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*Modeling of High-Explosive-Driven
Plasma Opening Switches*



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LA-9702-MS

UC-38

Issued: March 1983

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MODELING OF
HIGH-EXPLOSIVE-DRIVEN PLASMA OPENING SWITCHES

by

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ABSTRACT

High-explosive-driven plasma opening switches have been modeled in one dimension using the Lagrangian MHD code RAVEN. These calculations have been made in both cylindrical and planar geometry. Simple compression can account for observed resistance increases at early times (time-of-flight of the high-explosive detonation products across the plasma conducting channel). Our results suggest that some improvements in switch performance might be achieved through a judicious choice of gases in the plasma channel and by lowering the pressure in the channel.

I. INTRODUCTION

Explosive-driven magnetic flux compression generators are being used to produce large currents to vaporize, ionize, and implode cylindrical metal foils. With foil masses and dimensions of interest, the current rise-time from these generators is too long to insure efficient implosions. To optimize operation it is desirable to have a low-inductance, fast-opening switch that can interrupt many megamperes of current through a ballast inductor in a fraction of a microsecond.

In 1977 Pavlovskii et al.¹ reported on a high-explosive (HE)-driven opening switch with which they achieved a resistance change of 0.2Ω in $0.45 \mu\text{s}$, interrupting a current of 7.3 MA and switching 4 MA into a 30-nH load. The qualitative circuit, as pictured by Pavlovskii et al., is shown in Fig. 1. From the dimensions cited in Ref. 1, we calculate a current density of 0.12 MA/cm for the Pavlovskii switch.

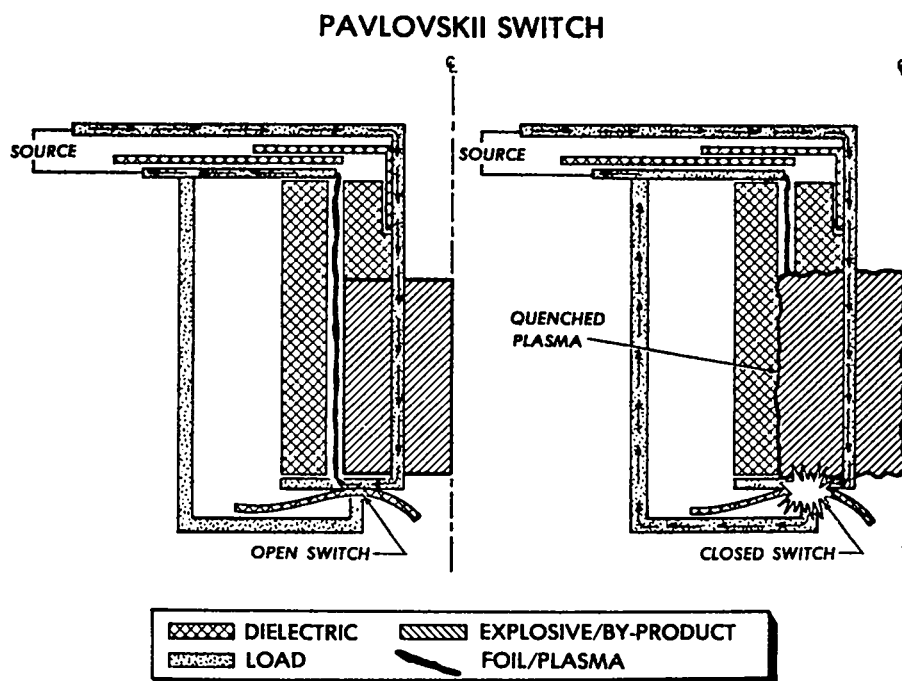


Fig. 1. Qualitative depiction of the Pavlovskii switch before and during switching. The high explosive expands to close off the plasma channel and blow through the dielectric at the bottom to make an alternate current path available.

There have been several experimental efforts in the United States with systems that are comparable to Pavlovskii's, but thus far they have not achieved results that are as good. Turman and Tucker² report that with a cylindrical system 4.5 cm in diameter and 3 cm in length they achieved resistance increases of as much as 30 to 40 m Ω with currents of as much as 2 MA (or current densities of 0.15 MA/cm). More recently Turman, Tucker, and Skogmo³ have reported resistance increases of nearly 90 m Ω in a cylindrical switch that is 20 cm in diameter, 10 cm in length, and carrying a current density of 0.06 MA/cm. Goforth, at the Los Alamos National Laboratory, reports resistance increases of 40 to 60 m Ω using a planar geometry switch.⁴ The foil in his system was 3.2 by 5 cm and expanded into a gap 0.3-cm high. He has used currents as high as 1 MA or a current density of 0.4 MA/cm.

A number of physical processes are occurring simultaneously or nearly simultaneously in these switches. It is the purpose of the present study to isolate the effect of one-dimensional (1-D) compression. An earlier attempt to simulate these explosive opening switches is reported by Baker.⁵ From his study Baker finds that simple, 1-D compression of the current-carrying channel cannot explain the resistance increases that have been observed. Indeed, he predicts resistance increases of only 1 to 2 m Ω . Baker did not, however, carry his analysis to the planar geometry where the effects of divergence on the HE detonation might be less severe. Using a J-W-L equation of state for the HE PETN, Baker finds a pressure that is two orders of magnitude below the Chapman-Jouguet pressure at the time the plasma channel is beginning to compress. In addition, he does not discuss the role of direct cooling of the plasma channel by heat transfer to the HE detonation products. This heat transfer is frequently mentioned as an important step in the switching process.

II. THE MODEL

In the present work we report on our effort to use the 1-D MHD code RAVEN⁶ to model these explosive-driven opening switches in both cylindrical and planar geometry. RAVEN is a Lagrangian code that uses the Braginskii⁷ formalisms for electrical and thermal conductivities.

For this study it was necessary to add to RAVEN a high explosive equation of state. In this equation of state the fraction of the material that has burned in any zone is determined by a detonation velocity. The internal energy is determined from

$$\begin{aligned} E &= C_v T && \text{if not burning} \\ &= B_f E_r && \text{if burning,} \end{aligned}$$

where B_f is the fraction burned and E_r is the specific energy released, a quantity that is HE dependent. The pressure is determined from

$$\begin{aligned} p &\propto [V_0 - V(t)] && \text{before burning} \\ p &= B_f (C_p - C_v) T / V(t) && \text{during and after burning.} \end{aligned}$$

Values for the parameters that we used in our model are listed in Table I.

For electrical conductivity we have used 100 mho/m in the detonation products, the value cited by Pavlovskii. For thermal conductivity we have examined two options: using the values cited in various tables for the predominant detonation products (this often involved extrapolations beyond the temperatures listed) and using the formalism cited by Spitzer⁸ assuming the products to

TABLE I
HE PARAMETERS USED IN CALCULATIONS^a

	<u>Cyclotol</u>	<u>PBX 9501</u>	
Specific Heat (constant volume)	8.6×10^6	8.6×10^6	erg/(°K-g)
Specific Heat (constant pressure)	2.11×10^7	2.17×10^7	erg/(°K-g)
Internal Energy Released	9.2×10^{10}	1.02×10^{11}	erg/g
Specific Volume (solid density)	0.57	0.543	cm ³ /g
Detonation Velocity	8.25×10^5	8.8×10^5	cm/s

^aThe choice of cyclotol for the cylindrical geometry was arbitrary, the choice of PBX 9501 for the planar switch was to match Goforth's experiments.

be singly ionized. In most calculations we use the second approach, which gives slightly higher values. In fact, neither approach makes any difference in the calculations. We varied the values for thermal conductivity upward by a factor of one thousand without affecting the results.

The electrical circuit available in RAVEN at the time of this study is shown in Fig. 2. This circuit will not permit us to exactly simulate Pavlovskii's experiment. We can introduce an external load on the parallel circuit leg (L_p), but this inductance will not include the return conductor as it does in the Pavlovskii experimental set-up.

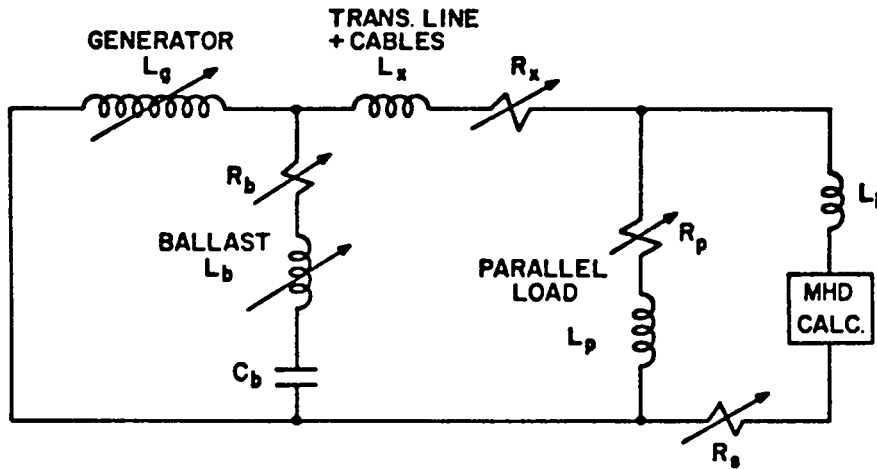


Fig. 2. The electrical circuit available in RAVEN. In our simulations we are using the generator (L_g), the ballast inductor ($L_b = 40$ nH), the transmission time ($L_x = 30$ nH), the load inductance ($L_1 = 1$ nH), and the parallel load. The parallel load is used to simulate the alternate current path so that R_p is dropped to zero at peak compression.

We have run cases in both cylindrical and planar geometry. The cylindrical system closely approximates the geometry of Pavlovskii. As is shown in Fig. 3, the cylinder has a radius of 12.5 cm. The plasma column has an initial radius of 10 cm and may expand into a 0.5-cm radius gap. The central conductor has a radius of 4 cm. We have taken the cylinder to be 10 cm long. Following the work of Baker,⁵ we note that after the aluminum foil has vaporized the aluminum represents little more than an impurity in the channel, which we therefore take to be atmospheric pressure oxygen at 2.5 eV, an arbitrary temperature that is high enough so that the oxygen will conduct.

Our planar geometry model is taken from the experimental work of Goforth.⁴ It is a 3-cm by 3-mm cavity that is 5 cm in length.

III. RESULTS AND CONCLUSIONS

The results of the cylindrical calculations are shown in Figs. 3 through 6. Results of the planar calculations are shown in Figs. 7 and 8. In these calculations we hold the outer insulator fixed in space and time. As a consequence, the results of the calculations after the initial compression are certainly not believable.

Figure 4 shows a resistance increase that is quite comparable to the initial rises reported by Turman and Tucker. This rise is due solely to the compression of the current-carrying channel. Figure 5 shows the temperature of the plasma as a function of time. Since the resistivity of the ionized plasma is inversely proportional to the temperature to the $3/2$ power, the sharp spike in the temperature leads to the rapid drop in the resistance once the cross-sectional area is no longer decreasing. Our resistance drop is sharper than that in the experimental results of Turman and Tucker, suggesting that our temperature increase may occur too rapidly. This would not be surprising since 1-D simulations tend to overestimate temperatures at the peak of compression. It should also be noted that our outer insulator is not permitted to move during the calculation and this may also lead to an overestimate of the rate of compression and the rate of temperature increase.

Our calculated secondary resistance increases are smaller than the first, and this does not agree with experiment. In addition, these secondary increases would probably not occur if the outer insulator were free to move. We suspect that later resistance increases are due more to mixing between the hot plasma and relatively cool detonation products than they are to compression. Such mixing is not included in our 1-D Lagrangian calculations.

Cell Boundaries

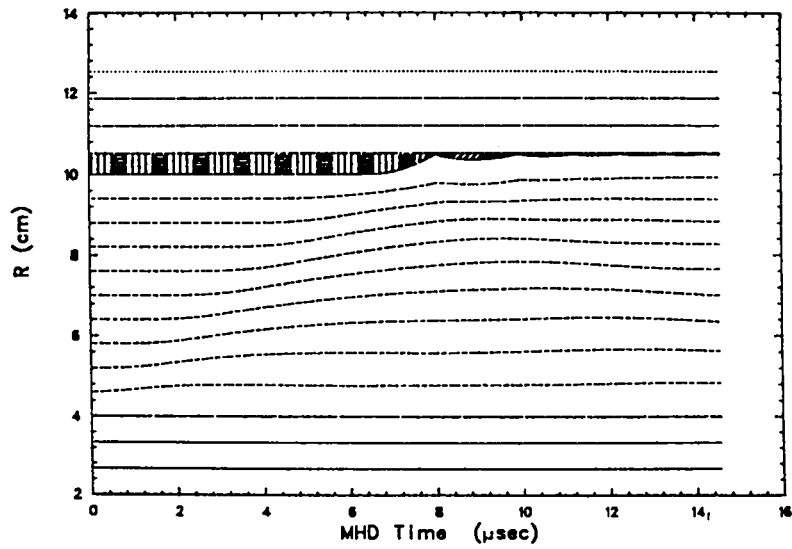


Fig. 3. Radii of each zone in the cylindrical calculation as a function of time. In each figure in which the zones are shown, the following code will apply:

——— central conductor -·-·-·- plasma channel
 ····· high explosive -·-·-·- external insulator.

Resistance

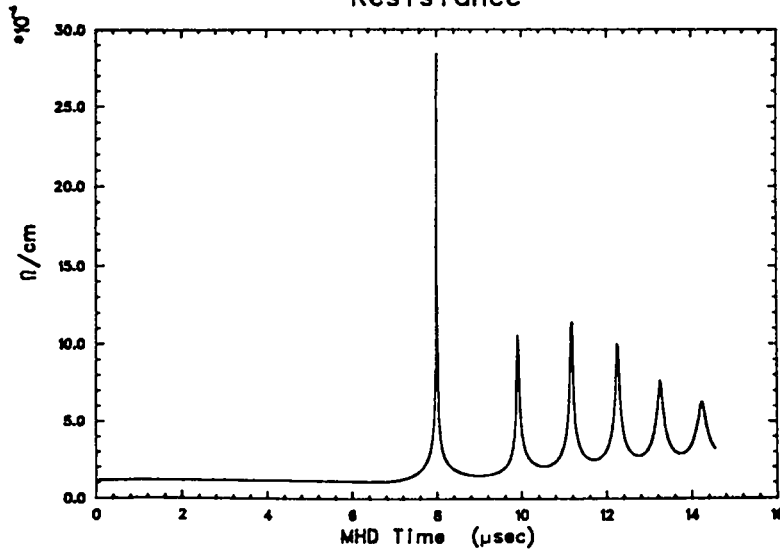


Fig. 4. Calculated resistance along the plasma channel as a function of time for the cylindrical geometry.

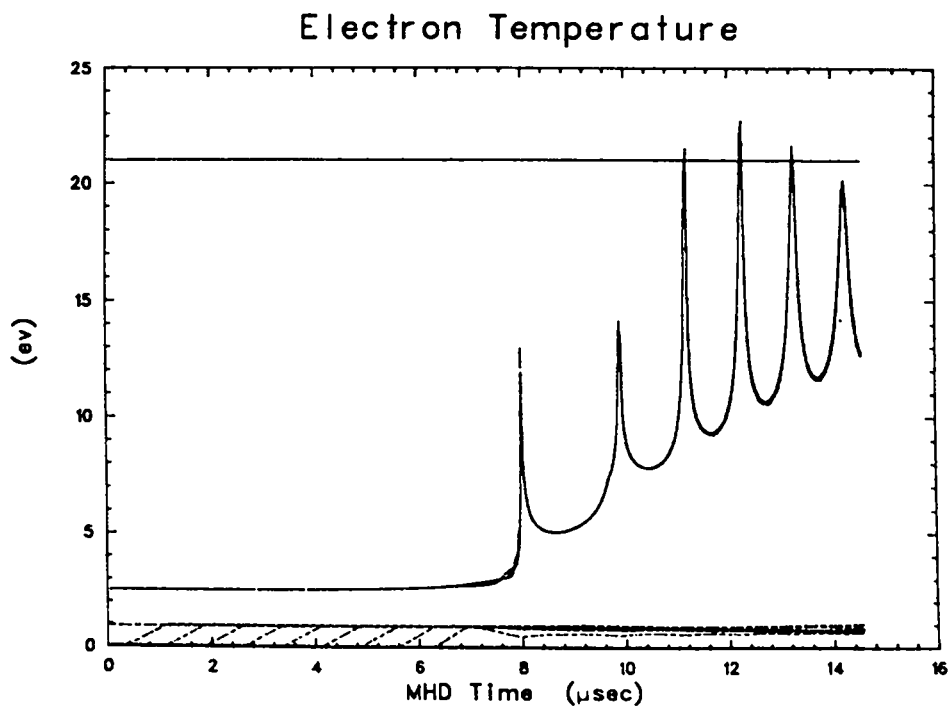


Fig. 5. Calculated temperatures for the cylindrical geometry as a function of time.

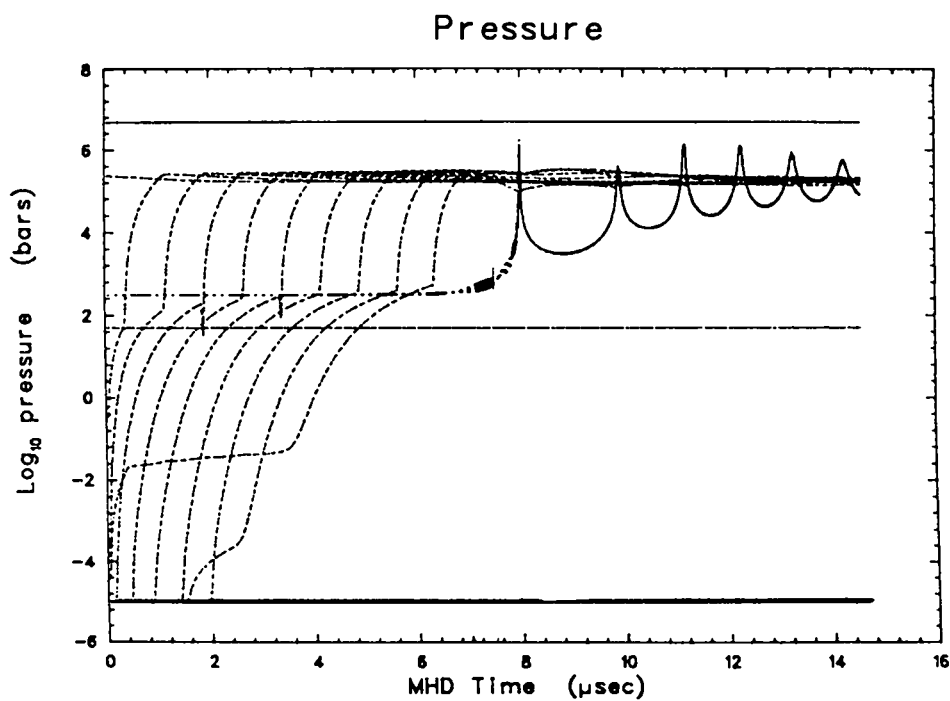


Fig. 6. Calculated pressures for the cylindrical geometry as a function of time.

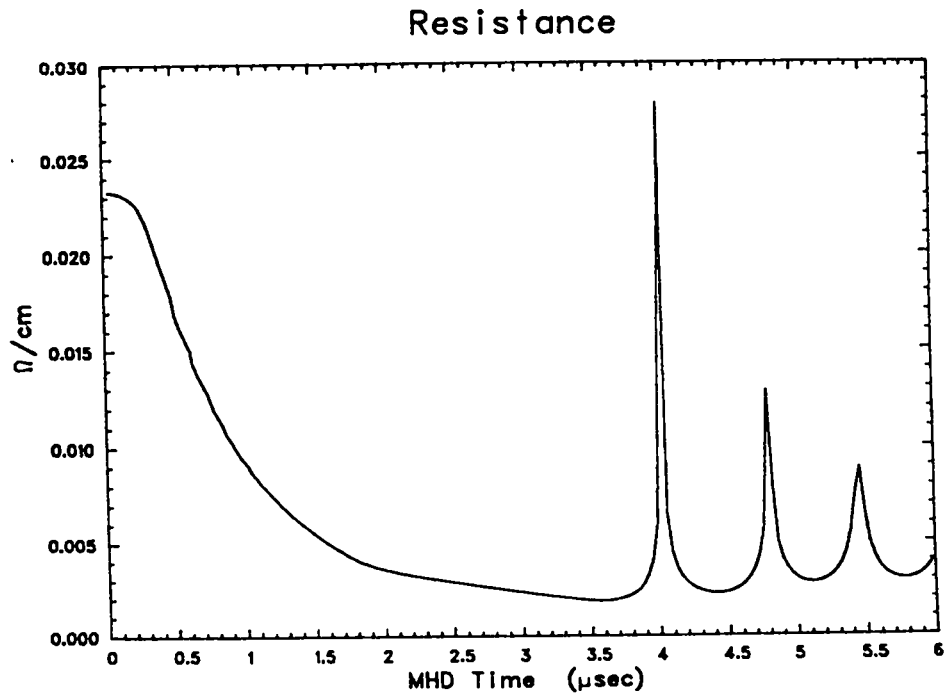


Fig. 7. Calculated resistance along the plasma channel as a function of time for planar geometry.

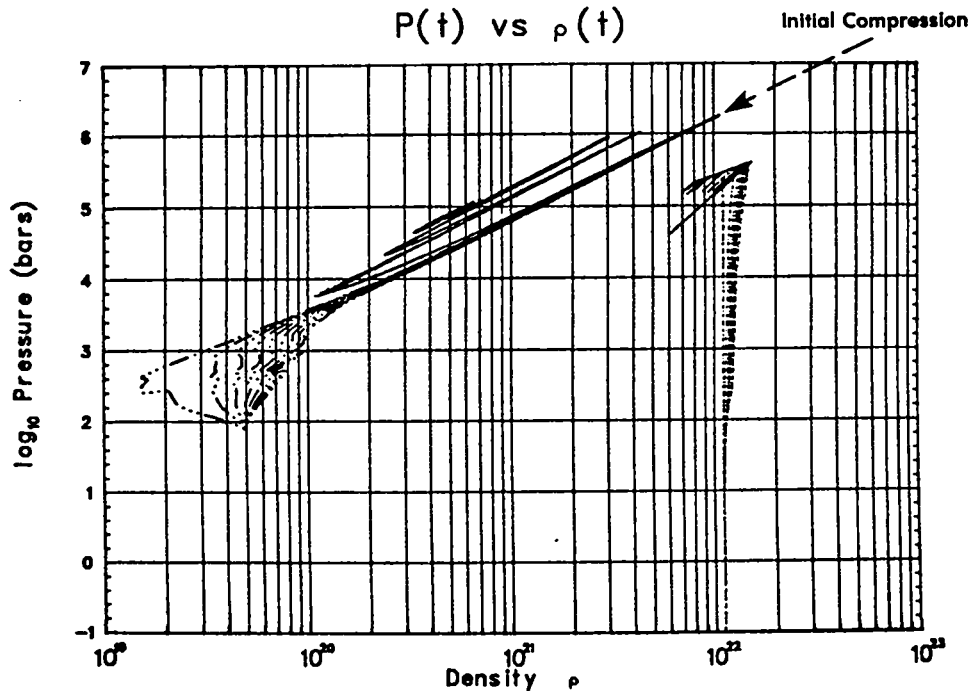


Fig. 8. Pressure vs density. For a Lagrangian code this is the equivalent of pressure vs (volume)⁻¹. Dashed line points to that portion of the plot that was traced during the initial compression.

We have considered the effects of thermal diffusion to both the HE detonation products and the outer insulator within the limit of the Lagrangian calculation and find that they have essentially no effect during the time scales that we are considering. For the external insulator we have used the thermal diffusion coefficient of Teflon, 2.6×10^4 erg-cm/(cm²-s-°K). For the HE detonation products, as mentioned, for most of our calculations we have used the formalism of Spitzer assuming the products are singly ionized. At 1 eV this gives a value of 1.4×10^5 erg-cm/(cm²-s-°K). We conclude that, without mixing, thermal diffusion to the insulator and the detonation products will not play an important role in increasing the resistance of these switches.

The calculated pressures in the cylindrical case are shown in Fig. 6. The pressures in the HE reach the Chapman-Jouget pressure of 320 kbars for cyclotol and deteriorate only slightly prior to compressing the plasma. These pressures in the HE are nearly two orders of magnitude higher than those calculated by Baker.⁵ We believe that essentially all of the other differences between our results and his stem from these pressure results.

In Fig. 7 we show the results for the planar geometry with a current density of 0.2 MA/cm. We initially found this resistance increase, which is significantly higher than that measured by Goforth, rather surprising. Our preliminary analysis indicated that these switches would be strongly limited by current density because the higher current, and subsequent Joule heating, would increase the ionization level of the plasma.

The resistance of one of these switches may be written

$$R = \frac{\eta l}{A} , \tag{1}$$

where η is the resistivity, A is the cross-sectional area, and l is the length. If we are dealing with an ionized plasma ($\bar{Z} > 1$),

$$\eta \propto T^{-3/2}, \quad (2)$$

where T is the temperature. We can also relate the cross-sectional area to the temperature if the compression is adiabatic:

$$PV^\gamma = \text{constant} \quad (3)$$

and if we make the ideal gas assumption

$$PV = nkT. \quad (4)$$

For an ideal gas $\gamma = 5/3$, so

$$V \propto T^{-3/2}. \quad (5)$$

Since the only portion of the volume that is changing is the cross-sectional area,

$$A \propto T^{-3/2}. \quad (6)$$

Equation (1) would, therefore, indicate that under the conditions of this analysis R would be a constant! Prior to compression in the planar case we find a temperature of nearly 10 eV. This temperature would result in an average ionization level, \bar{Z} , greater than 1.5.

One's first thought in response to the result shown in Fig. 7 is that the compression of the gas must not be adiabatic. However, Fig. 8 shows the plot of pressure versus density for this run, for a Lagrangian code this provides the same information as a PV plot. The dashed line on this plot indicates the portion that

corresponds to the time of the first compression. This is a reasonably straight line, indicating an adiabatic compression. However, the slope is not 1.67 but closer to 1.35. This lower value of γ indicates some sink of energy not included in our initial analysis. Additional thought indicates that this sink is likely to be ionization of the oxygen plasma.

Armed with this knowledge we can estimate γ based on the level of ionization. The internal energy can be written

$$U = 3/2 P + \sum_Z \epsilon_Z n_Z, \quad (7)$$

where ϵ_Z is the sum of the ionization energies up to the level z and n_Z is the number of ions in level z , that is,

$$\epsilon_Z = \sum_{z=0}^Z I_z. \quad (8)$$

We note, however, that we can write pressure

$$P \equiv (\gamma - 1)U/V, \quad (9)$$

where U is the internal energy and V is the volume. Therefore, simple substitution yields

$$\gamma - 1 = P/(3/2 P + \sum_Z \epsilon_Z n_Z), \quad (10)$$

so

$$\begin{aligned} \gamma &= 1 + 2/3 \left(\frac{1}{1 + 2/3 \frac{\sum \epsilon_Z n_Z}{P}} \right) \\ &= 1 + 2/3 \left[\frac{1}{1 + 2/3 \frac{\sum \epsilon_Z n_Z}{(n_i + n_e)kt}} \right], \end{aligned} \quad (11)$$

where n_i is the total ion population. Then in terms of \bar{z}

$$\gamma = 1 + \frac{2}{3} \left[\frac{1}{1 + \frac{2}{3} \frac{\sum \epsilon_z (n_z/n_i)}{(1 + \bar{z})kt}} \right]. \quad (12)$$

In the limits of no energy in ionization ($\epsilon_z = 0$) or very high temperature ($t \rightarrow \infty$), Eq. (12) limits to $\gamma = 5/3$ as it should. Between these limits the Saha equation must be solved based on the density and temperature in the plasma. We have solved the Saha equation iteratively for values of constant temperature and density to find the values of n_z and, hence, \bar{z} in oxygen and substituted these values into Eq. (12). The results of these calculations are shown in Fig. 9.

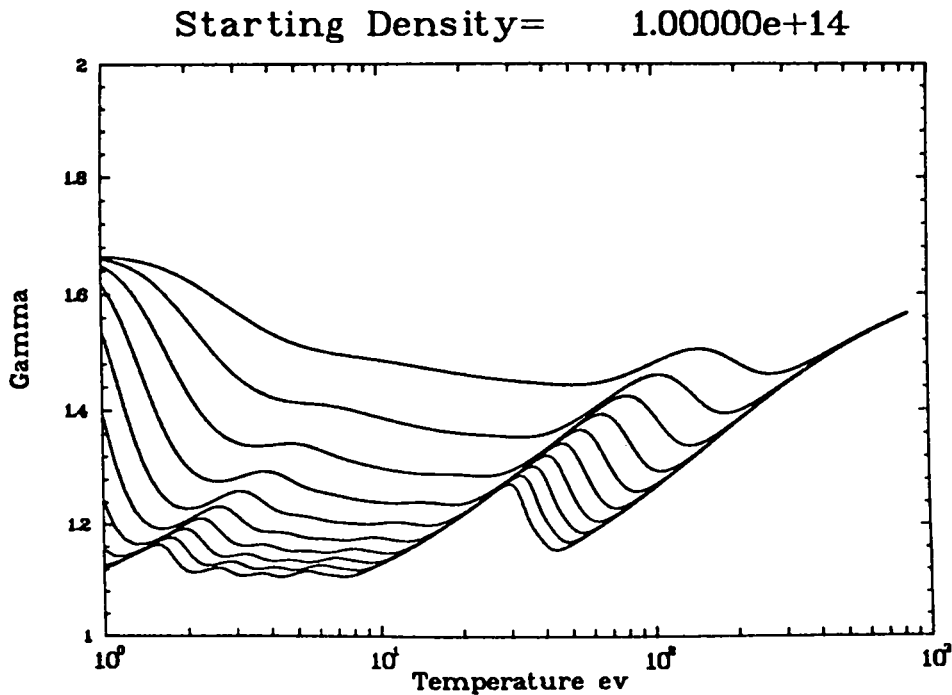


Fig. 9. Gamma as functions of temperature in oxygen. Each line is a factor of ten in density higher than the line below it. The bottom line is an ion density of 10^{14} cm⁻³, the top line is 10^{23} cm⁻³.

From Fig. 9 we reach three important conclusions. First, it is possible for these switches to be ionized ($\bar{Z} > 1$) and still have $\gamma < 1.67$, so that the resistance can increase, even if weakly, as a function of temperature. This should mean that these switches can be used at higher current densities, and hold these higher current densities for longer times, than we had initially anticipated. This conclusion appears to hold for a very wide range of temperatures. This temperature range is, however, a function of the gas in the plasma. Our second conclusion, therefore, is that something may be gained by carrying out an analysis for a variety of plasma gases in order to determine which gas will provide the best switch.

Finally, we note that although the performance of these switches in the compression regime should be relatively independent of the temperature prior to compression, over a wide temperature range, Fig. 9 predicts a strong dependence on pressure or density. This pressure dependence may be the most sensitive parameter available for improving the performance of these switches.

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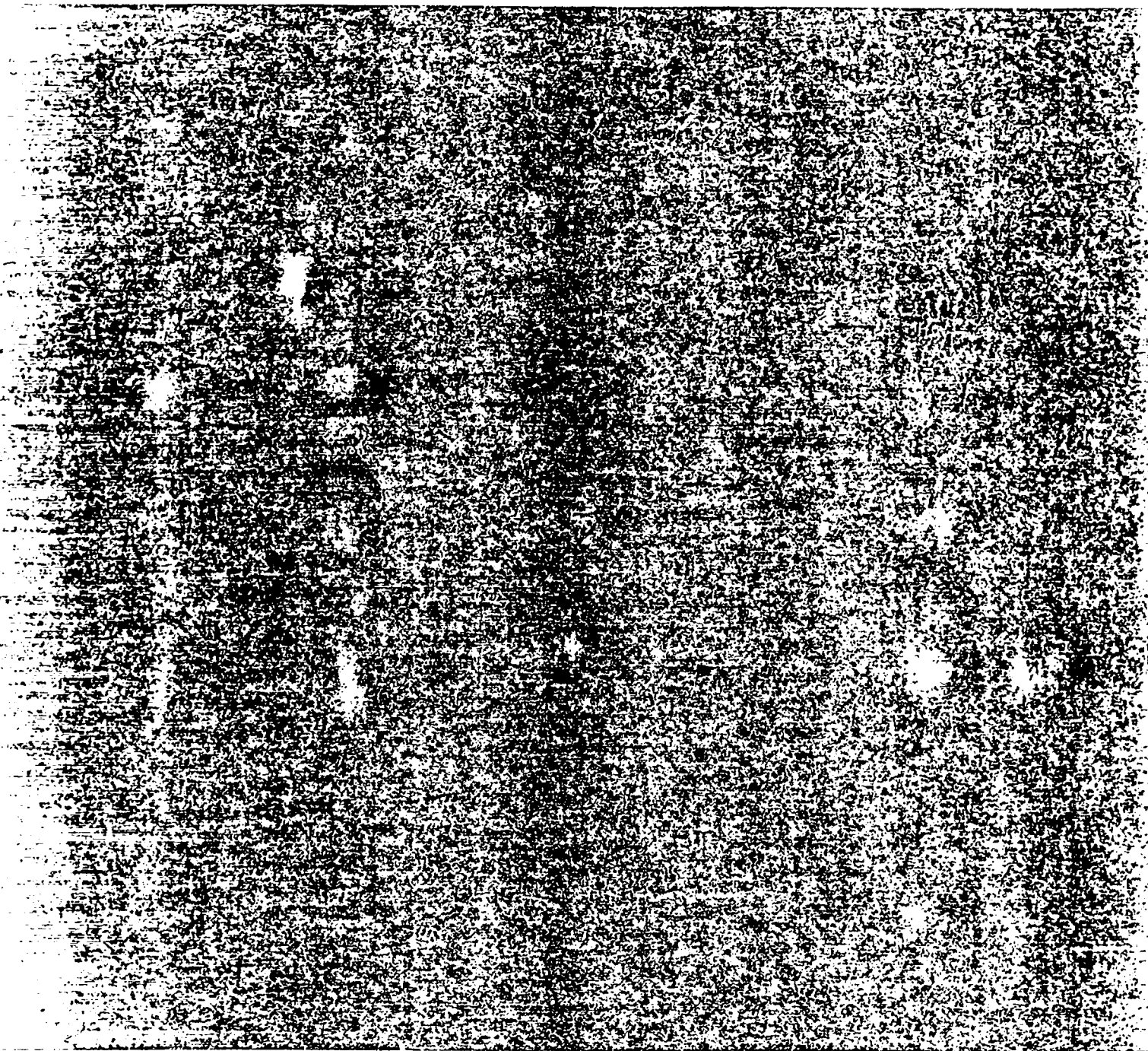
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