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SLIPPING INTERFACES: A POSSIBLE SOURCE OF *S* RADIATION FROM EXPLOSIVE SOURCES

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ABSTRACT

We consider the problem of reflection and refraction of purely compressional waves incident on an interface separating identical solid half-spaces in which the condition of continuity of shear displacement at the boundary is generalized to one that allows slippage. The problem is solved using the Cagniard-de Hoop technique. It is found that the generation of reflected *P* and *S* waves, as well as transmitted *S* waves, is most effective in the case of perfectly unbonded half-spaces. We discuss the implications of this model for the generation of *S* waves by block movement in the vicinity of an underground explosion.

INTRODUCTION

Considerable effort has been devoted to the understanding of *SH* radiation from underground explosive sources. A variety of possible mechanisms have been proposed, as reviewed by Aki and Tsai (1972). They include (1) mode conversion at irregular interfaces (e.g., Aki and Larner, 1970), (2) cracking in the vicinity of the shot point (e.g., Kisslinger *et al.*, 1961), (3) prestress relaxation associated with the creation of the cavity (e.g., Archambeau, 1972, 1973), and (4) triggering of an actual earthquake with large scale faulting, a model favored by Aki and Tsai (1972).

It is likely that a combination of these mechanisms acts to produce the observed *SH*- and Love-wave radiation from underground explosions. The relative contributions of various mechanisms will of course depend on local characteristics of the site. Nevertheless, it may be noted that the two last mechanisms listed above involve the release of strain energy stored in the medium prior to the experiment, and therefore require the presence of sufficient prestress. On the other hand, the two first mechanisms can operate even in the absence of significant initial stress, as evidenced by the observations of Kisslinger *et al.* (1961). These authors observed *SH* radiation generated by small scale explosions detonated in mud, and suggested that near-source cracking was the most likely cause of shear-wave radiation.

In order to better understand the influence of site characteristics and prestress on *SH*- and Love-wave generation, an improved assessment of the phenomena which do not require prestress is needed. Cracking and block motions in the vicinity of explosions fall in this category; a review of these problems is provided by Bache and Lambert (1976).

The purpose of this paper is to demonstrate that the presence of imperfectly bonded interfaces and joints near a purely dilatational source could account for far-

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field transverse components of displacement and thus, with a suitable geometry, for *SH*- and Love-wave radiation.

The problem is greatly simplified because we restrict ourselves to linear boundary conditions, and solve the problem in the far-field, first-motion approximation. It appears that if the bonding is sufficiently weak, and if the geometry is favorable, at least a portion of the observed shear radiation could be explained by this mechanism.

1. Statement of the problem, and boundary conditions. Consider the simple geometry depicted on Figure 1: a purely dilatational point source is located at a height h above the plane interface between two elastic half-spaces. To simplify the problem, we restrict our attention to the case when the two half-spaces are identical. The boundary conditions usually adopted in seismology involve continuity of tractions and displacements, and in this case, the problem reduces to that of a point source in an infinite space.

We shall now relax the boundary conditions, and request continuity of tractions and of normal displacement but allow a jump in the tangential displacement. In other words, slippage is allowed between the two half-spaces. Sezawa and Kanai (1940), Kanai (1961), and Murty (1975, 1976) proposed a boundary condition which

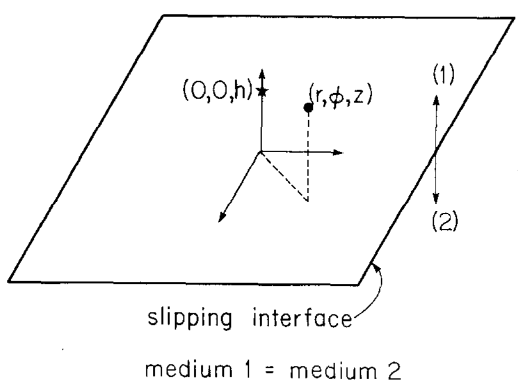

FIG. 1. A purely dilatational point source is located at a height h above the slipping interface separating identical half-spaces.

does allow such slippage to occur, and has the advantage of being linear; it is best described as the result of the following limiting process.

Suppose the two half-spaces are actually coupled through a viscous layer of thickness H and viscosity η. The geometry is described on Figure 2. Assuming a plane Couette flow in the viscous layer, the shear traction is related to the jump in tangential velocity by

$$\sigma_{rz} = \eta \frac{\partial \dot{u}_r}{\partial z} \cong \frac{\eta}{H} \left(\dot{u}_r(r, 0^+, t) - \dot{u}_r(r, 0^-, t) \right). \quad (0 \leqq \eta \leqq \infty) \quad (1.1)$$

We now allow the thickness H and the viscosity η to go to zero, but require that their ratio remain constant. For convenience in the subsequent analyses, we define

$$\frac{\eta}{H} = \frac{\Phi}{1 - \Phi} \frac{\mu}{\beta} \qquad (0 \leqq \Phi \leqq 1) \quad (1.2)$$

where μ is the rigidity and β the shear-wave velocity in the two half-spaces. Φ is a dimensionless bonding parameter. In this limit, the boundary condition becomes

$$\sigma_{rz} = \frac{\Phi}{1-\Phi}\frac{\mu}{\beta}[|\dot{u}_r(r, 0, t)|]^{+}_{-} \tag{1.3}$$

where $[|\dot{u}_r|]^{+}_{-}$ is the jump of tangential velocity across the interface

$$[|\dot{u}_r(r, 0, t)|]^{+}_{-} = \dot{u}_r(r, 0^+, t) - \dot{u}_r(r, 0^-, t).$$

When $\Phi = 0$, the tangential traction vanishes and the interface is perfectly lubricated; when $\Phi = 1$, no relative motion between the two half-spaces is allowed and the interface is perfectly bonded.

The main advantage of the boundary conditions (1.3) is its linearity, which greatly simplifies wave propagation problems. Imperfectly bonded interfaces have been modeled using nonlinear boundary conditions [e.g., Chez *et al.* (1978); Miller (1979)], but the difficulties are such that only plane waves have apparently been considered.

2. *Analysis of the problem.* The wave propagation problem is amenable to a Cagniard-de Hoop treatment, and simple closed form solutions may be found in the first motion approximation.

We consider the potential decomposition (e.g., Cagniard, 1962)

$$\left.\begin{aligned} u_r &= \frac{\partial\Phi}{\partial r} - \frac{\partial\Psi}{\partial z} \\ u_z &= \frac{\partial\Phi}{\partial z} + \frac{1}{r}\frac{\partial}{\partial r}(r\Psi) \end{aligned}\right\} \tag{2.1}$$

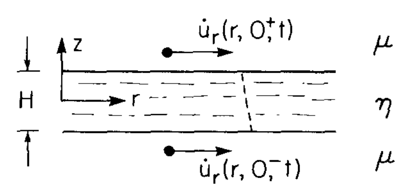

FIG. 2. The identical half-spaces are separated by an interface which is assumed a fluid-filled gap of viscosity η.

and the solutions of the form

$$\left.\begin{aligned} \Phi_i &= e^{st}x_i \\ \Psi_i &= e^{st}y_i \end{aligned}\right\} \qquad i = 1, 2 \tag{2.2}$$

where, for $z > 0$, the source medium,

$$\left.\begin{aligned} x_1 &= \int_0^\infty e^{-\nu_\alpha|z-h|}J_0(kr)\frac{kdk}{\nu_\alpha} + \int_0^\infty Q_1(k)J_0(kr)e^{-\nu_\alpha z}dk \\ y_1 &= \int_0^\infty S_1(k)\frac{\partial}{\partial r}J_0(kr)e^{-\nu_\beta z}dk \end{aligned}\right\} \tag{2.3}$$

and, for $z < 0$

$$\left.\begin{aligned} x_2 &= \int_0^\infty Q_2(k) J_0(kr) e^{\nu_\alpha z}\, dk \\ y_2 &= \int_0^\infty S_2(k) \frac{\partial}{\partial r} J_0(kr) e^{\nu_\beta z}\, dk \end{aligned}\right\}. \tag{2.4}$$

Here

$$\nu_\alpha = \left[k^2 + \frac{s^2}{\alpha^2}\right]^{1/2}, \qquad \nu_\beta = \left[k^2 + \frac{s^2}{\beta^2}\right]^{1/2},$$

and α, β, are the P- and S-wave velocities, respectively.

The unknown coefficients, $\{Q_i(k), S_i(k);\ i = 1, 2\}$ are to be determined by use of the boundary conditions. Continuity of tractions, normal displacement, slip condition (1.3) lead to the following system of simultaneous equations.

$$\begin{pmatrix} \nu_\alpha & \nu_\alpha & k^2 & -k^2 \\ -a & a & -2\mu\nu_\beta k^2 & -2\mu\nu_\beta k^2 \\ 2\mu\nu_\alpha & 2\mu\nu_\alpha & a & -a \\ b & -b & c & c \end{pmatrix} \begin{pmatrix} Q_1 \\ Q_2 \\ S_1 \\ S_2 \end{pmatrix} = m \begin{pmatrix} 1 \\ a/\nu_\alpha \\ 2\mu \\ 2\mu\left[(1 - \Phi) - \dfrac{s}{\beta\nu_\alpha}\Phi\right] \end{pmatrix} \tag{2.5}$$

where

$$\left.\begin{aligned} a &= 2\mu k^2 + \rho s^2 \\ b &= 2\mu\left[(1 - \Phi)\nu_\alpha + \frac{s}{\beta}\Phi\right] \\ c &= a(1 - \Phi) + \frac{2\mu s \nu_\beta}{\beta}\Phi \\ m &= k e^{-\nu_\alpha h} \end{aligned}\right\}. \tag{2.6}$$

The determinant in (2.5) is given by

$$\Delta(k) = 4\rho\nu_\alpha s^2 D(k)$$

$$D(k) = [(2\mu k^2 + \rho s^2)^2 - 4\mu^2\nu_\alpha\nu_\beta k^2](1 - \Phi) + \frac{2\mu\rho\nu_\beta}{\beta} s^3\Phi. \tag{2.7}$$

Notice that in the limit $\phi = 0$, $D(k)$ reduces to the Rayleigh denominator.

The solutions to (2.5) are

$$Q_1(k) = \frac{e^{-\nu_\alpha h}}{D(k)}[4\mu^2\nu_\beta k^3(1 - \Phi)] \tag{2.8a}$$

$$Q_2(h) = \frac{k e^{-\nu_\alpha h}}{\nu_\alpha D(k)}\left[\frac{2\mu\rho\nu_\beta}{\beta} s^3\Phi + (2\mu k^2 + \rho s^2)^2(1 - \Phi)\right] \tag{2.8b}$$

$$S_1(k) = S_2(k) = \frac{ke^{-\nu_\alpha h}}{D(k)}[2\mu(2\mu k^2 + \rho s^2)(1 - \Phi)]. \tag{2.8c}$$

In the case of a perfectly bonded interface, $\Phi = 1$, and the only nonvanishing coefficient is $Q_2(k)$, as expected.

By considering derivatives of these coefficients with respect to the bonding parameter Φ, one finds that no extrema occur for Φ [0, 1]. In other words, the most efficient conversion of P waves to S waves must occur for $\Phi = 0$, that is, a perfectly lubricated interface.

We can now substitute these coefficients in the relations (2.3) and (2.4), and replace k in terms of the ray parameter p (e.g., de Hoop, 1960)

$$k = -isp. \tag{2.9}$$

After application of the Schwartz reflection principle, the solution takes the form

$$x_i = \frac{2}{\pi} Im \int_0^{i\infty} Q_i(p) K_0(spr)\, dp \tag{2.10a}$$

$$y_i = \frac{2}{\pi} Im \int_0^{i\infty} S_i(p) \frac{\partial K_0(spr)}{\partial r}\, dp \tag{2.10b}$$

where

$$Q_1(p) = \frac{e^{-s\eta_\alpha|z-h|}}{\eta_\alpha} sp - \frac{4\mu^2(1-\Phi)sp^3\eta_\beta e^{-\eta_\alpha|z+h|}}{\mathscr{D}(p)} \tag{2.11a}$$

$$Q_2(p) = \frac{e^{-s\eta_\alpha(h-z)}}{\eta_\alpha \mathscr{D}(p)} sp \left[\frac{2\mu\rho\eta_\beta}{\beta}\Phi + (\rho - 2\mu p^2)^2(1-\Phi)\right] \tag{2.11b}$$

$$S_1(p) = \frac{e^{-s(\eta_\alpha h + \eta_\beta z)}}{\mathscr{D}(p)} 2\mu(1-\Phi)p(\rho - 2\mu p^2) \tag{2.11c}$$

$$S_2(p) = \frac{e^{-s(\eta_\alpha h - \eta_\beta z)}}{\mathscr{D}(p)} 2\mu(1-\Phi)p(\rho - 2\mu p^2) \tag{2.11d}$$

and where the following defintions hold

$$\eta_\alpha = \left(\frac{1}{\alpha^2} - p^2\right)^{1/2}; \qquad \eta_\beta = \left(\frac{1}{\beta^2} - p^2\right)^{1/2} \tag{2.12}$$

$$\mathscr{D}(p) = [(\rho - 2\mu p^2)^2 + 4\mu^2\eta_\alpha\eta_\beta p^2](1-\Phi) + \frac{2\mu\rho\eta_\beta}{\beta}\Phi. \tag{2.13}$$

Asymptotic approximations of (2.10) and first-motion approximations can then be obtained in the usual fashion (e.g., Helmberger, 1974; Langston, 1976). By substitution of the asymptotic results

$$K_0(spr) \cong \sqrt{\frac{\pi}{2spr}}\, e^{-spr}$$

$$\frac{\partial K_0(spr)}{\partial r} \cong -sp \sqrt{\frac{\pi}{2spr}}\, e^{-spr} \tag{2.14}$$

the solutions (2.10) are reduced to integrals of the form

$$\tilde{I}(r, z, s) = \sqrt{\frac{2}{\pi sr}}\, Im \int_0^{i\infty} e^{-st(r,z,p)} f(p)\, dp \tag{2.15}$$

from which the displacement must be retrieved via the formulas (2.1). This requires finding the spatial derivatives of I. Keeping only the far-field terms, we have

$$\frac{\partial \tilde{I}(r, z, s)}{\partial r} = -s \sqrt{\frac{2}{\pi sr}}\, Im \int_0^{i\infty} e^{-st(r,z,p)} p f(p)\, dp \tag{2.16}$$

$$\frac{\partial \tilde{I}(r, z, s)}{\partial z} = -s \sqrt{\frac{2}{\pi sr}}\, Im \int_0^{i\infty} e^{-st(r,z,p)} \frac{\partial t}{\partial z} f(p)\, dp. \tag{2.17}$$

The body-wave contributions, which we are interested in, are obtained by changing the variable of integration to t so that by use of Cauchy-Goursat's theorem on the complex t-plane we have

$$\frac{\partial \tilde{I}}{\partial r} = -s \sqrt{\frac{2}{\pi sr}}\, Im \int_0^{\infty} H(t - t_A) p f(p) e^{-st} \frac{dp}{dt}\, dt \tag{2.18}$$

$$\frac{\partial \tilde{I}}{\partial z} = -s \sqrt{\frac{2}{\pi sr}}\, Im \int_0^{\infty} H(t - t_A) \frac{\partial t}{\partial z} f(p) e^{-st} \frac{dp}{dt}\, dt. \tag{2.19}$$

The geometrical arrival time is defined by

$$\left.\begin{aligned} t_A &= t(r, z, p_0) \\ \left.\frac{\partial t}{\partial p}\right|_{p_0} &= 0 \end{aligned}\right\} \tag{2.20}$$

If $F(t)$ is the source function for the potentials, then the time domain equivalents to (2.18) and (2.19) are

$$\frac{\partial I(r, z, t)}{\partial r} = \frac{1}{\pi} \sqrt{\frac{2}{r}} \frac{\partial}{\partial t} \left\{ F'(t) * \frac{H(t)}{t^{1/2}} * Im\left(H(t - t_A) p f(p) \frac{dp}{dt} \right) \right\} \tag{2.21}$$

$$\frac{\partial I(r, z, t)}{\partial z} = \frac{1}{\pi} \sqrt{\frac{2}{r}} \frac{\partial}{\partial t} \left\{ F'(t) * \frac{H(t)}{t^{1/2}} * Im\left(H(t - t_A) \frac{\partial t}{\partial z} f(p) \frac{dp}{dt} \right) \right\} \tag{2.22}$$

where $*$ indicates a convolution. To compute the first-motion approximation we note that $t(r, z, p)$ is of the form

$$t = pr + \eta_\alpha g_\alpha(z) + \eta_\beta g_\beta(z) \tag{2.23}$$

where $g_\alpha(z)$ and $g_\beta(z)$ can be constructed from t in Table 1 for the various rays. Using the definition (2.20) and replacing t by t_A in the first-motion approximation, we get

$$\frac{dp}{dt} \cong i\left[2(t - t_A)\left(\frac{g_\alpha(z)}{\alpha^2\eta_\alpha^3} + \frac{g_\beta(z)}{\beta^2\eta_\beta^3}\right)\right]^{-1/2} \tag{2.24}$$

so that, by substitution in (2.21) and (2.22), and reduction

$$\frac{\partial I(r, z, t)}{\partial r} \cong \frac{p_0 f(p_0) F(t - t_A)}{\left[r\left(\frac{g_\alpha(z)}{\alpha^2\eta_\alpha^3(p_0)} + \frac{g_\beta(z)}{\beta^2\eta_\beta^3(p_0)}\right)\right]^{1/2}} \tag{2.25}$$

$$\frac{\partial I(r, z, t)}{\partial z} \cong \frac{1}{p_0}\frac{\partial t}{\partial z}\frac{\partial I}{\partial r}. \tag{2.26}$$

TABLE 1

REFLECTED AND TRANSMITTED RAYS

Ray Description	$t(r, z, p)$	$\partial t/\partial z$	$f(p)$
Incident P	$pr + \eta_\alpha \lvert z - h \rvert$	$\eta_\alpha sgn(z - h)$	$p^{1/2}/\eta_\alpha$
Reflected P	$pr + \eta_\alpha(z + h)$	η_α	$4\mu^2(1 - \Phi)p^{5/2}\eta_\beta/\mathscr{D}(p)$*
Transmitted P	$pr + \eta_\alpha(h - z)$	$-\eta_\alpha$	$2\mu\rho\Phi p^{1/2}\eta_\beta/[\beta\eta_\alpha\mathscr{D}(p)] + (1 - \Phi)p^{1/2}(\rho - 2\mu p^2)^2/[\eta_\alpha\mathscr{D}(p)]$
Reflected S	$pr + \eta_\alpha h + \eta_\beta$	η_β	$\mu(1 - \Phi)p^{3/2}(2\mu p^2 - \rho)/\mathscr{D}(p)$
Transmitted S	$pr + \eta_\alpha h - \eta_\beta$	$-\eta_\beta$	$\mu(1 - \Phi)p^{3/2}(2\mu p^2 - \rho)/\mathscr{D}(p)$

* $\mathscr{D}(p)$ is given by (2.13).

The displacement components for the various rays can be obtained by using these results in (2.1) and substituting the quantities from Table 1.

3. *Radiation patterns.* Suppose the displacements found in the previous section are rotated to a coordinate system that renders the slipping interface vertical. For the purpose of our block-motion study, such coordinates are Earth coordinates.

For a takeoff angle of 90° in Earth coordinates, all S radiation reflected from and transmitted through the interface is pure SH relative to the surface of the Earth. Displacement radiation patterns in that case are simple to understand and were computed for different values of the bonding parameter and the following values of the remaining parameters

Parameter	Value
α/β	1.72
ρ	3.2 gm/cm^3
h	1.0 km
$R = [r^2 + (z - h)^2]^{1/2}$	1000.0 km.

Amplitudes are measured perpendicular and parallel to the ray joining the receiver and the source. This gives rise, respectively, to the transverse and radial components of displacement as seen in Figure 3. The radiation patterns for a unit step-source function are given in Figure 4 where the solid and dotted lines are the transverse and radial components, respectively, and where positive, for the transverse component, is defined to be counterclockwise. The horizontal axis represents the edge view of the slipping interface and the units on this axis represent the ratio of the amplitude of the component to that of the incident P wave.

These patterns are seen to change shape with bonding. For $\phi = 0.0$, the angle of maximum displacement corresponds to angles of critical reflections. Similar lobed patterns were found by Burridge *et al.* (1964) for reflections from a free surface.

There is good conversion for bonding $\phi = 0.0$: the amplitude of the SH wave is at least 0.5 of the P wave. Conversion gets progressively smaller as the bonding is varied to higher values. For $\phi = 0.5$ and 0.9, the amplitude of the SH wave is,

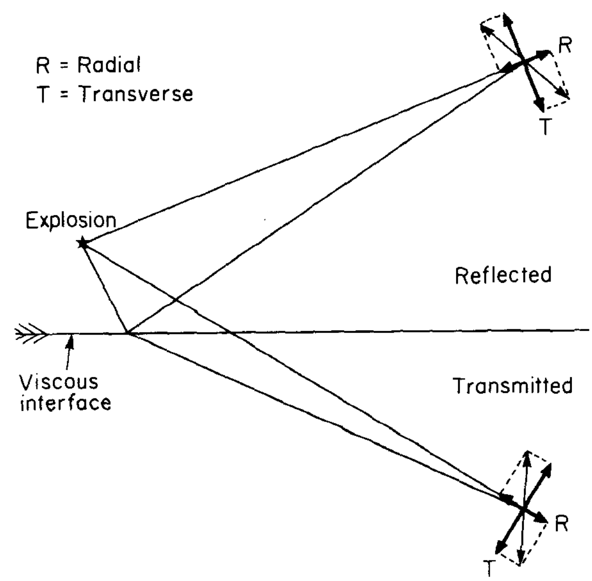

FIG. 3. Amplitudes are measured perpendicular (transverse component T) and parallel (radial component R) to the ray-joining source with receiver.

respectively, 0.1 and 0.01 of the P-wave amplitude. This is due to the fact that as the bonding becomes stronger the jump in shear displacement across the interface decreases, and vanishes when the bonding is perfect, i.e., $\phi = 1.0$. This is seen in Figure 5.

So it appears that a low value of the bonding parameter is needed for this mechanism to be an efficient shear-wave generator. It is also likely that such joints or faults have a nonlinear rheology. Recent developments in equivalent linearization techniques of nonlinear boundary conditions have accurately approximated the behavior of slip surfaces under Coulomb friction by use of the linearly viscous boundary condition for the case of incident periodic waves (e.g., Miller, 1979). The results of this linearization process are that bonding is frequency and amplitude dependent such that when the amplitudes and the frequencies are high, which gives rise to high local transient stresses, the fault surfaces tend to debond. If this property persists for the case of an arbitrary transient wave, the proximity of the large-scale faulting that results near underground nuclear explosions (e.g., Bache and Lambert,

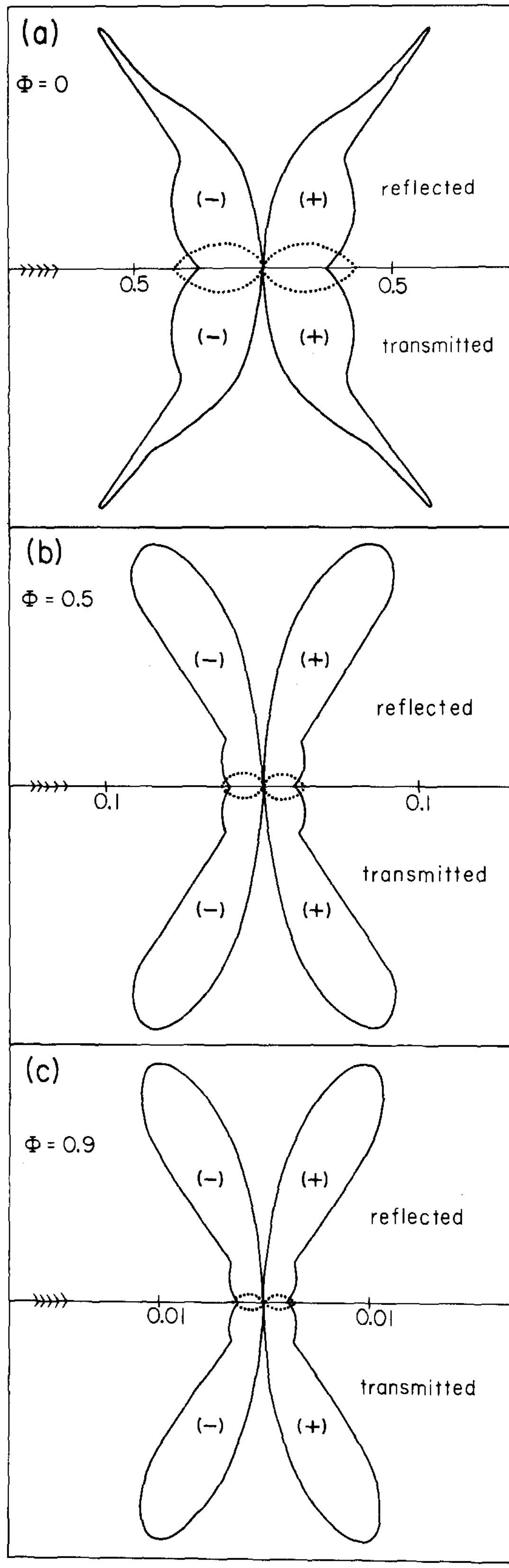

FIG. 4. Radiation patterns for a take-off angle of 90°. Solid line represents the transverse component while the dotted line represents the radial component.

1976) would suggest that bonding on these slipping surfaces is low and, therefore, that these surfaces constitute efficient shear-wave generators.

These radiation patterns are subject to change in shape and amplitude with changes in the values of the parameters in the expressions for the displacement. Most important of these is the perpendicular distance h from the source to the slipping interface. From expressions (2.25) and (2.26), together with the results of Table 1, it can be easily shown that the efficiency of P and S conversion increases with decreasing h. Thus, this linear theory predicts that if a slipping joint is sufficiently close to an underground explosion it could be an efficient converter of P to S waves. One must remember, however, that the approximation developed in the previous section is far-field and linear and therefore must break down for small h, so that this conclusion is only qualitative.

This faulting process due to the presence of a joint in the vicinity of an underground explosion gives rise to a different displacement pattern than that given either by the relief of prestress along a fault or the impinging of P waves on a slipless impedance mismatching surface, both of the same orientation as the joint. As

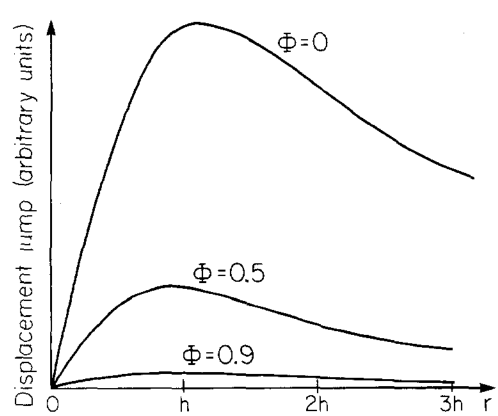

FIG. 5. Displacement jump across the slipping interface for various values of bonding parameter Φ.

appreciated in Figure 4, the SH displacement radiation pattern has a cylindrical symmetry with respect to the line through the source and perpendicular to the interface: it is symmetric with respect to the interface but antisymmetric with respect to the line perpendicular to the interface through the source. The release of prestress along a vertical fault, in this case, a vertical strike-slip fault, has a double-couple radiation pattern which, in the case of the SH component, is symmetric with respect to the interface and the line perpendicular to the interface through the source, as can be appreciated in Heaton (1979). Furthermore, the nodes occur at different angles than those for a slipping interface. On the other hand, the radiation pattern from the incidence of P waves on a slipless surface with impedance mismatch has, as does that for the slipping interface, symmetry with respect to the line perpendicular to the interface and through the source but it differs from the radiation pattern of the slipping joint in that it is, in general, asymmetric with respect to the interface as is seen in Ewing *et al.* (1957). Furthermore, an anomalously large lateral impedance mismatch is necessary to account for the SH waves that are observed radiating from explosions.

This distinction of radiation patterns may provide the best tool through measurement to determine whether block motion is indeed a significant contributor of *S* waves from large underground explosions. It does not matter that the actual rheology of the fault is nonlinear (e.g., Coulomb friction) because the radiation pattern of the first motion will have the same polarity as in this study. On the other hand if the test site is significantly prestressed, observation of block motion through the polarity of the radiation pattern may be difficult to observe.

We have derived a high frequency, first motion approximation for the scattered *P* and *S* waves generated by an incident spherical *P* wave on a slipping interface separating identical solids. Conversion is most favorable when the source is close to the interface, and when the interface is debonded. In the case of underground explosions, the former requirement is met in that there is observational evidence that large scale cracking occurs near the source; in support of the latter requirement, there is limited theoretical evidence for large amplitude, high-frequency incident waves, and more realistic (e.g., Coulomb) interface conditions. Our model provides, therefore, a possible theoretical interpretation for observed block motion near buried explosions, and lends support to the hypothesis that such block motion contributes to the anomalous *S*-wave radiation. In the absence of significant prestress, such block motion may be the dominant contributor.

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