(NASA-TM-82083) GRAND UNIFICATION AND POSSIBLE MATTER-ANTIMATTER DOMAIN STRUCTURE IN THE UNIVERSE (NASA) 43 p HC A03/MF A01 CSCL 03B

N81-23986

Unclas
G3/90 23479

NASA

Technical Memorandum 82083

Grand Unification and Possible Matter–Antimatter Domain Structure in the Universe

F. W. Stecker

JANUARY 1981

National Aeronautics and
Space Administration

Goddard Space Flight Center
Greenbelt, Maryland 20771

GRAND UNIFICATION AND POSSIBLE MATTER-ANTIMATTER DOMAIN STRUCTURE IN THE UNIVERSE

F. W. Stecker
NASA/Goddard Space Flight Center
Laboratory for High Energy Astrophysics
Greenbelt, MD 20771, U. S. A.

Invited Paper Presented at the
Tenth Texas Symposium
on Relativistic Astrophysics

To be published in the Proceedings

Our universe is a sorry little affair
unless it has in it something
for every age to investigate.
Seneca

INTRODUCTION

Are there cosmologically significant amounts of antimatter, even whole galaxies of antimatter, elsewhere in the universe? Some would argue that the conservative answer to this question is no. However, the truly conservative answer to this question at this point in time is that we do not know. We cannot tell from present observations that a galaxy in the Hercules cluster of galaxies, for example, is made of matter rather than antimatter. An alternative cosmology to the present orthodox matter-centric viewpoint can be constructed, based on the modern gauge theory paradigm. This viable alternative can be tested by observation and has potential advantages over the orthodox picture in explaining presently existing astrophysical data. Although one cannot expound an entire cosmology in one short presentation, I will try to give you some basic ideas and highlights of such a cosmology in which matter and antimatter play an equal role in the universe.

UNIFIED GAUGE FIELD THEORIES

We first review the basic theoretical concepts. We begin with a powerful formulation of localized quantized field theory developed by Schwinger,[1,2]. This formulation starts by considering states described by sets of commuting operators ζ on spacelike space-time surface σ and infinitesimal changes in the transformation function $< \zeta_1' \, \sigma_1 \mid \zeta_2'' \, \sigma_2 >$ and the unitary operators U_{12} which describe the evolution of the system from σ_2 to σ_1 ,

$$\delta<\zeta_1' \, \sigma_1 | \zeta_2'' \, \sigma_2> = <\zeta_2' \, \sigma_2 | \delta U_{12}^{-1} | \zeta_2'' \, \sigma_2> \tag{1}$$

where

$$\delta U_{12}^{-1} = i U_{12}^{-1} \, \delta W_{12} \tag{2}$$

and δW_{12} is an infinitesimal Hermitian operator (we set $\hbar = c = 1$).

The operator W_{12} has the general form

$$W_{12} = \int_{\sigma_2}^{\sigma_1} (dx) \; L[x] \tag{3}$$

where the Lagrangian density $L[x]$ is a function of the quantum fields $\phi^{(a)}(x)$ and their derivatives $\phi^{(a)}_{,\mu}(x)$. If the parameters of the system are unchanged, the variation can be defined in terms of generating operators at the endpoints

$$\delta W_{12} = \delta \int_{\sigma_2}^{\sigma_1} (dx) L[x] = G(\sigma_1) - G(\sigma_2) \tag{4}$$

The action integral is unchanged by infinitesimal variations inside the region bounded by σ_1 and σ_2 . This then leads to the equations of motion for the fields

$$\partial L / \partial \phi^{(a)} = \partial_\mu (\partial L / \partial \phi^{(a)}_\mu) \equiv \partial_\mu \pi^{\mu(a)} \tag{5}$$

with the generating functions given by

$$G(\sigma) = \int_\sigma d\sigma_\mu [\Pi^{\mu(a)} \delta\phi^{(a)} + L\delta x^\mu]. \tag{6}$$

The resulting unitary transformation operators are, in infinitesimal form

$$\begin{aligned} U &= 1 - iG \\ U^{-1} &= 1 + iG \end{aligned} \tag{7}$$

The various fields describing the forces of nature can be represented by the symmetries they possess in terms of the transformations of the quantum systems they produce which leave the Lagrangian invariant. The generators can be related to generalized charges. For example, in the case of QED, conservation of charge can be derived from the symmetry with respect to the phase transformation (called, for historical reasons, a global gauge transformation)

$$\delta\phi^{(a)} = -ieq^{(a)}\phi^{(a)}\delta\lambda \tag{8}$$

where $q^{(a)} = \pm 1, 0$ represents the sign of the charged field $\phi^{(a)}$.

The generator is then

$$\begin{aligned} G\ \delta\lambda &= -i\int_\sigma d\sigma_\mu \Pi^{\mu(a)} q^{(a)} \phi^{(a)} \delta\lambda \\ &= Q(\sigma)\delta\lambda \end{aligned} \tag{9}$$

where $Q(\sigma)$ is the total charge

$$Q(\sigma) = \oint d\sigma_\mu j^\mu \tag{10}$$

with

$$j^\mu = -ie\pi^{\mu(a)}q^{(a)}\phi^{(a)} \tag{11}$$

being the current.

The symmetry

$$\delta W_{12} = Q(\sigma_1) - Q(\sigma_2) = 0 \tag{12}$$

expresses the conservation of charge. The symmetry group of these QED gauge transformations is the one-parameter unitary group U(1).

The electromagnetic field $A_\mu(x)$ is introduced by requiring invariance under local phase (gauge) transformations $\lambda(x)$ and requiring that the derivatives of the charged fields transform in the same way as the fields themselves

$$\begin{aligned} \phi &\to \phi e^{-i\lambda(x)} \\ \phi^\dagger &\to \phi^\dagger e^{i\lambda(x)} \end{aligned} \tag{13}$$

This leads to the introduction of the gauge covariant derivative

$$\begin{aligned} D_\mu\phi &\equiv \phi,_\mu - ieA_\mu \\ D_\mu\phi^\dagger &\equiv \phi^\dagger,_\mu + ieA_\mu \end{aligned} \tag{14}$$

and the gauge transformation law for the electromagnetic field

$$A_\mu \to A_\mu - \frac{1}{e}\partial_\mu \lambda(x) \tag{15}$$

More complex gauge fields can be constructed from generators which preserve the form of the Lagrangian under more complex symmetry groups involving larger numbers of parameters, i.e., group spaces of higher dimension. These generators $T^a{}_{(i)b}$ obey Lie algebras, i.e.,

$$[T^a_{(i)b}, T^b_{(j)c}] = c_{ijk} T^a_{(k)c} \tag{16}$$

where c_{ijk} are the structure constants which define the Lie algebra. An example of importance to the unified field theory of electromagnetic and weak interactions, is the gauge group SU(2), the unitary group whose fundamental representation consists of two-dimensional (traceless) matrices of determinant + 1. For this group, the generators can be represented by the familiar Pauli spin matrices $T_i = \sigma_i/2$, so that

$$[T_i, T_j] = i\varepsilon_{ijk} T_k \tag{17}$$

with ε_{ijk} being the totally antisymmetric unit tensor of rank three.

The unitary transformations are then given by

$$U = e^{-i\vec{T}.\vec{\lambda}} \tag{18}$$

where the vector symbols and dot product refer to the abstract three parameter group space. This group of transformations is locally isomorphic to the group of rotations in three dimensional space.

The demand for local gauge invariance under SU(2) transformations, as in the case of QED, requires the introduction of a new gauge field B_μ and coupling constant g (instead of e) such that

$$D_\mu\phi = (\partial_\mu - ig\vec{T}\cdot\vec{B}_\mu)\phi \tag{19}$$

with the new fields transforming as

$$(\vec{T}\cdot\vec{B}_\mu) \rightarrow U(\vec{T}\cdot\vec{B}_\mu)U^{-1} - \frac{1}{g}(\partial_\mu U)U^{-1} \tag{20}$$

In the electroweak theory of Glashow, Weinberg and Salam (GWS)[3-5] the gauge group is the product SU(2) x U(1), where the SU(2) group, by analogy with spin, is called the weak isospin group and the U(1) group, by analogy with electromagnetism $U(1)_{EM}$ has a generator Y called "weak hypercharge". In the quantum gauge theory of strong interactions, QCD (quantum chromodynamics), the generalized charges are referred to as colors. In GWS, the four transformation parameters result in the four gauge bosons γ (photon), $W^{\pm}$, Z^0 the heavy bosons which carry the weak charged and neutral currents. For an SU(n) theory, there are n^2-1 free parameters. In QCD or color SU(3) there are $3^2-1 = 8$ gluons which carry the force. In the simplest grand unified theory[6], viz. SU(5), there are a total of 24 gauge bosons, γ, $W^{\pm}$, Z^0, the 8 gluons and 12 new superheavy bosons, $X^{4/3}$, $Y^{1/3}$ of all three colors together with their antiparticles. It is these bosons which are responsible for the "leptoquark" force which can transform quarks into leptons and <u>vice</u>

versa, violating baryon number and producing an excess of matter (or antimatter) out of the primordial thermal radiation.

SPONTANEOUS SYMMETRY BREAKING

Of course, in our world of "low temperature" physics much of the symmetry of the unified theories is badly broken, leaving only $SU(3)_c$ and $U(1)_{EM}$.

This is reflected in the large masses of all of the gauge bosons except γ and the gluons (which are massless) and the corresponding weakness of the weak and leptoquark interactions. The broken symmetries are incorporated into the theory by keeping the full symmetry in the Lagrangian but allowing the gauge bosons to obtain their masses "spontaneously" as the result of introducing new scalar (or "Higgs") fields which have a non-zero vacuum expectation value. One big advantage of the Higgs mechanism is that it allows the construction of a theory which is renormalizable, i.e., for which the calculations of observables give finite results.

The way the Higgs mechanism works is as follows. Consider for example, a real scalar field whose contribution to the Lagrangian takes the form

$$L_s = \frac{1}{2} (\partial_\mu \phi)(\partial^\mu \phi) - V(\phi) \tag{21}$$

where the potential term is an even function $V(\phi) = V(-\phi)$. Consider, e.g., a potential of the form

$$V(\phi) = \frac{1}{2} \mu^2 \phi^2 + \frac{1}{4} \lambda \phi^4 \tag{22}$$

where $\lambda > 0$ so that the energy is bounded from below. In the case $\mu^2 < 0$, $V(\phi)$ has minima at

$$\langle\phi\rangle = \pm \frac{(-\mu^2)^{1/2}}{\lambda} \equiv v \tag{23}$$

which gives, by definition, the vacuum expectation value for ϕ. The Lagrangian (21) gives the equation of motion

$$\partial_\mu \partial^\mu \phi + \partial V/\partial\phi = 0 \tag{24}$$

For small excitations of the field near v, i.e., $\phi = v + \delta\phi$

$$[\partial_\mu \partial^\mu + (\partial^2 V/\partial\phi^2)_v]\delta\phi = 0 \tag{25}$$

which is the Klein-Gordon equation for a boson of mass

$$m^2 = \frac{\partial^2 V}{\partial\phi^2} = 2\mu^2 = v\ (2\lambda)^{1/2} \tag{26}$$

If ϕ couples to fermions with a coupling of the Yukawa form

$$L_Y = f\phi\bar{\psi}\psi \tag{27}$$

the Higgs field ϕ gives fermions masses of order fv. Thus, without explicitly introducing masses into the Lagrangian, the Higgs mechanism produces masses in the theory which are proportional to v.

$$\begin{aligned} &m_f \sim fv \\ &m_\phi \sim hv,\ h = (2\lambda)^{1/2} \text{ is the Higgs self-coupling constant} \\ &m_B \sim gv,\ g \text{ is the gauge field coupling constant} \end{aligned} \tag{28}$$

For a more detailed discussion of this mechanism of spontaneous symmetry breaking, see, e.g., references (7, 8).

So far we have spoken of vacuum expectation values $\langle\phi\rangle$ of the scalar fields in a zero-temperature theory with the symmetries of the Lagrangian broken by the Higgs mechanism. The cosmological implications come in when we consider what happens as T increases to temperatures $T \gtrsim \langle\phi\rangle$. In this case some, or all, of the symmetry in the theory may be restored[9,10], i.e., $\langle\phi\rangle_T \to 0$ for $T > T_c$ (some critical temperature) and the corresponding masses go to zero. A direct analogy can be made here with the theory of superconductivity, where the Cooper pairs play the role of Higgs particles and the photon acquires an effective mass for $T < T_c$ which disappears at $T > T_c$ (the Meissner effect). In the finite temperature case, the Higgs fields have a thermal distribution of excitations and the vacuum expectation value is replaced by the operator Gibbs average. In the simple case of equation (22) the resulting potential acquires an effective quadratic term

$$-\mu_{eff}^2(T) \approx -\mu^2 + \sigma T^2 \qquad (29)$$

and critical temperature $T_c = |\mu|\sigma^{-1/2}$ where $\mu_{eff} = 0$ in the case $\sigma > 0$. In general, σ is a function of the coupling constants of the model.

BARYON PRODUCTION IN THE EARLY UNIVERSE

We now have most of the conceptual machinery in place for a discussion of the evolution of the early universe. To set the stage, we also require an idea of some of the time and temperature scales involved.

At the temperatures of interest here, if the dynamics of the universe is dominated by the energy density of the thermal radiation, the temperature of the universe $T \propto t^{-1/2}$, where t is the age of the universe. (The exception is when the expansion is dominated by the energy density of the Higgs field. That case will be discussed later.) More precisely[11], the expansion rate ($k = 1$)

$$\Gamma_H = \frac{\dot{R}}{R} = 1.66\ T^2\ N^{1/2} M_p^{-1}\ . \tag{30}$$

where the Planck mass, $M_p \equiv G^{-1/2} = 1.2 \times 10^{19}$ GeV and N is the number of helicity states of all particle species in the thermal radiation.

The critical temperature for symmetry breaking at the electroweak level, i.e., $SU(2)_L \times U(1)_Y \rightarrow U(1)_{EM}$ is usually considered to be of order $G_F^{-1/2} \approx 300$ GeV, but as one can see from equation (29), T_c depends on the specific parameters of the theory. In fact, it is possible[12] that $T_c \gg G_F^{-1/2}$ as we will discuss later. The characteristic temperature scale for grand unification is given by the energy scale at which the coupling constants for the electroweak gauge groups and strong gauge group become comparable[13]. This is given from renormalization group theory[14] to be of order $\sim 10^{15}$ GeV, above which for the SU(5) theory only one coupling constant, associated with this simple gauge group, exists.

The proton lifetime against leptoquark decay $\tau_p \propto m_X^4$. The experimental lower limit[15] on τ_p gives a lower limit on m_X, viz $m_X \gtrsim 10^{14}$ GeV, consistent with the value obtained from considering the energy (temperature) dependence of this coupling constants given by renormalization group theory. Thus, it is at this temperature level, $T \sim m_X$, that baryon generation processes will be of importance.

A scenario for baryon production through the decay of these superheavy gauge and Higgs bosons has been given by Weinberg[11]. He considered the decay of these "X-bosons" into two channels $X \to q l$ and $X \to \bar{q}\bar{q}$ with branching ratios r and $1-r$ respectively, together with the antiparticle decays $\bar{X} \to \bar{q}\bar{l}$ and $\bar{X} \to qq$ with branching ratios $\bar{r}$ and $1-\bar{r}$.

The three conditions for production of a baryon excess in the early universe are (1) baryon (quark) nonconservation, (2) nonconservation of C (charge conjugation) and CP (C x parity) and (3) thermal disequilibrium. We have seen that grand unification supplies condition (1). The expansion of the universe supplies condition (3). The need for condition (2) can clearly be seen in the Weinberg scenario. The baryon number generated in the X and $\bar{X}$ decays is

$$\Delta B = \frac{1}{2}\left[\frac{1}{3}r - \frac{2}{3}(1-r) - \frac{1}{3}\bar{r} + \frac{2}{3}(1-\bar{r})\right] = \frac{1}{2}(r-\bar{r}) \ . \qquad (31)$$

If CP is conserved, $r = \bar{r}$ and no baryon excess is generated. It should also be noted that the sign of the CP violation determines the sign of $r - \bar{r}$ and therefore the sign of ΔB. Thus, whether a baryon excess or an antibaryon excess is created by this process depends on the sign of the CP violation parameter.

The rates for leptoquark interactions and X-boson decay are given by

$$\Gamma_I \approx \frac{\alpha_X^2 T^5 N}{(T^2 + m_X^2)} \qquad (32)$$

and

$$\Gamma_D \approx \frac{\alpha_X m_X^2 N}{(T^2 + m_X^2)^{1/2}} \qquad (33)$$

where $\alpha_X \equiv g_X^2/4\pi$. These rates are to be compared with the Hubble expansion rate Γ_H given by equation (30). For $\Gamma_D > \Gamma_H$ the X-bosons decay. If that condition is met when $T \lesssim m_X$, inverse decay is blocked by the Boltzmann factor $e^{-m_X/T}$ and the X-bosons decay freely. The result is a baryon-to-photon ratio

$$\eta \equiv \frac{n_B}{n_\gamma} \approx 0.28 \frac{N_X}{N} \Delta B \tag{34}$$

where N_X/N is $\sim 10^{-2}$ to $\sim 10^{-1}$ and ΔB is given by equation (31).

From astrophysical observations, one obtains $10^{-10} \lesssim \eta \lesssim 10^{-8}$. Nanopoulos and Weinberg[16] conclude that the decays of the superheavy scalar bosons are most relevant for cosmological baryon production. They estimate that $10^{-8}\,\varepsilon \lesssim \Delta B \lesssim 10^{-6}\varepsilon$. The parameter ε, is a parameter characterizing the strength of CP violation, which Nanopoulos and Weinberg consider to be in the range $\sim 10^{-2}$ to 1, the sign being undetermined. All in all, these authors estimate $10^{-9} \lesssim |\eta| \lesssim 10^{-3}$ immediately after the era of baryon production. (Numerous other authors have also worked on the problem of estimating η . Their work cannot be reviewed here for lack of space.)

CP VIOLATION AND COSMOLOGICAL IMPLICATIONS

It follows from the discussions of the previous section that the sign of the baryon number excess, which determines whether matter or antimatter is created, depends on the sign of the CP violation parameter. In the scenarios usually considered, CP violation of one sign only is put into the model explicitly in the Lagrangian via complex Yukawa couplings between the fermions and scalar fields, i.e., L_Y of the form in equation (27) with f complex, or in complex self couplings of the scalar fields, i.e., λ complex in the potential term $\frac{1}{4}\lambda\phi^4$. However, it is also possible for the CP violation to arise from the mechanism of spontaneous symmetry breaking. Such a mechanism has been proposed to explain the smallness of the CP violation implied by the small electric dipole moment of the neutron[17]. Furthermore, if CP is broken spontaneously, the amount of CP violation is finite and calculable, whereas the presently popular baryon production scenarios invoke a "hard" CP violation, leading to infinite renormalizations of the CP parameter which thus become incalculable undetermined free parameters. With spontaneous CP violation the Lagrangian is CP invariant (f and λ real), but the scalar fields themselves take on complex vacuum expectation values which produce the CP violation. In this second case, the CP violation is not put in by hand *ad hoc*. We start out with a completely CP symmetric theory with the symmetry of the Lagrangian reflected in the state of the universe at the highest temperatures. This being the case, owing to the finite age of the universe t_u, regions separated by distances greater than $\sim ct_u$ are not, and never were during the course of the expansion, in causal contact. Thus, if spontaneous symmetry breaking of CP occurred at a time t_{CP}, it would have occurred independently and with random signs in regions separated by distances larger than $\sim ct_{CP}$. We will call these "seed domains" and consider how they arise and scenarios for their subsequent growth and evolution. This domain structure is not unlike the domain structure generated when a piece of ferromagnetic material cools without the presence of an external magnetic field. In that case, each of the domains contain atoms having their magnetic moments aligned in a given direction. On the average, there will be no preferred direction on a global scale. Analogously, one may expect that spontaneous symmetry breaking processes in the early big-bang will most likely break baryon symmetry in localized regions of the universe but will preserve the overall global matter-antimatter symmetry of the initial state. Thus, present ideas of unified gauge theories with spontaneous CP symmetry breaking can lead naturally to an overall baryon-symmetric cosmology[18]. Senjanovic and Stecker[19] have considered mechanisms of spontaneous soft CP violation within the context of the specific grand unified theories involving the SU(5) and SO(10) gauge groups. They discuss two distinct classes of models, viz., those with only one source of CP violation independent of temperature for SU(5) and those in which the CP violation at the super-heavy mass scale for SO(10) has nothing to do with the observed CP violation at "low temperatures" in the K^0-$\bar{K}^0$ system. They conclude that independently of the particular model, the domain picture of the universe emerges naturally in theories of soft CP violation.

In the minimal SU(5) model with only one Higgs multiplet, CP violation has to be put in "by hand" in the Lagrangian in the form of complex Yukawa couplings, since the vacuum expectation value of the Higgs field can always be redefined to be real by means of a gauge transformation. Choosing such a hard CP violation, yields a baryon-photon ratio which is unacceptably small compared to that determined by astrophysical observation[20,21]. It is therefore necessary for consistency to increase the number of 5-dimensional Higgs multiplets. Increasing this number to three results in a realistic grand unified theory based on SU(5) which allows for soft CP violation at high temperatures. Two of the Higgs fields acquire vacuum expectation values with a relative phase which cannot be transformed away, since they carry the same U(1) quantum number. Senjanovic and Stecker consider a Higgs sector with three 5-dimensional multiplets with the following pattern of symmetry breaking at the electroweak level ($T \lesssim 300$ GeV):

$$\langle\chi\rangle = \begin{pmatrix} 0 \\ 0 \\ 0 \\ 0 \\ \rho \end{pmatrix}, \quad \langle\phi_1\rangle = \begin{pmatrix} 0 \\ 0 \\ 0 \\ 0 \\ v_1 \end{pmatrix}, \quad \langle\phi_2\rangle = \begin{pmatrix} 0 \\ 0 \\ 0 \\ 0 \\ v_2 e^{i\theta} \end{pmatrix} \tag{35}$$

It can be shown that at $T \gg 300$ GeV the symmetry will still be broken, with $\langle\chi\rangle = 0$ but with $\langle\phi_1\rangle$ and $\langle\phi_2\rangle$ nonvanishing. This follows from having the coefficient μ_{eff} of the quadratic terms in the Lagrangian for $V(\phi_1,)$ and $V(\phi_2)$ of the form given by equation (28) with $\sigma < 0$ at $T \sim 300$ GeV. Then, noting that $\sigma = \sigma\,(T)$ is a slowly varying function of T, owing to the logarithmic temperature dependence of the coupling "constants" (obtained from renormalization group theory), it has been found that in some cases[12] $\sigma(T)$ becomes positive for $T_c \gtrsim m_X$.

Thus, spontaneous soft CP breaking at the electroweak level can be effective even at the grand unification temperatures when baryons are produced.

The Higgs potential as a function of θ can, in general, be written as

$$V(\theta) = A + B \cos\theta + C \cos 2\theta \tag{36}$$

where A, B, and C are independent of θ. Obviously, for an appropriate range of parameters, the mimimum of the Higgs potential lies at $\theta_0 \neq 0$ with $\cos\theta_0 = -B/4C$, so that <u>we always have two solutions</u>, θ_0 and $-\theta_0$.

The value of $r-\bar{r}$ is proportional to $\sin\theta$. Now since $\theta = \pm\,\theta_0$ (the solution of the minimization of the potential), one obtains from equation (31)

$$n_B/n_\gamma \propto \pm \sin\theta_0 \tag{37}$$

The renormalization group analysis suggests the possibility (intuitively expected) that at even higher temperatures $T > m_X \approx 10^{15}$ GeV, the symmetry was unbroken. Then as the temperatu decreased below the mass scale of the superheavy gauge bosons, we expect that separate domains were generated with θ_0 and $-\theta_0$ phases (by the analogy with ferromagnetic systems.) Therefore from equation (37) it is obvious that one is bound to expect domains with matter and antimatter excesses in the universe. Thus, a realistic theory of soft CP violation leads to the domain picture with matter and antimatter being randomly distributed throughout the universe. Senjanović and Stecker also consider the development of domains at $T \sim m_X$ in a recently suggested model[22], based on SO(10) grand unified theory (see ref. 19).

DOMAIN GROWTH AND HORIZON GROWTH

The above discussion suggests that the initial domains were formed at a time when the temperature of the universe was comparable to the masses of the superheavy gauge or Higgs bosons involved in the symmetry breaking. The initial domains could then have acted as nuclei for triggering growth to much larger sized regions.

One particularly promising mechanism for domain growth to an astronomically relevant scale has been suggested by Sato quite recently[23]. This mechanism depends on the fact that the expansion of the universe can be drastically altered from the standard radiation-dominated relationship if the energy density of the Higgs field is larger than that of the thermal radiation.

To see how the energy density of the Higgs fields affects the Einstein equations, note that the Higgs fields define the vacuum state of the universe. Thus, given the Einstein equations

$$R_{\mu\nu} - \frac{1}{2} Rg_{\mu\nu} = 8\pi T_{\mu\nu} \;, \tag{38}$$

with the total energy-momentum tensor divided into a radiation part and a vacuum part, i.e., $T_{\mu\nu} = T^r_{\mu\nu} + T^v_{\mu\nu}$. The radiation part

$$T^r_{\mu\nu} = (p_r + \varepsilon_r)u_\mu u_\nu - p_r g_{\mu\nu} \;, \tag{39}$$

and the vacuum part

$$\langle 0|T_{\mu\nu}|0\rangle \equiv T^v_{\mu\nu} = (p_v + \varepsilon_v)u_\mu u_\nu - p_v g_{\mu\nu} \;. \tag{40}$$

The first term in equation (40) must vanish in order to preserve the Lorentz invariance of the vacuum for all coordinate systems regardless of relative motion. Thus,

$$p_v + \varepsilon_v = 0 \;. \tag{41}$$

Equation (38) thus becomes (using equation (41))

$$R_{\mu\nu} - 1/2\, Rg_{\mu\nu} = 8\pi T^r_{\mu\nu} + 8\pi\varepsilon_v g_{\mu\nu} \tag{42}$$

which is of the form

$$R_{\mu\nu} - 1/2\, Rg_{\mu\nu} = 8\pi T_{\mu\nu} + \Lambda g_{\mu\nu} \tag{43}$$

so that we may identify[24]

$$\Lambda = 8\pi\varepsilon_V. \tag{44}$$

This is the form of Einstein's equations with non-zero "cosmological constant", except Λ is now a temperature dependent parameter.

In the early, high temperature universe, using the Robertson-Walker metric, equation (43) becomes

$$\left(\frac{\dot{R}}{R}\right)^2 = \frac{\kappa}{R} + \frac{\Lambda}{3} + \frac{8\pi G\varepsilon}{3} \simeq \frac{8\pi G}{3}(\varepsilon_r + \varepsilon_V) \tag{45}$$

for $\varepsilon_r \gg \varepsilon_V$ with $\varepsilon_r \propto T^4$, equation (45) yields the standard result $T \propto t^{-1/2}$, i.e., equation (29). However, when $\varepsilon_V \gg \varepsilon_r$ the result is

$$\frac{\dot{R}}{R} = \left(\frac{8\pi G}{3}\right)^{\frac{1}{2}} \varepsilon_V^{\frac{1}{2}}. \tag{46}$$

For temperatures not near the critical temperatures for symmetry breaking $\varepsilon_V(T) \simeq$ const. and it follows from equation (46) that the universe expands exponentially. This rapid expansion is a result of the large negative pressure of the vacuum.[25-27]. The result is an exponential stretching of the domains of CP coherence[23] from their initial size $\sim ct_X$, provided a first order (discontinuous) phase transition is involved. In the Sato scenario, the universe then supercools below T_c to a T_{c1} whereupon the transition becomes second order (continuous) or possibly driven,(cf. Reference 28) whereupon a rapid universal phase transition releases an energy density ε_V. The universe then reheats to temperatures where X-particles are produced, which subsequently decay to give baryon and antibaryon asymmetries on a macroscopic scale. These exponentially stretched domains of baryon and antibaryon excess may evolve[29] further leading to the formation of matter and antimatter galaxies in separate regions of the universe[30,31]. The scenario for the evolution of a baryon symmetric cosmology based on grand unification is shown in Figure 1.

OTHER THEORETICAL CONSIDERATIONS

The symmetry breaking mechanisms which we have been discussing can lead to the formation of various topological structures such as monopoles, strings and domain walls, which could affect the dynamics and isotropy of the universe. The problem of monopole formation has received the most attention since, for simple grand unification scenarios, the production of these particles would result in the universe having a mass density many orders of magnitude higher than astronomical observations allow[32]. Some suggestions for solving the monopole problem involve the exponential stretching process discussed in the last section[27] and multiple phase transition (symmetry breaking) scenarios[33]. The breaking of discrete symmetries can lead to domain wall formation, and it has been argued that such walls, if formed, must disappear at an early stage in order to be consistent with the observed homogeneity of the universe[34]. Clearly, the exponential stretching mechanism which has been invoked to solve the monopole problem could also alleviate the wall problem while providing a mechanism for domain growth. Vilenkin[35] has considered the dynamics of walls and strings and discussed several mechanisms for wall disappearance, one of which again involves multiple symmetry breaking[33]. He has also found that domain walls do not reflect light but do repel nonrelativistic particles. Such a repulsion might play a role in keeping matter and antimatter apart at some stage in the early universe. There is also the possibility that some models may provide a continuous solution set for the vacuum expectation value CP parameter θ, so that CP in this case is not a "discrete" symmetry[18].

The alternative of postulating intrinsic hard CP breaking, unvarying over all time and all space in the universe, leaves us in a rather bleak position from the epistomological point of view. For then our invokation of the whole apparatus of grand unified theory has only resulted in our replacing one parameter ($\eta = n_B/n_\gamma$) by another, viz., ε, equally mysterious if not more so.

GALAXY FORMATION

Various workers have tried to trace the growth of the regions of matter and antimatter by coalescence and Leidenfrost effects up to the era marking the decoupling of the matter and antimatter from the blackbody radiation field[29]. These studies have shown that baryon symmetric cosmology can lead more readily to galaxy formation than can the standard totally asymmetric cosmology[30, 31].

At a time of the order of 10^6-10^7 years after the big-bang, the cosmic plasma was cool enough to combine into neutral atoms. Starting at this point in the evolution of the universe, the questions of large scale structure and galaxy formation arise. Models of galaxy formation from "primordial turbulence" have always been attractive as a way of accounting for galaxy formation as well as for observed parameters such as the angular momenta and spatial distribution of galaxies. However, in that work, turbulence was introduced in ad hoc manner and, furthermore, such turbulence would be strongly damped out in the cosmic plasma because of the very high viscosity of the blackbody radiation field which remains coupled to the plasma until the neutralization ("recombination") epoch.

In the baryon symmetric cosmology scenario, this viscous dissipation is constantly fought by continuing radiation pressure from annihilation on the boundaries of matter and antimatter regions, which regenerates the turbulence. Radiation pressure from the annihilation, being directed generally away from the boundary regions, can drive mass fluid motions of the domains as well as causing further coalescence until the domains reach the size of galaxy clusters.

At the recombination epoch, two important changes were caused in the cosmic fluid motions. The viscosity dropped drastically and the turbulent fluid motions became supersonic. These changes occurred because the sound speed dropped sharply from its value in the cosmic plasma of $3^{-1/2}c$ (because the momentum was transferred by radiation) to the thermal velocity of the neutral gas. Thus, whereas the cosmic plasma behaved as a viscous incompressible fluid, both "small-scale" turbulence and density fluctuations could start to build up in the decoupled atomic fluid and later contract to form galaxies. In this scenario annihilation pressure can provide a continuous source of generating turbulence. This model for galaxy formation gives reasonable values for rotational velocities of galaxies and domain sizes (of galaxy cluster or supercluster size) for the present epoch.

THE COSMIC γ-RAY BACKGROUND RADIATION

One of the most significant consequences of globally baryon symmetric big-bang cosmology lies in the prediction of an observable cosmic background of γ-radiation from the decay of π^0-mesons produced in nucleon-antinucleon annihilations throughout the history of the universe. This is also perhaps at present the most encouraging aspect of this cosmology, since it satisfactorily explains the observed energy spectrum of the cosmic background γ-radiation as no other proposed mechanism does (with the possible exception of hypothetical point sources).

For high redshifts z, when pair production and Compton scattering become important, it becomes necessary to solve a cosmological photon transport equation in order to determine the γ-ray background spectrum. For a differential photon energy spectrum, we find this equation to be of the form

$$y\frac{\partial I}{\partial y} + \varepsilon\frac{\partial I}{\partial \varepsilon} = 2I + \frac{y^2 \Omega \nu}{[1+\Omega(y-1)]^{1/2}}\Big[A(\varepsilon)I - \int_{\varepsilon}^{b(\varepsilon)} d\varepsilon' B(\varepsilon|\varepsilon') I(\varepsilon', y) \quad (47)$$

$$-\xi^2 \, \Omega n_c y^3 v(T(y)) \frac{\sigma_A(T(y))}{\pi r_e^2} G_A(\varepsilon)\Big]$$

where $I = I(\varepsilon, y)$ is the annihilation γ-ray flux, and $y = 1 + z$, $\varepsilon = E_\gamma / m_e c^2$. The parameter $\nu = (n_c c/H_0)(\pi r_e^2)$, H_0 is the Hubble constant, r_e is the classical electron radius and σ_A the annihilation cross section, and $G_A(\varepsilon)$ is the source annihilation γ-ray function. The function $A(\varepsilon)$ is proportional to the total cross section for absorption and scattering of γ-rays by pair production and Compton interactions. The scattering function $B(\varepsilon|\varepsilon')$ is proportional to the probability that a γ-ray of energy ε' will Compton scatter to energy ε. The upper limit is

$$b(\varepsilon) = \begin{cases} \varepsilon/(1 - 2\varepsilon), & \varepsilon < 1/2 \\ \infty, & \varepsilon > 1/2 \end{cases} \tag{48}$$

The function $I_A(E_\gamma, y = 1)$ obtained by numerical solution of equation (47) corresponds to the present-day (z=0) γ-ray background spectrum predicted from these calculations to arise from matter-antimatter annihilations in the universe.

Figure 2 shows the observational data on the γ-ray background spectrum. The dashed line marked X is an extrapolation of the X-ray background component. The theoretical curve marked "annihilation" is the calculated annihilation spectrum[36]. The excellent agreement between theory and data is apparent. This striking evidence has been a prime motivation for studying BSDC. Other recent attempts to account for the γ-ray background radiation spectra by diffuse processes give spectra which are, in one way or another, inconsistent with the observations, generally by being too flat at the higher energies.

It is possible that the γ-ray background is made up of a superposition of point sources (see the paper of Fabian, these proceedings). However, since only one extragalactic source has been seen at energies above ~1 MeV, this remains a conjecture. Such a hypothesis must be tested by determining the spectral characteristics of extragalactic sources and comparing them in detail with the characteristics of the background spectrum. It presently appears, e.g., that Seyfert galaxies may have a characteristic spectrum which cuts off above a few MeV, so that they could not account for the flux observed at higher energies.

ANTIMATTER IN THE COSMIC RADIATION

Two groups have now reported the detection of antiprotons in the cosmic radiation[37,38]. Their results indicate a tantalizingly curious flux level. The level reported is a factor of ~ 4 to ~ 10 (energy dependent) higher than calculations of secondary production in interstellar cosmic-ray collisions predict. This is shown in Figure 3, after Szebelski et al.[39] (but with an erroneous calculation removed). Such calculations are, of course, dependent on the mean path length λ(g/cm^2) of matter traversed by the cosmic-rays. Measurements of the fluxes of secondary nuclei and positrons from π^+ decay produced by the same mechanism, give a value for the path length of $\lambda \approx 5$ g cm^{-2} (see Figure 4). The $\bar{p}$ flux reported, if of secondary origin, would, of course, require a value for λ a factor of ~ 4 to ~ 10 higher. (After this paper was given, a new determination of the $\bar{p}$ flux was recently reported (Buffington, et al.,Ap.J.in press) at low energies where the flux from secondary production is expected to be orders of magnitude lower than the measured value (see Figure 3)). This inconsistency may point to a primary extragalactic origin for the cosmic-ray antiprotons. Based on studies of galactic γ-rays, it is now generally believed that the bulk of the cosmic radiation is of galactic origin[40] except at the highest energies[41,42]. Since the γ-ray background observations require that matter and antimatter regions be separated on at least a galactic scale, a small extragalactic cosmic-ray flux containing $\bar{p}$'s would be consistent with a baryon symmetric domain cosmology (BSDC). An extragalactic cosmic-ray component with a flux $I_{ex}/I_{gal} \simeq \xi_{NG} \sim 10^{-5} - 10^{-4}$ is expected from leakage out of normal galaxies, based on rough energetic arguments[43]. These arguments also give an estimated flux from active galaxies[43] $\xi_{AG} \sim 10^{-3}$. If we assume that half of this flux is from antimatter sources, one gets a crude estimate for a BSDC primary $\bar{p}/p$ flux ratio in the cosmic rays $\bar{p}/p \sim 5 \times 10^{-4}$. This is interestingly quite close to the measured levels[38,39]. Present upper limits on the $\bar{\alpha}/\alpha$ ratio[44] are consistent with the $\bar{p}/p$ limit with the possible exception of a measurement in the 4.33 GeV/c range of $\bar{\alpha}/\alpha \lesssim 1 \times 10^{-4}$. (Buffington, et al., Ap. J., in press) find $\bar{\alpha}/\alpha < 2.2 \times 10^{-5}$ in the energy range 130-370 MeV/nucleon.) However, this latter upper limit is consistent with $\bar{\alpha}/\alpha = \xi_{NG}/2 = (\sim 5 \times 10^{-6} - 5 \times 10^{-5})$. Note that we can only argue that $\bar{\alpha}/\alpha = \bar{p}/p$ for cosmic ray production in normal galaxies, since we are comparing extragalactic fluxes with fluxes produced by processes in our own galaxy. It is conceivable that cosmic ray α's produced in the cores of active galaxies are broken up by collisions with matter or photons. Thus, the observed $\bar{p}$'s could come from active antimatter galaxies without accompanying $\bar{\alpha}$'s, but with the expected $\bar{\alpha}/\alpha \sim 10^{-5}$ from normal galaxies. In this case, future cosmic ray experiments may soon see $\bar{\alpha}$'s. Another possible cause for a lack of $\bar{\alpha}$'s would arise if the $\bar{p}$'s are from an early "bright phase" of cosmic ray acceleration associated with galaxy formation. It is possible for primordial helium to be photodisintegrated in the BSDC. This mechanism has been suggested to account for the recent observations of low He abundances in less evolved galaxies[45] and may in itself be an argument for BSDC. Thus galaxies, in their "bright phase", may have had very little helium to accelerate.

At present there are only upper limits on the fraction of antinuclei in the cosmic rays, consistent with the extragalactic primary hypothesis discussed above. Of course, a convincing detection of such antinuclei (Z>1) would strongly support BSDC, since they would not be readily produced in collision processes. In this regard, it is interesting to note that one reported cosmic-ray event, first interpreted to be a monopole, may have been a heavy antinucleus[46,47].

"CELL" STRUCTURE OF THE UNIVERSE

Not only do galaxies form clusters, but also these clusters of galaxies are not uniformly distributed; they cluster into superclusters. Between the superclusters are large voids--regions with a very low (possibly zero) space density of galaxies[48-50]. The existence of these holes, which is difficult to understand in the context of standard big-bang cosmology, is the kind of structure which can arise from a BSDC. The cosmic background γ-radiation originating from supercluster boundary annihilations should exhibit angular fluctuations which can best be studied with a high-resolution detector such as the 100 MeV spark chamber detector proposed for a future satellite "Gamma Ray Observatory".

The astronomical observations of the non-uniform "cell structure" distribution of galaxies also gains credence with the evidence of nonuniformity, which comes from studies of the origin and propagation of ultrahigh energy cosmic rays (UHCR). The lifetime of UHCRs should be cut short by their interaction with the background radiation. The result should be a high-energy cutoff in their energy spectrum which is not in accord with observation as shown in Figure 5. Various hypothesis have been proposed to account for the lack of a cutoff and detailed calculations have been made. After careful consideration of all the evidence it appears that the explanation lies in a true nonuniformity of the sources of these particles with the observed UHCRs coming mainly from within the local supercluster of which our galaxy is a member[41,42]. The obvious inference is that immediately beyond the region of the local supercluster there is a dearth of UHCR sources. Making the logical assumption that UHCRs are produced in galaxies or radio sources, we would then infer a real dearth of galaxies between the superclusters, supporting the domain structure viewpoint.

FUTURE TESTS USING HIGH ENERGY COSMIC RAY NEUTRINOS

Since a neutrino is not its own antiparticle, it is possible to determine whether a given source of cosmic-ray neutrinos is made of matter or antimatter. Several suggestions have been made recently for using high-energy neutrino astronomy to look for antimatter elsewhere in the universe[51-53]. These suggestions are all based on the fact that cosmic ray pp and pγ interactions favor the secondary production of π^+'s over π^-'s, whereas for $\bar{p}\bar{p}$ and $\bar{p}\gamma$ interactions the situation is reversed. The subsequent decay of the pions results in equal amounts of ν_μ's and $\bar{\nu}_\mu$'s of almost equal energies. However, π^+ decay leads to ν_e production, whereas π^- decay leads to $\bar{\nu}_e$ production. A production mechanism of particular importance in this context because of its large inherent charge asymmetry involves the photoproduction of charged pions by ultrahigh energy cosmic rays interacting with the universal 3K blackbody background radiation. The most significant reactions are

$$p + \gamma \rightarrow n + \pi^+ \qquad (49)$$
$$\bar{p} + \gamma \rightarrow \bar{n} + \pi^-$$

which occur in the astrophysical context principally through the Δ resonance channels because of the steepness of the ultrahigh energy cosmic ray spectrum.

There is a significant and potentially useful way of distinguishing ν_e's from $\bar{\nu}_e$'s, namely through their interactions with electrons. The $\bar{\nu}_e$'s have an enhanced cross section through formation of weak intermediate vector bosons such as the W^-, via $\bar{\nu}_e + e^- \rightarrow W^-$, the Glashow resonance effect[54]. For electrons at rest in the observer's system, the resonance occurs for cosmic $\bar{\nu}_e$'s of energy $E^W = M_W^2/2m_e = 6.3 \times 10^3$ TeV in the GWS model.

If one entertains the possibility of higher mass intermediate vector bosons[55], B^- and resonance energies $E^B = M_B^2/2m_e > E^W$ a feasible test for cosmic antimatter may be at hand.

The cosmic and atmospheric fluxes for $\bar{\nu}_e$'s, based cosmic ray production calculations[56] are shown in Figure 6. Assuming that there is no significant enhancement in the flux from production at high redshifts, the integral $\bar{\nu}_e$ spectrum from $\gamma\bar{p}$ interactions is expected to be roughly constant at 10^{-18} to 10^{-17} $\bar{\nu}_e$'s cm^{-2} sr^{-1} up to an energy of $\sim 2 \times 10^7$ TeV (2×10^{19} eV) above which it is expected to drop steeply. Figure 6 shows the estimated upper limit (UL) and lower limit (LL). It is expected that the largest competing background flux of $\bar{\nu}_e$'s will be prompt $\bar{\nu}_e$'s from the decay of atmospherically produced charmed mesons. The estimated upper and lower limits for this flux are also shown, and the position of the W^- resonance is indicated by an arrow. It can be seen that a cosmic $\bar{\nu}_e$ signal may be heavily contaminated by prompt atmospheric $\bar{\nu}_e$'s at the resonance energy E^W. The cosmic flux is expected to dominate the higher energies so that the existence of higher mass bosons B^- may be critical to any proposed test for cosmic antimatter using diffuse fluxes.

An acoustic deep underwater neutrino detector may provide the best hope for testing for cosmic antimatter by studying the diffuse background neutrinos. The practical threshold for such devices appears to be in the neighborhood of 10^3 - 10^4 TeV[57], where the W^- resonance occurs. For higher mass resonances B^-, the relevant neutrino resonance energy $E^B \propto M_B^2$ and the effective detection volume $V_{eff} \propto M_B^6$. Considering that the incident flux is expected to be roughly constant up to energies $\sim 2 \times 10^7$ TeV, one gains much in looking for higher mass Glashow resonances at higher energies. Acoustic detectors of effective volume >> 10 km^3 (10^{10} tons) may be economically feasible and consequently event rates of $\sim 10^2$ - 10^4 yr^{-1} may be attained in time.

The asymmetry in the production of charged pions in matter versus antimatter sources is reflected in cosmic-ray pp and $\overline{pp}$ interactions as well as $p\gamma$ and $\bar{p}\gamma$ interactions. Through the principal decay mode, this asymmetry is again reflected in a ν_e - $\bar{\nu}_e$ asymmetry and thus in the characteristics of events produced in deep underwater neutrino detectors. For ν-sources, these effects may be measurable at energies $\sim$ 1-10 TeV with optical detectors. The details of this possibility have been discussed by Learned and Stecker[51].

The possibility that $p\gamma$ and $\bar{p}\gamma$ interactions in quasars and active galaxies would produce significant fluxes of $\bar{\nu}_e$'s , detectable through the W^- resonance, has been suggested by Berezinsky and Ginzburg[52] as a way of looking for cosmic antimatter. Hopefully, this interesting suggestion will be explored in more detail as our understanding of the nature of cosmic ray production in compact objects increases.

It should be kept in mind that any positive observational data supporting the existence of large amounts of antimatter in the universe will be evidence of the spontaneous nature of CP violation at high energies, in accord with our earlier discussion. Indeed, as we have seen, various astrophysical data such as the cosmic γ-ray background spectrum, cosmic-ray $\bar{p}$ flux measurements, recent determinations of a low primordial He abundance[45,58], and galaxy clustering, can be interpreted as favoring this point of view.. The reader is referred to References 39 and 59 for a discussion of other aspects of BSDC. A longer review is planned for the near future.

REFERENCES

1. SCHWINGER, J., 1951, Phys. Rev. 82; 914-927.
2. SCHWINGER, J., 1953, Phys. Rev. 91; 713-728.
3. GLASHOW, S. L., 1961, Nuc. Phys. 22: 579-588.
4. WEINBERG, S., 1967, Phys. Rev. Lett. 19: 1264-1266.
5. SALAM, A., 1968, In Elementary Particle Theory, Ed. N. Svartholm (Almquist and Forlag, Stockholm): 367-377.
6. GEORGI, H. and S. L. GLASHOW, Phys. Rev. Lett. 32: 438-441.
7. ABERS, E. S. and B. W. LEE, 1973, Phys. Lett. 9C: 1-141.
8. BÉG, M. A. B. and A. SIRLIN, 1974, Ann Rev. Nucl. Sci. 24: 379-449.
9. WEINBERG, S., 1974, Phys. Rev. D9: 3357-3378.
10. LINDE, A. D., 1979, Rep. Prog. Phys. 42: 389-437.
11. WEINBERG, S. 1979, Phys. Rev. Lett. 42: 850-853.
12. MOHAPATRA, R. N. G. and SENJANOVIĆ, 1979, Phys. Lett. 89B: 57-60
13. GEORGI, H., H. R. QUINN and S. WEINBERG, 1974, Phys. Rev. Lett. 33: 451-454.
14. BURAS, A. J., E. ELLIS, M. K. GAILLARD and D. V. NANOPOULOS, 1978, Nuc. Phys. B135: 66-92.
15. LEARNED, J., F. REINES and A. SONI, 1979, Phys. Rev. Lett. 43: 907-910.
16. NANOPOULOS, D. V. and S. WEINBERG, 1979, Phys. Rev. D20: 2484-2493.
17. MOHAPATRA, R. N. and G. SENJANOVIĆ, 1978, Phys. Lett. 79B: 283-286.
18. BROWN, R. W. and F. W. STECKER, 1979, Phys. Rev. Lett. 43: 315-318.
19. SENJANOVIĆ, G. and F. W. STECKER, 1980, Phys. Lett. 96B: 285-288.
20. BARR, S., G. SEGRÉ and H. A. WELDON, Phys. Rev. D20: 2949-2498.
21. YILDIZ, A. and P. COX, 1980, Phys. Rev. D21: 906-909.
22. HARVEY, J. A., P. RAMOND and D. B. REISS, 1980, Caltech Preprint 68-758.

23. SATO, K., 1981, Phys. Lett. 99B: 66-70.

24. ZEL'DOVICH, YA. B. and I. D. NOVIKOV 1971, Relativistic Astrophysics (Chicago: University of Chicago Press) pg. 29.

25. BLUDMAN, S. and M. A. RUDEPMAN, 1977, Phys. Rev. Lett. 38: 255-257.

26. KOLB, E. W. and S. WOLFRAM, 1980, Astrophys. J. 239: 428-432.

27. GUTH, A. H., 1981, Phys. Rev. D23: 347-356.

28. WITTEN, E., 1980, Harvard preprint HUTP80-A040.

29. OMNES, R., 1972, Phys. Rep. 3C: 1-56.

30. STECKER, F. W. and J. L. PUGET, 1972, Astrophys. J. 178: 57-76.

31. See, e.g., Ref. 27 and references therein and EINHORN, M.B.,, D. L. STEIN and D. TOUSSAINT, 1980.

33. LANGACKER, P. and S. Y. PI, 1980, Phys. Rev. Lett. 45: 1-4.

34. ZEL'DOVICH, YA. B., I. YU. KOBZAREV and L. B. OKUN, 1974, Zh. Eksp. Teor. Fiz. 67: 3-11.

35. VILENKIN, A., 1981, Phys. Rev. D23: 852-857.

36. STECKER, F. W., 1978, Nature 273: 493-497.

37. GOLDEN, R. L., et al., 1979, Phys. Rev. Lett. 43: 1196-1199.

38. BOGOMOLOV, E. A., et al., 1979, Proc. 16th Intl. Cosmic Ray Conf. Kyoto 1: 330-335.

39. SZABELSKI, J., J. WDOWCZYK, and A. W. WOLFENDALE, 1980, Nature 285: 386-388.

40. STECKER, F. W., 1975, Phys. Rev. Lett. 35: 188-191.

41. STECKER, R. W., 1978, Comments on Astrophys. 7: 129-136.

42. CUNNINGHAM, G., et al., 1980, Astrophys. J. 236: L71-L75.

43. GINZBURG, V. L. and S. I. SYROVATSKII, 1964, Origin of Cosmic Rays (MacMillan: New York).

44. BADHWAR, G. D., et al., 1978, Nature 274: 137-139.

45. STECKER, F. W., 1980, Phys. Rev. Lett. 44: 1237-1240.

46. HAGSTROM, R., 1977, Phys. Rev. Lett. 38: 729-732.

47. PRICE, P. B., et al., 1978, Phys. Rev. D18: 1382-1421.

48. JOEVEER, M. and J. EINASTO., 1978, in The Large Scale Structure of the Universe (ed. M. S. Longair and J. Einasto) I.A.U. Symp. No. 79 (Dordrecht: Reidel Pub. Co.) pp 241-251.

49. GREGORY, S. A. and L. A. THOMPSON, 1978, Astrophys. J. 222: 784-799.

50. CHINCARINI, G. and H. J. ROOD, 1979, Astrophys. J. 230: 648-654.

51. LEARNED, J. G. and F. W. STECKER, 1980, Proc. Neutrino 79 Intl. Conf. Bergen, Norway (ed. A. Haatuft and C. Jarlskog) 2: 461-470.

52. BEREZINSKY, V. S. and V. L. GINZBURG, 1981, Mon. Not. Roy. Astr. Soc. 194: 3-14.

53. STECKER, F. W. and R. W. BROWN, 1981, Proc. 1980 DUMAND Summer Workshop, U. Hawaii, in press.

54. GLASHOW, S. L. 1960, Phys. Rev. 118: 316-317.

55. BROWN, R. W. and F. W. STECKER, 1981, Proc. 1980 DUMAND Summer Workshop, U. Hawaii, in press.

56. STECKER, F. W., 1979, Astrophys. J. 228: 919-927

57. BOWEN, T. and J. G. LEARNED, 1979, Proc. 16th Intl. Cosmic Ray Conf., Kyoto 10: 386-391.

58. STECKER, F. W., 1981, Phys. Rev. Lett. 46: 517.

59. STECKER, F. W., 1980, Astrophys. J. 235: L1-L3.

FIGURE CAPTIONS

Figure 1. Framework for evolution of a baryon symmetric domain cosmology.

Figure 2. Data on the cosmic γ-ray background radiation from Apollo 15 and the SAS-2 satellite. Also shown are upper limits obtained from high altitude balloon experiments.

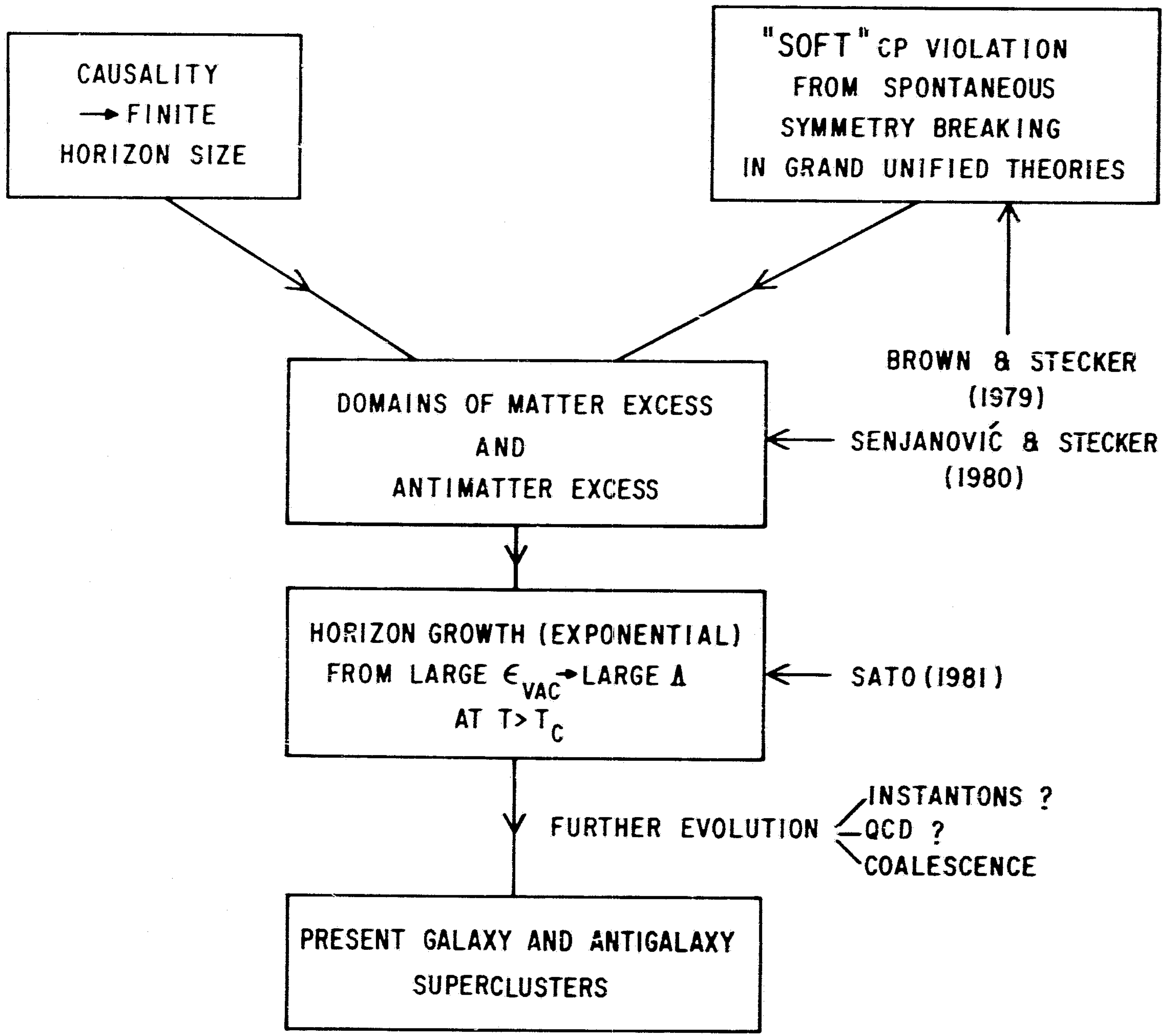
Figure 3. Measured cosmic ray antiproton/proton ratios and theoretical predictions for secondary antiproton production in interstellar cosmic ray collisions (after Szebelski et al., Ref. 39). The new, low energy determination of Buffington, Schindler and Pennypacker, Astrophysical Journal,in press), which is orders of magnitude above the flux expected from secondary production, is also shown. All data are consistent with a primary extragalactic $\bar{p}$ flux with $\bar{p}/p \approx$ const. plus a secondary $\bar{p}$ component at higher energies.

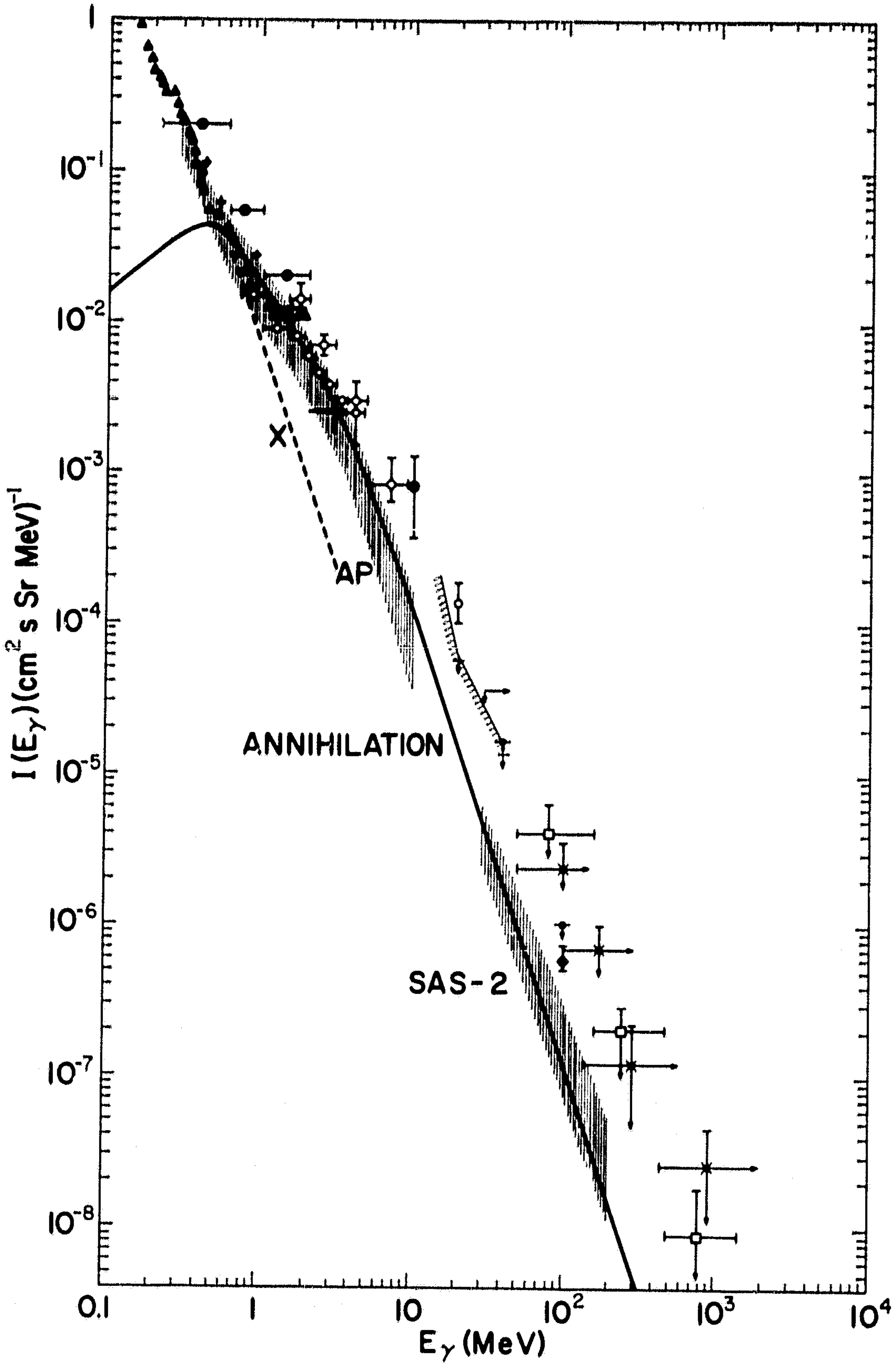
Figure 4. Path lengths implied by $\bar{p}$, e^+ and secondary nucleus production (Ref. 39).

Figure 5. The ultrahigh energy cosmic ray spectrum (xE^3) from extensive air shower data together with the expected spectrum (solid curve) for a uniform cosmic ray source distribution showing the expected high energy cutoff (Ref. 42).

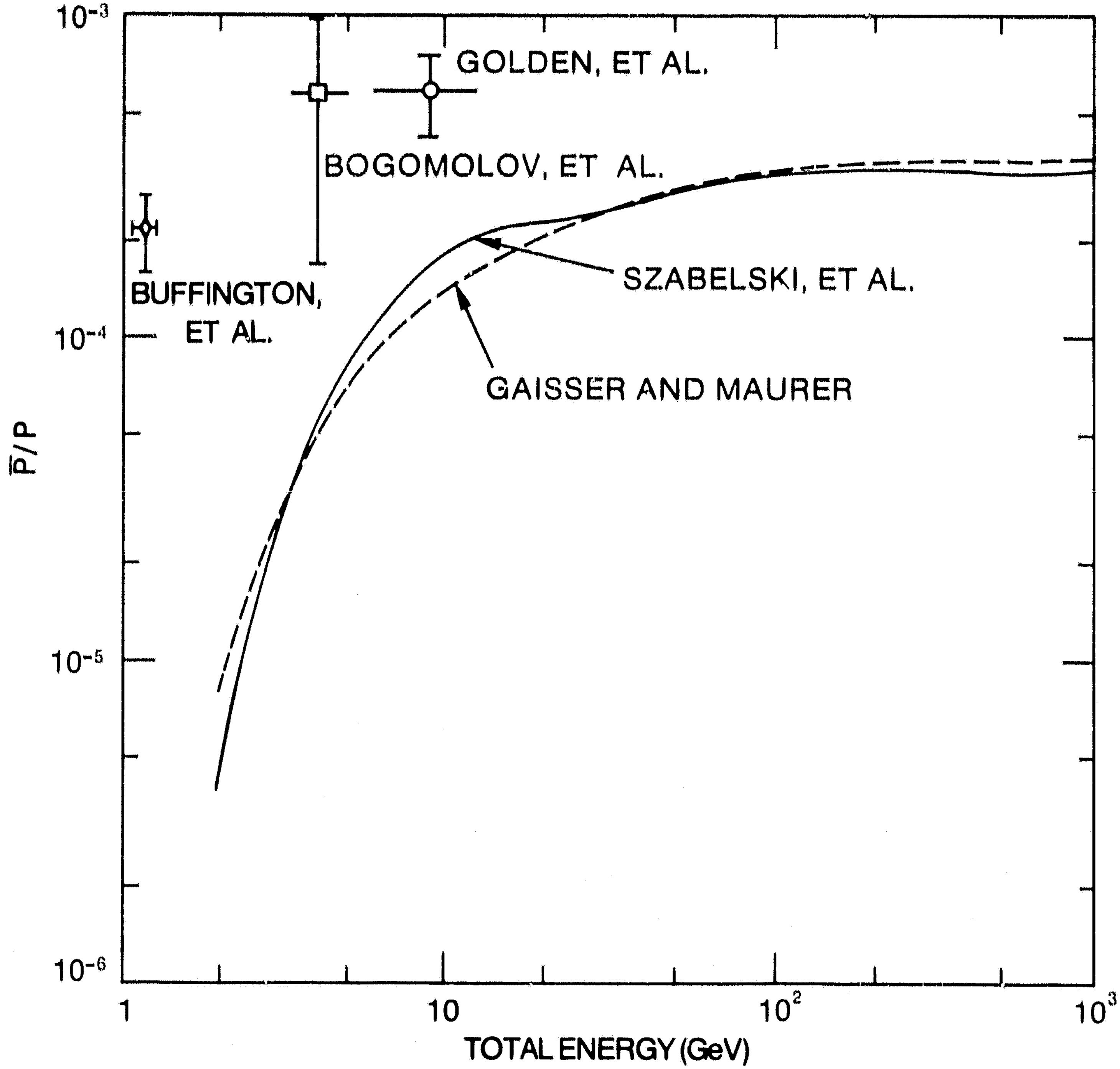
Figure 6. Cosmic and atmospheric $\bar{\nu}_e$ fluxes (see text).

SIMPLEST BARYON SYMMETRIC BIG-BANG SCENARIO

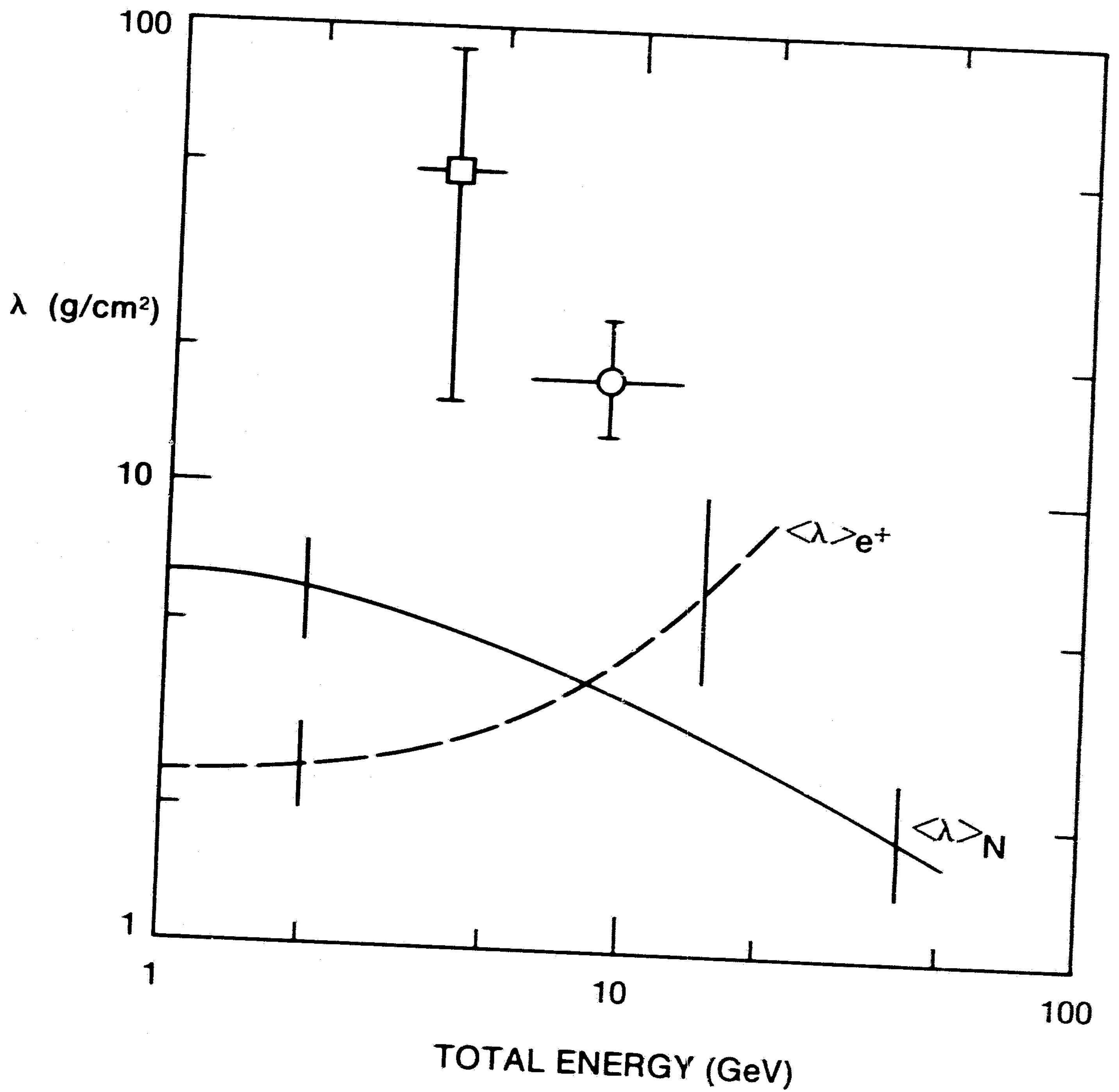


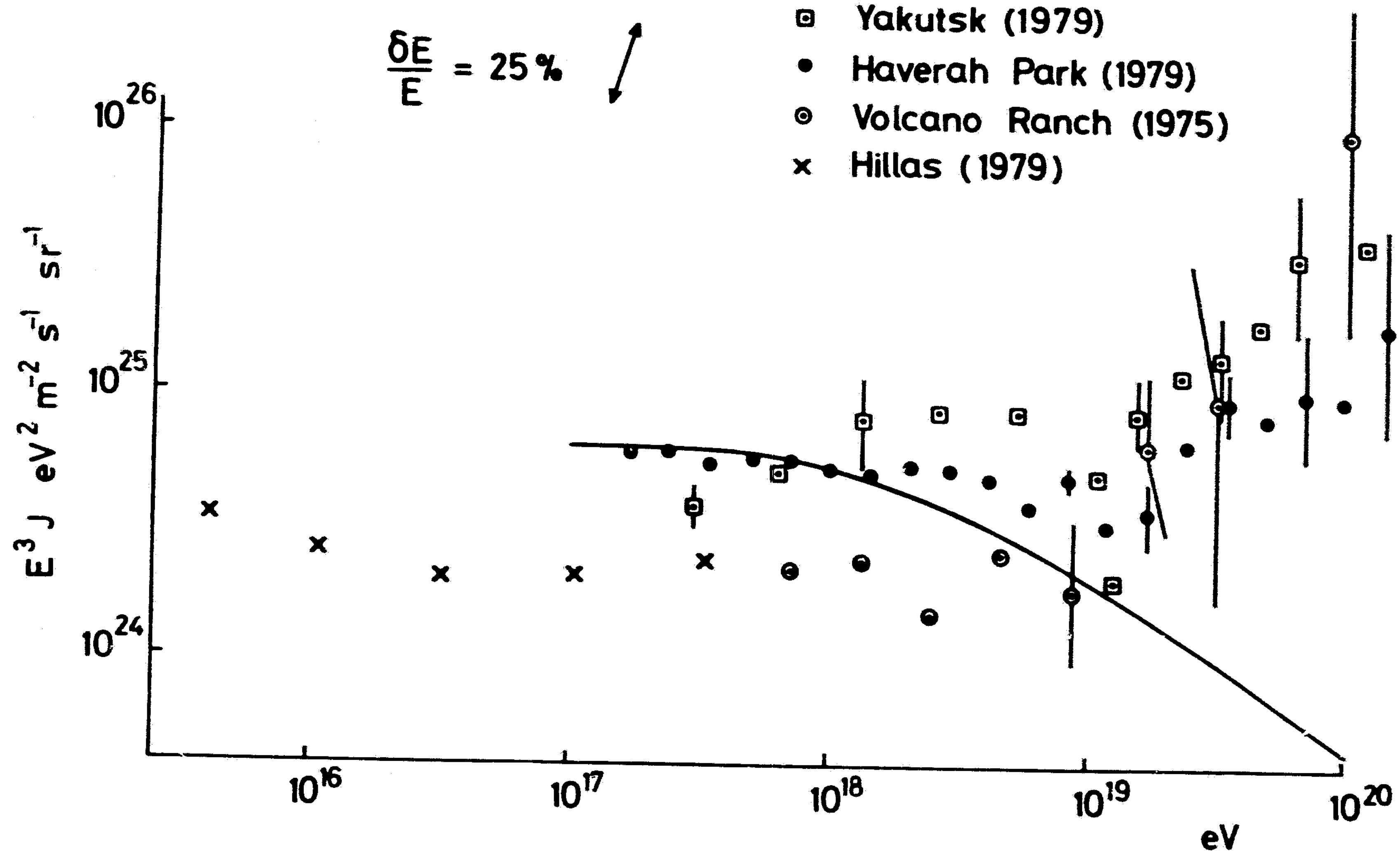


I (Eγ)(cm² s Sr MeV)⁻¹
Eγ (MeV)
X
AP
ANNIHILATION
SAS-2

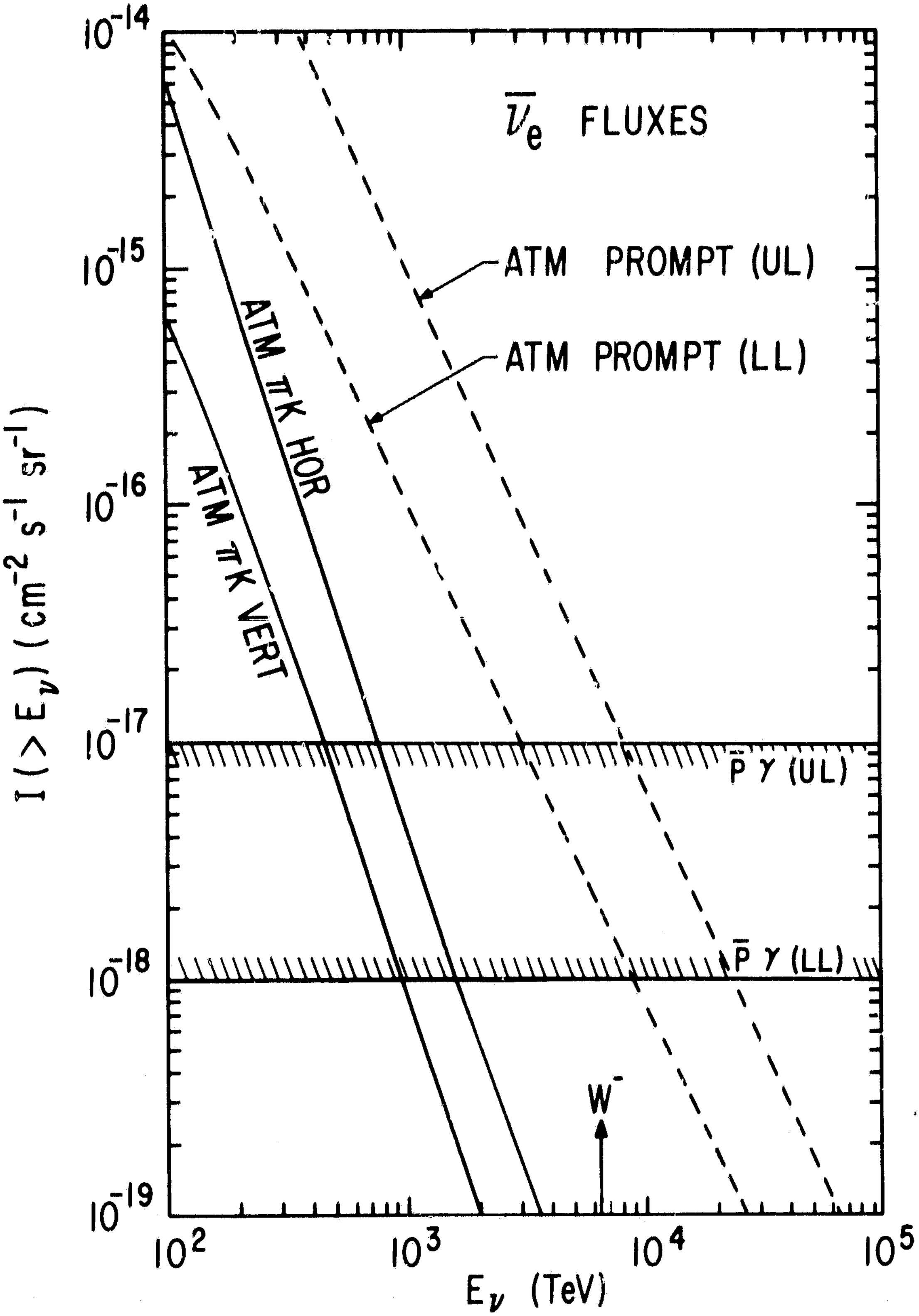


GOLDEN, ET AL.
BOGOMOLOV, ET AL.
BUFFINGTON, ET AL.
SZABELSKI, ET AL.
GAISSER AND MAURER
$\overline{P}/P$
10^{-3}
10^{-4}
10^{-5}
10^{-6}
1
10
10^{2}
10^{3}
TOTAL ENERGY (GeV)


100
10
1
λ (g/cm²)
⟨λ⟩e+
⟨λ⟩N
1
10
100
TOTAL ENERGY (GeV)



$\frac{\delta E}{E} = 25\%$
Yakutsk (1979)
Haverah Park (1979)
Volcano Ranch (1975)
Hillas (1979)
$E^3 J$ $eV^2 m^{-2} s^{-1} sr^{-1}$
10^{26}
10^{25}
10^{24}
10^{16}
10^{17}
10^{18}
10^{19}
10^{20}
eV



$\bar{\nu}_e$ FLUXES
ATM PROMPT (UL)
ATM PROMPT (LL)
ATM πK HOR
ATM πK VERT
$\bar{p}\,\gamma$ (UL)
$\bar{p}\,\gamma$ (LL)
W^-
10^{-14}
10^{-15}
10^{-16}
10^{-17}
10^{-18}
10^{-19}
$I(>E_\nu)$ $(\mathrm{cm^{-2}\,s^{-1}\,sr^{-1}})$
10^2
10^3
10^4
10^5
E_ν (TeV)