which bubble compression is not of a severe character (collapse), leading to the emission of an acoustical pulse, comparable in magnitude to the pulse produced at breakdown. To increase the degree of compression it is necessary to decrease the value of the gas content parameter, which can be done by increasing the external pressure.

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STRUCTURE OF A SHOCK WAVE FRONT IN A POROUS SOLID

S. Z. Dunin and V. V. Surkov

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The investigation of the nature of wave propagation in substances with a disruption of continuity is important for several reasons: The study of shock heating of a porous material in high-intensity waves makes it possible to deduce the equation of state of the continuous material under anomalous conditions (megabar pressures and temperatures of the order of the melting point) [1]; the majority of materials are not continuous in nature, and the wave propagation process is largely determined by the actual structure of the solid.

In the investigation of shock waves in solids with a disruption of continuity it is essential to take the following considerations into account. First, as in the analysis of a shock front in gases with retarded excitation of certain degrees of freedom [1], the structure of the shock transition in porous solids must be investigated with regard for the inertial properties of the medium [2-4]. Thus, in the shock loading of a porous solid to pressures in the tens of kilobars (such that the influence of heating of the substance can be neglected), the structure of the wave front is affected by the pore-selection dynamics [2-4]. An investigation of this type indicates that the pressure in the substance depends not only on the density of the substance, but also on its derivatives. Second, a number of theoretical and experimental studies [3-8] suggest an appreciable influence of the viscous properties of the porous material on the nature of the propagation and attenuation of shock waves. Third, estimates [2-4, 9] show that the porosity changes significantly only when the entire mass of the solid substance enters into the ductile state.

In the present study we discuss the characteristics of low-intensity shock wave propagation, where the influence of heating of the substance can be neglected (tens of kilobars), but the actual nature of wave propagation is largely determined by the behavior of the porous solid in the ductile state, viz., the pore-selection dynamics exerts a strong influence on the wave structure.

1. The shock profile is investigated in the example of a plane stationary wave propagating with velocity D. In this case all physical quantities (density, particle velocity, etc.) turn out to depend only on one variable $\zeta = x - Dt$, and the equations of mass and momentum conservation are easily integrated. Considering media of low porosity, we can neglect the dependence of the stress deviator on the porosity factor [10] and regard it as constant, with a value close to the yield point of the solid. Then in a coordinate system attached to the shock wave the equations are written in the form

$$\rho_0 D = \rho(D - v),\ p - p_0 = \rho_0 v D, \tag{1.1}$$

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where p, v, and ρ are the pressure, particle velocity, and density of the porous medium and depend on one variable ξ (the pressure p contains a viscous term proportional to $dv/d\xi$); p_0 is the pressure in the elastic medium; and we consider the initial density ρ_0 to be equal to its value in the undisturbed medium. To obtain the pressure–density relation in the plastic wave it is necessary to analyze the ductile-flow dynamics of the pores.

The medium is homogeneous in the initial state. We partition the entire substance into identical unit cells, each containing a single pore, and assume that all pores are spherical with radius a_0. It is practical to adopt a sphere as the equivalent unit cell. The initial cell radius b_0 must be such that the total mass of all the cells per unit mass is equal to unity, i.e., $4\pi N\rho_m(b_0^3 - a_0^3)/3 = 1$, were N is the number of cells per unit mass and ρ_m is the density of the solid (matrix). The volume variations of the isolated cell at the leading shock front is characterized by variation of the porosity of the medium.

We consider shock waves for which the width Δ of the leading front is much greater than the characteristic cell dimension λ. Then the relative variation of the macroscopic parameters of the medium (such as the average density or particle velocity) over the space scale λ of the isolated cell is of the order $(\lambda/\Delta) \ll 1$. It may be assumed in this approximation that the given cell participates in two independent motions at the wave front; as a unit whole with the particle velocity v of the medium and subject to compaction under the action of the pressure in the wave.

We define the macroscopic porosity parameter α as the ratio of the total volume of the isolated unit cell with coordinates x, t to the volume of the solid within the unit cell. Assuming that the stress distribution near the center of the pore remains spherically symmetrical during the deformation process and that the cell retains a shape close to a hollow sphere, we obtain

$$\alpha = b^3/(b^3 - a^3), \tag{1.2}$$

where b and a are the outside and inside radii of the equivalent sphere at point x at time t. Now in the coordinate system attached to the center of the pore the equation of motion of the medium describing the compaction of the cell is written in the form

$$\rho_m du/dt = \partial\sigma_r/\partial r + 2(\sigma_r - \sigma_\theta)/r, \tag{1.3}$$

where u is the particle velocity of the medium toward the center of the pore and σ_r and $\sigma_\theta = \sigma_\varphi$ are the radial and tangential components of the local stress tensor. The following condition holds on the surface of the pore:

$$\sigma_r|_{r=a} = 0.$$

The density ρ_m of the solid medium is constant in the investigated pressure range. Then the variation of the density ρ of the medium is solely attributable to variation of the porosity. It follows from (1.2) that

$$\rho = \rho_m/\alpha. \tag{1.4}$$

The relationship between the initial position of point r_0 and its coordinate r at time t can be derived from the conditions of mass conservation of the cell and incompressibility of the solid component of the medium:

$$r^3 - r_0^3 = a^3 - a_0^3, \quad b^3 - a^3 = b_0^3 - a_0^3. \tag{1.5}$$

Differentiating the first relation in (1.5) with respect to the time, we obtain the particle velocity from (1.2) and (1.5):

$$u = \dot{r} = \frac{a_0^3\dot{\alpha}}{3(\alpha_0 - 1)r^2}. \tag{1.6}$$

The solid phase of the medium obeys the flow condition for viscoplastic media [11]:

$$\sigma_r - \sigma_\theta = Y + 2\eta(\partial u/\partial r - u/r), \tag{1.7}$$

where Y and η are the yield point and viscosity coefficient of the solid. Applying expressions (1.6) and (1.7) and the fact that the pressure p_m in the solid component (matrix) of the medium is described by the equation

$$p_m = -(\sigma_r + 2\sigma_\theta)/3,$$

we rewrite Eq. (1.3) in the form

$$\rho_m du/dt = -\partial p_m/\partial r + 2Y/r. \tag{1.8}$$

The viscosity coefficient in this case enters only into the boundary conditions:

$$p_m = \frac{2Y}{3} + \frac{4\eta}{3}\left(\frac{\partial u}{\partial r} - \frac{u}{r}\right)\Big|_{r=a}. \tag{1.9}$$

We integrate Eq. (1.8) over the radius from a to r with regard for expressions (1.4), (1.6) and condition (1.9). As a result, we have

$$p_m(r, t) = \frac{2Y}{3} + 2Y\ln\frac{r}{a} + \frac{4\eta\rho_m\dot{\rho}}{3\rho(\rho_m - \rho)} - \frac{\rho_m^2 a_0^2}{3(\alpha_0 - 1)\rho^2}\left[\left(-\ddot{\rho} + \frac{2\dot{\rho}^2}{\rho}\right)\left(\frac{1}{a} - \frac{1}{r}\right) - \frac{a_0^3\rho_m\dot{\rho}^2}{6(\alpha_0 - 1)\rho^2}\left(\frac{1}{a^4} - \frac{1}{r^4}\right)\right]. \tag{1.10}$$

Averaging the pressure distribution (1.10) in the vicinity of the pore over the volume of the unit cell and using relations (1.2) together with the condition of incompressibility of the solid, we obtain

$$p = p_s(\rho) + p_v(\rho, \dot{\rho}) + p_d(\rho, \dot{\rho}, \ddot{\rho}), \tag{1.11}$$

where

$$p_s(\rho) = \frac{2Y}{3}\ln\frac{\rho_m}{\rho_m - \rho}; \qquad p_v(\rho, \dot{\rho}) = \frac{4\eta\dot{\rho}}{3(\rho_m - \rho)};$$

$$p_d(\rho, \dot{\rho}, \ddot{\rho}) = \frac{\rho_m a_0^2}{(\alpha_0 - 1)^{2/3}\rho^{8/3}}\left[\left\{\frac{\rho_m^{2/3} - (\rho_m - \rho)^{2/3}}{2} - \frac{\rho}{3(\rho_m - \rho)^{1/3}}\right\}(-\rho\ddot{\rho} + 2\dot{\rho}^2) + \left\{\frac{1}{\rho_m^{1/3}} - \frac{1}{(\rho_m - \rho)^{1/3}} + \frac{\rho}{3(\rho_m - \rho)^{4/3}}\right\}\frac{\dot{\rho}^2\rho_m}{6}\right].$$

The first two terms in (1.11) coincide with the corresponding terms obtained in [3, 4], but the last term differs insofar as the authors of the cited works used the static relation between the average pressure and the pressure on the surface of the cell (in dynamic analysis this relation must take account of the dynamics of the behavior of the porous medium).

The expression for p can also be augmented with a term accounting for the viscous resistance of the medium during motion of the cell as a unit whole. Estimates show that this term is of an order of magnitude $1/(\alpha - 1)$ smaller than the term describing viscous friction associated with motion of the substance toward the center of the pore.

2. Relations (1.1) and (1.11), in which it is required to transform to the variable ζ, describe the profile of a stationary plastic wave. The elastic wave pressure corresponding to transition of the substance into the ductile state is

$$p_0 = \frac{2Y}{3}\ln\frac{\rho_m}{\rho_m - \rho_0}.$$

The equation for the structure of the shock front is conveniently written in the following dimensionless form, which is solvable for the function α:

$$\frac{1}{(\alpha_0 - 1)^{2/3}}\left\{\left[\frac{\alpha^{2/3} - (\alpha - 1)^{2/3}}{2} - \frac{1}{3(\alpha - 1)^{1/3}}\right]\frac{\alpha''}{\alpha} + \left[\frac{1}{\alpha^{1/3}} - \frac{1}{(\alpha - 1)^{1/3}} + \frac{1}{3(\alpha - 1)^{4/3}}\right]\frac{\alpha'^2}{6\alpha}\right\}$$
$$= \frac{\alpha_0 - \alpha}{\alpha_0^2} + \frac{2k^2}{3}\ln\frac{\alpha_0(\alpha - 1)}{\alpha(\alpha_0 - 1)} - \frac{4kR\alpha'}{3\alpha(\alpha - 1)}, \tag{2.1}$$

where we have introduced the dimensionless variable $\xi = \zeta/a_0$ (the prime signifies differentiation with respect to ξ) and the dimensionless parameters k and R:

$$k = \frac{1}{D}\sqrt{\frac{Y}{\rho_m}}, \quad R = \frac{\eta}{a_0 l \sqrt{Y\rho_m}}.$$

The quantity R^{-1} is the analog of the Reynolds number for solid media.

For a wave propagating in the positive ξ direction the boundary equations for Eq. (2.1) have the form

$$\alpha' \to 0, \quad \alpha \to \alpha_0 \quad \text{as} \quad \xi \to +\infty. \tag{2.2}$$

The solutions of (2.1) that do not pass through the singular point $\alpha = 1$, $\alpha' = 0$ describe, in the general case, nonlinear damped or periodic oscillations. It is essential to note that, owing to the irreversibility of the stress–strain diagram for porous solids [12], expressions (1.11) and (2.1) are applicable only in the loading phase, when $\alpha' > 0$ and $\rho' < 0$. During unloading (stress relief), the medium behaves as an elastic or elasto-plastic medium (in which case the ductile flow region is very small [12]) and is not described by relations (1.11) and (2.1). Below, we investigate the behavior of the solution up to the first turning point (satisfying the condition $\alpha' = 0$) on the integral curve describing the solution of (2.1).

3. We analyze separately the influence of the parameters R and k on the shock profile. If the viscosity of the solid phase is close to zero, we can put R = 0. Equation (2.1) is integrable at once in this case. Taking (2.2) into account, we obtain

$$\alpha'^2 = \frac{2(\alpha_0 - 1)^{2/3}}{3}\left[\frac{(\alpha_0 - \alpha)^2(2\alpha + \alpha_0)}{2\alpha_0^2} - k^2\left\{\alpha_0 - \alpha + \ln\frac{\alpha_0 - 1}{\alpha - 1} + \alpha^2 \ln\frac{\alpha_0(\alpha - 1)}{\alpha(\alpha_0 - 1)}\right\}\right]\Big/ h(\alpha), \tag{3.1}$$

$$h(\alpha) = (\alpha - 1)^{-1/3}/3 + \left[(\alpha - 1)^{2/3} - \alpha^{2/3}\right]/2.$$

The denominator in (3.1) is greater than zero for any values of $\alpha > 1$, and so expression (3.1) is meaningful when the numerator is also greater than zero. In particular, the requirement that $\alpha'^2 > 0$ as $\alpha \to \alpha_0$ yields the condition

$$k^2 \leqslant k_0^2 = \frac{3(\alpha_0 - 1)}{2\alpha_0} \quad \text{or} \quad D \geqslant D_{\min} = \sqrt{\frac{2Y\alpha_0}{3\rho_m(\alpha_0 - 1)}}. \tag{3.2}$$

The quantity $D_{\min}$ is the miminum velocity of propagation of shock waves in the investigated porous media [13].

If $\alpha = \alpha_1 > 1$ at the turning point, then the vanishing of α' is equivalent to vanishing of the numerator in (3.1). Consequently, zero-valuedness of the numerator in (3.1) determines the relationship between k and the quantity α at the turning point and, with regard for relations (1.1) and (1.4), the curve of maximum deviations of the density and pressure for R = 0 from their initial values. A comparison of the latter curve with the static compression curve shows that for a given pressure amplitude the minimum porosity α_1 turns out to be smaller in the dynamic case. The pores are compressed at the shock front in such a way as to make their radius smaller than the equilibrium value obtained as the result of static compression.

We now consider a shock wave of moderate intensity, in which the variations of the density and porosity are small, $(\alpha_0 - \alpha_1) \ll 1$. Knowing that α' vanishes twice for $\alpha = \alpha_0$ and $\alpha = \alpha_1$, we make use of the smallness parameter, expanding the right-hand side of expression (3.1), first with respect to $(\alpha_0 - \alpha)$ and then with respect to $(\alpha - \alpha_1)$:

$$\alpha'^2 = \frac{2}{3}(\alpha_0 - 1)^{2/3}\left[\frac{1}{\alpha_0^2} + \frac{k^2}{(\alpha_1 - 1)^2}\right](\alpha_0 - \alpha)^2(\alpha - \alpha_1)/h(\alpha_1). \tag{3.3}$$

Now the following relation, deduced from the condition $\alpha' = 0$ for $\alpha = \alpha_1$, must be satisfied:

$$\frac{2\alpha_1 + \alpha_0}{2\alpha_0^2} - \frac{k^2}{(\alpha_1 - 1)} = 0 \quad \text{or} \quad D = \sqrt{\frac{2Y\alpha_0^2}{\rho_m(\alpha_1 - 1)(2\alpha_1 + \alpha_0)}} \geqslant D_{\min}. \tag{3.4}$$

The presence of the singular point $\alpha = 1$ restricts the region of convergence of the series in expansion with respect to the small parameter. Therefore, the given approximation is valid under the condition

$$(\alpha_0 - \alpha_1)/(\alpha_1 - 1) < 1 \quad \text{or} \quad \alpha_0 + 1 < 2\alpha_1. \tag{3.5}$$

Integrating (3.3) and setting the constant of integration, which determines the position of the reference point, equal to zero, we obtain the equation for the wave profile in the form

$$\alpha - \alpha_1 = (\alpha_0 - \alpha_1)\,\mathrm{th}^2(\zeta/\Delta), \tag{3.6}$$

where

$$\Delta = \frac{2a_0\alpha_0}{(\alpha_0 - 1)^{1/3}}\sqrt{\frac{3(\alpha_1 - 1)h(\alpha_1)}{(\alpha_0 + 4\alpha_1 - 2)(\alpha_0 - \alpha_1)}}.$$

Equation (3.6) describes a solitary symmetrical wave. However, the solution is applicable only in the loading phase in the interval of values of ζ from 0 to ∞. For $\zeta < 0$ and up to the next turning point, the pressure–density relation is determined by the elastic behavior of the porous substance. Thus, for negative values of ζ the wave profile has a complex oscillatory character. The effective width Δ of the weak shock front turns out to be independent of the properties of the solid phase and is determined solely by the geometry of the pore space. The initial requirement $\Delta \gg \lambda > a_0$ is satisfied under the condition $\sqrt{\alpha_0 - \alpha_1} \ll 1$.

4. We now investigate the case in which the viscosity of the solid medium is large and the inertial terms in (2.1) can be neglected. The equation for the structure of the wave front can be obtained formally from (2.1) by letting a_0 tend to zero. As a result, we have

$$\frac{4k\bar{R}}{3\alpha(\alpha - 1)}\frac{d\alpha}{d\zeta} = \frac{\alpha_0 - \alpha}{\alpha_0^2} + \frac{2k^2}{3}\ln\frac{\alpha_0(\alpha - 1)}{\alpha(\alpha_0 - 1)}, \tag{4.1}$$

where $\bar{R} = \eta/\sqrt{\rho_m Y}$. Making use of the fact that $d\alpha/d\zeta > 0$ at the loading wave front, we infer that $D \geq D_{min}$, where D_{min} is given by expression (3.2). We solve Eq. (4.1) for a weak shock wave. Expanding (4.1) with respect to the small parameter as in Sec. 3, we obtain the expression

$$\frac{d\alpha}{d\zeta} = \frac{(\alpha_0 - \alpha)(\alpha - \alpha_1)(2\alpha_1 - 1)\sqrt{3}}{2\bar{R}\alpha_0\sqrt{2\alpha_1(\alpha_1 - 1)}}. \tag{4.2}$$

Relations analogous to (3.4) and (3.5) must be satisfied in this case. Integrating (4.2) and setting the constant of integration equal to zero, we find

$$\alpha = \frac{\alpha_0 \exp(\zeta/\Delta) + \alpha_1}{\exp(\zeta/\Delta) + 1}, \tag{4.3}$$

where

$$\Delta = \frac{2\bar{R}\alpha_0\sqrt{2\alpha_1(\alpha_1 - 1)}}{(\alpha_0 - \alpha_1)(2\alpha_1 - 1)\sqrt{3}}.$$

The minimum porosity α_1 is attained for $\zeta \to -\infty$. Oscillations at the front are absent, and the wave profile has a monotonic character. The shock adiabat for the substance coincides with the static compression curve. The condition $\Delta \gg \lambda$ holds if $\bar{R}/(\alpha_0 - \alpha_1) \gg b_0$.

5. We determine the asymptotic behavior of the solution of (2.1) for a weak shock wave, taking inertial and viscous terms into account. We make an order-of-magnitude comparison of the terms entering into the inertial component. Inasmuch as $\alpha'' \sim (\alpha_0 - \alpha)/\xi^2$, $\alpha'^2 \sim (\alpha_0 - \alpha)^2/\xi^2$, and for a weak shock wave $(\alpha_0 - \alpha) \ll 1$, we can neglect the second term, which is proportional to α'^2, in comparison with the first. Then, expanding the coefficients of the derivatives and the free term and introducing the variable $\gamma = \alpha - \alpha_0$, we arrive at the asymptotic relation

$$f(\alpha_0)\gamma'' - kRg(\alpha_0)\gamma' - (1 - k^2/k_0^2)(\gamma/\alpha_0) = 0, \tag{5.1}$$

where

$$f(\alpha_0) = (\alpha_0 - 1)^{-2/3}h(\alpha_0); \quad g(\alpha_0) = (4/3)(\alpha_0 - 1)^{-1}.$$

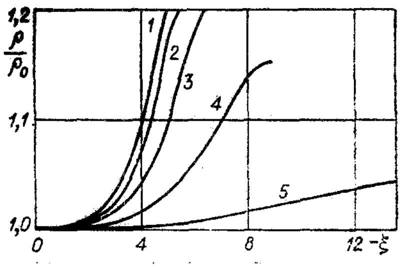

Fig. 1

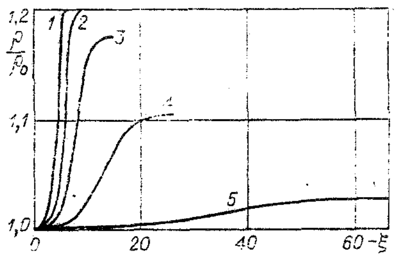

Fig. 2

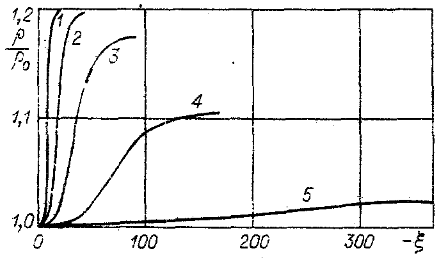

Fig. 3

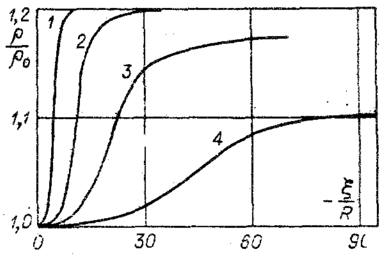

Fig. 4

For $\alpha_0 > 1$ the functions f and g are strictly greater than zero. Solving Eq. (5.1) subject to the boundary conditions (2.2), we obtain

$$\alpha_0 - \alpha = c \exp(-\zeta/\Delta), \tag{5.2}$$

where

$$\Delta = \frac{2a_0 f(\alpha_0)}{-kRg(\alpha_0) + \sqrt{[kRg(\alpha_0)]^2 + (1 - k^2/k_0^2)\, 4f(\alpha_0)/\alpha_0}};$$

c is a constant of integration.

We analyze two limiting cases. If the viscosity of the solid substance is small, we can expand (5.2) with respect to the small parameter R to obtain an expression for the characteristic space scale Δ in which the amplitude undergoes an e-fold variation:

$$\Delta = a_0 \sqrt{\frac{\alpha_0 f(\alpha_0)}{(1 - k^2/k_0^2)}} \left\{1 + \frac{kRg(\alpha_0)}{2\sqrt{f(\alpha_0)(1 - k^2/k_0^2)/\alpha_0}}\right\}. \tag{5.3}$$

Equation (5.3) is consistent with the results obtained in Sec. 3 for the case R = 0.

If the parameter R is not small, then relation (5.2) can be simplified on the basis of the fact that for a weak wave $k \sim k_0$. As a result, we have

$$\Delta = \frac{\alpha_0 g(\alpha_0) k\bar{R}}{(1 - k^2/k_0^2)} \left\{1 + \frac{(1 - k^2/k_0^2) f(\alpha_0)}{\alpha_0 g^2(\alpha_0) k^2 R^2}\right\}.$$

The space scale Δ turns out to be practically independent of the pore radius a_0 and is determined mainly by the parameter $\bar{R}$ (see Sec. 4).

6. Equation (2.1) is integrated numerically in the general case. The results of calculations of the shock profiles for certain values of the parameters R and k are given in Figs. 1-4. The initial porosity of the substance is $\alpha_0 = 1.2$, for which $k_0 = 0.5$. Curves 1-5 in Figs. 1-4 correspond to k = 0.1, 0.2, 0.3, 0.4, and ~~0.475~~ 0.475. For R = 0 (Fig. 1) the width of the shock front is of the order of a few pore radii a_0. The wave front in Figs. 2 and 3, corresponding to R = 0.24 and 1.6, is longer. Curves 5 are well described in this case of expression (4.3) obtained for a weak shock with neglect of the inertial terms in Eq. (2.1).

Figure 4 gives the results of the calculations for $R > 10$. In this interval of R the inertial terms can be neglected everywhere except in the small region where $\alpha \to 1$, in which case the wave profile is described with sufficient accuracy by expression (4.1). The parameter R for real media varies between very wide limits, roughly from 10^{-2} to 10^3-10^4. This result is attributable to the fact that the viscosity η and characteristic dimension of the pores for various substances, according to the experimental data [4-7], differ by several orders of magnitude.

7. The foregoing analysis of the structure of the shock front in a viscoplastic porous medium with allowance for the pore-flow dynamics indicates that the fundamental laws governing the propagation of shock waves are determined by the complex dependence of the pressure on the density and its derivatives. This dependence is obtained from an analysis of the dynamic behavior of unit cells of the medium containing the pores at the wave front.

The width of the front and the profile of the wave depend on three dimensionless parameters: α_0, $k = \sqrt{Y/(\rho_m D^2)}$, $R = \eta/(a_0\sqrt{Y\rho_m})$, where the existence of plastic shock waves is possible only under the condition $k \le k_0$ or $D \ge D_{min}$.

The investigation of the analytical solution describing a weak shock wave with viscosity neglected shows that the width of the front is $\sim a_0/\sqrt{\alpha_0 - \alpha_1}$ and is determined mainly by the geometry of the pore space. If the dynamic terms can be neglected in the equations, the characteristic width of the weak shock front depends on the number $\bar{R} = \eta/\sqrt{Y\rho_m}$ and is of the order of magnitude $\bar{R}/(\alpha_0 - \alpha_1)$.

The strong shock profile obtained by numerical methods depends largely on the shock velocity and on the properties of the porous medium (parameters R and α_0). For $R \gtrsim 10$ we neglect inertial effects in the ductile flow of the pores, and the structure of the shock front is determined mainly by the viscoplastic properties of the porous medium and the shock velocity.

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