Piezoelectric Rayleigh Wave Excitation by Bulk Wave Scattering

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Abstract—The excitation of a Rayleigh surface wave on a piezoelectric crystal by the scattering of bulk waves from strips deposited on the crystal's surface is considered. Using small scatterer approximations to obtain a lower limit for the strength of excitation that can be obtained by the scattering method, the Rayleigh wave excitation resulting from the mass of the deposited strips is determined. Numerical calculations for CdS, LiNbO₃, and isotropic solids indicate that this method of excitation is feasible. Simple, small scatterer estimates for the order of magnitude of the excitation resulting from the conductivity of strips deposited on CdS and LiNbO₃ suggest that the Rayleigh wave can be more strongly excited through the mass of the strips than through their conductivity.

In order to evaluate the Rayleigh wave excitation, an expression has been derived that describes the excitation of elastic surface waves, guided by configurations composed of infinite, parallel layers of arbitrary piezoelectric and elastic materials, when the source of these waves are prescribed, two-dimensional force and current distributions. This expression for the surface wave excitation coefficient depends only on the source and the properties of the surface wave, and is expected to prove useful in rating a wide variety of exciting structures. For the scattering method of excitation considered here, the effect of the incident bulk wave on the strips can be represented by approximately equivalent sources, which excite the surface wave.

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I. INTRODUCTION

Devices whose operation involves the propagation of elastic surface waves, e.g., Rayleigh wave delay lines, are receiving increased attention because of their possible application in miniaturized circuits. Of particular interest for use in these devices are guiding configurations for the surface waves that contain piezoelectric crystals, since direct electrical excitation of the surface waves is then possible. One aspect of the design of such devices is the selection of an effective means for exciting the surface waves. In making this selection, it is important to be able to estimate the surface wave excitation coefficients for various types of exciting structures and their dependence on the parameters of the structure. However, analytic expressions approximating the surface wave excitation coefficients have been obtained only for a few simple structures used to excite Rayleigh waves [1]–[3].

In Section II, a simple expression is derived that describes the excitation of elastic surface waves, guided by configurations composed of infinite, parallel layers of arbitrary piezoelectric and elastic materials, when the sources of these waves are prescribed, two-dimensional force and current distributions. The plane-stratified surface wave configurations treated here are those composed of infinite, homogeneous, parallel slabs or layers of lossless piezoelectric crystals and elastic materials, which can be perfectly conducting or insulating. The surface waves are assumed to be excited by
prescribed, harmonic volume forces and currents that are uniform, to within a linear phase progression, along a rectangular coordinate parallel to the layers. Using the complete dynamic equations (elastic and electromagnetic equations, including piezoelectric coupling) the excitation coefficient of a surface wave excited by such sources is found, without approximation, as an inner product of the source distribution with the field components of the surface wave, normalized to the power carried by the surface wave. The method used to obtain this form for the excitation coefficient is similar to that employed by Collin [4] for electromagnetic surface waves on cylindrical structures.

For many methods of surface wave excitation, input or terminal conditions such as terminal voltage or current or incident fields are specified rather than volume forces and currents, and the surface wave excitation coefficient cannot be computed directly from the expression found in Section II. However, exciting structures with specified input or terminal conditions can often be approximately represented by two-dimensional source distributions, and the results of Section II permit a simple evaluation of the surface wave excitation coefficients. Moreover, the inner product form for the excitation coefficient should clarify the choice of parameters for most effective excitation.

Such approximately equivalent sources are used in Section III to determine the degree of excitation of Rayleigh surface waves generated when bulk waves are scattered by strips deposited on the surface of a solid. In connection with this method of excitation, it is first argued that if the surface is unperturbed, a plane wave propagating in the solid and incident on the surface will not excite the Rayleigh wave, even if the wavenumber parallel to the surface of the plane wave is equal to that of the Rayleigh wave, as can happen in anisotropic crystals.

The bulk wave scattering method for exciting Rayleigh waves is of interest because of its possible application for high frequencies and high powers. Using ion or electron beams to etch strips out of a crystal’s surface or out of a layer of previously deposited material, it should be possible to surpass the high frequency limit (about 1 GHz) that light places on present types of exciting structures, whose fabrication requires the depositing of interdigital strips by photoetching on piezoelectric crystals. Also, since the initial electromagnetic conversion is accomplished by a bulk wave transducer, the scattering method of excitation overcomes the amplitude limit set by arcing in the present interdigital structures.

With the help of small obstacle approximations, the effect of the incident bulk wave on the strips is represented in terms of approximately equivalent sources. Using these sources, expressions have been obtained for the Rayleigh wave excitation resulting from the mass and from the conductivity of deposited metal strips. Numerical calculations for the excitation, resulting from the mass of the strips, of Rayleigh waves on the basal plane of CdS, on Z-cut LiNbO₃, and on isotropic solids indicates that the scattering method of excitation is feasible. Since the calculations are based on small obstacle approximations, in which mass loading is assumed light, they should be viewed only as establishing a lower limit for the obtainable efficiency. Evaluation of the efficiency that can be obtained with greater mass loadings requires the use of higher order approximations. Simple estimates for conducting strips on CdS and LiNbO₃ further indicate that the Rayleigh wave can be more strongly excited through the mass of the strips than through their conductivity.

II. DETERMINATION OF THE SURFACE WAVE EXCITATION COEFFICIENT

A lossless, plane-stratified, piezoelectric configuration capable of guiding elastic surface waves is assumed to have as its direction of stratification the x₁ axis of a rectangular coordinate system, so that the configuration is invariant under translations along x₁ and x₂. Fig. 1 depicts an example of such a configuration composed of infinite, lossless slabs of metal, glass, and a piezoelectric crystal. The fields in this configuration are excited by a two-dimensional distribution of harmonic sources having time dependence eᵢ맞. The source distribution is composed of a volume force \( \mathbf{F}(x_2, x_3)e^{-\beta x_1} \), an electric current \( j(x_2, x_3)e^{-\beta x_1} \), and, for problems involving the use of equivalent sources, a magnetic current \( \mathbf{M}(x_2, x_3)e^{-\beta x_1} \), where the phase constant \( \beta \) along x₁ is assumed real.

The sources are assumed to be confined to some finite range \( \beta_0 < x_2 < \alpha_0 \) about \( x_2 = 0 \).

With the time dependence \( e^{\text{i}\omega t} \), the time-averaged stress and strain fields, total charge, and particle velocity field (\( j \omega \times \mathbf{E} \)) radiated by the source. These fields may exist in any material layer or slab of the configuration, and satisfy the elastic equations

\[
\frac{\partial}{\partial x_1} T_\alpha = j \omega p V_j - F_\alpha e^{-\beta x_1},
\]

\[
\frac{1}{2} \left( \frac{\partial}{\partial x_1} V_j + \frac{\partial}{\partial x_1} V_i \right) = j \omega S_{ij},
\]

where \( \rho \) is the mass density of the layer, and the summation convention for repeated indices is used. In vacuum regions, \( T_{ij} \) vanishes and \( S_{ij} \) and \( V_i \) need not be considered. Within insulating layers, the total radiated electric field \( E(r) \), displacement field \( D(r) \), and magnetic intensity field \( H(r) \), with time dependence \( e^{\text{i}\omega t} \), satisfy the Maxwell equations

\[
\nabla \times \mathbf{H} = j \omega \mathbf{D} + \mathbf{J} e^{-\beta x_1},
\]

\[
\nabla \times \mathbf{E} = -j \omega \mu_0 \mathbf{H} - \mathbf{M} e^{-\beta x_1}.
\]

The permeability tensor \( \mu \) is used in (2) to allow for layers composed of magnetically anisotropic material. In perfectly conducting layers \( \mathbf{E}, \mathbf{D}, \) and \( \mathbf{H} \) vanish.

In the piezoelectric crystal layers the constitutive relations take the form

\[
S_{ij} = s_{ijkl} T_{kl} + d_{ijk} E_k
\]

\[
D_i = d_{ijk} T_{jk} + \varepsilon_{ij} E_j,
\]
where the $s_{ijkl}$ are the crystal compliances for zero electric field, the $d_{ijkl}$ are the piezoelectric moduli, and the $\epsilon_{ij}$ are the dielectric constants for zero stress. For nonpiezoelectric layers, the constitutive relations have the form (3) with $d_{ijkl}=0$. At the interfaces between the different layers, the following boundary conditions hold:

- $T_{ij}$ continuous
- $V_j$ continuous unless $T_{ij}=0$ on the interface
- $E_1$ and $E_2$ continuous
- $H_1$ and $H_2$ continuous unless $E_1=E_2=0$ on the interface.

Finally, since all sources have $x_1$ dependence $e^{-j\kappa x_1}$, and since the plane-stratified configuration is uniform along $x_1$, the radiated fields will also have this $x_1$ dependence.

Contributions to the total far fields radiated by the sources come from surface waves guided by the stratification. These surface waves decay exponentially for $|x_1| \to \infty$ and propagate along the planes of stratification as $e^{-j(x_1+\kappa x)}$, where for each index $p$, $\eta_p$ is the real wavenumber of the corresponding surface wave. For each $p$, $\eta_p$ is the solution of a resonance equation depending on $\xi$. Thus, $\eta_p$ is a function of $\xi$ and is, in general, real only for a certain range of $\xi$. It is assumed here that $\xi$ is such that at least some of the $\eta_p$’s are real, therefore permitting far field contributions to arise from propagating surface waves. For the purposes of this analysis, all surface waves propagate as $e^{-j\eta_p x_3}$, the direction of phase progression along $x_3$ of the $p$th surface wave being determined by the sign of $\eta_p$.

In addition to the true propagating surface waves, which propagate unattenuated along the stratification and decay for $|x_1| \to \infty$, in restricted regions of space the source may excite leaky surface waves, which decay along the stratification and grow exponentially for large $|x_1|$. Although the leaky waves in certain cases may have a low attenuation [5] and thus give a significant contribution to the total radiated fields over a substantial region of space, only the excitation of the true propagating surface waves is considered here. The restriction to surface waves that decay for $|x_1| \to \infty$ also excludes from consideration any direction of propagation, i.e., value of $\xi$, at which a surface wave degenerates into a propagating bulk wave of a semi-infinite substrate [5].

The field components $T_p(r), V_p(r), E_p(r), H_p(r)$ of the $p$th surface wave, where $T_p$ is the dyadic representation of the stress tensor, satisfy the homogeneous dynamic equations, i.e., (1) and (2) with $F=J=M=0$, the constitutive relations (3), and the boundary conditions (4). These field components can be written in the form

$$
\begin{bmatrix}
\mathbf{T}_p(r) \\
\mathbf{V}_p(r) \\
\mathbf{E}_p(r) \\
\mathbf{H}_p(r)
\end{bmatrix} = \begin{bmatrix}
\mathbf{S}_{p}(x_3) \\
\mathbf{D}_{p}(x_3) \\
\mathbf{P}_{p}(x_3) \\
\mathbf{J}_{p}(x_3)
\end{bmatrix} e^{-j(E_{00}+\eta_p x_3)},
$$

(5)

where the script letters represent the $x_3$ dependent “polarization tensors” of the surface wave fields, i.e., the relative amplitudes of the various field components of the surface wave as a function of $x_3$. The normalization of the surface wave is arbitrary. In order to determine the excitation coefficient of the surface wave (5) due to a source, it will first be shown that any two such surface waves having different real wavenumbers $\eta_p$ carry power independently in the $x_3$ direction.

Let $\eta_p$ and $\eta_q$ be the $x_3$ wavenumbers of two surface waves and consider the divergence of the vector field

$$
\mathbf{n}^{\text{ss}} = - \mathbf{T}_p \cdot \mathbf{V}_q - \mathbf{T}_q \cdot \mathbf{V}_p + \mathbf{E}_p \times \mathbf{H}_q + \mathbf{E}_q \times \mathbf{H}_p.
$$

(6)

By direct expansion and using (1) and (2), for $F=J=M=0$, and (3), it can be shown that $\nabla \cdot \mathbf{n}^{\text{ss}} = 0$ in the absence of loss. The lossless requirement is imposed so that: a) the compliance tensor has the symmetry $S_{ijkl} = S_{klij}$; b) the piezoelectric moduli have the symmetry $d_{ijkl} = d_{klij}$; c) the dielectric tensor has the symmetry $\epsilon_{ij} = \epsilon_{ji}$; and d) the permeability tensor has the symmetry $\mu_{ij} = \mu_{ji}$. Since $\nabla \cdot \mathbf{n}^{\text{ss}} = 0$ and since (4) implies that the $x_3$ component of $\mathbf{n}^{\text{ss}}$ is continuous across the interfaces between the different layers, Gauss’s theorem requires that the surface integral of $\mathbf{n}^{\text{ss}}$ over any closed surface vanish. In particular, the surface integral will vanish for the infinitely long cylinder of Fig. 2. This cylinder is parallel to the $x_2$ axis and has rectangular cross section with sides parallel to the $x_1$ and $x_3$ axes.

Because the surface wave fields vanish for $|x_3| \to \infty$, no contribution to the surface integral arises from the ends of the cylinder at $|x_3| \to \infty$. Furthermore, since the fields of both surface waves vary as $e^{-j\kappa x_3}$, $\mathbf{n}^{\text{ss}}$ is independent of $x_1$ and the contributions to the surface integral from the two sides of the cylinder at $x_1=0$, I cancel. Thus, defining

$$
\mathbf{N}^{\text{ss}} = \frac{1}{4} \int_{-\infty}^{\infty} \left[ - \mathbf{E}^* \cdot \mathbf{V}_p(x_3) - \mathbf{E} \cdot \mathbf{V}^*_p(x_3) \right] dx_3
$$

(7)

as the end correction to the surface integral, we obtain

$$
\mathbf{N}^{\text{ss}} = \frac{1}{4} \int_{-\infty}^{\infty} \left[ - \mathbf{E}^* \cdot \mathbf{V}_p(x_3) - \mathbf{E} \cdot \mathbf{V}^*_p(x_3) + \mathbf{E} \times \mathbf{J}_p(x_3) + \mathbf{E}^* \times \mathbf{J}^*_p(x_3) \right] dx_3.
$$
and with the help of (5), the integration over the sides of the cylinder at \( x_3 = \alpha, \beta \) yields

\[ 0 = 4 \left[ e^{-j(q_\eta - \eta_\pi)\alpha} - e^{-j(q_\eta - \eta_\pi)\beta} \right] N_{2pq}, \tag{8} \]

where \( N_{2pq} \) is the \( x_3 \) component of \( N_{pq} \) of (7).

If \( \eta_{pq} \neq \eta_{q}, \) (8) can hold, for arbitrary \( \alpha \) and \( \beta, \) only if \( N_{2pq} = 0. \) This fact implies that the total real power carried in the \( x_2 \) direction by several surface waves, having distinct values of \( \eta_{pq} \) is the sum of the real powers carried by the individual surface waves. In what follows, it will be assumed that the independent surface waves have distinct values of \( \eta_{pq} \) or can be defined such that \( N_{2pq} = 0. \)

Consider now the vector

\[ n = - \mathbf{T}_p \times \mathbf{V} - \mathbf{E}_p \times \mathbf{H}_p \times H + E \times H_p^*, \tag{9} \]

where \( \mathbf{T}, \mathbf{V}, \mathbf{E}, \) and \( \mathbf{H} \) are the total fields radiated by the sources in (1) and (2) and \( \mathbf{T}_p, \mathbf{V}_p, \mathbf{E}_p, \) and \( \mathbf{H}_p \) are the surface wave fields of (5) for any particular value of \( p. \) Since all the fields in (9) have the \( x_3 \) dependence \( e^{-j\eta_\pi x_3}, \) it is seen that \( n \) is independent of \( x_1. \) By direct expansion and using (1), (2), (3), and (5), it can be shown that

\[ \mathbf{V} \cdot n = (\mathbf{U}_p^r \cdot \mathbf{F} - \mathbf{E}_p^r \cdot \mathbf{J} - \mathbf{I}_p^r \cdot \mathbf{M}) e^{j\eta_\pi x_3}. \tag{10} \]

Assuming the sources to be entirely confined to the region \( \beta < x_3 < \alpha_0; \) both sides of (10) are now integrat ed over the volume of the cylinder of Fig. 2 with \( \alpha > \alpha_0 \) and \( \beta < \beta_0. \) Because both sets of fields in (9) satisfy the boundary conditions (4), \( n \) is continuous across the interfaces between the layers and hence the volume integral of \( \mathbf{V} \cdot n \) can be converted into a surface integral. As before, the contributions to the surface integral from the ends of the cylinder at \( |x_3| = \infty \) vanish, while those from the sides at \( x_1 = 0, 1 \) cancel. Thus, the integration of (10) leads to

\[ \int_{-\infty}^{\infty} n_2(\alpha, x_3) dx_3 - \int_{-\infty}^{\infty} n_2(\beta, x_3) dx_3 = U_p, \tag{11} \]

where

\[ U_p = \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} (\mathbf{U}_p^r \cdot \mathbf{F} - \mathbf{E}_p^r \cdot \mathbf{J} - \mathbf{I}_p^r \cdot \mathbf{M}) e^{j\eta_\pi x_3} dx_2 dx_3. \tag{12} \]

In order to obtain the surface wave excitation coefficients from (11), it is necessary to consider a Fourier integral representation, over the transform variable \( \eta, \) of the radiated fields. The Fourier integral representation of interest here is the one whose integrand is \( e^{-j\eta x_2} \) times the \( x_3 \) dependent transform of the fields. By appropriately deforming the path of integration for the Fourier integral representation into the complex \( \eta \) plane, for all \( x_3 > 0 \) the total radiated fields can be expressed as a sum over real, surface-wave pole residues (discrete spectrum) plus the contributions arising from the integration over the deformed path (continuous spectrum). Similarly, by a different deformation of the path of integration, for all \( x_3 < 0 \) the total radiated fields can be written as a sum of real, surface-wave pole residues and the contribution from the deformed path.

Because of the radiation condition for \( |x_3| \to \infty, \) the real, surface-wave poles contributing for \( x_3 > 0 \) are those for which the corresponding surface waves carry real power in the positive \( x_3 \) direction, i.e., those for which \( N_{2pq} > 0, \) where \( N_{2pq} \) is the \( x_3 \) component of the vector defined in (7) with \( q = p \) and represents the total, real, power flow in the \( x_2 \) direction, per unit length along \( x_3, \) of the \( p \) surface wave defined in (5). For \( x_3 < 0, \) the contributing, real, surface-wave poles are those for which \( N_{2pq} < 0. \) Denoting those real, surface-wave poles that contribute for \( x_3 > 0 \) by \( p^+ \) and those contributing for \( x_3 < 0 \) by \( p^-, \) and letting the subscripts \( c^+ \) and \( c^- \) refer to the contributions to the fields from the continuous spectra for \( x_3 > 0 \) and \( x_3 < 0, \) respectively, then for \( x_3 \gtrless 0 \)

\[ \left[ \mathbf{T}; \mathbf{V}; \mathbf{E}; \mathbf{H} \right] = \left[ \mathbf{T}_{c^\pm}; \mathbf{V}_{c^\pm}; \mathbf{E}_{c^\pm}; \mathbf{H}_{c^\pm} \right] + \sum_{p=0} a_{p^\pm} \left[ \mathbf{T}_{p^\pm}; \mathbf{V}_{p^\pm}; \mathbf{E}_{p^\pm}; \mathbf{H}_{p^\pm} \right]. \tag{13} \]

In (13), the \( a_{p^\pm} \) s represent the surface wave excitation coefficients that are to be determined and \( \mathbf{T}_{p^\pm}, \mathbf{E}_{p^\pm}, \mathbf{H}_{p^\pm}, \mathbf{V}_{p^\pm} \), etc., are independent, the second and third terms in (14)

\[ 4 a_p | N_{2pq} | + e^{i\eta_{pq} N_{2pq}(Z)} - e^{i\eta_{pq} N_{2pq}(\beta)} = U_p, \tag{14} \]

where \( a_p \) is one of the unknown excitation coefficients in (13) and \( N_{2pq}(x_3) \) is the \( x_3 \) component of

\[ N_{2pq} = \int_{-\infty}^{\infty} \left( -\mathbf{3}_{p^\pm} \cdot \mathbf{V}_{c^\pm} - \mathbf{T}_{c^\pm} \cdot \mathbf{J}_{c^\pm} \right) e^{j\eta_{pq} x_3} dx_3. \tag{15} \]

Note that \( \mathbf{T}_{c^\pm}, \mathbf{E}_{c^\pm}, \mathbf{H}_{c^\pm}, \mathbf{V}_{c^\pm} \) vary as \( e^{-j\eta_{pq} x_3} \) so that \( N_{2pq} \) is a function of \( x_2 \) only. Since \( a_p, N_{2pq}, \) and \( U_p \) are constants and \( \alpha > \alpha_0 \) and \( \beta < \beta_0 \) are independent, the second and third terms in (14) must independently be constants in order that (14) be satisfied. But \( \mathbf{T}_{c^\pm}, \mathbf{E}_{c^\pm}, \mathbf{H}_{c^\pm}, \mathbf{V}_{c^\pm} \) at \( x_3 = \alpha, \) and hence \( N_{2pq}(\alpha) \) in (14), contain a continuous distribution of \( \eta \) dependences of the form \( e^{-j\eta_{pq} x_3}, \) with \( \eta \neq \eta_{pq}, \) and therefore cannot annul the factor \( e^{i\eta_{pq} x_3} \) in (14) for all \( \alpha > \alpha_0. \) Thus \( N_{2pq}(\alpha) \) must vanish for all \( \alpha > \alpha_0. \) Similarly, \( N_{2pq}(\beta) \) must vanish for all \( \beta < \beta_0 \) and hence from (14)

\[ a_p = U_p / 4 \left| N_{2pq} \right|, \tag{16} \]

where \( N_{2pq} \) is the \( x_3 \) component of the total, real Poynting vector of the surface wave (5), as defined in (7) for \( q = p. \)

Expression (16) provides a simple means for evaluating the excitation coefficients of the surface waves (5) propagating in plane-stratified, lossless, piezoelectric configurations when the excitation is by a two-dimensional source distribution having the \( x_1 \) dependence \( e^{-j\eta_{pq} x_1}. \) From (7), (12), and (16), it is seen that the actual radiated surface wave fields in (13),
e.g., the stress field \( \mathbf{\sigma}_e \mathbf{J}_e \), are independent of the choice of normalization used in defining the surface wave "polarization tensors" in (5). Also, from the form (13) for the radiated fields, and in view of the power orthogonality of the surface waves, the radiated power per unit length in the \( x_1 \) direction, carried in the \( x_2 \) direction by the \( p \)th surface wave, is seen with the help of (5), (7), and (16) to be

\[
N_p = \left| U_p \right|^2 / 10 N_i^{pp}, \quad (17)
\]

In the next section, this simple expression will be used to determine the efficiency of Rayleigh wave excitation by bulk wave scattering.

III. Rayleigh Wave Excitation by Bulk Wave Scattering

Many methods of surface wave excitation employ structures with specified input or terminal conditions, such as terminal voltage or current or incident fields, rather than specified source distributions. For these methods, the surface wave excitation coefficient cannot be obtained directly from (16). However, it is frequently possible to approximately represent the exciting structure by easily calculated sources. In what follows, such approximately equivalent sources are used to determine the efficiency with which the Rayleigh surface waves on the basal plane of CdS, on Z-cut LiNbO₃, and on isotropic solids are excited when beams propagating in the bulk of the crystal are scattered by parallel strips deposited on the solids' surfaces. In this context, it is first shown that a plane wave incident on an unperturbed surface will not excite the Rayleigh wave, even if the wavenumber parallel to the surface of the plane wave is equal to that of the Rayleigh wave, as can happen in anisotropic crystals.

The method of Rayleigh wave excitation employing bulk wave scattering is depicted in Fig. 3, in which a two-dimensional beam is incident on infinitely long, metallic strips deposited on the crystal's surface parallel to the \( x_1 \) axis. The beam is assumed to be of infinite extent along \( x_1 \) and to have central fields that are those of a plane wave having no phase variation along \( x_1 \). To within the approximations used here, the surface wave excitation consists of a mechanical contribution resulting from the mass of the strips and, in the case of CdS and LiNbO₃, an electrical contribution resulting from the conductivity of the strips.

The approximations used here in determining the mechanical contribution to the surface wave excitation are applicable for sufficiently small \( t \) in Fig. 3 and for all angles of incidence of the beam. Numerical computations, used to indicate a lower limit for the efficiency of this method of excitation, are performed only for the normally incident, longitudinal, plane wave beams in CdS and LiNbO₃ and as a function of incidence angle for longitudinal and shear beams in isotropic solids whose two Lamé coefficients are equal. Substantial increases in efficiency can be expected from the use of larger mass loadings, the analysis of which requires higher order approximations than those used here.

Approximations used to calculate the electrical contribution to the surface wave excitation on the basal plane of CdS and on Z-cut LiNbO₃ require \( t \) and \( d \) in Fig. 3 to be sufficiently small, and the incident plane wave to be such that the total electric field has no \( x_1 \) component. For the normally incident, longitudinal plane wave, these approximations indicate that there is no electrical contribution to the excitation. Moreover, simple estimates for the order of magnitude of the electrical contribution to the excitation based on these approximations suggest that for all other angles of incidence, the Rayleigh wave can be more strongly excited through the mechanical contribution.

A. On the Possibility of Rayleigh Wave Excitation at an Unperturbed Interface

With the study of Rayleigh waves on anisotropic crystals, the possibility was recognized that there could exist crystal cuts such that some propagating plane wave in the bulk of the crystal would have a wavenumber parallel to the surface that was identical to that of the Rayleigh wave. Recognizing this, one is led to speculate whether such a plane wave, when incident on an unperturbed surface, could excite the Rayleigh wave. By rephrasing this speculation in a physically meaningful way in terms of the source of the plane wave, it is easily seen from the surface wave excitation coefficient that the plane wave excitation is not possible. Since the proof relies on the surface wave excitation coefficient, it applies only to Rayleigh waves whose fields decay away from the surface.

Assume a distribution of sources is located in a plane parallel to and at a distance \( h \) below the surface of a semi-infinite crystal. By suitably choosing these sources, they will radiate a beam towards the surface whose central fields are those of any desired plane wave. Taking the source distribution in the plane to be sufficiently wide, the amplitude of the central, propagating, plane wave, beam fields that are incident on the interface can be made to vary only slightly as \( h \) increases over many elastic wave lengths. If the central, plane wave fields of the beam can be viewed as exciting the Rayleigh wave on the unperturbed surface, then that portion of the excitation coefficient due to the plane wave will vary only slightly as \( h \) increases by many wavelengths. However, using (16) to express the excitation coefficient directly in terms of the source, it is seen that as \( h \) increases by many
wavelengths, the excitation coefficient will approach zero, because of the exponential dependence of the Rayleigh wave fields on depth. Thus, a propagating plane wave incident on an unperturbed surface cannot excite the Rayleigh wave, no matter what its wavenumber parallel to the surface.

B. Rayleigh Wave Excitation Resulting from Mass Loading by Strips

The approximations made in determining the surface wave excitation due to the mass loading of the strips are based on the assumption that the mass deposited per unit area, are sufficiently small so that the motion of the surface is essentially the same as for the unloaded surface. Under this assumption and neglecting any rigidity within the strips, there exists a surface force density under the strips that is the negative of the deposited mass per unit area times the acceleration of the unloaded surface. Thus, in (12)

$$F(x_2, x_3) = -j\omega V(x_3) \sigma(x_2) \delta(x_3),$$

where \(V(x_3)\) is the velocity of the unloaded surface, \(\sigma(x_2)\) is the mass per unit area deposited on the surface, and \(\delta(x_3)\) is the Dirac delta function.

Under the approximations used here, the excitation due to all the strips in Fig. 3 is the sum of the excitations due to the individual strips. Thus, for identical strips, as in Fig. 3, the excitation due to all strips can easily be found from that due to the strip centered at \(x_2 = 0\).

Assuming the strips to be composed of material of density \(\rho_0\), in (18) \(\sigma(x_2) = \rho_0 t\) for \(-d/2 < x_2 < d/2\) and vanishes outside this region, when considering the single strip at \(x_2 = 0\). If the incident plane wave has wavenumber \(k_1\) in the \(x_2\) direction, then in (18) \(V(x_3) = V_0 e^{-j k_1 x_3}\). Using (18), the mechanical contribution \(U^m_p\) to the quantity \(U_p\) in (12) due to the single strip at \(x_2 = 0\) is therefore given by

$$U^m_p = -j2\omega \rho_0 d V_0 \cdot U^m_p(0) \sin \left(\frac{(\nu_p - \eta_1) d/2}{(\nu_p - \eta_2)}\right).$$

From (19) it is seen that the surface wave will be most strongly excited when the strip width \(d\) is such that \((\nu_p - \eta_1) d/2\) is an odd multiple of \(\pi/2\). For a normally incident plane wave \(\eta_1 = 0\), so that the minimum value of \(d\) that gives optimum excitation in this case is \(\lambda_p/2\), where \(\lambda_p = 2\pi/\nu_p\) is the wavelength of the surface wave.

As previously discussed, (18) and hence (19) are only valid for \(t\) sufficiently small. For too large values of \(t\), the mass loading of the surface significantly reduces the surface motion so that the actual force density is less than that given in (18). In order to estimate the range of \(t\) for which (18) and (19) are valid, first consider the case of normal incidence. In this case, the amplitude of the velocity of the unloaded surface will be \(2A_1\), where \(A_1\) is the amplitude of the particle velocity of the incident plane wave. If an additional layer of the crystal material of thickness \(t_0\) is deposited on the surface, the particle velocity at \(x_3 = 0\) will have amplitude \(A_1(1 + e^{-k_1 t_0})\), where \(k_1\) is the wavenumber along \(x_3\) of the incident plane wave. For small \(t_0\), the change relative to the velocity amplitude of the unloaded surface in the particle velocity at \(x_3 = 0\) resulting from the addition of crystal material is \((-j k_1 t_0)\). Since the additional crystal layer acts as a mass loading for small \(t_0\), the magnitude of the relative change in the surface velocity due to loading the surface with a uniform layer of thickness \(t\) and mass density \(\rho_0\) is \(k_1 \rho_0 / \rho\), where \(\rho\) is the mass density of the crystal. The case of plane waves incident at an angle is complicated by the fact that both shear and longitudinal plane waves are generated, in general, at the interface. However, considerations similar to those above show that the magnitude of the relative change in the surface velocity will be less than \(k_1 \rho_0 / \rho\), where \(k\) is the largest of the wave numbers along \(x_3\) of the plane waves involved.

Tseng and White [6] have shown that a Rayleigh surface wave can propagate on the basal plane of CdS and have computed its wavenumber \(\eta_1\) and the velocity and electric field components of the \(\eta_1\) dependent "polarization tensors." Using their calculated values, in the present notation the \(x_2\) and \(x_3\) components of particle velocity at \(x_3 = 0\) are

$$U_p(x_2) = j1.02 \omega A_1; \quad U_p(x_3) = 1.90 \omega A_1,$$

where \(A_1\) is an arbitrary constant (the \(x_1\) component of the velocity is zero everywhere for the surface wave). Since \(|U_p(x_2)| < |U_p(x_3)|\), it is seen from (19) that the surface wave will be more strongly excited by the normally incident longitudinal wave than by a normally incident shear wave having the same amplitude of particle velocity. For the above reason, only the normally incident, longitudinal wave will be considered.

With the help of the first four equations of Tseng and White [6] and the fact that \(\nabla \cdot D = 0\), it is easily shown that the harmonic, longitudinal wave propagating as \(e^{j \omega \nu_s x_3}\) along the \(C\) axis in CdS, i.e., perpendicular to the basal plane, has for the amplitudes of its nonvanishing field components the following:

$$V_3 = A_1; \quad T_{11} = T_{12} = -A_1 \frac{k_1}{\omega} (c_{13}^B + c_{33}^S / \epsilon_3^S);$$

$$T_{33} = -A_1 \frac{k_1}{\omega} (c_{33}^B + c_{33}^S / \epsilon_3^S) \epsilon_3^S; \quad E_3 = A_1 \frac{k_1}{\omega} (c_{33}^B + c_{33}^S / \epsilon_3^S),$$

where

$$k_1 = \omega \sqrt{\frac{\rho}{c_{33}^B + c_{33}^S / \epsilon_3^S}}.$$

In (21) and (22) the \(c_{ij}^B\) are the elastic stiffness constants for vanishing electric field, the \(\epsilon_3^S\) are the dielectric constants for vanishing strain, and the \(\epsilon_{ij}\) are the piezoelectric constants. When the positive root is taken in (22), the resultant wave propagates in the positive \(x_3\) direction, i.e., is incident
on the interface in Fig. 3. The free surface boundary conditions for the unloaded surface can be satisfied by the addition of a reflected wave, identical to the incident wave (21) except for the choice of the negative root in (22). Thus, the particle velocity $V^0$ at the surface will be along $x_3$ with amplitude $2A_i$.

In order to calculate the power scattered per unit length along $x_3$ into the surface wave from (17), it is necessary to know $N_p$ as defined in (7) for $p = q$. Using the “polarization tensor” components found by Tseng and White [6] and the crystal constants of CdS [7], [8] necessary to determine the remaining components, it can be shown that $N_p = 11.1 \times 10^{-10} |A_i|^2$ (watts/meter). Thus, neglecting the electric contribution and assuming in (19) that $d = \lambda_p / 2$ and that a longitudinal wave, with particle velocity amplitude $A_i$, is normally incident on the interface, from (17) and (20) the power scattered into the surface wave by a single strip is found to be

$$N_p = 0.324 \times 10^{-10} \frac{\omega^2 \mu_{pq}^2 |A_i|^2}{\eta_p^2}.$$  \hspace{1cm} (23)

To set $N_p$ of (23) in perspective, it is compared to the power carried by the incident plane wave through an area of unit length along the strip and width $\lambda_p$ along $x_3$, i.e., to the power $\lambda_p N_i = -\lambda_p T_{33} V_{23}^* + T_{23}$ and $V_{23}^*$ as given in (21). While this choice of normalization for $N_p$ has little meaning for the case of a single strip, it leads to an ideal but approachable scattering efficiency for the case of many strips. Using (21) and (22), $N_i = |A_i|^2 \rho_{po}/2k$ and defining $\text{eff}(1) = N_p / \lambda_p N_i$, it can be shown that

$$\text{eff}(1) = 0.384 \left(\frac{\kappa \rho_{po}}{\rho}\right)^{\frac{1}{3}}.$$  \hspace{1cm} (24)

For comparison, a similar calculation has been carried out for the excitation of the Rayleigh wave propagating at $30^\circ$ to the $X$ direction on Z-cut LiNbO$_3$ by the normally incident longitudinal wave. Using the values of $\eta_p$, $\Omega_p(0)$, and $\lambda_{eff}$ determined by Campbell and Jones [9], [10], and the crystal constants of LiNbO$_3$ necessary to find $V^0$ and $N_i$ for the incident longitudinal wave [11], it has been found that $\text{eff}(1) = 0.395 \left(\kappa_3 \rho_{po}/\rho\right)^{\frac{1}{3}}$. The similarity in $\text{eff}(1)$ for CdS and LiNbO$_3$, despite the dissimilarity of their stiffness tensors, suggests that $\text{eff}(1)$ for normally incident longitudinal waves will be near the value (24) for a wide variety of crystals.

The quantity $\kappa \rho_{po}/\rho$ in (24) has been shown to approximate the relative change in the velocity of the surface due to the mass loading. It is expected that the force (18) will give a reasonably good approximation for the surface wave excitation for mass loadings producing relative surface velocity changes up to about 0.10, which result in values of $\text{eff}(1)$ up to about 0.004. Because of the quadratic dependence of $\text{eff}(1)$ on $t$, values of $\text{eff}(1)$ well above 0.004 should be possible using mass loadings above that for which $\kappa \rho_{po}/\rho = 0.10$.

Consider now a beam that is normally incident on an array of $n$ identical strips, as in Fig. 3, with center-to-center spacing $\lambda_p$. For simplicity, the beam fields are idealized to be those of a longitudinal plane wave over a region, of width $n\lambda_p$ along $x_3$, that is centered on the array of strips, and to be zero outside this region. Since the centers of the strips are $\lambda_p$ apart, the surface wave excitation due to each strip is in phase for the normally incident fields. Thus, neglecting the scattering of the surface wave due to one strip by subsequent strips, the surface wave fields due to $n$ strips will be $n$ times those due to one strip and the power carried by the surface wave will be greater by a factor $n^2$. Since the power incident in this idealized beam is $n\lambda_p N_i$, where $N_i$ is the incident power flux, the efficiency $\text{eff}(n)$ with which $n$ strips convert the incident power into surface wave power is $\text{eff}(n) = n \text{eff}(1)$.

For large enough $n$, i.e., for wide beams, the incident power in that region of a physically obtainable beam in which the fields go continuously to zero can be made small compared to the power in the central region of the beam. Therefore, while $\text{eff}(n)$ was calculated on the basis of an idealized beam whose fields are discontinuous at the beam edges, it does give a reasonable measure of the efficiency that can be obtained for a large number of strips. Note, however, that for too large values of $n$, e.g., for $n$ such that $\text{eff}(n)$ is above 0.10, the dependence of $\text{eff}(n)$ on $n$ will deviate from linearity because the scattering, by all subsequent strips, of the surface wave excited by one strip becomes significant.

From (24), it is seen that efficiencies on the order of 0.10 can be obtained using 25 to 30 strips of width $d = \lambda_p / 2$ and thickness $t$ such that $\kappa \rho_{po}/\rho = 0.10$. As previously stated, substantially higher single strip efficiencies should be possible using values of $t$ beyond that quoted above. This would permit the use of less strips, and consequently increase bandwidth, to obtain the same value of $\text{eff}(n)$. However, the analysis for larger values of $t$ would require consideration of higher order approximations.

Improvement in the efficiency is possible using beams incident at an angle. In particular, the bidirectionality inherent in the excitation by normally incident beams is avoided. For an obliquely incident beam, the strip spacing $D$ in Fig. 3 is chosen such that $\eta_p D = \eta_p D + 2\pi m$ for some integer $m$, which ensures that the Rayleigh waves excited in the plus $x_3$ direction by each strip add in phase. In practice, $D$ and consequently $m$ should be kept small so as to obtain the maximum number of strips per unit of beamwidth, and hence maximum efficiency.

In order to demonstrate the dependence of the efficiency on the angle of incidence, $\text{eff}(1)$ has been calculated for obliquely incident longitudinal and shear waves in isotropic solids. For the reason given above, it is assumed that $m = 1$ and that for each angle of incidence $\theta$ in Fig. 3, the strip spacing $D$ is chosen so as to maintain the relation $\eta_p D - \eta_p D + 2\pi$ against the variation in $\eta_p$ with $\theta$. Similarly, to optimize (19) for each $\theta$, the strip width $d$ is chosen such that $\left(\eta_p - \eta_p D + 2\pi / \eta_p D = \pi\right)$. At each angle, $\text{eff}(1)$ is defined as the ratio of the power carried by the surface wave, per unit length along $x_3$, to the power of the incident plane wave falling on an area
in the $x_2=0$ plane of unit width along $x_1$ and width $D$ along $x_2$. As in the case of unit incidence, the efficiency for $n$ strips is equal to $n \varepsilon_{\text{eff}}(1)$, so long as this value is not too large. The results of the calculations for $\varepsilon_{\text{eff}}(1)$, as depicted in Fig. 4, are based on the Rayleigh wave properties given by Ewing, Jardetzky, and Press [12] for solids having equal Lamé coefficients. For an incident longitudinal wave, the value of $\varepsilon_{\text{eff}}(1)/(k_1 t p_0 / p)^2$, where $k_1$ is the wavenumber of longitudinal plane waves in the medium, is plotted as a function of $\theta$ over the range $0^\circ$ to $90^\circ$. Similarly, $\varepsilon_{\text{eff}}(1)/(k_2 t p_0 / p)^2$ is plotted for incident shear waves polarized such that the particle velocity is in the plane of incidence, where $k_2$ is the wavenumber of shear waves. The narrow peak in this plot at about $36^\circ$ occurs when the reflected longitudinal wave at the unperturbed surface is just beyond cutoff. The overall increase of $\varepsilon_{\text{eff}}(1)$ up to about $\theta=75^\circ$ is primarily due to the factor $(\eta_0 - \eta_1)$ in (19), while the sudden decrease of $\varepsilon_{\text{eff}}(1)$ to zero at $\theta\approx90^\circ$ occurs because of the vanishing of the particle velocity at the unperturbed surface for grazing incidence.

C. Rayleigh Wave Excitation Resulting from Strip Conductivity

If the electrical properties of the deposited strips of Fig. 3 differ from those of free space, scattering into the Rayleigh surface wave will result from this perturbation when the incident plane wave produces an electric field above the unperturbed surface. For simplicity in estimating the degree of this scattering, the strips are assumed to be perfectly conducting. Considering each strip separately, the incident plane waves may be viewed as inducing electric currents on the surface of these conducting strips, which in turn excite the surface wave, as described by (12). Since the Rayleigh surface wave propagating in the $x_2$ direction on the basal plane of CdS has only $x_1$ and $x_2$ components of electric field, an incident plane wave with no phase variation along $x_1$ will excite this surface wave only through the transverse currents induced on the strips.

In what follows, currents that are approximately equivalent to the induced currents are obtained using a “small obstacle” or “quasi-static” approach [13] that is valid when the incident and scattered electric fields have no components along the strips and the cross-sectional dimensions of the strips are small compared to the elastic wavelengths. For the crystal cut used here, and for incident plane waves having no phase variation and no particle velocity component along $x_3$, it can be argued from the dynamic equations (1) and (2), the constitutive relations (3) for CdS, and the boundary conditions, that the total electric field satisfies the above requirements.

As in the case of the mechanical contribution, the electrical contribution to the surface wave excitation is calculated for a single strip centered at $x_2=0$. To within the approximations used here, the scattered fields due to several strips are the sum of the fields scattered by the individual strips. Also, the effect of the mass of the strip on the induced currents is neglected.

Assuming the transverse dimensions $t$ and $d$ of the strip to be small compared to elastic wavelengths, to first order the electrical contribution $U_\ast$ to the quantity $U_\ast$ in (12) for the single strip centered at $x_2=0$ is

$$U_\ast = -J^0 \cdot \varepsilon_\ast(0^\circ),$$

where the transverse vector $J^0$ represents the total induced current per unit length along $x_1$. When performing the integration in (12) to obtain $U_\ast$, no ambiguity as to the value of $J^0 \cdot \varepsilon_\ast$ on the $x_2=0$ face of the strip arises, since the $x_1$ component of $\varepsilon_\ast$ is continuous at $x_1=0$ and the $x_2$ component of the surface current vanishes on this face of the strip. Associated with the induced surface currents, there is a surface charge density on the conducting strips. If $P^0$ is the dipole moment per unit length along $x_1$ of the induced surface charge density, then $J^0 = j_0 P^0$.

For $t$ and $d$ small compared to elastic wavelengths, at any instant the scattered electric field close to the strip is approximately equal to the electrostatic field that would be produced by the instantaneous surface charge. In other words, if $p(x_3, x_2)$ is the amplitude of the induced charge density, then the scattered electric field near the strip is approximately $e^{j\omega t} E(x_2, x_3)$, where $E^\circ$ is defined by the electrostatic equations $\nabla \times E^\circ = 0$ and $\nabla \cdot D^\circ = p$ and the constitutive relations $D^\circ = \epsilon_0 E^\circ$ for $x_3>0$ and $D^\circ = \epsilon_1 E^\circ$ for $x_3<0$, the last of which is obtained from (3) by neglecting the term $d_{ijkl} T_{jk}$. Therefore, if $e^{j\omega t} E(x_2, x_3)$ is the field that would exist near the interface if the strip were not present, then $E^\circ + E^\ast$ will approximately satisfy the boundary condition of vanishing tangential electric field on the strip. Conversely, a static surface charge distribution on the strip, whose field, together with $E^\circ$, satisfies the boundary condition on the strip, should be approximately equal to the amplitude of the surface charge induced in the dynamic case. In particular, the dipole moment $P^\circ$ of the static distribution will be approximately equal to $P^0$.

The static charge distribution is generally difficult to find. However, only the dipole moment $P^\circ$ of the distribution is of interest here, and for $t<<d$ it can readily be estimated.
Recalling that \( d \) is small compared to elastic wavelengths, the \( x_3 \) component of the displacement field on the unperturbed surface will be uniform in the vicinity of the strip. Neglecting edge effects for \( t \ll d \), this component of the displacement field will induce a surface charge \( \pm \varepsilon_0 E_0 \delta(0, 0^+) \) on the top and bottom faces of the strip centered at \( x_3 = 0 \). Thus, the dipole moment per unit length along \( x_1 \) of the surface charge induced by the \( x_3 \) component of the displacement field is approximately given by \( P_s^* = d \varepsilon_0 E_0 \delta(0, 0^+) \).

Again, for \( t \) and \( d \) small compared to elastic wavelengths, the \( x_2 \) component of the electric field on the unperturbed crystal surface will be uniform in the vicinity of the strip. For \( t \ll d \), the dipole moment of the static charge distribution induced by the \( x_3 \) component of the electric field can be estimated by assuming the strip to be of zero thickness and neglecting edge effects for \( x_2 \) below the strip and \( x_2 \) above the unloaded crystal surface. However, the order of magnitude of the electrical contribution occurring at other angles of incidence can easily be estimated by using the surface wave field components themselves for \( E^0 \) and \( V^0 \) in (27).

Using the values of \( \varepsilon_p \) and \( \Psi_p \) found by Tseng and White [6] for \( E^0 \) and \( V^0 \), taking for \( \rho_0 \) the mass density of CdS and assuming \( t \ll d \), then \( U_{p*}/U_{p\text{m}} \sim 0.01d/t \), where the term \( E_3^0 \delta_{p3}^* \) in (27) is neglected for \( t \ll d \) because \( \delta_{p3}^* \sim -j \varepsilon_{p2} \). Recalling that the derivation of the mechanical contribution to the surface wave excitation is valid for \( k_d \rho_0 / \rho_0 \) less than about 0.1, the maximum value of \( t \) for which (27) approximates \( U_{p*}/U_{p\text{m}} \sim 0.1/\eta_p \). In the derivation of the electrical contribution, the assumption that \( d \) be small compared to elastic wavelengths requires that \( d \) be no greater than \( 1/\eta_p \). With these maximum values \( d/t \sim 10 \) so that \( U_{p*}/U_{p\text{m}} \sim 0.2d/t \), so that the mechanical contribution to the Rayleigh wave excitation can easily be made substantially stronger than that due to the electrical effect. A similar computation for the Rayleigh wave propagating at 30° to the \( X \) direction on Z-cut LiNbO\(_3\), as determined from the properties of the Rayleigh wave [10] and the crystal constants [11], yields the value \( U_{p*}/U_{p\text{m}} \sim 0.02d/t \), so that the mechanical contribution to the Rayleigh wave excitation can easily be made substantially stronger than the electrical contribution.

### IV. CONCLUSIONS

The Rayleigh wave excitation resulting from bulk wave scattering by the mass of strips deposited on a crystal’s surface has been considered. Using small scatterer approximations to obtain a lower limit for the strength of excitation, the scattering method is found to be feasible. A simple, although crude, estimate of the order of magnitude of the scattering resulting from the conductivity of strips deposited on piezoelectric crystals has shown the conductivity to be less important than the mass of the strips. While more careful analysis could be made of the scattering due to conductivity, it is felt that the study of higher mass loadings, which requires higher order approximations, is more significant, since the improvement through higher mass loadings is potentially much greater.

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Microsound Surface Waveguides

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Invited Paper

Abstract—An important requirement for the development of surface wave microsound systems is the ability to guide the energy. The theoretical and experimental progress towards this aim is reviewed. Some preliminary results on topographic guides are presented. Measurement techniques make use of phase-sensitive laser probes to detect the CW surface waves. The technique permits very accurate determination of dispersion characteristics.

I. INTRODUCTION

ACOUSTIC waves are typically five orders of magnitude slower than electromagnetic waves, a property which has for long been exploited in the use of acoustic delay lines. Until fairly recently it was believed that material losses would frustrate the extension of such applications to frequencies much in excess of 100 MHz. In the last few years a new class of materials, having remarkably low losses well into the microwave region has been discovered; moreover, some quite unglamorous materials, such as silicon, whose acoustic attenuation has been known for a long time [1], appeared, on reexamination, to be more suitable for high-frequency application than had been generally appreciated. As a result, microwave acoustic delay lines are now a reality. A number of workers have recently pointed out [2]–[8] that the slowness of acoustic waves leads to another attractive feature, and one which may ultimately prove of even greater import to microelectronic technology. Components whose size is of the order of a wavelength—i.e., most typical microwave components—can be realized in a volume which, in principle, could be 15 orders of magnitude smaller than that required by their electromagnetic counterparts. Since at present the component density in integrated microwave circuits is dominated, to an overwhelming extent, by the size of the passive components, the use of microsound circuits could lead to a dramatic increase in this density. The hopes for realizing such a technology have been further greatly increased by the rapid improvement in transducer technology, notably for surface wave generation [9], [10] and by the development of surface wave acoustic amplifiers [11], [12].

The possibility of using surface acoustic waves, at present, appears to offer the greatest promise for the realization of microelectronic circuits. The main reason is precisely that which has lead to the silicon planar technology—that one can get at the surface.1 In fact, it might prove possible to build microsound circuits on a silicon slice, which would also include active elements. This possibility would be greatly aided if the microsound circuit could be constructed using methods which are compatible with the surface oriented techniques used in silicon device fabrication.

1 The superior performance of surface as compared with bulk transducers, may be attributed to this basic advantage; the surface geometry permits a connection of the electrodes to reduce the total input capacitance, which is not readily duplicated in a bulk transducer.

REFERENCES