Burn Characteristics of Marginal Deuterium-Tritium Microspheres

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(Received 5 August 1974)

Long mean free paths for ions in the tail of the distribution may allow escape, quenching the burn of marginal ($\rho R < 10^{-2}$ g/cm^2) deuterium-tritium microspheres, possibly explaining the lack of success in experiments to date.

The thermonuclear burn of laser-fusion pellets may be usefully studied by consideration of the burn of hot compressed microspheres.[1] Such microspheres may be thought of as being the inner cores of deuterium-tritium (DT) droplets which have been compressed and heated by laser-driven ablation of the outer region of the pellets.[2-4] Or they may be thought of as the tamped gas fill in structured "microballoon" targets such as have been widely discussed in previously classified laser-fusion programs.[5] In these applications the microballoon is a hollow spherical shell of glass or metal, filled by diffusion at elevated temperature with DT gas. In this case the compression is from the initial gas density and may or may not bring the DT up to its normal liquid density. Such targets have been tested at Los Alamos Scientific Laboratory, KMS Industries, Inc., and other laboratories, but were not completely reported because of (previous) classification rules. In these tests the neutron yields fell below expectations.

From the study of Ref. 1, we know that the fractional burnup f_{ro} may be expected to follow from the fusion reactivity $\langle \sigma v \rangle$ as

$$f_{ro} = (\langle \sigma v \rangle / 8 C_s M_i) \rho R \tag{1}$$

for conditions in which $f_{ro} \ll 1$. In evaluating the expressions in Eq. (1), the density ρ, radius R, and sound speed $C_s = (2\gamma T/M_i)^{1/2}$ are evaluated at the initial conditions. We use $\gamma = \frac{5}{3}$ and M_i is the ion mass. For $\langle \sigma v \rangle$ we use the average over the Maxwellian ion distribution $f(v)$, integrated out to infinity[6]:

$$\langle \sigma v \rangle_\infty = \int_0^\infty v \sigma(v) f(v) v^2 \, dv . \tag{2}$$

Justification of the integration out to infinity requires that the mean free path for test ions at speed v, $\lambda(v)$, be smaller than dimensions of interest over ranges of v important to the integration. Unfortunately the cross section σ rises very rapidly with increasing v (below the maximum near 100 keV, deuteron energy) so that important ranges of v in the integration may correspond to energies well above the thermal energy in $f(v)$ and to very long mean free paths $\lambda(v)$.

At low energies the cross-section behavior is due mostly to Gamow's barrier-penetration formula[7]:

$$\sigma \sim \exp(-v^*/v)/v^2 ,$$

where $v^* = (2\pi)^2 e^2/h = 1.375 \times 10^9$ cm/sec. It is convenient to actually perform the integrations with a center-of-mass energy variable $E = M_R v^2/2$, while experimental data is usually reported in deuteron energy $E_D = M_D v^2/2$. Thus $E_D^* = 1.97 \times 10^6$ eV $= (1.403 \times 10^3)^2$ eV, while an empirical best fit is $(E_D^*)^{1/2} = A = 1.453 \times 10^3$ (eV)$^{1/2}$. Using this last value and $E_D = (M_D/M_R)E = \frac{5}{3}E$, we obtain

$$\sigma \sim \exp(-\sqrt{\tfrac{3}{5}} A/\sqrt{E})/E .$$

Then changing integration variable from dv to dE, the integrand in Eq. (2) is proportional to

$$K(E) = \exp(-\sqrt{\tfrac{3}{5}} A/\sqrt{E}) \exp(-E/kT) \tag{3}$$

which has its maximum at

$$E_m = [\, 0.5 \sqrt{\tfrac{3}{5}} A kT \,]^{2/3} , \tag{4}$$

which for $kT = 1$ keV is 6.8 keV.

The mean free path for 90° deflection for a test ion, on the other hand, is[8]

$$\lambda = M_i^2 v^4 / 8\pi n e^4 \ln\Lambda .$$

Substituting $E_c = M_i v^2/2$ and $n = \rho/M_i$, this expression can be evaluated as

$$E_c = 559 \text{ keV } (\ln\Lambda/10)^{1/2} (\rho\lambda)^{1/2} ,$$

where $\ln\Lambda = 10$ is typical, applying to a 1.0-μg microsphere of DT at $\rho R = 10^{-4}$ g cm^{-2} and $kT = 1$ keV. (That the total mass, the areal mass density ρR, and the temperature are the important initial-value parameters is clear from Ref. 1.) We suppose that test ions with $\lambda \sim R$ ought to diffuse quickly out of the hot core into the surrounding blow-off material or into the microballoon material. Such ions are then lost to the thermonuclear burn. In the case specified, we would expect to lose ions above 5.6 keV, while the maximum in the integral for $\langle \sigma v \rangle$ is found to

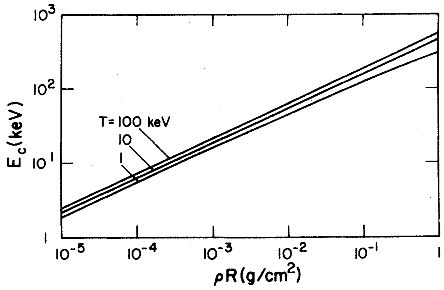

FIG. 1. Cutoff energy E_c for ion loss versus areal mass density ρR.

be at 6.8 keV. Thus, significant quenching is expected.

Rather than follow test ions in a Monte Carlo process, I have adopted $\lambda = R$ as the criteria for lost ions, and have truncated the Maxwellian tails at the corresponding cutoff E_c's in performing the $\langle \sigma v \rangle$ integration. A better treatment would also include slowing down of the test ions as well as reactions in flight by the lost ions, but our criteria appear to be adequate to estimate this important and overlooked factor in standard treatments. The cutoff E_c's obtained are shown in Fig. 1 for a 1.0-μg microsphere. The weak temperature dependence (in $\ln\Lambda$) is indicated; the still weaker mass (or density in $\ln\Lambda$) dependence is insignificant for 0.1- to 10.0-μg masses. It should be noted that for cases in which the cutoff E_c's do not greatly exceed the temperature T, then the fluid assumptions underlying Ref. 1 and this calculation are not valid. In these cases a proper Monte Carlo treatment is required. In cases in which these conditions are met, we need to consider the rate at which collisions are able to fill out the distribution function. This requires a Fokker-Planck calculation which is not in hand.

Using a best fit[9] to empirical data[6] for the cross section σ, I have numerically integrated $\langle \sigma v \rangle$ up to various cutoff energies.[10] It is observed that at high E_c the several curves saturate to the traditional $\langle \sigma v \rangle_\infty$ values, while at smaller E_c they fall rapidly and become approximately independent of temperature.[10] This approximate temperature independence is a remarkable and important result. In order to get a better understanding of the important values of E_c, than simply that from Eq. (4), I have plotted in Fig. 2 the E_c which reduces $\langle \sigma v \rangle$ to 0.5 and to 0.1 of its asymptotic value. The $(kT)^{2/3}$ dependence for small temperatures is apparent.

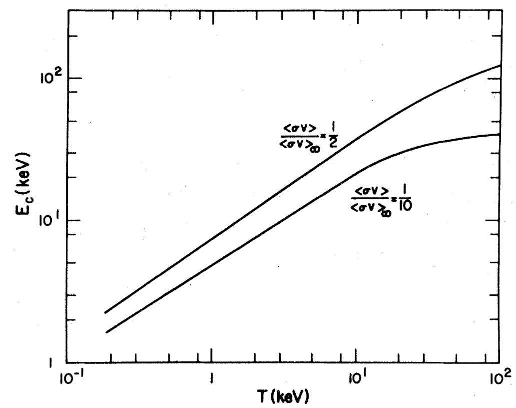

FIG. 2. Cutoff E_c required to reduce the reactivity $\langle \sigma v \rangle$ to 0.5 and 0.1 of its asymptotic value versus temperature T.

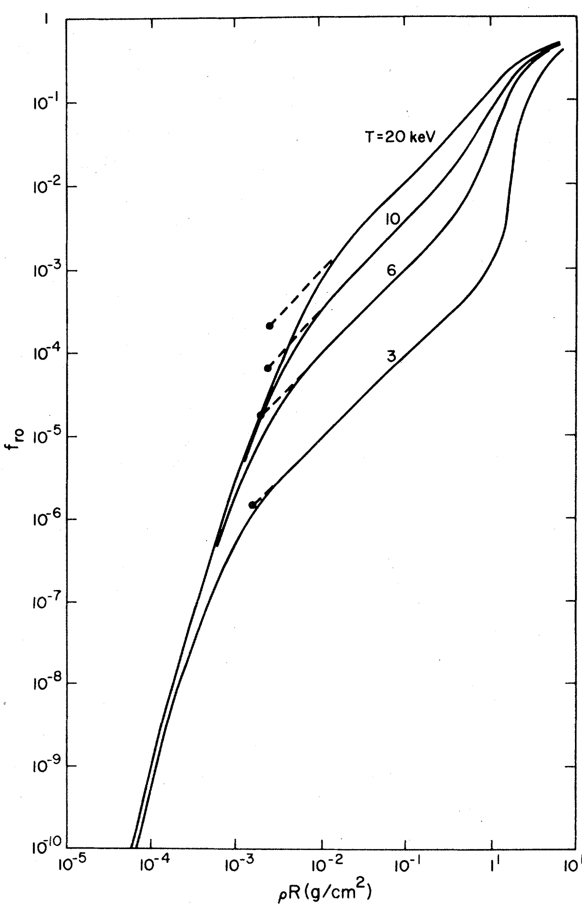

FIG. 3. Burnup fraction f_{ro} of a 1.0-μg sphere of DT versus areal mass density ρR for various temperatures T. This is a composite of the present results and a figure from Ref. 1.

In Ref. 1, it is shown that the burnup fraction f_{ro} may be expected to be one-half the characteristic expansion time $\tau_e = R/4C_s$, where $C_s = (2\gamma T/M_i)^{1/2}$ divided by the characteristic burn time $\tau_r = 1/n\langle\sigma v\rangle$.

I have therefore taken the computed data (τ_e/τ_r) and joined it onto the complete burn study result from Fig. 10(a) of Ref. 1. In some cases a small (1.0 to 1.1×) multiplicative adjustment was made in order to join the two. (The computer simulation data in Ref. 1 followed the dashed lines and terminated as shown.) In the composite figure, Fig. 3, we see that below $\rho R \approx 10^{-2}$ g cm^{-2}, we have a weakly temperature-dependent, fast ion-loss regime. From $\rho R \approx 10^{-2}$ to 1.0 g cm^{-2}, we have scaling as predicted in Eq. (1). Above $\rho R \approx 1.0$ g cm^{-2}, we have bootstrap heating to higher burnup. That the lost-ion regime is separated from the nonlinear burn regime by a factor of 100 in ρR is important in the justification of the analysis done here.

To date there have been no true thermonuclear neutrons observed in experiments,[11] which in some cases has been hard to reconcile with Eq. (1). The faster than linear small-ρR falloff of f_{ro} in these marginal experiments is adequate and probably a correct explanation. We should note that above $\rho R \approx 10^{-2}$ g cm^{-2} this difficulty goes away so our new results do not provide any obstacle in the path of useful laser-fusion applications. There is also a bright side to the present marginal ρR experimental difficulty: When true thermonuclear neutrons are observed from small microspheres, they will serve as proof of compression, an important milestone to laser fusion.

Dr. E. J. Linnebur provided helpful comments during this work as well as performing some of the initial computations. This work was done under the auspices of the U. S. Atomic Energy Commission.

[1] G. S. Fraley, E. J. Linnebur, R. J. Mason, and R. L. Morse, Phys. Fluids 17, 474 (1974).

[2] J. S. Clarke, H. N. Fisher, and R. J. Mason, Phys. Rev. Lett. 30, 89, 249 (1973).

[3] J. Nuckolls, L. Wood, A. Thiessen, and G. Zimmerman, Nature (London) 239, 139 (1972).

[4] J. Nuckolls, J. Emmett, and L. Wood, Phys. Today 26, No. 8, 46 (1973).

[5] The use of a number of contiguous shells of various low- and high-Z materials in laser and e-beam initiated fusion targets has been under study at the Los Alamos and Livermore Laboratories for many years.

[6] J. L. Tuck, Nucl. Fusion 1, 202 (1961).

[7] G. Gamow and C. L. Critchfield, *Theory of Atomic Nucleus and Nuclear Energy Sources* (Clarendon, Oxford, England, 1951), Chap. X.

[8] L. Spitzer, Jr., *Physics of Fully Ionized Gases* (Interscience, New York, 1956). To obtain, use $\lambda = \omega t_D$ and substitute (5-21), (5-13), and (5-4) in sequence.

[9] B. H. Duane, Battelle Northwest Laboratory Annual Report No. BNWL-1685, 1972 (unpublished), p. 75.

[10] D. B. Henderson, LASL Report No. LA-5603-MS, 1974 (unpublished).

[11] G. H. McCall, F. Young, A. W. Ehler, J. F. Kephart, and R. P. Godwin, Phys. Rev. Lett. 30, 1116 (1973).