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Non-Hamiltonian symmetries of a Boussinesq equation

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For a certain class of Hamiltonian systems there exist infinite series of non-Hamiltonian symmetries. Some properties of these series are illustrated using a Boussinesq equation. It is shown that the recursion operators generated by these non-Hamiltonian symmetries are powers of the original recursion operator. A class of recursion formulas for the constants of the motion (not for the corresponding symmetries) is given.

I. INTRODUCTION

For a certain class of Hamiltonian systems there exist so-called recursion operators for symmetries. Repeated application of such a recursion operator yields a series of symmetries. Often it is possible to construct in this way infinite series of Hamiltonian symmetries (corresponding to constants of the motion) and infinite series of non-Hamiltonian symmetries. The most well-known example is the Korteweg-de Vries equation, where the Lénard operator generates an infinite series of Hamiltonian symmetries and an infinite series of non-Hamiltonian symmetries. In this paper we use a Boussinesq equation to illustrate some properties of these series, in particular the series of non-Hamiltonian symmetries. Similar results can be obtained for various other equations, see Ten Eikelder.1,2 In this paper we work within the framework of differential geometry. For definitions of various concepts (symmetry, recursion operator for symmetries, etc.) see, for instance, Ref. 2, where also notations and conventions are given.

II. SYMMETRIES OF A BOUSSINESQ EQUATION

We study a Boussinesq equation of the form

\begin{equation}
\begin{align*}
v_t &= w_x, \\
wh &= v w_x + \lambda v_{xxx},
\end{align*}
\end{equation}

(1)

We consider (1) as an evolution equation in a topological vector space \( \mathcal{W} \) of pairs of smooth functions \((v,w), \) which decay, together with their \( x \) derivatives, sufficiently fast for \( |x| \rightarrow \infty. \) The spaces \( \mathcal{W} \) and \( \mathcal{W}^* \) are constructed such that their duality map \( \langle \cdot, \cdot \rangle \) is the \( L^2 \) inner product. A possible choice is \( \mathcal{W} = \mathcal{J}_p \times \mathcal{J}_p \) and \( \mathcal{W}^* = \mathcal{V}_p \times \mathcal{V}_p, \) where the function spaces \( \mathcal{J}_p \) and \( \mathcal{V}_p \) are described in Ref. 1. In terms of \( u = (v,w) \in \mathcal{W} \) we can write (1) as

\[ \dot{u} = X(u) \left( \frac{d}{dt} u(t) \right). \]

A Hamiltonian form of (1) is well known. Let the function (functional) \( F_0 \) on \( \mathcal{W} \) be given by

\[ F_0 = \int_{-\infty}^{\infty} \left( \frac{1}{6} v^6 - \frac{1}{2} \lambda v_x^2 + \frac{1}{2} w_x^2 \right) dx \]

and let the symplectic form \( \Omega \) on \( \mathcal{W} \) be (represented by the linear mapping \( \Omega: \mathcal{W} \rightarrow \mathcal{W}^* \)) given by

\[ \Omega = \begin{pmatrix} 0 & \partial^{-1} \\ \partial^{-1} & 0 \end{pmatrix}. \]

Then the vector field \( X \) can be written as \( X = \Omega^{-1} dF_0 \) (\( d \) = exterior derivative), so (1) is a Hamiltonian system.

The invariance of (1) for translations along the \( t \) and \( x \) axis and for a scale transformation yields the following elementary symmetries:

\begin{equation}
\begin{align*}
X_0 &= X = \begin{pmatrix} w_x \\ v_x \end{pmatrix}, \\
Y_0 &= \begin{pmatrix} w_x \\ v_x \end{pmatrix}, \quad Z_0 = \begin{pmatrix} 2v + x v_x \\ 3w + x w_x \end{pmatrix} + 2t X_0.
\end{align*}
\end{equation}

A recursion operator for symmetries of (1), written in terms of the "coordinates" of a modified Boussinesq equation, has been given by Fordy and Gibbons.3 In terms of the "original coordinates" \( v \) and \( w \) this operator reads

\[ \Lambda = \Lambda A \Lambda \left( \mathcal{L}_A \right) = \text{Lie derivative in direction of } A. \]

This also can be verified by a straightforward computation.

Let \( A \) and \( B \) be vector fields on \( \mathcal{W} \) such that \( \mathcal{L}_A \Lambda = a \Lambda \) and \( \mathcal{L}_B \Lambda = b \Lambda \) for \( a, b \in \mathbb{R}. \) Define \( A_k = \Lambda^k A \) and \( B_k = \Lambda^k B, \) for \( k = 0,1,2,\ldots. \) Using the fact that the Nijenhuis tensor of \( \Lambda \) vanishes, it is easily shown (see, for instance, Ref. 2151 J. Math. Phys. 27 (8), August 1986 0022-2488/86/082151-03$02.50 © 1986 American Institute of Physics 2151
2) that the Lie bracket \([A_k, B_l] \) is given by
\[
[A_k, B_l] = laB_{k+l} - kbA_{k+l} + \Lambda^{k+l} \{A, B \}. \tag{3}
\]

A simple computation shows that
\[
\mathcal{L}_{X_k} \Lambda = 0, \quad \mathcal{L}_{Y_k} \Lambda = 0, \quad \mathcal{L}_{Z_k} \Lambda = 3\Lambda,
\]
\[
[X_n, Y_o] = 0, \quad [Z_o, X_n] = 2X_o, \quad [Z_o, Y_o] = Y_o. \tag{4}
\]

Substitution in (3) yields the only nonvanishing Lie brackets between the elements of the series \(X_k, Y_k, \) and \(Z_k\) are given by
\[
[Z_k, X_l] = (3l + 2)X_{k+l}, \quad [Z_k, Y_l] = (3l + 1)Y_{k+l}, \quad [Z_k, Z_l] = 3(1 - k)Z_{k+l}. \tag{5}
\]

Since the Nijenhuis tensor of \(\Lambda\) vanishes, it immediately follows that
\[
\mathcal{L}_{X_k} \Lambda = 0, \quad \mathcal{L}_{Y_k} \Lambda = 0, \quad \mathcal{L}_{Z_k} \Lambda = 3\Lambda^{k+1}, \quad k = 0, 1, 2, \ldots. \tag{6}
\]

The first relation corresponds to the well-known fact that \(\Lambda\) is also a recursion operator for symmetries of the "higher-order Boussinesq equations" \(\dot{u} = X_k\). The second relation shows that \(\Lambda\) is also a recursion operator for the equations \(\dot{u} = Y_k\).

Next we discuss some properties of the series of symmetries \(Z_k\). For every non-Hamiltonian symmetry \(Z\) a nonvanishing recursion operator for symmetries is given by \(\Omega^{-1} \mathcal{L}_{Z} \Omega\). If \(Z\) is a Hamiltonian symmetry this expression yields \(0\) [because \(\mathcal{L}_{Z} \Omega = d(\Omega Z) = 0\)]. Note that the recursion operators obtained in this way are always the product of a canonical operator \(\Omega^{-1}\) (also called Hamiltonian operator or implicative operator) and a closed operator \(\mathcal{L}_{Z}(\Omega)\) (also called symplectic operator). Most interesting recursion operators have such a factorization, see, for instance, Magri, Fuchssteiner and Fokas, or Gel'fand and Dorfman. In Ref. 2, we computed recursion operators for the massive Thirring model by this method.

The symmetries \(Z_0\) and \(Z_1\) turn out to be non-Hamiltonian. The corresponding recursion operators are found to be
\[
\Omega^{-1} \mathcal{L}_{Z} \Omega = 3I \quad (I = \text{identity mapping}: \mathcal{W} \rightarrow \mathcal{W}), \quad \Omega^{-1} \mathcal{L}_{Z} \Omega = 6\Lambda. \tag{7}
\]

So the recursion operator \(\Lambda\) can be reconstructed from the symmetry \(Z_1\). From (6) and (7) it is easily shown by induction that
\[
\mathcal{L}_{Z}^{k} \Omega = 3^k(k + 1)!\Omega \Lambda^k, \quad k = 0, 1, 2, \ldots. \tag{8}
\]

Since the Lie derivatives and the exterior derivative commute, this relation yields a very simple proof of the well-known fact that all the two-forms \(\Omega \Lambda^k\) are closed. This property implies that
\[
Z_0 = d(\Omega Z_0) = d(\Omega \Lambda^k Z_0) = \mathcal{L}_{Z_0}(\Omega \Lambda^k)
\]
\[
= (\mathcal{L}_{Z_0}(\Omega)) \Lambda^k + \Omega \mathcal{L}_{Z_0}(\Lambda^k) = (3k + 3)\Omega \Lambda^k \neq 0,
\]
\[
k = 0, 1, 2, \ldots. \tag{9}
\]

Thus we have proved that all the symmetries \(Z_k\) are non-Hamiltonian and that the corresponding recursion operators are powers of \(\Lambda\) (up to a multiplicative constant).

Because \(X_0\) is a Hamiltonian symmetry \(\mathcal{L}_{X_0} \Omega = 0\). A simple computation shows that \(Y_0 = \Omega^{-1} dG_0\) with \(G_0 = \int_0^\infty \nu \nu d\nu\), so \(Y_0\) is also a Hamiltonian symmetry. A computation similar to (9) shows that all the symmetries \(X_k, Y_k\) (\(k = 0, 1, 2, \ldots\)) are Hamiltonian vector fields, i.e., there exist two series of constants of the motion \(F_k\) and \(G_k\) such that
\[
X_k = \Omega^{-1} dF_k, \quad Y_k = \Omega^{-1} dG_k, \quad k = 0, 1, 2, \ldots. \tag{10}
\]

The corresponding symmetries commute, so all these constants of the motion are in involution. The existence of the series \(F_k\) is a standard property in this case, see, for instance, Ref. 6. It follows from (10) that
\[
\Omega^{\lambda} \mathcal{X} = dF_k,
\]
which can be considered as "pre-Hamiltonian" forms for \(X = X_0\). The original Hamiltonian form is obtained for \(k = 0\), while formally \(k = -1\) with \(F_{-1} = \int_0^\infty \nu d\nu \), yields the second Hamiltonian form of the Boussinesq equation.

We now give a class of recursion formulas for the constants of the motion \(F_k\) and \(G_k\). The Hamiltonian vector field corresponding to the function \(\mathcal{L}_{Z_k} F_k\) on \(\mathcal{W}\) is
\[
\Omega^{-1} d\mathcal{L}_{Z} F_k = \Omega^{-1} dZ_k F_k
\]
\[
= \mathcal{L}_{Z_k}(\Omega^{-1} dF_k) - (\mathcal{L}_{Z_k} \Omega^{-1}) dF_k
\]
\[
= [Z_0, X_k] + \Omega^{-1}(\mathcal{L}_{Z_0} \Omega^{-1}) dF_k
\]
\[
= (3k + 2)X_{k+1} + (3l + 3)\Lambda X_k
\]
\[
= (3k + 3l + 5)\Omega^{-1} dF_{k+1},
\]
where we used (5) and (9). This yields the recursion formulas
\[
F_{k+1} = \frac{1}{3k + 3l + 5} \mathcal{L}_{Z_k} F_k = \frac{1}{3k + 3l + 5} (dF_k, Z_k). \tag{11}
\]

In a similar way we get
\[
G_{k+1} = \frac{1}{3k + 3l + 4} \mathcal{L}_{Z_k} G_k = \frac{1}{3k + 3l + 4} (dG_k, Z_k). \tag{12}
\]

Note that in these recursion formulas it is not necessary to reconstruct a functional from its derivatives. The part of \(Z\) with "coefficient" \(\nu = 0\) is \(2X_1\) [see (2)], so this term can be omitted in (11) and (12).

The symmetry \(Z_1\) is given by
\[
Z_1 = \left( \begin{array}{c} Z_{1,1} \\ Z_{1,2} \end{array} \right) + 2tX_1,
\]
where
\[
Z_{1,1} = 12w + 2w_x \partial^{-1} v + 2w_x \partial^{-1} w
\]
\[
+ 40w_xw_x + x(4w_xw_x + 8w_xw_x),
\]
\[
Z_{1,2} = 4w + 2w_x \partial^{-1} v + 2w_x \partial^{-1} v + 58w_xw_x
\]
\[
+ 45w_x^2 + 48w_xw_x + 9w_x^2 + 2w_x \partial^{-1} w
\]
\[
+ x(4w_xw_x + 12w_xw_x)
\]
\[
+ 24w_xw_x + 8w_xw_x + 4w_xw_x)
\]
turn out to hold for the other symmetries $Z_k$. For $l = 1$ we obtain from (11) and (12) the recursion formulas

$$F_{k+1} = \frac{1}{3k+8} \int_{-\infty}^{\infty} \left( \frac{\delta F_k}{\delta v} Z_{1,1} + \frac{\delta F_k}{\delta w} Z_{1,2} \right) dx,$$

(13)

$$G_{k+1} = \frac{1}{3k+7} \int_{-\infty}^{\infty} \left( \frac{\delta G_k}{\delta v} Z_{1,1} + \frac{\delta G_k}{\delta w} Z_{1,2} \right) dx.$$

(14)

Starting with $F_0$ and $G_0$ these relations enable us to generate the series $F_k$ and $G_k$. In fact it is also possible to begin with $F_{-1}$ and $G_{-1}$ = $\int_{-\infty}^{\infty} v dx$.

A constant of the motion that depends explicitly on $t$ is $J = \int_{-\infty}^{\infty} (xv + tw) dx$. Constants of the motion of this type always exist if a conserved density (in this case $v$) has a flux that is also conserved, see Broer and Backerra. The Hamiltonian symmetry corresponding to $J$ is formally given by

$$Z_{-1} = \Omega^{-1} J = \begin{pmatrix} 0 \\ 1 \end{pmatrix}.$$

It can be shown that (11) and (12) also hold for $l = -1$ and $k > 0$. This yields the relations

$$F_{k-1} = \frac{1}{3k+2} \mathcal{L} Z_{-1} F_k = \frac{1}{3k+2} \int_{-\infty}^{\infty} \frac{\delta F_k}{\delta w} dx,$$

$$G_{k-1} = \frac{1}{3k+1} \mathcal{L} Z_{-1} G_k = \frac{1}{3k+1} \int_{-\infty}^{\infty} \frac{\delta G_k}{\delta w} dx,\quad k = 0, 1, 2, \ldots .$$

While (13) and (14) allow us to go upwards in the series of constants of the motion, these two relations allow us to go downwards.