J. Fluid Mech. (1967), *vol.* 29, *part* 2, *pp.* 373–383
Printed in Great Britain

An underwater explosive shock gun

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(Received 20 December 1966)

A theoretical study is made of a simple design for an underwater shock gun, which consists of a chamber in the form of a hollow circular cone, with a spherical sector of explosive charge fitted into the apex. When the explosive is initiated at the apex, the resulting sector of a spherical blast wave will be diffracted by expansion waves moving inwards after the leading shock has emerged from the rim of the cone.

The progress of the expansion wave-fronts is calculated, and the results show a surprising inability of the diffraction process to 'eat into' the full-strength sector of spherical blast. It is found to be possible to design such a gun so that it is capable of projecting a high intensity shock-pressure beam over a considerable range, using only a very small explosive charge.

1. Introduction

Reports were received in 1957 that an underwater explosive gun, 'firing a shock wave instead of a bullet' had been seen somewhere in Europe being used for underwater fishing, and it was reputed to be capable of stunning a large fish at a considerable range. It was thought that, if these reports were true, there may be other applications of such a device. In this paper, a possible simple design, capable of theoretical study, and which appears to have all the essential properties of the reported device, is proposed and studied.

The 'chamber' of the proposed gun consists merely of a rigid hollow circular cone (see figure 1) containing an explosive charge in the form of a spherical sector fitting into the apex of the cone, the point of initiation being at the apex. The dimensions of the conical chamber and the mass of the explosive charge must satisfy certain conditions which are derived theoretically.

The principle behind the proposed gun is explained very simply, for if the conical chamber were rigid and of infinite extent, the radial distribution of pressure behind the leading shock at any time would be precisely the same (except for the boundary layer on the wall of the cone) as that obtained from detonating the complete spherical charge of which the conical sector forms a part.

When the conical chamber is of finite size, however, the flow behind the leading shock, after it leaves the rim of the cone, is diffracted by expansion waves moving inwards from the initially undisturbed fluid. The rate at which diffraction occurs near the leading shock must be studied, therefore, in order to determine the minimum size of cone which is necessary to meet the required performance.

The results obtained show a surprising inability of the diffraction process to 'eat into' the full-strength sector of spherical blast wave produced by a sector of a spherical charge; and thus appear to bear out the reported possibilities of such a device.

The shock diffraction problem studied here is related to the problem of the irregular reflexion at an air–water interface of the leading blast shock from an underwater explosion, as described, for instance, by Keil (1948) and Rosenbaum & Snay (1956). At the air–water interface, the oblique shock may be reflected regularly as a centred simple wave or Prandtl–Meyer expansion, but the reflexion becomes critical at some stage when the shock is nearly normal to the surface, when the flow into the expansion wave (in the appropriate quasi-stationary frame of reference) becomes sonic. Beyond this stage, the rarefaction wave diffracts or 'eats into' the incident shock, in much the same way as in the problem considered here.

This paper was written in the autumn of 1957, and was published as a Ministry of Supply report in 1958. Subsequent attempts to trace the origin of the reports which initiated the work described in the paper proved unsuccessful. The possibility of military applications of the device has prevented open publication of the paper until now.

The original paper, in its limited publication, aroused some interest, particularly in the United States, and a variety of possible applications has been examined. Experimental measurements have also been made, in the United States, with a device designed in accordance with the theory, and the results essentially confirm the theory. A description of these experiments will, it is hoped, appear soon in an American Journal.

The original paper was also one of the first publications to draw attention to, and to make use of, the geometrical scaling of explosions described above, as distinct from the size scaling of explosions. Laporte & Cole (1957) and Campbell (1958) also described conical or sector shock-tubes, driven by compressed gas, for generating spherical or cylindrical shocks under laboratory conditions; and, subsequently, Filler (1960) made measurements in a conical shock tube driven by a high explosive charge in the apex. More recently the U.S. Naval Weapons Laboratory† has announced the construction, due for completion in October 1966, of a conical shock-tube 2500 ft. in length, for the full-scale simulation of nuclear blasts up to 20 kilotons, and at atmospheric conditions appropriate to altitudes up to 100,000 ft. above sea level.

The original 1957 manuscript has been revised and the presentation improved, but, substantially, the paper is printed here as it was written in 1957.

2. The properties of weak shocks in water

Theoretical studies of the flow properties behind shock waves moving through water have been undertaken by many authors, assuming various forms of equation of state for the medium. For the present purpose, however, it is both

† U.S. Naval Weapons Laboratory, Dahlgren, Virginia. 1965. *The Conical Shock Tube Facility*.

convenient and sufficient to use simple approximate analytical relationships between the flow parameters, relationships which are valid for relatively weak shocks.

Consider, therefore, a shock wave moving with a velocity U into water at rest, in which p_0 denotes the (constant) hydrostatic pressure, ρ_0 the density and a_0 the local speed of sound. Immediately behind the shock the particle velocity will be u, say, the pressure p, the density ρ and the local speed of sound a. Now conservation of mass demands that

$$\rho(U-u) = \rho_0 U, \tag{2.1}$$

and conservation of momentum gives

$$p+\rho(U-u)^2 = p_0+\rho_0 U^2. \tag{2.2}$$

The flow parameters would be specified exactly by a third equation, which would ensure that energy is conserved, and by the appropriate equation of state. In water, however, it is found that the influence of entropy changes is negligible, unless the pressure is extremely high, so that, in water, pressure may generally be considered as a function of density alone, and not dependent on entropy. Thus, the caloric equation of state of water, cf. Courant & Friedrichs (1948), may be taken as

$$p/p_0-1 = A[(\rho/\rho_0)^\gamma-1], \tag{2.3}$$

where A and γ are independent of entropy and have the values $A = 3000$, $\gamma = 7$. In this case, therefore, equations (2.1), (2.2) and (2.3) can be solved to give the three flow parameters U, ρ and u in terms of p. Thus, if we write

$$\epsilon = (p/p_0-1)/A\gamma, \tag{2.4}$$

then

$$\frac{U^2}{a_0^2} = \frac{\epsilon(1+\gamma\epsilon)^{1/\gamma}}{(1+\gamma\epsilon)^{1/\gamma}-1}, \tag{2.5}$$

$$\rho/\rho_0 = (1+\gamma\epsilon)^{1/\gamma}, \tag{2.6}$$

$$u^2/a_0^2 = \epsilon[1-(1+\gamma\epsilon)^{-1/\gamma}], \tag{2.7}$$

and

$$a^2/a_0^2 = (1+\gamma\epsilon)^{1-1/\gamma}, \tag{2.8}$$

where a_0, of course, is given by

$$a_0^2 = (\partial p/\partial\rho)_0 = A\gamma p_0/\rho_0. \tag{2.9}$$

3. The decay of spherical shocks in water

Experimental and theoretical results concerning the decay of shock waves originating from spherical underwater explosions are discussed by Cole (1948). For the present purpose it is convenient to use existing results for an expression for the peak overpressure of the form

$$p-p_0 = k(W^{\frac{1}{3}}/R)^\alpha, \tag{3.1}$$

where W is the mass of the spherical charge, R is the shock radius, and k and α are constants depending on the composition of the explosive (k has dimensions).

Formulae of this type have been fitted to experimental data over limited ranges of pressure, and values of the constants for TNT, loose tetryl and pentolite are given by Cole (1948, p. 242). With TNT, for example, if peak overpressure is measured in lb. wt./in.², W in pounds and R in feet, it is found that

$$k = 2{\cdot}16\times 10^4 \quad \text{and} \quad \alpha = 1{\cdot}13.$$

In this case, the experimental fit was obtained over the pressure range 20,000–500 lb. wt./in.² and we assume that the same fit will be approximately valid also for lower peak overpressures.

A length d may be defined, therefore, for any underwater explosion, such that the peak overpressure at a shock radius R is given by

$$p/p_0 - 1 = A\gamma(d/R)^\alpha, \tag{3.2}$$

where

$$d = (k/A\gamma p_0)^{1/\alpha}\, W^{\frac{1}{3}}. \tag{3.3}$$

The shock parameters of all underwater explosions, within some limited range of overpressure, are then given by (2.5)–(2.8) with

$$U = dR/dt \tag{3.4}$$

and

$$\epsilon = (d/R)^\alpha. \tag{3.5}$$

4. The diffraction problem

Let us now consider a conical sector of a spherical charge, lying in the apex of a hollow, right, circular, rigid cone of slant height b and semi-angle β (see figure 1), which is detonated at the vertex. If the cone were infinitely large, the ensuing flow would be precisely the same (neglecting the boundary layer on the walls of the cone) as the flow caused by a complete spherical charge of the same radius when detonated at the centre. When the cone is of finite size, however, the leading shock, after it leaves the rim AA' of the cone, is diffracted by expansion waves moving inwards from the initially undisturbed fluid.

Referring now to figure 1, we consider the flow at time t when the leading shock PP' has reached a radius R and the expansion wave-front PD has reached the point P on the shock, so that PP' is the surviving portion of the spherical shock front which is undisturbed. After a further time δt the shock will have reached a radius $R+\delta R$. Now we know that the expansion wave-front will propagate everywhere with a speed equal to the local speed of sound relative to the fluid. The particle at P at time t will move to S, a radial distance of $u\,\delta t$, after a time δt. At time $t+\delta t$, therefore, the point of intersection of the expansion wave-front and the undiffracted shock will be at Q, a point on the undisturbed shock front distant $a\,\delta t$ from S. This, of course, assumes that no other (compressive) disturbance reaches the shock at a position in advance of Q.

From the geometry of figure 1, therefore, we may write

$$(\delta R - u\,\delta t)^2 + R^2(\delta\theta)^2 = a^2(\delta t)^2,$$

which, in the limit when $\delta t \to 0$, gives the differential equation

$$\left(\frac{dR}{dt} - u\right)^2 + R^2\left(\frac{d\theta}{dt}\right)^2 = a^2 \tag{4.1}$$

for the rate of advance of the expansion wave-front along the shock.

After substituting for $U = dR/dt$, u and a from (2.5), (2.7) and (2.8), it is found that

$$R^2\left(\frac{d\theta}{dR}\right)^2 = \frac{(1+\gamma\epsilon)^{1+1/\gamma}-1-(\gamma+1)\,\epsilon}{\epsilon(1+\gamma\epsilon)^{2/\gamma}}.$$

This may be expanded to give, after substituting $\epsilon = (d/R)^\alpha$ from (3.5),

$$R^2\left(\frac{d\theta}{dR}\right)^2 = \frac{(\gamma+1)}{2}\left(\frac{d}{R}\right)^\alpha - \frac{(\gamma+1)(\gamma+5)}{6}\left(\frac{d}{R}\right)^{2\alpha} + O\left(\frac{d}{R}\right)^{3\alpha}. \qquad (4.2)$$

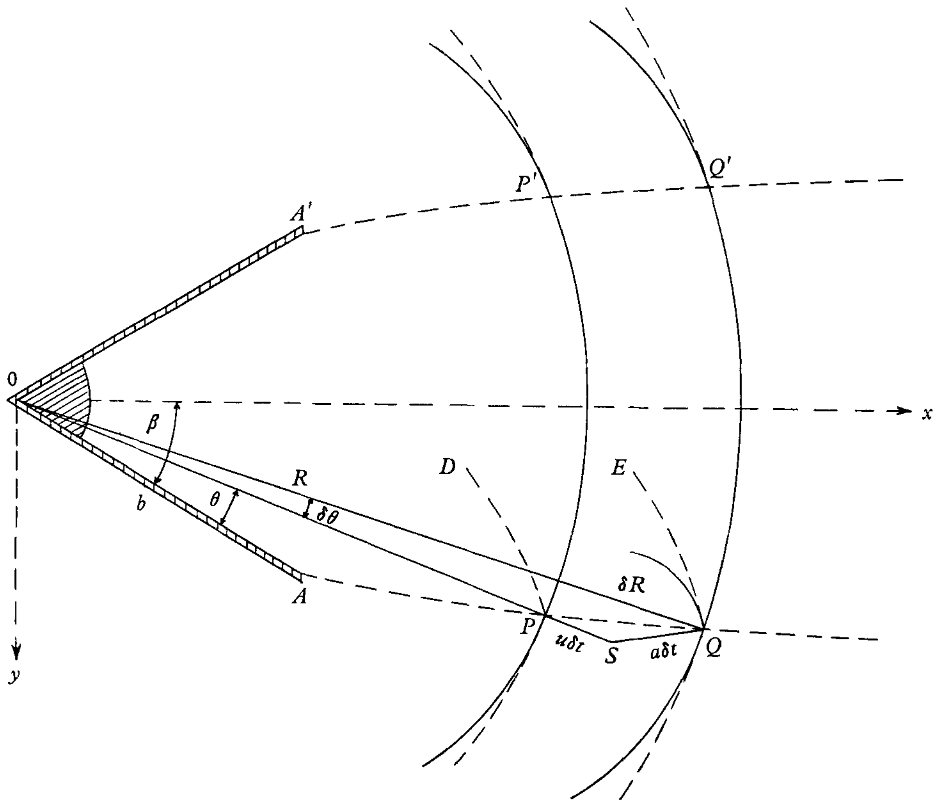

FIGURE 1. Diffraction of the leading shock from an underwater explosive shock gun.

The locus described must originate at the rim of the cone, A, so that $R = b$ when $\theta = 0$. With this boundary condition, (4.2) integrates, therefore, to give

$$\alpha\theta = (2\gamma+2)^{\frac{1}{2}}\left[\left(\frac{d}{b}\right)^{\frac{1}{2}\alpha} - \left(\frac{d}{R}\right)^{\frac{1}{2}\alpha}\right] - \frac{(2\gamma+2)^{\frac{1}{2}}(\gamma+5)}{18} \times \left[\left(\frac{d}{b}\right)^{\frac{3}{2}\alpha} - \left(\frac{d}{R}\right)^{\frac{3}{2}\alpha}\right] + O\left[\left(\frac{d}{b}\right)^{\frac{5}{2}\alpha} - \left(\frac{d}{R}\right)^{\frac{5}{2}\alpha}\right],$$

and when $\gamma = 7$, this reduces simply to

$$\frac{\alpha\theta}{4} = \left[\left(\frac{d}{b}\right)^{\frac{1}{2}\alpha} - \left(\frac{d}{R}\right)^{\frac{1}{2}\alpha}\right] - \frac{2}{3}\left[\left(\frac{d}{b}\right)^{\frac{3}{2}\alpha} - \left(\frac{d}{R}\right)^{\frac{3}{2}\alpha}\right] + O\left[\left(\frac{d}{b}\right)^{\frac{5}{2}\alpha} - \left(\frac{d}{R}\right)^{\frac{5}{2}\alpha}\right], \qquad (4.3)$$

as the required locus. It may be referred to the Cartesian co-ordinates (x, y) of figure 1, by the relations

$$x = R\cos(\beta-\theta), \quad y = R\sin(\beta-\theta). \tag{4.4}$$

If the conical sector of explosive has a mass m, the mass W of the sphere, of which the sector forms a part, is

$$W = 2m/(1-\cos\beta). \tag{4.5}$$

The length d, therefore, is given by

$$d = D(1-\cos\beta)^{-\frac{1}{3}}, \tag{4.6}$$

where

$$D = \left(\frac{k}{A\gamma p_0}\right)^{1/\alpha}(2m)^{\frac{1}{3}}. \tag{4.7}$$

In order to ensure that the cone will be large enough to contain the explosive, we must have

$$\tfrac{4}{3}\pi\rho_E b^3(1-\cos\beta) > 2m, \tag{4.8}$$

where ρ_E is the density of the explosive. But when the explosive charge nearly fills the conical chamber, the shock pressure at the rim AA' of the cone will be too large for the expansion (4.2) to be valid.

5. Analysis of results

Let us examine first the nature of the locus of positions along the leading shock which are reached by the expansion wave-front. If we write

$$\phi = \beta - \frac{4}{\alpha}\left(\frac{d}{b}\right)^{\frac{1}{2}\alpha} + \frac{8}{3\alpha}\left(\frac{d}{b}\right)^{\frac{3}{2}\alpha} + O\left(\frac{d}{b}\right)^{\frac{5}{2}\alpha}, \tag{5.1}$$

then the locus (4.3) may be written as

$$\beta-\theta = \phi + \frac{4}{\alpha}\left(\frac{d}{R}\right)^{\frac{1}{2}\alpha} - \frac{8}{3\alpha}\left(\frac{d}{R}\right)^{\frac{3}{2}\alpha} + O\left(\frac{d}{R}\right)^{\frac{5}{2}\alpha}. \tag{5.2}$$

It is somewhat easier to present the main features of this curve if it is expressed in terms of the Cartesian co-ordinates, defined by (4.4). Now when R is large it is found that

$$\begin{aligned}\frac{R\cos\phi}{x} = {} & 1 + \frac{4}{\alpha}\tan\phi\left(\frac{d\cos\phi}{x}\right)^{\frac{1}{2}\alpha} + \frac{8}{\alpha^2}[1+(2-\alpha)\tan^2\phi]\left(\frac{d\cos\phi}{x}\right)^{\alpha} \\ & + \frac{8}{3\alpha^3}\tan\phi\,[20-18\alpha^2-\alpha^2+3(2-\alpha)(4-3\alpha)\tan^2\phi]\left(\frac{d\cos\phi}{x}\right)^{\frac{3}{2}\alpha} + O\left(\frac{d\cos\phi}{x}\right)^{2\alpha},\end{aligned} \tag{5.3}$$

so that the locus may be presented in the form

$$\begin{aligned}\frac{y}{x} = {} & \tan\phi + \frac{4}{\alpha}\sec^2\phi\left(\frac{d\cos\phi}{x}\right)^{\frac{1}{2}\alpha} + \frac{8(2-\alpha)}{\alpha^2}\sec^2\phi\,\tan\phi\left(\frac{d\cos\phi}{x}\right)^{\alpha} \\ & + \frac{8}{3\alpha^3}[8-6\alpha-\alpha^2+3(2-\alpha)(4-3\alpha)\tan^2\phi]\sec^2\phi\left(\frac{d\cos\phi}{x}\right)^{\frac{3}{2}\alpha} + O\left(\frac{d\cos\phi}{x}\right)^{2\alpha}.\end{aligned} \tag{5.4}$$

It will be seen, therefore, from either (5.2) or (5.4) that if ϕ is negative the expansion wave-front moving up the shock will always reach the axis $y = 0$ and the undisturbed portion of the shock wave will disappear. If ϕ is positive, however, the wave-front will not reach the axis $y = 0$, unless other (compressive) disturbances reach the shock, and the undisturbed part of the shock will propagate in precisely the same manner as the leading shock from the equivalent spherical charge.

Thus $\phi \geqslant 0$ is the condition necessary to ensure that some portion of the undisturbed spherical shock survives indefinitely; and in the limiting case when $\phi = 0$ the locus is given by

$$\frac{y}{x} = \frac{4}{\alpha}\left(\frac{d}{x}\right)^{\frac{1}{2}\alpha} + \frac{8}{3\alpha^3}(8-6\alpha-\alpha^2)\left(\frac{d}{x}\right)^{\frac{3}{2}\alpha} + O\left(\frac{d}{x}\right)^{\frac{5}{2}\alpha}. \tag{5.5}$$

The parameter ϕ, defined in (5.1), is obviously of major interest and its behaviour under varying conditions will provide the general principles on which the design of an underwater shock gun could be based.

Remembering that d is a function of β, given by (4.6), it is found that

$$\frac{\partial\phi}{\partial\beta} = 1 + \frac{2\sin\beta}{3(1-\cos\beta)}\left(\frac{d}{b}\right)^{\frac{1}{2}\alpha} - \frac{4\sin\beta}{3(1-\cos\beta)}\left(\frac{d}{b}\right)^{\frac{3}{2}\alpha} + O\left(\frac{d}{b}\right)^{\frac{5}{2}\alpha}, \tag{5.6}$$

so that $\partial\phi/\partial\beta$ is positive. This means that for a given mass m of a particular explosive and a given slant height of cone b, ϕ will increase monotonically with β. It follows, therefore, that if β is equal to or greater than a certain critical angle β^* which makes $\phi = 0$ in (5.1), and is given by the solution of

$$\begin{aligned}\beta^* - \frac{4}{\alpha}\left[\frac{D}{b}(1-\cos\beta^*)^{-\frac{1}{3}}\right]^{\frac{1}{2}\alpha} + \frac{8}{3\alpha}\left[\frac{D}{b}(1-\cos\beta^*)^{-\frac{1}{3}}\right]^{\frac{3}{2}\alpha}& \\ + O\left[\frac{D}{b}(1-\cos\beta^*)^{-\frac{1}{3}}\right]^{\frac{5}{2}\alpha} &= 0,\end{aligned} \tag{5.7}$$

then, under the assumptions of the theory, there will always survive some portion of undisturbed spherical shock which will propagate to infinity like the shock from the equivalent spherical charge.

Again, differentiating (5.1) with respect to m, it is found that

$$3m\frac{\partial\phi}{\partial m} = -2\left(\frac{d}{b}\right)^{\frac{1}{2}\alpha} + 4\left(\frac{d}{b}\right)^{\frac{3}{2}\alpha} + O\left(\frac{d}{b}\right)^{\frac{5}{2}\alpha}, \tag{5.8}$$

so that $\partial\phi/\partial m$ is negative. This means that for a given explosive composition and a given size and shape of cone, ϕ will decrease monotonically as m increases. It follows, therefore, that if m does not exceed a certain critical mass m^* which makes $\phi = 0$ in (5.1) and is given by the solution of the equation

$$\left.\begin{aligned}\beta - \frac{4}{\alpha}\left(\frac{d^*}{b}\right)^{\frac{1}{2}\alpha} + \frac{8}{3\alpha}\left(\frac{d^*}{b}\right)^{\frac{3}{2}\alpha} + O\left(\frac{d^*}{b}\right)^{\frac{5}{2}\alpha} = 0,& \\ \text{where} \qquad d^* = (k/A\gamma p_0)^{1/\alpha}(1-\cos\beta)^{-\frac{1}{3}}(2m^*)^{\frac{1}{3}},&\end{aligned}\right\} \tag{5.9}$$

then, again, under the assumptions of the theory, there will always survive some portion of undisturbed spherical shock which will propagate to infinity like the shock from the equivalent spherical charge.

Finally, if (5.1) is differentiated with respect to b, it is found that

$$b\frac{\partial\phi}{\partial b} = 2\left(\frac{d}{b}\right)^{\frac{1}{2}\alpha} - 4\left(\frac{d}{b}\right)^{\frac{3}{2}\alpha} + O\left(\frac{d}{b}\right)^{\frac{5}{2}\alpha}, \tag{5.10}$$

so that $\partial\phi/\partial b$ is positive. This means that for a given mass m of a particular explosive and a given cone angle 2β, ϕ will increase monotonically with b. It follows, therefore, that if b is not not less than a certain critical length b^* which makes $\phi = 0$ in (5.1) and is given by the equation

$$\beta - \frac{4}{\alpha}\left(\frac{d}{b^*}\right)^{\frac{1}{2}\alpha} + \frac{8}{3\alpha}\left(\frac{d}{b^*}\right)^{\frac{3}{2}\alpha} + O\left(\frac{d}{b^*}\right)^{\frac{5}{2}\alpha}, \tag{5.11}$$

then, again, under the assumptions of the theory, there will always survive some portion of undisturbed spherical shock which will propagate to infinity like the shock from the equivalent spherical charge.

This is an important result so far as the design of such an underwater explosive shock gun is concerned. For if we require a specified shock pressure at a specified distance, then the equivalent spherical charge can be obtained at once from (3.1), and if we are restricted to a certain charge mass m, the maximum cone angle 2β which will meet the requirements is specified by the relation (4.5). Equation (5.11) will then give the minimum slant height of cone b which will ensure that the shock will propagate in the required manner.

The solution of the equations (5.7), (5.9) and (5.11) giving the critical quantities β^*, m^*,and b^* may be obtained in series form. If we write

$$n = 3\alpha/(6+2\alpha), \quad \text{so that} \quad \alpha = 6n/(3-2n), \tag{5.12}$$

the series solutions are found to be

$$\beta^* = 2^{2-n}\alpha^{2n/3-1}\left(\frac{D}{b}\right)^n\left[1 - \frac{(20n-3)}{27}2^{1-2n}\alpha^{4n/3-1}\left(\frac{D}{b}\right)^{2n} + O\left(\frac{D}{b}\right)^{4n}\right], \tag{5.13}$$

$$m^* = 2^{2-6/n}\alpha^{3/n-2}\left(\frac{k}{A\gamma p_0}\right)^{-3/\alpha} b^3\beta^{3/n}\left[1 + \frac{(20n-3)}{12(3-2n)}\beta^2 + O(\beta^4)\right], \tag{5.14}$$

and

$$b^* = 2^{2/n-1}\alpha^{2/3-1/n}\beta^{-1/n}D\left[1 - \frac{(20n-3)}{36(3-2n)}\beta^2 + O(\beta^4)\right]. \tag{5.15}$$

6. Numerical calculations

To illustrate the significance of the results outlined in § 5, let us consider a specified $\frac{1}{12}$ oz. charge of TNT. If $(p-p_0)$ is measured in lb. wt./in.^{2}, R in feet and charge mass in pounds, the shock decay constants are $k = 2{\cdot}16 \times 10^4$, $\alpha = 1{\cdot}13$. Using these values it is found, from (4.7), that $D = 0{\cdot}0208$ ft.

The critical angle β^* as a function of b, or alternatively, b^* as a function of β obtainable from (5.7) and (5.11), or from the series expansions (5.13) and (5.15), is given in figure 2. The equivalent spherical charge masses W (in oz.) are also indicated along the curve.

The locus of the boundary of the undisturbed shock given by (4.3) and (4.4), or by the expansion (5.4), is shown in figure 3 for three values of cone semi-angle when $b = 1$ foot. The corresponding peak overpressures behind the shock at various distances are shown on each curve. The critical angle β^* in this case is $31° 53'$.

A similar curve is displayed in figure 4 for the case $b = 2$ ft, when the critical angle β^* is found to be $24° 5'$.

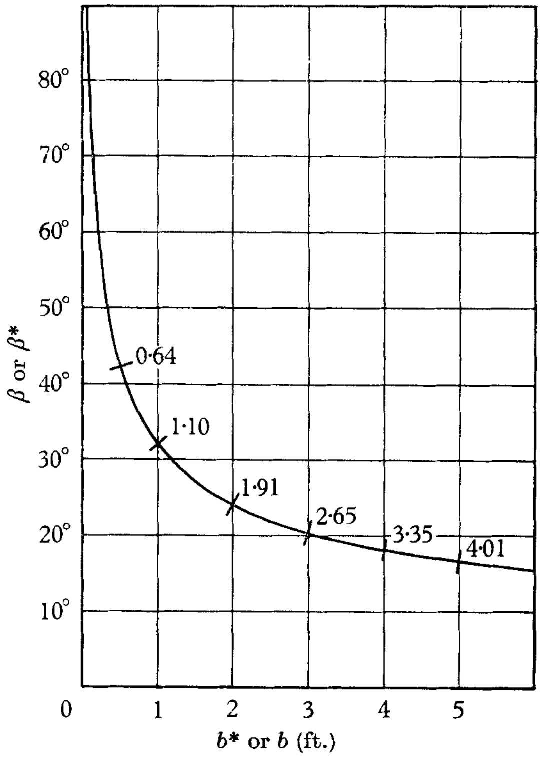

FIGURE 2. Charge: $\frac{1}{12}$ oz. of TNT. Equivalent spherical charge masses W (oz.) are indicated.

7. Conclusions

Whatever shape of chamber is used in the design of an underwater shock gun, it is apparent that the requirements for the emerging shock to propagate to comparatively large distances with the least possible attenuation, are that it shall emerge sufficiently weak and with appropriate directional properties to allow it to withstand diffraction effects.

A study of the diffraction of a spherical wave leaving a conical chamber has illustrated the main features of the problem. The results could be verified experimentally as follows. To obtain a shock beam from a mass m of explosive charge, whose peak pressure is the same as that of a shock propagated from a spherical charge mass W, the cone angle β is determined at once by (4.5) and the minimum slant height of cone b^* is given by (5.15). The theoretical beam width $(2y)$ at any distance is then given by (5.4).

It must be stressed that the theory given here determines only how the shock front survives the diffraction process, and does not deal with the complete spherical shock wave behind the shock front. It is clear from figure 1 that even though the sector QQ' of the shock front is undisturbed, the rarefaction waves EQ can penetrate into the conical sector QOQ'. Thus, although a target, struck by the undisturbed shock front QQ' will feel the same initial peak pressure as if

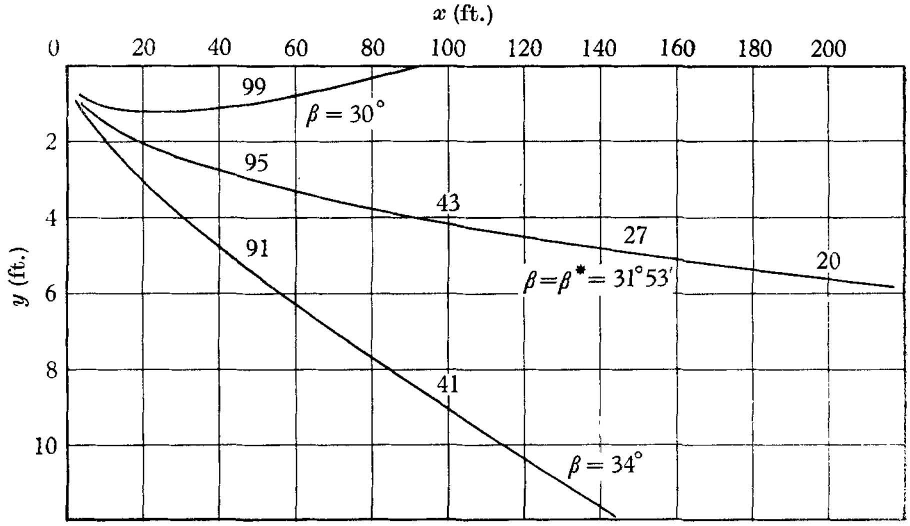

FIGURE 3. Charge: $\frac{1}{12}$ oz. of TNT, $b = 1$ foot. Peak overpressures in lb.wt./in.^{2} are indicated.

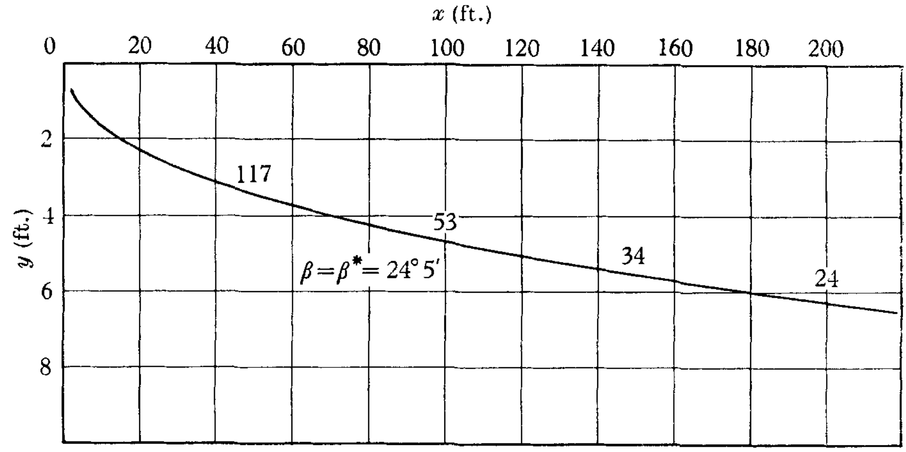

FIGURE 4. Charge: $\frac{1}{12}$ oz. of TNT, $b = 2$ ft. Peak overpressures in lb.wt./in.^{2} are indicated.

the complete spherical explosion had occurred, it will not experience the complete pressure-time pulse, but instead, a pulse which begins as this but is soon cut off by the expansion waves.

It is natural to ask whether similar results are obtained in the case of air. The corresponding theory has been developed for an ideal polytropic gas, and the calculations have been made for air ($\gamma = 1{\cdot}4$) using a suitable expression, instead of (3.1), to represent the peak over-pressure as a function of shock radius in a spherical explosion in air. The results obtained are very similar to those for water except that, ceteris paribus, the slant height of cone required in air, at sea-level ambient conditions, is very close to ten times as great as in water. Again, there is no information about the penetration of rarefaction waves behind the undisturbed shock front, and it is not possible to say, for instance, whether or not this is far more serious in air than in water.

The response of targets to shock waves from explosions depends, amongst other things, on the shape and duration of the pressure-time pulse behind the shock front and on the relative values of the fundamental pressures and times associated with the explosion and the distortion of the target. In some cases peak-pressure may be the dominant factor which determines whether or not a given degree of damage is caused, whilst in other cases the duration or the impulse per unit area of the wave (i.e. the area under the pressure-time curve) may become important. A general theory of the response of targets to (undisturbed) shock waves from spherical explosions in air has been given by Coombs & Thornhill (1960).

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