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Microwave-Induced Thermoacoustic Effect in Dielectrics and Its Coupling to External Medium—A Thermodynamical Formulation

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Abstract—**A thorough formulation of electromagnetic wave interaction with biological systems is presented. The thermodynamic process of the microwave-induced thermoacoustic generation is clearly defined. Couplings of the acoustic and thermal energies to the surrounding medium are included through consideration of discontinuities of thermodynamical variables and microwave exposure. Contrary to prior analyses, it is shown that acoustic waves may be generated by pulsed microwaves, even in the absence of inhomogeneity of microwave absorption, owing to discontinuities of thermodynamical variables and microwave exposure conditions across the interface. General equations for the thermoacoustic waves are derived, and the validity of the first-order linear approximation is estimated in terms of its percentage error. For a system with water as the absorbing dielectric interfacing with air of 1 atmosphere pressure, the first-order approximation becomes invalid for a peak specific absorption rate greater than 13 kW/gm.**

I. INTRODUCTION

THERMOACOUSTIC WAVES generated by pulsed microwaves have been cited as a mechanism for microwave hearing [1]–[6]. This effect has been experimentally shown to require an intact cochlea and 8th cranial nerve with a perceptional threshold near that of bone conduction. These facts imply that acoustic waves are transduced in the animal's skull by pulsed microwaves, then transmitted to the Organ of Corti by bone conduction. More recently, it has been demonstrated that acoustic waves are transduced in dielectric objects simulating the ocular lens when exposed to pulsed microwaves [7]. The putative effect has also been cited as the operant mechanism for cellular damage in studies of the murine ocular lens *in vitro* [8]. Several theories on the microwave auditory mechanism have been reviewed and compared by Lin [6]. Among them, it was found that only the thermoelastic mechanism could produce elastic waves of magnitude large enough to explain the experimental observations.

Manuscript received October 12, 1983; revised March 8, 1984. This work was supported in part by the Walter Reed Army Institute of Research through the U.S. Army Medical R&D Command under the U.S. Naval Sea Systems Command Contract N00024-83-C-5301.

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The theoretical derivation of the thermoelastic mechanism was based on the linear elastic equation with an external excitation derived from the thermal expansion that was produced by the absorption of incident microwave energy. The supposition was that an inhomogeneous microwave distribution in a dielectric body produced an inhomogeneous thermal expansion and thereby caused a mechanical stress. The pressure from the elastic excitation was then assumed to be the product of the spatial variation of the microwave energy and some thermal coefficients including the compressibility, thermal expansion coefficient, and specific heat. This approach fails to consider the proper balance of the distribution of the absorbed microwave energy among the internal thermal energies and the bulk kinetic energy. Indeed, since the thermodynamic process of the microwave acoustic effect was not formally identified, all thermal coefficients are not adequately defined. Furthermore, until now all theoretical analyses of the acoustic wave generation have been limited to either isolated systems or completely constrained systems [3]–[6], [9], both of which exclude the possibility of coupling to the external medium. If these conditions were taken verbatim, confirmatory experimental measurements would be difficult, if not impossible.

In this paper, a thorough formulation of the microwave-induced thermoacoustic effect in dielectrics and its coupling to external medium, based on thermodynamical considerations, is presented. Couplings of the acoustic and thermal energies to the surrounding medium are included through discontinuities of the thermodynamical variables and conditions of microwave exposure. Results of these couplings may provide a method for experimental evaluation of the microwave-generated pressure waves in small dielectric objects for which direct measurement is difficult. Solutions are derived for the first-order linear approximation. The percentage error is calculated, and the upper limit of the specific absorption rate consistent with its validity is also estimated. The ratio of radiated acoustic energy to absorbed microwave energy is also derived for a one-dimensional system.

II. THE FORMULATION

In this section, we shall present a general formulation of microwave-induced thermoacoustic effect. The objective is to derive the governing equations for the thermodynamical and mechanical variables, which include the mass density ρ or the specific volume $V = 1/\rho$, the pressure p, or, for an anisotropic medium, the stress tensor p_{ij}, the thermodynamic internal energy U_I, the bulk kinetic energy U_k, which pertains to the acoustic waves, the bulk velocity $\vec{v}$, the absolute temperature T, and the microwave specific absorption rate (SAR, the rate of energy absorption per unit mass) P. All these quantities depend on time t as well as on space $\vec{x}$. All thermodynamic extensive quantities are defined as per-mass quantities. We shall show that the thermoacoustic effect may essentially be described by four equations: the conservation of mass, the conservation of momentum, the conservation of energy, and the thermodynamic equation of state.

When a dielectric system is exposed to microwave radiation, microwave energy is absorbed into the dielectric through the work of the electric field on the electric polarization: $\mathscr{E}\,d\wp$. By means of molecular collisions, the energy is then partially converted into the internal energy U_I and partially into the bulk kinetic energy U_K via thermal expansion and other elastic vibrations. The process is complex and irreversible; however, we shall assume that the time domain which concerns us is much greater than the time it takes to reach a quasi-steady state and that every microscopically small part of the system is in local equilibrium such that thermodynamic variables, such as temperature, pressure, density, etc., may be defined. Then, from the first law of thermodynamics, which states that the net heat added to the system is equal to the sum of the work done by the system δw, and the increase in the internal and bulk energies [10], one has

$$\delta Q = d(U_I + U_K) + \delta w - \mathscr{E}\,d\wp \tag{1}$$

where d represents a perfect differential and δ an imperfect differential. Q is the total heat input which includes both conductive and convective heat flows, and $\mathscr{E}$ and $\wp$ are, respectively, the electric field and electric polarization. Note that we have separated the energy density into two terms: the first term U_I is the internal thermal energy density which depends on ρ and p, and the second term U_K, which is equal to $v^2/2$, is the bulk kinetic energy density; the velocity $\vec{v}$ is the bulk velocity of the acoustic vibration, which is to be distinguished from molecular thermal motions.

The time domain of our interest is the width of the microwave pulse, which is in the order of microseconds and is too short for any heat conduction and convective flow to take place. Therefore, δQ in (1) may be ignored. As to the last term on the right-hand side, the time cycle of the electric field is in the range of subnanoseconds, which is much shorter than the time domain of our concern. Therefore, we may average both $\mathscr{E}$ and $\wp$ over the microwave cycle. The resulting $\mathscr{E}\,d\wp$ represents the microwave energy absorbed by the dielectric, which we shall denote by δQ_p. Thus, (1) may be rewritten as

$$\delta Q_p = d(U_I + U_K) + \delta w. \tag{2}$$

In order to convert (2) into a differential equation with respect to space and time, we present a more general form of the equation of conservation. Let g be a thermodynamic extensive quantity and Sg be the rate of production of g, which may be due to internal production or external input, then the equation of conservation for g is

$$\partial_t g + \nabla \cdot (g\vec{v}) = Sg \tag{3}$$

where the first term on the left-hand side represents the rate of increase of g per unit volume, and the second term is the rate of outgoing flow of g from a unit volume. In terms of the convective derivative $d_t \equiv \partial_t + \vec{v} \cdot \nabla$, the above equation may also be written as

$$d_t g + g\nabla\cdot\vec{v} = Sg. \tag{4}$$

Let g be the density ρ, then (4) gives the equation of continuity

$$d_t\rho + \rho\nabla\cdot\vec{v} = 0. \tag{5}$$

On the other hand, if g is taken to be the momentum density ρv_i, the corresponding Sg is the force density, which is given by $-\partial_j p_{ij}$ ($\partial_j \equiv \partial/\partial x_j$), (4) then becomes

$$d_t(\rho v_i) + \rho v_i \nabla\cdot\vec{v} = -\partial_j p_{ij}$$

which is essentially a general form of Newton's second law. Note that we have adopted the summation convention, i.e., repeated indices are to be summed through the degrees of freedom. With the help of (5), the above equation may also be written as

$$\rho\, d_t v_i + \partial_j p_{ij} = 0 \tag{6}$$

which is Euler's equation [11]. Similarly, if g is taken to be the energy density $\rho(U_I + U_K)$, the corresponding Sg is the microwave power absorption per unit volume $\rho P(\vec{x}, t)$ plus the rate of work by the stress tensor, which is $-\nabla\cdot(\vec{v}\cdot\vec{\vec{p}})$. Again, using (5), the conservation law for energy reads

$$\rho\, d_t(U_I + U_K) = -\nabla\cdot(\vec{v}\cdot\vec{\vec{p}}) + \rho P.$$

Since $U_K = v^2/2$, by (6), $\rho\, d_t U_K + v_i\,\partial_j p_{ij} = 0$, the above equation may then be written as

$$\rho\, d_t U_I = -p_{ij}\,\partial_j v_i + \rho P. \tag{7}$$

Thus, the first term on the right-hand side of the above equation is the part of the mechanical work done to the system that contributes to the thermal internal energy. This fact will be used later to derive the expression of the internal energy in terms of other thermodynamical variables.

Equations (5)–(7) all together contain four unknowns, viz., ρ, $\vec{v}$, $\vec{\vec{p}}$, and U_I. One more equation is needed to complete the system, which will be furnished by the thermodynamic equation of state

$$U_I = U(\rho, \vec{\vec{P}}).$$

Here, we have ignored the dependency of U_I on the electric polarization $\boldsymbol{P}$ since the cycle of the microwave is much shorter than the time domain of our interest, the microwave pulse width, which is in the order of microseconds. As a thermodynamical system, the dielectric has two independent thermodynamic variables, which may be any pair among the variables ρ, T, $\vec{\vec{p}}$, and U, and all other variables may be considered as functions of the pair. Since ρ and $\vec{\vec{p}}$ are the variables in (5)–(7), they are selected as the independent variables. A differential of U_I may be expressed as

$$dU_I = (\partial_\rho U_I)_{p_{ij}}\, d\rho + (\partial_{p_{ij}} U_I)_\rho\, dp_{ij}. \tag{8}$$

We would like to express the partial derivatives of U_I with respect to ρ and p_{ij} in terms of known empirical quantities such as specific heat and compressibility. To do this, we use the Jacobian method of relating one partial derivative to another. In the following derivation, we shall first drop the indices i, j of the stress tensor p_{ij}, and later restore them after the result is derived. The Jacobian $J(u,v)_{x,y}$ is defined as

$$J(u,v)_{x,y} = \det\begin{pmatrix} (\partial_x u)_y & (\partial_y u)_x \\ (\partial_x v)_y & (\partial_y v)_x \end{pmatrix}.$$

It is straightforward to show that the ratio of $J(u,v)_{x,y}$ to $J(r,s)_{x,y}$ is equal to $J(u,v)_{r,s}$, independent of which pair of (x, y) is used for differentiation. Thus, we shall drop the subscripts whenever a ratio of Jacobians is concerned. It may also be easily shown that $J(u,v)_{x,y} = -J(v,u)_{x,y} = J(v,u)_{y,x}$, and $J(u,v)/J(x,v) = (\partial_x u)_v$. Using these identities, one may then derive from (8) the following equation:

$$\begin{aligned} dU_I &= (\partial_\rho U_I)_p\, d\rho + (\partial_p U_I)_\rho\, dp \\ &= \frac{J(U_I,p)}{J(\rho,p)}\, d\rho + \frac{J(U_I,\rho)}{J(p,\rho)}\, dp \\ &= \frac{J(U_I,p)/J(T,p)}{J(\rho,p)/J(T,p)}\, d\rho + \frac{J(U_I,\rho)/J(T,\rho)}{J(p,\rho)/J(T,\rho)}\, dp \\ &= \frac{J(U_I,p)/J(T,p)}{J(\rho,p)/J(T,p)}\, d\rho + C_v\frac{J(T,\rho)/J(p,T)}{J(p,\rho)/J(p,T)}\, dp \\ &= \left(-\frac{C_p}{\beta_p}\frac{1}{\rho} + \frac{p}{\rho^2}\right) d\rho + \left(\frac{C_v\kappa_T}{\beta_p}\right) dp \end{aligned} \tag{9}$$

where C_v and C_p are the specific heat (per mass) at constant volume and constant pressure, respectively, β_p is the isobaric thermal expansion coefficient, and κ_T the isothermal compressibility. In deriving the last equality of the above equation, the following relationships have also been used:

$$\beta_p = (1/V)(\partial_T V)_p = (1/V)J(V,p)/J(T,p)$$

$$\kappa_T = -(1/V)(\partial_p V)_T = -(1/V)J(V,T)/J(p,T)$$

and, since from the first law of thermodynamics, $\delta Q = dU_I + p\,dV = dU_I + p d(1/\rho)$

$$C_v = (\partial_T U_I)_\rho = J(U_I,\rho)/J(T,\rho)$$

$$\begin{aligned} C_p &= (\partial_T U_I)_p + [\partial_T(p/\rho)]_p \\ &= J(U_I,p)/J(T,p) - (p/\rho^2)J(\rho,p)/J(T,p). \end{aligned}$$

It is straightforward to generalize the above result to include the stress tensor $\vec{\vec{p}}$. Including the indices i, j of the stress tensor, (9) takes the form

$$\begin{aligned} dU_I = {}& \left[-C_p/(\rho\beta_p) + (1/\rho^2)p_{ij}\,\partial_j v_i/\nabla\cdot\vec{v}\right] d\rho \\ & + (C_v/\beta_p)(\kappa_T)_{ij}\, dp_{ij} \end{aligned} \tag{10}$$

where C_v and β_p are as expressed above, except that now $\vec{v}$ must also be included in the independent variables and be kept constant in the partial derivatives; C_p is given by

$$\begin{aligned} C_p &= (\partial_T U_I)_{p,\vec{v}} + \left[(\partial_T)\left(p_{ij}\,\partial_j v_i/(\rho\nabla\cdot\vec{v})\right)\right]_{p,\vec{v}} \\ &= J(U_I, p)/J(T,p) - \left[(1/\rho^2)\,p_{ij}\,\partial_j v_i/\nabla\cdot\vec{v}\right] \\ &\quad \cdot J(\rho,p)/J(T,p) \end{aligned}$$

and $(\kappa_T)_{ij}$ is the compressibility tensor defined as

$$\begin{aligned} (\kappa_T)_{ij} &= -(1/V)(\partial V/\partial p_{ij})_T \\ &= -(1/V)J(V,T)/J(p_{ij},T). \end{aligned}$$

In deriving the above results, we have used the generalized first law of thermodynamics for mechanically anisotropic media

$$\delta Q = dU_I + (p_{ij}\,\partial_j v_i/\nabla\cdot\vec{v})\,dV.$$

To understand the above equation, we consider δQ as the total heat input to a material element of unit mass during time dt, during which the total energy is increased by dU. The total work done by the element is equal to $\nabla\cdot(\vec{v}\cdot\vec{p})\,dt$ times the volume, which is $1/\rho$. On the other hand, from (5), the increment in volume during dt is $dV = d(1/\rho) = (1/\rho)\nabla\cdot\vec{v}dt$. Therefore, the work done associated with a volume increment dV is $[\nabla\cdot(\vec{v}\cdot\vec{p})/\nabla\cdot\vec{v}]dV$, which may be separated into two terms: $(p_{ij}\,\partial_j v_i/\nabla\cdot\vec{v})\,dV + (v_i\,\partial_j p_{ij}/\nabla\cdot\vec{v})\,dV$, where the second term contributes to the bulk kinetic energy U_K, while the first term contributes to the thermodynamical internal energy U_I. It is the internal energy to which the thermodynamical variables such as the compressibility and specific heat are related.

It is remarked that we have ignored the internal energy associated with the anisotropicity of the strain tensor, simply to reduce complexity. In other words, we have assumed that the energy associated with different directions of strain are approximately equal, and, therefore, the strain tensor may be represented by the increment in the specific volume, or equivalently, the increment in the density, as it is expressed in the first term in the right-hand side of (10).

To summarize, we list the four equations of the system, viz. (5)–(7) and (10) as follows:

$$d_t\rho + \rho\nabla\cdot\vec{v} = 0 \tag{11}$$

$$\rho\, d_t v_i + \partial_j p_{ij} = 0 \tag{12}$$

$$P_{ij}\,\partial_j v_i + \rho\, d_t U_I = \rho P \tag{13}$$

$$\begin{aligned} d_t U_I = \left[-C_p/(\rho\beta_p) + (1/\rho^2)\,p_{ij}\,\partial_j v_i/\nabla\cdot\vec{v}\right] d_t\rho \\ + (C_v/\beta_p)(\kappa_T)_{ij}\,d_t p_{ij}. \end{aligned} \tag{14}$$

The convective derivative $d_t = \partial_t + \vec{v}\cdot\nabla$ may be cumbersome. It is possible to transfer the convective derivative into the simple derivative ∂_t by the transformation from the Eulerian specification to the Lagrangian specification [11]. The former refers to the specification of using the space vector $\vec{x}$ to describe a material element, whereas the latter attaches a fix vector to each material element. So, if we specify each dielectric element by its initial value of $\vec{x}$, denoted by $\vec{q}, \vec{q} = \vec{x}(t=0)$, then a description of, say, pressure, may be expressed by $p(\vec{x},t)$, referring to the pressure at time t and space point $\vec{x}$, or expressed by $p(\vec{q})$, referring to the pressure at time t and the space point occupied by the element which was at the space point $\vec{q}$ when $t=0$. Denote by $\vec{u}(\vec{q},t)$ the displacement of the $\vec{q}$-element from time 0 to time t, then $\vec{x}(\vec{q},t) = \vec{q} + \vec{u}(\vec{q},t)$. The velocity $\vec{v}(\vec{x},t)$, or $\vec{v}(\vec{q},t)$, is defined as $(\partial\vec{u}/\partial t)_{\vec{q}}$, from which one obtains, for any function $f(\vec{q},t)$ or $f(\vec{x},t)$

$$(\partial_t f)_{\vec{q}} = (\partial_t f)_{\vec{x}} + (\vec{v}\cdot\nabla_{\vec{x}})f$$

which is exactly the convective derivative of f. Since $\partial/\partial x_i = (\partial q_j/\partial x_i)\partial/\partial q_j$ (using the summation convention) and $u_i = x_i - q_i$, one has

$$(\partial q_k/\partial x_i)(\delta_{jk} + \partial u_j/\partial q_k) = \delta_{ij}.$$

Thus, denoting by A the matrix $A_{ij} = \delta_{ij} + \partial u_j/\partial q_i$, one has the following identities:

$$\begin{aligned} (A^{-1})_{ij} &= \partial q_j/\partial x_i \\ \partial/\partial x_i &= (A^{-1})_{ij}\,\partial/\partial q_j. \end{aligned}$$

With these identities, the transformation to the Lagrangian specification brings (11)–(14) into the following form:

$$(\partial_t f)_{\vec{q}} = (\partial_t f)_{\vec{x}} + (\vec{v}\cdot\nabla_{\vec{x}})f$$

$$\partial_t\rho + \rho(A^{-1})_{ij}\,\partial_j v_i = 0 \tag{15}$$

$$\rho\,\partial_t v_i + (A^{-1})_{jk}\,\partial_k p_{ij} = 0 \tag{16}$$

$$\rho\,\partial_t U_I = P_{ij}(A^{-1})_{jk}\,\partial_k v_i = \rho P(\vec{q},t) \tag{17}$$

$$\begin{aligned} \partial_t U_I = \left[-C_p/(\rho\beta_p) + (1/\rho^2)\,p_{ij}\,\partial_j v_i/\nabla\cdot\vec{v}\right]\partial_t\rho \\ + (C_v/\beta_p)(\kappa_T)_{ij}\,\partial_t p_{ij} \end{aligned} \tag{18}$$

where $\partial_t = (\partial/\partial t)_{\vec{q}}$ and $\partial_i = (\partial/\partial q_i)_t$.

Equation (15) can be immediately integrated without involving the other three equations. It is straightforward to verify that the following is its solution:

$$\rho(\vec{q},t) = \rho(\vec{q},0)/\det(A) \tag{19}$$

where $\det(A)$ is the determinant of the matrix A. Equations (18) and (19) may be substituted into (16) and (17) to eliminate U_I, then we are left with (16) and (17) with two unknowns, viz., $\vec{u}$ and p_{ij} (note that $\vec{v} = \partial_t\vec{u}$), to solve. However, they are nonlinear equations, especially the thermodynamical variables C_v, C_p, β_p, and κ_T may be complicated functions of ρ and p. In the next section, we shall further simplify these equations by making the first-order linear approximation. It is remarked here that, instead of using (14) or (18), one may substitute it by another thermodynamic equation of state, which may be any interrelationship among the variables. The principal content of most biological dielectrics is water, for which many empirical equations of state are available [12].

III. THE LINEAR APPROXIMATION

In this section, we consider the first-order approximation and reduce the nonlinear equations (15)–(18) to linear ones. The order of magnitude here refers to the variation of all the quantities, including the displacement $\vec{u}, \vec{v} = \partial_t\vec{u}, \rho, p$, and all thermodynamical coefficients, from their initial values. We shall assume that initially the system is in an equilibrium state with density ρ_0, pressure p_0, temperature T_0, and $\vec{u}=\vec{v}=0$; in general, ρ_0, p_0, and T_0 may still depend on the space vector $\vec{q}$. The quantities $\vec{u}$ and $\vec{v}$ are therefore first-order quantities. In the first-order approximation, the matrix A in (15)–(17) may then be taken to be unity. Note also that, to the first order, there is no difference between the Eulerian specification and the Lagrangian specification, so one may use either the set of equations (11)–(14) or the set (15)–(18). Substituting (18) for the first term in (17), (15)–(17) then become

$$\partial_t\rho+\rho\nabla\cdot\vec{v}=0 \tag{20}$$

$$\rho\,\partial_t v_i+\partial_j p_{ij}=0 \tag{21}$$

$$-(C_p/\beta_p)\,\partial_t\rho+(\rho C_v/\beta_p)(\kappa_T)_{ij}\,\partial_t p_{ij}=\rho P \tag{22}$$

where, in obtaining the last equation, (20) has been used to cancel two terms involving $\partial_j v_i$. Equations (20) and (21) may be used to eliminate $\vec{v}$ and yield

$$\partial_t\big[(1/\rho)\,\partial_t\rho\big]-\partial_i\big[(1/\rho)\,\partial_j p_{ij}\big]=0. \tag{23}$$

Finally, (22) and (23) may work together to eliminate ρ to yield the equation for p_{ij}

$$\partial_t\big[(C_v/C_p)(\kappa_T)_{ij}\,\partial_t p_{ij}\big] - \partial_i\big[(1/\rho)\,\partial_j p_{ij}\big]=\partial_t\big[(\beta_p/C_p)P\big]. \tag{24}$$

Note that, apart from the assumption that the matrix A (which was defined in the paragraph before (15)) is unity, we have not used other first-order approximations to derive (20)–(24). In general, if the strain tensor $\partial_i u_j$ is isotropic so the matrix A is a multiple of unity, then all the above equations are derivable from (15)–(18), even in nonlinear cases, except for a strain factor $A^{-1}=(\rho/\rho_0)^{1/3}$, which must be multiplied by the second terms of (20) and (21) and other equations derived therefrom. It is also worthwhile noting that, for isotropic media, one has the following relationship between ρ and p, without any approximation:

$$-(C_p/\beta_p)\,\partial_t\rho+(\rho C_v\kappa_T/\beta_p)\,\partial_t p=\rho p. \tag{25}$$

This result is obtained from (17) by substituting, in its left-hand side, the first term with the expression from (18), and the second term with $(p/\rho)\,\partial_t\rho$ obtained from (15).

In the first-order approximation, $\partial_t p_{ij}$ and $\partial_j p_{ij}$ are the first-order terms, so one may approximate the factors $(1/\rho)$, (C_v/C_p), (β_p/C_p), and $(\kappa_T)_{ij}$ in (24) by their respective equilibrium values to yield the linear wave equation for the stress tensor p_{ij}

$$\big[(\rho C_v/C_p)(\kappa_T)_{ij}\big]_0\,\partial_t^2 p_{ij}-\partial_i\,\partial_j p_{ij}=(\rho\beta_p/C_p)_0\,\partial_t P \tag{26}$$

where $[\cdots]_0$ denotes that the quantity inside the brackets is evaluated at its equilibrium value. In the principal axis of a mechanically anisotropic dielectric, p_{ij} is diagonal and, from its definition described in the last section, so is $(\kappa_T)_{ij}$. Then, the above equation gives the wave equation for the stress in each principal axis, with the wave velocity $c_i=[C_p/(\rho C_v(\kappa_T)_i)]_0^{1/2}$ in the ith principal axis.

For mechanically isotropic systems, (20)–(22) may also be simplified to obtain linear wave equations for $\partial_t\rho$, p, and $\vec{v}$. They are, with c denoting $[C_p/(\rho C_v\kappa_T)]_0^{1/2}$, the acoustic wave velocity in the dielectric

$$\partial_t^2(\partial_t\rho)-c^2\nabla^2(\partial_t\rho)=\big[c^2\rho\beta_p/C_p\big]_0\nabla^2 P \tag{27}$$

$$\partial_t^2 p-c^2\nabla^2 p=\big[\rho c^2\beta_p/C_p\big]_0\,\partial_t P \tag{28}$$

$$\partial_t^2\vec{v}-c^2\nabla^2\vec{v}=-\big[c^2\beta_p/C_p\big]_0\nabla P. \tag{29}$$

Equation (28) comes directly from (26), from which one obtains (27) with the help of (25). As to the equation for $\vec{v}$, it is derived from (21) and (28) with the assumption that $\vec{v}$ and P vanish initially. Equation (27) is similar to what has been suggested by other authors [3]–[6], [9], except that here we have all the thermodynamical quantities well defined. It is also interesting to note that, as one can see from these equations, acoustic waves are generated either by an inhomogeneity or temporal variation of the product of the microwave specific absorption rate and some thermodynamical variables. The inhomogeneity or the temporal variation may be intrinsic in the microwave source or the dielectric absorption, or may be simply due to some discontinuities in the thermodynamical variables. As we shall see in the next section, acoustic waves are generated in dielectrics even when the absorption of microwave pulses is spatially homogeneous, which is generally the case for small dielectric objects, such as ocular lenses. It is also remarked that the isotropic assumption may be an over simplification for biological tissues which are generally anisotropic thermally, as well as elastically, especially for tissues composed of cells with a large aspect ratio, which is the case for ocular lenses.

Equations (27)–(29) are the first-order thermoelastic wave equations for isotropic media without viscous damping under microwave radiation. These equations must be supported by the respective boundary conditions, which are implicated in (20)–(22). To extract these boundary conditions for a system with specific interface that separates one medium from another, we first integrate both sides of (20) across a thin layer of the boundary surface and, by Stoke's theorem, obtain the boundary condition for $\hat{n}\cdot\vec{v}$, where $\hat{n}$ is a unit vector normal to the boundary surface. With this condition, (21) then gives the condition for $(1/\rho)\hat{n}\cdot\nabla p$. Similarly, integrating both sides of (21) across the surface yields the condition for p, which, together with the condition for $(1/\rho)\hat{n}\cdot\nabla p$ and the relationship expressed by (25), gives the boundary conditions for $\partial_t\rho$. Finally, expressing $\nabla\cdot\vec{v}$ in terms of $\partial_t\rho$ with the help of (20), the boundary condition for $\partial_t\rho$ gives a condition for $\nabla\cdot\vec{v}$. Listed below are these boundary conditions.

1) For ρ, the following two quantities must be continuous:

$$c^2\big[\partial_t\rho+(\rho\beta_p/C_p)P\big] \quad \text{and } c^2\hat{n}\cdot\big[(1/\rho)\,\partial_t\nabla\rho+(\beta_p/C_p)\nabla P\big]. \tag{30}$$

2) For p, the following two quantities must be continuous:

$$p \text{ and } (1/\rho)\hat{n}\cdot\nabla p. \tag{31}$$

3) For $\vec{v}$, the following two quantities must be continuous:

$$\hat{n}\cdot\vec{v} \text{ and } c^2\rho\left[\nabla\cdot\vec{v}-\left(\beta_p/C_p\right)P\right]. \tag{32}$$

With these boundary conditions, solution of the above wave equations is mathematically straightforward. In the next section, we shall consider a one-dimensional case and obtain some numerical results for the microwave-induced thermoacoustic effect in water, which is the main dielectric in biological materials.

IV. THE ONE-DIMENSIONAL CASE

To illustrate the implications of the formulation described in the last two sections, we consider a dielectric system that has only one degree of freedom in mechanical vibration, such as water in a solid tube which is closed at one end and open to the atmospheric pressure at the other end. Instead of q, we shall use x to denote the spatial coordinate in the Lagrangian specification, so x is the position where the concerned material element is located at time $t=0$. For a one-dimensional system, (19) gives $\rho(x,t)=\rho_0(x)/A(x,t)$, where $A(x,t)=\partial_x u$, $u(x,t)$ is the displacement of the material element, and $\rho_0(x)=\rho(x,0)$ describes the initial density of the system; note that ρ_0 may be x-dependent. Replacing A^{-1} by ρ/ρ_0, (15)–(17) then become

$$\rho_0\partial_t\rho+\rho^2\partial_x v=0 \tag{33}$$

$$\rho_0\partial_t v+\partial_x p=0 \tag{34}$$

$$\rho_0\partial_t U_I+p\partial_x v=\rho_0 P(x,t). \tag{35}$$

Equation (18) may then be used to substitute for the first term on the left-hand side of (35). Upon multiplying these equations by suitable factors and differentiating both sides with respect to t or x, then subtracting one equation from another, they may be decoupled to yield the equations for the variables ρ, p, and v, respectively. Boundary conditions may also be obtained from similar considerations, as we did in the previous section. Below we list the results.

1) The equation and boundary conditions for $\partial_t(\rho_0/\rho)$ are

$$\left[\partial_x\rho_0^{-1}\partial_x\left(\rho C_p/(\rho_0 C_v\kappa_T)\right)-\partial_t^2\right]\partial_t(\rho_0/\rho) = \partial_x\rho_0^{-1}\partial_x\left(\beta_p/(C_v\kappa_T)P\right) \tag{36}$$

with the following two quantities continuous across the boundary:

$$\left[\left(\rho C_p/(\rho_0 C_v\kappa_T)\right)\partial_t(\rho_0/\rho)-\left(\beta_p/(C_v\kappa_T)\right)P\right]$$

and

$$(1/\rho_0)\partial_x\left[\left(\rho C_p/(\rho_0 C_v\kappa_T)\right)\partial_t(\rho_0/\rho)-\left(\beta_p/(C_v\kappa_T)\right)P\right]. \tag{37}$$

2) The equation and boundary conditions for p are

$$\left[\partial_x\rho_0^{-1}\partial_x-\partial_t\left(\rho_0 C_v\kappa_T/(\rho C_p)\right)\partial_t\right]p(x,t) = -\partial_t\left(\rho_0\beta_p/(\rho C_p)P\right) \tag{38}$$

with the following two quantities continuous across the boundary:

$$p(x,t) \text{ and } \rho_0^{-1}\partial_x p(x,t). \tag{39}$$

3) The equation and boundary conditions for $v(x,t)$ are

$$\left[\partial_x\left(\rho C_p/(\rho_0 C_v\kappa_T)\right)\partial_x-\rho_0\partial_t^2\right]v(x,t) = \partial_x\left(\beta_p/(C_v\kappa_T)P\right) \tag{40}$$

with the following two quantities continuous across the boundary:

$$v(x,t) \text{ and } (1/\rho_0 C_v\kappa_T)\left(\rho C_p\partial_x v-\rho_0\beta_p P\right). \tag{41}$$

The above equations are exact without any approximation; note also that all thermodynamical coefficients depend, in general, on density and pressure and are, therefore, space and time dependent.

In the first-order approximation, the above equations reduce to common linear wave equations; they are given by (27)–(29) with ∇ replaced by ∂_x. To illustrate the implications of these equations and their respective boundary conditions, we consider a thin solid tube of arbitrary shape in the cross section filled with a dielectric liquid to height l, with one end of the tube closed and the other end open to one atmosphere pressure of air. Let A denote the cross section of the tube and x denote the position of a liquid element along the length of the tube; the closed end of the tube is taken to be $x=0$ and the open end $x=l$. A microwave pulse of duration τ is incident from the side perpendicular to the axis of the tube. The tube is assumed to be thin enough with respect to the microwave wavelength that the induced electric field is uniform throughout the cross section of the tube. It is also assumed that the microwave absorption by the liquid exhibits no spatial variation along the axis of the tube. Therefore, the system is one-dimensional and the specific absorption rate (SAR), $P(t)$, is independent of x. We shall calculate the total acoustic energy coupled to the air per unit microwave energy absorption by the liquid. For this purpose, we need only to obtain the solution for v. With the configuration described above, the boundary conditions are, from (32)

$$v_1(0,t)=0, \quad \partial_x v_1(0,t)=0 \tag{42}$$

$$v_1(l,t)=v_2(l,t) \tag{43}$$

and

$$c_1^2\rho_{10}\left[\partial_x v_1(x=l,t)-\left(\beta_p/C_p\right)_1 P(t)\right] = c_2^2\rho_{20}\partial_x v_2(x=l,t)$$

where the subscripts 1 and 2 label, respectively, the dielectric liquid and the air, and $c_{1,2}$ denotes the corresponding value of $(C_p/(\rho C_v\kappa_T))^{1/2}$ in either medium, which is the velocity of the acoustic wave in the medium.

It is more convenient to work in the frequency space using the Fourier transformation. The Fourier transformation between any function $g(t)$ and its counterpart $g(\omega)$ is defined as

$$g(t) = (2\pi)^{-1/2} \int_{-\infty}^{\infty} g(\omega) e^{-i\omega t} d\omega \qquad (44)$$

$$g(\omega) = (2\pi)^{-1/2} \int_{-\infty}^{\infty} g(t) e^{i\omega t} dt \qquad (45)$$

where, since there will be no ambiguity, the same notation is used for the function in either frequency or time. Note also that the frequency ω refers to the acoustic frequency, which is to be distinguished from the microwave carrier frequency. If $P(t)$, the SAR, is equal to P_0 during the pulse width τ, then, in the frequency space, it is given by

$$P(\omega) = -P_0(2\pi)^{1/2}[1 - e^{i\omega\tau}] / [2\pi i(\omega - i0)] \qquad (46)$$

where $-i0$ in the denominator means that the singularity is immediately below the real axis of the complex ω-plane. The solutions for the velocity in both media are given by

$$\begin{pmatrix} v_1(x,\omega) \\ v_2(x,\omega) \end{pmatrix} = \frac{P(\omega)}{\omega - i0} \frac{(c\beta_p/C_p)_1}{\cos(k_1 l) - i\tan\phi \sin(k_1 l)} \cdot \begin{pmatrix} \sin(k_1 x) \\ \sin(k_1 l)\exp[ik_2(x-l)] \end{pmatrix} \qquad (47)$$

where $k_i = \omega/c_i$, the acoustic wavenumber in either medium, and $\tan\phi = (\rho c)_2/(\rho c)_1$, the ratio of the acoustic impedances of the air to the liquid.

The coupling of the energy from the liquid to the air may be calculated easily using the above result. Noting that the pressure in either side of the interface is equal, owing to the first boundary condition in (39), at any time dt, the work done by the liquid to the air per unit surface of the interface is $p_1 v_1\, dt$. Since $v_1(x,t)$ is a first-order quantity, so, in the first-order approximation of $p_1 v_1\, dt$, one may replace p_1 by p_0, the initial equilibrium atmospheric pressure. Thus, the total acoustic energy coupled into the air induced by the absorption of a microwave pulse by the liquid is

$$E_{\text{air}} = Ap_0 \int_{-\infty}^{\infty} v_1(l,t)\, dt$$

where A is the area of the cross section of the tube. From the Fourier transformation (45), the integral $\int_{-\infty}^{\infty} v_1(l,t)\, dt$ is equal to $(2\pi)^{1/2}$ times the value of $v_1(l,\omega)$ evaluated at $\omega = 0$, which may be easily obtained from (47). Thus, the above equation gives

$$E_{\text{air}} = (2\pi)^{1/2} lAp_0(\beta_p/C_p)_1 P(\omega = 0).$$

On the other hand, the total microwave energy absorption by the liquid is

$$E_{\text{abs}} = lA \int_{-\infty}^{\infty} P(t)\, dt = (2\pi)^{1/2} lA\rho_1 P(\omega = 0).$$

Comparing the above two equations, one obtains the total acoustic energy coupled to the air induced by the absorption of a unit microwave energy by the liquid

$$\eta_{\text{air}} = E_{\text{air}}/E_{\text{abs}} = p_0(\beta_p/(\rho C_p))_1. \qquad (48)$$

It is interesting to note that the coupling efficiency of the acoustic energy to a nonabsorbing medium due to the absorption of a microwave pulse by a liquid is proportional to the pressure in the nonabsorbing medium and, apart from that, is independent of its other thermodynamic properties. It is also independent of the microwave pulse width, so that total acoustic energy coupled to the air is proportional to the total microwave absorption cross section. This is simply due to the linear approximation. Note also that the coupling efficiency calculated above does not tell the coupling efficiency of the microwave energy absorption to the acoustic energy generated in the liquid, which may be calculated using the same approach through the solutions for $p_1(x,t)$ and $v_1(x,t)$. Although previous studies have calculated the coupling of microwave to elastic wave in the absorbing medium [1], [2], they are limited to either a totally constrained or totally isolated body without coupling to the surrounding medium, and are also based on the simplified assumption of the temperature-induced stress due to inhomogeneous microwave absorption in the dielectric, which failed to consider the proper distribution of the absorbed energy between the internal energy and the bulk kinetic energy.

The pressure waves in both media may also be calculated using the same approach. For p, the boundary conditions are, from (31)

$$\partial_x p_1(x=0,t) = 0$$

$$p_1(l,t) = p_2(l,t)$$

and

$$(1/\rho_{10})\,\partial_x p_1(x=l,t) = (1/\rho_{20})\,\partial_x p_2(x=l,t).$$

In the frequency domain, the solutions for the pressures are

$$p_1(x,\omega) = i\left(\frac{c^2\rho\beta_p}{C_p}\right)_1 \frac{P(\omega)}{\omega - i0} \cdot \left[\frac{\cos(k_1 x)}{\cos(k_1 l) - i\tan\phi\sin(k_1 l)} - 1\right] \qquad (49)$$

$$p_2(x,\omega) = -\left(\frac{c^2\rho\beta_p}{C_p}\right)_1 \frac{P(\omega)}{\omega - i0} \cdot \frac{\tan\phi\sin(k_1 l)\exp[ik_2(x-l)]}{\cos(k_1 l) - i\tan\phi\sin(k_1 l)}. \qquad (50)$$

To obtain the pressures in the time domain, we perform the Fourier transformation on the above results and employ the contour integration in the complex ω-plane. The results are summations of the time-series of the form $\exp(i\omega_n t)$, with ω_n being the zeros of the factor $[\cos(\omega l/c_1) - i\tan\phi\sin(\omega l/c_1)]$, corresponding to the poles of $p_i(x,\omega)$. A straightforward analysis of this factor shows that these poles are located at the following points in the ω-plane:

$$\omega_n = (c_1/2l)(n\pi - ir_0) \tag{51}$$

where n is any odd number if $\tan\phi$ is smaller than 1, and any even number if it is greater than one, and r_0 is the solution of the equation $\coth(r_0)+(-1)^n \operatorname{csch}(r_0) = \tan\phi$. The pressure waves in the time domain are then

$$p_1(x,t) = \sum{}_{\substack{n = \text{odd integers if } \tan\phi < 1 \\ n = \text{even integers if } \tan\phi > 1}} (-1)\frac{n+1}{2}\sqrt{2\pi}\left(\frac{\rho\beta_p c^3}{C_p l}\right)_1 \frac{P(\omega_n)}{\omega_n}\cos\left(\frac{\omega_n x}{c_1}\right)\exp(-i\omega_n t) \tag{52}$$

$$p_2(x,t) = \sum\begin{pmatrix} n = \text{odd integers if } \tan\phi < 1 \\ n = \text{even integers if } \tan\phi > 1 \end{pmatrix} - i\sqrt{2\pi}\left(\frac{\rho\beta_p c^3}{2C_p l}\right)_1 \frac{P(\omega_n)}{\omega_n}\tan(2\phi)\exp\left[i\omega_n\left(\frac{x-l}{c_2} - t\right)\right] \tag{53}$$

where, for $\phi > \pi/4$, $[\cos(2\phi)]^{1/2}$ is taken to be positive imaginary.

For the liquid–air system, only poles with an odd value of n exist and, since $\tan\phi$ is about 3×10^{-4}, r_0 is negligibly small. Note that the amplitudes of the pressure waves are linear in the peak SAR. However, due to the dependence on the poles ω_n, they are not linear in the total absorbed energy, which is in contrast to the coupling efficiency η_{air}. For each frequency component ω_n, the corresponding amplitude exhibits some resonance phenomena according to the value of the pulse width. For the liquid–air system, the fundamental acoustic frequency is approximately $c_1/(4l)$, corresponding to a quarter-wave in the length of the liquid, for which the resonance values of the pulse width are given by odd-integer multiples of $2l/c_1$. Similar resonance behavior has also been derived previously for totally constrained or totally isolated systems [6].

V. REMARKS AND DISCUSSIONS

A thorough thermodynamical formulation of a microwave-induced thermoacoustic effect has been presented. It is shown that acoustic waves are generated if there is any discontinuity, either spatial or temporal, in the radiation density, the absorption coefficient, or the thermodynamic variables such as the mass density ρ, the isobaric or isovolumetric specific heat, C_p and C_v, respectively, the isothermal compressibility κ_T, or the isobaric thermal expansion coefficient β_p. Therefore, even for a small biological subject, such as an eye lens, in which the microwave power deposit is practically uniform, acoustic waves can still be generated by the microwave absorption.

Coupling of the acoustic energy from a microwave absorbing medium to a nonabsorbing medium is also formulated through the thermodynamic equations and the boundary conditions. Since direct measurement in small biological objects is difficult, these results may be useful for the evaluation of the pressure waves inside small biological objects by measuring the pressure waves (or other thermodynamical quantities) in the surrounding medium. The result of the linear approximation indicates that the percentage of the absorbed microwave energy being coupled to a nonabsorbing medium is, apart from its pressure in the equilibrium state, independent of the thermodynamic properties of the nonabsorbing medium. Taking the dielectric liquid to be water, then $\beta_p = 2.8\times10^{-4}{}^\circ\text{C}^{-1}$, $C_p = 4.186\times10^7$ erg/gm-°C, and $\rho = 1$ gm/cm^3. With the equilibrium pressure of the air p_0 being 1 atm, which is 1.01×10^6 dyne/cm^2, then (48) gives $\eta_{\text{air}} = 6.7\times10^{-6}$. This also gives the lower limit of the conversion efficiency of the microwave energy to acoustic energy in the absorbing liquid, since part of the generated acoustic energy in the liquid is converted into thermal energy through viscosity and conducted away to the surrounding medium.

The formulation indicates that the microwave-induced thermoacoustic effect is, in general, nonlinear. At low radiation density, one may make the linear approximation by ignoring all second-order and higher order terms. It is worthwhile to evaluate the validity of the linear approximation. This can be done by calculating the generated pressure wave and comparing it to the initial equilibrium value. Taking the fundamental frequency component in (52), the corresponding amplitude for $p_1(x,t)$ is equal to, for a resonance pulse width, $(8l/\pi^2)(c\beta_p/C_p)_1 P_0$ (assuming a negligible value of $\tan\phi$, which is valid for the liquid–air system). With a 1-cm length of water in the tube being the dielectric liquid and with air of 1 atm at the open end, a peak SAR of $P_0 = 24$ kW/gm results in a pressure wave of amplitude equal to 1.82×10^5 dynes/cm^2 in the dielectric liquid, which is 18 percent of the initial equilibrium pressure (1 atm). Since, in making the linear approximation, the factor $\beta_p/(\rho C_p)$ in the right-hand side of (38) has been assumed constant, an 18 percent change in the pressure produces approximately equal percentage change in $1/\rho$, and therefore produces at least so much error by making the linear approximation; the error is indeed higher if the variations of β_p and C_p with respect to pressure are also accounted for. Furthermore, in making the above estimate, we have only included the fundamental frequency of the pressure wave; the total pressure wave could be much higher. Taking 10-percent error as the maximum tolerable level, this implies that the peak SAR of 13 kW/gm is about the upper limit for the validity of the first-order approximation, beyond which a nonlinear approach must be taken.

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