence from hypercharge four-point interaction.

Figure 13: Quadratic divergence from SU(2) four-point interaction.