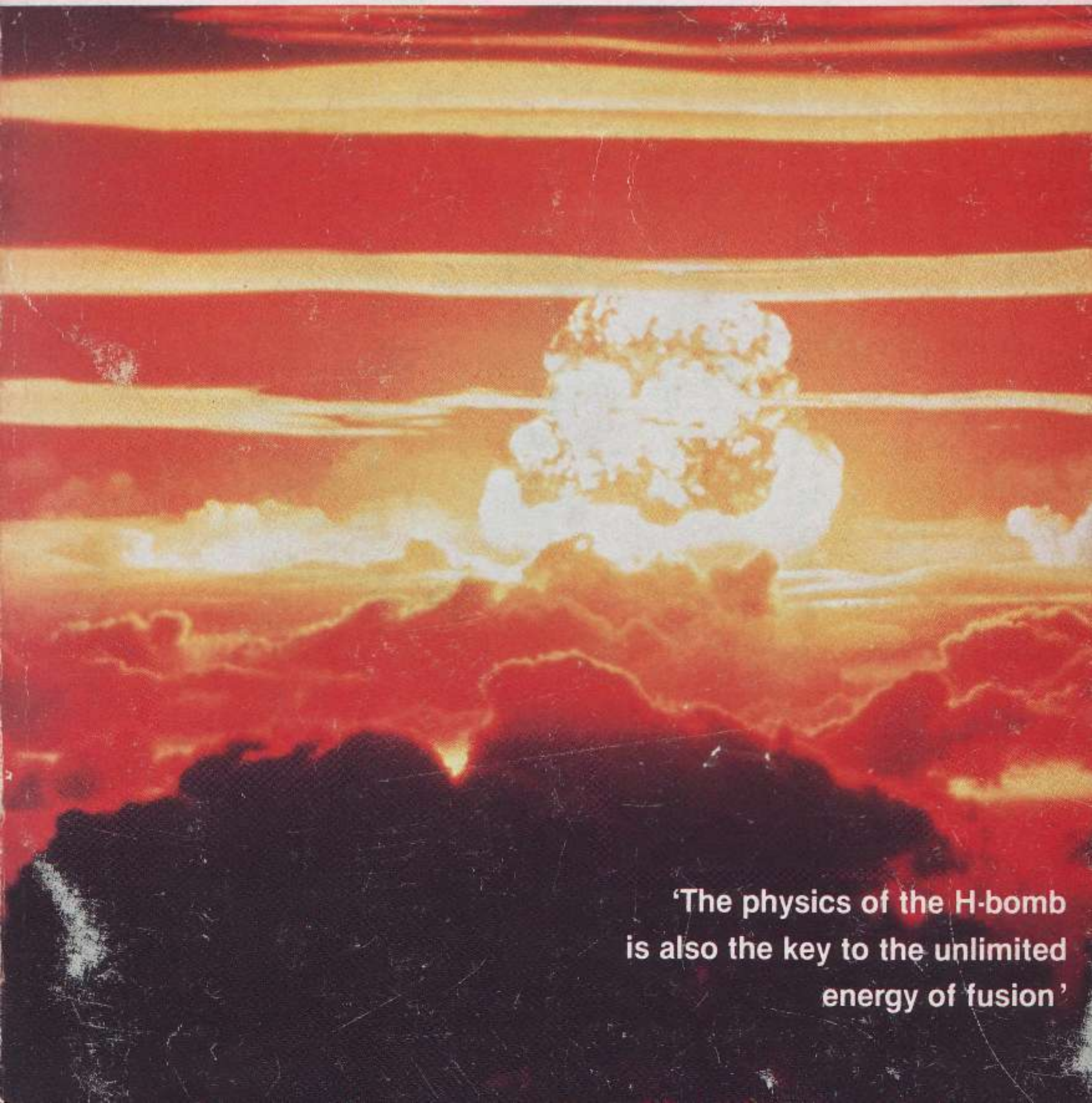
FRIEDWARDT WINTERBERG

The Physical Principles Of Thermonuclear Explosive Devices



'The physics of the H-bomb
is also the key to the unlimited
energy of fusion'

The Physical Principles Of Thermonuclear Explosive Devices

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The Physical Principles of Thermonuclear Explosive Devices
by Friedwardt Winterberg
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The Physical Principles Of Thermonuclear Explosive Devices

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Foreword

It is appropriate that this, the first volume in the Fusion Energy Foundation series of books on the frontiers of science and technology, should have as its subject the basic physics of thermonuclear explosions. The tremendous energy release from fusion reactions defines the most advanced aspects of strategic and economic policy for the last decades of the twentieth century.

It has been an unavoidable fact of life since the Manhattan Project that fusion science has been closely associated with research on thermonuclear weapons. The H-bomb first brought the process of fusion to public attention, two decades before the scientific advances that have brought fusion energy to the threshold of energy breakeven. The solution to the originally formidable problems of substantial energy release from weapons systems has been an integral part of the achievement of controlled fusion, particularly in inertial-confinement fusion.

Some form of inertial-confinement fusion will ultimately produce the most efficient and useful output of fusion energy, in all likelihood. Inertial fusion therefore defines one of the most important frontiers of civilian and military science about which citizens must be adequately informed.

This is not possible at present, however, because many of the basic ideas and results in inertial fusion are still classified. For many years we of the Fusion Energy Foundation have fought to change this situation.

It is intolerable that the theoretical underpinnings of the science that, in the judgment of most knowledgeable individuals, will determine the future of the world should be classified. It is intolerable that the leading edge of human thought has deliberately been made inaccessible, not only to the layman, who needs to be informed, but also to the working scientist, to whom this situation poses the unacceptable choice of surrendering an independent role as a civilian scientist or being cut off from access to the principal data that relate to this most important area of research.

Dr. Winterberg's book is not only a valuable contribution toward properly informing citizens on these crucial civilian and military scientific frontiers. It also reminds us of the hydrodynamic tradition involving such pioneering figures as Ludwig Prandtl, which has proven so fruitful in modern science and technology.

If this book contributes to informing the national citizenry as well as to restoring the vigor of that scientific tradition; if it contributes to the development of new theoretical and experimental approaches to solving the problems of the fundamental processes of fusion, then it will have achieved its primary purposes. The physics of the H-bomb is also the key to the unlimited energy of fusion.

Dr. Morris Levitt
Executive Director

Dr. Uwe Parpart
Director of Research

FUSION ENERGY FOUNDATION
New York City

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Preface

Thermonuclear weapons are the most destructive technology ever devised by man. Why, then, some may ask, would I write such a book? Would the publication of it not accelerate the race toward more countries acquiring these weapons, ultimately leading to a thermonuclear holocaust, with casualties in excess of 600 million people, by conservative estimates, and the destruction of a large part of our cities and industries in the civilized world? The answer to this question would be yes if there ever was such a secret to guard, without which no thermonuclear weapons could be built.

This assumption, however, is totally false. The possibility of thermonuclear reactions has been known since 1928, and the discovery of nuclear fission ten years later provided a match for their ignition. The ignition of thermonuclear reactions requires temperatures of $\sim 10^8$ K, which can be produced with an exploding fission bomb. All the remaining questions as to how fission and thermonuclear explosives would have to be arranged are technical details. The only reason that not every country with the desire to own thermonuclear weapons does not have them is not the lack of knowledge, which can be obtained from the published literature, but the extreme expense and the need for a highly developed industry to produce the required nuclear explosives, including the materials both for the fission trigger and the fusion bomb.

Fortunately, however, there is another aspect of thermonuclear energy that is known as controlled thermonuclear fusion. It requires replacing the fission trigger with some other means of igniting the fuel. Recent progress toward this goal gives us a real hope that we are close to a breakthrough. The prospect of controlled thermonuclear energy promises an abundant and clean source of energy that could last for millions of years. This underlines the extreme significance this energy source holds for the future. However, because of the close scientific connection between inertial-confinement fusion, which is one of the most promising approaches, and thermonuclear weapons, the government has put this kind of research work under the wraps of secrecy.

But since it is a delusion to believe that a secret exists, this self-imposed secrecy has the effect that research performed in the government laboratories cannot be checked by the general scientific community. If secrecy persists, therefore, errors are likely to be made that will cost the taxpayer millions of dollars. The resulting failure could retard progress toward controlled fusion by many years. Ironically, this eventually could lead to the very energy war fought with thermonuclear weapons that the government wants to avoid.

Some people may argue that perhaps we should declassify inertial-confinement fusion research but keep weapons research secret. This suggestion, however, is completely unfeasible, since certain aspects of inertial-confinement fusion research closely resemble the problems in weapons research, making a separation of the two problems impossible.

In science one is not permitted to make a prejudicial value judgment before the data are put on the table. A request like that would resemble an order given by the Church to scientists who study the scriptures to restrict their studies to those scriptures arbitrarily called the only holy or canonical ones by ancient Church Congresses and to ignore all other apocryphal gospels or scriptures. Science cannot be divided like this, and any real scholar would disregard such an order. A scientist cannot be ordered to ignore any questions of inertial-confinement fusion for peaceful energy production that may also have weapons applications. To comply with this indivisibility of science I have included in my book the topic of thermonuclear microexplosions and some exotic concepts making use of thermonuclear explosives that can even advance basic physics.

Wernher Von Braun once said that we have invented rockets, not to destroy our planet, but to explore the universe. Similarly, we may say that we have discovered thermonuclear energy, not to destroy our planet, but to advance mankind toward a peaceful, galactic culture.

—Dr. Friedwardt Winterberg

Introduction

The physical principles for the explosive release of thermonuclear energy are known to at least five countries: the United States, the Soviet Union, Britain, France, and China. Each of these countries has tested thermonuclear explosive devices with the goal of establishing a thermonuclear weapons capability. In spite of this apparently widespread knowledge, the physical principles by which these devices operate have been guarded by each of these countries as one of their greatest secrets. This secrecy has no parallel in the entire history of military science, and there is no rationale for this state of affairs. True, the existence of thermonuclear weapons is a threat to all mankind that is unmatched in our entire history, but it is not clear how secrecy can reduce this threat. Experience shows that every country with the will to develop thermonuclear weapon devices and possessing the necessary industrial infrastructure has independently succeeded in attaining this knowledge. The reason given by the U.S. government for this secrecy is that disclosure of this knowledge would accelerate the proliferation of nuclear weapons. But in view of the already existing proliferation, this does not appear to be a very credible argument. The facts suggest that proliferation is the result of a political decision, rather than anything else.

I am an advocate of total nuclear disarmament and the eventual establishment of a world government. At the present time such a goal seems to be quite Utopian. Perhaps mankind will have to learn the horrors of an all-out thermonuclear war before realizing the necessity of a world government in the nuclear age.

I have never been exposed to alleged secrets by any government—although I have been working in the field of thermonuclear research for almost 30 years and have laid the foundations for the ignition of thermonuclear microexplosions by intense electron and ion beams for controlled, peaceful energy production. I am therefore in a position to state that *there are no secrets surrounding thermonuclear explosive devices* and that *all the basic physics is accessible in the open, published scientific literature.* Therefore, in reading this work that exposes such "secrets" as ordinary physics, the scientists working in the secret nuclear-weapons laboratories may judge for themselves if they actually hold such great secrets.

My purpose is not to be sensational but, rather, to demystify the secret of the H-bomb. For it is not the *secret* of the H-bomb that protects us from thermonuclear annihilation but, rather, the correct political decisions by our leaders. To cover up their own political inability it is, of course, understandable that governments try to make their people believe it is secrets that protect them. I hope that the publication of this book will not only contribute in demystifying the whole business of secrets, but also make the public aware that a belief in security by secrets is dangerous, wishful thinking.

This book is written in the language of the physicist, but much of its information should be also comprehensible to the interested engineer. Even the layman should be able to understand the physical principles shown in the illustrations.

CHAPTER ONE

The Historical Origin Of Thermonuclear Explosive Devices

Thermonuclear reactions were first proposed in 1928 as the energy source of the sun and the stars.[1] For many years, man-made release of thermonuclear energy was considered in the realm of science fiction. This situation was dramatically changed, however, with the discovery of nuclear fission by Hahn and Strassmann in 1938. Since that time the thought of using an exploding fission bomb as a match to ignite a much larger thermonuclear explosion must have occurred to many physicists. It is certain that such ideas were entertained not only in Los Alamos. For example, in a 1946 publication[2] of the renowned Austrian theoretical physicist Hans Thirring, the concept of such a thermonuclear superbomb is discussed critically. In 1950 Ulrich Jetter,[3] a German physicist, proposed using ^{6}LiD as the explosive for a "dry" hydrogen bomb, several years before this same concept was adopted by the United States and the Soviet Union in their thermonuclear weapons devices.[4]

Since simply placing a thermonuclear explosive beside an exploding fission bomb gives a negligible yield, another crucial problem had to be solved: *finding a configuration for the geometric arrangement* of the fission and thermonuclear explosives that would give a large yield. In fact, configurations do exist in which just one small fission bomb can ignite an arbitrarily large amount of thermonuclear explosive. In published accounts of the history of the hydrogen bomb, frequent reference is made to the Teller-Ulam configuration.

Without being privy in any way to this configuration, I shall show that there are supersonic flow configurations, first studied by the famous German aerodynamicist Prandtl and his students Meyer and Busemann, that have all the properties required for a thermonuclear trigger configuration. The studies by these scientists go back to the year 1908. One configuration of particular interest was analyzed by Busemann in 1942.[5] In it a spherical convergent shock wave is produced by a supersonic flow into a convergent nozzle of a specially chosen wall curvature. In another geometrically similar configuration, soft X-rays from an exploding fission bomb are focused by a multilayered curved wall. Both these configurations can be used to ignite a thermonuclear explosion of arbitrarily large magnitude with one exploding fission bomb.

CHAPTER TWO

Thermonuclear Explosives

Thermonuclear reactions in stars proceed at a rather slow pace and are thus unsuitable for man-made thermonuclear energy release, whether it takes place explosively in thermonuclear bombs or more slowly in magnetic plasma-confinement devices. The most frequently discussed fusion reactions involve two hydrogen nuclei forming isotopes of helium plus free neutrons and energy.

The *ignition temperature* of a thermonuclear reaction is defined as the lowest temperature above which the energy production by thermonuclear reactions exceeds the energy loss by radiation.

Of all known thermonuclear reactions the DT (deuterium-tritium) reaction

$$D + T \longrightarrow {}^4He + n + 17.6 \text{ MeV} \tag{1}$$

between the heavy (D) and superheavy (T) hydrogen isotopes has the lowest ignition temperature, around 5×10^7 K. Next in line is the DD (deuterium-deuterium) thermonuclear reaction

$$D + D \begin{cases} \longrightarrow T + H + 4 \text{ MeV} \\ \longrightarrow {}^3He + n + 3.25 \text{ MeV}, \end{cases} \tag{2}$$

with both branches occurring with about the same probability. The ignition temperature of the DD thermonuclear reaction is about 10 times larger than for the DT reaction, that is, around 5×10^8 K. All other thermonuclear reactions have much higher ignition temperatures. It is for this reason that the DT and DD reactions were first studied. The somewhat inappropriate name "hydrogen bomb," now used for all thermonuclear explosive devices, has its roots here.

Since, as we shall show in the next chapter, the temperature in a fission explosion can reach the ignition temperature of the DT but not of the DD thermonuclear reaction, it was long believed that a thermonuclear weapons program would depend on a large tritium inventory. This, however, has two very serious drawbacks. First, the tritium would have to be generated in a nuclear reactor out of the lithium isotope ^{6}Li through neutron bombardment by the reaction

$$\mathrm{n} + {}^6\mathrm{Li} \longrightarrow \mathrm{T} + {}^4\mathrm{He} + 4.8\ \mathrm{MeV}. \qquad (3)$$

This is very costly. Second and worse, tritium has a beta decay half-life of 12.3 years transforming it into ^{3}He. To sustain a large tritium inventory, therefore, a large, continuously operating tritium-production facility to compensate for the continuously decaying tritium would be required.

To overcome these difficulties, Jetter[3] proposed in a remarkable paper to use ^{6}LiD directly as the thermonuclear explosive by letting the tritium generation occur through neutron reactions during the thermonuclear explosion itself. To accomplish this, Jetter suggested using the following closed-chain reaction:

$$\begin{array}{ccc} \mathrm{T} + \mathrm{D} & \longrightarrow & {}^4\mathrm{He} + \mathrm{n} + 17.6\ \mathrm{MeV} \\ \uparrow & & \downarrow \\ 4.8\ \mathrm{MeV} + \mathrm{T} + {}^4\mathrm{He} & \longleftarrow & {}^6\mathrm{Li} + \mathrm{n} \end{array} \qquad (4)$$

In this reaction one of the chain links is a neutron reaction, not a thermonuclear reaction that only takes place between charged particles. The large quantities of neutrons initially required to keep the reaction going at a large pace could therefore in principle be supplied by a fission explosive. However, since by diffusion some neutrons are always lost, the chain reaction (4) must be coupled with a neutron multiplier. As neutron multipliers, Jetter proposed using fissile material such as ^{238}U (an isotope of uranium) or ^{9}Be (an isotope of beryllium).

Since the chain reaction (4) is not a pure thermonuclear reaction, with one chain link going over a neutron reaction, what has been previously said about the ignition temperatures of thermonuclear reactions does not apply here. However, the presence of some higher-Z material increases the radiation losses, and the ignition temperature is therefore expected to be higher than for the DT thermonuclear reaction. Estimates suggest an increase in the ignition temperature up to 10^8 K. Such a temperature increase would make the ignition of this reaction by an exploding fission bomb dubious. Before coming to any conclusions here though, we must also determine the heating effect from the neutrons released in the fission explosion and reacting with the ^{6}Li. We shall see below that the temperature can be considerably increased over the temperature in the center of a fission explosion by *gasdynamic energy focusing.* For this reason a tenfold increase in the ignition temperature does not pose a problem.

Although tritium and deuterium exist in liquid or solid form only at very low temperatures, ^{6}LiD is solid at room temperature, which gives it a distinct advantage as an explosive for thermonuclear weapons devices. However, if used in conjunction with a ^{238}U neutron multiplier (which at the same time can serve as a tamp for the thermonuclear explosion), it produces extensive radioactive fallout from the release of fission products. A ^{6}LiD bomb is therefore always quite *dirty*. A pure DT or DD explosion, in contrast, is *clean*. DD explosives, in particular, may have interesting nonmilitary technical applications.

CHAPTER THREE

The Temperature and Energy Flux Occurring In a Fission Explosion

The energy released in a fission explosion goes in part into particle energy and in part into radiation. The particle energy density is given by

$$\varepsilon_p = (f/2)nkT,$$

where n is the particle-number density, T the temperature, f the number of degrees of freedom, and k the Boltzmann constant.

For the linear dimensions of a fission explosive the fireball is optically opaque, and the energy density of the radiation field is that of the blackbody radiation, given by $\varepsilon_r = aT^4$, with $a = 7.67 \times 10^{-15}$ erg/cm^3-K^4. The density of metallic uranium or plutonium is about 18 g/cm^3 with an atomic number density in the solid state equal to $n_s \simeq 5 \times 10^{22}$ cm^{-3}. Therefore, if $\varepsilon_r > \varepsilon_p$, the energy density of the blackbody radiation will be predominant over the particle-energy density. This will happen if

$$T > (7.7 \times 10^6)(fn/n_s)^{1/3}\ \mathrm{K}, \tag{5}$$

where we have assumed that the fissile material may be precompressed by an explosive above solid density with $n > n_s$. In a uranium (or plutonium) plasma of many million kelvins, the internal atomic excitation f can become quite large. Let us assume that $f = 100$, even though the true value of f may be actually smaller. This then would imply that for blackbody radiation to predominate, $T > 3.7 \times 10^7$ K.

The energy per unit volume set free in a fission chain reaction is given by

$$\varepsilon = n\varepsilon_f F, \tag{6}$$

where

$$\varepsilon_f = 180 \text{ MeV} = 2.9 \times 10^{-4} \text{ erg}$$

is the kinetic energy released in the fission reaction and F is the burnup fraction. Let us assume that $F = 0.1$ and furthermore that $n = n_s$. This would lead to

$$\varepsilon = 1.45 \times 10^{17} \text{ erg/cm}^3.$$

By equating

$$\varepsilon = \varepsilon_p + \varepsilon_r,$$

and solving for T we obtain $T = 6.6 \times 10^7$ K. Therefore, for a fission explosion with the chosen parameters the energy density of the blackbody radiation greatly exceeds the particle-energy density because of the T^4 dependence of ε_r. This conclusion is not changed if the fissile material is precompressed to approximately tenfold higher densities, which can be done with high explosives. A tenfold increase in density would change the temperature above which blackbody radiation becomes predominant by the factor $10^{1/3} \simeq 2.16$, from $T = 3.7 \times 10^7$ K to $T = 8 \times 10^7$ K, but at the same time it would also change the temperature of the fission plasma by the factor $10^{1/4} \simeq 1.8$, from $T = 6.6 \times 10^7$ K to $T \simeq 1.2 \times 10^8$ K.

From the foregoing considerations it follows that the maximum temperatures in fission explosions can reach the ignition temperature of the DT reaction, but not of the DD or any other pure thermonuclear reaction.

Because of the predominance of ε_r over ε_p the pressure of an exploding fission bomb is dominated by the blackbody radiation and is given by

$$p_r = aT^4/3 \simeq \varepsilon/3 \simeq 5 \times 10^{16}\ \mathrm{dyn/cm^2},$$

and the maximum radiative energy flux P is given by the Stefan-Boltzmann law

$$P = \sigma T^4, \tag{7}$$

where

$$\sigma = ac/4 = 5.75 \times 10^{-5}\ \mathrm{erg/cm^2\text{-}sec\text{-}K^4}.$$

For $T = 6.6 \times 10^7$ K, one finds that

$$P \simeq 10^{27}\ \mathrm{erg/cm^2\text{-}sec} = 10^{20}\ \mathrm{W/cm^2}.$$

Another number of considerable importance is the maximum neutron flux in a fission explosion. To estimate this flux we consider a critical mass of fissionable material of the order $\sim 10^4$ g. The fission assembly has here a radius of about 5 cm. With a burnup ratio of 10%, about $N \simeq 3 \times 10^{24}$ neutrons would be set free. On the other hand, the hydrodynamic disassembly time for the fission explosion is of the order $\tau \sim r/c_s$, where

$$c_s = \sqrt{(4/3)aT^4/\rho}$$

is the velocity of sound in a plasma dominated by blackbody radiation. With $\rho = 18.7\ \mathrm{g/cm^3}$, the density of uranium or plutonium, and $T = 6.6 \times 10^7$ K, the temperature of a fission explosion given above, we find $c_s \simeq 10^8$ cm/sec, and hence $\tau \sim 5 \times 10^{-8}$ sec. The average neutron flux at the surface of area $4\pi r^2$ is thus given by

$$\phi = N/4\pi r^2\tau \simeq 2 \times 10^{29}\ \mathrm{cm^{-2}\text{-}sec^{-1}}.$$

This neutron flux lasts for the time $\sim\tau$.

CHAPTER FOUR

The Fission Bomb As a Trigger

The use of a fission bomb as the trigger for a thermonuclear explosion raises the following serious problem: To ignite a thermonuclear explosive by a fission bomb, the thermonuclear material must be exposed in some way to the large energy flux of the fission explosion. This requirement makes it more difficult to use plutonium bombs as a means of triggering a thermonuclear explosion because plutonium from production reactors always contains a small fraction of the plutonium isotope ^{240}Pu. This plutonium isotope is generated by radiative, nonfissioning neutron capture in already-converted ^{239}Pu. The fraction of ^{240}Pu is therefore much larger for plutonium generated in fission power-plant reactors with much longer fuel-irradiation times.

The negative effect of ^{240}Pu is that it produces a large neutron background by spontaneous fission processes, so that the plutonium bomb must be assembled from a subcritical into a critical configuration very quickly, much faster than for a ^{235}U bomb. This can be done only by imploding the subcritical parts of the assembly with a velocity of about 10 km/sec, a velocity that is just marginally attainable with explosive shape-charge lenses. Such shape-charge lenses, however, require surrounding the plutonium bomb with large quantities of high explosives. The reaction products of these high explosives form a physical obstacle to transporting the large energy flux from the fission explosion to the thermonuclear material.

One way in which this problem can be eased is to remove the ^{240}Pu by isotope separation. Although possible in principle, this is very expensive, and it is not known to me personally whether this method is being used in the United States or the Soviet Union. It is therefore understandable that all five countries that have developed thermonuclear weapons employed ^{235}U triggers, at least at first, which do not have the problem of neutron background. In contrast to plutonium, ^{235}U can be detonated by simply shooting a cylinder of subcritical ^{235}U into a preformed cylindrical hole of another subcritical sphere of ^{235}U. The problem of the reaction products from the conventional explosive charge surrounding the exploding fission bomb does not arise here. The first neutrons to start the chain reaction in the fission bomb are supplied by an (αn) neutron source, for example, polonium-beryllium.

The use of polonium for this purpose is suggested by the fact that it is a pure α-emitter, which allows activating the source in the moment of greatest criticality, for example, by separating the polonium from the beryllium with a gold foil, which is a good α-absorber, and which could be mechanically disrupted by the incoming subcritical ^{235}U projectile. The use of radium, for example, instead of polonium would have the disadvantage that a radium-beryllium neutron source would also produce neutrons through the (γn) reaction. Unlike α-particles, the γ-quanta cannot easily be shielded to prevent a neutron background from (γn) reactions.

The necessity to trigger a thermonuclear detonation with ^{235}U bombs is the largest obstacle for smaller countries developing a thermonuclear weapons capability, because this condition requires building large, expensive isotope-separation facilities.

However, since the critical mass for plutonium is so much smaller than that for ^{235}U, a country possessing a plutonium-separation facility would be in the position to build rather small thermonuclear weapons devices. This may be of importance in the quest for the neutron bomb, which is a small thermonuclear weapon. The Soviet Union's strong opposition to this weapon could be explained by the lack of such a plutonium isotope-separation facility in their possession.

The potential use of ^{233}U as a trigger has the disadvantage that the breeding of this isotope from ^{232}Th (thorium) is accompanied by the generation of some ^{236}U, which is a hard γ-emitter with a rather long lifetime. To have such a substance in a weapons device is a serious problem because it would require a massive biological radiation shield surrounding the weapon.

CHAPTER FIVE

The Ignition Problem

There are two effects that can in principle be used to ignite a thermonuclear reaction by a detonating fission bomb. The first of these two effects is the high temperature of the fission explosion. The second, no less important, effect is the large neutron flux. The neutrons produce large quantities of heat, both by inelastic scattering and neutron-induced nuclear reactions of heat, and therefore enhance ignition. A more important effect of the neutrons is that during a nuclear explosion they can convert *in situ* ^{6}Li into tritium.

One necessary condition for the ignition of a thermonuclear explosive is that it be heated above the ignition temperature, which is 5×10^7 K for the DT reaction and roughly 10 times larger for the DD reaction. For the ^{6}LiD chain, this temperature is estimated to be around 10^8 K.

In addition, for all thermonuclear reactions there is a critical condition that has to be satisfied. This condition is known as the *Lawson criterion.* For its derivation it is important that all the thermonuclear reactions take place between light nuclei at a temperature that is above the temperature for complete ionization. The thermonuclear material at these temperatures will therefore emit a rather small amount of radiation and can in most cases be considered optically transparent. The energy to heat the thermonuclear material, therefore, goes primarily into kinetic particle energy with an energy density equal to

$$\varepsilon_k = (3/2)(Z + 1)nkT.$$

The Lawson criterion can be derived in the case of inertial confinement in the following simple way: Consider a thermonuclear assembly heated up to the ignition temperature. The energy produced by thermonuclear reactions in the time τ during which the thermonuclear material is inertially confined is given by an expression of the form

$$\varepsilon_t = n^2 f(T)\tau,$$

where $f(T)$ is some known function of the temperature T. If more thermonuclear energy is to be set free than the energy needed to heat the thermonuclear plasma to a temperature above the ignition temperature generally, the inequality

$$\varepsilon_t > \varepsilon_k \tag{8}$$

must be satisfied. This inequality leads directly to the Lawson criterion, which can be written as follows:

$$n\tau > g(T), \tag{9}$$

where $g(T)$ is another known function of the temperature T and is proportional to $T/f(T)$. In the case of the DT reaction, $g(T)$ has a minimum at $\sim 10^8$ K with $n\tau \gtrsim 10^{14}$ sec/cm^3. At the ignition temperature of the DT reaction, $\sim 5 \times 10^7$ K, the value of $n\tau$ is larger and closer to 10^{15} sec/cm^3. For the DD reaction the minimum of $g(T)$ occurs at a temperature roughly ten times larger with the minimum value $n\tau \gtrsim 10^{16}$ sec/cm^3. To obtain a thermonuclear detonation requires substantially larger values for $n\tau$ than the critical ones given by the Lawson condition.

Since $n \propto \rho$ and $\tau \propto r/v$, where ρ is the fuel density, r its radius, and $v \propto T^{1/2}$ its thermal expansion velocity, another form of the

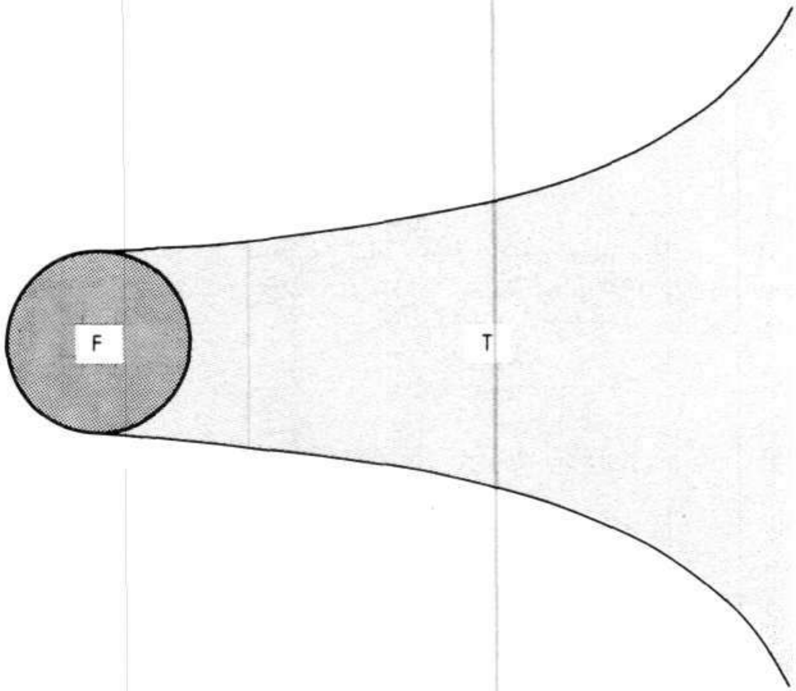

Figure 1. General schematic arrangement of fissile and thermonuclear explosives in which a fissile explosive F is to ignite the thermonuclear explosive T.

Lawson criterion is $\rho r \geq a$. For the DT reaction, $a \simeq 0.1$ g/cm^2.

It appears that the most simple approach to achieve ignition would be to place the thermonuclear explosive in direct contact with an exploding fission bomb (see Figure 1). In this hypothetical configuration, some of the thermonuclear material would serve as a fuse to ignite a much larger thermonuclear explosion. The major drawback of this approach is that it could only work (if at all) with DT as the thermonuclear explosive because the temperature of the fission explosion is not high enough to ignite other thermonuclear reactions. The temperature would also be short of what is required for ignition if the thermonuclear explosive is ^{6}LiD and the heating is done by neutrons through the reaction

$$n + {}^6\mathrm{Li} \longrightarrow {}^4\mathrm{He} + \mathrm{T} + 4.8\ \mathrm{MeV}.$$

The energy density that builds up during the inertial confinement time $\tau \simeq 5 \times 10^{-8}$ sec of the fission bomb over which the neutron flux lasts is given by

$$\varepsilon = n\sigma\phi\tau(E + E_n), \tag{10}$$

where $\sigma \simeq 3 \times 10^{-25}$ cm^2 is the neutron-reaction cross section, E = 4.8 MeV is the neutron-reaction energy, and $E_n \simeq 1$ MeV is the kinetic energy of the incoming fission neutron causing the reaction.

Using the above given values

$$\phi = 2 \times 10^{29}\ \text{cm}^{-2}\text{-sec}^{-1}, \qquad \tau = 5 \times 10^{-8}\ \text{sec},$$

and the atomic number density for ^{6}LiD,

$$n \simeq 5 \times 10^{22}\ \text{cm}^{-3},$$

one finds

$$\varepsilon = 1.4 \times 10^{15}\ \text{erg/cm}^3.$$

The kinetic particle energy density with four electrons to be counted for each ^{6}LiD molecule is $\varepsilon_k = (9/2)nkT$. Equating ε with ε_k, we get $T = 5 \times 10^7$ K. This result shows that the heating effect of the neutrons, although significant, is not sufficient to reach the ignition temperature for ^{6}LiD, which is around 10^8 K.

A more serious problem in the approach of just placing the thermonuclear explosive beside the fission explosion is that, because of the large radiation pressure coming from the exploding fission bomb, the thermonuclear explosive will be simply blown aside long before a large thermonuclear yield can be reached. As can be easily seen, this will be especially true for the thermonuclear fuse, which would be connected through a horn to a larger thermonuclear assembly. The diameter d of such a fuse must be of the same order of magnitude as the diameter of the fission bomb, that is, of the order $\sim$10 cm. The radiation pressure of the exploding fission bomb is given by

$$p_r = (aT^4/3) = \varepsilon/3 \simeq 5 \times 10^{16}\ \text{dyn/cm}^2.$$

This radiation pressure will exert a force on the fuse equal to $F \simeq p_r d^2 = 5 \times 10^{18}$ dyn. The mass of the thermonuclear fuse is of the order

$$M \simeq \rho d^3 \simeq 1.7 \times 10^2 \text{ g},$$

where $\rho \simeq 0.17$ g/cm^3 is the density of liquid DT. The acceleration of the fuse material by the radiation pressure is given by

$$b = F/M \simeq 3 \times 10^{16} \text{ cm/sec}^2.$$

Therefore, the fuse material will be displaced by the distance d in the time

$$\tau_d \simeq \sqrt{2d/b} \simeq 2.6 \times 10^{-8} \text{ sec}.$$

After this time, which is about half the hydrodynamic disassembly time (the inertial confinement time of the fission explosion), the fuse has been blown away. But in case of the DT reaction, the product $n\tau_d$ at $n = 5 \times 10^{22}$ cm^{-3} is equal to $n\tau_d \simeq 10^{15}$ sec/cm^3, which is only slightly larger than the Lawson product for the DT reaction at the ignition temperature of 5×10^7 K. These estimates suggest that this simple approach, even with DT as the thermonuclear explosive, would be marginally feasible at best.

The problem of igniting a thermonuclear explosion has thus been aptly compared to the problem of igniting a cigarette in a storm with a burning match. Long before the cigarette can catch fire, the match will be blown out.

Rather than placing the thermonuclear material outside the thermonuclear fission bomb, one may put it inside, as shown in Figure 2. The advantage of this configuration is that it increases the inertial confinement time of the thermonuclear material, which is compressed by an ingoing spherical implosion wave. This compression effect can be quite important because the thermonuclear reaction rate is proportional to the square of the density. A compression may be also indispensable for more exotic thermonuclear explosives.

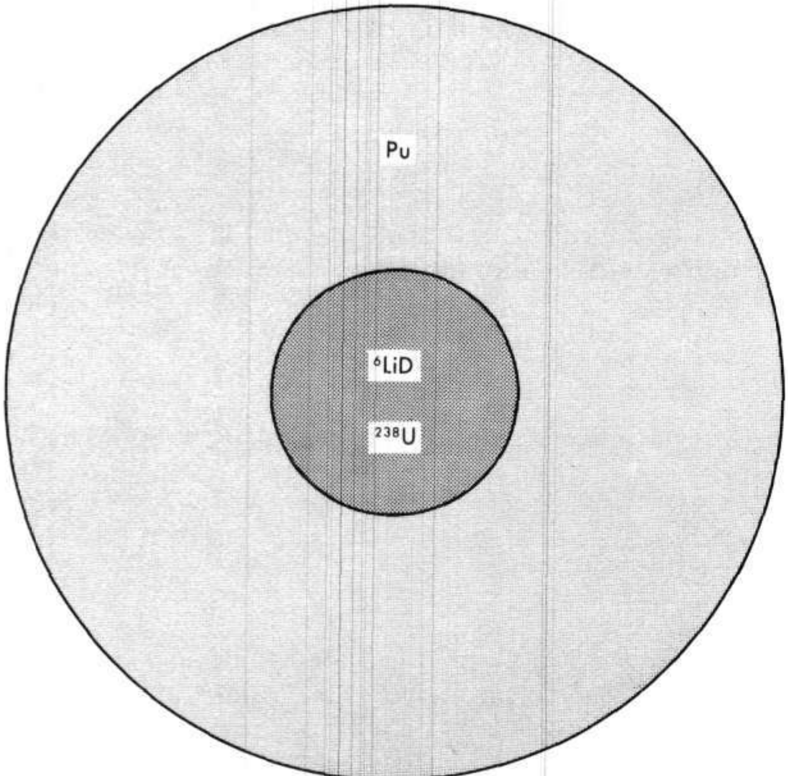

Figure 2. In the booster principle, the thermonuclear material, a mixture of ^{6}LiD and ^{238}U, is placed in the center of the fission plutonium bomb. This increases the total amount of neutrons released and hence the yield.

Another advantage is that the implosion process also increases the temperature. To ignite ^{6}LiD, for example, a temperature around 10^8 K is needed, which is above what is possible with fission explosions using uncompressed fissile materials.

There are, however, two important disadvantages of this configuration. First, it will not allow igniting an arbitrarily large amount of thermonuclear explosive; second, it substantially increases the critical mass of the fission assembly. This second disadvantage, at least, is compensated for by the new possibility that implanting a thermonuclear explosive in the fission bomb introduces a neutron multiplier into the system that couples the thermonuclear fusion process to the fission process.

If, for example, the thermonuclear material is DT, it by itself would act as a very efficient neutron multiplier through the ample generation of 14-MeV fusion neutrons. Through such a neutron multiplier the fission chain reaction can be greatly accelerated, but only after the neutron-producing thermonuclear reaction has been initially ignited. This effect therefore promises a substantial increase in the yield over what is attainable with a pure fission explosive. For a weapon, liquid DT would be impractical as a thermonuclear neutron multiplier. ^{6}LiD would also be unsuitable because it does not produce excess neutrons.

A better choice would be to use a mixture of ^{6}LiD and ^{6}LiT, but here the unstable tritium would again create a distinct disadvantage. More interesting is the possibility of mixing ^{6}LiD with uranium in some matrix, either homogeneously or, better, heterogeneously. The DT fusion neutrons have an energy of $\sim$14 MeV and can therefore fission ^{238}U nuclei much more effectively than the $\sim$1-MeV fission neutrons. Another possibility for a neutron multiplier is mixing ^{9}Be with the ^{6}LiD thermonuclear explosive.

These various concepts fall under what is called the *thermonuclear booster principle.* The advantage of such hybrid fission-fusion bombs over pure fission bombs is probably so great that we have reason to believe that practically all smaller nuclear weapons now incorporate this thermonuclear booster principle. One additional advantage of the booster principle is that it works even with "ordinary" plutonium, that is, plutonium from which ^{240}Pu has not been removed by isotope separation. The reason for this is simply that the high explosive imploding the plutonium bomb does not here interfere in the interaction between the fission and fusion processes.

In summary, we can say that a naive study of the ignition problem already leads to the very useful booster principle. Solving the more ambitious problem of igniting a thermonuclear explosion of arbitrarily large size with just one fission bomb requires more sophisticated concepts, some of which are analyzed in the following chapters.

CHAPTER SIX

The Polyhedron Configuration

One of the problems in igniting thermonuclear explosives other than DT is that the temperature of an exploding fission bomb falls short of what is required to ignite the DD thermonuclear reaction, and even the ^{6}LiD reaction. To overcome this limitation in temperature, we may put several fission bombs on the corners of a polyhedron. If the fission bombs in this configuration are detonated simultaneously, they will generate a convergent shock wave, which would become spherically symmetric in the limiting case of an infinite number of fission bombs. In practice, the number of fission bombs is limited. The smallest possible number is four, which would have to be placed at the four corners of a tetrahedron. A minimum of six bombs may actually be required or the departure from spherical symmetry may become too large.

For a spherical convergent shock wave, which can be described by Guderley's self-similar solution,[6] the temperature is predicted to rise in proportion as $r^{-0.8}$, where r is the distance of the wave front from the center of convergence. It is likely that with six fission bombs a threefold increase in temperature can be achieved. This would make the ignition of the ^{6}LiD reaction possible.

We shall now show that spherical implosion symmetry can be improved by a simple trick permitting reduction of the number of bombs required to four. In the next chapter we shall show that by extrapolation of the same "trick" the number of required bombs can be reduced to just one.

Suppose that the detonation waves from four fission explosions behave like small-amplitude sound waves. Then the configuration shown in Figure 3, displaying a cross section in one plane of the tetrahedral assembly, would give a much better approximation to a spherical implosion wave. In it the four fission bombs are placed each at one focus of four ellipsoidal mirrors, formed by a heavy substance (e.g., lead), reflecting the waves. The second foci of these

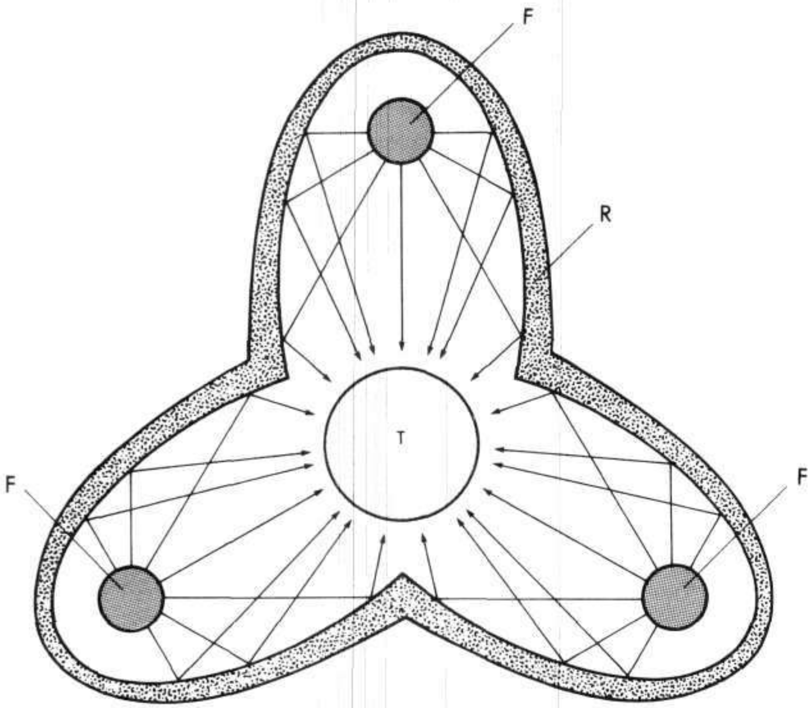

Figure 3. In the polyhedron-bomb ignition principle, several fission bombs are arranged around the thermonuclear explosive so that the shock waves from their simultaneous explosions will be reflected by the wave reflector R onto the thermonuclear explosive T positioned at the center of the tetrahedron. Shown here is a cross section in which only three of the four fission bombs appear.

four ellipsoidal mirrors coincide at one point positioned at the center of the tetrahedron. In this case the superposition of all the reflected detonation waves would result in an almost perfect spherical implosion wave. In reality, however, the behavior of large-amplitude detonation waves differs from that of small-amplitude sound waves, so that the shape of the large-amplitude focusing mirrors deviates from that of a mathematical ellipsoid. The mathematical shape of these large-amplitude shock wave focusing mirrors can be derived with the help of a 1908 theory by Prandtl and Meyer. We therefore give the correct wall shape the name Prandtl-Meyer ellipsoid, even though it is different from a mathematical ellipsoid.

The shape of the Prandtl-Meyer ellipsoid is closely related to the particular wall shape of a convergent nozzle to obtain a convergent conical flow first studied by Busemann.[5] The expected good spherical implosion symmetry thus obtained should make it possible to increase the temperature in the center of convergence substantially over the temperature of an exploding fission bomb and reach the ignition temperature of the DD thermonuclear reaction. We shall show how this configuration can be used to ignite an arbitrarily large pure-deuterium explosion.

CHAPTER SEVEN

Ignition by Implosion With Only One Fission Bomb

The concept of shock wave focusing we have outlined suggests using just one fission bomb for the ignition process, as shown in Figure 4. This fission bomb is put into one focus and the ignition point for the thermonuclear explosive into the other focus of a Prandtl-Meyer ellipsoid. The exact mathematical theory for treating this problem has been given by Busemann.[5] Here we shall instead use a simplified approximate treatment, which is physically more meaningful. In Figure 4, lines 1 and 2 represent the rays of the outgoing and ingoing shock waves, respectively, within an oval cavity representing a Prandtl-Meyer ellipsoid. By the intersection of the incoming wave at one particular point P with the wall of the Prandtl-Meyer ellipsoid a simple, or Riemann, wave is emitted from P under the Mach angle

$$\mu = \arcsin (1/M), \tag{11}$$

where M is the Mach number of the hypersonic flow associated with the divergent detonation wave at this particular point P. This raises the question of whether there exists a unique wall shape for which all Mach lines or characteristics of the reflected spherical waves meet in exactly one point.

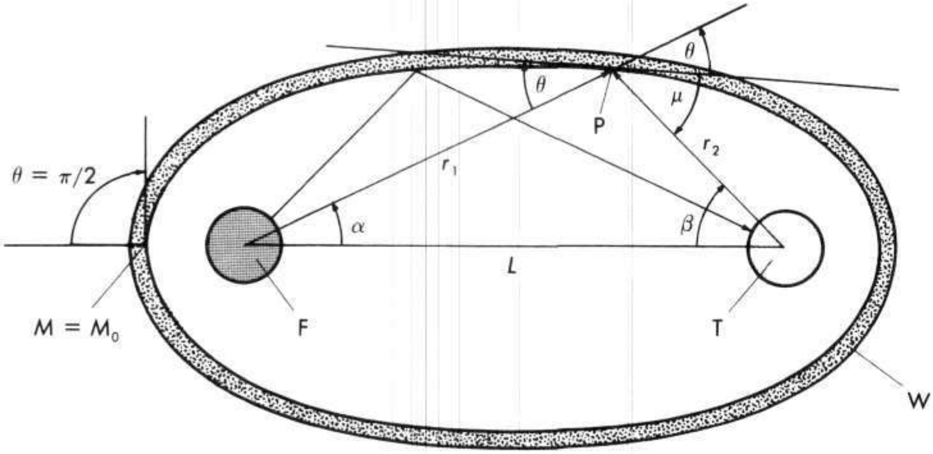

Figure 4. In the Prandtl-Meyer ellipsoid shock wave focusing principle, only one fission bomb F is used to induce a spherical implosion onto the thermonuclear explosive T by shock wave reflection from the wall inside the ellipsoidlike cavity. P is the point where the incoming wave intersects with the wall; M is the Mach number of the hypersonic flow associated with the diverging wave at P; θ is the angle between the wall slope and the incoming ray; r_1 and r_2 are the rays of the shock wave.

In general, if the Mach lines of the reflected waves intersect at different locations, the envelope of these intersecting Mach lines forms a stationary shock front. In the special case where all the Mach lines of the reflected waves meet in just one point, the result will be a spherical convergent shock wave in the vicinity of this point. A geometric configuration with this property indeed exists. It has two distinct points, the first the point from which the outgoing detonation wave originates and the second the point into which all the reflected waves converge. These two points are thus the two foci of the Prandtl-Meyer ellipsoid. In the limit of small-amplitude waves, corresponding to the acoustic approximation, the Prandtl-Meyer ellipsoid becomes a mathematical ellipsoid.

To determine the wall shape of the Prandtl-Meyer ellipsoid we need, in addition to the relationship between μ and M, a number of other relations. One of these relates the angle θ that the incoming ray makes with the wall slope and the Mach number M. This relationship is given by the integrated Prandtl-Meyer expression for a supersonic flow along a curved wall:

$$\theta - \theta_0 = \nu(M) - \nu(M_0), \tag{12}$$

where

$$\nu(M) \equiv \left(\sqrt{\frac{\gamma+1}{\gamma-1}} \arctan \sqrt{\frac{\gamma-1}{\gamma+1}(M^2-1)}\right) - \arctan \sqrt{M^2-1}. \tag{13}$$

In Eq. (13), γ is the specific heat ratio. For a plasma dominated by blackbody radiation, $\gamma = 4/3$. Furthermore, if M_0 is the Mach number at the vertex point of the ellipsoid next to the focus where the fission explosive is placed, it follows that $\theta_0 = \pi/2$. One may furthermore make the assumption that $M_0 \simeq 1$, since in reaching the vertex point the fireball has not yet appreciably expanded to become supersonic. Under this assumption $\nu(M_0) = 0$ and hence

$$\theta = \sqrt{\tfrac{5}{3}} \arctan \sqrt{\tfrac{3}{5}(M^2-1)} - \arctan \sqrt{M^2-1} + \pi/2. \tag{14}$$

The shape of the Prandtl-Meyer ellipsoid is determined by the condition that the outgoing shock waves from the first focus must meet in the second focus. This condition can best be expressed in bipolar coordinates r_1, r_2 together with the angles α and β as shown in Figure 4. Obviously, then

$$\alpha + \beta = F(M), \tag{15}$$

with

$$F(M) \equiv \theta(M) + \mu(M). \tag{16}$$

The change in density along a variable Mach-number flow is given by

$$\rho/\rho_0 = [(1 + \tfrac{1}{2}(\gamma - 1)M_0^2)/(1 + \tfrac{1}{2}(\gamma - 1)M^2)]^{1/(\gamma-1)}. \tag{17}$$

Inserting into this $M_0 = 1$ and $\gamma = 4/3$, we find

$$\rho/\rho_0 = [7/(6 + M^2)]^3. \tag{18}$$

Finally, we have to relate ρ to r_1 and, respectively, r_2. If the flow is completely isentropic, ρ must be a symmetric function of r_1 and r_2. We note that for a large-amplitude spherical expansion wave, $\rho \propto 1/r^2$. It is therefore reasonable to assume that we may put

$$\frac{\rho}{\rho_0} = r_0^2\left(\frac{1}{r_1^2} + \frac{1}{r_2^2}\right). \tag{19}$$

In Eq. (19), the radius r_0 is equal to the distance of the focus from the adjacent vertex point, where it was assumed that $M = 1$. Eliminating ρ/ρ_0 from Eqs. (18) and (19) gives

$$\frac{7}{6 + M^2} = r_0^2\left(\frac{1}{r_1^2} + \frac{1}{r_2^2}\right). \tag{20}$$

Introducing the distance between the two foci L, we obtain the two additional geometric relations

$$r_1 \cos\alpha + r_2 \cos\beta = L \tag{21}$$

and

$$r_1 \sin\alpha = r_2 \sin\beta. \tag{22}$$

The set of four equations (15), (20), (21), and (22) with the five unknowns r_1, r_2, α_1, α_2, and M then reduces to one relation between two unknowns, for which we may choose r_1 and r_2, giving the wall curve of the Prandtl-Meyer ellipsoid in bipolar coordinates. In practice, the shape of the cavity will be different, since part of the cavity wall will be ablated. This ablation effect will require an asymmetric cavity and the correct form of this cavity must be determined by much more complex computer calculations and, probably, many nuclear-explosive tests.

One problem arises immediately. If we use this configuration to ignite a thermonuclear explosion, the volume the thermonuclear explosive occupies cannot be much larger than the volume of the fission bomb. The reason for this is that the volume around the second focus, where the high temperature is reached, cannot be much larger than the volume occupied by the fission explosive. This problem is of no consequence, however, if the configuration is only used to ignite a *fuse* placed in the second focus, from which a thermonuclear detonation wave can proceed into an arbitrarily large amount of thermonuclear explosive.

The pressure in the second focus of the Prandtl-Meyer ellipsoid where the thermonuclear fuse is placed is the same as the pressure at the surface of the exploding fission bomb, provided the wave focusing is completely isentropic. At the temperature of $T \simeq 6.6 \times 10^7$ K this pressure is equal to the radiation pressure

$$p_r = 5 \times 10^{16} \text{ dyn/cm}^2.$$

The temperature creating this radiation pressure is approximately the same as the ignition temperature of the DT reaction. Therefore, in placing DT material in the second focal point of the Prandtl-Meyer ellipsoid, the particle-number density of the compressed and heated plasma is computed by putting $2nkT = p_r$. For the given example this gives

$$n \simeq 2.7 \times 10^{24} \text{ cm}^{-3} \simeq 50\, n_s,$$

indicating a large compression of the thermonuclear material above solid-state density.

A similarly elegant focusing method for the neutrons is not possible. In interacting with the wall material, the neutrons emitted from the fission explosion are uniformly scattered. If the wall material is a good neutron reflector, however (a property of almost any heavy material for fast neutrons), the neutron flux in the ellipsoidal cavity can be higher than would be the case without this effect. But a large neutron flux is necessary if ^{6}LiD is used as the thermonuclear explosive.

To ignite a ^{6}LiD fuse one must in addition surround the ^{6}LiD material by a shell of fissile material. This fissile material, in conjunction with the large neutron flux from the fission bomb explosion, will locally increase the neutron flux at the focus where the thermonuclear material is placed, which is necessary to convert the ^{6}Li into tritium.

Focusing of the soft X-rays from the blackbody radiation of the exploding fission bomb is not possible with this configuration. However, since these rays arrive first at the position of the thermonuclear fuse, the fuse must be shielded against them to prevent substantial preheating. In the following chapter we shall present a configuration making use of a conical shield, which can also shield a thermonuclear assembly ignited at the base of the cone. In case the soft X-ray precursor should pose a serious problem, one could even use the soft X-rays for ignition with the concepts of the X-ray resonance mirror, explained in Chapter 18, and ignition by ablative implosion, explained in Chapter 15.

An X-ray resonance mirror uses many foils of increasing atomic number and should be able to reflect X-rays up to 10 keV. If the ignition is done by soft X-rays, the shape of the focusing mirror is a mathematical ellipsoid consisting of many layers of atomic material, from light up to heavy elements, to reflect the largest possible energy spectrum of soft X-rays. In practice, the ellipsoidal X-ray resonance mirror may be combined with the Prandtl-Meyer ellipsoid first to ignite the fuel with soft X-rays and then to increase inertial confinement by the subsequent shock wave. Since a shock-wave-focusing Prandtl-Meyer ellipsoid consists of a thick tamp and the X-ray resonance mirror consists of many foils, the compromise configuration may be to put a soft X-ray resonance ellipsoidal mirror inside the Prandtl-Meyer shock-wave-focusing mirror.

CHAPTER EIGHT

Other Ignition Configurations

The configuration analyzed in the preceding chapter is a mirror-type shock wave focusing system, in effect. This suggests that the same result might be accomplished by a shock wave lens. It is easy to show that such a shock wave lens is, in fact, possible.

From the theory of plane shock waves it follows that for a given temperature T behind the shock front the propagation velocity of the shock wave is proportional to $A^{-1/2}$, where A is the atomic weight of the material through which the shock propagates. From this there follows a relative refractive index n for the shock wave propagating from a first medium of atomic weight A_1 to a second medium with atomic weight A_2

$$n = (A_2/A_1)^{1/2}. \tag{23}$$

If, for example, $A_1 = 1$—assuming that the first medium is filled with hydrogen gas—then a refractive index of $n = 2$ could be achieved by choosing a second medium with $A_2 = 4$ (e.g., helium gas). The still undetermined parameter of the density ratio ρ_2/ρ_1 for both mediums must be chosen to avoid a pressure jump that requires $p_2 = p_1$. This is necessary because if $p_1 \gg p_2$ the second medium would be blown away, or if $p_1 \ll p_2$ most of the incident shock wave energy would be reflected. Since $p = (1/2)\rho v^2$, the condition $p_2 = p_1$ then simply implies that

$$\rho_2/\rho_1 = A_2/A_1. \tag{24}$$

A configuration for shock wave focusing by such a lens is shown in Figure 5.

Next let us consider another useful focusing method that differs from the previous ones in that it permits the focusing of a cylindrical convergent wave. Such a configuration is especially useful for the ignition of a growing thermonuclear wave starting from the end of a fuel cylinder, as shown in Figure 6.

The fission explosive is placed at the vertex point of a flat, solid cone. The configuration is surrounded by a conical tamp as shown, with a gap between the outer cone surface and the inner tamp surface. The width of this gap space is chosen to be equal to the

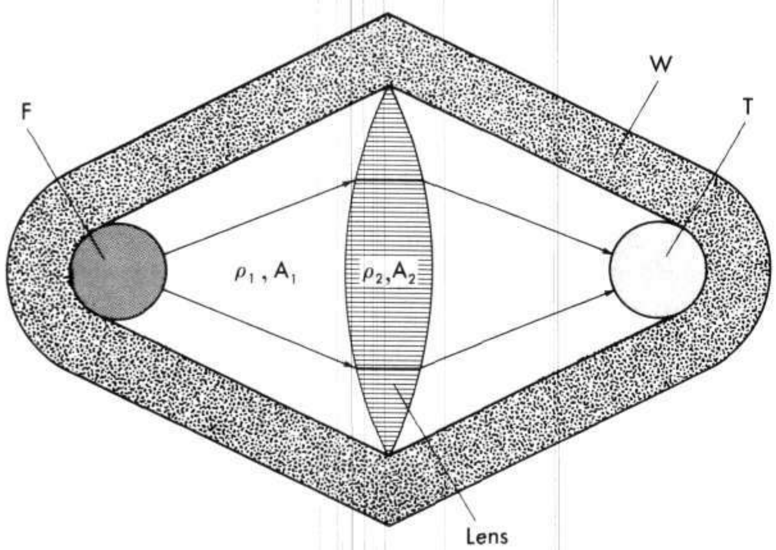

Figure 5. Shock wave focusing by a shock wave lens: the divergent shock wave from the fission explosive F is made convergent again by passing it through a medium of larger density acting as a lens. T is the thermonuclear explosive; W is the wall; ρ denotes density and A atomic weight.

diameter of the fission explosive. After the fission explosive has been detonated, a supersonic conical blast wave will then propagate from the vertex point of the cone to its base. At the base the supersonic blast wave can be rotated according to the theory of Prandtl-Meyer[8] by the angle

$$\beta \leq \beta_{\max}, \tag{25}$$

where

$$\beta_{\max} = \frac{\pi}{2}\left(\sqrt{\frac{\gamma + 1}{\gamma - 1}} - 1\right). \tag{26}$$

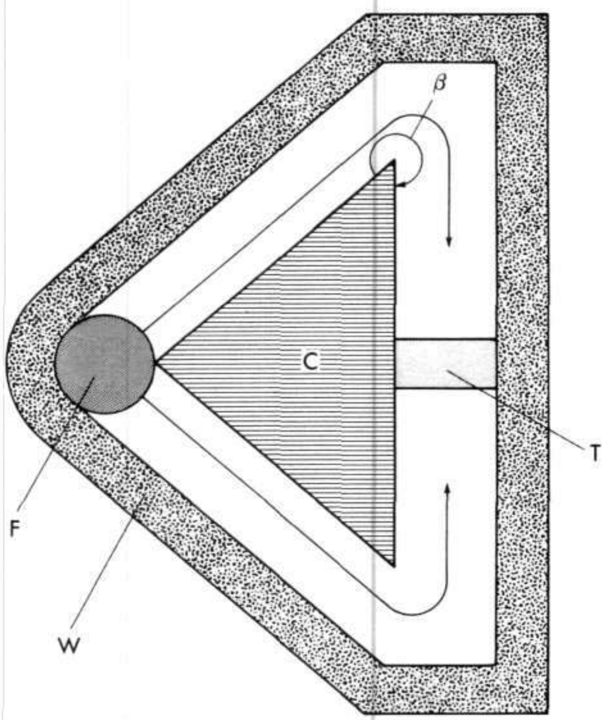

Figure 6. Cylindrical implosion by supersonic Prandtl-Meyer flow around the corner of a cone: the supersonic flow starting from the fission explosive F is here guided along the wall W and the outer surface of the cone C to generate a cylindrical implosion onto the thermonuclear explosive T.

Since the specific heat of the blast wave produced by the fission bomb is primarily from blackbody radiation, we may set the specific heat ratio $\gamma = 4/3$ and obtain

$$\beta_{max} \simeq 148^\circ.$$

This large number shows that one can make the cone rather flat.

After the conical blast wave has been rotated by the angle β, it will propagate as a cylindrical convergent wave along the base surface of the cone toward its axis. The thermonuclear fuel, having here the shape of a cylinder, would have to be positioned as shown in Figure 6.

The significance of this last configuration is that it shields the thermonuclear fuse optimally against the soft X-ray precursor of the exploding fission bomb. A disadvantage here is the cylindrical implosion geometry that results, rather than the optimal spherical configuration.

CHAPTER NINE

Multishell Velocity Amplification

As was previously pointed out, one of the problems in igniting a thermonuclear explosion in pure DD is the high ignition temperature, which is approximately 5×10^8 K, that is, ten times larger than the ignition temperature for the DT reaction. Such a high temperature should be attainable with the tetrahedral, four-fission-bomb ignition concept in conjunction with a Guderley convergent shock wave. The fact that the temperature in the Guderley case rises as $r^{-0.8}$ excludes the use of one fission bomb as the trigger.

There is, however, another configuration whereby a much larger temperature rise with decreasing radius is possible. This configuration makes use of the fact that in a series of elastic collisions of bodies decreasing in mass, the velocity toward the smaller mass is increased. If the mass ratio

$$m_{n+1}/m_n = \alpha$$

between elastically colliding bodies is constant, and the velocity of the first body with mass m_0 is v_0, then the nth body assumes the velocity after collision of

$$v_n = [2/(1 + \alpha)]^n v_0. \qquad (27)$$

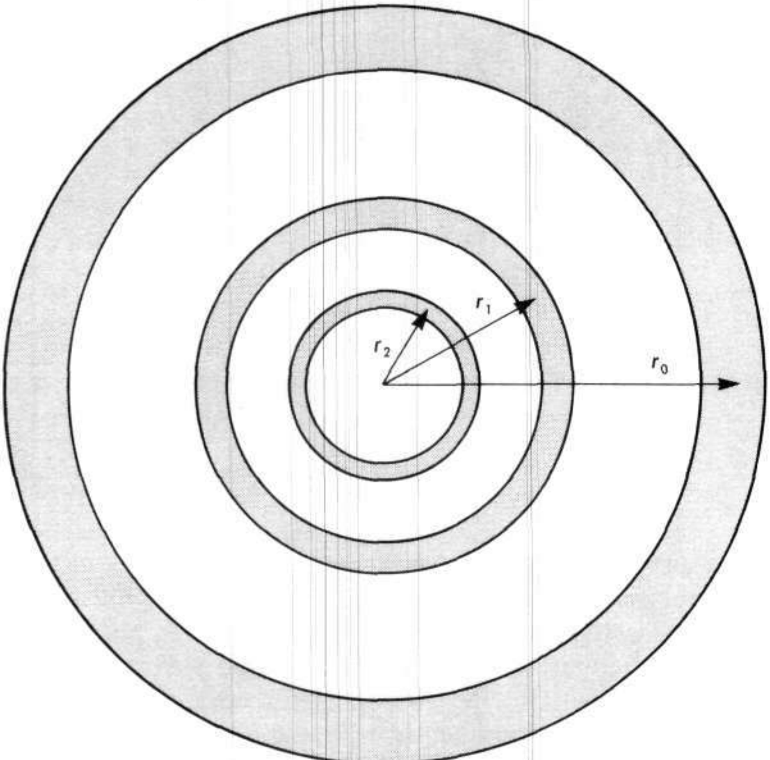

Figure 7. The implosion velocity can be raised through the subsequent collision of several concentric shells of decreasing mass.

We now make the special assumption that the subsequently colliding bodies are concentric spherical shells of radius r_n, having the mass

$$m_n = 4\pi\rho r_n^{\ 2} \cdot \Delta r_n, \tag{28}$$

where ρ is the density of the shell material and Δr_n is the shell thickness (see Figure 7). We further assume that the assembly of concentric shells is geometrically self-similar:

$$\Delta r_n = \varepsilon r_n, \qquad \varepsilon = \text{const} < 1. \tag{29}$$

We thus have

$$m_n = 4\pi\rho\varepsilon r_n^{\ 3} \tag{30}$$

and

$$\frac{m_n + 1}{m_n} = \left(\frac{r_n + 1}{r_n}\right)^3 = \alpha, \tag{31}$$

or

$$r_{n+1} = \alpha^{1/3} r_n. \tag{32}$$

We then put

$$r_n = r_0 f(n), \tag{33}$$

with $f(0) = 1$. Inserting this into Eq. (32) results in the functional equation for $f(n)$:

$$f(n + 1) = \alpha^{1/3} f(n). \tag{34}$$

With $f(0) = 1$, this functional equation has the solution

$$f(n) = \alpha^{n/3}. \tag{35}$$

We thus have

$$r_n = r_0 \alpha^{n/3}. \tag{36}$$

Eliminating from Eqs. (36) and (27) the parameter n we can express v_n in terms of r_n, that is, v as a function of r. Putting

$$a \equiv -\frac{\log\,[2/(1 + \alpha)]^3}{\log \alpha}, \tag{37}$$

we find

$$v_n/v_0 = (r_0/r_n)^a, \tag{38}$$

or, simply,

$$v \propto r^{-a}. \tag{39}$$

The temperature generated by a shell with a velocity v rises in proportion to v^2, hence

$$T \propto r^{-2a}. \tag{40}$$

We take as an example $\alpha = 1/2$ and find $a = 1.25$, hence $2a = 2.5$. It thus follows that a convergent multishell "wave" gives a temperature rise as a function of r with an exponent roughly three times larger than in the Guderley case. Therefore, by starting with a spherical wave at $r_0 = 5$ cm, a threefold velocity amplification is already reached at $r = 2$ cm. The number of shells needed is five since for $\alpha = 1/2$ and $v_n/v_0 = 3$ one computes from Eq. (27) that $n = 4$. In particular, $r_1 \simeq 4$ cm, $r_2 \simeq 3$ cm, $r_3 \simeq 2.5$ cm, and $r_4 \simeq 2$ cm. The threefold velocity amplification is accompanied by a $3^2 \sim 10$-fold rise in temperature. Actually, the temperature rise will be less dramatic because the collisions between the shells will be inelastic. This necessitates making r_0 larger and increasing the number of shells. Alternatively, one may make the inner shells out of fissile material that would become critical upon compression. This would compensate for any losses and could result in an even more rapid temperature rise.

CHAPTER TEN

Thermonuclear Detonation Waves

Once a thermonuclear reaction is ignited, a thermonuclear deflagration can spread to adjacent thermonuclear fuel provided a number of conditions are satisfied. It is interesting to note that these conditions are quite different from those for a detonation propagating through a chemical explosive.[7] The reason for this is that in a thermonuclear detonation wave the range of the fusion products is large in comparison to the shock discontinuity of the wave.

To derive the conditions for a thermonuclear detonation, we shall assume a rotational symmetric configuration as shown in Figure 8. Initially the trigger energy E_0 is deposited into the thermonuclear fuel occupying the volume V_0. If the fuel heated thus is a fully ionized hydrogen plasma with the specific heat ratio $\gamma = 5/3$, its equation of state is given by

$$p = 2nkT. \tag{41}$$

From the theory of plane shock waves[7] it follows that the rapid heating of the radially confined volume V_0 to the temperature T results in a shock wave propagating to the right with the velocity

$$v_0 = \sqrt{32kT/3M}, \tag{42}$$

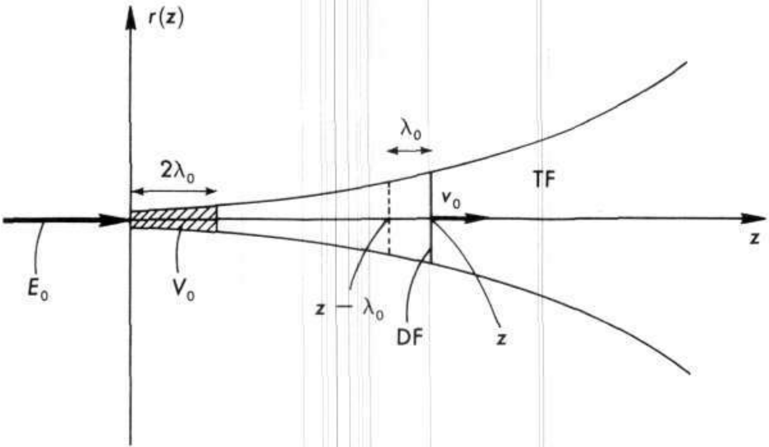

Figure 8. The physics of a growing thermonuclear detonation wave in a rotational symmetric configuration. The trigger energy E_0 is deposited into the thermonuclear fuel occupying the volume V_0.

where M is the average mass of the plasma ions. Furthermore, if the atomic number density in front of the shock is n_0 and behind it n_1, then

$$n_1 = 4n_0. \tag{43}$$

The fluid velocity v behind the shock front is

$$v = \tfrac{3}{4}v_0 \tag{44}$$

and it follows that

$$(1/2)Mv^2 = 3kT. \tag{45}$$

This means that half of the energy E_0 deposited into V_0 goes into heat and half of it into kinetic fluid energy. From this the value of T is obtained by the relation

$$E_0/V_0 = 6n_0kT. \tag{46}$$

For the shock wave to become a detonation wave, the energy supplied by thermonuclear reactions must keep its temperature at the constant value T. If the cross section of the fuel assembly as a function of z is given by $S(z)$, the required energy supply to drive the detonation wave is hence

$$\varepsilon(z) = 6n_0 kTv_0 S(z), \tag{47}$$

or, from Eq. (42),

$$\varepsilon(z) = 24\sqrt{2/3}\, n_0(kT)^{3/2}M^{-1/2}S(z). \tag{48}$$

The thermonuclear energy production per unit volume and time is given by

$$\dot{E} = (1/4)n_1^2\langle\sigma v\rangle\varepsilon_0 = 4n_0^2\langle\sigma v\rangle\varepsilon_0, \tag{49}$$

where ε_0 is the energy set free in one thermonuclear reaction and $\langle\sigma v\rangle$ is the product of the particle velocity and the nuclear cross section, averaged over a Maxwellian velocity distribution at the temperature T. [In the special case of identical particles reacting with each other, as in the DD reaction, the reaction rate is twice as large and the factor 144 in Eqs. (60) and (61) below must be changed to 72.]

For the following it is important to note that the detonation wave is driven primarily by charged fusion products, which both in the case of the DT thermonuclear reaction and of the ^{6}LiD chain are α-particles. In the case of the DD thermonuclear reaction, there are three kinds of charged fusion products: protons, tritons, and ^{3}He particles.

The flux ϕ of the charged fusion products in a plasma of temperature T and density n is approximately attenuated by the law

$$\phi = \phi_0 e^{-z/\lambda_0}. \tag{50}$$

Here

$$\lambda_0 = a(kT)^{3/2}/n \tag{51}$$

is the average range of the charged fusion products, with the constant a depending on the reaction energy of the α-particles generated by the thermonuclear reaction. Only those coming from the distance

$$L = \int_{-\infty}^{z} e^{-(z-z')/\lambda_0}\, dz' = \lambda_0 \tag{52}$$

behind the detonation front can contribute to its required energy supply. Here we must put $n = 4n_0$ into the expression for the range equation to account for the fourfold density increase behind the detonation front. Because of the six spatial directions only the fraction 1/6 of these charged fusion products will go in the forward direction. Furthermore, only the fraction f of the energy released in a particular thermonuclear reaction goes into charged fusion products. The energy supplied by the charged fusion products to the detonation wave is therefore given by

$$\varepsilon_c(z) = \frac{f}{6}\int_{z-\lambda_0}^{z} \dot{E}S(z)\, dz = \frac{2}{3} f n_0^{\,2}\langle\sigma v\rangle\varepsilon_0 \int_{z-\lambda_0}^{z} S(z)\, dz. \tag{53}$$

The condition for a self-sustained detonation wave is

$$\varepsilon(z) = \varepsilon_c(z), \tag{54}$$

and thus

$$S(z) = \frac{\sqrt{3/2}}{36}\,\frac{f n_0\langle\sigma v\rangle\varepsilon_0 M^{1/2}}{(kT)^{3/2}}\int_{z-\lambda_0}^{z} S(z)\, dz. \tag{55}$$

This integral equation has the solution

$$S(z) = S_0 e^{z/z_0}, \qquad S_0 = \text{const}, \tag{56}$$

where z_0 is determined by the transcendental equation

$$z_0(1 - e^{-\lambda_0/z_0}) = 36\sqrt{2/3}[(kT)^{3/2}/f n_0\langle\sigma v\rangle\varepsilon_0 M^{1/2}]. \tag{57}$$

The surface equation of the thermonuclear assembly is given by

$$r(z) = \sqrt{S(z)/\pi} = \sqrt{S_0/\pi}\, e^{z/2z_0}. \tag{58}$$

Using the range equation (51) with $n = 4n_0$, we can write for Eq. (57)

$$1 - e^{-x} = cx, \tag{59}$$

where

$$x \equiv \lambda_0/z_0,$$

$$c = c(T) \equiv \frac{144\sqrt{2/3}}{fa\varepsilon_0 M^{1/2}\langle\sigma v\rangle}. \tag{60}$$

The temperature dependence in the expression for c results from the temperature dependence of $\langle\sigma v\rangle$ only. From Eq. (59) it is clear that there can be a growing wave with $z_0 > 0$ only if $c < 1$. For $c = 1$, the solution is $x = 0$, corresponding to a wave of constant cross section moving along a cylinder. For $c > 1$, the solution becomes negative and the wave can only propagate along an assembly of decreasing cross section. An example of an assembly of decreasing cross section is a spherical convergent detonation wave. Therefore, a thermonuclear detonation wave that cannot propagate along an assembly of increasing or constant cross section may still be able to propagate in the form of a convergent spherical detonation. In any case, it is remarkable that the condition that determines whether or not the wave can be a growing one does not depend on the density of the thermonuclear explosive.

The largest positive root of Eq. (59) occurs for the smallest possible value of c, which is reached for the largest value of $\langle\sigma v\rangle = \langle\sigma v\rangle_{\max}$. But only if $\langle\sigma v\rangle$ is larger than the minimum value given by

$$\langle\sigma v\rangle_{\min} = 144\sqrt{\tfrac{2}{3}}/fa\varepsilon_0 M^{1/2}, \tag{61}$$

which is obtained putting $c = 1$ in Eq. (60), can a wave propagate along a thermonuclear assembly of increasing cross section. With this definition of $\langle \sigma v \rangle_{\min}$ we can also write for Eq. (60)

$$c = \langle \sigma v \rangle_{\min} / \langle \sigma v \rangle. \tag{62}$$

The smallest value of c, of course, is

$$c_{\min} = \langle \sigma v \rangle_{\min} / \langle \sigma v \rangle_{\max}.$$

In case of the DT thermonuclear reaction one has

$$f = 0.2,$$

$$\varepsilon_0 = 17.6 \text{ MeV} = 2.82 \times 10^{-5} \text{ erg},$$

$$M = 2.5 \, M_H,$$

where M_H is the proton mass. The constant in the range formula for the charged fusion products, which are α-particles in this case, is

$$a \simeq 2.5 \times 10^{34} \text{ cm}^{-2}\text{-erg}^{-3/2}.$$

The maximum value for $\langle \sigma v \rangle$ here is

$$\langle \sigma v \rangle_{\max} \simeq 10^{-15} \text{ cm}^3/\text{sec}.$$

This maximum value is reached at a temperature of $T \simeq 8 \times 10^8$ K $\simeq 70$ keV. The smallest value still permitting a wave of constant cross section is given by Eq. (61):

$$\langle \sigma v \rangle_{\min} \simeq 4.05 \times 10^{-16} \text{ cm}^3/\text{sec}.$$

This smallest value is reached at a temperature of $T \simeq 1.7 \times 10^8$ K $\simeq 20$ keV, at which a wave of constant cross section is therefore still possible.

For the smallest value of c one has $c_{\min} \simeq 0.405$. For $c = c_{\min}$, Eq. (59) has the maximum possible root

$$x = x_{\max} \simeq 2.2.$$

To appreciate this number, let us discuss the case of a wave propagating in solid DT where $n_0 = 5 \times 10^{22}\ \mathrm{cm}^{-3}$. For $n = n_1 = 4n_0 = 2 \times 10^{23}\ \mathrm{cm}^{-3}$ and $T = 8 \times 10^8$ K, we find from Eq. (51) $\lambda_0 = 4.5$ cm and hence

$$z_0 = \lambda_0 / x_{\max} \simeq 2\ \mathrm{cm}.$$

According to Eq. (58), the radius of an exponential DT horn can increase e-fold over the distance of ~4 cm or 10-fold over the distance of ~9 cm, which shows that a rapidly growing wave is here possible. This fact suggests attaching a long cylinder to a flat ignition cone, both of which are made of DT. In the conical section the wave will grow, and in the cylinder section it will propagate easily with a constant cross section.

In the case of the DD thermonuclear reaction we have to consider the range of three kinds of charged fusion products: protons, tritons, and ^{3}He particles, each of which possesses a different energy. If we average the ranges over the energy of these fusion products, we obtain for the average of the constant entering the range equation (51)

$$a = 4.0 \times 10^{34}\ \mathrm{cm}^{-2}\text{-}\mathrm{erg}^{-3/2}.$$

For the maximum value of $\langle \sigma v \rangle$ we must choose a temperature at which the bremsstrahlung losses do not become too large. Choosing $T \simeq 300$ keV, we have $\langle \sigma v \rangle_{\max} \simeq 5 \times 10^{-17}\ \mathrm{cm}^3/\mathrm{sec}$. For the energy ε_0 we must average over the two equally probable branches of the DD reaction given by Eq. (2). The result is

$$\overline{\varepsilon}_0 = 3.62\ \mathrm{MeV} = 5.8 \times 10^{-6}\ \mathrm{erg}.$$

We also have $f = 0.66$ and $M = 2M_H$, and with these numbers find that $c_{\min} = 1.4 > 1$. It therefore seems to follow that a growing thermonuclear burn wave in deuterium is impossible.

This conclusion, however, is only true as long as we neglect the formation of tritium and helium-3 in the DD thermonuclear reaction. Then if

$$\langle\sigma v\rangle_{\mathrm{DT}} \gg \langle\sigma v\rangle_{\mathrm{DD}} \quad \text{and} \quad \langle\sigma v\rangle_{\mathrm{D^3He}} \gg \langle\sigma v\rangle_{\mathrm{DD}},$$

both the tritium and helium-3 fusion products of the DD reaction will burn rapidly behind the detonation front, adding the energy of their charged fusion products from both the DT and D^3He thermonuclear reactions to drive the detonation front. We therefore have to deal with three additional fusion products, two α-particles and a proton, coming from these secondary thermonuclear reactions. One α-particle is the result of the DT thermonuclear reaction, Eq. (1); the other α-particle and the proton come from the reaction

$$\mathrm{D} + {}^3\mathrm{He} \longrightarrow {}^4\mathrm{He} + \mathrm{p} + 18.4\ \mathrm{MeV}. \tag{63}$$

With these additional fusion products the constant of the range equation averaged over the energies is now

$$\bar{a} \simeq 6.9 \times 10^{34}\ \mathrm{cm}^{-2}\text{-}\ \mathrm{erg}^{-3/2}.$$

To account for the energy of the secondary charged fusion products, we simply equate $f\varepsilon_0$ in the expression for c with the total energy ε_t of all the charged fusion products for both the primary and secondary reactions, averaged over both branches of the DD reaction. We easily find we must put

$$f\varepsilon_0 \longrightarrow \varepsilon_t = 13.4\ \mathrm{MeV} = 2.14 \times 10^{-5}\ \mathrm{erg}.$$

The approximate expression for c then reads

$$c \simeq \frac{72\sqrt{2/3}}{\bar{a}\varepsilon_t M^{1/2}\langle\sigma v\rangle_{\mathrm{DD}}}. \tag{64}$$

We choose a temperature of $T \simeq 300$ keV $\simeq 3.6 \times 10^9$ K, at which $\langle\sigma v\rangle_{DD}$ is still much smaller than both $\langle\sigma v\rangle_{DT}$ and $\langle\sigma v\rangle_{D^3He}$. At this temperature,

$$\langle\sigma v\rangle_{DD} \simeq 5 \times 10^{-17}\ \mathrm{cm^3/sec},$$

and $c \simeq 0.4$. The value of $c \simeq 0.4$ according to Eq. (59) gives the same value for x as in the case of the DT reaction. However, the temperature is much higher here.

This result shows that a rapidly growing thermonuclear detonation wave is possible even in liquid deuterium because of the large energy release by the secondary thermonuclear reactions. The only serious problem here is how to reach the much larger ignition temperature, which is about 10 times higher than for the DT reaction. This high temperature may require more than one fission bomb for ignition.

For a thermonuclear detonation propagating in ^{6}LiD the situation is even more complex. As can be seen from Eq. (4), in addition to thermonuclear reactions neutron reactions are also involved here. To analyze this situation we first assume that all the ^{6}Li in front of the detonation wave is converted into tritium and ^{4}He by the reaction

$$\mathrm{n} + {}^6\mathrm{Li} \longrightarrow \mathrm{T} + {}^4\mathrm{He}.$$

The conversion of ^{6}Li into T and ^{4}He takes place through the lower branch of Eq. (4). The neutrons that cause the conversion are supplied by the upper branch of Eq. (4), produced by DT thermonuclear reactions in the detonation wave. Under the second assumption, that no neutrons are lost, we can then show that the first assumption of complete conversion of ^{6}Li into T and ^{4}He is justified.

In reality, however, there will always be some neutron losses through the finiteness of the geometry, and the second assumption of complete conversion cannot be valid. But these losses can be compensated by surrounding the thermonuclear explosive with some neutron-multiplying reflector, notably ^{238}U but also ^{9}Be. This situation is schematically illustrated in Figure 9. Some neutrons produced by the DT reactions taking place within the detonation wave are shot into the as-yet-unburned explosive, where they convert the ^{6}Li into T and ^{4}He. Some of the neutrons will be lost from the assembly, but other neutrons will be reflected back. Still other neutrons will make neutron-multiplying reactions with some of the neutrons reflected back.

The reaction mean free path of the neutrons penetrating into the ^{6}Li positioned in front of the detonation wave is given by

$$\lambda_n = 1/n_0\sigma_n, \tag{65}$$

where $\sigma_n \simeq 3 \times 10^{-25}$ cm^2 is the reaction cross section for the 14-MeV DT thermonuclear fusion neutrons coming from the burn zone in the wave. In the case of a plane detonation wave, only one-sixth of the neutrons produced in the burn zone will be emitted in the forward direction, the other five-sixths going in the remaining five spatial directions. If only that many neutrons were available the

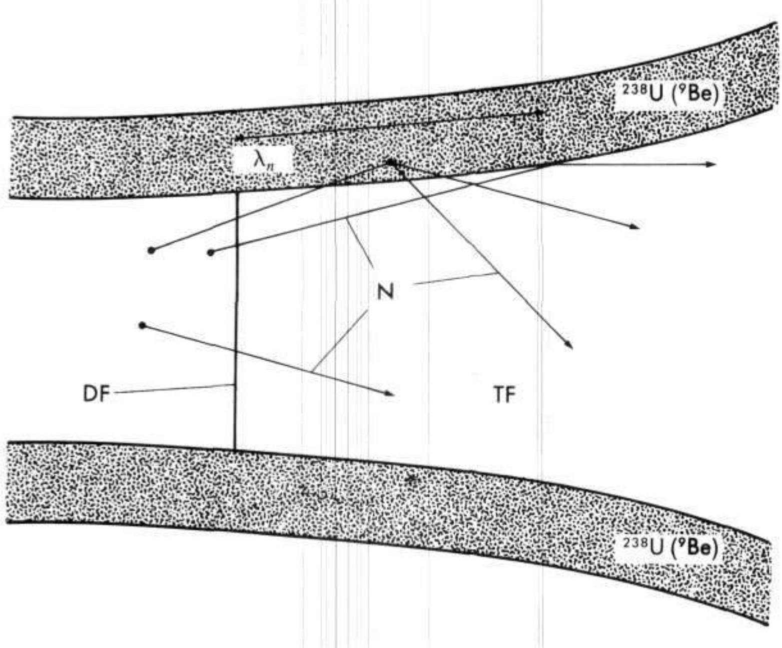

Figure 9. The dry H-bomb principle using ^{6}LiD as the explosive: the tritium breeding is accomplished here during the explosion process by the neutron precursor from the thermonuclear detonation wave acting with a ^{238}U or ^{9}Be neutron multiplier and transforming the ^{6}LiD explosive TF into TD. DF is the detonation front.

detonation would soon die out from a lack of neutrons to support the closed chain shown in Eq. (4). Both ^{238}U and ^{9}Be are good neutron reflectors. Even a 100% neutron albedo would be insufficient, however, because this still would mean that only one-half of all neutrons would go in the forward direction. Therefore, the reflector must have an albedo greater than 2. This is only possible, of course, if neutron-multiplying reactions take place within the reflector.

The neutron cross section for 14-MeV neutrons in ^{238}U is about $\sigma_f \simeq 10^{-24}$ cm^2 with an approximately twofold neutron yield per fission. ^{238}U also has a large $(n, 2n)$ cross section, which is

$$\sigma_{n,2n} \sim 1.6 \times 10^{-24}\ \mathrm{cm}^2,$$

peaking at $\sim$10 MeV. Since the atomic number density in ^{238}U is 5×10^{22} cm^{-3}, it follows that a ^{238}U blanket a few centimeters thick would make a neutron-multiplying reflector of the required albedo. Similar numbers result if ^{9}Be is used instead.

The neutron flux entering the explosive positioned in front of the wave is then given by

$$\phi = (n_1^2/4)\langle\sigma v\rangle\lambda_n = 4n_0\langle\sigma v\rangle/\sigma_n. \qquad (66)$$

Inserting into this

$$n_0 = 5 \times 10^{22}\ \mathrm{cm}^{-3} \quad \text{and} \quad \langle\sigma v\rangle \simeq 10^{-15}\ \mathrm{cm}^3/\mathrm{sec},$$

one finds

$$\phi \simeq 2 \times 10^{32}\ \mathrm{cm}^{-2}\text{-}\mathrm{sec}^{-1}.$$

The rate at which the ^{6}Li is converted into T and ^{4}He is given by

$$dn_0/dt = n_0\sigma_n\phi = 4n_0^2\langle\sigma v\rangle. \qquad (67)$$

From this the characteristic conversion time follows:

$$\tau_c = n_0(dn_0/dt)^{-1} = (4n_0\langle\sigma v\rangle)^{-1}. \qquad (68)$$

For the example given above, $\tau_c \simeq 5 \times 10^{-9}$ sec. This time must be compared with the time τ_D it takes for the detonation wave to propagate over one neutron-reaction mean free path. Then, if

$$\tau_c \ll \tau_D,$$

all the ^{6}Li in front of the wave will be converted into T and ^{4}He before entering the wave. The time τ_D is given by

$$\tau_D = \lambda_n / v_0, \tag{69}$$

where v_0 is given by Eq. (42). For the given numbers

$$n_0 = 5 \times 10^{22}\ \text{cm}^{-3} \quad \text{and} \quad \sigma_n = 3 \times 10^{-25}\ \text{cm}^2,$$

we find $\lambda_n = 67$ cm. For $T = 70$ keV, where

$$\langle \sigma v \rangle \simeq 10^{-15}\ \text{cm}^3/\text{sec},$$

we have

$$v_0 = 5.3 \times 10^8\ \text{cm/sec},$$

and hence

$$\tau_D = 1.3 \times 10^{-7}\ \text{sec}.$$

In our case this condition is well satisfied, and the first assumption is verified. Since in general the second assumption is not satisfied by a neutron-multiplying reflector, the correctness of the model is established.

Two remarks are in order. In our estimates we have assumed that the atomic and molecular number densities for ^{6}Li, D, T, and ^{6}LiD are the same and equal to 5×10^{22} cm^{-3}. This, of course, is only approximately true. The other remark is that a detonation wave in ^{6}LiD differs from one in DT because one-third of the material entering the wave is ^{4}He. The presence of ^{4}He reduces the range of the charged DT fusion products and hence the range constant. The range is inversely proportional to the stopping power, which is proportional to the charge Z of the plasma ions. Therefore, the average of $1/Z$ for a plasma composed of two kinds of ions, one with charge Z_1 and relative fraction α and the other with charge Z_2 and relative fraction β, is given by

$$\frac{1}{\bar{Z}} = \frac{1}{\alpha Z_1 + \beta Z_2}. \tag{70}$$

In our case $Z_1 = 1$, $\alpha = 2/3$, $Z_2 = 2$, $\beta = 1/3$, and we find

$$(\bar{Z})^{-1} = 3/4,$$

by which the range constant must be multiplied. Since the range constant enters Eq. (60) in the denominator, this increases the value of c in comparison to the value of c for the DT reaction by the factor 4/3. The value

$$c_{\min} \simeq 0.405$$

for the DT reaction is thus here increased to

$$c_{\min} \simeq 0.54.$$

For this value the root of Eq. (59) is now

$$x = x_{\max} = 1.4,$$

and thus almost 60% smaller. To compute from this the value of $z_0 = \lambda_0/x_{\max}$ we must also correct the value of λ_0, which is reduced from its value of $\lambda = 4.5$ cm to

$$\lambda_0 = 4.5 \times (3/4) = 3.4 \text{ cm}.$$

Thus $z_0 = 2.4$ cm, which is only slightly larger than the value $z_0 = 2$ cm computed for the DT reaction. It thus follows that in ^{6}LiD a rapidly growing thermonuclear detonation wave is possible. This is a fact of great military importance because an H-bomb using ^{6}LiD as the explosive is "dry" and does not require a bulky refrigeration system.

CHAPTER ELEVEN

Various Configurations To Reach Large Thermonuclear Yields

With the results of the preceding chapters we are now in a position to make various designs that can be used in a practical thermonuclear explosive device. All the proposed systems have in common three components, shown schematically in Figure 10, which are: (1) the ignition module I containing the fission trigger, (2) the conical transition module C enlarging the thermonuclear deflagration front, and (3) the main thermonuclear explosive module T. All components are, of course, surrounded by a tamp.

The kind of ignition module used, whether it be the one with the Prandtl-Meyer ellipsoid or the one using the flat cone, depends on a number of technical considerations. An ignition module using the Prandtl-Meyer ellipsoid should be the most efficient configuration, but it is quite bulky. Therefore, for more compact thermonuclear explosive devices the flat cone may be advantageous, even though its efficiency is lower. (An even more efficient and compact ignition system is discussed separately in Chapter 18.)

In the case of a pure deuterium explosion, where the ignition temperature is high, one will probably have to use the Prandtl-Meyer ellipsoid configuration. The most simple configuration would be a pure DT bomb, because it would not require surrounding the ignition point with fissile material as would be required for the ^{6}LiD

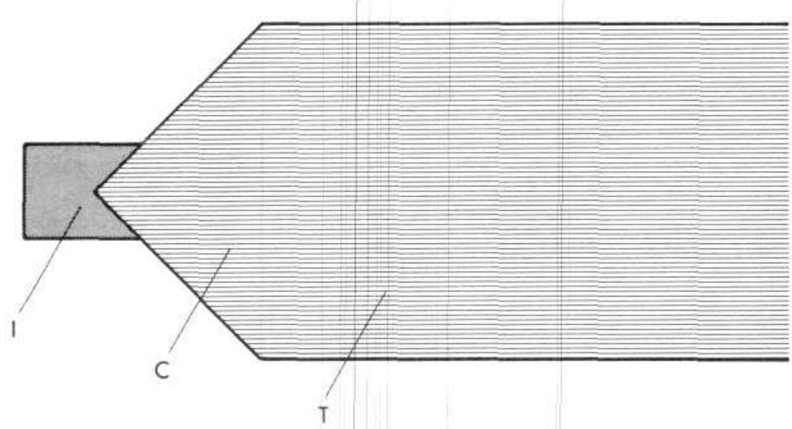

Figure 10. Schematic arrangement of the ignition module I containing the fission trigger, the conical transition module C that enlarges the thermonuclear deflagration front, and the main thermonuclear explosive module T—to reach an arbitrarily large thermonuclear yield.

H-Bomb

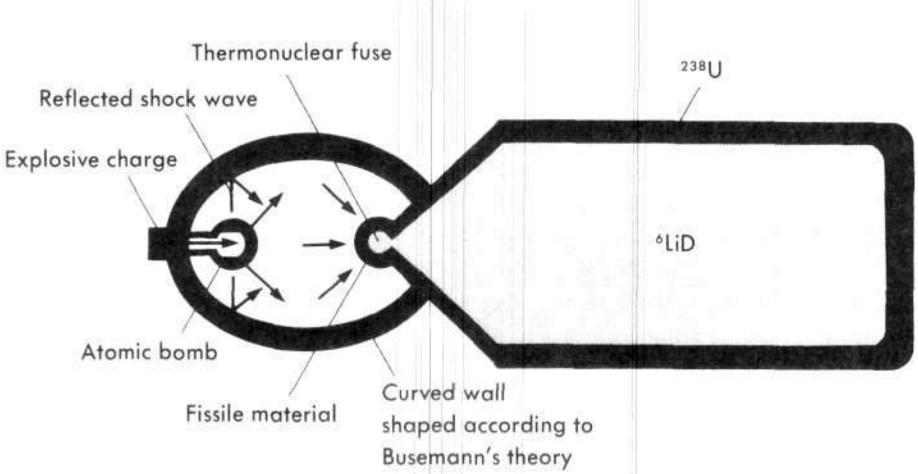

Figure 11. Design of dry H-bomb that I proposed in 1952. In this concept, ignition is accomplished using a Prandtl-Meyer ellipsoid with a thermonuclear exponential horn and a cylinder.

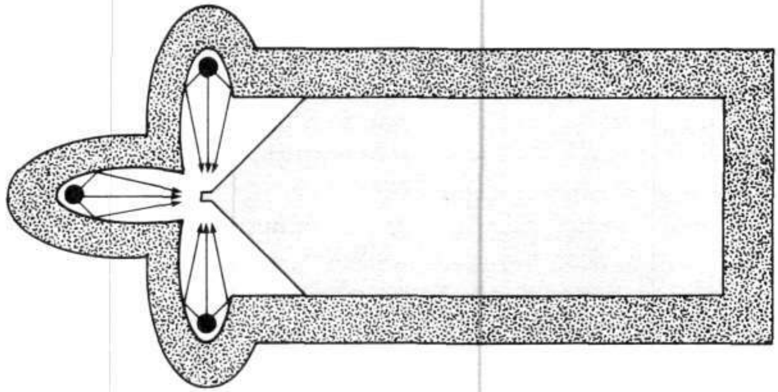

Figure 12. Ignition of a pure deuterium explosion (without any tritium) using several fission bombs. This is the tetrahedron configuration along with a Guderley convergent shock.

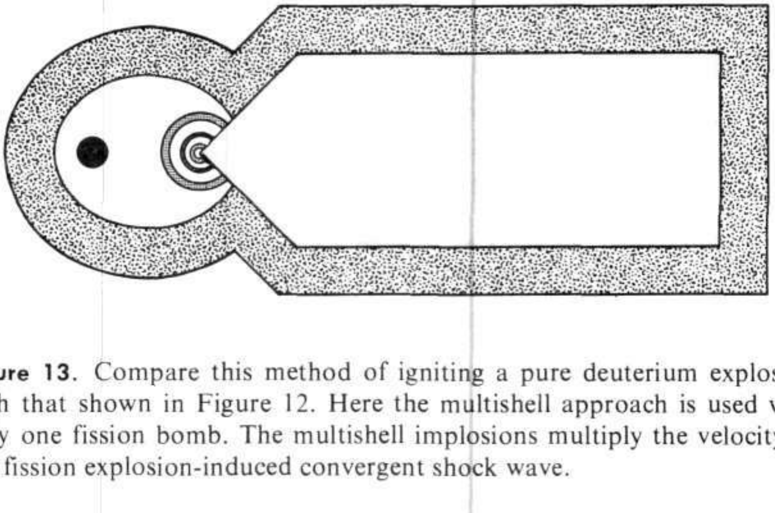

Figure 13. Compare this method of igniting a pure deuterium explosion with that shown in Figure 12. Here the multishell approach is used with only one fission bomb. The multishell implosions multiply the velocity of the fission explosion-induced convergent shock wave.

bomb. Figure 11 shows one way in which a ^{6}LiD bomb would work.

More difficult is the ignition of a pure DD bomb. One method is to use DT *in the ignition region and* in that part of the transition cone located near the ignition point, which then, as the detonation front propagates down the transition cone, changes slowly into pure deuterium by continuously changing the D/T ratio along the cone.

A method for igniting a pure deuterium explosion without using any tritium is shown in Figure 12, using the tetrahedron configuration in conjunction with a Guderley convergent shock.

Another method for the ignition of a pure deuterium explosion is shown in Figure 13. Here the multishell approach is used, requiring only one fission explosive.

All these configurations have in common that their thermonuclear yield can be conveniently varied by varying the length of the main thermonuclear fuel cylinder.

For the MIRV concept (multiple independently targeted reentry vehicles), requiring a cluster of hydrogen bombs on the top of a rocket launcher, one can give each bomb a hexagonal cross section, as it is used for incendiary cluster bombs. For MIRV configurations each separate fusion bomb can be chosen to be shaped rather long, like a hexagonal pencil.

CHAPTER TWELVE

The Neutron Bomb

In the so-called neutron bomb, the intention is to deliver as large a portion of the energy as possible as neutrons. The purpose is obvious: Neutrons have a large lethal biological effect but do not produce a blast wave. The reason for this behavior is the long path length of neutrons in matter, which can be several meters at solid densities. Hence, a hypothetical nuclear reaction producing neutrons exclusively would be "ideal." In practice, though, there is no known self-sustaining nuclear reaction producing neutrons exclusively, and even the optimal neutron bomb will always produce some blast.

One of the best candidates for a neutron bomb is a DT thermonuclear reaction, since there 80% of the energy released goes into neutrons and only 20% into α-particles. But even there the neutron yield is less than 80% because some portion of the energy is always released by the fission trigger. The problems connected with liquid DT can be largely avoided by giving the chemical composition of the bomb the following combination:

$$(1/3)^{6}\mathrm{LiT} + (2/3)^{7}\mathrm{LiD}.$$

In this case one-third of the material can burn as a DT reaction, setting free one neutron, which can react with the one-third of ^{6}Li to make one-third of T to react with the remaining one-third of D. Here, however, the overall relative neutron yield is less than 80% because the relative one-third fraction of the reaction energy for

$$\mathrm{n} + {}^{6}\mathrm{Li} \longrightarrow \mathrm{T} + {}^{4}\mathrm{He} + 4.8\ \mathrm{MeV},$$

that is, $4.8/3 = 1.6$ MeV, goes into charged fusion products, reducing the relative neutronic energy release from 80% to 74.4%. The actual neutron yield, however, is still smaller since some of the neutrons are always lost through the surface of the thermonuclear assembly.

These losses can be reduced, of course, by a neutron reflector. The use of ^{238}U for this purpose is unsuitable. Although it can have an albedo greater than 1 it will reduce the overall relative neutron yield because of the fission processes in ^{238}U, with most of the energy released in the form of charged fission products rather than neutrons. Instead of ^{238}U one may use a neutron reflector ^{9}Be, which undergoes the reaction

$$^{9}\mathrm{Be} + \mathrm{n} \longrightarrow {}^{8}\mathrm{Be} + 2\mathrm{n} - 1.666\ \mathrm{MeV}. \tag{71}$$

Beryllium-8 subsequently decays into two α-particles with an energy release of 0.096 MeV. The overall reaction is thus endothermic, with an absorbed energy of 1.57 MeV. In fact, the neutron-multiplying reaction being endothermic here does not increase the explosive yield as in the case of a ^{238}U neutron reflector. The cross section for the reaction (71) at a neutron energy of 14 MeV, that is, the energy of DT fusion neutrons, is

$$\sigma_{n,2n} \simeq 0.65 \times 10^{-24}\ \mathrm{cm}^2.$$

This cross section is sufficiently large for our purpose. Besides ^{9}Be one may use other substances with a large $(n,2n)$ cross section, for example, niobium, which has about the same cross section as ^{9}Be for 14-MeV neutrons.

Figure 14 shows the principle of such a neutron bomb. It differs from the ordinary H-bomb only in replacing ^{6}LiD with ^{6}Li$_2$DT and the ^{238}U shell with ^{9}Be.

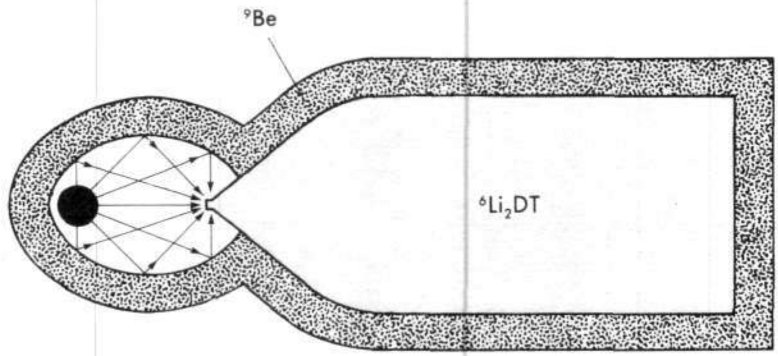

Figure 14. In the neutron bomb, the thermonuclear fuel is $^6Li_2DT$, replacing the ^6LiD of the ordinary H-bomb. The ^9Be shell serves as a neutron-multiplying reflector so that most of the energy goes into neutrons, with little blast effect.

A substantial reduction in the size of the neutron bomb, but also of any other thermonuclear explosive, is possible with the autocatalytic principle explained in the following chapter. Because one goal of the neutron bomb is to make it as small as possible, the autocatalytic concept is especially important here. However, since the autocatalytic principle also promises to raise the overall yield it is the method to be used for all advanced thermonuclear explosive concepts.

To reduce the blast as much as possible and thereby increase the relative neutron yield, one may use a fission fizzle for the neutron bomb as a trigger for the fusion reaction (see Chapter 18).

CHAPTER THIRTEEN

Autocatalytic Thermonuclear Detonation Waves

Until now we have assumed that the thermonuclear fuel cylinder is uncompressed during its burn. But since the thermonuclear reaction rate increases with the square of the density, even a modest compression will greatly increase the reaction rate. If, for example, only an $\sim$10-fold density increase is achieved the reaction rate will be increased $\sim$100-fold. The problem, of course, is how the fuel can be compressed by even a modest amount.

The solution to this problem is the *autocatalytic process*. The name "autocatalytic" is adopted from the Los Alamos "primer," consisting of unclassified notes by the late E.U. Condon. In this process, the energy delivered in the detonation wave itself can be used to compress the yet-unignited thermonuclear fuel. Such a process, provided it can be realized, shall be called *autocatalytic*. The significance of this process is that it permits us to build small thermonuclear explosive devices, as they are used in the neutron bomb.

There are two possible alternatives for realizing this. The first is explained in Figure 15, showing a detail of the assembly near the position of the detonation front DF. Ahead of the front is the unburned thermonuclear fuel TF and behind it the thermonuclear burn zone BZ. The entire assembly is surrounded by a tamp T. Furthermore, the fuel assembly is separated by a small gap G from the tamp T.

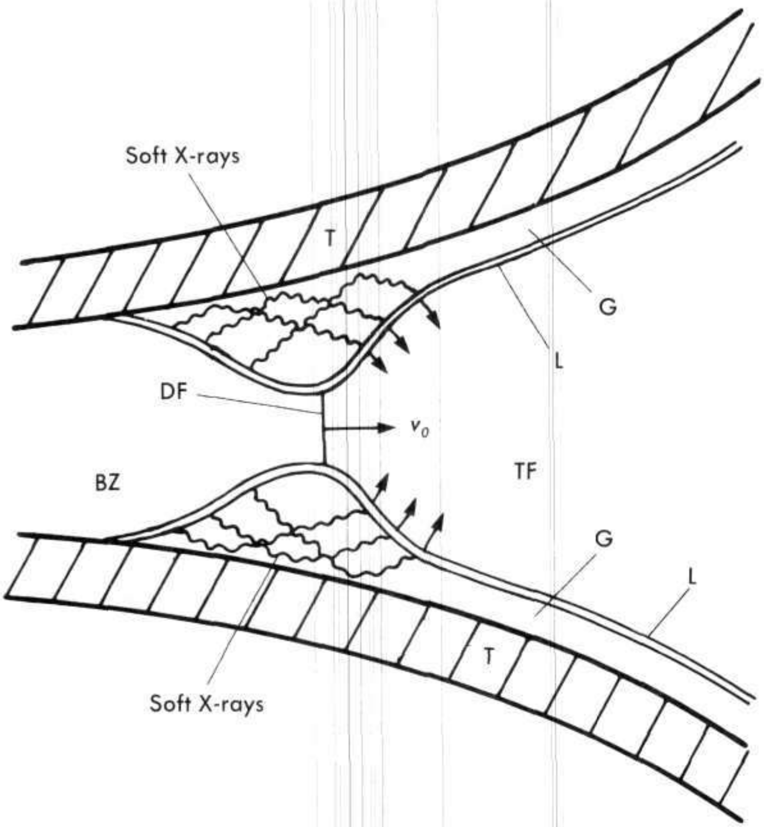

Figure 15. Autocatalytic thermonuclear detonation using a soft X-ray precursor to precompress the thermonuclear fuel TF before the detonation front DF from the burn zone BZ reaches it. The soft X-rays travel through the gap G between the tamp T and the liner L.

Now let us contemplate what use can be made of the X-rays emitted as bremsstrahlung from the thermonuclear burn zone BZ. Part of this radiation will be absorbed in the tamp T, and part of it reflected back into the thermonuclear plasma. Furthermore, because of the particular geometry, part of the radiation will propagate along the channel formed by the gap separating the fuel from the tamp. Finally, some part of the radiation will be absorbed in the thermonuclear fuel TF located ahead of the detonation front DF. If that portion of the bremsstrahlung radiation that is transported along the gap is large, and if it is absorbed in a surface layer of some fraction of the unburned fuel located in front of the detonation wave, this portion of the fuel will be compressed by ablative implosion. The emission of the radiation into the gap space G takes place on a much shorter time scale than the closure of the gap by thermal plasma expansion, justifying the assumption of radiation flow within the gap.

The photon energy of the bremsstrahlung radiation for a hydrogen plasma at thermonuclear temperatures will be of the order of many kiloelectron volts. At these energies the reflection of the photons from the surface of the tamp is small. The situation is changed, however, if the thermonuclear fuel is surrounded by a high-atomic-number liner L. Such a high-atomic-number liner, in contact with the hot, multi-keV thermonuclear plasma, will be heated to high temperatures and thereby become a secondary source of much softer X-rays with a higher reflection coefficient. The transport of soft X-rays along the gap space therefore acts as a high-intensity photon precursor, precompressing the fuel in front of the detonation wave. A detailed analysis of this precursor is obviously very difficult and can be done by extensive computer modeling only. Such a detailed analysis is beyond the scope of this book, which emphasizes only the physics aspect.

To make a simple estimate of the precursor effect we consider the energy flux ϕ of soft secondary X-rays emitted from the burn zone BZ through action of the liner L. The detonation front is located at $z' = z$, and the diffusion of these X-rays for $z' > z$ within the gap G separating the liner L from the tamp T is then governed by the equation

$$d\phi/dz' = -(1 - R)(S_g/V_g)\phi. \tag{72}$$

In Eq. (72), R is the average reflection coefficient for the soft X-rays, and S_g and V_g are the surface and volume per unit length, respectively, along the z-axis of the gap G separating the liner from the tamp. For Eq. (72) to be valid the detonation front is assumed to be at rest. This is not true, of course, but Eq. (72) is still a good approximation as long as the detonation velocity is small compared to the velocity of light. In our case $v_0/c \sim 10^{-2}$, and the approximation is well justified.

Assuming that

$$L \equiv V_g/S_g(1 - R) = \text{const},$$

we obtain from Eq. (72)

$$\phi(z') = \phi(z)e^{-(z'-z)/L}. \tag{73}$$

The X-ray flux at the position z is calculated in the following way. The energy emitted by bremsstrahlung from the DT plasma is given by[9]

$$\varepsilon_r = 1.42 \times 10^{-27} n^2 \sqrt{T} \text{ erg/cm}^3\text{-sec}. \tag{74}$$

At an energy of $\sim$10 keV, corresponding to a thermonuclear temperature of $\simeq 10^8$ K, the bremsstrahlung has an absorption length in material of density ρ (in grams per cubic centimeter) equal to $\lambda_a \simeq 10^{-2}/\rho$ cm. Hence, with a liner having a density $\rho \sim 10$ g/cm^3 the absorption length is $\sim 10^{-3}$ cm, which has to be smaller than the thickness of the liner.

To calculate the energy of the softer secondary radiation emitted from the liner we proceed as follows: The energy in the form of bremsstrahlung radiation emitted through the surface of unit length and radius r is given by $\simeq \varepsilon_r r/2$. Assume, for example, that for the emitted bremsstrahlung the fuel density is equal to solid density, that is, $n = n_0 = 5 \times 10^{22}$ cm^{-3}; furthermore, $T = 10^8$ K, and $r = 10^{-1}$ cm. It then follows that

$$\varepsilon_r r/2 \simeq 2 \times 10^{21} \text{ erg/cm}^2\text{-sec}.$$

This energy flux has to be set equal to the energy flux from the liner heated to the temperature T_l. Since the liner will be optically opaque, it radiates according to the Stefan-Boltzmann law

$$\varepsilon_l = \sigma T_l^4,$$

where $\sigma = 5.75 \times 10^{-5}$ erg/cm^2-sec-K. Equating

$$\varepsilon_r(r/2) \simeq \varepsilon_t$$

we find that

$$T_t \simeq 2 \times 10^6 \text{ K} \simeq 200 \text{ eV},$$

corresponding to an $\simeq$50-fold increase in the wavelength of the secondary radiation entering the gap space G.

The photon energy flux delivered from the burn zone and reaching the detonation front is given by

$$E_p \simeq \tfrac{1}{2}\varepsilon_r SL, \tag{75}$$

where we have assumed that an equal amount of energy goes into the forward and backward directions and that the velocity of the detonation front is small compared to the velocity of light. From Eq. (75) we then obtain for the photon energy flux $\phi(z)$ at the position of the detonation front

$$\phi(z) \simeq E_p/S_g L = \tfrac{1}{2}\varepsilon_r(S/S_g). \tag{76}$$

Now, since $S = \pi r^2$ and $S_g \simeq 2\pi r$,

$$\phi(z) \simeq \tfrac{1}{4}\varepsilon_r r. \tag{77}$$

With the same parameters of n_0 and T as before to compute ε_r, we find that

$$\phi(z) \simeq 5 \times 10^{20} \text{ erg/cm}^2\text{-sec} = 5 \times 10^{13} \text{ W/cm}^2.$$

However, since the thermonuclear fuel in that part of the burn zone located immediately behind the detonation front is highly compressed, the value of $\phi(z)$ should actually be much higher. Assuming, for example, an average density of the burn zone 10 times solid density, the power level will be increased $\sim$100-fold to 5×10^{15} W/cm^2.

The Lawson criterion, Eq. (9), requires that

$$\rho \propto r^{-1} \propto S^{-1/2}. \tag{78}$$

The compression energy per unit length of the fuel is

$$E_c \propto pS, \tag{79}$$

where

$$p \propto \rho^{5/3}$$

is the required ablation pressure for isentropic compression. Combining Eqs. (78) and (79) we find

$$E_c \propto S^{1/6} \propto r^{1/3}. \tag{80}$$

The compression energy delivered by the photons to the fuel surface is proportional to $\phi(z)S_g$. But since

$$\phi(z) \propto \rho^2 r \quad \text{and} \quad S_g \propto r,$$

it follows that

$$\phi(z)S_g \propto (\rho r)^2 = \text{const.}$$

Because the required compression energy is proportional to $r^{1/3}$ the fuel must be slightly overcompressed initially to compensate for the rather insensitive dependence of E_c on r.

Finally, it can be shown that the time dependence of the power profile of the photon flux with respect to some fuel at a fixed position to be compressed by this photon flux is qualitatively similar to the one required for high-density compression, which is given by[10]

$$E_c(t) \simeq E_c(0)(1 - t/t_0)^{-15/8}. \tag{81}$$

The photon energy flux ϕ at a fixed fuel location $z' = z_0$ as a function of time is obtained by putting $z = v_0 t$ into Eq. (73). This means that at the time $t = 0$ the detonation front is at the position $z = 0$. We obtain

$$\phi(t) = \phi(z)e^{-z_0/L}e^{v_0 t/L} = \phi(z_0)e^{v_0 t/L}. \tag{82}$$

Here we introduced the definition

$$\phi(z_0) \equiv \phi(z)e^{-z_0/L},$$

for the X-ray flux at the position $z = z_0$ at the time $t = 0$. Now, because $E_p(t)$ is proportional to $\phi(t)$, we can compare the expressions for $E_c(t)$ and $\phi(t)$. For this we first expand Eq. (81), yielding

$$E_c(t) \propto \left(1 + \frac{15}{8}\frac{t}{t_0} + \frac{15 \times 7}{2 \times 8^2}\left(\frac{t}{t_0}\right)^2 + \cdots\right), \tag{83}$$

and then we expand Eq. (82):

$$\phi(t) \propto (1 + v_0 t/L + v_0^2 t^2/2L^2 + \cdots). \tag{84}$$

Comparison shows that both power profiles are congruent up to the first order if

$$15/8t_0 = v_0/L. \tag{85}$$

With this adjustment the second-order term in the expansion (84) becomes $15^2 t^2/(2 \times 8^2 t_0^2)$, which is about a factor 2 larger than the required value of the coefficient according to Eq. (83).

From experimental data for 100-eV X-rays one deduces that for perpendicular incidences $R \lesssim 10^{-2}$ and for glancing incidence $R \simeq 0.1$. For X-ray fluorescent resonance scattering, the reflection coefficient may actually be substantially larger. But in spite of this possibility we shall conservatively assume that $R \ll 1$ and thus put $L \simeq V_g/S_g$. Using, for example, $V_g/S_g \simeq 10^{-2}$ cm, and $v_0 \simeq 10^8$ cm/sec, we find from Eq. (85)

$$t_0 = (\tfrac{15}{8})L/v_0 \simeq 2 \times 10^{-10} \text{ sec},$$

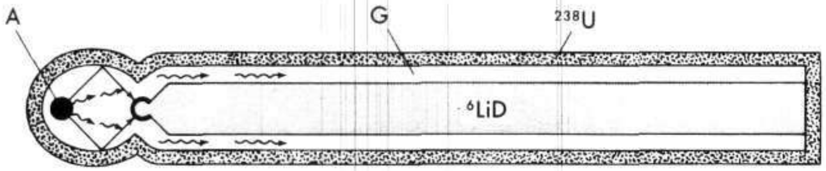

Figure 16. H-bomb using the autocatalytic principle, in which the atom bomb A sends soft X-rays through the gap G between the ^{238}U liner and the ^{6}LiD thermonuclear fuel.

which turns out to be of the right order of magnitude as required for isentropic high-density compression.

Figure 16 shows one possible way in which this concept can be incorporated into a complete weapons design.

The second alternative to generate a precursor precompressing the thermonuclear fuel in front of the detonation front is by using the large neutron flux released from the burn zone. The principle of this idea is shown in Figure 17, where F is the material surrounding the fuel assembly that serves as a tamp but, in addition, also reacts with neutrons. Here there is no need for a gap space as required for the soft X-ray precursor. The material F can consist of some light material with a large neutron cross section, for example, ^{6}Li, ^{9}Be, or ^{10}B. Boron-10 has the advantage that it is abundant and cheap. Of course, F may also consist of fissile material.

The number of neutrons produced in the burn zone per unit time and volume is given by

$$\dot{Q} = (n_1^2/4)\langle\sigma v\rangle = 4n_0^2\langle\sigma v\rangle. \qquad (86)$$

To estimate the neutron flux on the surface of the detonation front, we shall assume that the neutrons are emitted from a fuel cylinder of equal height and diameter. That this assumption is reasonable can be seen as follows:

The inertial confinement time of the compressed fuel cylinder is of the order $\tau \simeq r/a_s$, where

$$a_s = \sqrt{10kT/3M}$$

is the adiabatic sound velocity. Therefore, the cylinder will remain inertially confined behind the detonation front over the distance

$$\Delta z = v_0\tau = \sqrt{32/10}\, r = 1.8r \simeq 2r.$$

Setting $\Delta z = 2r$, the volume of the cylinder with a height Δz is $V = 2\pi r^3$ and its surface is $S = 6\pi r^2$. If the neutron-emitting volume were a sphere of radius r, the neutron flux on its surface would be

$$\phi(r) = \dot{Q}V/S = \dot{Q}r/3,$$

where $V = (4\pi/3)r^3$ and $S = 4\pi r^2$. In the case of the fuel cylinder $V/S = r/3$, and it is therefore reasonable to assume that the neutron flux at the detonation front has the same value as for a sphere with the same radius as the cylinder.

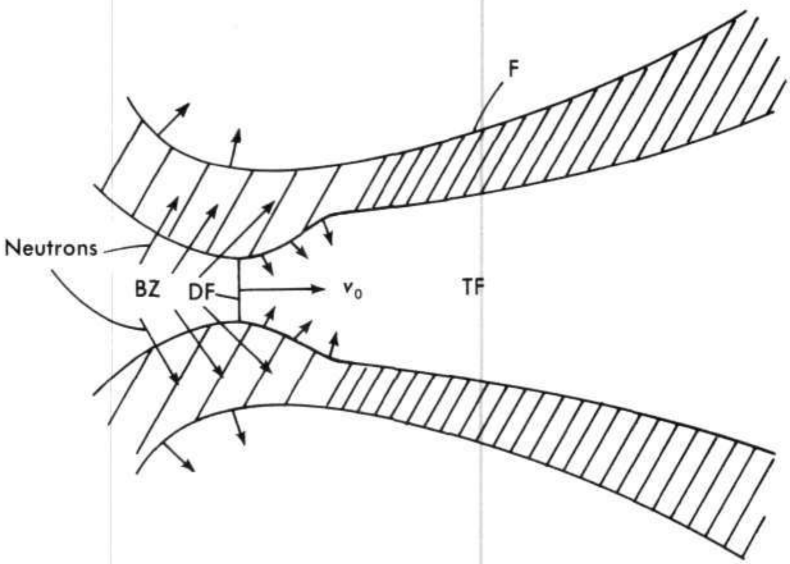

Figure 17. Autocatalytic thermonuclear detonation using a neutron precursor to precompress the fuel TF ahead of the detonation front DF moving with velocity v_0. BZ is the burn zone.

Since the detonation front is positioned at $z' = z$,

$$\phi(z) = \dot{Q}r/3.$$

The neutron flux at the position $z' > z$ is then given simply by the inverse square law

$$\phi(z') = \phi(z)\left(\frac{r}{z' - z + r}\right)^2. \tag{87}$$

As before in the case of the soft X-rays, we consider the neutron flux as a function of time at the fixed position $z' = z_0$ by putting $z = v_0 t$ into Eq. (87):

$$\begin{aligned} \phi(t) &= \frac{\phi(z)}{(1 + z_0/r)^2} \frac{1}{[1 - v_0 t/(z_0 + r)]^2} \\ &= \frac{\phi(z_0)}{[1 - v_0 t/(z_0 + r)]^2}. \end{aligned} \tag{88}$$

Here, as before in the case of X-rays, we introduce the definition

$$\phi(z_0) \equiv \frac{\phi(z)}{(1 + z_0/r)^2},$$

for the neutron flux at the position z_0 and time $t = 0$. In the material F the neutrons are reacting at the rate $n_f \sigma_f \phi(t)$, and if an energy ε_f is set free for each reaction in the form of a short-range charged particle, then, with ε_n the kinetic energy of the reacting neutron, it will result in a heat source of strength

$$\dot{q} = n_f \sigma_f \phi(t)(\varepsilon_f + \varepsilon_n). \tag{89}$$

Therefore, the total amount of heat released at the position z_0 up to the time $t = 0$ is given by

$$q = \int_{-\infty}^{0} \dot{q}\, dt = n_f \sigma_f (\varepsilon_f + \varepsilon_n)\phi(z_0)(z_0 + r)/v_0. \tag{90}$$

The maximum amount of heat q_{max} is released at the position $z_0 = 0$ located at the detonation front:

$$\begin{aligned} q_{\max} &= n_f \sigma_f (\varepsilon_f + \varepsilon_n) \phi(z) r / v_0 \\ &= \tfrac{1}{3} n_f \sigma_f (\varepsilon_f + \varepsilon_n) \dot{Q} r^2 / v_0 \qquad (91) \\ &= \tfrac{4}{3} n_f \sigma_f (\varepsilon_f + \varepsilon_n) n_0^2 \langle \sigma v \rangle r^2 / v_0. \end{aligned}$$

From this value one can compute the maximum attainable temperature T_{max} in F by equating q_{max} with aT_{max}^4, where $a = 7.67 \times 10^{-15}$ erg/cm^3-K^4.

To evaluate Eq. (91) we take as an example a DT plasma with

$$\langle \sigma v \rangle \simeq 10^{-15}\ \mathrm{cm^3/sec}, \qquad \varepsilon_n = 14.1\ \mathrm{MeV} = 2.25 \times 10^{-5}\ \mathrm{erg},$$

and

$$n_0 = 10 n_s = 5 \times 10^{23}\ \mathrm{cm^{-3}}.$$

Furthermore, we shall assume

$$n_f \simeq 5 \times 10^{22}\ \mathrm{cm^{-3}}, \qquad v_0 \simeq 10^8\ \mathrm{cm/sec}, \qquad \text{and } r = 10^{-1}\ \mathrm{cm}.$$

This leads to

$$q_{\max} \simeq 1.6 \times 10^{45} \sigma_f (\varepsilon_f + 2.25 \times 10^{-5})\ \mathrm{erg/cm^3}.$$

We then consider two cases of special interest.

In the first case F consists of ^{10}B undergoing the reaction

$$^{10}\mathrm{B} + \mathrm{n} \longrightarrow {}^{7}\mathrm{Li} + {}^{4}\mathrm{He},$$

with a reaction energy of

$$\varepsilon_f = 3.0\ \mathrm{MeV} = 4.8 \times 10^{-5}\ \mathrm{erg},$$

and a neutron cross section

$$\sigma_f \simeq 1.5 \times 10^{-24}\ \mathrm{cm}^2$$

at 14 MeV. We thus obtain

$$q_{\max} \simeq 6.5 \times 10^{16}\ \mathrm{erg/cm}^3,$$

and from this

$$T_{\max} \simeq 6 \times 10^7\ \mathrm{K}.$$

In the second case, the substance F consists of fissile material, but preferably not the expensive fissile materials ^{235}U, ^{239}Pu, or ^{233}U but, rather, the inexpensive materials ^{238}U or ^{232}Th. Since the neutrons released in the DT reaction have an energy of about 14 MeV they are well above the threshold for fast neutron fission of both substances. At 14 MeV the neutron cross section for ^{238}U is $\sigma_f \simeq 1.2 \times 10^{-24}$ cm^2, and with a reaction energy of

$$\varepsilon_f \simeq 180\ \mathrm{MeV} = 2.9 \times 10^{-4}\ \mathrm{erg},$$

we find

$$q_{\max} \simeq 6 \times 10^{17}\ \mathrm{erg/cm}^3,$$

and hence $T_{\max} \simeq 9 \times 10^7$ K. At these temperatures the maximum pressure is primarily determined by the radiation pressure $p_r = aT_{\max}^4/3$, which for the two given examples is $p_r \simeq 2 \times 10^{16}$ dyn/cm^2 and $p_r \simeq 2 \times 10^{17}$ dyn/cm^2. These maximum pressures are sufficiently high to ensure large fuel precompression.

In using ^{235}U, ^{239}Pu, or ^{233}U the fission cross section would be roughly three times larger. This would increase $q_{\max}$ and p_r by the same factor but $T_{\max}$ only by the factor $3^{1/4} \simeq 1.3$. Therefore, no advantage is seen in going to these very expensive materials except to increase the fission yield. In the neutron bomb, where a low fission yield is desired, the best substance is probably ^9Be because its $(n,2n)$ reaction simultaneously acts as a neutron multiplier.

The power flow in the implosion process must be proportional to $\dot{q}$ and hence to $\phi(t)$ as given by Eq. (88). As before, expanding Eq. (88) up to the second order:

$$\phi(t) \propto \left(1 + 2\frac{v_0 t}{z_0 + r} + 3\left(\frac{v_0 t}{z_0 + r}\right)^2 + \cdots\right), \tag{92}$$

and comparing this expansion with the expansion for $E_c(t)$, we again find that isentropic compression is possible up to first order with

$$15/8t_0 = 2v_0/(z_0 + r). \tag{93}$$

It remains for us to check the consistency of the obtained result in regard to the fuel burnup. The total number of neutrons coming out of the fuel cylinder of surface area $S = 6\pi r^2$ during the inertial confinement time $\tau = 2r/v_0$ is given by

$$N = \phi(z)S\tau = 16\pi r^4 n_0^2\langle\sigma v\rangle/v_0. \tag{94}$$

Obviously, this number cannot be larger than the number of nuclear reactions, which is equal to $(1/2)nV$, whereby each pair of nuclei produces one neutron. Since the precompressed fuel receives an additional fourfold compression in passing through the detonation front, the maximum possible number of neutrons that can be produced is given by

$$N_0 = \tfrac{1}{2}4n_0 2\pi r^3 = 4\pi r^3 n_0. \tag{95}$$

From the condition $N \leq N_0$ it follows that

$$4rn_0\langle\sigma v\rangle/v_0 \leq 1. \tag{96}$$

If this inequality is not satisfied the burnup is incomplete. With the parameters

$$r = 10 \text{ cm}, \qquad n_0 = 10n_s = 5 \times 10^{23} \text{ cm}^{-3},$$

$$\langle\sigma v\rangle = 10^{-15} \text{ cm}^3/\text{sec}, \qquad v_0 = 10^8 \text{ cm/sec},$$

we find that

$$4rn_0\langle\sigma v\rangle/v_0 = 200.$$

However, this number could easily be two times smaller since v_0 may be larger and $\langle\sigma v\rangle$ smaller than assumed. But the fact that the number is so large indicates a large fuel burnup.

We would like to comment on the assumption that the energy released in the neutron-absorbing substance is in thermodynamic equilibrium with the blackbody radiation. Whether this is true will, of course, depend upon the optical thickness of the neutron absorber relative to the radial dimension of the thermonuclear assembly, which is of the order $r \gtrsim 10$ cm in our case. Only if the optical thickness is small compared to this dimension is the blackbody assumption justified.

For incompletely ionized matter the optical cross section is $\sim 10^{-18}$ cm^2, and hence the optical thickness at solid densities is of the order $\sim 10^{-4}$ cm. In this case our assumption would be fully justified. For completely ionized matter the optical thickness can become larger by many orders of magnitude. The temperature for complete ionization is given by

$$T_i \simeq 2.1 \times 10^5 Z^{2.42}.$$

For uranium, with $Z = 92$, we find $T_i = 1.2 \times 10^{10}$ K.

In the case of complete ionization and optically thin plasmas, all the energy goes into kinetic particle energy and the maximum kinetic temperature is computed from

$$q_{\max} = (3/2)(Z + 1)nkT_k.$$

Then, only if $T_k > T_i$ is the matter completely ionized. For uranium and the above given value of

$$q_{\max} = 6 \times 10^{17} \text{ erg/cm}^3,$$

one finds $T_k = 6.2 \times 10^8$ K, which is well below the temperature for complete ionization. Therefore, in the case of uranium as a neutron absorber the plasma is opaque.

In the case of a boron plasma the situation is different. Here the total ionization temperature is $T_i = 10^7$ K, but the maximum kinetic particle temperature, with

$$Z = 5, \quad n = 1.4 \times 10^{23}\ \text{cm}^{-3},$$

and

$$q_{\max} = 6.5 \times 10^{16}\ \text{erg/cm}^3,$$

is $T_k = 3.7 \times 10^9$ K, which is larger than T_i. In this case, then, we have a fully ionized boron plasma, and te optical thickness is here entirely determined by bremsstrahlung. The bremsstrahlung losses of a fully ionized plasma are given by

$$P = bZ^3n^2T^{1/2}\ \text{erg/cm}^3\text{-sec},$$

with

$$b = 1.42 \times 10^{-27}\ \text{erg-cm}^3/\text{sec-K}^{1/2},$$

and the optical thickness is given by

$$\lambda = acT^4/P = 1.62 \times 10^{23}\ T^{3.5}/Z^3n^2\ \text{cm}.$$

Using $T = T_k = 3.7 \times 10^9$ K this results in $\lambda \simeq 2 \times 10^8$ cm, and it follows that the boron plasma is optically transparent. The kinetic plasma pressure is now twice as large as the radiation pressure because here

$$p = (2/3)q_{\max} \simeq 4 \times 10^{16}\ \text{dyn/cm}^2.$$

Since one may wish to surround the boron with some heavier high-Z material to improve inertial confinement, however, this high-Z material is expected to mix with the low-Z boron, which can reduce the optical thickness by many orders of magnitude.

In the smallest conceivable design for an ordinary H-bomb or a neutron bomb, the thermonuclear fuel is simply precompressed by the soft X-rays from an exploding fission bomb without making use

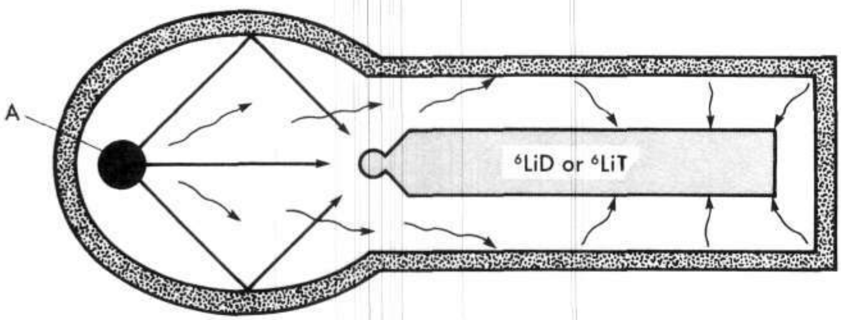

Figure 18. Compact H-bomb or neutron bomb in which the thermonuclear fuel is simply precompressed by the soft X-rays without using the autocatalytic principle.

of the autocatalytic principle. This concept is shown in Figure 18. This design is meaningful only for rather small fuel cylinders.

CHAPTER FOURTEEN

The Question of Nonfission Ignition

Up to the present time a thermonuclear detonation has only been successfully ignited by a fission bomb. There are, however, at least in principle, several alternative possibilities to ignite a thermonuclear detonation of arbitrary magnitude by other energy sources. One class of these alternative ignition concepts makes use of intense beams of photons (lasers), electrons, ions, and even solid projectiles. All these concepts are currently under study in the context of the quest for controlled thermonuclear energy release by inertial confinement fusion.

If any one of these methods is successful, it could be used to ignite a thermonuclear explosion of arbitrary size by letting the thermonuclear reaction spread to adjacent thermonuclear material through the process of autocatalytic thermonuclear detonation. However, all published experiments so far conducted indicate that the apparatus to produce beams with the required intensity and energy is much too large to be incorporated into a portable weapons device.

A much smaller beam generator could be built, at least in principle, that draws its energy from a large chemical explosion. The theoretical estimates presented below suggest that an energy of $\sim 10^7$ J would be sufficient to ignite a thermonuclear microexplosion, which then could spread and grow into a large thermonuclear explosion of arbitrary size. A trigger energy of $\sim 10^7$ J could be easily stored in several kilograms of high explosives. Therefore, if an efficient way can be found to transform a substantial part of this energy into a short laser pulse by some sort of chemical laser, the laser-ignited thermonuclear weapon would become a real possibility. Similarly, methods by which this energy could be transformed into a superpowerful electric spark, consisting of an intense stream of charged particles, may lead to a portable electric bomb trigger. For this second possibility, storage of electric energy in a capacitor using ferroelectric substances might become important.

The energy stored in high explosives can also be converted into kinetic projectile energy, which in turn could be used to ignite a thermonuclear reaction. However, since for this purpose projectile velocities up to several hundred kilometers per second are required, the conversion of stored chemical energy into kinetic projectile energy is not a trivial problem. (See Chapter 17 for a discussion of a more recent, low-velocity concept using a magnetized plasma target as an intermediate ignition stage.) One method by which this could be done is an explosively driven magnetic rail gun. Another more direct method makes use of the velocity-amplification effect described above, in which a sequence of concentric shells is imploded on each other. There the rise in temperature is given by Eq. (40).

Consider the case shown in Figure 19, where the explosive is arranged on a spherical shell surrounded by a heavy tamp. If we assume that the initial radius is ~ 1 m, and the initial temperature of the explosive is $\sim 10^4$ K, it would follow that in going from $r_0 = 10^2$ cm down to $r_1 \simeq 2.5$ cm the temperature rises to $\sim 10^8$ K. This corresponds to an implosion velocity of several hundred kilometers per second for the innermost shell. We shall now present a further sophistication of this idea, however, which promises a substantial reduction of the required impact velocity, provided that the impact mass is that much larger to give the same trigger energy.

For any ignition source using particle beams, the energy density ε and energy flux density ϕ_ε are related to each other by

$$\phi_\varepsilon = v\varepsilon, \tag{97}$$

where v is the particle velocity. Thermonuclear ignition requires,

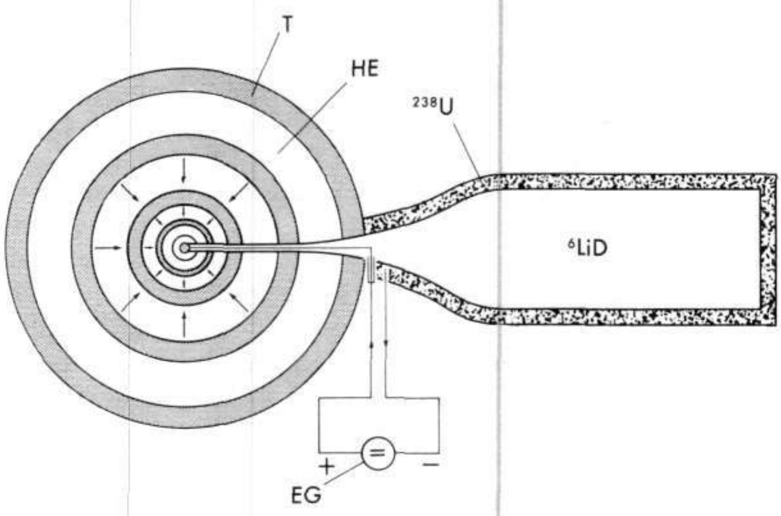

Figure 19. Nonfission ignition initiated by chemical high explosive HE: use is made of multishell velocity-amplifying implosion onto a thermonuclear fuse ^{6}LiD that is magnetized by a high-current pinch discharge.

first,

$$\phi_\varepsilon \geq 10^{14}\ \mathrm{W/cm^2}$$

delivered to an area of approximately 1 $\mathrm{cm^2}$ and lasting for about 10^{-8} sec, and second, a total ignition energy of $E_0 \gtrsim 10^6$ J. From Eq. (97) it follows that low velocities must be compensated by large values of ε to reach the same value of ϕ_ε.

For all photon and relativistic particle beams, of course, $v = c$. In this limiting case ε can thus assume the smallest possible value compatible with the required value of ϕ_ε. This fact suggests that we start the ignition process with a projectile having a comparatively low velocity but large ε value. Here, even though the energy flux density ϕ_ε can be comparatively small, the energy density ε may still be quite large. If this kinetic projectile energy is then converted and stored as kinetic energy of many much lighter particles moving with the velocity v_p, the energy flux density can be amplified by the factor v_p/v. For photons with $v_p = c$, the power amplification would reach its absolute maximum c/v. This fact suggests converting the initial projectile energy into photons, which would have to be stored and confined inside a small cavity possessing a highly opaque wall before being released onto the thermonuclear target.

The conversion of the initial projectile energy into photons could be done by letting the projectile make an impact on a low-density gas of high atomic weight placed inside the cavity. Upon impact with the hypervelocity projectile, a strong shock wave is produced, which becomes an intense source of photons. Because of the isotropic nature of the radiation thus emitted the cavity will be filled with blackbody radiation. The isotropy of the radiation leads to a 1/4 smaller energy flux, given by the Stefan-Boltzmann law

$$\phi_\varepsilon = (c/4)\varepsilon, \tag{98}$$

rather than $\phi_\varepsilon = c\varepsilon$, as is the case for a laser beam. But even then, the large value of $c/4$ as compared to v for the incoming hypervelocity projectile still leads to a very large power amplification factor, given by $c/4v$.

Figure 20 shows how this idea can be realized practically. A hypervelocity projectile is impacting on a conical cavity filled with a low-density, high-atomic-weight gas, producing in it a strong shock wave that emits many photons. The incoming projectile further compresses the photon gas thus formed, thereby raising its temperature. This rise in temperature reduces the opacity and increases the power absorbed by other materials—including the thermonuclear fuel, which is ignited at the vertex of the cone. The rapid decrease in opacity from the rising temperature of the photon gas furthermore ensures that most of the photon energy is delivered toward the end of the compression process, where the power peaks strongly.

If a power of 10^{14} W/cm^2 = 10^{21} erg/cm^2-sec is to be reached, the required projectile impact velocity is computed by putting

$$(c/4v)(1/2)\rho v^3 = (c/8)\rho v^2 = 10^{21}\ \text{erg/cm}^2\text{-sec}.$$

Assuming $\rho \simeq 10$ g/cm^3, one finds $v \simeq 2.5$ km/sec, which is an astonishingly small velocity.

The actual velocity, however, must be larger, principally for two reasons: (1) Since it is necessary to establish a sufficiently dense

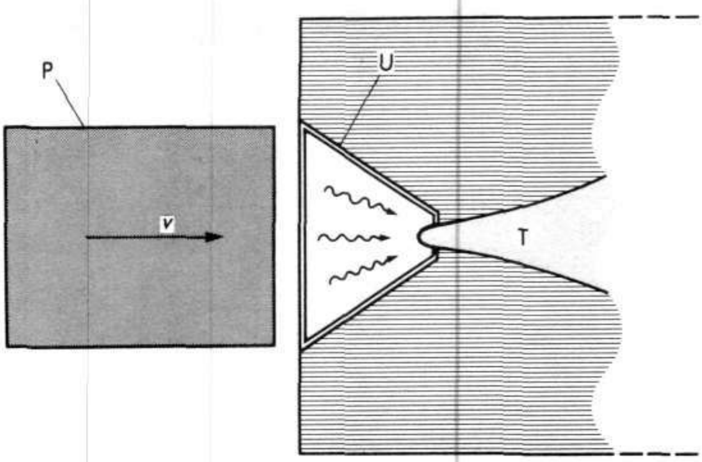

Figure 20. Nonfission ignition using hypervelocity impact of a projectile P traveling with velocity v and colliding with a conical cavity U filled with a low-density, high-atomic-weight gas. The cavity fills with blackbody radiation from the impact and ignites the thermonuclear fuel T at the vertex of the cone.

photon gas, which can only be done at a temperature in excess of several million kelvins, a much larger impact velocity is required. (2) Since the cavity wall is not completely opaque against diffusion losses by photons, these losses must be reduced by a sufficiently large implosion velocity. We shall now give a more quantitive account of these ideas.

Upon hypervelocity impact the high-atomic-number gas entrapped inside the cavity is initially shock-heated. For an assumed specific heat ratio $\gamma = 5/3$, the shock velocity U is related to the impact velocity v by $U = (4/3)v$.[7] Therefore, substituting the value of $U = (4/3)v$ into the well-known relations for strong shocks,[7] the temperature of the shock-heated gas, which is transformed into a Z-times-ionized plasma, is

$$T = [A/3R(Z + 1)]v^2. \tag{99}$$

Here A is the atomic weight and Z is the degree of ionization. Furthermore, $R = 8.31 \times 10^7$ erg/g-K is the universal gas constant.

The degree of ionization Z depends itself upon T. A useful approximate relation between Z and T is given by

$$Z \simeq 6.6 \times 10^{-3} T^{0.41}. \tag{100}$$

For $Z \gg 1$, Eqs. (99) and (100) can be combined to give T as a function of v and A:

$$T \simeq 4 \times 10^{-5} A^{0.71} v^{1.42}, \tag{101}$$

or

$$v \simeq 1.3 \times 10^{3} T^{0.7} / A^{0.5}. \tag{102}$$

The minimum required temperature needed to reach

$$\phi_{\varepsilon} = 10^{14}\ \mathrm{W/cm^2} = 10^{21}\ \mathrm{erg/cm^2\text{-}sec}$$

can now be computed from the Stefan-Boltzmann law, putting

$$\phi_{\varepsilon} = \sigma T^4,$$

where

$$\sigma = 5.75 \times 10^{-5}\ \mathrm{erg/cm^2\text{-}sec\text{-}K^4}.$$

We find $T = 2 \times 10^6$ K. The impact velocity to reach this temperature is computed by Eq. (102), and, assuming $A = 200$, $v \simeq 23$ km/sec. The size of the required cavity volume V can again be computed from the Stefan-Boltzmann law, but here setting

$$E_0 = aT^4 V,$$

where $a \equiv 4\sigma/c$ and E_0 is the ignition energy. For $E_0 = 10^{14}$ erg and $T = 2 \times 10^6$ K, it follows that $V \simeq 10^3$ cm.

This rather large volume is very inconvenient, especially if one chooses to use the multishell configuration of Figure 10. It is therefore desirable to go to higher temperatures, even though this requires higher velocities. A reasonable compromise appears to be a temperature of $T \simeq 5 \times 10^6$ K, for which $V \simeq 25$ cm^3. Here $v \simeq 46$ km/sec, which is still substantially less than the required 200 km/sec for ignition by direct projectile impact.

During the implosion process the temperature of the photon gas will rise by adiabatic compression in proportion to $V^{-1/3}$. Thus, an $\sim$10-fold reduction in the cavity volume will result in an $\sim$2-fold rise in the temperature from 5×10^6 K up to $\sim 10^7$ K. At this temperature $\phi_c \simeq 6 \times 10^{16}$ W/cm^2, which is reached at the end of the implosion process. This large energy flux density means that the target from which the thermonuclear deflagration spreads can be made much smaller, for example, with a radius of only $\sim$0.1 cm.

If we make the density of the high-atomic-weight gas sufficiently low, its pressure can be neglected against the radiation pressure. The gas density must still be chosen to be sufficiently high for shock waves to occur, ensuring that most of the energy goes into the radiation field. This requires that the mean free path be small compared to the cavity dimension. For a high-atomic-weight gas at $T \sim 10^7$ K, the collision cross section is approximately given by $\sigma_{\text{coll}} \sim 10^{-16}$ cm^2. If the collision mean path shall not exceed $\sim 1/100$ of the cavity dimension, then $n \geq 10^{18}$ cm^{-3}. In this case one also has $\tau_c < \tau$, where

$$\tau_c = \lambda_{\text{coll}}/v_{\text{thermal}}$$

is the collision time and $\tau \simeq r/v$ the cavity implosion time. If $\tau_c < \tau$, the kinetic particle energy is rapidly converted into photon energy by inelastic collisions, with a decay time of the excited atoms less than $\sim 10^{-8}$ sec. This is much shorter than either τ or τ_c.

The pressure of the Z-times ionized gas is given by

$$p_g = (Z + 1)nkT.$$

Making this pressure small compared to the radiation pressure $p_r = (a/3)T^4$ leads to the condition

$$n < (d/3k)[T^3/(1 + 6.6 \times 10^{-3}T^{0.41})]. \tag{103}$$

where we have expressed Z as a function of T by Eq. (100). For $T \simeq 10^7$ K, we find that $n \leq 10^{21}$ cm^{-3}.

Another limitation on the gas density is determined by the requirement that the opacity be sufficiently low to make it optically transparent. The opacity of high-atomic-weight material is given by[11]

$$\kappa = 4.34 \times 10^{25} \rho T^{-3.5} (g/t). \tag{104}$$

In Eq. (104), g and t are the Gaunt and the guillotine factors, respectively. If the density ρ is not too high, we can put $g \simeq 1$. For the chemical composition of stellar atmospheres (Russel mixture), $t \simeq 10$. In the case of uranium or other high-atomic-weight elements, however, the level density is so large that it is very likely justified to use its minimum value $t = 1$ for the guillotine factor. Under this assumption, together with the definition of the optical path length

$$\lambda_{\text{opt}} = (\rho\kappa)^{-1},$$

we find for $T \simeq 10^7$ K and for a gas density of $\rho \lesssim 0.3$ g/cm^3 that

$$\lambda_{\text{opt}} > r \sim 1 \text{ cm},$$

where r is the linear cavity dimension. For $A \simeq 200$ this implies that $n \lesssim 10^{21}$ cm^{-3}. To make $\lambda_{\text{opt}} \sim 10r$ would thus require choosing $n \sim 10^{20}$ cm^{-3}, which would make the mean free path $\sim 10^{-4}$ cm. This is still sufficiently short to produce shock waves within the cavity.

To obtain a value for the diffusion velocity by which the energy is lost through the cavity wall we must solve the radiative heat-transfer equation. Since, according to Eq. (104), the absorption coefficient $\kappa\rho$ increases with ρ^2, where ρ is the density of the wall material, we should use uranium, gold, or some other high-density material for optimal photon confinement. Because the photons are expected to diffuse just a short distance into the wall during the implosion time, it is sufficient to coat the cavity wall with these dense substances. Electronic heat conduction losses are small by comparison and can be neglected in the heat-transfer equation.

To set up the heat-transfer equation the ratio of the kinetic to the radiative energy density within the wall material must be known. We know that this ratio is small inside the cavity because the density of the high-atomic-weight gas is rather low, even though this is not expected to be true in the wall material itself. There this ratio is given by

$$f = \frac{(3/2)(Z + 1)nkT}{aT^4} = \frac{3(Z + 1)nk}{2aT^3}. \tag{105}$$

Assuming $T = 10^7$ K, resulting in $Z \simeq 5$, according to Eq. (100), we find for uranium with $n = 5 \times 10^{22}\ \text{cm}^{-3}$ that $f \simeq 0.1$. Because f is so small the heat-conduction equation can be approximated fairly well by

$$\rho c_v \frac{\partial T}{\partial t} = \frac{\partial}{\partial x} \frac{\lambda_p c}{3} \frac{\partial}{\partial x}(aT^4). \tag{106}$$

The coordinate x is directed perpendicularly into the wall. c_v is the specific heat at constant volume given by

$$c_v = (3/2)(Z + 1)(R/A) = 3.74 \times 10^6\ \text{erg/g-K},$$

and λ_p is the photon path length, which, according to Eq. (104), is given by

$$\lambda_p = bT^{3.5}/\rho^2, \tag{107}$$

with $b = 2.3 \times 10^{-26}$ in cgs units. We thus have

$$\frac{\partial T}{\partial t} = \frac{abc}{3\rho^3 c_v} \frac{\partial}{\partial x} T^{3.5} \frac{\partial}{\partial x} T^4. \tag{108}$$

To obtain the velocity v_d of the heat-conduction wave we substitute into Eq. (108)

$$\partial T/\partial t = v_d \partial T/\partial x. \tag{109}$$

Then, after integrating Eq. (108) over x, we find

$$v_d = (abc/3\rho^3 c_v)T^{2.5}(\partial T^4/\partial x), \tag{110}$$

for which we can also write

$$v_d = \frac{4}{6.5}\frac{abc}{3\rho^3 c_v}\frac{\partial T^{6.5}}{\partial x} \simeq 0.2\frac{abc}{\rho^3 c_v}\frac{\partial T^{6.5}}{\partial x}. \tag{111}$$

To make a rough estimate for v_d we can approximate

$$\partial T^{6.5}/\partial x \sim T^{6.5}/x,$$

and thus have

$$v_d \simeq 0.2(abc/\rho^3 c_v)(T^{6.5}/x). \tag{112}$$

The distance x by which the heat-conduction wave has advanced is obtained from Eq. (112) with $v_d = dx/dt$, resulting in the differential equation.

$$\frac{dx}{dt} \simeq 0.2\frac{abc}{\rho^3 c_v}\frac{T^{6.5}}{x}, \tag{113}$$

which we can rewrite as follows:

$$dx/dr = -B/2vx, \tag{114}$$

where

$$B \equiv 0.4\frac{abc}{\rho^3 c_v}T^{6.5} \simeq \frac{5.7\times 10^{-37}T^{6.5}}{\rho^3}. \tag{115}$$

In Eq. (114) we have put

$$dx/dt = (dx/dr)(dr/dt) = -v\,dx/dr,$$

where r is the cavity radius and v the implosion velocity. In the limiting case of no losses through the wall, the blackbody radiation follows the isentropic change of state given by VT^3 = const, or

$$rT = r_0 T_0, \tag{116}$$

where r_0 and T_0 define the initial cavity radius and temperature, respectively, at the beginning of the implosion. The cavity radius being a function of time, it can be equated to the height of the cone during the implosion process, itself a function of time, because a cone changes its volume depending on this quantity in the same proportion as the volume of a sphere depending on its radius. With the help of Eq. (116), we can write for Eq. (115):

$$B = B_0 (r_0/r)^{6.5}, \tag{117}$$

where

$$B_0 \equiv 5.7 \times 10^{-37} T_0^{6.5}/\rho^3. \tag{118}$$

By inserting the value of B given by Eq. (117) into Eq. (114), the variation of B in the approximation that the losses through the wall are small is taken into account. We thus have

$$\frac{dx}{dr} \simeq -\frac{B_0}{2vx}\left(\frac{r_0}{r}\right)^{6.5}. \tag{119}$$

Integrating Eq. (119) from $r = r_0$, $x = 0$ to $r = r_1$, x, we find

$$x^2 = \frac{B_0 r_0}{5.5v}\left[\left(\frac{r_0}{r_1}\right)^{5.5} - 1\right]. \tag{120}$$

We shall take as an example the situation where the final energy accumulated inside the cavity is $\sim 10^{14}$ erg, with the cavity temperature assumed to be $\simeq 10^7$ K. This condition will be reached at the final cavity radius $r = r_1$. The initial cavity radius is $r = r_0$ and the initial cavity temperature is $T_0 \simeq 5 \times 10^6$ K, the latter requiring an impact velocity of ~46 km/sec. We can easily compute that this example requires $r_0 \sim 1$ cm and $r_1 \sim 0.5$ cm. We then obtain from

Eq. (120) $x \simeq 0.07$ cm. It thus follows that at a final radius of $r_1 \simeq 0.5$ cm only the fraction $f \simeq 3x/r_1 \simeq 40\%$ of the accumulated energy is lost through the wall. This implies an overall efficiency of about 60%.

The amount of radiation incident upon the thermonuclear target assumed to be spherical in shape with a radius r_p and positioned at the vertex point of the conical cavity is given by the Stefan-Boltzmann law

$$P = 4\pi r_p^2 \sigma T^4. \tag{121}$$

To obtain the variation of P with time Eq. (116) requires $T = T_0 (r_0/r)$ and, furthermore, $r = r_0(1 - t/t_0)$, where $t_0 = r_0/v$ is the cavity implosion time. Then

$$P(t) = 4\pi r_p^2 \sigma T_0^4 (1 - t/t_0)^{-4}. \tag{122}$$

We take the example $T_0 = 5 \times 10^6$ K, requiring $v \simeq 46$ km/sec, and a target radius $r_p \simeq 0.1$ cm. This yields

$$P(0) = 4.35 \times 10^{21} \text{ erg/sec} = 435 \text{ TW}.$$

If the cavity collapses from a radius of $r_0 \sim 1$ cm to a radius of $r_1 \sim 0.5$ cm, the temperature rises in the time

$$(r_0 - r_1)/v \simeq 10^{-7} \text{ sec}$$

from $T \sim 5 \times 10^6$ K to $T \sim 10^7$ K, and the power delivered to the target rises from 435 TW to $\simeq$7000 TW. This is a respectable power amplification.

Obviously, not all the incident energy is immediately absorbed in the target, but the photons randomly reflected back from the target surface into the cavity are trapped within the imploding cavity and thus can hit the target many times over during the entire implosion process. The fraction of energy absorbed by the target, of course, is removed from the energy field of the blackbody radiation. This absorption process changes the dependence of T upon r, giving a smaller-than-isentropic rise of T as r decreases. In the optimal case, all the radiative energy would be delivered to the target. If all the kinetic projectile energy of 10^{14} erg is eventually absorbed by the target, the average power of absorbed energy would be only 100 TW. The actual power delivered to the target peaks at the end of the implosion process, however, as mentioned before. This peaking in the delivered power is principally determined by two effects: (1) A rise in T increases the optical path length in proportion to $T^{3.5}$, and thus in proportion to the fraction of energy absorbed. (2) The incident power rises as T^4. Both effects combined will therefore result in a power rising in proportion to $T^{7.5}$. Since T rises as r^{-1}, it therefore follows that most of the radiative energy accumulated during the implosion of the cavity will be delivered toward the end of the implosion process with a much higher power than the average power.

Returning to the onionlike configuration of imploded shells described in Figure 19, it is obvious that this same configuration can also be used to generate and compress blackbody radiation. The only change required is to fill the gaps between the shells with tenuous, high-atomic-weight gas of the right density. The much smaller initial velocities needed here make the concept thus modified much more likely to work as a means for thermonuclear ignition. A velocity of only 50 km/sec could be reached with a small number of shells, which, however, would have to be more massive to attain the same amount of kinetic impact energy.

After ignition has been achieved in the center of the shell structure, the problem remains of guiding a thermonuclear burn wave from the center to the thermonuclear material placed outside. We shall describe one method by which this could be achieved. This method involves a strong magnetic field created by a large current, produced by an auxiliary setup in the form of an explosive generator. The current would first have to pass through a cable into the center of the onionlike structure and from there would break down into a radial pinch channel burning all the way through to the outer surface of the entire assembly. The channel itself could be produced by an exploding lithium wire. To ensure a fast pinch discharge the inductance of the channel should be as small as possible. This dictates using a coaxial cable, wherein the inner conductor is the exploding lithium wire forming the pinch and the return current passes through the outer cylinder.

The pinch current I produces an azimuthal magnetic field H_ϕ near its surface given by

$$H_\phi = 2I/rc. \tag{123}$$

If strong enough, this azimuthal magnetic field can radially confine the charged fusion products, and hence could guide a thermonuclear burn wave propagating along the pinch channel. To derive the necessary conditions for this to happen, we compute the Larmor radius of the charged fusion products. If their mass is m, their velocity v and charge Ze, their Larmor radius is

$$r_L = mvc/ZeH_\phi = rvmc^2/2ZeI. \tag{124}$$

In order to confine the fusion products within the reacting volume requires that $r_L < r$, and thus

$$I > I_0 = mc^2v/2Ze. \tag{125}$$

Introducing the energy E of the charged fusion products from $E = (1/2)mv^2$, we find that

$$I_0 = (c^2/Ze)\sqrt{mE/2} \text{ esu}. \tag{126}$$

It is a remarkable fact that this critical confinement current depends only on the energy of the fusion products and not on the target size. We assume that the thermonuclear fuel in the center of the onionlike structure and along the pinch channel is DT, because it is the easiest fuel to ignite. The fusion α-particles have here an energy $E = 3.5$ MeV and one obtains

$$I_0 = 1.3 \times 10^6 \text{ A}.$$

Finally, we shall show that the assumed trigger energy of $E_0 \simeq 10^{14}$ erg is a reasonable estimate. For this we recognize the fact that the plasma volume to be heated up to the ignition temperature is given by

$$V_0 \simeq 2\lambda_0 S_0 = 2a(kT)^{3/2}S_0/n_0, \tag{127}$$

where S_0 is the initial cross section of the pinch channel positioned in the center of the spherical assembly. According to Eq. (46) the trigger energy, which is here independent of n_0, is thus given by

$$E_0 \simeq 12aS_0(kT)^{5/2}. \tag{128}$$

Since it appears technically feasible to focus an electric high-current discharge down to a cross section of $S_0 \sim 10^{-4}$ cm, we find for T = 70 keV that

$$E_0 \simeq 1.2 \times 10^{14}\,\text{erg} = 1.2 \times 10^7\,\text{J}.$$

It may even be possible to ignite a thermonuclear explosion by the configuration shown in Figure 21. Here the detonation waves from a high explosive (HE) are reflected from the wall of the Prandtl-Meyer ellipsoid and are thereby refocused onto the multi-shell configuration S_1, S_2, ... Between these shells there is a tenuous, high-atomic-weight buffer gas to generate blackbody radiation. The shells must consist of dense material to be optically

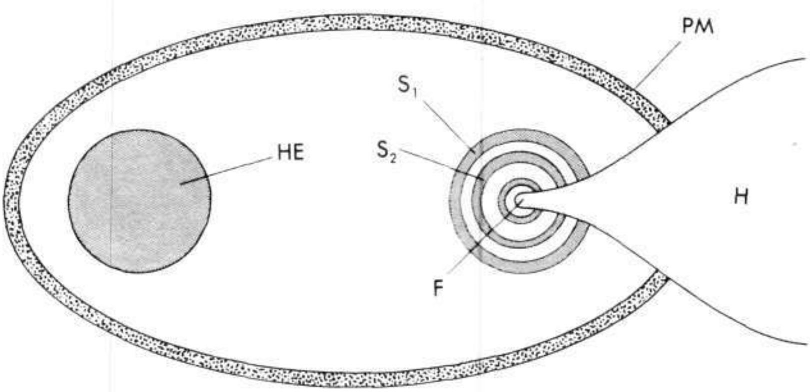

Figure 21. Chemical ignition of a thermonuclear explosion using high explosive HE in one focus of the Prandtl-Meyer ellipsoid PM. The thermonuclear fuse F positioned in the other focus is surrounded by several shells S_1, S_2, . . . , with tenuous gas between the shells to generate intense blackbody radiation. After the fuse is lit, a thermonuclear deflagration will advance into the horn H.

opaque. The blackbody radiation set up by the implosion of one shell onto the tenuous gas between this shell and the next inner shell then ablates and implodes this inner shell onto the tenuous gas between this imploding shell and the next following inner shell. This sequence continues until the innermost shell implodes onto the gas surrounding the thermonuclear fuse. If the implosion velocity of the *innermost shell is larger than* ~50 km/sec this would create enough blackbody radiation to make thermonuclear ignition possible.

The rather low impact velocity of ~50 km/sec required should also make it possible to reach such a velocity by shape charges. In this case the shape charge could propel a projectile to this velocity and, just as shown in Figure 20, impact onto a conical cavity. Since with shape charges velocities of ~90 km/sec have been reached, this alternative possibility does not seem to be totally unreasonable.

The general considerations and very simple estimates presented in this chapter demonstrate that the possibility of a nonfission-triggered fusion bomb cannot be ruled out. In support of this statement, we refer also to the concept discussed in Chapter 17.

CHAPTER FIFTEEN

Thermonuclear Microexplosions

An introduction to the physical principles of thermonuclear explosive devices would be incomplete did it not give at least a cursory overview of the different approaches to igniting thermonuclear microexplosions. A thermonuclear microexplosion is an explosive release of thermonuclear energy many orders of magnitude smaller than from a thermonuclear weapon. To achieve this goal, the fission trigger must be replaced by some other means of producing the required ignition temperature, but in a much smaller volume and without the large energy release.

The possibility of thermonuclear microexplosions is itself a consequence of two facts: first, the fact that the minimum volume to make a thermonuclear explosion is given by the range of the charged fusion products λ_0, and, second, scaling. At solid densities and thermonuclear temperatures the range is of the order of a few centimeters and, according to Eq. (51), is inversely proportional to the density ρ of the thermonuclear explosive. We can therefore write for this range

$$\lambda_0 = a/\rho, \qquad a = \text{const.} \tag{129}$$

On the other hand, the ignition energy E_{ign} is proportional to the mass of the explosive, and we thus can write

$$E_{\text{ign}} = b\rho\lambda_0^3, \qquad b = \text{const.} \tag{130}$$

Combining Eqs. (129) and (130) shows that

$$E_{\text{ign}} = a^3 b/\rho^2, \tag{131}$$

for which we can also write

$$E_{\text{ign}}/E_{\text{ign}}^{(0)} = (\rho_s/\rho)^2, \tag{132}$$

where $E_{\text{ign}}^{(0)}$ is the minimum ignition energy for solid target density $\rho = \rho_s$. For the DT reaction we easily find that the minimum ignition energy is about 10^8 J. To obtain a useful gain the output energy should be at least ~ 100 times larger, that is, 10^{10} J = 10^{17} erg, which corresponds to the explosive power of approximately 2 tons of TNT. Equation (132) shows that the ignition energy can be substantially reduced by increasing the density of the thermonuclear explosive. However, since for $\rho > \rho_s$ the compression to higher densities also requires energy, the input energy is in reality larger than the value given by Eq. (132).

The time τ to deposit the energy in the target is given by $\tau \simeq r/v$, where v is the thermal expansion velocity of the thermonuclear plasma of radius r, which is a function only of the temperature. Putting $r \simeq \lambda_0$, we find that

$$\tau/\tau^{(0)} = \rho_s/\rho, \tag{133}$$

where $\tau^{(0)}$ is the time the energy has to be delivered for $\rho = \rho_s$.

The power P of the ignition pulse to be delivered to the thermonuclear explosive is given by $P = E_{\text{ign}}/\tau$, and hence

$$P/P^{(0)} = \rho_s/\rho, \tag{134}$$

where $P^{(0)}$ is the power for $\rho = \rho_s$.

Finally, the power density W of the ignition pulse to be delivered to the thermonuclear target is proportional to

$$P/r^2 \propto P/\lambda_0^2,$$

and we find

$$W/W^{(0)} = \rho/\rho_s, \tag{135}$$

where $W^{(0)}$ is the power density for $\rho = \rho_s$.

The four expressions (132) through (135) are the basic scaling laws for thermonuclear microexplosions. These scaling laws show that both the trigger energy and the power of the ignition pulse can be greatly reduced if the thermonuclear explosive is precompressed. For $\rho > \rho_s$ only the specific power given by Eq. (135) must be larger, which is a direct result of the fact that the time duration for the ignition pulse decreases as ρ increases.

These scaling laws do not take into account the fact that additional energy is required to compress the target. The amount of this energy can be made quite small, however, if the compression is done isentropically, keeping the temperature during this precompression as low as possible. To make an isentropic compression, the energy must be supplied in a certain time-dependent manner.[10]

Ideally, the pressure in the cold thermonuclear explosive is that of a degenerate Fermi gas, with a specific heat ratio $\gamma = 5/3$. If the implosion is spherically symmetric, the temperature rises as

$$T \propto n^{\gamma-1} \propto r^{-3(\gamma-1)}.$$

For an isentropic compression the implosion velocity must be equal to the thermal velocity

$$v_{\text{imp}} = v_{\text{th}} \propto T^{1/2} \propto r^{-3(\gamma-1)/2}.$$

Putting $v_{\text{imp}} = -dr/dt$, we obtain the differential equation

$$dr/dt = -\text{const}\, r^{-3(\gamma-1)/2} \tag{136}$$

Integration of Eq. (136) yields

$$r = r_0(1 - t/t_0)^m, \tag{137}$$

where $m \equiv 2/(3\gamma - 1)$ and $r_0 = r(0)$. The value $r = 0$ is reached for $t = t_0$, $r = 0$. From Eq. (137) we find

$$v_{\text{imp}} = -dr/dt \propto (1 - t/t_0)^{-q}, \tag{138}$$

where $q \equiv 3(\gamma - 1)/(3\gamma - 1)$.

The work done to compress the thermonuclear material during the implosion has the rate

$$dA/dt = 4\pi r^2 pv, \tag{139}$$

where the pressure $p \propto nv^2 \propto v^2/r^3$, hence $dA/dt \propto v^3/r$. With Eqs. (137) and (138) we thus find

$$dA/dt \propto (1 - t/t_0)^{-s}, \tag{140}$$

where $s \equiv (9\gamma - 7)/(3\gamma - 1)$. For $\gamma = 5/3$, $s = 2$. This is approximately the same as the more correct value $s = 15/8$ used in Eq. (81), which is obtained from a more exact theoretical treatment.[10] Since dA/dt must be set equal to the power P of the ignition pulse,

$$P(t) = P_0(1 - t/t_0)^{-2}, \tag{141}$$

where P_0 is the power at $t = 0$. Equation (141) shows that $P \to \infty$ for $t \to t_0$. This is not surprising since at $t = t_0$ the density of the compressed material will go to infinity. Therefore, to reach only a finite density $\rho > \rho_s$ the compression pulse can be cut off for $t < t_0$.

The particular power profile given by Eq. (141) for isentropic compression minimizes the work required to achieve a desired density. Using such a power profile, we find that the compression energy is roughly 1/10 the energy required to heat the thermonuclear material to the ignition temperature. Without the particular power profile given by Eq. (141) the energy for compression would be about 10 times larger and thus of the same order as the energy required to heat the entire explosive to the ignition energy. However, in a highly compressed explosive this ignition energy is reduced by the factor $(\rho_s/\rho)^2$.

Since the range of the charged fusion products at a density $\rho \sim \rho_s$ is of the order of a few centimeters, an $\sim$100-fold compression reduces this length down to a few $\sim 10^{-2}$ cm. The ignition pulse should thus ideally be focused down to such a small linear dimension. At thermonuclear temperatures the thermal velocity is $v_{th} \sim 10^8$ cm/sec and the ignition pulse must therefore be delivered in a time of the order $\sim 10^{-9}$ sec. For an ignition energy of $E_{ign} \sim 10^6$ J, the power of the ignition pulse would be $P \sim 10^{15}$ W. More detailed calculations show that a power roughly 10 times less should be sufficient. A further reduction in the power required for ignition is possible if the energy is first converted into blackbody radiation by imploding a small cavity. This concept was discussed in the previous chapter, but it could also be applied here, using the beam energy to implode a small cavity inside of which the thermonuclear explosive is placed.

The required large ignition energy of $\sim 10^6$ J together with the power and focusing requirements suggests that the ignition pulse should be produced by powerful beams, where the beams can be composed of particles ranging from zero rest mass (photons) up to projectiles (called macroparticles). In the one extreme case of photons, the beam particles move with the velocity of light; in the other extreme case of macroparticles, the velocity is in the range of 100 to 200 km/sec. In the table on the next page, the different beam options are summarized.

Of these different options only laser, relativistic-electron, and light ion beams are presently studied experimentally. Heavy ion beams produced in conventional particle accelerators are the most likely next candidate to join this group, with macroparticles presently under serious discussion. The table also lists microparticle beams, which are to be composed of submicron-size particles, typically with a linear dimension of $\sim 10^{-6}$ cm. No work on microparticle beams is presently planned, however, because the technology to produce such microparticle beams is quite difficult and little work has been done in the past. This situation may change, though, because the advan-

Comparison of Beam Options for Inertial Fusion

Driver	Major advantage	Major disadvantage	v/c
Laser beam	very good beam focusing	low efficiency, very expensive	1
Relativistic electron beam	can be cheaply generated	poor coupling to target	$\simeq 1$
Light ion beam	can be cheaply generated	difficult to focus	$\simeq 1/10$
Heavy ion beam	can be produced by conventional accelerator technology	expensive, accelerator several kilometers long	$\simeq 1/3$
Microparticle beam	very good beam-energy deposition	unexplored, unproven accelerator technology	$\simeq 3 \times 10^{-3}$
Superconducting projectile (macroparticle)	simple target compression and very good inertial confinement	accelerator several kilometers long	$\simeq 10^{-3}$

tages of microparticles over heavy ion beams are substantial. One reason for this is that they would have a smaller velocity and could be more easily concentrated. The comparatively low velocity of $\sim 10^8$ cm/sec also makes the energy-deposition problem rather easy in this case.

The last column of the table gives the ratio v/c of the particle velocity v to the velocity of light c. The available beam options therefore range from highly relativistic velocities down to velocities of a few hundred kilometers per second. The lowest velocities are thus not much larger than the velocities of meteors, which are less than 70 km/sec.

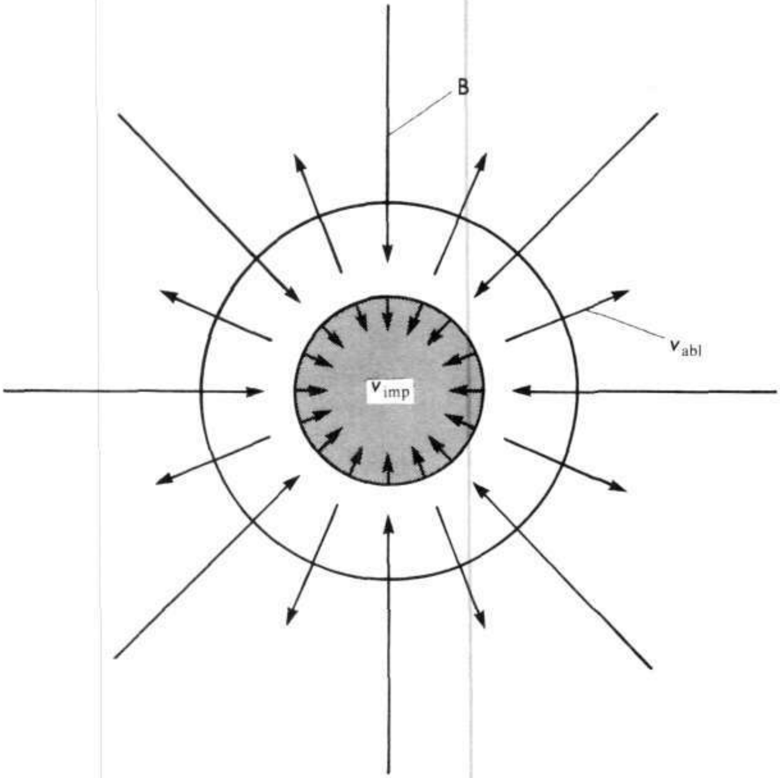

Figure 22. Ablation implosion of thermonuclear target bombarded by beams B (either laser or charged particles) from many sides. v_{imp} is the implosion velocity, and v_{abl} is the ablation product velocity.

Figure 22 shows the beam ignition principle with several beams used to bombard the thermonuclear explosive from several sides. The beams hitting the spherical thermonuclear target surface make it ablate. The recoil from this ablation causes a rocketlike implosion of the target.

The ignition concept using just one beam is explained[12] in Figure 23. Here the same principle of shock wave focusing by a curved wall that we encountered in the concept of igniting a thermonuclear explosion by a fission bomb is used. In the case of a laser beam the focusing can be done by a parabolic mirror.

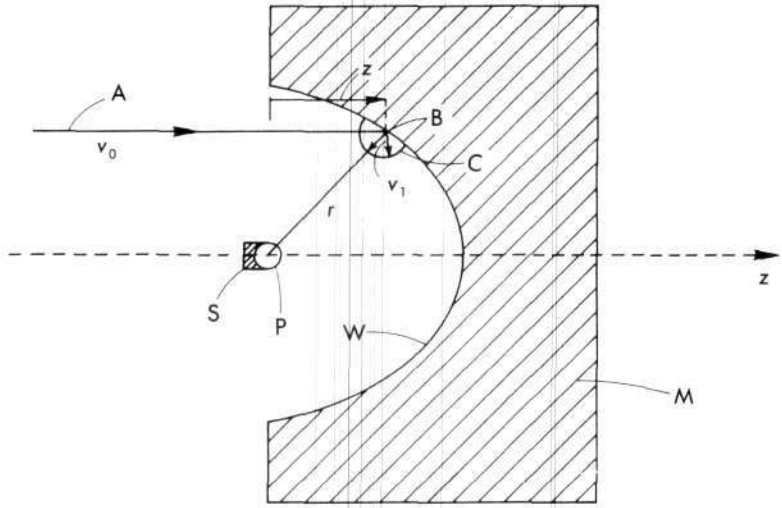

Figure 23. Ignition of fusion pellet P using just one beam of lasers or charged particles to ignite a thermonuclear explosion with shock-wave focusing by a curved wall. A is a ray of the beam striking the curved wall W of a solid block M at point B. C is the spherical shock wave emitted from the impact point B, and S is the radiation shield. The beam is viewed as a bundle of rays with a total cross section equal to the wall diameter.

Next we shall explain how the different beams are produced. Because of the very large energy and power needed this is not a trivial problem and requires the invention of entirely new techniques. We first explain how the intense energetic laser beams can be produced and then go through the table's list, up to the acceleration of macroparticles.

Laser Beams

The laser principle is essentially a quantum-mechanical effect. A laser consists of a medium whose atoms are in an excited, long-lived, metastable state. The excitation into this atomic state can be accomplished by a variety of techniques, including an intense light

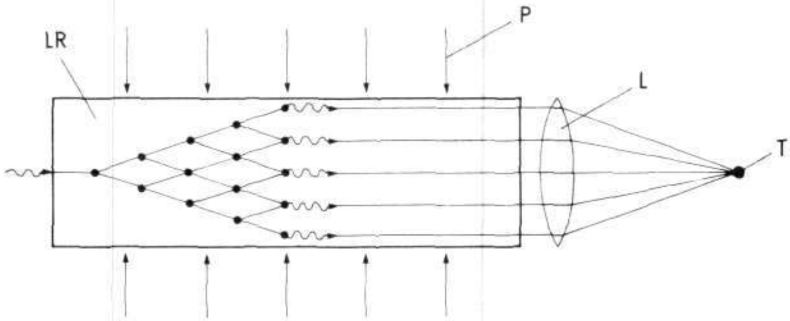

Figure 24. Laser-beam generation and bombardment of thermonuclear target T. P is the pump energy depositing its energy into laser rod LR. Photon avalanche in LR leads to a coherent light beam, which is focused onto T by lens L.

flash (optical pumping), an electric discharge (electric pumping), or a chemical reaction (chemical pumping). This list, although not complete, includes the most important pumping methods. In the case of an electric discharge, the pumping can be done in a variety of ways. Besides using an ordinary gas discharge the pumping can also be accomplished by an intense electron or ion beam.

By the quantum-mechanical laser effect, the excited atoms go into their lower laser state if hit by a photon of the same frequency and direction as the one emitted by that transition. This effect, called *stimulated emission*, causes the atom to decay under photon emission into a lower level, which otherwise would take much longer since the excited level is metastable. Therefore, if one photon of the right frequency is initially released at the left side of the laser rod shown in Figure 24, a growing photon avalanche will move to the right through the laser rod. If the photon density becomes large enough during this avalanche, all excited atoms will be hit by a photon and thereby give off a photon of the same frequency and direction.

To estimate the volume of the laser rod to produce an energy $E \sim 10^6$ J $= 10^{13}$ erg we assume that the photon energy given off by each excited atom is

$$\varepsilon \sim 0.1 \text{ eV} \sim 10^{-13} \text{ erg}.$$

If the number density of the excited atoms in the laser rod is n_0, the total laser volume is determined by putting

$$n_0 V \cdot \varepsilon = E.$$

For the given values of ε and E this leads to $n_0 V \simeq 10^{26}$. The maximum value n_0 can attain is $n_0 \sim 10^{22}$ cm^{-3}, which is the density of the solid state. In reality, however, n_0 is about 10^3 times smaller than this value. Therefore, assuming $n_0 \sim 10^{19}$ cm^{-3}, we find $V \simeq 10$ m^3. This large volume explains the large dimensions of the lasers presently developed to reach thermonuclear ignition.

If the laser has the form of a cylinder, the intense light pulse emerging from the end of the cylinder can be simply focused by the optical lens L onto the thermonuclear target T. The fact that the laser light is coherent and thus parallel makes the focusing very easy. To reach a power of $\sim 10^{14}$ W at 10^6 J, the pulse length must be $\sim 10^{-8}$ sec. The pulse shape of Eq. (141) can be obtained by stacking a large number of laser pulses, which, of course, requires more than one laser rod and a more complicated optical system than that shown in Figure 24.

One of the great disadvantages of using laser beams for thermonuclear ignition is the low conversion efficiency of the primary pumping energy into photon energy. In the case of optical pumping this efficiency can be less than 0.1%. For electric-discharge lasers the efficiency can be much higher, but for the short pulse lengths required here this efficiency is not more than a few percent.

Relativistic Electron Beams

Intense relativistic electron beams[13] can be produced using the Marx generator concept (see Figure 25). Here a high voltage dc source generated by a high-voltage transformer T and rectifier RT charges up a bank of capacitors C in parallel, which are discharged in series over the spark-gap switches SG. The resistors R prevent the capacitors from being shortened over the spark-gap switches. The switching in series adds up all the voltages of the capacitors. For example, if the high-voltage transformer generates 10 kV and there are 100 capacitors, and if these capacitors are switched in series, the final voltage will be 10^6 V.

The high voltage at the end of the column of capacitors switched in series is fed over a low-inductance transmission line TL into the diode D. To ensure a low inductance, a strip or coaxial line must be used. The high voltage suddenly applied to the diode pulls electrons out of the cathode by field emission. These electrons are then accelerated toward the anode, and on their way form an intense megavolt (hence relativistic) electron beam. If some low-density gas is put into the diode, the electric space charge is neutralized and the beam focuses down to a small diameter by magnetic pinch forces.

This focusing effect can then be used to bombard a thermonuclear target by placing it on the anode. To reach a beam power of $\sim 10^{14}$ W, at a beam voltage of $\sim 10^6$ V, would require the enormous current of $\sim 10^8$ A. To avoid the problems resulting from such a large current it is planned to bombard the target with many different beams, for example, with 100 beams each having a current of $\sim 10^6$

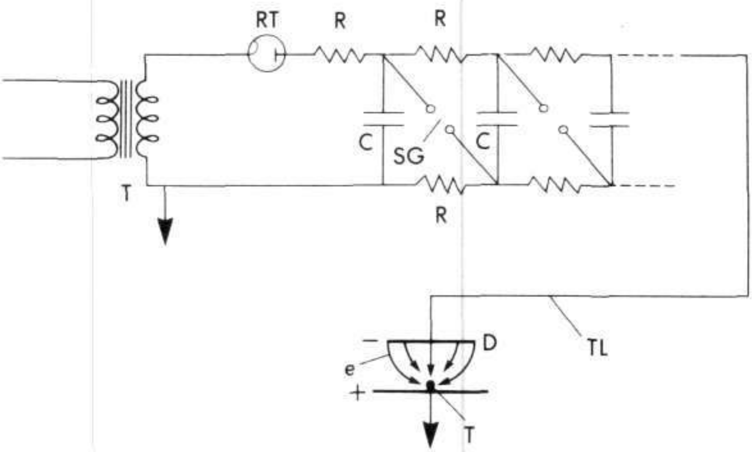

Figure 25. Generation of intense relativistic electron beam by the Marx generator principle. T is a high-voltage transformer; RT is a rectifier; C is a capacitor; SG is a spark-gap switch; R is a resistor; TL is a low-inductance transmission line; and D is a high-voltage diode. The thermonuclear target T is placed on the grounded anode. *e* is an electron trajectory of the electron beam focusing within D.

A. This, of course, requires that each beam pass through the anode into a drift space, with all the beams meeting at the position of the thermonuclear target.[14,15]

The electric energy density in ordinary capacitors is of the order $\sim 10^{-1}$ J/cm^3. Hence, to store an energy of $\sim 10^6$ J again requires a volume of ~ 10 m^3. Substantially smaller volumes can be obtained with water capacitors, but in this case the charging has to be done very fast or the water suffers dielectric breakdown. The rapid charging in principle could be done by magnetohydrodynamic power conversion of a thermonuclear microexplosion itself.

One problem of the electron-bombardment method is that relativistic electrons are difficult to stop in the small thermonuclear target.

Light Ion Beams

Light ion beams[16] are generated by a slight modification of the existing technique to produce relativistic electron beams. The trick is to prevent the electrons from crossing the diode gap by applying a strong magnetic field in a direction perpendicular to it (see Figure 26). Then, if $H > E$, where H is the magnetic and E the electric field (both measured in cgs units), the electron flow from the cathode to the anode is greatly suppressed, because the electrons undergo a drift motion perpendicular to E and H. If now the anode is covered with a plasma to serve as an ion source, an intense ion beam is formed and accelerated from the anode to the cathode. Assuming that $H = 2 \times 10^4$ G, which can be easily reached with electromagnets, this then implies that

$$E \lesssim 2 \times 10^4 \text{ esu} = 6 \times 10^6 \text{ V/cm}.$$

Therefore, with a diode gap of a few centimeters and with a voltage of several megavolts applied to it, an intense multimegavolt ion beam can be produced.

One immediate advantage of using light ion beams produced in this way is that the beam-stopping problem is largely eliminated. Another advantage is the low nonrelativistic ion velocities of approx-

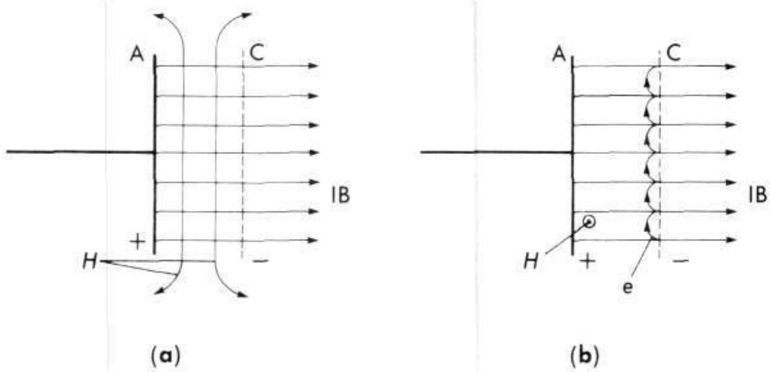

Figure 26. The principle of the magnetically insulated diode: (a) with the magnetic lines of force H in the plane of the figure and (b) with the magnetic lines of force perpendicular to the plane of the figure. Light ion beams IB are generated by magnetically confining the electrons to the cathode. e is the cycloidal trajectory of an electron.

imately $c/10$, which permit the ions to be bunched by at least a factor of ~ 10 if the voltage pulse on the diode is properly programmed.[17] The described beam generation and bunching method presently used for intense light ion beams can, of course, also be used for intense heavy ion beams.

Heavy Ion Beams

The expression "heavy ion fusion" has been coined to embrace proposals to use a semi-intense heavy ion beam of near-relativistic particle velocity for thermonuclear ignition.[16,18,19] Here the beam voltages are of the order of ~ 10 GeV, rather than $\sim 10^6$ eV. Accordingly, to produce the same beam power of $\sim 10^{14}$ W, currents smaller by a factor $\sim 10^4$, from $\sim 10^8$ A down to $\sim 10^4$ A, are sufficient. If accelerated to an energy of ~ 10 GeV a light ion would have relativistic velocities and the same problem of beam energy deposition encountered in the case of relativistic electron beams would occur. This does not happen, however, for heavy ions, for example, uranium ions, which for an energy of ~ 10 GeV are still nonrelativistic. The accelerator to produce these heavy ion beams can make use of conventional techniques, such as the drift tube type shown in Figure 27, where the beam draws its energy from electromagnetic energy stored in microwave cavities. The accelerating voltages are here $\sim 2 \times 10^6$ V/m. Hence, to reach 10 GeV requires an accelerator ~ 5 km long. In an early proposal for heavy ion fusion[16] it was suggested that a magnetically insulated electrostatic energy storage device be used from which the beam would draw its energy. Such a device promises the attainment of gigavolt potentials in one step and would make the many-kilometers-long accelerator unnecessary.

To focus the beam at the end of the accelerator a magnetic lens can be used. The focusing ability, although not as good here as for lasers, while better than for light ion beams, is nevertheless fully sufficient.

Microparticle Beams

The particles in microparticle beams[20,21] have a size of $\sim 10^{-6}$ cm with a mass approximately 10^6 times larger than atoms and hydrogen or about 10^4 times heavier than uranium ions. Consequently the velocities are even smaller here, typically $\sim 10^8$ cm/sec. This is an energy of ~ 10 keV per atom, which is exactly right for thermonuclear ignition. Microparticle beams would therefore actually be

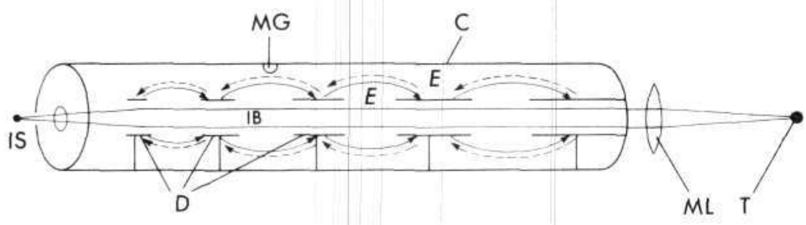

Figure 27. The principle of a conventional ion accelerator proposed for heavy-ion fusion. IS is the ion source; E is the time-varying oscillating electric-field vector between the drift tubes D within the cavity C. MG is a microwave power supply and ML a magnetic lens focusing the ion beam IB onto the thermonuclear target T.

even better than heavy ion beams. Because the technology to make such beams is highly underdeveloped, no efforts in this direction are under way at the present time, although this situation is expected to change.

Macroscopic Projectiles

Rather than accelerating a beam of particles one may just accelerate one large particle or projectile to a velocity in excess of ~100 km/sec. This concept has been called *impact fusion*. The only credible approach known for reaching the high velocities required is by magnetic acceleration of a superconducting solenoid.[22,23] This is explained in Figure 28. The projectile here is accelerated by a traveling magnetic wave produced by a lumped parameter transmission line with a variable LC product (L is the inductance and C is the capacitance). Again, as in heavy ion fusion, the greatest obstacle is the length of the accelerator, which will be many kilometers.

Other Ignition Concepts

Besides these six ignition concepts, all of which are based on the conversion of kinetic beam energy into heat upon impact, from time

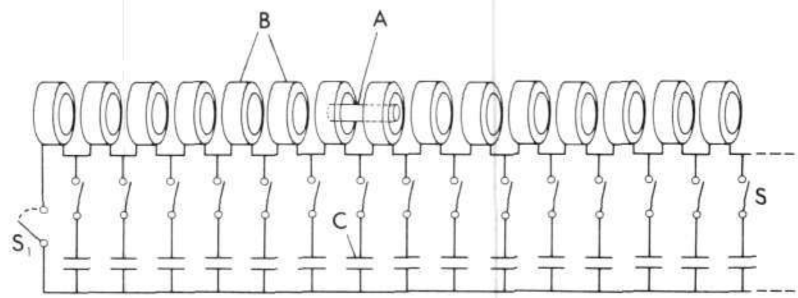

Figure 28. Magnetic traveling-wave accelerator: The projectile A is a small superconducting solenoid that is accelerated magnetically through external field coils B. C are capacitors and S_1, . . . , S are switches to be closed as the projectile moves down the accelerator tube.

to time proposals have been made using different concepts. One proposal quite popular at its time was to use exploding wires. It is doubtful, however, if exploding wires alone can reach ignition conditions.

One way to increase the performance of exploding-wire heating is to combine it with some other pulsed source of energy, such as implosions induced by chemical high explosives. This possibility, which was discussed in Chapter 14, is unlikely to lead to a thermonuclear reactor, however. Combining exploding wires with beam energy sources is also possible. But, again, the complexity associated with a practical target configuration that could be incorporated into a reactor makes this concept equally doubtful as a potential contender.

There is one alternative approach, though, which is singularly promising. It is the idea of imploding blackbody radiation inside a small cavity, which in turn would implode a thermonuclear target placed inside the cavity. We discussed this concept in the last chapter in the context of a chemically triggered H-bomb. But here we would like to use the same principle to ignite a microexplosion.

To implode a cavity with the required velocity of $\gtrsim 50$ km/sec still requires some other source of energy. Any one of the six beam options, ranging from lasers up to macroparticles, can be used as such a source. However, because the implosion velocity of the cavity is so much smaller than for direct thermonuclear ignition and because the cavity has a much larger diameter, both the required beam power and beam-focusing ability are here greatly reduced. In contrast to direct ignition, the beam power is about 10^{12} W, to be focused down to a radius of ~ 1 cm. The total energy, of course, is unchanged and of the order $\sim 10^6$ to 10^7 J.

Figure 29 shows how this can be done in the case of laser or charged particle beams. Several beams (B) are projected onto an ablator (A), evaporating it at a high temperature. The resulting high ablation velocity leads to a rocketlike implosion of the dense pusher P forming the cavity wall. The large implosion velocity v shock-heats a low-density, high-atomic-weight gas inside the cavity. The heated gas becomes an intense source of blackbody radiation, which is further compressed, and at peak power finally bombards the thermonuclear target T uniformly from all sides. In Figure 30 the same

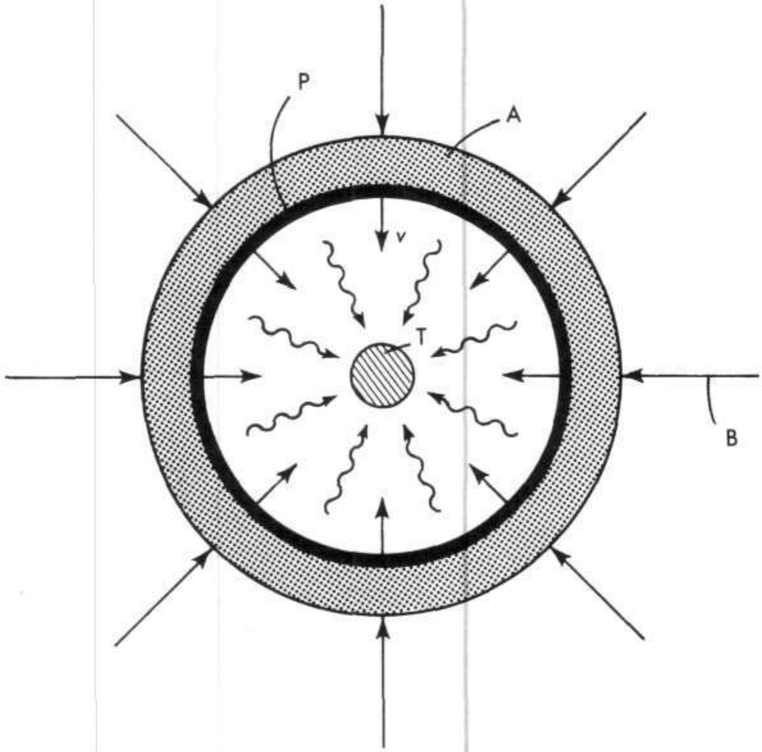

Figure 29. Implosion of blackbody radiation by an ablatively driven shell. B is an incoming laser or charged-particle beam; A is the ablator; P the pusher; and T the thermonuclear pellet inside a cavity filled with blackbody radiation.

concept is shown as applied to impact fusion by a hypervelocity projectile. Here the cavity is conical. If macroparticles are used to implode the cavity, the rather low velocity of ~50 km/sec opens up the possibility for other means of acceleration besides the traveling-wave acceleration of a superconductor.

If thermonuclear microexplosions can be ignited by any one of these described methods, this would have far-reaching consequences in at least two areas:

- An unlimited, clean source of energy ultimately depending only on deuterium as a source material would become possible. Magnetic

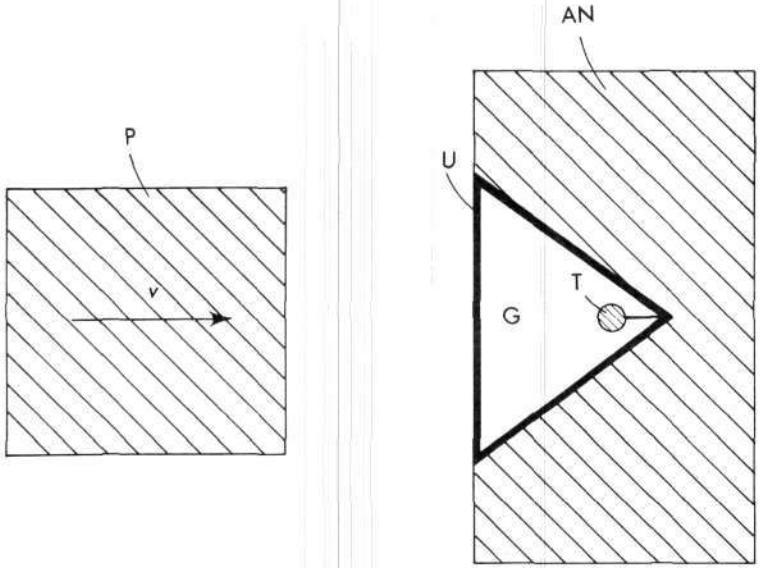

Figure 30. Implosion by hypervelocity impact of blackbody radiation trapped inside a conical cavity. P is the hypervelocity projectile moving with velocity v to hit the high-atomic-weight gas G inside the cavity. U is a thin but dense high-atomic-weight material (uranium, for example) covering the inner surface of the cavity, which is surrounded by the anvil AN. T is the thermonuclear target to be imploded by the blackbody radiation.

fusion, which depends on lithium as a source material to make tritium, is "dirty" in comparison. The microexplosions would be confined within a large container,[16] whose walls would have to be protected by a liquid layer using, for example, a liquid metal. If the container were permeated by a magnetic field some of the explosively released energy could be converted directly into electric energy. Figure 31 shows my 1969 design of this concept.

- The prospect for an efficient rocket-propulsion system by which large payloads could be moved at great speed ($\sim 10^3$ km/sec) within the solar system would also become possible. Here the microexplosions would take place in the focus of a concave magnetic reflector. The magnetic field required for the reflector can be generated by superconducting magnetic field coils.[24] A magnetohydrodynamic loop could there pick up a fraction of the explosively released energy

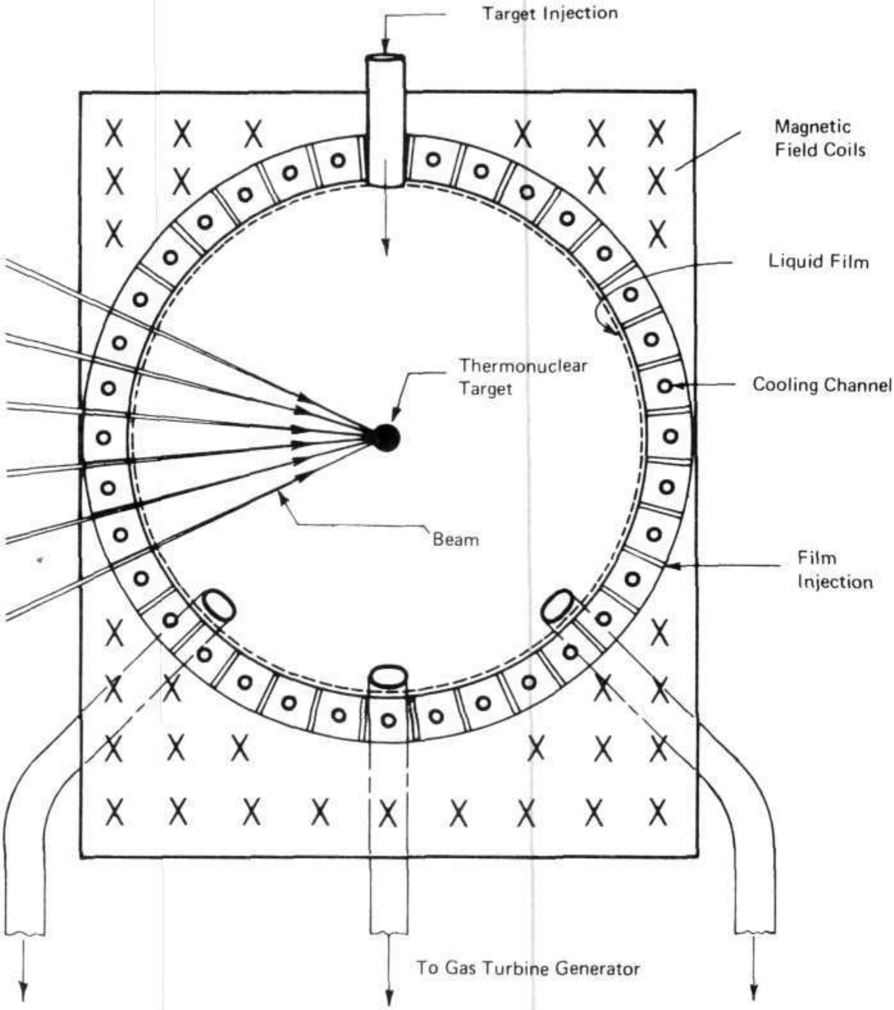

Figure 31. Schematic drawing of a thermonuclear reactor based on the confinement of a chain of microexplosions inside a spherical chamber. The beams are laser or particle beams.

to charge up the capacitors for the energy pulse that triggers the following microexplosion. Figure 32 shows my 1970 design of this propulsion concept.

Such a propulsion system would one day make even interstellar

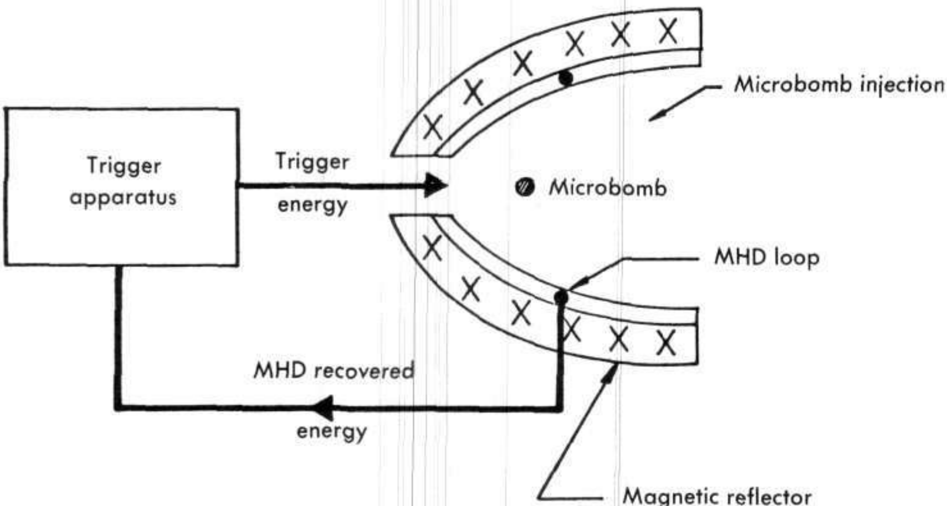

Figure 32. Schematic drawing of a nuclear microexplosion unit to be used for an efficient rocket-propulsion system by which large payloads could be moved at great speed within the solar system.

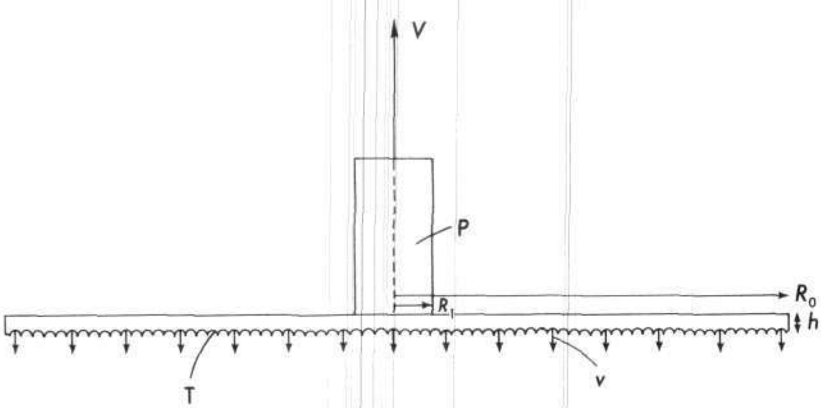

Figure 33. An interstellar spaceship (schematic) using the propulsion unit shown in Figure 32. This ship could weigh millions of tons and travel at one-tenth the velocity of light. T is one of the 10,000 microexplosion propulsion units, and P is the payload.

space flight at velocities approximately one-tenth the velocity of light a reality, and eventually lead to the colonization of the entire galaxy.[25-28] In Figure 33 an interstellar space ship weighing millions of tons and propelled with about 10,000 microexplosion propulsion units is shown. The dimensions of this spacecraft are such that it could carry a large crew of thousands of people and a payload of millions of tons to a neighboring solar system in a travel time of several decades.

CHAPTER SIXTEEN

Thermonuclear Lenses And Shape Charges

In this chapter we shall discuss some unusual thermonuclear configurations. For conventional high explosives, detonation wave-shaping techniques have been developed into an art. In the so-called explosive-lens technique, for example, it is possible to obtain perfectly plane, convergent cylindrical and spherical detonation waves. In other related configurations known under the name *shape charges*, very fast jets with velocities greatly exceeding the detonation velocity can be produced.

We shall first discuss the explosive-lens techniques. In one of these wave-shaping techniques two different detonation velocities are used. This technique is not very practical under circumstances where thermonuclear explosives are involved because of the greatly differing temperatures for essentially all thermonuclear explosives. This is especially true if DT and D are used as the two explosives.

An alternative proven technique for wave shaping is to place holes or inert bodies in the path of the detonation front around which the wave has to pass.[29] This technique also works with just one explosive. It is therefore ideally suited for wave shaping of thermonuclear detonation fronts. As we shall see below, wave shaping of thermonuclear detonations may have great importance not only for certain technical applications but also for basic science.

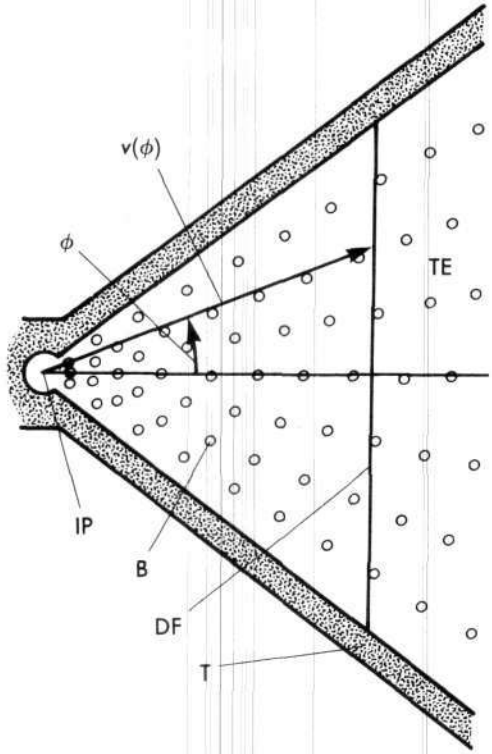

Figure 34. Thermonuclear plane wave lens by which the detonation wave can be shaped by bubbles or inert obstacles B placed in the path of the detonation front DF moving from the ignition point IP of the thermonuclear explosive TE. T is the tamp.

The wave-shaping technique is most easily explained for the example of a thermonuclear plane-wave lens. This configuration is shown in Figure 34. A thermonuclear detonation front starting from the ignition point IP propagates into the thermonuclear explosive TF, conical in shape and surrounded by the tamp T. In the path of the detonation front, hollow bubbles or solid inert bodies B are placed as shown. The density $\rho(B)$ of these obstructions as a function of the radial distance r from the axis of the conical assembly is now chosen in such a way as to give the detonation front velocity v as a

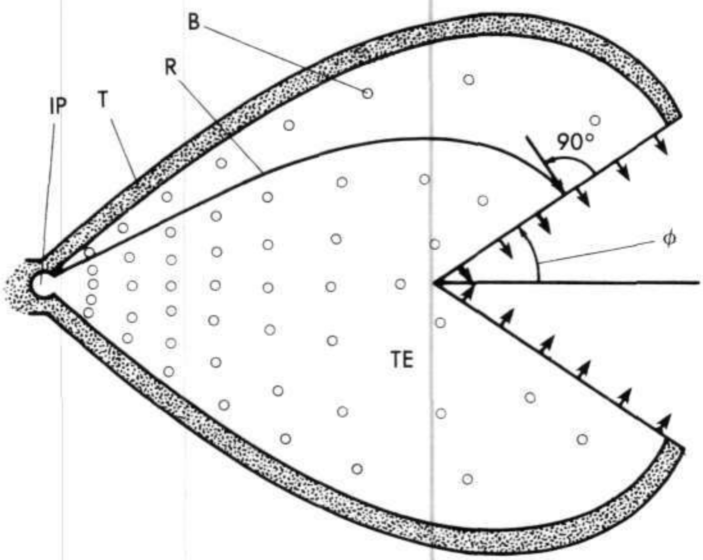

Figure 35. A wave-shaping lens for conical implosion where the detonation lens produces a convergent conical wave. IP is the ignition point of the thermonuclear explosive TE; B are the bubbles placed in the wave path; T is the tamp; and R is a ray of the detonation wave.

function of the polar angle ϕ the following dependence:

$$v(\phi) = \text{const}/\cos\phi, \tag{142}$$

because it is exactly this dependence that will result in a plane wave emerging from the ignition point IP. Since it is reasonable to assume that the wave speed is inversely proportional to the density of the obstacles, we thus would here have to put

$$\rho(B) = \text{const}(\cos\phi). \tag{143}$$

Other distributions $\rho(B) = f(r,\phi)$ obviously permit the attainment of other wave shapes. In Figure 35, for example, a wave-shaping lens is shown that is designed to result in a conical implosion. In this lens configuration all rays R of the detonation wave obey the constant time relation

$$t = \int ds/v = \text{const},$$

where ds is the line element along the ray and v is the local detonation velocity, with the integral taken from the ignition point to the surface of the cone. In Figure 36, finally, it is shown how a spherical implosion can be obtained by wave shaping. To obtain a cylindrical implosion a more complicated three-dimensional lens configuration is needed if the detonation wave starts from a point.

Next we shall discuss the concept of shape charges in the context of thermonuclear explosives. A good shape charge requires two things: (1) a well-shaped wave and (2) a liner propelled by the wave. The most simple shape charge is shown in Figure 37. A thermonuclear detonation starting from the ignition point IP enters the conical plane wave lens PW. After leaving the lens the plane wave thus produced propagates into the thermonuclear explosive TE. This

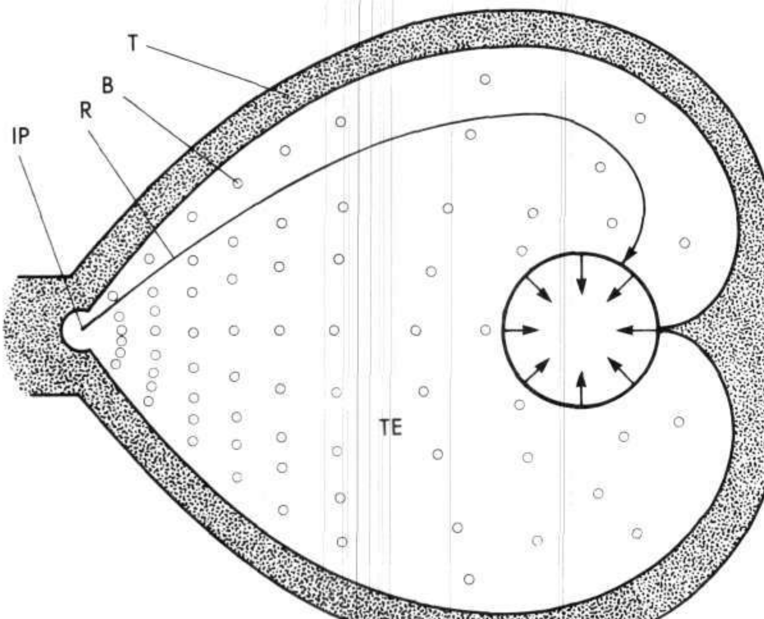

Figure 36. A spherical implosion can be obtained by wave shaping. The detonation lens here produces convergent spherical waves. IP is the ignition point; TE is the thermonuclear explosive; T is the tamp; R is a ray of the detonation wave; and B are the bubbles.

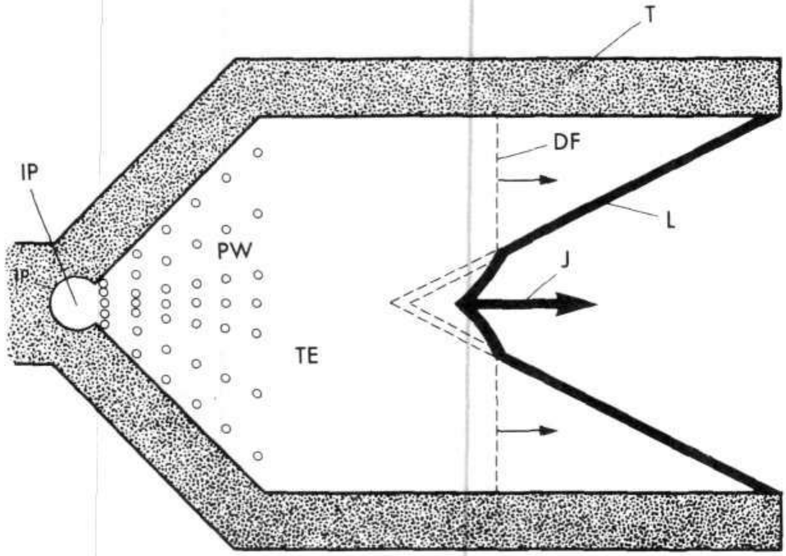

Figure 37. Simple thermonuclear shape charge: IP is the ignition point; PW is a plane wave lens; TE is the thermonuclear explosive; DF is the detonation front; L is a metallic liner; J is a jet of the collapsing liner; and T is the tamp.

wave will then collapse the conical metal liner L toward its axis, resulting in a fast forward metal jet J. Figure 37 shows the partially collapsed liner at the moment the detonation wave has reached the position indicated by the line *DF*. This most simple shape charge leads to maximum jet velocities only twice as fast as the detonation velocity. Since the density of the metal liner can be roughly 10 times larger than the density of the thermonuclear explosive, the energy density of the jet is roughly 40 times larger than in the exploding fusion material.

A shape charge leading to much higher velocities uses a conical implosion onto a conical liner. How a conical implosion can be made by wave shaping was shown in Figure 35. In this case the resulting jet velocity is inversely proportional to $\sin\phi$, where ϕ is the angle shown in Figure 35. From this it seems to follow that by decreasing the cone angle ϕ, the jet velocity could become arbitrarily large. However, a lower practical cutoff for ϕ has been found to be arc$\phi \simeq 0.1$. Therefore, the maximum attainable jet velocity is about 10 times larger than the detonation velocity. For deuterium the detonation velocity is about 1/100 the velocity of light, and it therefore follows that a jet roughly 10% the velocity of light could be reached.

We shall now briefly discuss three applications for three thermonuclear lens and shape-charge techniques.

(**1**) For thermonuclear rocket propulsion, explosive yields in the 10^2 to 10^3 ton TNT-equivalent range are of considerable interest, even though yields of this magnitude make the magnetic reflector concept shown in Figure 32 impractical. For yields in this range the much simpler, although less efficient, pusher-plate concept shown in Figure 38(a) is much better suited. This concept was originally developed for the Orion fission-bomb nuclear-pulse rocket.

In the pusher-plate concept, however, efficiency is lost not only by the absence of a magnetic field reflector, but, more importantly, also by the fact that the explosion must take place a safe distance

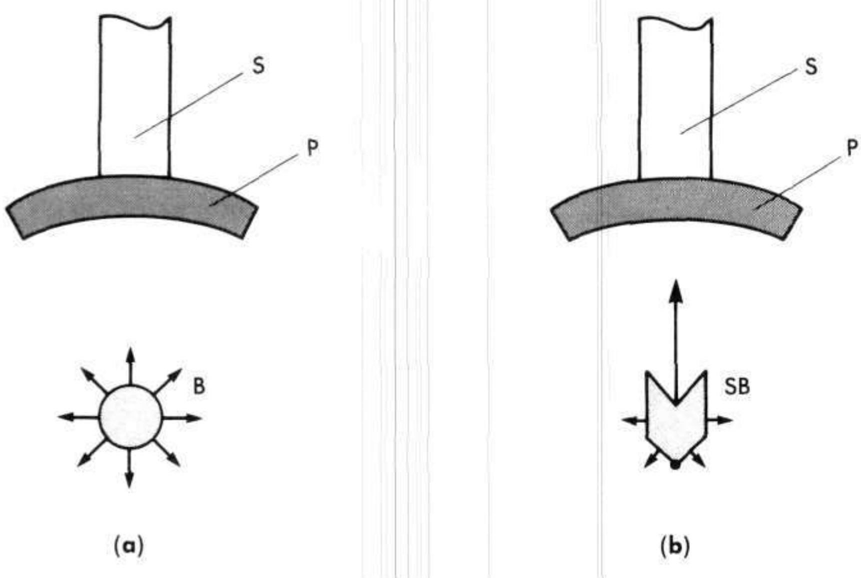

Figure 38. Simple Orion-type pusher-plate configuration for spherically exploding fission bomb (a). The same pusher-plate configuration is shown in (b) with asymmetrically exploding thermonuclear shape-charge bomb SB. S is a shaft connecting the pusher plate to the payload.

away from the plate, ruling out the paraboloidal reflector configuration shown in Figure 32. The reason for this loss in efficiency, of course, is the loss of radial confinement, which greatly reduces the amount of momentum that can be transferred from the explosion to the pusher plate. This loss in momentum transfer can be greatly reduced if thermonuclear shape charges are used. An arrangement applying the shape-charge technique is shown in Figure 38(b).

(**2**) The pressure in a plane thermonuclear detonation wave is given by

$$p = 8n_0kT, \tag{144}$$

where n_0 is the initial density of the explosive and T the combustion temperature. In a thermonuclear explosion the temperature can reach the value of $T \simeq 10^9$ K. Assuming that $n_0 = 5 \times 10^{22}$ cm^{-3}, corresponding to liquid deuterium, one finds that

$$p = 5.5 \times 10^{16}\ \text{dyn/cm}^2.$$

If a detonation wave with this strength implodes a spherical liner using the lens configuration shown in Figure 36, the velocity of the inner liner surface will rise as $r^{-3/2}$, where r is the cavity radius.[7] Hence, with an approximately 20-fold decrease in r and an initial implosion velocity of

$$v = \sqrt{3kT/M} \simeq 3 \times 10^8\ \text{cm/sec}$$

in deuterium at $T = 10^9$ K, one would reach velocities approaching the velocity of light if the initial cavity radius is chosen to be ~ 10 cm, and if it implodes down to ~ 0.5 cm, a quite realistic assumption. The stagnation pressure of this inward motion, which is initially $\sim 5 \times 10^{16}$ dyn/cm^2, would rise as r^{-3} up to $\sim 5 \times 10^{20}$ dyn/cm^2. The attainment of these large pressures in such comparatively large volumes, simulating conditions in white dwarf stars, should be of considerable interest in basic physics.

(**3**) If the described implosion technique is used to implode and amplify a magnetic field, magnetic field strengths up to $H \sim 10^{11}$ G could be obtained, if we equate the stagnation pressure of $p = 5 \times 10^{20}$ dyn/cm^2 with the magnetic pressure $H^2/8\pi$.

CHAPTER SEVENTEEN

Some Recent Developments

In this chapter we shall present several proposals that have only recently emerged. Two of these proposals combine fusion by hypervelocity impact with the fast-liner approach. Impact fusion, which was explained in Chapter 15, is the idea of triggering a thermonuclear microexplosion by the impact of a hypervelocity projectile onto a thermonuclear target. It greatly suffers from the required minimum velocity of ~200 km/sec, which is difficult to reach. A reduction of the needed velocity down to ~50 km/sec is possible, however, with the concept of imploding a cavity filled with blackbody radiation.

We shall now show that a further reduction of the velocity down to ~20 km/sec seems possible, even for arbitrarily large yields. In this concept impact fusion is combined with the fast-liner approach to nuclear fusion and, furthermore, with the concept of target staging. *Target staging* is the concept whereby the energy released in one microexplosion is used to ignite a larger, second-stage microexplosion.

In the fast-liner approach to nuclear fusion, a magnetized plasma having a density less than solid density is imploded together with the magnetic field by some external power source.[30–33] As the driving power source a large capacitor bank is normally used. The discharge of the bank generates a strong $j \times H$ force, acting on a metallic liner by either direct or induced currents. For example, if the liner has the form of a cylindrical shell, an axial current j_z discharged along the liner surface produces an azimuthal magnetic field H_ϕ with the resulting force directed radially inward, or a fast-rising axial magnetic field $H_z(t)$ induces in the liner surface an azimuthal current j_ϕ, again with the resulting force directed radially inward.

These fast-liner concepts require implosion velocities of $\gtrsim 20$ km/sec. Velocities of this magnitude can be reached with large capacitor banks if the liner to be imploded is sufficiently thin. But since the liner has to be rather thin in this case to reach these velocities, the time to hold the plasma together is not very long. As a result, the plasma density must be rather low and the dimensions of the imploded magnetized plasma rather large to satisfy the Lawson $n\tau$ breakeven condition. For this reason, several authors suggested producing the needed implosion velocities by hypervelocity projectile accelerators rather than by magnetic-implosion techniques.[34,35] This would permit much higher plasma densities and hence much smaller plasma dimensions. But even here the target density would be lower than solid density, which would result in a comparatively low thermonuclear gain. Therefore, unless the target is still quite large, no net energy can be produced.

This drawback can be overcome in the novel approach suggested here. In this approach an ~ 20 km/sec projectile again serves to ignite a thermonuclear reaction in a magnetized, less-than-solid-density target, but the energy released in this low-density, low-yield target is now used to ignite a second high-density, high-yield target. This two-stage target promises very large final thermonuclear yields even with a comparatively low initial impact velocity. Furthermore, the required velocities are here so small that they can be cheaply produced by magnetic acceleration or isentropic-light gas guns.[36]

The basic principle of our approach is explained in Figure 39. In Figure 39(a) an incoming projectile having a conical hole implodes a cylindrical target chamber, which is positioned in the conical hole of an anvil A. The space inside the target chamber is filled with DT gas of relatively low density and serves as the first-stage target I. The DT gas is furthermore permeated by a magnetic field of $H_0 \sim 10^5$ G, which can be set up if the target chamber consists of a one-turn pulsed solenoid. Alternatively, it may also be produced by using for the target chamber a small, hollow, superconducting coil, or by strong ferromagnets like gadolinium with a saturation field strength of $\sim$60 kG, if the upper and lower ends of the target chamber consist of permanent magnets.

Just prior to the moment shown in Figure 39(a) where the incoming projectile strikes the target chamber, a short-pulse, relatively low-energy beam passes through an opening O into the target chamber, preheating the DT gas to about $T_0 \sim 10^6$ K. The resulting

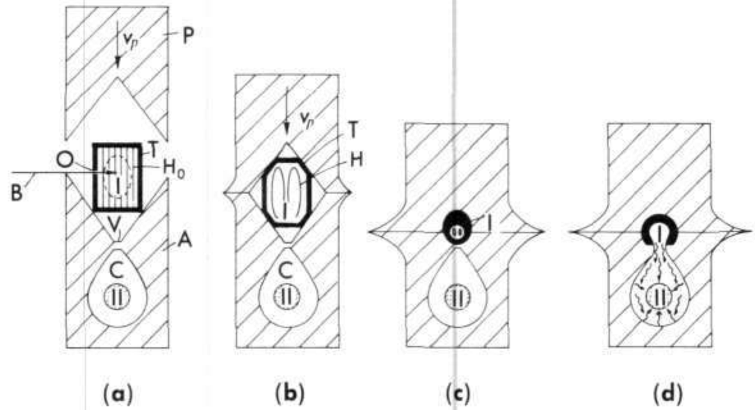

Figure 39. In the two-stage magnetic booster impact-fusion target, the first target I is the booster target and the second target of thermonuclear fuel II is the high-yield target. P is the incoming projectile moving with velocity v_p and B is a laser or charged-particle beam that passes through opening O into T, the booster target chamber. V is the conical vertex position, H_0 is the initial magnetic field inside T. (a) An incoming projectile implodes a low-density DT gas that has been magnetized and preheated by the laser or charged-particle beam. (b) Magnetic field reversal closes the field lines and the target I is highly compressed. (c) The magnetic field rises to maximum compression and the target chamber is at its minimum diameter. (d) The DT plasma in I reaches its ignition temperature and ruptures the cavity wall at V, releasing large amount of energy into the adjacent chamber C. (e) The radiation and hot plasma released into C implode ablatively and ignite the high-yield target II.

magnetized DT plasma confined in the target chamber is then highly compressed by the rapidly moving projectile, making the chamber shrink in size. This leads to a rapid rise of the magnetic field trapped inside the target chamber and, by currents induced in the plasma and the wall, results in magnetic field reversal closing the field lines. This moment is shown in Figure 39(b). After field reversal has taken place the heat-conduction losses of the DT plasma to the walls of the target chamber are restricted to those perpendicular to the magnetic field, greatly increasing the energy confinement time.

If all energy losses are neglected, the plasma temperature and magnetic field in the target chamber will rise as

$$T/T_0 = H/H_0 = (\ell_0/\ell)^2, \tag{145}$$

where ℓ_0 and ℓ, respectively, are the linear dimensions of the target chamber at the beginning of the implosion process and some time later. Field reversal occurs at $H \gtrsim 2H_0$, that is, for

$$\ell \lesssim \ell_0/\sqrt{2}.$$

Since a temperature of $T_i \simeq 10^8$ K is needed to reach thermonuclear ignition, the cavity must be imploded to a minimum diameter

$$\ell_{\rm min} \simeq \ell_0/10.$$

Therefore, apart from the initial phase of the implosion process, the magnetized plasma is in a state of complete field reversal most of the time. At the final minimum diameter of the target chamber, the magnetic field has risen to its maximum value $H_{\rm max} \simeq 10^7$ G. This state of maximum compression is shown in Figure 39(c).

After reaching the ignition temperature T_i the DT plasma confined inside the target chamber makes a thermonuclear excursion, greatly increasing its total energy. The rapid rise in the internal energy caused by the thermonuclear reactions leads to a rapid rise in the plasma pressure, which eventually ruptures the target chamber wall at its weakest point. If this weakest point is chosen to be at the vertex point V of the conical part of the chamber formed by the anvil, a large amount of energy in the form of radiation and hot plasma will be released into the adjacent chamber C inside the anvil, into which a second-stage dense thermonuclear target II is placed. This event is shown in Figure 39(d). The energy flowing into the chamber C will then implode ablatively and ignite target II. It is this second-stage target II that releases most of the energy. Only in this two-stage target arrangement are substantially smaller impact velocities permitted without paying the price of a small gain.

These general ideas will now be supported by some more quantitative estimates:

(**1**) Let us assume that the projectile has a density $\rho_p \simeq 10$ g/cm^3 and moves with $v_p = 2 \times 10^6$ cm/sec, leading upon impact to a stagnation pressure

$$p_s = (1/2)\rho_p v_p^{\,2} = 2 \times 10^{13} \text{ dyn/cm}^2,$$

which must be equal to the final plasma pressure in the target chamber $p = 2nkT_i$ at $T_i \simeq 10^8$ K. We then find that $n \simeq 10^{21}$ cm^{-3}. The initial density of the DT gas inside the target chamber is smaller by the factor

$$(\ell_{\min}/\ell_0)^3 \simeq 10^{-3}.$$

Therefore, the initial density of the gaseous DT target is $n_0 \simeq 10^{18}$ cm^{-3}.

(**2**) Let us assume that the initial and final diameter of the target chamber are $\ell_0 \simeq 4$ cm and $\ell_{\min} \simeq 0.4$ cm, respectively. The volume at maximum compression is thus

$$\ell_{\min}{}^3 \simeq 6 \times 10^{-2} \text{ cm}^3,$$

and the total number of atoms in the chamber is $N \sim 6 \times 10^{19}$.

(**3**) Because the implosion of the chamber is three-dimensional, the time in which the DT gas is heated from $T_0 \simeq 10^6$ K up to $T_i \simeq 10^8$ K, under the assumption that the compression is completely isentropic, is given by

$$\tau_A \simeq \ell_0/2v_p. \tag{146}$$

In our example we find $\tau_A \simeq 10^{-6}$ sec. This time has to be smaller than the radiative loss time τ_R from bremsstrahlung and the heat-conduction loss time τ_c in the presence of a strong transverse magnetic field. Otherwise our assumption that the DT gas is isentropically compressed is invalid. The loss time for bremsstrahlung is given by

$$\tau_R \simeq (3 \times 10^{11}) T^{1/2}/n, \tag{147}$$

and for heat conduction by

$$\tau_c \simeq (2.5 \times 10^{-2})(\ell_{\min} H)^2 T^{1/2}/n. \tag{148}$$

In our example with

$$n = 10^{21}\ \text{cm}^{-3}, \qquad H = 10^7\ \text{G}, \qquad T = 10^8\ \text{K},$$

and

$$\ell_{\min} = 0.4\ \text{cm},$$

we find

$$\tau_R \simeq 3 \times 10^{-6}\ \text{sec} \quad \text{and} \quad \tau_c \simeq 4 \times 10^{-6}\ \text{sec}.$$

The assumption of isentropic compression is therefore reasonably well satisfied.

(**4**) To heat a plasma composed of $N \simeq 6 \times 10^{19}$ ions to $T_0 = 10^6$ K requires the energy

$$E_0 = 3NkT \simeq 2.4 \times 10^{10}\ \text{erg} = 2.4\ \text{kJ}.$$

This relatively small energy required for preheating can be easily supplied by a short-pulse laser or charged-particle beam. Because the beam pulse has to enter the chamber through a small opening in the target chamber, a laser beam seems to be better suited for this purpose. Furthermore, since the initial density of the DT gas is rather low, an infrared gas laser of high efficiency can be used.

(**5**) To heat the DT gas by isentropic compression to the ignition temperature $T_i \simeq 10^8$ K requires that its internal energy be raised to

$$E_i = 3NkT_i \simeq 2.4 \times 10^{12} \text{ erg}.$$

Assuming pessimistically that only about 1% of the kinetic projectile energy goes into this internal energy, its energy would therefore have to be 2.4×10^{14} erg. With a projectile velocity of 2×10^6 cm/sec, the projectile mass is then $m_p = 120$ g. The remaining 99% of the projectile energy would not be lost but would serve to inertially confine the target. Under this assumption most of the energy is used for inertial confinement and not for ignition, similar to laser or charged-particle beam fusion.

(**6**) After the DT plasma has reached the thermonuclear ignition temperature $T_i \simeq 10^8$ K the part of the thermonuclear energy set free in the form of α-particles is dissipated within the DT plasma because the Larmor radius of these α-particles at $H = 10^7$ G is $r_L \simeq 0.03$ cm, and is thus more than 10 times smaller than the diameter of the imploded chamber, which is $\ell_{\min} \simeq 0.4$ cm. As a result the DT plasma undergoes a thermonuclear excursion, raising its temperature to much higher values as long as the inertial confinement lasts. The inertial confinement time is of the order

$$\tau_i \simeq h/v_p, \tag{149}$$

where h is the thickness of the dense material composing both the projectile and the anvil tamping the chamber. The value of h can be estimated putting $h^3\rho_p = m_p$, which in our example, with $\rho_p = 10$ g/cm^3 and $m_p = 120$ g, gives $h = 2.3$ cm. It thus follows that $\tau_i \simeq 10^{-6}$ sec. The fuel burnup time, on the other hand, is given by

$$\tau_b \simeq (n\langle \sigma v \rangle)^{-1}. \tag{150}$$

In a thermonuclear excursion the temperature rises until $\langle \sigma v \rangle$ has reached its maximum, which is

$$\langle \sigma v \rangle_{\max} \simeq 10^{-15} \text{ cm}^3/\text{sec}$$

for the DT reaction and is reached at a temperature of $\sim 8 \times 10^8$ K. Using our value of $n = 10^{21}$ cm^{-3}, we thus find that $\tau_b \simeq 10^{-6}$ sec.

(**7**) Since $\tau_i \simeq \tau_b$ we may assume a large fuel burnup, for example, 50%. The total energy released into α-particles, each having a kinetic energy of 2.8 MeV, for a DT plasma of 6×10^{19} ions, is thus given by

$$E_\alpha = \tfrac{1}{2}(\tfrac{1}{2}6 \times 10^{19})(4.5 \times 10^{-6}) = 3.4 \times 10^{14}\ \text{erg} = 34\ \text{MJ}. \qquad (151)$$

This energy suddenly released inside the target chamber will raise the plasma pressure to the order

$$p \sim E_\alpha/\ell_{\min}{}^3 \simeq 5 \times 10^{15}\ \text{dyn/cm}^2,$$

which is about 100 times larger than the magnetic pressure at 10^7 G. As a result the hot plasma will convectively mix with the wall material. Because of this mixing effect most of the energy will go into blackbody radiation. The temperature T_b of this blackbody radiation is determined by

$$aT_b{}^4 = E_\alpha/\ell_{\min}{}^3, \qquad (152)$$

where

$$a = 7.67 \times 10^{-15}\ \text{erg/cm}^3\text{-K}^4,$$

and one finds that $T_b \simeq 3 \times 10^7$ K. If we assume that the chamber is permitted to expand approximately 3-fold from the high pressure before the weak point at the vertex position breaks, the temperature would go down to $T_b' = T_b/3 \simeq 10^7$ K.

(**8**) After rupture of the weak vertex point the photon energy flux into the cavity C inside which the high-yield thermonuclear target is placed is given by

$$P = \sigma T_b'^4, \qquad (153)$$

where

$$\sigma = ac/4 = 5.75 \times 10^{-5}\ \text{erg/cm}^2\text{-sec-K}^4.$$

With $T_b' = 10^7$ K, we find that

$$P = 5.75 \times 10^{23}\ \text{erg/cm}^2\text{-sec} = 5.75 \times 10^{16}\ \text{W/cm}^2.$$

If the cross section of the opening formed at the breaking point through which the energy can flow is of the order $\ell_{min}{}^2 \sim 10^{-1}$ cm^2, the power flux through this opening is $\sim 5 \times 10^3$ TW. Of the total initially available energy, equal to about 30 MJ, only one-third, that is, ~ 10 MJ, is available as blackbody radiation. The remaining ~ 20 MJ goes into work expanding the target chamber ~ 3-fold in its diameter, but the available ~ 10 MJ is more than enough to implode a high-density, high-yield thermonuclear target. The wavelength of blackbody radiation at $T_b' \simeq 10^7$ K is sufficiently short to ensure good coupling to the target for densities up to $\sim 10^4$ times solid densities.

Finally, we shall suggest a configuration in which the ignition is performed by detonation wave lenses and shape charges using conventional explosives. This configuration is shown in Figure 40. Here, the magnetic target chamber is surrounded by several shape-charge lenses, which implode and compress the magnetized low-density DT booster target I inside the chamber. As before, this booster target is preheated by a short laser pulse of modest energy. An ordinary shape charge can produce a jet velocity of ~ 10 km/sec.[29] However, the multiple shape-charge arrangement as shown in Figure 40 produces a quasispherical implosion in addition. If this implosion is approximated by a Rayleigh-type cavity implosion,[7] the implosion velocity would rise as $\ell^{-3/2}$.

In reality, the velocity will rise less quickly, because of the rising plasma pressure in the imploding cavity. But even if we take this slowing-down mechanism into account, it appears very likely that a velocity at least twice as large as already attainable with an ordinary shape charge should be reached, which would be the required 20 km/sec. Furthermore, rather than letting the blackbody radiation emerging from the booster target I simply hit a spherical second-stage target, one can use it to ignite a thermonuclear detonation wave, as shown in Figure 40.

The trigger energy in this last concept requires perhaps not more than ~ 100 kg of TNT, and the proposed configuration, if

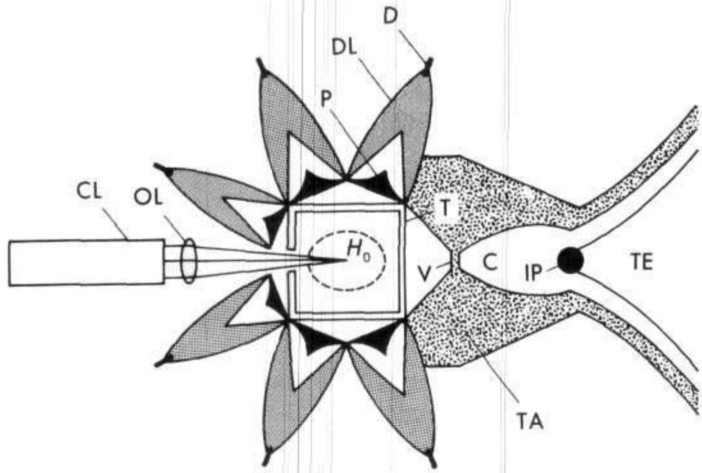

Figure 40. Two-stage magnetic booster target configuration wherein the thermonuclear fuel is ignited by chemical detonation wave lenses. CL is a chemical laser; OL an optical lens; T the booster target chamber; DL are detonation lenses; D are detonators; P metallic plates deflecting the shape-charge jets from DL; V the vertex position; C the cavity; TA the tamp; IP the ignition point for autocatalytic thermonuclear detonation propagating into a horn filled with thermonuclear explosive TE.

technically feasible, would therefore be of great interest for nuclear rocket propulsion using the simple pusher-plate concept. The trigger mechanism would already be incorporated here in the microbomb, and no beam-generating machine would have to be carried on board the spacecraft. The yield of such a chemically ignited microbomb could be chosen much larger than for a typical microexplosion, but still smaller than for a true fission explosion as proposed for the Orion concept. However, the explosion would still be so large that a thermonuclear shape charge, as described in the previous chapter, should be used to increase the momentum transfer.

Another recent development of considerable interest is the result that the effective value of $\langle \sigma v \rangle$ for the $H^{11}B$ thermonuclear reaction is increased in a reacting plasma over the value obtained by averaging over a Maxwellian velocity distribution, which increases the chance for this thermonuclear reaction to occur.[37] This is because the fast thermonuclear reaction products increase the high-velocity tail of the ion distribution over what one would obtain in a pure Maxwellian distribution. In addition, there is a temperature shift between the electrons and ions. In particular, it has been found that for a temperature of $\sim 3 \times 10^9$ K the thermonuclear reaction rate in a $H^{11}B$ plasma should just exceed the bremsstrahlung losses with a value for $\langle \sigma v \rangle$ given by

$$\langle \sigma v \rangle \simeq 7.4 \times 10^{-16} \text{ cm}^3/\text{sec}. \tag{154}$$

We shall test this result against our formula for the propagation of a thermonuclear detonation wave developed in Chapter 10, to see if a thermonuclear detonation in $H^{11}B$ is possible. This is even more chancy than just a simple burn.

The minimum value that $\langle \sigma v \rangle$ must at least assume to make a detonation wave possible was given by Eq. (61). For the $H^{11}B$ plasma we must put

$$M = (1/2)(M_H + M_B),$$

where $M_B = 11 M_H$ is the mass of the boron atom. Hence $M = 6M_H$. Furthermore, $f = 1$ and

$$\varepsilon_0 = 2.93 \text{ MeV} = 4.7 \times 10^{-6} \text{ erg}.$$

The effective value of $1/Z$ in the range equation is here $1/3$. With this value and the energy of $\varepsilon_0 = 2.93$ MeV for the $H^{11}B$ reaction we find for the range constant

$$a = 0.75 \times 10^{34} \text{ cm}^{-2}\text{-erg}^{-3/2}.$$

By inserting this into Eq. (61), we then find

$$\langle \sigma v \rangle_{\min} \simeq 1.05 \times 10^{-15} \text{ cm}^3/\text{sec}. \tag{155}$$

The value given by Eq. (154) is only slightly smaller than the minimum required value given by Eq. (155), and taking possible inaccuracies into account, it therefore cannot be excluded that a detonation wave in $H^{11}B$ is marginally possible.

Even then, however, this is only true if there is no energy loss by bremsstrahlung, assuming the correctness of the result quoted above. Therefore, to make it possible at all, the bremsstrahlung losses must somehow be reduced. Such a reduction takes place, in fact, at very high plasma densities because of quantum-mechanical effects. The above-quoted calculation, however, was made at a rather low plasma density, at which a large temperature shift between electrons and ions exists. It is thus not clear if these results also apply to the high-density case. The possibility of a detonation wave in $H^{11}B$ therefore seems not to be excluded but requires a deeper analysis.

CHAPTER EIGHTEEN

Ignition by A Fission Fizzle

In this final chapter we present a concept that, if technically feasible, demonstrates better than any other concept using the energy released by a fission chain reaction to trigger thermonuclear ignition the close connection and significance these concepts have both for weapons and commercial power production. The result obtained from thermonuclear microexplosion research was that a power of $\simeq 10^{14}$ W focused to less than ~ 1 cm^2 is sufficient to compress ablatively and ignite a thermonuclear microexplosion target. Assuming a power density of

$$P = 10^{15}\ \mathrm{W/cm^2} = 10^{22}\ \mathrm{erg/cm^2\text{-}sec},$$

just to be on the safe side, we may ask what the temperature of a blackbody need be to radiate at this power. The answer, obtained from the Stefan-Boltzmann law $P = \sigma T^4$, is that $T \simeq 3.6 \times 10^6$ K.

The energy density of the radiation at this temperature is given by

$$\varepsilon_r = (4\sigma/c)T^4 = 1.3 \times 10^{12}\ \mathrm{erg/cm^3},$$

and is approximately 1 to 2 orders of magnitude larger than the energy density of high explosives, although about 5 orders smaller than the energy density reached in a fission explosion with a burnup of approximately 10%, which reaches a temperature of about 6.6×10^7 K. Hence, to reach a temperature of only $\sim 3.6 \times 10^6$ K could be accomplished with a very low yield fission explosion, called a *fission fizzle*.

At a temperature of 3.6×10^6 K the kinetic particle energy density $\varepsilon_k = (f/2)nkT$ exceeds the radiative energy density ε_r. At this temperature the plasma is roughly five times ionized and hence $f = 18$. For uranium at solid-state densities, $n = 5 \times 10^{22}$ cm^{-3}, and hence $\varepsilon_k \simeq 6 \times 10^{13}$ erg/cm^3. By comparison, the energy density in a fully developed fission explosion with a temperature of 6.6×10^7 K is determined by blackbody radiation and is $\varepsilon_r \simeq 1.4 \times 10^{17}$ erg/cm^3. It follows that the energy density in a fission fizzle reaching a temperature of 3.6×10^6 K is reduced by the factor 2.3×10^5 compared to a fully developed fission explosion. Therefore, assuming that a fully developed fission explosion sets free an energy equivalent to 2×10^4 tons of TNT, the energy set free in the fission fizzle would be equivalent to ~ 100 kg of TNT.

Since the linear dimensions of a critical fission explosive are just a few centimeters, the same applies to a fission fizzle. Because the fission fizzle will reach a temperature of 3.6×10^6 K, the typical wavelength of this radiation, given by $h\nu = kT$, is $\lambda = c/\nu = 1.3 \times 10^{-7}$ cm. This wavelength is much larger than the typical wavelength at the much higher temperature of a fully developed fission explosion. As a consequence, radiation at such a long wavelength, corresponding to that of soft X-rays, can be reflected from a metal surface if its frequency is matched to the K_α line of the discrete X-ray spectrum.

The frequency, and hence wavelength, of the K_α line is determined by Moseley's law:

$$\lambda = (1.2 \times 10^{-5})(Z - 1)^{-2} \text{ cm}, \tag{156}$$

where Z is the atomic number of the substance in question. One then finds that for $\lambda = 1.3 \times 10^{-7}$ cm resonance reflection would occur for a substance with $Z \sim 11$, that is, for sodium. The spectrum of the blackbody radiation, however, at 3.6×10^6 K is a Planck distribution with a large width covering longer but also substantially shorter wavelengths. To make a resonance mirror covering the many frequencies of such a spectrum one may use many thin layers of substances going from light elements with $Z \sim 10$ up to heavy elements with $Z \sim 100$. According to the Moseley law this would cover a broad X-ray spectrum.

The possibility to reflect the soft X-rays from a fission fizzle by such a multilayered resonance mirror can be used to refocus the radiation onto a thermonuclear microexplosion target or a thermonuclear fuse. The first case is of interest if one wishes to apply this concept to commercial power production and the second case for a neutron bomb with a greatly reduced energy output of the fission trigger. In contrast to the ignition by a focused shock wave using the Prandtl-Meyer ellipsoid, however, the ignition must be done here by ablative implosion, and the geometric shape of the Prandtl-Meyer ellipsoid must be replaced by the geometric shape of a real ellipsoid. The need for ablative implosion is a consequence of the fact that, because of their zero rest mass, photons may carry large quantities of energy but carry much less momentum than a fission-induced shock wave. It therefore follows that the momentum transfer to the target must be accomplished by ablative rather than by direct implosion.

To make a fission fizzle is easy. All that is required is to make the fission explosive less critical than is required for a fully developed explosion.

For a weapons application of this concept, the fission trigger can be assembled in the conventional way by high explosives. If commercial power production is the goal, the fission trigger could be assembled by hypervelocity guns accelerating two subcritical masses at each other at high speed. If the mutual collision velocity is a few million centimeters per second, which could be realized by electromagnetic guns or even isentropic light gas guns, then the fission material could be compressed to substantially higher than solid-state densities.

If, for example, an approximately 5-fold compression could be achieved, the critical mass would be reduced 25-fold. The yield of such a compressed fission explosion would be reduced by the same factor, from ~ 100 kg of TNT-equivalent down to ~ 4 kg. This is an

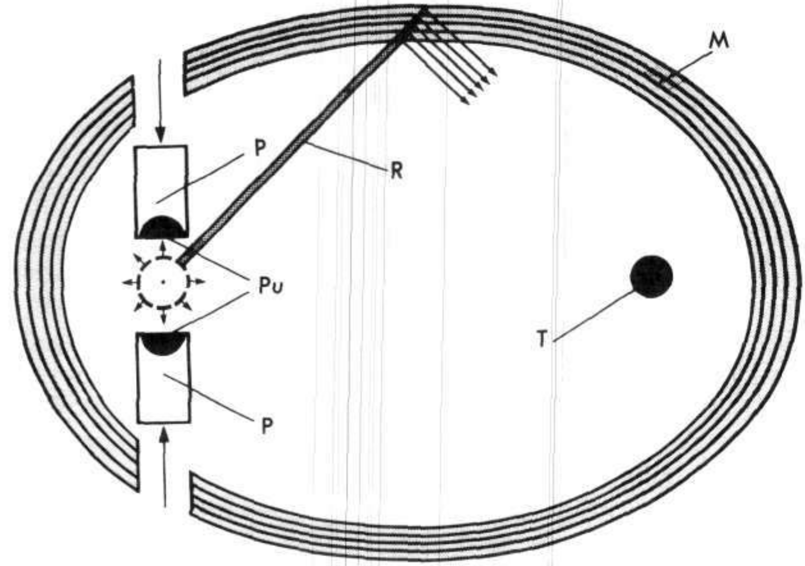

Figure 41. In this configuration two hypervelocity projectiles P both assemble the subcritical plutonium hemispheres Pu and serve as neutron reflectors. M is a multilayered K_α-line ellipsoidal resonance mirror and T is the thermonuclear microexplosion target. R is a photon ray from the fission fizzle, where the different ray frequencies are reflected from the different layers of M.

energy not very much larger than what is contemplated for beam-ignited microexplosions.

To make the critical mass as small as possible one may place two subcritical masses inside tamps serving as neutron reflectors, which are shot onto each other head on, with the position of the collision in one focus of the soft X-ray reflecting ellipsoid, and the microexplosion target in the second focus. This configuration is shown in Figure 41. The two semispherical subcritical masses are placed as shown inside the two heavy slugs, which have been accelerated to a few times 10^6 cm/sec. If their combined mass is about 1 kg, then the kinetic projectile energy would be $\sim 10^{15}$ erg. If the total yield, therefore, of the ensuing thermonuclear microexplosion is $\sim 10^{17}$ erg, the gain would be large enough to make this a worthwhile energy-generating concept but not so large that it could not be confined in a container serving as a reactor vessel.

In this regard, the concept is distinctly different from the Pacer concept, where the yield is much larger, requiring the explosion to be confined in an underground cavity.

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About the Author

Dr. Friedwardt Winterberg, a pioneer in inertial-confinement fusion, is considered the father of impact fusion for his early work on thermonuclear ignition by hypervelocity impact. Now a research professor at the Desert Research Institute of the University of Nevada System, he has long been at the forefront of research on the implementation of nuclear energy for spaceflight. His concept of a rocket engine propelled by a sequence of thermonuclear microexplosions, in fact, was the inspiration for the Project Daedalus starship study of the British Interplanetary Society. He received the 1979 Hermann Oberth gold medal of the Hermann Oberth—Wernher von Braun International Space Flight Foundation for his work on thermonuclear propulsion.

Born in Berlin, Germany in 1929, Dr. Winterberg became fascinated with spaceflight as a youth, by reading the writings of spaceflight pioneer Hermann Oberth. He taught himself calculus at the age of 14 in order to understand Oberth's mathematical theories. After receiving his doctorate in physics under Werner Heisenberg in 1955, he helped design the research reactor of the Society to Advance Nuclear Energy for Naval Propulsion in Hamburg, West Germany.

Dr. Winterberg's published work includes more than 130 scientific papers and articles in many books. In 1963 he published the first proposal for the ignition of a thermonuclear microexplosion by a beam of microparticles accelerated in conventional particle accelerators. And in 1967 he began a series of papers on the use of intense electron and ion beams for fissionless thermonuclear ignition, culminating in a 1969 paper describing the magnetically insulated diode as a means of producing ultraintense ion beams.

"I have never believed that it can serve any purpose to keep physics facts that can be obtained from any college textbook a 'secret.'"

—Friedwardt Winterberg

The Physical Principles Of Thermonuclear Explosive Devices

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"Our policy is not to comment on things like this for publication, because if it does contain something classified, then we are revealing what is classified."

—Robert Duff
Director of Classification, DOE

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