On energy and entropy influxes in the Green–Naghdi Type III theory of heat conduction

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The energy-influx/entropy-influx relation in the Green–Naghdi Type III theory of heat conduction is examined within a thermodynamical framework à la Müller–Liu, where that relation is not specified a priori irrespectively of the constitutive class under attention. It is shown that the classical assumption, i.e. that the entropy influx and the energy influx are proportional via the absolute temperature, holds true if heat conduction is, in a sense that is made precise, isotropic. In addition, it is proved that influx proportionality cannot be postulated in general, because a counterexample can be given in the case of transversely isotropic conduction.

1. Introduction

A key assumption in the standard continuum theory of heat propagation is that, for \( \mathbf{q}, \mathbf{h} \) and \( r, s \) the influx vectors and the external sources of, respectively, energy and entropy, the energy inflow \( (\mathbf{q}, r) \) is proportional to the entropy inflow \( (\mathbf{h}, s) \), the proportionality factor being the absolute temperature \( \vartheta \):

\[
(q, r) = \vartheta (h, s), \quad \vartheta > 0. \tag{1.1}
\]

Whether or not this assumption is generally tenable has been the subject of considerable debate (see e.g. the papers by Müller [1] and Liu [2,3]). The issue can be taken up in the framework of one or another continuum theory of heat conduction: Podio-Guidugli [4] uses the framework of the standard theory, which leads
to a parabolic heat equation, Bargmann & Steinmann [5] that of the Type III theory of Green & Naghdi [6], which allows for propagation of heat waves with finite speed (we summarize these two theories in §2). In this paper, we use the latter framework, in a manner that differs from Bargmann and Steinmann’s in a number of points. As in Podio-Guidugli [4] for the standard theory, we prove: in §3, that for thermodynamic consistency (1.1) must hold true when heat conduction is, in a sense that we make precise, isotropic; in §4, that (1.1) may be violated in the transversely isotropic case. In our proofs, we exploit certain simple mathematical tools that are described in appendices A and B.

2. Heat conduction theories

Thermal and deformational phenomena can be considered coupled, as is done, e.g. in thermoelasticity. For our present purposes, it is sufficient to consider the simplest instance, namely, heat propagation in a rigid conductor. Consequently, we need not distinguish between reference and current configurations; we also leave all spatial dependences tacit.

The standard continuum theory of heat conduction is based on the following two laws:

— *(energy balance)*

\[ \dot{\varepsilon} = -\text{div } q + r, \]

(2.1)

where \( \varepsilon \) denotes the volume density of the internal energy and \( \dot{\varepsilon} \) its time rate;

— *(entropy imbalance)*

\[ \dot{\eta} \geq -\text{div } h + s, \]

(2.2)

with \( \eta \) the entropy. It is customarily assumed that (1.1) holds. With this and the notion of Helmholtz free energy:

\[ \psi = \varepsilon - \vartheta \eta, \]

the following *free-energy growth inequality* is arrived at

\[ \dot{\psi} + \eta \dot{\vartheta} + \vartheta^{-1} q \cdot \nabla \vartheta \leq 0. \]

(2.3)

When a set of state variables is chosen—say, \( \{\vartheta, \nabla \vartheta, \dot{\vartheta}\} \)—this inequality is used to derive restrictions on the choice of the state functions delivering \( \psi, \eta \) and \( q \), in the manner devised by Coleman & Noll [7] in their classic paper.

Green & Naghdi [8] modified the classic path, in that they took the following entropy law as their point of departure:

— *(entropy balance)*

\[ \xi = \dot{\eta} + \text{div } h - s, \]

(2.4)

and they chose not to assume that the *internal entropy production* \( \xi \) is non-negative. As to internal energy, they accepted the balance equation (2.1), implicitly excluding any internal production of energy; as to entropy and energy inflows, the classic proportionality expressed by (1.1). Consequently, they arrived at the following version of (2.4) in terms of Helmholtz free energy:

\[ \dot{\psi} + \eta \dot{\vartheta} + h \cdot \nabla \vartheta + \xi = 0. \]

(2.5)

Moreover, in their Type III theory, they assumed that the constitutive mappings delivering \( \psi, \eta \) and \( h \), depend on the following list of state variables:

\[ S := \{\alpha, \dot{\alpha}, \nabla \alpha, \nabla \dot{\alpha}\}, \]

(2.6)
where the thermal displacement $\alpha$, a field whose precise interpretation in statistical mechanics is still wanted, is somehow indirectly characterized by postulating that its time derivative equals an empirical temperature $T$

$$\dot{\alpha}(t) := T(t).$$

As anticipated in §1, Müller and Liu questioned the general applicability of assumption (1.1) and proposed a thermodynamic format within which to test whether or not it fits a given constitutive class. This format amounts to considering the energy balance (2.1), with $r = 0$, as an internal constraint to be appended to the entropy imbalance (2.2), with $s = 0$, after multiplication by the so-called ‘Lagrange multiplier’ $\lambda$, a positive scalar to be constitutively specified

$$\dot{\eta} + \text{div} \ h - \lambda(\dot{\varepsilon} + \text{div} \ q) \geq 0, \quad \lambda > 0. \quad (2.7)$$

When the constitutive mappings involved in (2.7) depend on the Green–Naghdi list (2.6) of state variables, one can try and use a Coleman–Noll type procedure to see whether the multiplier $\lambda$ may be shown to be, if not equal, at least proportional to coldness $1$, that is, the inverse $\vartheta - 1$ of the absolute temperature.

3. Isotropic conduction implies proportionality of energy and entropy influxes

On choosing the state variables as in (2.6), inequality (2.7) reads:

$$(\partial_a \eta - \lambda \partial_a \varepsilon) \dot{\alpha} + (\partial_a \eta - \lambda \partial_a \varepsilon) \dot{\alpha} + (\partial \alpha \eta - \lambda \partial \alpha \varepsilon) \cdot \nabla \dot{\alpha}$$

$$+ (\partial \alpha \eta - \lambda \partial \alpha \varepsilon) \nabla \dot{\alpha} + \text{div} \ h - \lambda \text{div} \ q \geq 0,$$

$$(3.1)$$

where

$$\text{div} \ h - \lambda \text{div} \ q = (\partial_a h - \lambda \partial_a q) \cdot \nabla \alpha + (\partial_a h - \lambda \partial_a q) \cdot \nabla \dot{\alpha}$$

$$+ (\partial \alpha h - \lambda \partial \alpha q) \cdot \nabla^2 \alpha + (\partial \alpha h - \lambda \partial \alpha q) \cdot \nabla^2 \dot{\alpha}. \quad (3.2)$$

We apply the Coleman–Noll procedure [7] to the inequality that is obtained by combining (3.1) and (3.2), that is, we require that it be satisfied whatever the local continuation $[\alpha, \nabla \alpha, \nabla \varepsilon, \nabla^2 \alpha]$ of each admissible process. This requirement is satisfied if and only if

$$\begin{align*}
\text{sym}(\partial \alpha h - \lambda \partial \alpha q) &= 0, \\
\text{sym}(\partial \alpha h - \lambda \partial \alpha q) &= 0,
\end{align*} \quad \text{sym}(\partial \alpha h - \lambda \partial \alpha q) = 0, \quad \partial \alpha \eta - \lambda \partial \alpha \varepsilon = 0$$

$$\begin{align*}
\partial \alpha \eta - \lambda \partial \alpha \varepsilon &= 0, \\
\partial \alpha \eta - \lambda \partial \alpha \varepsilon &= 0,
\end{align*} \quad \text{sym}(\partial \alpha h - \lambda \partial \alpha q) = 0, \quad \partial \alpha \eta - \lambda \partial \alpha \varepsilon = 0, \quad \partial \alpha \eta - \lambda \partial \alpha \varepsilon = 0,$$

and, moreover, inequality (2.7) reduces to

$$(\partial_a \eta - \lambda \partial_a \varepsilon) \dot{\alpha} + (\partial \alpha a \eta - \lambda \partial \alpha \varepsilon) \nabla \dot{\alpha}$$

$$+ (\partial_a h - \lambda \partial_a q) \cdot \nabla \alpha \geq 0.$$

$$(3.4)$$

1 Here is a quotation from Green & Nagdi [8], p. 180: ‘The temperature $T$ (on the macroscopic scale) is generally regarded as representing (on thermoelectric scale) some ‘mean’ velocity magnitude or ‘mean’ (kinetic energy)$^{1/2}$. With this in mind, we introduce a scalar $\alpha = \alpha(X, t)$ through an integral of the form

$$\alpha = \int_{t_0}^{t} T(X, r) \, dr + \alpha_0,$$ 

$$(7.3)$$

where $t_0$ denotes some reference time and the constant $\alpha_0$ is the initial value of $\alpha$ at time $t_0$. In view of the above interpretation associated with $T$ and the physical dimension of the quantity defined by (7.3), the variable $\alpha$ may justifiably be called thermal displacement magnitude or simply thermal displacement. Alternatively, we may regard the scalar $\alpha$ (on the macroscopic scale) as representing a ‘mean’ displacement magnitude on the molecular scale and then $T = \dot{\alpha}$.1
Conditions (3.3)\textsubscript{1,2} are expedient to prove our main result:

For each chosen value of the independent variables \( \alpha, \dot{\alpha} \), let the energy and entropy influxes and the Lagrangian multiplier be delivered by isotropic constitutive mappings \( \hat{q}(\alpha, \dot{\alpha}, \cdot, \cdot), \hat{h}(\alpha, \dot{\alpha}, \cdot) \), \( \hat{\lambda}(\alpha, \dot{\alpha}, \cdot, \cdot) \), the first two vector-valued, the third delivering positive scalars. Then, the energy influx is proportional to the entropy influx via the Lagrangian multiplier, which depends neither on the thermal displacement gradient \( \nabla \alpha \) nor on the temperature gradient \( \nabla \dot{\alpha} \):

\[
\hat{h} = \lambda \, \hat{q}, \quad \lambda = \hat{\lambda}(\alpha, \dot{\alpha}) > 0. \tag{3.5}
\]

Our proof is achieved as follows. For simplicity, we let the dependence on \( \alpha \) and \( \dot{\alpha} \) be tacit, and write \( \mathbf{u} \) for the third state variable in the list (2.6), and \( \mathbf{v} \) for the fourth. On making use of the representation formula (A 3) from appendix A, we set:

\[
\tilde{h}(\mathbf{u}, \mathbf{v}) = \hat{h}_1 \mathbf{u} + \hat{h}_2 \mathbf{v}, \quad h_j = \hat{h}_j(|\mathbf{u}|, |\mathbf{v}|, \mathbf{u} \cdot \mathbf{v}) \quad (j = 1, 2) \tag{3.6}
\]

and

\[
\hat{q}(\mathbf{u}, \mathbf{v}) = q_1 \mathbf{u} + q_2 \mathbf{v}, \quad q_j = \hat{q}_j(|\mathbf{u}|, |\mathbf{v}|, \mathbf{u} \cdot \mathbf{v}) \quad (j = 1, 2),
\]

whence

\[
\mathbf{h} - \lambda \, \mathbf{q} = (h_1 - \lambda q_1) \mathbf{u} + (h_2 - \lambda q_2) \mathbf{v}; \tag{3.7}
\]

provisionally, we also set

\[
\lambda = \tilde{\lambda}(|\mathbf{u}|, |\mathbf{v}|, \mathbf{u} \cdot \mathbf{v}).
\]

With this, condition (3.3)\textsubscript{1} reads:

\[
(h_1 - \lambda q_1) \mathbf{I} + (\partial_1 h_1 - \lambda \partial_1 q_1)|\mathbf{u}|^{-1} \mathbf{u} \otimes \mathbf{u} + (\partial_3 h_2 - \lambda \partial_3 q_2) \mathbf{v} \otimes \mathbf{v} + (\partial_3 h_1 + \partial_1 h_2)|\mathbf{u}|^{-1} - \lambda (\partial_3 q_1 + \partial_1 q_2)|\mathbf{u}|^{-1}) \text{sym}(\mathbf{u} \otimes \mathbf{v} + \mathbf{v} \otimes \mathbf{u}) = 0, \tag{3.8}
\]

and has to be satisfied for all \( \mathbf{u} \) and \( \mathbf{v} \) (here \( \partial_i \) (\( i = 1, 2, 3 \)) denotes differentiation of a function with respect to the \( i \)-th of its arguments). An application of the first of the two algebraic lemmas proved in appendix B permits us to conclude that:

\[
\begin{align*}
\partial_1 (h_1 - \lambda q_1) + (\partial_1 \lambda)q_1 &= 0, \\
\partial_3 h_2 - \lambda \partial_3 q_2 &= 0.
\end{align*} \tag{3.9}
\]

and

\[
\begin{align*}
\partial_3 h_1 - \lambda \partial_3 q_1 + (\partial_1 h_2 - \lambda \partial_1 q_2)|\mathbf{u}|^{-1} &= 0.
\end{align*}
\]

Quite analogously, condition (3.3)\textsubscript{2} yields

\[
\begin{align*}
\partial_2 (h_2 - \lambda q_2) + (\partial_2 \lambda)q_2 &= 0, \\
\partial_3 h_1 - \lambda \partial_3 q_1 &= 0 \tag{3.10}
\end{align*}
\]

and

\[
\begin{align*}
\partial_3 h_2 - \lambda \partial_3 q_2 + (\partial_2 h_2 - \lambda \partial_2 q_2)|\mathbf{v}|^{-1} &= 0.
\end{align*}
\]

Now, (3.9)\textsubscript{1,2} and (3.10)\textsubscript{1,2} imply that the Lagrange multiplier \( \lambda \) cannot depend on \( |\mathbf{u}|, |\mathbf{v}| \); moreover, (3.9)\textsubscript{3} and (3.10)\textsubscript{1} allow to conclude that it does not depend on \( \mathbf{u} \cdot \mathbf{v} \) as well, which establishes the second of (3.5). Consequently, again with the use of (3.9)\textsubscript{1} and (3.10)\textsubscript{1}, we arrive at the following precise statement of (3.5):

\[
\hat{h}(\alpha, \dot{\alpha}, \nabla \alpha, \nabla \dot{\alpha}) = \hat{\lambda}(\alpha, \dot{\alpha}) \, \hat{q}(\alpha, \dot{\alpha}, \nabla \alpha, \nabla \dot{\alpha}). \tag{3.11}
\]

A direct consequence of (3.11) is that, when two material bodies are in ideal thermal contact, that is, by definition, when neither the energy nor the entropy flux suffers a jump at a point of a
common interface oriented by the normal field $n$:

$$\|q \cdot n\| = \|h \cdot n\| = 0, \quad (3.12)$$

then the Lagrange multiplier is also continuous at the interface:

$$\|\lambda\| = 0. \quad (3.13)$$

At this point, in the words of Müller, ‘... there is a price to pay for not having introduced the temperature so far in rational thermodynamics with Lagrange multipliers’ [9, p. 168]. The desideratum is a proof that

$$\lambda \propto \vartheta^{-1}, \quad (3.14)$$

with $\vartheta^{-1}$ the coldness; hence, in particular, that $\hat{\lambda}$ is a universal function. This can be achieved, as Müller [1] himself proposed by considering a situation when a material body of whatsoever constitutive nature is put in ideal thermal contact with an ideal gas, for which the kinetic theory permits to show that (3.14) indeed holds, provided $\vartheta$ is taken proportional to the empirical temperature $T$. We leave it to the reader to decide whether or not such an argument is convincing. For us, in that it relies on importing a result from a theory tacitly regarded as more foundational in nature, it is for sure suggestive of assuming (3.14) right away, without any need to imply, let alone accept, a vassalage between theories.

4. A counterexample to proportionality: transversely isotropic conduction

In this section, we show that assumption (1.1) is not tenable in general. Within the framework of a standard theory of heat conduction modified by the introduction of a Lagrangian multiplier, this result has been announced by Liu [10], who proposed to consider the case of transversely isotropic materials to exhibit a counterexample to proportionality of energy and entropy influxes; the proof he offered is faulty, although easily amendable [4]. We here exploit the same counterexample within the Green–Naghdi theory of Type III.

For $e \in V$ any chosen unit vector, consider the class of vector-valued mappings over $V \times V$ of the following form:

$$\tilde{f}(u, v; e) = \varphi_1 u + \varphi_2 v + \varphi_3 e, \quad \varphi_i = \hat{\varphi}_i(|u|, |v|, u \cdot v, u \cdot e, v \cdot e) \quad (i = 1, 2, 3) \quad (4.1)$$

(once again, we have left the dependence on $a$ and $\dot{a}$ tacit). Clearly, each mapping $\tilde{f}$ in this class is transversely isotropic with respect to the axis $e$, in the sense that it satisfies

$$Q\tilde{f}(u, v) = \tilde{f}(Qu, Qv), \quad \forall u, v \in V, \quad \forall Q \in \text{Orth}(e), \quad (4.2)$$

where $\text{Orth}(e)$ denotes the continuous group of all rotations about the span of $e$.

On using a representation of type (4.1) for both $\tilde{h}$ and $\tilde{q}$:

$$\tilde{h}(u, v) = h_1 u + h_2 v + h_3 e, \quad h_i = \tilde{h}_i(|u|, |v|, u \cdot v, u \cdot e, v \cdot e) \quad (i = 1, 2, 3) \quad (4.3)$$

and

$$\tilde{q}(u, v) = q_1 u + q_2 v + q_3 e, \quad h_i = \tilde{q}_i(|u|, |v|, u \cdot v, u \cdot e, v \cdot e) \quad (i = 1, 2, 3),$$

whence

$$h - \lambda q = (h_1 - \lambda q_1) u + (h_2 - \lambda q_2) v + (h_3 - \lambda q_3) e, \quad (4.4)$$

with the following provisional assumption as to the multiplier $\lambda$:

$$\lambda = \hat{\lambda}(|u|, |v|, u \cdot v, u \cdot e, v \cdot e). \quad (4.5)$$
Now, condition (3.3)\(_1\) becomes:

\[
(h_1 - \lambda q_1)I + (\partial_1 h_1 - \lambda \partial_1 q_1)|u|^{-1} u \otimes u + (\partial_3 h_2 - \lambda \partial_3 q_2) v \otimes v
+ (\partial_4 h_3 - \lambda \partial_4 q_3) e \otimes e
+ (\partial_3 h_1 + \partial_1 h_2)|u|^{-1} - \lambda(\partial_3 q_1 + \partial_1 q_2)|u|^{-1}) \text{sym}(u \otimes v + v \otimes u)
+ (\partial_4 h_1 + \partial_1 h_3)|u|^{-1} - \lambda(\partial_4 q_1 + \partial_1 q_3)|u|^{-1}) \text{sym}(u \otimes e + e \otimes u)
+ (\partial_4 h_2 + \partial_2 h_3 - \lambda(\partial_4 q_2 + \partial_2 q_3)) \text{sym}(v \otimes e + e \otimes v) = 0,
\]

for all \(u\) and \(v\); condition (3.3)\(_2\) yields a completely analogous identity. By applying Lemma 2 in appendix B to each of these two identities, we recover relations (3.9) and (3.10), as well as the following list of additional relations:

\[
\begin{align*}
\partial_4 h_3 - \lambda \partial_4 q_3 &= 0, \\
\partial_4 h_1 - \lambda \partial_4 q_1 + (\partial_1 h_3 - \lambda \partial_1 q_3)|u|^{-1} &= 0, \\
\partial_4 h_2 - \lambda \partial_4 q_2 + \partial_2 h_3 - \lambda \partial_3 q_3 &= 0,
\end{align*}
\]

and

\[
\begin{align*}
\partial_5 h_3 - \lambda \partial_5 q_3 &= 0, \\
\partial_5 h_2 - \lambda \partial_5 q_2 + (\partial_2 h_3 - \lambda \partial_2 q_3)|v|^{-1} &= 0, \\
\partial_5 h_1 - \lambda \partial_5 q_1 + \partial_3 h_3 - \lambda \partial_3 q_3 &= 0.
\end{align*}
\]

Now, we have shown that (3.9) and (3.10) imply that \(\lambda\) cannot depend on the first three arguments in the preliminary representation (4.5); moreover, owing to (3.9)\(_1\) and (3.10)\(_1\), we have the following rearrangements of, respectively, conditions (4.7)\(_2\) and (4.8)\(_2\) and conditions (4.7)\(_3\) and (4.8)\(_3\):

\[
\begin{align*}
(\partial_4 \lambda)q_1 + \partial_1 (h_3 - \lambda q_3)|u|^{-1} &= 0, \\
(\partial_5 \lambda)q_2 + \partial_2 (h_3 - \lambda q_3)|v|^{-1} &= 0,
\end{align*}
\]

and

\[
\begin{align*}
(\partial_4 \lambda)q_2 + \partial_2 h_3 - \lambda \partial_3 q_3 &= 0, \\
(\partial_5 \lambda)q_1 + \partial_3 h_3 - \lambda \partial_3 q_3 &= 0.
\end{align*}
\]

The last two relations imply that

\[
(\partial_4 \lambda)q_2 = (\partial_5 \lambda)q_1,
\]

a condition which poses no restrictions on the choice of \(q\) if and only if \(\lambda\) is independent of the fourth and fifth argument in (4.5). But, if \(\lambda\) may depend only on \(\alpha\) and \(\dot{\alpha}\), conditions (4.10) coalesce and, together with (4.7)\(_1\), (4.8)\(_1\) and (4.9), imply that \((h_3 - \lambda q_3)\) may also depend only on \(\alpha\) and \(\dot{\alpha}\).

All in all, we have that

\[
h_1 - \lambda q_1 = 0, \quad h_2 - \lambda q_2 = 0, \quad h_3 - \lambda q_3 = f \quad \text{with } f = \tilde{f}(\alpha, \dot{\alpha}), \quad \lambda = \tilde{\lambda}(\alpha, \dot{\alpha}).
\]

Accordingly, an influx discrepancy in the direction of the transverse isotropy remains:

\[
h - \lambda q = (h_3 - \lambda q_3)e = f e,
\]

where

\[
\tilde{h}_3(\alpha, \dot{\alpha}, |\nabla \alpha|, |\nabla \dot{\alpha}|, \nabla \alpha \cdot \nabla \dot{\alpha}, \partial_e \alpha, \partial_e \dot{\alpha})
\]

\[
= \tilde{\lambda}(\alpha, \dot{\alpha}) \tilde{q}_3(\alpha, \dot{\alpha}, |\nabla \alpha|, |\nabla \dot{\alpha}|, \nabla \alpha \cdot \nabla \dot{\alpha}, \partial_e \alpha, \partial_e \dot{\alpha}) + \tilde{f}(\alpha, \dot{\alpha}).
\]
5. Conclusions

The classical assumption of proportionality between entropy and energy influxes via the temperature has been accepted even in the case of non-classical continuum theories, such as Green–Naghdi’s, in spite of the objections raised by Müller [1] and Liu [2,3] and of the fact that it has been found inappropriate in several circumstances (for instance, when one deals with thermodynamics of diffusion [3,10] and inconsistent with the kinetic theory of ideal gases (see [3,10] and references therein). As this assumption bears directly on the mathematical description of heat propagation, we have focused on rigid heat conductors within the framework of Green–Naghdi’s Type III theory, and demonstrated its exclusively constitutive nature by showing that:

(i) when heat conduction is isotropic (that is, in Green–Naghdi’s theory, when energy and entropy influxes are assumed to depend isotropically on the gradients of both the thermal displacement and its time rate), then entropy and energy influxes must be proportional, via a multiplier that needs not be the temperature, although it may depend on it;
(ii) when heat conduction is transversely isotropic, then proportionality cannot be postulated in general, because an additional entropy flux along the direction of transverse isotropy is possible.

We believe that these results carry on without any substantial change for material classes more general than that of rigid conductors. What postulation about entropy and energy influxes fits one or another non-isotropic type of heat propagation is an issue that remains open.

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Appendix A. Representation of an isotropic vector function of two vector arguments

Here, we derive a representation formula for isotropic vector-valued mappings of two vectorial arguments; although the result is well known [11,12], our proof is, to our knowledge, new. As a premiss, we recall that a scalar-valued function \( f \) of two vectorial arguments is isotropic, i.e. such that

\[
 f(u, v) = f(Qu, Qv), \quad \forall u, v \in \mathcal{V}, \quad \forall Q \in \text{Orth},
\]

where \( \mathcal{V} \) is a suitable vector space and Orth the full orthogonal group, if and only if

\[
 f(u, v) = \varphi(|u|, |v|, u \cdot v). \quad (A 1)
\]

**Lemma A.1.** Let \( \hat{f} \) be an isotropic mapping of \( \mathcal{V} \) into itself, i.e. such that

\[
 Q\hat{f}(u, v) = \hat{f}(Qu, Qv), \quad \forall u, v \in \mathcal{V}, \quad \forall Q \in \text{Orth}. \quad (A 2)
\]

Then, \( \hat{f} \) has the following representation:

\[
 \hat{f}(u, v) = \varphi_1(|u|, |v|, u \cdot v) u + \varphi_2(|u|, |v|, u \cdot v) v. \quad (A 3)
\]

**Proof.** Write (A 2) for \( Q \), a point of a smooth curve through \( \text{I} \), on the manifold Orth:

\[
 t \mapsto Q(t), \quad Q(t_0) = \text{I}, \quad \dot{Q}(t_0) = :W \in \text{Skw}, \quad (A 4)
\]

and differentiate at \( t = t_0 \), so as to get:

\[
 W\hat{f}(u, v) = \hat{F}_1(u, v)Wu + \hat{F}_2(u, v)Wv
\]

and

\[
 \hat{F}_1(u, v) := \partial_1 \hat{f}(u, v), \quad \hat{F}_2(u, v) := \partial_2 \hat{f}(u, v), \quad \forall u, v \in \mathcal{V}, \quad \forall W \in \text{Skw}. \quad (A 5)
\]
This condition is equivalent to
\[ W \cdot (a \otimes \hat{f}(u, v) - \hat{F}_1^T(u, v)a \otimes u - \hat{F}_2^T(u, v)a \otimes v) = 0, \] (A 6)
\[ \forall u, v, a \in V, \forall W \in \text{Skw}, \text{ and then} \]
\[ a \times \hat{f}(u, v) - \hat{F}_1^T(u, v)a \times u - \hat{F}_2^T(u, v)a \times v = 0, \quad \forall u, v, a \in V. \] (A 7)
Now, take the vector product with \( w := u \times v \) and use the identity \((a \times b) \times c = (a \cdot c)b - (b \cdot c)a\) so as to obtain:
\[ (a \cdot w)\hat{f}(u, v) - (\hat{f}(u, v) \cdot w)a = (a \cdot \hat{F}_1(u, v)w)u + (a \cdot \hat{F}_2(u, v)w)v \] (A 8)
\[ \forall u, v, a \in V. \text{ Because} \]
\[ (a \cdot \hat{F}_1(u, v)w)u = (u \otimes \hat{F}_1(u, v)w)a = (F_1^T(u, v)u \otimes w)a \] (A 9)
and
\[ (a \cdot \hat{F}_2(u, v)w)v = (v \otimes \hat{F}_2(u, v)w)a = (F_2^T(u, v)v \otimes w)a, \]
we have that
\[ (a \cdot w)\hat{f}(u, v) - (\hat{f}(u, v) \cdot w)a = \hat{F}_1^T(u, v)u \otimes w + \hat{F}_2^T(u, v)v \otimes w)a, \] (A 10)
\[ \forall u, v, a \in V. \text{ Given the arbitrariness of } a, \text{ it suffices to choose } a \cdot w = 0 \text{ to deduce that} \]
\[ \hat{f}(u, v) \cdot w = 0, \quad \forall u, v \in V \Leftrightarrow \hat{f}(u, v) = f_1(u, v)u + f_2(u, v)v. \] (A 11)
To conclude, one observes that this provisional form of \( \hat{f}(u, v) \) is compatible with (A 2) if \( f_1 \) and \( f_2 \) are isotropic. Thus, an application of (A 1) suffices.

**Appendix B. Two algebraic lemmas**

**Lemma B.1.** Let the algebraic equality
\[ \alpha I + \beta u \otimes u + \gamma v \otimes v + \delta (u \otimes v + v \otimes u) = 0 \] (B 1)
hold for all \( u, v \in V \). Then,
\[ \alpha = \beta = \gamma = \delta = 0. \] (B 2)

**Lemma B.2.** Let the algebraic equality
\[ \alpha I + \beta u \otimes u + \gamma v \otimes v + \delta (u \otimes v + v \otimes u) \]
\[ + \varphi e \otimes e + \chi (u \otimes e + e \otimes u) + \psi (v \otimes e + e \otimes v) = 0 \] (B 3)
hold for all \( u, v \in V \). Then,
\[ \alpha = \beta = \gamma = \delta = \varphi = \chi = \psi = 0. \] (B 4)

**References**

1. Müller I. 1971 The coldness, a universal function in thermoelastic bodies. *Arch. Ration. Mech. Anal.* 41, 319–332. (doi:10.1007/BF00281870)
2. Liu I-S. 1972 Method of Lagrange multipliers for exploitation of the entropy principle. *Arch. Ration. Mech. Anal.* 46, 131–148. (doi:10.1007/BF00250688)
3. Liu I-S. 1973 On the entropy supply in a classical and relativistic fluid. *Arch. Ration. Mech. Anal.* 40, 111–117. (doi:10.1007/BF00249878)
4. Podio-Guidugli P. 2013 On heat propagation in rigid conductors. In preparation.
5. Bargmann S, Steinmann P. 2007 Classical results for a non-classical theory: remarks on thermodynamics relations in Green–Naghdi thermo-hyperelasticity. *Contin. Mech. Thermodyn.* 19, 59–66. (doi:10.1007/s00161-007-0045-x)
6. Green AE, Naghdi PM. 1993 Thermoelasticity without energy dissipation. *J. Elast.* 31, 189–208. (doi:10.1007/BF00044969)
7. Coleman BD, Noll W. 1963 The thermodynamics of elastic materials with heat conduction and viscosity. *Arch. Ration. Mech. Anal.* 13, 167–178. (doi:10.1007/BF01262690)
8. Green AE, Naghdi PM. 1991 A re-examination of the basic postulates of thermomechanics. *Proc. R. Soc. Lond. A* 432, 171–194. (doi:10.1098/rspa.1991.0012)
9. Müller I. 1985 *Thermodynamics*. Boston, MA; London, UK; Melbourne, Australia: Pitman.
10. Liu I-S. 2009 On entropy flux of transversely isotropic elastic bodies. *J. Elast.* 96, 97–104. (doi:10.1007/s10659-009-9200-x)
11. Wang C-C. 1970 On representation for isotropic functions. *Arch. Ration. Mech. Anal.* 36, 166–197. (doi:10.1007/BF00272241)
12. Wang C-C. 1971 Corrigendum. *Arch. Ration. Mech. Anal.* 43, 392–395.