Zero-dimensional field theory

E.N. Argyres\textsuperscript{a}, A.F.W. van Hameren\textsuperscript{a,b,†}, R.H.P. Kleiss\textsuperscript{b,*}, and C.G. Papadopoulos\textsuperscript{a,+}

\textsuperscript{a} Institute of Nuclear Physics, NCSR “Demokritos”, 15310, Athens, Greece
\textsuperscript{b} University of Nijmegen, Nijmegen, The Netherlands

Abstract

A study of zero-dimensional theories, based on exact results, is presented. First, relying on a simple diagrammatic representation of the theory, equations involving the generating function of all connected Green’s functions are constructed. Second, exact solutions of these equations are obtained for several theories. Finally, renormalization is carried out. Based on the anticipated knowledge of the exact solutions the full dependence on the renormalized coupling constant is studied.

\textsuperscript{†} andrevh@sci.kun.nl, \textsuperscript{*} kleiss@sci.kun.nl, \textsuperscript{+} Costas.Papadopoulos@cern.ch.
1 Introduction

In this paper we study several aspects of zero-dimensional quantum field theory. Such
theories may serve as a model (the static ultra-local limit) of more realistic quantum
field theories, and as a useful didactic object in their own right, since zero-dimensional
theories, for which the path integral is actually a simple integral, allow for many explicit
and exact solutions that cannot be obtained in higher dimensions. As recent examples,
we may quote ’t Hooft [1] and Bender et al. [2]. Questions of particular interest here
are the behavior of theories in high orders of perturbation theory (either many loops, or
large number of external legs), and of the relation between the diagrammatic perturbation
expansion and the full solution. The layout of this paper is as follows. We start by
a diagrammatic (re)derivation of equations that govern the set of all connected Green’s
functions of the theory. We show how for a general scalar theory with arbitrary interactions
the Green’s functions may be obtained order by order. We point out how the Schwinger-
Dyson equation, although derivable from purely diagrammatic arguments, in fact describes
a much larger class of solutions. Next, we discuss the representation of these solutions as
path integrals over contours in the complex $\varphi$-plane. Exact solutions for theories with
interactions up to $\varphi^4$ are obtained including explicitly perturbative and non-perturbative
contributions. We show how a classification of the allowed contours in the complex $\varphi$-
plane can immediately determine properties of the non-perturbative character of these
theories. Renormalization, which for these theories is equivalent to imposing restrictions
to diagrams, is also studied. The wave function renormalizations are fully determined and
their dependence on the renormalized coupling constant of the theory is presented and
discussed.
2 Basic equations

In this section we derive equations for an arbitrary zero-dimensional field theory. The derivation is based entirely upon the diagrammatic representation of the theory. A theory is diagrammatically defined by a sequence of vertices that are weighted by the ‘coupling’ constants taken for convenience as $-\lambda_k$, for the $k$-th vertex. In fact, the two-point coupling $\lambda_2 = m^2 \equiv \mu$ can be eliminated by the introduction of the propagator which means that every line of a diagram accounts for a factor $1/\mu$. Moreover a loop in a graph is counted by an additional parameter, $\bar{h}$. A solution of a zero-dimensional theory is determined by a sequence of objects $C_n$, $n = 0, 1, 2, \ldots$ that represent the connected, $n$-point Green’s functions, i.e. the sum of all connected diagrams with $n$ external lines. One can define the generating function of the Green’s functions as

$$\phi(x) = \sum_{n=0}^{\infty} \frac{x^n}{n!} C_{n+1}. \quad (1)$$

We want to write down an equation for $\phi$, and in order to do so, we represent it with a diagram:

$$\phi(x) = \begin{array}{c}
\end{array}.$$ 

Derivatives of $\phi$ with respect to $x$ are represented by extra lines,

$$\phi'(x) = \begin{array}{c}
\end{array}, \quad \phi''(x) = \begin{array}{c}
\end{array},$$

and so on.

2.1 The Schwinger-Dyson equation

Let us consider a theory with a $k$-th vertex. In order to write an equation for $\phi$ we start with the bare vertex and we attach $k-1$ blobs

$$h^0 \frac{\phi^{k-1}}{(k-1)!}.$$

The factor $1/(k-1)!$ is due to the $k-1$ identical blobs. Considering a one loop attachment we similarly have

$$h^1 \frac{\phi'}{2!} \frac{\phi^{k-3}}{(k-3)!}.$$
The factor $1/(k - 2)!$ is again due to the $k - 2$ identical blobs, whereas the $1/2!$ is due to the symmetry factor of blob with two lines. Following this reasoning we can proceed with higher powers of $\hbar$. For instance at the two-loop level we get two terms

$$
\hbar^2 \frac{\phi'^2 \phi^{k-5}}{2!2! (k - 5)!} + \hbar^2 \frac{\phi'' \phi^{k-4}}{3! (k - 4)!}.
$$

Finally the last term, i.e the term with the largest number of loop attachments, will simply read

$$
\hbar^{k-2} \frac{\phi^{(k-2)}}{(k - 1)!}.
$$

The result looks like

$$
\phi = \sum_{k} \frac{\phi^{(k)}}{(k - 1)!}.
$$

The equation reads

$$
x = \mu \phi + \lambda_k \left( \frac{\phi^{k-1}}{(k - 1)!} + \hbar \frac{\phi'}{2!} \frac{\phi^{k-3}}{(k - 3)!} + \hbar^2 \frac{\phi'^2 \phi^{k-5}}{2!2! (k - 5)!} + \hbar^2 \frac{\phi'' \phi^{k-4}}{3! (k - 4)!} + \ldots + \hbar^{k-2} \frac{\phi^{(k-2)}}{(k - 1)!} \right).
$$

For an arbitrary theory a sum over $k$ should be understood. It represents a non-linear differential equation for $\phi$, the Schwinger-Dyson (SD) equation, which has been derived by the direct application of the Feynman rules.

In order to be more specific let us consider a theory with only a 3-point and a 4-point vertex. Following the abovementioned reasoning a diagrammatic equation for $\phi$ looks like

$$
\phi = \sum_{k} \frac{\phi^{(k)}}{(k - 1)!}.
$$

This reads

$$
\phi(x) = \frac{x}{\mu} - \frac{\lambda_3}{2\mu} \left[ \phi(x)^2 + \hbar \phi'(x) \right] - \frac{\lambda_4}{6\mu} \left[ \phi(x)^3 + 3\hbar \phi(x)\phi'(x) + \hbar^2 \phi''(x) \right].
$$

(3)
This equation generates equations for the Green’s functions if the power series of $\phi(x)$ is inserted. The first three are

\[
C_1 = -\frac{1}{6\mu} C_2 (\lambda_4 C_1 + 3\lambda_3) - \frac{\hbar}{2\mu} C_2 (\lambda_4 C_1 + \lambda_3) - \frac{\hbar^2}{6\mu} \lambda_4 C_3 ,
\]

\[
C_2 = -\frac{1}{2\mu} (2\lambda_3 C_1 C_2 - 2 + \lambda_4 C_1^2 C_2) - \frac{\hbar}{2\mu} (\lambda_3 C_3 + \lambda_4 C_1 C_3 + \lambda_4 C_2^2) - \frac{\hbar^2}{6\mu} \lambda_4 C_4 ,
\]

\[
C_3 = -\frac{1}{2\mu} (2\lambda_3 C_1 C_3 + 2\lambda_3 C_2^2 + \lambda_4 C_3 C_2^2 + 2\lambda_4 C_2^2 C_1) - \frac{\hbar}{2\mu} (3\lambda_4 C_2 C_3 + \lambda_3 C_4 + \lambda_4 C_1 C_4) - \frac{\hbar^2}{6\mu} \lambda_4 C_5 .
\]

(4)

The SD equation is invariant under certain redefinition of the parameters involved. It is not difficult to prove that if $\phi(\mu, \lambda_k, \hbar; x)$ is a solution, also $c^3 \phi(c^{\alpha-2\beta} \mu, c^{\alpha-k\beta} \lambda_k, c^{\alpha-n} \hbar; c^{\alpha-n} x)$ is a solution for any $c, \alpha, \beta$. This scaling property is also a consequence of the fact that $\phi(\mu, \lambda_k, \hbar; x)/\sqrt{\hbar/\mu}$ is a dimensionless function of the scaled variables $y = x/\sqrt{\hbar/\mu}$ and $g_k = \lambda_k \hbar^{k/2-1}/\mu^{k/2}$. The scaling property can be expressed with the following equations, derived, for instance, by differentiating with respect to $c$ and taking $c = 1$:

\[
\left( x \frac{\partial}{\partial x} + \mu \frac{\partial}{\partial \mu} + \hbar \frac{\partial}{\partial \hbar} + \lambda_k \frac{\partial}{\partial \lambda_k} \right) \phi = 0 ,
\]

\[
\left( 1 + x \frac{\partial}{\partial x} + 2\mu \frac{\partial}{\partial \mu} + k \lambda_k \frac{\partial}{\partial \lambda_k} \right) \phi = 0 .
\]

These equations are equivalent to the usual topological relations that relate the number of external lines $E$, the number of internal lines $I$, the number of $k$-vertices $V_k$, and the number of loops $L$, appearing in any diagram,

\[
kV_k = E + 2I \quad V_k = I + 1 - L .
\]

A sum over $k$ should be understood in the general case.

### 2.2 Stepping equations

In the diagrammatic construction, one assumes that every Green’s function can be written as a sum of diagrams, consisting of vertices connected by lines (propagators). The power of $1/\mu$ in a diagram is equal to the number of propagators, and hence the operation $-\partial/\partial \mu$ on this diagram corresponds to cutting a single propagator in all possible places in that diagram. There are two possibilities for the result: the chosen propagator may be either a part of a loop, in which case the diagram remains connected when we cut this line, or part of the ‘tree skeleton’, such that cutting it makes the diagram disconnected:

\[
\text{---} = \text{---} + \text{---} \text{---} .
\]
In the first case the cut diagram remains connected but gains two external lines at the 
price of one loop (i.e. on power of $\bar{h}$); in the second place, the cut diagram falls apart into 
two connected diagrams:

$$\frac{\partial}{\partial \mu} \bigcirc = \bigcirc + \bigcirc \bigcirc .$$

Putting in the correct symmetry factors, we can express this procedure by

$$S2: \frac{\partial}{\partial \mu} \phi(x) + \phi(x) \phi'(x) + \frac{\bar{h}}{2} \phi''(x) = 0 ,$$

(6)

where the second term comes from diagrams that fall apart under cutting, and the third one from loops that are cut open. We call Eq.(6) the Step-2 equation (S2) since it describes a procedure in which the number of external legs is increased in steps of 2.

Like SD (Eq.(3)), the Stepping equation S2 implies relations between various $C$’s. The lowest few of these read

$$C_3 = -\frac{2}{\bar{h}} \left( C_1 C_2 + \frac{\partial}{\partial \mu} C_1 \right) ,$$

$$C_4 = -\frac{2}{\bar{h}} \left( C_1 C_3 + C_2^2 + \frac{\partial}{\partial \mu} C_2 \right) ,$$

$$C_5 = -\frac{2}{\bar{h}} \left( C_1 C_4 + 3C_2 C_3 + \frac{\partial}{\partial \mu} C_3 \right) ,$$

$$C_6 = -\frac{2}{\bar{h}} \left( C_1 C_5 + 4C_2 C_4 + C_3^2 + \frac{\partial}{\partial \mu} C_4 \right) ,$$

(7)

and so on. Note that S2 is completely independent of the interaction potential, and therefore perforce contains information independent of that contained in the SD. It follows that there must be solutions to SD that do not obey S2 and that these solutions cannot be represented by Feynman diagrams.

It is possible to combine SD and S2 in the following manner. Taking the first equation in Eq.(4), we express $C_3$ in $C_2$ and $C_1$, and solve for $C_2$:

$$C_2 = \frac{2\lambda_4 \frac{\partial C_1}{\partial \mu} - \frac{1}{\bar{h}} \left( 6\mu C_1 + 3\lambda_3 C_1^2 + \lambda_4 C_1^3 \right)}{3\lambda_3 + \lambda_4 C_1} .$$

Inserting this into the second equation of Eq.(5), we find a differential equation for $C_1$ alone:

$$0 = 4\bar{h}^2 \lambda_4^3 \left( \frac{\partial C_1}{\partial \mu} \right)^2 - 2\bar{h}^2 \lambda_4^3 \frac{\partial^2 C_1}{\partial \mu^2} (\lambda_4 C_1 + 3\lambda_3)$$

$$- \bar{h} \frac{\partial C_1}{\partial \mu} \left[ 9\lambda_3^2 (3\lambda_3 + \lambda_4 C_1) + 6\mu \lambda_4^2 C_1 - 36\mu \lambda_3 \lambda_4 \right]$$

$$- 9\mu \lambda_3 C_1^2 (3\lambda_3 + \lambda_4 C_1) + 3\lambda_4^2 C_1^2 - 54\mu^2 \lambda_3 C_1 - 27\bar{h} \lambda_3^2 .$$

(8)
By inserting the series expansion
\[ C_1 = \sum_{k \geq 1} \alpha_k h^k , \]
we can then successively determine the coefficients:
\[
\begin{align*}
\alpha_1 &= -\frac{1}{2\mu^2} \lambda_3 , \\
\alpha_2 &= -\frac{1}{24\mu^3} (15\lambda_3^3 - 16\mu\lambda_3^2) , \\
\alpha_3 &= -\frac{1}{48\mu^4} (90\lambda_3^5 - 185\mu\lambda_3^4 + 66\mu^2\lambda_3^2\lambda_4^2) , \\
\alpha_4 &= -\frac{1}{1152\mu^5} (9945\lambda_3^7 - 30270\mu\lambda_3^5\lambda_4 + 42480\mu^2\lambda_3^3\lambda_4^2 - 4352\mu^3\lambda_3\lambda_4^3) ,
\end{align*}
\]
and so on.

Whereas S2 is independent of the interaction potential, we can also derive stepping equations by deleting vertices rather than cutting lines. For example, let us depict all possible ways in which a selected \( \varphi^3 \) vertex (denoted by a dot) can occur in a connected graph:

\[
\begin{array}{c}
\text{Deleting this vertex gives us} \\
\end{array}
\]

or, in terms of \( \phi(x) \), the following Step-3 equation (S3):
\[
\text{S3:} \quad \frac{\partial}{\partial \lambda_3} \phi + \frac{1}{6} \hbar^2 \phi''' + \frac{1}{2} \hbar \left[ \phi'' + (\phi')^2 \right] + \frac{1}{2} \phi^2 \phi' = 0 .
\]

A similar treatment holds for \( \varphi^4 \) vertices: the possible ways in which such a vertex can occur is given by

\[
\begin{array}{c}
\end{array}
\]
and the result of deleting is given by
\[- \frac{\partial}{\partial \lambda_4} \rightarrow \oplus \quad + \quad \oplus + \quad \oplus \oplus + \quad \oplus \oplus \oplus .\]

The corresponding Step-4 equation (S4) is
\[
S4: \quad \frac{\partial}{\partial \lambda_4} \phi + \frac{1}{24} h^3 \phi''' + \frac{1}{12} h^2 [2 \phi \phi'' + 5 \phi \phi'''] + \frac{1}{4} h [\phi^2 \phi'' + 2 \phi (\phi')^2] + \frac{1}{6} \phi^3 \phi' = 0 . \tag{10}
\]

### 2.3 The charged scalar field

Up to now we dealt with diagrammatic construction of zero-dimensional field theories involving only one field. As an illustrative extension, we consider a theory with two fields, \( \text{i.e.} \), a complex, or charged, scalar field. The Green’s functions are labeled with two integers, and the generating function has two expansion parameters \( x \) and \( \bar{x} \). Let us introduce the notation
\[
\partial := \frac{\partial}{\partial x} , \quad \bar{\partial} := \frac{\partial}{\partial \bar{x}} ,
\]
then
\[
\phi(x, \bar{x}) = \sum_{n,m=0}^{\infty} \frac{x^n \bar{x}^m}{n! m!} C_{n,m+1} .
\]

To write down the SD equation, we introduce two kind of lines, distinguishable by an arrow. The generating function is represented by
\[
\phi(x, \bar{x}) = \rightarrow \oplus \oplus , \quad \bar{\phi}(x, \bar{x}) = \rightarrow \oplus .
\]

An incoming external line represents a \( \bar{\partial} \), and an outgoing line represents a \( \partial \). Notice that
\[
h \bar{\partial} \phi = h \bar{\partial} \bar{\phi} = \rightarrow \oplus .
\]

We also introduce a four point vertex with two incoming and two outgoing lines, so that the SD equation we want \( \phi \) to satisfy is given by
\[
\rightarrow \oplus = x \rightarrow + \rightarrow \oplus + \rightarrow \oplus + \rightarrow \oplus ,
\]
or
\[
\phi = \frac{x}{\mu} - \frac{\lambda}{2\mu} \phi^2 \bar{\phi} - \frac{\lambda h}{\mu} \phi \partial \phi - \frac{\lambda h}{2\mu} \bar{\phi} \bar{\partial} \phi - \frac{\lambda h^2}{2\mu} \bar{\partial} \bar{\partial} \phi . \tag{11}
\]
Notice that incoming and outgoing lines are not equivalent, which is represented by the symmetry factors.

Also for the charged scalar field, we can write down stepping equations. For the first one, we use that in the diagrammatic interpretation

\[ \text{Diagram} = \text{Diagram} + \text{Diagram} + \text{Diagram}, \]

leading to

\[ \frac{\partial}{\partial \mu} \phi = -\hbar \partial_\lambda \phi - \phi \partial_\phi - \bar{\phi} \partial \phi \quad (12) \]

The diagrammatic derivation of the stepping equation involving the derivative with respect to \( \lambda \), although equally straightforward, is rather cumbersome, leading to many terms which we refrain from listing here.
3 Solutions to the equations

3.1 The integral representation

The SD equation is highly non-linear. Let us consider the general term connected with the coupling $\lambda_k$ in Eq.(2):

$$Q_k = \phi^{k-1} + \frac{\phi'}{2! (k-3)!} + \frac{\phi''}{3! (k-4)!} + \sum_{k=5}^{\infty} \frac{\phi^{k-2}}{(k-1)!} \cdot$$

It can obviously be organized such that it can be written as

$$Q_k = \sum_{m=1}^{k-1} \sum_{\{\bar{a}_{k-1}; m\}} \frac{\hbar^{k-m-1}}{(1!)^{a_1} (2!)^{a_2} a_2! \cdots (k-1)!^{a_{k-1} a_{k-1}!}} (\phi)^{a_1} (\phi')^{a_2} \cdots (\phi^{(k-2)})^{a_{k-1}}, \quad (13)$$

where $\sum_{\{\bar{a}_{k-1}; m\}}$ stands for the summation with $a_1, a_2, \ldots, a_{k-1}$ running over all positive integers under the restrictions that

$$a_1 + 2a_2 + 3a_3 + \ldots + (k-1)a_{k-1} = k-1 \quad \text{and} \quad a_1 + a_2 + \ldots + a_{k-1} = m.$$  

This sum can be interpreted following the time-honored formula of Faà di Bruno [3]:

$$\frac{d^n}{dx^n} f(g(x)) = \sum_{m=0}^{n} f^{(m)}(g(x)) \sum_{\{\bar{a}; m\}} (n; a_1, \ldots, a_n) \{g'(x)\}^{a_1} \{g''(x)\}^{a_2} \cdots \{g^{(n)}(x)\}^{a_n}$$

where

$$(n; a_1, \ldots, a_n) = \frac{n!}{(1!)^{a_1} (2!)^{a_2} a_2! \cdots (n!)^{a_n} a_n!}.$$  

The identifications $g'(x) = \phi(x)$ and $f^{(m)}(g(x)) = \hbar^{-m} f(g(x))$ with the solution

$$g(x) = \int dx \phi(x) \quad \text{and} \quad f(g(x)) = R(x) = \exp \left( \frac{1}{\hbar} \int dx \phi(x) \right), \quad (14)$$

lead to an equation for $R$, which, including all possible vertices, reads

$$\sum_{k=3}^{\infty} \frac{\lambda_k}{(k-1)!} \hbar^{k-1} R^{(k-1)} + \mu \hbar R' - (x - \lambda_1) R(x) = 0.$$  

This is a linear equation, and a solution can be represented by an integral

$$R_G(x) = \int_G d\varphi \exp \left\{ \frac{1}{\hbar} \left[ x \varphi - S(\varphi) \right] \right\}, \quad (16)$$

where

$$S(\varphi) = \lambda_1 \varphi + \frac{1}{2} \mu \varphi^2 + \sum_{k=3}^{\infty} \frac{\lambda_k}{k!} \varphi^k.$$
and where $\Gamma$ is a contour in the complex $\varphi$-plane, such that the difference between the values of the integrand in the end-points is zero. This is the well known path integral representation, with the action $S$.

A remark is in order. Although the SD equations resulted from a purely diagrammatic construction, their solutions, expressed through the path integral representation, include non-perturbative ones that cannot be realized in a weak coupling expansion, as we will see below.

Secondly, we note that Eq. (15) can be used to write the original SD equation for $\phi$ compactly as

$$x = \lambda_1 + \mu \phi + \sum_{k \geq 3} \frac{\lambda_k}{(k-1)!} \left( \hbar \frac{\partial}{\partial x} + \phi \right)^{k-2} \phi \ .$$

(17)

For a general interacting theory, differentiating $R_\Gamma$ with respect to $\lambda_k$ in Eq. (16), the stepping equation in terms of $\phi$ can be rewritten as

$$\frac{\partial \phi}{\partial \lambda_k} = -\frac{1}{k!} \frac{\partial}{\partial x} \left( \phi + \hbar \frac{\partial}{\partial x} \right)^{k-1} \phi$$

(18)

and in case only a $k$-vertex and the tadpole $\lambda_1$ is present, combining with the SD a simpler form is obtained

$$\frac{\partial \phi}{\partial \lambda_k} = -\frac{1}{k \lambda_k} \left( (x - \lambda_1) \phi' + \phi - 2 \mu \phi \phi' - \hbar \mu \phi'' \right) \ .$$

(19)

Moreover for a charged scalar field the stepping equation in terms of $\phi$ can be written in a compact form

$$\frac{\partial \phi}{\partial \lambda} = -\frac{1}{4 \bar{\lambda}} (\bar{\phi} + \hbar \partial)^2 (\phi + \hbar \partial) \phi \ .$$

(20)

Finally, the linear SD equation for $\varphi^3 + \varphi^4$-theory becomes simply

$$\frac{1}{6} \lambda_4 \hbar^3 R'''(x) + \frac{1}{2} \lambda_3 \hbar^2 R''(x) + \mu \hbar R'(x) - x R(x) = 0 \ ,$$

(21)

and we see that $R(x)$ admits 3 linearly independent solutions (2 if $\lambda_4 = 0$). Hence $\phi(x)$ has a 2-parameter family of solutions (a 1-parameter family if $\lambda_4 = 0$). In the sequel we will show how to get exact explicit solutions for a number of scalar theories.

### 3.2 Results for pure $\varphi^3$-theory

In this section we derive results for the pure $\varphi^3$-theory, with action

$$S(\varphi) = \frac{1}{2} \mu \varphi^2 + \frac{1}{6} \lambda \varphi^3 \ .$$

(22)
This theory is interesting because as we will see the solution for the generating function can be expressed directly in terms of known special functions. Defining

\[ y = \frac{x}{\sqrt{\hbar \mu}}, \quad \xi = \frac{\lambda \sqrt{\hbar}}{6 \mu^{3/2}} \]

the SD equation becomes

\[ 3\xi R''(y) + R'(y) - yR(y) = 0 \]  

which admits the following general solution

\[ R(y) = e^{-y/6\xi} [c_1 \text{Ai}(t) + c_2 \text{Bi}(t)] \]

where

\[ t = (3\xi)^{-1/3} \left( \frac{1}{12\xi} + y \right). \]

Ai and Bi are the Airy functions (cf. [3]). The solution for the generating function of connected Green’s functions is given by

\[ \phi(x) = \sqrt{\frac{\hbar}{\mu}} \left( -\frac{1}{6\xi} + 2^{1/2} t_0^{1/4} \frac{\text{Ai}'(t) + K \text{Bi}'(t)}{\text{Ai}(t) + K \text{Bi}(t)} \right) \]

with \( t_0 = t(y = 0) \). The constant \( K \) is not determined by the SD equation: in fact it could have been even a function of \( \xi \).

For solutions that admit a diagrammatic representation extra information can be obtained by combining SD and stepping equations. For instance, the scaling and stepping equations of the previous section result to a \( K \) that is independent of \( \xi \). Moreover by combining SD and S2 an equation involving only \( C_1 \):

\[ 2\mu^2 C_1 + \lambda \mu C_1^2 + h \lambda^2 \frac{\partial}{\partial \mu} C_1 + h\lambda = 0 \]

(26)

can be obtained. The series of substitutions

\[ v = \frac{h \lambda^2}{\mu^3}, \quad C_1 = -\frac{\lambda h}{2 \mu^2} f(v), \quad w = \frac{1}{3v}, \quad f(v) = -\frac{2k'(w)}{vk(w)}, \quad k(w) = w^{1/3}e^{-w} \psi(w) \]

leads to the Bessel equation

\[ w^2 \psi''(w) + w \psi'(w) - \left( w^2 + \frac{1}{9} \right) \psi(w) = 0 \]

The special solution choice \( \psi(w) = K_{1/3}(w) \) gives the following tadpole and its asymptotic expansion:

\[ f(v) = \frac{2}{v} \left( \frac{K_{2/3}(w)}{K_{1/3}(w)} + 1 \right), \]

\[ C_1 \sim -\frac{2\mu}{\lambda} - \frac{h \lambda^2}{2 \mu^2} \left( 1 - \frac{5}{4} v + \frac{15}{4} v^2 - \frac{1105}{64} v^3 + \frac{1695}{16} v^4 + \cdots \right) . \]

(27)
This tadpole, therefore, has a non-perturbative contribution. The more generic choice 
\( \psi = k_1 I_{1/3}(w) + k_2 I_{-1/3}(w) \), with \( k_1 \neq -k_2 \), gives

\[
f(v) = -\frac{2}{v} \left( \frac{k_1 I_{-2/3}(w) + k_2 I_{2/3}(w)}{k_1 I_{1/3}(w) + k_2 I_{-1/3}(w)} - 1 \right),
\]

\[
C_1 \sim -\frac{\hbar \lambda^2}{2\mu^2} \left( 1 + \frac{5}{4} v + \frac{15}{4} v^2 + \frac{1105}{64} v^3 + \frac{1695}{16} v^4 + \cdots \right),
\] (28)

which is the standard perturbative result \[4\]. The coefficients \( k_1 \) and \( k_2 \) drop out for the perturbative expansion: they simply account for non-perturbative contributions that are not computable perturbatively!

A remark is in order here: although all the terms in the perturbative series for \( C_1 \) have strictly the same complex phase and according to the traditional wisdom the series is not Borel summable, the exact result is well defined, indicating that a suitable generalization of the Borel transform will produce the right answer \[5\].

Another interesting aspect is the large \( n \) behavior of the Green's functions, where \( n \) refers to the number of external legs, a problem that is traditionally seen as relevant to the unitarity of the \( S \) matrix \[6\]. This can be traced from the analytical structure of the solution for the generating functions in the complex \( x \)-plane. As is evident from the fact that the solution, Eq. (24), for the generating function of all connected and disconnected graphs is an entire function, the corresponding Green's function \( Z_n \) grows slower than \( n! \); in fact it grows like \( (n!)^{2/3} \). On the other hand the \( C_n \), the connected graphs, exhibit a factorial growth, since their generating function \( \phi(x) \) possesses poles at finite complex values of \( x \).

### 3.3 Results for pure \( \varphi^4 \)-theory

In this section we derive the lowest Green's functions for the pure \( \varphi^4 \) theory, with action

\[
S(\varphi) = \frac{1}{2} \mu \varphi^2 + \frac{1}{24} \lambda \varphi^4 .
\] (29)

Defining

\[
y = \frac{x}{\sqrt{\hbar \mu}}, \quad \xi = \frac{\lambda \sqrt{\hbar}}{24 \mu^2}
\]

we get for the SD equation

\[
4\xi R''(y) + R'(y) - y R(y) = 0 .
\] (30)
There are three solutions, which can be represented as follows:

\[ R_1(y) = \sum_{n=0}^{\infty} \frac{y^{2n}}{n!} (32\xi)^{-n/2} U(n; (8\xi)^{-1/2}) \]

\[ R_2(y) = \sum_{n=0}^{\infty} (-1)^n \frac{y^{2n}}{n!} (32\xi)^{-n/2} \frac{V(n; (8\xi)^{-1/2})}{\Gamma(n + \frac{1}{2})} \]

\[ R_3(y) = \sum_{n=0}^{\infty} \frac{y^{2n+1}}{(2n+1)!} (4\xi)^{-n/2} i^n H_n(i(16\xi)^{-1/2}) \]  \hspace{1cm} (31)

where \(U(\nu; x)\) and \(V(\nu; x)\) are the parabolic cylinder functions, and \(H_n\) is the \(n^{th}\) Hermite polynomial (cf. [3]). The general solution is a linear combination with arbitrary coefficients.

As we can immediately see contrary to what is argued in many standard textbooks the odd Green’s functions do not necessarily vanish.

On the other hand one can study the S2 equation as well. For this we have to distinguish two possible cases: the ‘standard’ one, with \(C_1\) and the higher odd Green’s functions vanishing, and the case where \(C_1 \neq 0\).

Let us consider the first case with a zero tadpole. In this case we cannot, of course, directly use the results derived above, since these deal with \(C_1\). The S2 becomes somewhat simpler, and in particular

\[ C_4 = -\frac{2}{\hbar} \left( C_2^2 + \frac{\partial}{\partial \mu} C_2 \right) . \]

On dimensional grounds we see that we can write

\[ C_2 = \frac{1}{\mu} \beta(v) \hspace{1cm} v = \frac{\lambda \hbar}{\mu^2} , \]

where \(v\) is dimensionless. Inserting all this into the first nonzero term (that with \(x^1\)) in SD, we find the following equation for \(\beta\):

\[ 4v^2\beta'(v) + v\beta(v)^2 + (2v + 6)\beta(v) - 6 = 0 . \]

The substitutions

\[ \beta(v) = 4v \frac{g'(v)}{g(v)} , \hspace{1cm} g(v) = v^{-1/4} e^w \psi(w) , \hspace{1cm} w = \frac{3}{4v} , \]  \hspace{1cm} (32)

lead then to

\[ w^2 \psi''(w) + w\psi'(w) - \left( w^2 + \frac{1}{16} \right) \psi(w) = 0 . \]  \hspace{1cm} (33)

which has the modified Bessel functions for its solutions. The general solution can always be written as

\[ \psi(w) = k_1 I_{1/4}(w) + k_2 I_{-1/4}(w) . \]
It is instructive to consider the perturbative form of these results, that is, the limit where \( \hbar \) becomes infinitesimally small, or \( w \) goes to infinity. Since \( I_{1/4} \) and \( I_{-1/4} \) have the same asymptotic expansion, a generic choice of \( k_{1,2} \) will lead to a single perturbative expansion. The single exception is the choice \( k_{1} = -k_{2} \) which leads to \( \psi(w) \propto K_{1/4}(w) \), with asymptotic expansion

\[
\beta(v) = \frac{3}{v} \left( \frac{K_{3/4}(w)}{K_{1/4}(w)} - 1 \right) \sim 1 - \frac{1}{2}v + \frac{2}{3}v^2 - \frac{11}{8}v^3 + \frac{34}{9}v^4 + \cdots .
\]

This is the standard perturbative expansion, in which the propagator starts with \( \frac{1}{\mu} \), and has loop corrections in powers of \( \hbar \):

\[
C_{2} = C_{2}^{(1)} = \frac{1}{\mu} \left( 1 - \frac{\lambda \hbar}{2\mu^2} + \frac{2\lambda^2\hbar^2}{3\mu^4} + \cdots \right) . \tag{34}
\]

The alternating signs are of course due to the fact that the Feynman rules prescribe a factor \(-\lambda\) for each vertex in our Euclidean model. The asymptotic expansion in all other cases is equal to that for the choice \( k_2 = 0 \), for which we find

\[
\beta(v) = -\frac{3}{v} \left( \frac{I_{-3/4}(w)}{I_{1/4}(w)} + 1 \right) \sim -\frac{6}{v} + 1 + \frac{1}{2}v + \frac{2}{3}v^2 + \frac{11}{8}v^3 + \frac{34}{9}v^4 + \cdots ,
\]

which gives a nonstandard expansion:

\[
C_{2} = C_{2}^{(2)} = -\frac{6\mu}{\lambda \hbar} + \frac{1}{\mu} \left( 1 + \frac{\lambda \hbar}{2\mu^2} + \frac{2\lambda^2\hbar^2}{3\mu^4} + \cdots \right) . \tag{35}
\]

Note the occurrence of a ‘non-perturbative’ term \( 1/\lambda \) here: the rest of the expansion has an apparent opposite sign of the coupling constant. An other way to look at this solution is by examining the saddle point equation, \( \delta S/\delta \phi = 0 \): the abovementioned solution corresponds to the saddle point \( \phi_c = \sqrt{-6\mu/\lambda} + \mathcal{O}(x) \).

In the case \( C_1 \neq 0 \) we can write, again on dimensional grounds,

\[ C_1 = \alpha(v) \sqrt{\frac{\mu}{\lambda}} , \quad C_2 = \frac{1}{\mu} \beta(v) , \]

with \( v \) as before. The first term (with \( x^0 \)) in SD now gives us a relation between \( \alpha \) and \( \beta \):

\[
\beta(v) = \frac{1}{\nu \alpha(v)} ((6 - v)\alpha(v) + 4v^2 \alpha'(v)) ,
\]

and then the second term (\( x^1 \)) gives

\[
16v^2 \alpha(v) \alpha''(v) - 32v^3 \alpha'(v)^2 + (32v - 24)\alpha(v)\alpha'(v) - 3\alpha(v)^2 = 0 .
\]

Using \( w \) as before, we may now substitute

\[
\alpha(v) = \frac{e^w \sqrt{v}}{\psi(w)} ,
\]

\[
\beta(v) = \frac{3}{v} \left( \frac{K_{3/4}(w)}{K_{1/4}(w)} - 1 \right) \sim 1 - \frac{1}{2}v + \frac{2}{3}v^2 - \frac{11}{8}v^3 + \frac{34}{9}v^4 + \cdots .
\]

This is the standard perturbative expansion, in which the propagator starts with \( \frac{1}{\mu} \), and has loop corrections in powers of \( \hbar \):

\[
C_{2} = C_{2}^{(1)} = \frac{1}{\mu} \left( 1 - \frac{\lambda \hbar}{2\mu^2} + \frac{2\lambda^2\hbar^2}{3\mu^4} + \cdots \right) . \tag{34}
\]

The alternating signs are of course due to the fact that the Feynman rules prescribe a factor \(-\lambda\) for each vertex in our Euclidean model. The asymptotic expansion in all other cases is equal to that for the choice \( k_2 = 0 \), for which we find

\[
\beta(v) = -\frac{3}{v} \left( \frac{I_{-3/4}(w)}{I_{1/4}(w)} + 1 \right) \sim -\frac{6}{v} + 1 + \frac{1}{2}v + \frac{2}{3}v^2 + \frac{11}{8}v^3 + \frac{34}{9}v^4 + \cdots ,
\]

which gives a nonstandard expansion:

\[
C_{2} = C_{2}^{(2)} = -\frac{6\mu}{\lambda \hbar} + \frac{1}{\mu} \left( 1 + \frac{\lambda \hbar}{2\mu^2} + \frac{2\lambda^2\hbar^2}{3\mu^4} + \cdots \right) . \tag{35}
\]

Note the occurrence of a ‘non-perturbative’ term \( 1/\lambda \) here: the rest of the expansion has an apparent opposite sign of the coupling constant. An other way to look at this solution is by examining the saddle point equation, \( \delta S/\delta \phi = 0 \): the abovementioned solution corresponds to the saddle point \( \phi_c = \sqrt{-6\mu/\lambda} + \mathcal{O}(x) \).

In the case \( C_1 \neq 0 \) we can write, again on dimensional grounds,

\[ C_1 = \alpha(v) \sqrt{\frac{\mu}{\lambda}} , \quad C_2 = \frac{1}{\mu} \beta(v) , \]

with \( v \) as before. The first term (with \( x^0 \)) in SD now gives us a relation between \( \alpha \) and \( \beta \):

\[
\beta(v) = \frac{1}{\nu \alpha(v)} ((6 - v)\alpha(v) + 4v^2 \alpha'(v)) ,
\]

and then the second term (\( x^1 \)) gives

\[
16v^2 \alpha(v) \alpha''(v) - 32v^3 \alpha'(v)^2 + (32v - 24)\alpha(v)\alpha'(v) - 3\alpha(v)^2 = 0 .
\]

Using \( w \) as before, we may now substitute

\[
\alpha(v) = \frac{e^w \sqrt{v}}{\psi(w)} ,
\]
to find that $\psi(w)$ again obeys the Bessel equation, Eq.(33). For the asymptotic expansions, again two distinct choices are possible. First, the choice

$$\psi(w) = \frac{1}{p} K_{1/4}(w)$$

gives

$$\alpha(v) = \frac{p e^w \sqrt{v}}{K_{1/4}(w)} , \quad \beta(v) = \frac{3}{v} \left( \frac{K_{3/4}(w)}{K_{1/4}(w)} - 1 \right) - \frac{p^2 e^{2w}}{K_{1/4}(w)^2} ,$$

and the following asymptotic forms for $C_{1,2}$:

$$C_1 \sim p \sqrt{\frac{3\mu}{2\pi \lambda} e^{2w}} , \quad C_2 \sim C_2^{(1)} + \frac{2p^2 we^{2w}}{\mu \pi} . \quad (36)$$

The alternative choice, for which we may take

$$\psi(w) = \frac{1}{p} I_{1/4}(w) ,$$

leads to

$$\alpha(v) = \frac{p e^w \sqrt{v}}{I_{1/4}(w)} , \quad \beta(v) = \frac{3}{v} \left( \frac{I_{3/4}(w)}{I_{1/4}(w)} + 1 \right) - \frac{p^2 e^{2w}}{I_{1/4}(w)^2} ,$$

and

$$C_1 \sim p \sqrt{\frac{3\pi \mu}{2\lambda}} , \quad C_2 \sim C_2^{(2)} - \frac{2\pi p^2 v}{\mu} \left( 1 - \frac{1}{4} v - \frac{13}{96} v^2 - \frac{73}{384} v^3 - \cdots \right) . \quad (37)$$

In contrast to the zero-tadpole case, there remains an arbitrary parameter in these solutions, $p$: it reflects the presence of the ‘non-perturbative’ tadpole-like contribution and has to be determined by additional requirements.

### 3.4 Results for $\varphi^3+\varphi^4$-theory

For the general zero-dimensional $\varphi^3+\varphi^4$-theory, the action is given by

$$S(\varphi) = \frac{1}{2} \mu \varphi^2 + \frac{1}{6} \lambda_3 \varphi^3 + \frac{1}{24} \lambda_4 \varphi^4 . \quad (38)$$

In the dimensionless variables

$$y = \frac{x}{\sqrt{\mu h}} , \quad g_3 = \frac{\lambda_3}{\mu} \sqrt{\frac{h}{\mu}} , \quad g_4 = \frac{\lambda_4 h}{\mu^2} ,$$

the SD equation becomes

$$\frac{1}{6} g_4 R''(y) + \frac{1}{2} g_3 R'(y) + R'(y) - y R(y) = 0 . \quad (39)$$
To solve this equation, let

$$R(y) = e^{-yg_3/g_4} F(y) .$$

Then $F$ satisfies the equation

$$\frac{1}{6} g_4 F'''(y) + \alpha F'(y) - (y + \beta) F(y) = 0 ,$$

where

$$\alpha = 1 - \frac{2g_3}{2g_4} , \quad \beta = \frac{g_3}{g_4} \left( 1 - \frac{g_3^2}{3g_4} \right) .$$

Finally, changing variables

$$y + \beta = \frac{\eta}{\sqrt{\alpha}} , \quad 4\xi = \frac{g_4}{6\alpha^2} = \frac{g_4}{6} \left( 1 - \frac{g_3^2}{2g_4} \right)^{-2} ,$$

Eq. (40) becomes

$$4\xi F'''(\eta) + F'(\eta) - \eta F(\eta) = 0 .$$

Eq. (41) is exactly Eq. (30) of the pure $\varphi^4$-theory, so that the solutions here are those of (31) with $\xi$ as given above and $y$ replaced by $\eta$.

### 3.5 Results for the charged scalar field

For the complex scalar field, the path integral solution is given by

$$R(x, \bar{x}) = \int d\varphi d\bar{\varphi} \exp \left\{ \frac{1}{\hbar} [x\varphi + \bar{x}\varphi - S(\varphi, \bar{\varphi})] \right\} , \quad S(\varphi, \bar{\varphi}) = \mu \varphi \varphi + \frac{\lambda}{4} (\varphi \varphi)^2 .$$

Due to charge conservation (O(2)-symmetry) one can easily show that $R$ only depends on the modulus $x\bar{x}$, so that it satisfies the following equation

$$\zeta R'''(\zeta) + 2 R''(\zeta) + \alpha^2 [R'(\zeta) - R(\zeta)] = 0 , \quad \zeta = \frac{x\bar{x}}{\mu \hbar} , \quad \alpha = \mu \left( \frac{2}{g \hbar} \right)^{1/2} .$$

The third order equation can be solved by power series expansion in $\zeta$ and two of its solutions are given by

$$R_1(\zeta) = \sum_{n=0}^{\infty} \frac{1}{i(n!)^2} \left( \frac{i\alpha}{\sqrt{2}} \right)^n H_n \left( \frac{i\alpha}{\sqrt{2}} \right) , \quad R_2(\zeta) = \sum_{n=0}^{\infty} \frac{(\zeta \alpha)^n}{n!} U(n + \frac{1}{2}, \alpha) ,$$

where $H_n$ stands for the $n^{th}$ order Hermite polynomial and $U(\nu, x)$ is the parabolic cylinder function. The third solution can be found by standard procedures but the actual result is rather cumbersome and we refrain from giving it explicitly.
Where $R_1$ is a purely non-perturbative solution, $R_2$ has an asymptotic series expansion which leads to the normal perturbation series for the connected Green’s functions: for instance the two point function is given by

$$C_{1,1} = \frac{1}{\mu} - \frac{g \bar{h}}{\mu^3} + \frac{5}{2} \frac{g^2 h^2}{\mu^5} + \mathcal{O}(g^3 h^3) \ .$$

(45) Although the standard integral representation of $R$ is given by (42), the differential equation in the variable $\zeta$ leads to another peculiar single contour integral representation

$$R(\zeta) = \int_{\Gamma} d\psi \exp \left( \zeta \psi - \ln \psi - \frac{\alpha^2}{\psi} + \frac{\alpha^2}{2 \psi^2} \right) \ ,$$

where the contour $\Gamma$ is from infinity to infinity such that the integral is convergent.

### 3.6 Contours in the integral representation

The integral representation of the solutions for the pure $\varphi^3$-theory, with $\psi = \sqrt{\mu/\bar{h}} \varphi$, can be written as

$$R(y; \xi) = K \int_{\Gamma} d\psi \exp \left( -\frac{1}{2} \psi^2 - \xi \psi^3 + y \psi \right) \ ,$$

(46) where $K$ is a constant which can depend on $\xi$. In case the moduli of the endpoints of the contour $\Gamma$ are taken to infinity the standard path-integral representation is recovered. In fact in this case the substitution

$$u = (3\xi)^{-1/3} \psi - \frac{1}{6\xi}$$

leads to

$$R(y; \xi) = K \exp \left( -\frac{1}{108\xi} - \frac{y}{6\xi} \right) \int_{\Gamma} du \exp \left( -\frac{1}{3} u^3 + tu \right) \ .$$

This integral can now be expressed in terms of the Airy functions

$$\int_{\Gamma_j} du \exp \left( -\frac{1}{3} u^3 + tu \right) = 2\pi \omega^j \text{Ai}(t \omega^j) \ ,$$

where $\omega = e^{i2\pi/3}$, $j = 0, 1, 2$ and the contours $\Gamma_j$ are depicted in Fig. II. Note that

$$\sum_j \omega^j \text{Ai}(t \omega^j) = 0 \ .$$

For a pure $\varphi^4$-theory, similar considerations allow us to express the functions $R_j$ defined in Eq.(31), $j = 1, 2, 3$ as follows

$$R_1 = \frac{1}{\sqrt{\pi}} (2\xi)^{1/4} \exp \left( -\frac{1}{32\xi} \right) \int_{-\infty}^{\infty} d\psi \exp \left( -\frac{1}{2} \psi^2 - \xi \psi^4 + y \psi \right)$$

$$R_2 = -\frac{i}{\pi^{3/2}} (2\xi)^{1/4} \exp \left( -\frac{1}{32\xi} \right) \int_{-\infty}^{\infty} d\psi \exp \left( -\frac{1}{2} \psi^2 - \xi \psi^4 + y \psi \right) \ (47)$$
Let us have a closer look at the various possible contours in the case of general $\varphi^3 + \varphi^4$-theory. Let us denote the various objects in the action as complex numbers:

$$\varphi = |\varphi|e^{i\omega}, \quad \lambda_3 = |\lambda_3|e^{in_3}, \quad \lambda_4 = |\lambda_4|e^{in_4}. $$

For simplicity and without loss of generality, we may keep $\mu$ real and positive. The direction in the $\varphi$ plane where the term $\lambda_4 \varphi^4$ goes to positive infinity as $|\varphi| \to \infty$ are given by

$$\omega \in \Omega_4(k), \quad \Omega_4(k) = \left(\pi k \frac{\pi}{8} - \frac{\eta_4}{4}, k \pi + \frac{\pi}{8} - \frac{\eta_4}{4}\right), \quad k = 0, 1, 2, 3. $$

Similarly ‘allowed’ directions for the $\lambda_3 \varphi^3$ term are

$$\omega \in \Omega_3(k), \quad \Omega_3(k) = \left(\pi k \frac{\pi}{6} - \frac{\eta_3}{3}, k \pi + \frac{\pi}{6} - \frac{\eta_3}{3}\right), \quad k = 0, 1, 2. $$

Finally, the $\mu \varphi^2$ term goes to positive infinity for

$$\omega \in \Omega_2(k), \quad \Omega_2(k) = \left(\pi k - \frac{\pi}{4}, k \pi + \frac{\pi}{4}\right), k = 0, 1. $$

By inspection of these endpoints, already statements can be made about the (non)perturbative character of the theory corresponding to a given contour. To illustrate this, let us consider a pure $\varphi^4$-theory, i.e. with $\lambda_3 = 0$. Let the contour start at some $\varphi_1$, chosen at infinity with argument $\omega_1$, and end at some $\varphi_2$, also at infinity in some direction with argument $\omega_2$. These values each have to be in some interval $\Omega_4$: let $\omega_1$ be in $\Omega_4(n_1)$, and $\omega_2$
in $\Omega(n_2)$. We can sufficiently specify the contour by giving $n_1$ and $n_2$ so that for instance the contour $\Gamma_{20}$ for $n_4 = 0$ denotes the standard $\varphi^4$-theory, where we may take the real line for $\varphi = -\infty$ and end at $\varphi = +\infty$ (interchange of the endpoints corresponds to replacing $R$ by $-R$ and hence does not influence $\phi(x)$). In total, there are six contours that give a viable $\varphi^4$-theory: $\Gamma_{01}$, $\Gamma_{12}$, $\Gamma_{23}$, $\Gamma_{30}$, $\Gamma_{02}$ and $\Gamma_{13}$. Note that these are related to each other by phase shifts: in fact,

$$\Gamma_{30} = \Gamma_{12}(\eta_4 \rightarrow \eta_4 + 2\pi), \quad \Gamma_{23} = \Gamma_{12}(\eta_4 \rightarrow \eta_4 + 4\pi),$$

$$\Gamma_{12} = \Gamma_{12}(\eta_4 \rightarrow \eta_4 + 6\pi), \quad \Gamma_{13} = \Gamma_{02}(\eta_4 \rightarrow \eta_4 + 2\pi).$$

Therefore, only $\Gamma_{02}$ and $\Gamma_{01}$, say, give really different theories, all other cases being obtainable by an appropriate shift in $\eta_4$. All contours, as stated, correspond to viable theories as long as $\lambda_4$ is non-vanishing, but when we let $|\lambda_4| \rightarrow 0$ there are two possibilities. It may happen that $\Omega_4(n_1)$ overlaps with one of the $\Omega_2$ segments, and $\Omega_4(n_2)$ with the other $\Omega_2$ segment. In that case, the limiting theory is equal to the free theory, and the limit $|\lambda_4| \rightarrow 0$ is smooth: we may call this the perturbative limit. In the other case the limit is not smooth, and the path integral $R$ will diverge as $|\lambda_4| \rightarrow 0$: we call this the non-perturbative limit. Clearly, the limiting behavior depends on the argument $\eta_4$: for the contour $\Gamma_{02}$ (the ‘standard one’) one has

$$-\frac{3}{2}\pi < \eta_4 < \frac{3}{2}\pi, \quad \frac{5}{2}\pi < \eta_4 < \frac{11}{2}\pi : \text{perturbative},$$

$$\frac{3}{2}\pi < \eta_4 < \frac{5}{2}\pi, \quad \frac{11}{2}\pi < \eta_4 < \frac{13}{2}\pi : \text{non-perturbative},$$

and for the other contour $\Gamma_{01}$:

$$\frac{1}{2}\pi < \eta_4 < \frac{3}{2}\pi, \quad \frac{9}{2}\pi < \eta_4 < \frac{11}{2}\pi : \text{perturbative},$$

$$\frac{3}{2}\pi < \eta_4 < \frac{9}{2}\pi, \quad \frac{11}{2}\pi < \eta_4 < \frac{17}{2}\pi : \text{non-perturbative},$$

Figure 2: The regions in the complex $\varphi$-plane which correspond with $\Omega_4(k)$, $\Omega_3(k)$ and $\Omega_2(k)$ with $\eta_j = -j\alpha_j$. 

\begin{align*}
\Gamma_{30} &= \Gamma_{12}(\eta_4 \rightarrow \eta_4 + 2\pi), \\
\Gamma_{23} &= \Gamma_{12}(\eta_4 \rightarrow \eta_4 + 4\pi), \\
\Gamma_{12} &= \Gamma_{12}(\eta_4 \rightarrow \eta_4 + 6\pi), \\
\Gamma_{13} &= \Gamma_{02}(\eta_4 \rightarrow \eta_4 + 2\pi).
\end{align*}
Figure 3: The shaded areas correspond to combinations of $\eta_4$ and $\eta_3$ (in units of $\pi$) for which the limit $\lambda_4 \to 0$ (with $\lambda_3$ fixed) is perturbative, for contour $\Gamma_{01}$ (left plot) and contour $\Gamma_{02}$ (right).

For the pure $\phi^3$-theory, there are of course three contours, related to each other: $\Gamma_{20} = \Gamma_{01}(\eta_3 \to \eta_3 + 2\pi)$, $\Gamma_{12} = \Gamma_{01}(\eta_3 \to \eta_3 + 4\pi)$. For the limiting theory we find, for contour $\Gamma_{01}$:

\[-\frac{5}{4}\pi < \eta_3 < \frac{1}{4}\pi, \quad \frac{7}{4}\pi < \eta_3 < \frac{13}{4}\pi: \text{perturbative},\]

\[\frac{1}{4}\pi < \eta_3 < \frac{7}{4}\pi, \quad \frac{13}{4}\pi < \eta_3 < \frac{19}{4}\pi: \text{non-perturbative}.\]

In a theory with both $\phi^3$ and $\phi^4$ couplings, things become more interesting. Of course, as long as $\lambda_4$ is nonzero, we are allowed to let $\lambda_3$ go to zero without jeopardizing the perturbativity. On the other hand, we can only let $\lambda_4$ vanish with fixed $\lambda_3$ if the selected $\Omega_4$ and $\Omega_3$ intervals overlap. We give in Fig.3 and Fig.4 the values of $\eta_4$ and $\eta_3$ that correspond to a perturbative $\lambda_4 \to 0$ limit. Finally, we may study the combined limits $\lambda_4 \to 0$ followed by $\lambda_3 \to 0$. The regions of perturbativity are given in Fig.2. Clearly, these are more restricted since we are in this case requiring a common overlap of $\Omega_2$, $\Omega_3$ and $\Omega_4$. 

21
Figure 4: The shaded areas correspond to combinations of $\eta_4$ and $\eta_3$ (in units of $\pi$) for which the limit $\lambda_4 \rightarrow 0$ followed by $\lambda_3 \rightarrow 0$ is perturbative, for contour $\Gamma_{01}$ (left) and $\Gamma_{02}$. 
4 Renormalization

Even though zero-dimensional field theories have no infinities, we may still consider the effects of renormalization, which here take a graph-theoretical significance. Renormalizing the field so that the exact propagator is $1/\mu$, and the coupling constant so that the proper vertices assume their tree-order form, we are counting Green’s functions without self-energy and vertex insertions, that is, we are counting the skeleton diagrams of the theory. We shall restrict ourselves to theories that are known to be perturbatively renormalizable in the usual four-dimensional case.

4.1 Renormalization of pure $\varphi^3$-theory

In this case renormalization proceeds as usual with the introduction of the tadpole ($z_1$), the mass ($z_2$) and the vertex ($z_3$) counter terms, as well as the corresponding renormalization conditions that imply the dependence of these counter terms on the renormalized coupling constant. The renormalized action can be written as ($\mu = 1, \bar{h} = 1$)

$$S = \frac{1}{2} z_2 \varphi^2 + \frac{1}{6} g z_3 \varphi^3 + z_1 \varphi ,$$

(49)

and the SD equation takes the form

$$z_2 \varphi = (x - z_1) - \frac{G_3}{2} (\varphi^2 + \varphi') , \quad G_3 \equiv g z_3 .$$

(50)

Moreover using Eq.(19) we get,

$$3 G_3 \frac{\partial \varphi}{\partial G_3} = z_2 \varphi'' + 2 z_2 \varphi \varphi' - \varphi - (x - z_1) \varphi' ,$$

(51)

whereas Eq.(6) becomes,

$$\frac{\partial \varphi}{\partial z_2} = - \varphi \varphi' - \frac{1}{2} \varphi'' .$$

(52)

The renormalization conditions that have to be applied are

**Condition 4.1.1.** No tadpoles, i.e. $\varphi(x = 0) = 0$ ;

**Condition 4.1.2.** propagator $= \varphi'(0) = 1$ ;

**Condition 4.1.3.** vertex $= \varphi''(0) = -g$.

Application of these conditions to the SD equation and its derivative leads to the equations

$$z_1 = -\frac{1}{2} g z_3 , \quad z_2 = 1 + \frac{1}{2} g^2 z_3 .$$

(53)
So if we know $z_3$ as function of $g$, we know $z_1$ and $z_2$ as function of $g$, and we can consider $\phi$ to be a function of $g$ and $x$ only. Its derivative w.r.t. $g$ can, using $\partial/\partial z_1 = -\partial/\partial x$, be written as

$$\frac{\partial \phi}{\partial g} = -\phi\dot{z}_1 + \frac{\partial \phi}{\partial z_2}\dot{z}_2 + \frac{\partial \phi}{\partial G_3} \dot{G}_3 \tag{54}$$

where a dot denotes differentiation w.r.t. $g$. Because $\phi$ is a function of $x$ and $g$ only, the l.h.s. is zero in $x = 0$ by Condition 4.1.1, and evaluation of the r.h.s. leads to the equation

$$\frac{g}{2} \frac{d z_3}{d g} = \frac{2 z_3 - 2(1 + g^2) z_3^2}{-4 + (4 + g^2) z_3}. \tag{55}$$

It is straightforward to derive that for $g = 0$ the perturbative counter terms read

$$z_1(0) = 0, \quad z_2(0) = 1, \quad z_3(0) = 1.$$

Eq.(55) is an Abel equation of the second kind [7]. The perturbative solution, satisfying the above initial condition, is

$$z_3(g) = 1 - g^2 - \frac{1}{2} g^4 - 4 g^6 - 29 g^8 - \frac{545}{2} g^{10} + \cdots, \tag{56}$$

an expansion previously given by Cvitanović et al. [4].

Having, however, solved the SD equation for $\varphi^3$-theory (Section 3.2), we also have an exact, albeit implicit, solution of Eq.(55), making use of Condition 4.1.2:

$$[c_1 \text{Ai}'(t_0) + c_2 \text{Bi}'(t_0)] \left(\frac{2}{g z_3}\right)^{1/3} - \frac{z_2}{g z_3} [c_1 \text{Ai}(t_0) + c_2 \text{Bi}(t_0)] = 0, \tag{57}$$

where

$$t_0 = \left(\frac{2}{g z_3}\right)^{1/3} \left(\frac{1}{2} g z_3 + \frac{z_2^2}{2 g z_3}\right) = \frac{z_2^2}{(2 g^2 z_3^2)^{2/3}} \left(1 + \frac{g^2 z_3^2}{z_2^2}\right),$$

and Ai and Bi are the two independent solutions of the Airy equation $f''(t) = tf(t)$. The meaning of this equation Eq.(57) is that for a given $g$ and by using Eq.(53) as well as the functional form of Ai and Bi we can determine $z_3$. To show that Eq.(57) is an implicit solution of Eq.(55), let

$$F(g) = \frac{(2 g^2 z_3^2)^{1/3}}{z_2},$$

implying $t_0 F(g)^2 - 1 = (g z_3 / z_2)^2$, and differentiate Eq.(57) with respect to $g$ to get

$$\left(F'(g) - \frac{d t_0}{d g}\right) [c_1 \text{Ai}'(t_0) + c_2 \text{Bi}'(t_0)] + F(g) [c_1 \text{Ai}''(t_0) + c_2 \text{Bi}''(t_0)] \frac{d t_0}{d g} = 0.$$
Using the Airy equation and Eq.(57), we get

\[ F'(g) + \frac{d t_0}{d g} (t_0 F(g)^2 - 1) = 0. \]

Explicitly, this says

\[ 2 z_2 \frac{d(g z_3)}{d g} - 3 g z_3 \frac{d z_2}{d g} - 2 z_3 \frac{d(g z_3)}{d g} = 0. \]

By using (53), one easily sees that the above equation is an equivalent form of Eq.(55).

In Fig.5 we present the results of a numerical calculation of \( z_3 \) for the \( \Gamma_{10} \)-contour as function of \( g \), as described in the Appendix. The left graph shows the real and imaginary part of \( z_3 \) as function of real and positive values of \( g \). Notice that \( z_3(0) = 1 \) as demanded, and that the imaginary part does not stay zero for real \( g \). This is, of course, an artifact of the definition of the path integral over a complex contour, which is the \( \Gamma_{10} \)-contour for.
$\varphi^3$-theory in this case (Fig.2). The right graph combines the real and imaginary part in one curve in the complex $z_3$-plane.

Fig.2 shows what happens if we let $g$ run with real and positive values of $gz_3$, so that the actual coupling constant is real and positive.

### 4.2 Renormalization of pure $\varphi^4$-theory

In the case of $\varphi^4$-theory, the renormalized action is given by ($\mu = \bar{h} = 1$)

$$ S = \frac{1}{2}z_2\varphi^2 + \frac{1}{4!}g z_4 \varphi^4 . $$  \hspace{1cm} (58)

The SD equation becomes

$$ z_2 \phi = x - \frac{G_4}{6}(\phi^3 + 3\phi \phi' + \phi'') , \quad G_4 = gz_4 , $$  \hspace{1cm} (59)

and the stepping equation Eq.(19), leads to

$$ 4G_4 \frac{\partial \phi}{\partial G_4} = z_2 \phi'' + 2z_2 \phi \phi' - \phi - x\phi' , $$  \hspace{1cm} (60)

whereas Eq.(6) assumes the form of Eq.(52). The renormalization conditions require that

**Condition 4.2.1.** $\phi(x = 0) = 0$;

**Condition 4.2.2.** $\phi'(0) = 1$;

**Condition 4.2.3.** $\phi''(0) = -g$;

and application to the SD equation leads to the relation

$$ z_2 = 1 - \frac{1}{6}(3 - g)gz_4 . $$  \hspace{1cm} (61)

As in the case of $\varphi^3$-theory, $\phi$ can be considered to be a function of $x$ and $g$ only, and its derivative w.r.t. $g$ can be written as

$$ \frac{\partial \phi}{\partial g} = \frac{\partial \phi}{\partial z_2} z_2 + \frac{\partial \phi}{\partial G_4} G_4 . $$  \hspace{1cm} (62)

The l.h.s. is zero in $x = 0$ by Condition 12.1, and evaluation of the r.h.s. leads to

$$ \frac{dz_4}{dg} = \frac{-6z_4 + (6 - 9g + 3g^2)z_4^2}{6g - g(6 - 5g + g^2)z_4} , $$  \hspace{1cm} (63)

another Abel equation of the second kind. The perturbative solution is given by

$$ z_4(g) = 1 + \frac{3}{2}g + \frac{3}{4}g^2 + \frac{11}{8}g^3 - \frac{45}{16}g^4 + \frac{499}{32}g^5 + \cdots . $$  \hspace{1cm} (64)
We want to remark at this point that the statement by Cvitanović et al. that, in the case of \( \varphi^3 \)-theory, the coefficients of the series expansion of \( g - g z_3 \) count connected three-point diagrams with no self-energy or vertex insertions cannot be carried forward to \( \varphi^4 \)-theory: the coefficients of the series expansion of \( g - g z_4 \) do not count connected four-point diagrams with no self-energy or (four-point) vertex insertions. There are, for example, no such diagrams with three vertices.

To find the exact implicit solution of Eq.(63), we apply Condition 4.2.2 to the solution (31) of \( \varphi^4 \)-theory, resulting in

\[
\frac{c_1 t U(1; t)}{z_2} - \frac{c_2 t V(1; t)}{z_2 \Gamma(\frac{3}{2})} = c_1 U(0; t) + \frac{c_2 V(0; t)}{\Gamma(\frac{1}{2})}, \quad t = \left( \frac{3 z_2^2}{g z_4} \right)^{1/2}.
\]

Letting

\[
F_1(t) = c_1 U(1; t) - \frac{c_2 V(1; t)}{\Gamma(\frac{3}{2})} \quad \text{and} \quad F_0(t) = c_1 U(0; t) + \frac{c_2 V(0; t)}{\Gamma(\frac{1}{2})},
\]

the above equation becomes

\[
t F_1(t) = z_2 F_0(t).
\]

Using the properties of the parabolic functions we can easily show that

\[
F_1'(t) = \left( \frac{z_2}{2} - 1 \right) F_0(t) \quad \text{and} \quad F_0'(t) = -\frac{1}{2} \left( t + \frac{z_2}{t} \right) F_0(t),
\]

so that differentiation of Eq.(12.3) leads to

\[
\frac{dt}{dg} \frac{z_2}{t} F_0(t) + t \frac{dt}{dg} \left( \frac{z_2}{2} - 1 \right) F_0(t) = \frac{dz_2}{dg} F_0(t) - \frac{z_2 dt}{2 dg} \left( t + \frac{z_2}{t} \right) F_0(t).
\]

This equation can be written as

\[
\frac{d}{dg} \left( \frac{t}{z_2} \right) + \frac{dt}{dg} \left( \frac{1}{2} + \frac{t^2}{z_2^2} \frac{g - 3}{6 g z_4} \right) = 0,
\]

where we used relation (61). Finally, since \( t/z_2 = \sqrt{3/g z_4} \) by definition of \( t \), it is easily seen that the above equation becomes Eq.(63).

In Fig.7 we show the results of the numerical calculation of \( z_4(g) \), as described in the Appendix. We used the \( \Gamma_{20} \)-contour for \( \varphi^4 \)-theory (Fig.2). Starting at \( z_4(0) = 1 \), \( z_4(g) \) stays real and positive for real and positive values of \( g \), as expected. Moreover \( z_2 \) exhibits a zero, whose position \( g_* \) can be calculated analytically and is given by

\[
g_* = 3 - \frac{1}{4} \left( \frac{\Gamma(\frac{1}{4})}{\Gamma(\frac{3}{4})} \right)^2 \sim 0.81155.
\]

27
Figure 7: $z_4$ as function of $g$ (left) and $z_2$ as function of $g$ (right).

Figure 8: $C_6$ as function of $g$ (left) and $\log z_4$ as function of $-\log(2 - g)$ (right).
At this point the theory becomes ‘massless’, in the sense that the bare mass becomes zero, yet the Green’s functions do not exhibit singular behavior. In fact let us consider the 6-point function as an example. It can be explicitly calculated and it reads

\[ C_6 = 6z_4^{-1} - 6 + 9g + g^2. \]

It is easy to see that the expansion around \( g = 0 \) reproduces the known perturbative series. Moreover, the left graph Fig.8 presents \( C_6 \) as a function of \( g \).

We also see that \( z_4 \) increases with increasing \( g \), and explodes if \( g \) approaches 2. The graph on the right of Fig.8 suggests that there is a simple pole at \( g = 2 \). In fact, substitution of a Laurent series around \( g = 2 \) in Eq.(63) results in a solution with a simple pole:

\[ z_4(g) = -\frac{6}{g - 2} + 6 - 12(g - 2) + 54(g - 2)^2 - 399(g - 2)^3 + 3948(g - 2)^4 - \cdots. \]  

One can ask the question whether this series expansion corresponds to a solution with \( z_4(0) = 1 \), that is, the perturbative solution. In order to get the perturbative solution from the implicit solution (63), in combination with Eq.(61), we should take the constants \( c_1 \) and \( c_2 \) such that the limit of \( t \to \infty \) exists. Using the properties of the parabolic functions and their asymptotic expansions, we find that the perturbative solution has to satisfy

\[ z_2 = t^2 \left( \frac{B_3/4 \left( \frac{1}{4} t^2 \right)}{B_1/4 \left( \frac{1}{4} t^2 \right)} - 1 \right), \quad B_\nu \left( \frac{1}{4} t^2 \right) := \begin{cases} \frac{1}{\cos \nu \pi} \left[ I_{-\nu} \left( \frac{1}{4} t^2 \right) + I_{\nu} \left( \frac{1}{4} t^2 \right) \right] & \text{if } \text{Re} t < 0, \\ \frac{\pi}{2 \sin \nu \pi} \left[ I_{-\nu} \left( \frac{1}{4} t^2 \right) - I_{\nu} \left( \frac{1}{4} t^2 \right) \right] & \text{if } \text{Re} t > 0, \end{cases} \]  

(68)

together with Eq.(61). For \( g \) close to, but smaller than, \( g = 2 \) we see that \( z_2 < 0 \), so that \( \text{Re} t < 0 \), and it is easy to see that the solution in this case has a simple pole at \( g = 2 \).
However, the coefficients for large powers in the series expansion seem to behave as $(2n+2)!$, so that the series has radius of convergence equal to zero, and the numerical solution of a curve around $g = 2$ in the complex $g$-plane reveals that there is a branch point (Fig. 9). In any case, for $g \to 2^-$ the bare coupling becomes strong and the bare mass squared large and negative whereas the connected Green’s functions are still perfectly calculable; for instance $C_6'(g = 2) = 16$.

4.3 Renormalization of $\phi^3+\phi^4$-theory

The renormalization of the $\phi^3+\phi^4$-theory is more involved, but straightforward. The action is given by

$$S = \frac{1}{2}z_2\phi^2 + \frac{1}{3!}G_3\phi^3 + \frac{1}{4!}G_4\phi^4 + z_1\phi^* , \quad G_3 = g_3z_3 , \quad G_4 = g_4z_4 ,$$

and the SD equation assumes the form

$$z_2\phi = (x - z_1) - \frac{G_3}{2}(\phi^2 + \phi') - \frac{G_4}{6}(\phi^3 + 3\phi\phi' + \phi'') . \quad (69)$$

The stepping equations read

$$\frac{\partial \phi}{\partial G_3} = -\frac{1}{6}\phi''' - \frac{1}{2}\phi\phi'' - \frac{1}{2}\phi^2\phi' ,$$

$$\frac{\partial \phi}{\partial G_4} = -\frac{1}{24}\phi'''' - \frac{1}{6}\phi\phi''' - \frac{5}{12}\phi'\phi'' - \frac{1}{4}\phi^2\phi'' - \frac{1}{2}\phi\phi''^2 - \frac{1}{6}\phi^3\phi' ,$$

$$\frac{\partial \phi}{\partial z_2} = -\phi\phi' - \frac{1}{2}\phi'' ,$$

and the renormalization conditions are now

Condition 4.3.1. $\phi(x = 0) = 0$;

Condition 4.3.2. $\phi'(0) = 1$;

Condition 4.3.3. $\phi''(0) = -g_3$;

Condition 4.3.4. $\phi'''(0) = 3g_3^2 - g_4$.

Combining these conditions with the SD equation one easily gets

$$z_1 = \frac{1}{2}g_3G_4 - \frac{1}{2}G_3 , \quad z_2 = 1 - \frac{1}{6}(3g_3^2 - g_4 + 3)G_4 + \frac{1}{2}g_3G_3 , \quad (70)$$

so that $\phi$ becomes a function of $g_3, g_4$ and $x$ only, leading to the the four equations:

$$\frac{\partial \phi}{\partial g_i} \bigg|_{x=0} = -\phi'(0) \frac{\partial z_1}{\partial g_i} + \frac{\partial \phi}{\partial z_2} \bigg|_{x=0} \frac{\partial G_3}{\partial g_i} + \frac{\partial \phi}{\partial G_3} \bigg|_{x=0} \frac{\partial G_3}{\partial g_i} + \frac{\partial \phi}{\partial G_4} \bigg|_{x=0} \frac{\partial G_4}{\partial g_i} = 0$$
and

\[
\frac{\partial \phi'}{\partial g_i} \bigg|_{x=0} = -\phi''(0) \frac{\partial \phi_1}{\partial g_i} + \frac{\partial \phi'}{\partial z_2} \bigg|_{x=0} + \frac{\partial \phi'}{\partial G_3} \bigg|_{x=0} + \frac{\partial \phi'}{\partial G_4} \bigg|_{x=0} = 0
\]

with \(i = 3, 4\). The coefficients at \(x = 0\) can be inferred from the stepping equations. This way we have a system of four equations involving the partial derivatives of the functions \(G_3(g_3, g_4)\) and \(G_4(g_3, g_4)\) with respect to \(g_3\) and \(g_4\). Notice that the equations are linear with respect to the four partial derivatives but highly non-linear with respect to the functions \(G_3(g_3, g_4)\) and \(G_4(g_3, g_4)\). They can be solved perturbatively with the result

\[
G_3 = g_3 - g_3^3 h + \frac{3}{2} g_3 g_4 h + 4 g_3^5 h^2 - 6 g_3^4 g_4 h^2 + \frac{3}{4} g_3 g_4^2 h^2 - 4 g_3^3 h^3 - \frac{3}{2} g_3^2 g_4^3 h^3
\]

\[
+ \frac{19}{4} g_3^2 g_4^2 h^3 + \frac{11}{8} g_3 g_4^3 h^3 + 7 g_3^6 h^4 - \frac{93}{2} g_3^5 g_4 h^4 + 81 g_3^4 g_4^2 h^4 - \frac{100}{3} g_3^3 g_4^3 h^4
\]

\[
- \frac{45}{16} g_3^4 g_4 h^4 + 47 g_3^4 h^5 - \frac{807}{2} g_3 g_4 h^5 + 927 g_3^2 g_4^2 h^5 - \frac{2787}{4} g_3^3 g_4^3 h^5
\]

\[
+ \frac{1785}{16} g_3^4 g_4 h^5 + \frac{499}{32} g_3^4 h^5
\]

\[
G_4 = g_4 + 3 g_3^4 h - 6 g_3^2 g_4 h + \frac{3}{2} g_4^2 h - 6 g_3^6 h^2 + 5 g_3^5 g_4 h^2 + \frac{3}{2} g_3 g_4 h^2 + \frac{3}{4} g_4 h^2
\]

\[
+ 9 g_3^8 h^3 - 43 g_3^6 g_4 h^3 + \frac{151}{2} g_3^4 g_4^2 h^3 - 39 g_3^5 g_4^3 h^3 + \frac{11}{8} g_3 g_4^4 h^3 + 33 g_3^3 g_4 h^4
\]

\[
- 324 g_3^8 g_4 h^4 - 834 g_3^6 g_4^2 h^4 - \frac{1485}{2} g_3^5 g_4^3 h^4 - \frac{1585}{8} g_3^4 g_4^4 h^4 - \frac{45}{16} g_3 g_4^5 h^4
\]

\[
+ \frac{1029}{2} g_3^4 g_4^2 h^5 - 4610 g_3^4 g_4^3 h^5 + \frac{2787}{2} g_3^3 g_4^2 h^5 - 17020 g_3^3 g_4^3 h^5 + \frac{68595}{8} g_3^3 g_4^4 h^5
\]

\[
- \frac{10705}{8} g_3^3 g_4^5 h^5 + \frac{499}{32} g_3^3 g_4^5 h^5
\]

where the \(h\) dependence has been restored for convenience.

In the limit \(g_3 \to 0\), also \(G_3 \to 0\), and the equations reduce to

\[
\frac{\partial G_4}{\partial g_3} = 0, \quad \frac{\partial G_4}{\partial g_4} = \frac{2(2 - g_4)G_4^2}{6g_4 - (6 - 5g_4 + g_4^2)G_4},
\]

\[
\frac{\partial G_3}{\partial g_3} = \frac{G_4}{g_4}, \quad \frac{\partial G_3}{\partial g_4} = 0.
\]

Note that \(G_4(0, g_4)\) can be identified as \(g_4 z_4(g_4)\) where \(z_4\) is the same function as in pure \(\varphi^4\)-theory. Another interesting result is that the term linear in \(g_3\) in the expansion of \(G_3\) is given by

\[
G_3(g_3, g_4) = g_3 z_4(g_4) + \mathcal{O}(g_3^2).
\]
4.4 Renormalization of the charged scalar field

In the case of the charged scalar field we consider the integral representation

\[ R(x, \bar{x}) = \sqrt{\frac{\mu}{\bar{\mu}}} \int d\varphi d\bar{\varphi} \exp \left\{ -\frac{1}{\hbar} \left[ \mu z_2 \varphi \bar{\varphi} + \frac{1}{4} \lambda z_4 (\varphi \bar{\varphi})^2 - x \varphi - \bar{x} \bar{\varphi} \right] \right\}, \quad (71) \]

which satisfies the SD equation

\[ \zeta R'''(\zeta) + 2R''(\zeta) + \frac{2}{gz_4} [z_2 R'(\zeta) - R(\zeta)] = 0, \quad \zeta = \frac{x \bar{x}}{\mu \bar{\mu}}, \quad g = \frac{\lambda \bar{\mu}}{\mu^2}. \quad (72) \]

In the dimensionless variables

\[ u = \frac{x}{\sqrt{\mu \bar{\mu}}}, \quad \bar{u} = \frac{\bar{x}}{\sqrt{\mu \bar{\mu}}}, \quad \zeta = u \bar{u}, \quad \psi = \sqrt{\frac{\mu}{\bar{\mu}}} \varphi, \quad \bar{\psi} = \sqrt{\frac{\mu}{\bar{\mu}}} \bar{\varphi}, \]

Eq.(71) becomes

\[ R(u, \bar{u}) = \int d\psi d\bar{\psi} \exp \left( -z_2 \psi \bar{\psi} - \frac{1}{4} g z_4 (\psi \bar{\psi})^2 + u \bar{\psi} + \bar{u} \psi \right), \]

implying

\[ \frac{\partial R}{\partial g} = -\frac{1}{4} \frac{d(g z_4)}{dg} \frac{\partial^4 R}{\partial u^2 \partial \bar{u}^2} - \frac{dz_2}{dg} \frac{\partial^2 R}{\partial u \partial \bar{u}} \]

or, in terms of the \( \zeta \)-variable

\[ \frac{\partial R}{\partial g} = -\frac{1}{4} \frac{d(g z_4)}{dg} (\zeta^2 R'''' + 4 \zeta R'''' + 2 R'') - \frac{dz_2}{dg} (\zeta R'' + R'). \quad (73) \]

Now, the generating function of the connected Green's functions is given by

\[ \phi(x, \bar{x}) = \bar{\mu} \frac{\partial}{\partial \bar{x}} \ln R(\zeta) = \frac{x R'(\zeta)}{\mu R(\zeta)}, \]

and the renormalization conditions are

**Condition 4.4.1.** \( \frac{\partial \phi}{\partial x}(x = \bar{x} = 0) = \frac{1}{\mu}, \) implying \( R'(0) = R(0) ; \)

**Condition 4.4.2.** \( \frac{\partial^3 \phi}{\partial \bar{x} \partial x^2}(x = \bar{x} = 0) = -\frac{\lambda}{\mu^4}, \) implying \( R''(0) = \left( 1 - \frac{g}{2} \right) R(0). \)

By combining equations Eq.(72) and Eq.(73) and the renormalization conditions we get

\[ z_2 = 1 - \left( 1 - \frac{g}{2} \right) g z_4. \]
and
\[
\frac{dz_4}{dg} = \frac{2z_4 - (3g^2 - 5g + 2)z_4^2}{-2g + g(g^2 - 3g + 2)z_4},
\]
(74)

with perturbative expansion
\[
z_4(g) = 1 + \frac{5}{2}g + \frac{9}{4}g^2 + \frac{49}{8}g^3 - \frac{271}{16}g^4 + \frac{5025}{32}g^5 + \cdots.
\]
(75)

To get an exact implicit solution of Eq.(74), we go back to Eq.(72) and change variables to get
\[
\eta R'''(\eta) + 2R''(\eta) + \alpha^2[R'(\eta) - R(\eta)] = 0, \quad \eta = \frac{\zeta}{z_2}, \quad \alpha = \sqrt{\frac{2}{gz_4}}z_2.
\]
(76)

This equation has exactly the form of (43), and the perturbative solution is given by
\[
R(\eta) = \sum_{n=0}^{\infty} \frac{(\eta \alpha)^n}{n!} U(n + \frac{1}{2}; \alpha).
\]

Condition 1.4.1 implies the implicit exact solution of the form
\[
\alpha U(\frac{3}{2}; \alpha) = z_2 U(\frac{1}{2}; \alpha),
\]
(77)

where, of course, \( z_2 = 1 - (1 - g/2)gz_4 \).

To show that Eq.(74) is indeed an implicit solution of Eq.(74), differentiate Eq.(74) with respect to \( g \):
\[
[U(\frac{3}{2}; \alpha) + \alpha U'(\frac{3}{2}; \alpha)] \frac{d\alpha}{dg} = \frac{dz_2}{dg} U(\frac{1}{2}; \alpha) + z_2 U'(\frac{1}{2}; \alpha) \frac{d\alpha}{dg},
\]

and using parabolic cylinder functions properties together with Eq.(77) to get
\[
\frac{dz_2}{dg} = \left( \frac{z_2}{\alpha} + \alpha z_2 + \frac{z_2^2}{\alpha - \alpha} \right) \frac{d\alpha}{dg},
\]
(78)

with \( \alpha = \sqrt{\frac{2}{gz_4}}z_2 \). It is a straight forward calculation to show that this is indeed Eq.(74).

In the following we present a derivation of the initial condition for \( z_4(g = 0) \). Using the path integral expression of Eq.(71), we find the SD equation
\[
\mu z_2 \hbar \bar{\partial} R + \frac{\lambda z_4}{2} \hbar^2 \bar{\partial}^2 R - xR = 0.
\]

The generating function \( \phi = \hbar \bar{\partial} \ln R \) of the connected diagrams satisfies
\[
\phi(0,0) = 0, \quad (\bar{\partial} \phi)(0,0) = 0, \quad \phi(0,0) = 0, \quad (\bar{\partial} \phi)(0,0) = 0.
\]
as well as the SD equation

$$\phi = \frac{x}{\mu z} - \frac{\lambda z}{2\mu z} (\phi \phi^2 + 2h \phi \partial \phi + \phi \bar{\partial} \phi + h^2 \bar{\partial} \partial \phi) ,$$

with the renormalization conditions 4.4.1 and 4.4.2. These should hold for any value of $\bar{h}$, and for $\bar{h} = 0$, the SD equation becomes

$$\phi_0 = \frac{x}{\mu z(0)} - \frac{\lambda z(0)}{2\mu z(0)} \phi_0 \phi_0^2 ,$$

from which we derive for the perturbative solution that

$$(\partial \phi_0)(0, 0) = \frac{1}{\mu z(0)} \Rightarrow z_2(0) = 1 ,$$

$$(\partial^2 \phi_0)(0, 0) = -\frac{\lambda z(0)}{\mu^4} \Rightarrow z_4(0) = 1 .$$

Notice that the value $\bar{h} = 0$ is directly related to $g = 0$ since $g$ is proportional to $\bar{h}$. 34
5 Summary

In this paper we studied several aspects of zero-dimensional field theories. In the first place we derived a set of diagrammatic equations, including the well known Schwinger-Dyson equations as well as a set of ‘stepping’ equations generalizing some previous results. Then we showed how to solve these equations exactly in terms of known functions and we established integral representations of these solutions, best known as the ‘path integral’ representation. Explicit results were obtained for $\varphi^3$, $\varphi^4$, $\varphi^3 + \varphi^4$ and the charged scalar field theories. Subsequently, we studied the ‘renormalization’ of such theories in zero dimensions, which is equivalent to counting diagrams with restrictions imposed on the type of diagrams considered, for instance diagrams without any tadpoles, self-energy insertions or vertex insertions. We were able to get explicit results for the dependence of the bare quantities such as the mass, the coupling, and the tadpole counter terms, on the renormalized (physical) coupling constant. Examples of interesting observations are the facts that in $\varphi^4$ theory, the bare mass exhibits a zero at a finite value of the renormalized coupling constant $g = g_*$ (Eq.(66)), whereas at $g \to 2 - \epsilon$ the bare coupling becomes strong and the mass squared becomes large and negative. Yet in both cases the ‘physical’ connected Green’s functions remain finite and calculable.
Appendix

Consider general $\varphi^p$-theory, and suppose that all but one renormalization conditions have been implemented through functions $z_k(g, z_p)$, $k = 1, 2, \ldots, p - 1$ of two variables $z_p$ and $g$, like in Eq. (53) and Eq. (61). This means that we are considering a theory with an action

$$S(g, z_p; \varphi) = \frac{gz_p}{p!} \varphi^p + \sum_{k=1}^{p-1} \frac{z_k(g, z_p)}{k!} \varphi^k.$$  

Let $\Gamma$ be a contour in the complex $\varphi$-plane, such that $\text{Re} \varphi^n \to \infty$ at the endpoints, and define

$$Z_n(g, z_p) := \int_{\Gamma} d\varphi \varphi^n \exp\{-S(g, z_p; \varphi)\}.$$  

Such an integral is not defined for all complex values of $gz_p$. Let $gz_p = |gz_p|e^{i\eta}$ and denote by $e^{-i\eta}\Gamma$ the contour that is obtained from $\Gamma$ by clockwise rotation over $\eta$. For complex values of $gz_p$, we define

$$Z_n(g, z_p) := \int_{e^{-i\eta}\Gamma} d\varphi \varphi^n \exp\{-S(g, z_p; \varphi)\} = e^{-i(n+1)\eta} \int_{\Gamma} d\varphi \varphi^n \exp\left(-\frac{|gz_p|}{p!} \varphi^p - \sum_{k=1}^{p-1} \frac{z_k}{k!} e^{-ik\eta} \varphi^k\right).$$  

Integrals of this type can easily be calculated to high precision by numerical integration. One just has to choose $\Gamma$ such that it goes through one or more saddle points, so that the integrand oscillates as little as possible.

To formulate the renormalization problem further, let us denote the connected moments by $C_n$, so

$$C_1 = \frac{Z_1}{Z_0}, \quad C_2 = \frac{Z_2}{Z_0} - C_1^2, \quad C_3 = \frac{Z_3}{Z_0} - 3C_2C_1 - C_1^3, \ldots$$  

and so on. The problem is to solve $z_p$ as function of $g$ from the implicit function equation

$$C_p(g, z_p) = -g,$$  

which represents the final renormalization condition. This equation can be solved numerically. Given a value of $g$, we have to find the zero of the function

$$F(z_p) := C_p(g, z_p) + g,$$  

which can be found using Newton-Raphson iteration

$$z_p \leftarrow z_p - \frac{F(z_p)}{F'(z_p)}.$$  

36
By making small steps in the value of $g$, the solution $z_p(g)$ on a curve in the complex $g$-plane can be determined. At the start of each iteration, the question arises of which initial value of $z_p$ to choose, and the obvious answer is to choose the final value of the previous iteration, which should lie close to the new final value if the steps in $g$ are not too large.

As a check one can look whether the results obtained with this method satisfy (numerically) the available differential equations for $z_p(g)$ (Eq.(55) and Eq.(56)).
References

[1] G. ’t Hooft, Counting planar diagrams with various restrictions, Nucl.Phys.B538: 389-410 (1999).

[2] C.M. Bender, K.A. Milton and V.M. Savage, Solution of Schwinger-Dyson equations for PT symmetric quantum field theory, preprint hep-th/9907043.

[3] M. Abramowitz and I. Stegun, Handbook of Mathematical Functions, (Dover).

[4] P. Cvitanović, B. Lautrup and R.B. Pearson, Number and weights of Feynman diagrams, Phys.Rev.D18: 1939 (1978).

[5] J. Ellis, E. Gardi, M. Karliner and M. A. Samuel, Phys. Lett. B366 (1996) 268 hep-ph/9509312.

[6] H. Goldberg and M.T. Vaughn, Tree and nontree multiparticle amplitudes, Phys.Rev.Lett.66: 1267-1270 (1991).

[7] Daniel Zwillinger, Handbook of differential equations, (Academic Press).