This is a summary of a loosely connected collection of results on supersymmetry (SUSY) breaking and cosmology in the context of superstring theory. The requirement of a satisfactory inflationary cosmology puts strong constraints on the superstring vacuum state. Some of these are phenomenological, i.e. certain things must be true if the theory is to reproduce observations of the Cosmic Microwave Background. By far the strongest constraints come from the simple requirement that the theory have a reasonable probability of producing a large universe. One can gain a tentative understanding of why there are no more than four large spacetime dimensions, and of why supersymmetry must be broken, based on this requirement alone. The problem of dilaton domination of the energy density of the universe is discussed and a plausible resolution due to Lyth and Stewart mentioned. Some remarks are made about the pattern of low energy SUSY breaking. In particular, we point out several indications that low energy, as opposed to hidden sector SUSY breaking, will pose cosmological problems in string theory. If it is the correct mechanism then the

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theory will contain very light particles which will give rise to coherent forces of millimeter to centimeter range. Finally, some speculations based on the holographic principle are advanced which enable one to estimate the cosmological constant. If SUSY breaking comes from a hidden sector the calculation gives a result near the present observational bound. For low energy SUSY breaking the result is too small to be of any observational significance.
1. Introduction

In this talk I want to summarize a loosely connected collection of results that I have obtained in collaboration with M.Dine, M.Berkooz and others \[1\], over the last several years. The general theme of these investigations was that the effective field theory of superstrings (EFTS), informed by cosmological considerations, can give us the beginnings of an understanding of how the Universe chooses the correct vacuum state of superstring theory. If they succeed in doing nothing else, I hope that these arguments will convince everyone that the correct context of this question is cosmological. The problem is very different from that of solving for the ground state of a Lorentz invariant field theory that does not include gravity.

Our work was based on two fundamental assumptions. The primary one is the assumption that the EFTS is weakly coupled at a scale just below the string scale\[1\]. We have several pieces of evidence for this. The first is the observed unification of couplings at a value $g^2 \sim \frac{1}{4\pi}$. The second is the striking(?) resemblance of the perturbative spectra of some string vacua to the spectrum of elementary particles in the world. However, I consider the hierarchy problem, the existence of a wide range of dynamical scales in nature, to be the most compelling argument in favor of weak coupling. At present we have no other way of understanding this fundamental fact than by attributing it to the properties of marginally relevant gauge couplings, which are weak at high energies.

The fact that we are in a weakly coupled region leads directly to our second assumption, namely that we cannot expect cancellation between different exponentially small effects in the coupling. Since some of the most popular models of string phenomenology are based on just such a cancellation \[3\] I should give a definitive argument ruling these models out. I will not. Instead I will characterize my second assumption as religious. I cannot give strong arguments to justify it, but I believe in it strongly enough that I will let it guide my future actions.

Within the context of these assumptions, there is only one known way to solve the

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1 Another hidden assumption is that the string scale is indeed the fundamental short distance scale in string theory, and that other scales like the Planck scale, or the Shenker scale\[2\] have something do do with the number of degrees of freedom per string scale.
Dine-Seiberg problem of the runaway dilaton. The constraints of analyticity, and the perturbative part of S-duality, guarantee that the superpotential of the dilaton is well approximated by a pure exponential in the weak coupling region, and that the dominant exponential comes from low energy field theory effects (if any). The dilaton can nonetheless be stabilized in the region of weak bare field theory coupling, if its Kahler potential is not well approximated by the tree level formula. This is plausible, because Shenker argued that generic objects in string theory have nonperturbative corrections of order $e^{-c\bar\phi}$. The results of show that there are no such corrections to the superpotential, but the argument does not apply to the Kahler potential. This means that the perturbation series for the Kahler potential may break down even for values of the coupling for which the low energy effective field theory is weakly coupled. In a matrix model argument was given suggesting that $c \sim 1$. This is supported by an explicit calculation of such a nonperturbative effect by Polchinski.

A further argument for the absence of dimension dependent powers of $2\pi$ in the constant $c$, comes from an analysis of Feynman graphs in string theory and quantum field theory. The analog of the world sheet representation of a string amplitude in field theory is the Schwinger parameter representation of Feynman graphs. This is derived from the momentum space representation by doing all of the loop momentum integrals for fixed values of the Schwinger parameters. Each loop integral contributes a factor of $\frac{1}{(4\pi)^{d/2}}$ obtained from the explicit momentum space measure and the gaussian integral over momenta ($d$ is the spacetime dimension). The origin of this factor is basically the normalization of wave packets to one particle per unit volume, plus the fact that particles must meet at a point in order to interact. This factor multiplies the Schwinger parametric integrand. By contrast, in string theory we have only a single momentum integral multiplying the integral over moduli space at any genus. Presumably this is a consequence of the fact that strings do not have to meet at a point in order to interact. Thus, a priori the expansion parameter in string theory is $g^2$ rather than the coupling divided by a dimension dependent power of $4\pi$.

Of course, if one writes the integral over moduli space in terms of sums over conformal field theory states propagating in tubes (a string field theory like decomposition), one
can see a possible origin for multiple momentum integrals in string theory. The point of our argument is that, as a consequence of the more compact representation in terms of a partition function integrated over multigenus moduli space, the existence of inverse powers of \(2\pi\) is much less obvious than in field theory (even in field theory one could imagine that the parametric integrals gave compensating powers of \(2\pi\) in the numerator). Further, the intuitive field theoretic argument for the geometric factors is based on locality, which is maximally violated in the “fat” part of moduli space which gives rise to the large order behavior of string perturbation theory. What we really need to settle this issue is a refinement of Shenker’s estimate of the large order behavior.

The bottom line of all of this speculation is a simple model of dilaton stabilization by an exponential superpotential and a complicated nonperturbative kahler potential. It is amusing that this mechanism of stabilization of the dilaton can work, and give zero cosmological constant, only if SUSY is broken. This is a consequence of a more general principle. A nonzero superpotential (in the vacuum state) is the order parameter for R-symmetries larger than the \(Z_2\) which simply reflects the supercoordinates. With an exponential superpotential, any nonzero value of the coupling thus implies broken R-symmetry. The standard supergravity formula for the potential

\[
V = e^K (|D_i W|^2 - 3|W|^2)
\]

then implies that if R-symmetry is broken, SUSY must also be broken if the cosmological constant vanishes. In the weak coupling framework advocated here, SUSY breaking is an automatic consequence of the dynamical principles of dilaton stability and vanishing cosmological constant.

To my mind, this is the simplest of all possible models of SUSY breaking in string theory. Nonperturbative dynamics will generically lead to a nonvanishing superpotential (see below), which is guaranteed to be dominantly a simple exponential in the weak coupling regime. If the Kahler potential can stabilize the dilaton with vanishing cosmological constant in this regime, then SUSY is broken. The price we must pay for this attractive scenario is a certain impotence with regard to explicit computations. Our assumptions imply that in the regime of interest we cannot calculate the detailed form of the potential
on moduli space. In the cosmological context to which we now turn, this will have some advantages (of the same nature of the advantages of the proverbial fig leaf.). Explicit perturbative superstring computations, combined with exponential superpotentials, lead to potentials which are not compatible with slow roll inflationary cosmology.

2. Modular Cosmology

The study of the cosmological effects of superstring moduli is necessary because moduli typically lead to problems in cosmology. It is attractive because conventional slow roll inflation cries out for fields with the properties of moduli. The fundamental assumption of slow roll inflation is that cosmological friction, a term of order $M_P^{-1}$ in the equations of motion, can compete with the force coming from the gradient of the potential. In conventional effective field theory, in which field values of order the Planck scale are forbidden, potential gradients of order $M_P^{-1}$ can only be achieved by unnatural fine tuning of dimensionless parameters. Even if one ignores the effective field theory restriction on field space, as in theories of chaotic inflation[7], then one must fine tune what appear to be dimensionless parameters in order to obtain density fluctuations of small enough amplitude.

By contrast, in string theory the natural range of variation of the moduli is the Planck scale. $\nabla V$ is always of order $M_P^{-1}$. Furthermore, the potential vanishes to all orders in perturbation theory, and is nonperturbatively bounded[2] by something of order $e^{-\frac{8\pi^2}{k\sigma^2}}$, where $k$ is an integer which cannot be too large ($k$ is related to the size of the low energy gauge group, whose rank is bounded in string theory). If a nonperturbative potential is generated on moduli space then the Hubble friction will naturally be of the same order of magnitude as the restoring force and the amplitude of density fluctuations will naturally be much less than one. The order of magnitude estimate one obtains for the number of e-foldings of inflation is 1. We know too little about the lagrangian for the moduli to decide whether the phenomenological constraint of more than 60 e-foldings poses a fine tuning problem.

The notion that string theory moduli are inflatons[8] also leads to a simple explanation of why we only see three large space dimensions. Moduli space is noncompact, but has
finite volume. For any finite value of the moduli, the universe has finite volume and we should in principle seek a quantum mechanical wave function for even the zero modes of the moduli fields. It is only in the regions of moduli space corresponding to large volume universes that we obtain approximate decoherence of states with different values of the homogeneous modes.

Let us try to sketch out the quantum mechanics of the homogeneous modes by first neglecting both the inhomogeneous modes and gravity. This is the approximation of quantum mechanics on moduli space. For moduli an order of magnitude or so larger than the string scale, we might argue that the effect of inhomogeneous modes would be just to renormalize the metric and potential on moduli space. This renormalization should go to zero as we approach the region of moduli space corresponding to large volume of more than three dimensions, because all low energy interactions are infrared free. In these extreme regions of moduli space then, the lagrangian of the moduli is well approximated by free motion in the classical metric. This system has a normalizable eigenfunction which is a constant on moduli space. Since the volume of moduli space is finite, the probability of being at large radius is predicted to fall like a power of the radius. Thus, any macroscopic universe is highly improbable.

This argument fails in four (or fewer) spacetime dimensions. Well understood effects in low energy nonperturbative field theory can lead to a nonperturbative superpotential on moduli space even in the limit of infinite three dimensional volume. Once such a potential is generated, the coupling of gravity, via the phenomenon of inflation, can invalidate the above argument. In dimension higher than four, the only scale for the potential of scalar fields is the Planck scale. Thus there is no low energy effective field theory regime in which we can show that inflation occurs. Further, even if we could justify the use of classical field theory, Planck scale inflation would lead to enormous density fluctuations, and a highly inhomogeneous universe filled with black holes. The inflationary escape from the prediction of a small universe is not available in high dimension string theory.

In order to prove that inflation occurred, one would have to have more detailed information about the wave functional of the inhomogeneous modes of the field. One would want to show that the energy density was dominated by the potential energy of the con-
stant mode over a region of space of size greater than the horizon volume corresponding to that potential energy. Thus we have not shown that inflation is inevitable in four dimensions, only that it does not occur in more than four, and that in its absence, string theory predicts the size of the world to be very small.

Our basic contention then is that string theory predicts a very small probability for the universe to be many orders of magnitude larger than the string scale unless there are regions of moduli space in which inflation occurs. On the other hand, we will see below that the EFTS approximation probably breaks down if there are too many e-foldings of inflation. Furthermore, dimensional analysis (and some knowledge of what happens in noncompact regions of moduli space) also suggests that string theory cannot lead to too many e-foldings of slow roll inflation. The order of magnitude estimate is \( O(1) \).

2.1. Density Fluctuations and SUSY Breaking

We have noted above that the string theorist’s expectations for nonperturbative modular physics lead to a plausible model of inflation. For moduli we expect a potential of the form

\[
V = M_P^4 \sum e^{-c_i \text{Re} S} V_i(\Phi) \equiv M_i^4 V_i. \tag{2.1}
\]

Where \( S = \frac{8\pi^2}{g^2} + i\theta; \quad 0 \leq \theta \leq 2\pi. \) is the dilaton/axion field and \( \Phi \) is a notation for all of the dimensionless moduli fields including \( S \). The \( V_i \) are supposed to be slowly varying functions of \( \text{Re} S \) in the weak coupling region. In [5] it was argued that truly stringy nonperturbative contributions were characterized by \( c_i = \text{even integer} \), while nonperturbative low energy field theory can give fractional values of \( c_i \). Assuming that the potential corresponding to the smallest \( c_i \equiv c_1 \) can lead to inflation, the order of magnitude amplitude of density fluctuations is

\[
\frac{\delta \rho}{\rho} \sim \left( \frac{M_1}{M_P} \right)^2. \tag{2.2}
\]

Observation thus suggests \( M_1 \sim 10^{-2.5} M_P \).

This is much larger than the highest allowed value for the square root of the SUSY breaking order parameter \( F \). If we assume that SUSY provides the resolution of the weak interaction gauge hierarchy problem then \( \sqrt{F} < 10^{-8} M_P \). The simplest resolution of this
discrepancy is to assume that the leading order potential $V_1$ has a supersymmetric minimum. In the inflationary cosmological context which we are discussing, this has dramatic consequences. The supergravity formula for the potential (1.1) is generically negative for supersymmetric vacua. Finding a supersymmetric stationary point requires us to solve $N$ complex equations for $N$ complex unknowns. Requiring the superpotential to vanish as well is an additional equation, which will not generally be satisfied.

In the inflationary context, the resulting negative cosmological constant is more than just a phenomenological embarrassment. Inflation has made the spatial curvature terms in Einstein’s equations completely negligible. As a consequence, there are no solutions of the postinflationary cosmological equations in which scalar fields come to rest at a minimum with negative cosmological constant. Instead, if the system tries to relax to such a stationary point, the universe recollapses on microscopic time scales, and inflation is not successfully completed. Thus, the only supersymmetric minima into which the universe can settle after inflation, are nongeneric ones where the superpotential vanishes.

String theory can provide us with such nongeneric superpotentials, but as might be expected, they can only occur at special points in moduli space. First let us examine the simplest nonperturbative field theory mechanism for generating a superpotential on moduli space, gaugino condensation. This leads to a superpotential of the form

$$W = e^{-CSw(M)}$$

(2.3)

Here $M$ are the nondilatonic moduli, and the function $w$ is determined by a one loop vacuum polarization calculation. $C$ is inversely proportional to the beta function of the hidden sector gauge group.

More generally, the low energy gauge theory which determines the leading contribution to the superpotential on moduli space will contain massless matter fields. In many cases, the low energy contribution to the superpotential will favor infinitely large values of matter field VEVs. Generically however, there will be terms of dimension four or higher in the tree level superpotential which will stabilize these VEVs at a finite value. The result of integrating out the matter fields will be a superpotential for the moduli of the form

$$(e^{-CSw(M)})^pF(M) = \Lambda_{\text{low}}^pF(M)$$

(2.4)
The second form of this equation is rewritten in terms of the nonperturbative scale $\Lambda_{\text{low}}$ of the low energy gauge theory, which depends implicitly on the moduli. The power $p$ is determined by the balance between the nonperturbative superpotential for the matter fields, and the tree level term, and is always positive. Finally, the function $F$ incorporates the moduli dependence of the coefficient of the leading term in the tree level superpotential; the term which stabilizes the massless modes of the matter fields.

A mechanism for obtaining a vanishing superpotential becomes apparent if we imagine that at some point $M_0$ in moduli space a number of other charged chiral superfields become massless. The effective value of $\Lambda_{\text{low}}$ is then determined by renormalization group matching\[10\] to a theory above the scale of the masses of these fields but below the string scale. If there are enough new massless states at $M_0$, then $\Lambda_{\text{low}}$ vanishes as the masses go to zero. If there is no compensating divergence in the function $F$ then the superpotential will vanish at $M_0$. If it vanishes rapidly enough, $M_0$ will be a supersymmetric minimum of the potential.

Even without going into details of specific models, we can make some general remarks about this mechanism. It probably cannot work for (2, 2) compactifications, where general theorems restrict the number of massless chiral multiplets charged under the low energy gauge group. The extant examples of this mechanism\[11\] occur in vacua obtained from a tree level (2, 2) vacuum by shifting fields to cancel the Fayet-Iliopoulos term of a tree level anomalous $U(1)$ gauge symmetry\[2\].

Returning to our model of inflation, we note that since the inflationary superpotential depends only exponentially on the dilaton, the vacuum energy at the minimum is zero independent of the dilaton field. This is an example of the lesson we learned in the first section. The dilaton cannot be frozen by a SUSY preserving superpotential with vanishing cosmological constant.

It is reasonable to assume, though by no means guaranteed, that the nondilatonic

\[2\] In passing we note that these vacua have been argued to be exact quantum eigenstates of string theory. They are supersymmetric, have zero vacuum energy and are part of a continuous moduli space. It is a challenge for string theorists to find an argument that explains why the world is not sitting in such a state.
moduli all obtain mass from the potential $V_1$. This means that their masses are of order $10^{-5}M_P$ and their nominal reheat temperature of order $10^{-7.5}M_P$. This is high enough to give us reason to worry about generating enough gravitinos at reheating to cause problems with nucleosynthesis. Perhaps we should not take this too seriously. We have been careless about numerical factors, and have ignored the difference between the Planck scale and the string scale (whose ratio might come into the estimate of the reheat temperature raised to a rather large power). What is clear is that if the nondilatonic moduli are frozen by the potential which generates inflation, they do not pose a reheating problem for the universe. Temperatures high enough for baryogenesis and nucleosynthesis are achieved by thermalization of the decay products of these moduli.

2.2. Dilatonic Dilemmas

We now come to the cosmological problems associated with the dilaton. The first of these is the cosmological version of the Dine-Seiberg problem. Brustein and Steinhardt\cite{12} argued that even if a potential for the dilaton was generated by SUSY breaking (and we have argued above that there is no alternative in the weak coupling region), the energy in the dilaton zero mode at early times would generally send it flying over this tiny barrier into the extreme weak coupling region. While the conclusion is far from obvious even in a general model\cite{3}, our present scenario provides a rather neat and definitive resolution of the problem. The dilaton will be stabilized by the second term in (2.1), which, as in section 1, is also the term responsible for SUSY breaking. This potential will have a minimum at $S_0$. However, during inflation, the first term in (2.1) will give a much larger effective potential for the dilaton as long as the moduli $M$ are far from their minimum. Since the potential is an exponentially varying function of the dilaton, the dilaton will not be one of the inflaton fields. In other words, during inflation it will quickly be drawn into a local minimum $S(M)$ which depends on those of the other moduli which are undergoing frictionally dominated motion. If the end of this “dilaton groove”, $S(M_0)$ lies on the strong coupling side of the minimum $S_0$, then the Brustein-Steinhardt problem is solved. The dilaton is gently deposited on the strong coupling side of its true minimum at the end of inflation. Its post

\footnote{See the appendix of 13}
inflationary energy is of the order of the barrier height between the minimum and extreme weak coupling.

We do not have the tools at present to calculate the dilaton potentials and so we cannot say that string theory definitely avoids the Brustein-Steinhardt problem. However, we see that some simple qualitative properties of the potential are sufficient. We cannot with equal ease solve the problem of dilaton domination of the energy density of the universe\cite{14}. The dilaton mass, since it is associated with hidden sector SUSY breaking, is of order $\frac{F}{M_P} \sim 10^{-16} M_P$. Its reheat temperature is thus of order $0.01 MeV$ and one is hard put to understand the success of conventional nucleosynthesis calculations. The most promising solution to this problem has been suggested by Lyth and Stewart\cite{15}. They argued that in many simple models, there would be a period of $O(10)$ e-foldings of thermal inflation, with Hubble constant of order the weak scale. Randall and Thomas\cite{16} had argued that such an era of intermediate scale inflation could redshift away the dilaton (and gravitino) energy densities, without affecting the observable properties of the Cosmic Microwave Background Anisotropy. The advantage of the Lyth-Stewart idea is that during thermal inflation the dilaton is attracted to a point quite close to its true minimum, so that there is no post-inflation Polonyi problem. Lyth and Stewart also suggest that their scheme may be able to accommodate an adequate amount of baryogenesis.

Another, more speculative idea, is the assumption that the order of magnitude estimate of the dilaton reheat temperature is simply wrong by two orders of magnitude. Then the dilaton could decay before nucleosynthesis. If there are renormalizable baryon number violating operators in the supersymmetric standard model, dilaton decay could also be the agent of baryogenesis. In such a scenario, the small size of the baryon to photon ratio is primarily explained by the large ratio between the dilaton mass and its reheat temperature. Such a model might be compatible with conventional cosmology, but the required violation of R-parity might deprive us of a natural candidate for dark matter.

Let us also note an issue that has not been addressed by any of the attempts to avoid dilaton domination of the universe\cite{13}. In all extant models, there is a period, stretching roughly from the time the scale of the energy density is $10^{11}$ GeV, until the weak scale, during which the dilaton or some other coherent scalar does dominate the energy density.
Energy density fluctuations on scales smaller than the initial horizon, grow by a factor of almost $10^{11}$ during this era. Since they initially had amplitude of order $10^{-5}$, they go nonlinear long before the weak scale is reached and the mechanisms for getting rid of the dilaton go into effect. The consequences of these nonlinear gravitational phenomena have not been worked out, and until they are, any picture of this era will be accompanied by a large question mark.

Finally, let me note that all resolutions of the problem of dilaton domination work only for hidden sector SUSY breaking. In a theory with low energy SUSY breaking only, the dilaton will be extraordinarily light. Indeed, for reasonable values of the low energy SUSY breaking scale ($1 - 100$ TeV), it is light enough to be on the edge of detection by Cavendish experiments. In such a model, the dilaton is stable, and its energy density cannot be inflated away until long after nucleosynthesis. Thus, all string theory models with low energy SUSY breaking have a serious unsolved cosmological problem.

3. The Cosmological Constant

A major lacuna in any discussion of cosmology is of course the question of why the cosmological constant is so small. We have provided an explanation above of why stationary points of the effective potential with negative cosmological constant (and in particular, typical supersymmetric minima) cannot be accessed after inflation. There is no corresponding mechanism for vacua with positive cosmological constant.

I believe that the answer to this question goes beyond the low energy effective field theory description of string theory that we have used in this lecture. The main feature which distinguishes DeSitter space from flat space is the existence of a horizon. The conventional low energy field theory picture of DeSitter space consists of an ever growing number of causally disconnected regions, each with its own independent set of field theoretic degrees of freedom. Recent work of Lowe et. al. [17] suggests that such a picture is a very bad approximation to string theory even when the cosmological constant is small and spacetime is locally flat. The discussion of [17] was oriented towards black holes, but their actual calculations referred primarily to the low energy effective description of a certain
set of (nice slice) coordinate systems in flat spacetime. These are a model of smooth coordinates which cross the horizon and interpolate between the external Schwarzschild coordinates and some good set of coordinates in the interior of a large black hole. A similar picture is sure to be approximately valid in DeSitter space. The major technical problem in trying to apply these ideas to DeSitter space, is that DeSitter space is not a solution of the classical equations of string theory.

The message of [17] is that the low energy (with respect to nice slice coordinates) effective description of string theory is not a local field theory. There are low energy states which are highly nonlocal in nice slice coordinates. They correspond to long strings stretched from one side of the horizon to another. Furthermore, there is no subset of local measurements which shows independence of the degrees of freedom on one side of the horizon from those on the other. Susskind argues that this is evidence for the holographic nature of string theory. The theory has no bulk degrees of freedom.

If one accepts these arguments, there are several clear implications. First of all, if an effective field theory calculation predicts a positive cosmological constant in string theory, then it is not self consistent. The effective field theory is not a valid description of low energy string theory on scales much larger than the DeSitter horizon. (It also breaks down very near the horizon in static coordinates which cover only a single horizon volume). Secondly, if there are really no bulk degrees of freedom, no effective field theory calculation of the cosmological constant is correct in string theory.

These ideas are very crude and we surely do not yet have an accurate idea of what they mean. I would like nonetheless to give the first outline of how one might calculate the cosmological constant in holographic string theory. I will begin by assuming that string theory describes our present universe as a roughly spherical object of radius $R$ (by which I mean only that its volume is approximately the three halves power of its area). Furthermore, the true degrees of freedom of the system are assumed to be distributed on a two dimensional submanifold of area $R^2$. This is the transverse submanifold of some light cone frame. Part of the magic of the theory, which we are very far from understanding, is that we can choose to locate the degrees of freedom on any submanifold with the same area. The density of degrees of freedom of the system on the submanifold is assumed to
obey the Bekenstein rule: one degree of freedom per Planck area.

On the other hand, we know that the Universe can be described by local field theory, to some degree of approximation. Indeed, one might argue that since we have tested local field theory up to energies of order \( e.g. \) 1 TeV, we know that the universe must contain at least \((1 \text{ TeV})^3 R_{\text{hor}}^3\) degrees of freedom \((R_{\text{hor}}\) is the horizon scale). This is incompatible with the holographic principle, but the argument is not really valid because the holographic degrees of freedom are related to local measurements in an extremely nonlocal way. We can only falsify the holographic idea by simultaneous measurements throughout the universe. If we tried to build a network of accelerators to make this test at the TeV scale, we would instead form a collection of black holes which swallowed up the apparatus.

I believe however that it is reasonable to claim that the degrees of freedom represented by the microwave background radiation today are independent. There is no problem of imagining a network of experiments to probe them, without forming black holes. We may ask, for what radius \( R_U \) of the universe are the microwave oscillators equal in number to a Planck density of degrees of freedom spread over an area of order \( R_U^2 \). The answer is \( R_U \sim 10^{35} R_{\text{hor}} \).

We can now try to calculate the cosmological constant. The total light cone energy of the system (which will be the cosmological term) is given by summing up the zero point fluctuations of these transverse degrees of freedom. We will also assume supersymmetry, broken at a scale \( \sqrt{F} \). The contributions of bosons and fermions with energies and transverse momenta greater that \( \sqrt{F} \) will cancel. The total energy is then

\[
E = F^{3/2} R_U^2. \tag{3.1}
\]

with everything measured in Planck units. But we claim that for this value of \( R_U \), the system can be reinterpreted as an effective local field theory spread over the volume. The cosmological term in the light cone energy is thus reinterpreted as the volume integral of a cosmological constant of order

\[
\Lambda \sim \frac{F^{3/2}}{R_U} \tag{3.2}
\]

\(^4\) The argument in the previous two paragraphs is an interpretation of remarks of L. Susskind.
If $\sqrt{F} \sim 10^{-8} M_P$, as we would expect in hidden sector models of SUSY breaking, then this gives a cosmological constant of order the current bound if $R_U \sim 10^{96} M_P^{-1}$. This is about $10^{35}$ times as large as the radius of our horizon volume, precisely the radius fixed by consistency between the microwave background entropy and a holographic picture. TeV scale SUSY breaking gives a cosmological constant $10^{-24}$ times smaller than the current bound, for the same radius $R_U$.

One amusing feature of this calculation is that it leads us to expect a nonzero value for the cosmological constant. This is of course favored by some fashionable cosmological models. In the above calculation, one obtains a phenomenologically interesting value of the cosmological constant only for hidden sector SUSY breaking.

There are many features of the preceding argument which are obscure. In particular, one may ask what is special about the present moment in cosmological history. Does our calculation imply that the cosmological constant and/or the number of degrees of freedom in the universe changes with time? I do not have any clear idea of how to answer either these or a host of other questions.

4. Patterns of SUSY Breaking

The model of SUSY breaking advocated in Section 1 implies that the SUSY breaking order parameter is the dilaton F term. In previous work on perturbative string theory such an assumption was shown to lead to a highly constrained pattern of SUSY breaking. Within the framework of [5] however, most of this predictive power is lost. Consider for example the squark masses. The Kahler potential will have terms of the form $Q_i^{\dagger} (M_Q^{2})^{ij} Q_j$ where $i$ and $j$ are generation indices. The matrix $M_Q^{2}$ will be dilaton dependent. Within the framework of [5] one can make no general statements about it in the coupling regime of interest. In particular, there are no apriori relations between the matrix evaluated at the dilaton VEV, and its mixed second derivative (wrt the dilaton) evaluated at the same point. These two matrices determine the squark wave function renormalization and the doublet squark masses respectively. Thus without imposing discrete symmetries, there is no reason for the squark mass matrix to be naturally proportional to the unit matrix. Similar remarks are valid for up and down squark mass matrices.
Next consider the $\mu$ and $B$ parameters of the supersymmetric standard model. The absence of the $\mu$ term in the tree level superpotential might be the consequence of a stringy symmetry, or of an R symmetry under which the operator $H_u H_d$ carries charge zero. In either case there is nothing to forbid a term $G H_u H_d + h.c.$, where $G$ is a function of the dilaton field, in the Kahler potential. The dilaton F term will then generate both a $\mu$ and a $B$ term of the right order of magnitude. They will be related to the first and second derivatives of $G$ at the dilaton VEV and there is no particular connection between them. One can also obtain a contribution to the $\mu$ term from a tree level coupling of $H_u H_d$ in the gauge kinetic function of the hidden gauge group$^{[19]}$.

If the symmetry which forbids $H_u H_d$ in the tree level superpotential is not an R-symmetry, or is an R-symmetry under which this operator carries nonzero charge, then the mechanism described above does not work. The $\mu$ and $B$ terms could only be obtained by coupling to a standard model singlet, which carries the relevant quantum number, in the tree level potential.

The only one of the weak coupling predictions of dilaton dominated SUSY breaking which survives in the scheme of $^3$ is the prediction of universal gaugino masses. The gauge kinetic functions are indeed unrenormalized apart from extremely small terms, and their dilaton dependence is determined at tree level.

It is also appropriate here to correct some statements that were made in $^3$ about the mass of the model independent axion. There it was claimed that the axion mass could naturally be very small. In fact, there is no principle known to me which could prevent the occurrence of a term in the Kahler potential of the form $e^{-c\sqrt{ReS}} f(S, S^*)$. According to the rules of $^3$, this term must be considered of order one. In general it will give the axion a mass of order the gravitino mass. No discrete symmetry can alter this conclusion. A gauged anomalous $U(1)$ forces the Kahler potential to be a function of only $ReS$, plus terms involving charged fields. However, in all known examples, the anomalous $U(1)$ is spontaneously broken to a discrete subgroup by the VEV of a charged field only slightly smaller than the string scale$^5$. Thus, the axion mass can at most be suppressed by a few

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$^5$ If this does not occur, the whole dilaton supermultiplet is Higgsed and disappears from the spectrum at a scale much higher than that of SUSY breaking. None of our discussion would be
powers of the charged field VEV over the string scale. It will never be light enough to serve as a QCD axion, or as dark matter.

5. Conclusions

Let us summarize our conclusions as a series of tentative answers to popular questions:

1. Why Is SUSY Broken? - The answer is essentially cosmological. Inflation is necessary to obtain a large universe. Postinflationary universes do not settle into vacua with negative cosmological constant. If R-symmetry is broken (and this happens in a very large class of strongly coupled supersymmetric gauge models even when SUSY is preserved) then supersymmetric vacua have negative cosmological constant. The criterion for zero cosmological constant is instead that SUSY be broken, and that there be a certain relation between the SUSY breaking scale and the R-symmetry breaking scale. In hidden sector models of SUSY breaking, this relation is automatically satisfied in order of magnitude.

In models of low energy SUSY breaking, the vanishing of the cosmological constant is more of a puzzle. It is traditional to resolve it by simply adding a constant to the superpotential. I do not think this is a satisfactory answer. A constant in the superpotential breaks R-symmetry, and we should understand the scales at which symmetries are broken in a dynamical manner, rather than putting them in by hand. In [13] a technically natural solution to this problem was proposed. The negative contribution to the potential from a strongly coupled supersymmetry preserving sector could have the same order of magnitude as the positive contribution from low energy SUSY breaking if the beta function of the SUSY preserving gauge group was 1.5 times as large as that of the SUSY breaking one. Although this is technically natural, it seems somewhat contrived. The “natural in order of magnitude” cancellation of the cosmological constant might be advanced as another argument which favors hidden sector SUSY breaking (recall that in models with low energy SUSY breaking, the dilaton is extremely light and dominates the energy density of the universe in an unacceptable manner).
2. How is the Dilaton Stabilized? - In string theory, any nonperturbative dynamics will produce a superpotential for the dilaton which is a sum of exponentials in the weak coupling region. The hierarchy of scales in nature is a consequence of these exponentials. In particular, we emphasize the hierarchy which is necessary to successful inflation: if the energy density during a large number of final e-foldings of inflation is not substantially smaller than the Planck scale, then large energy density fluctuations cause all of the matter in the universe to recollapse into a set of black holes. Thus, from a cosmological point of view, the dilaton must be stabilized in the weak coupling regime. Apart from cancellation between different exponentially small effects (which we eschew for religious reasons) there is one hypothetical mechanism for stabilizing the dilaton at weak coupling. If it is true that the perturbative estimate of the Kahler potential is wrong in the region where low energy field theory is weakly coupled, then it is also possible that the dilaton potential has a stable weak coupling minimum with zero cosmological constant even for a purely exponential superpotential. As above, this is only possible if SUSY is broken.

3. Why Do We Live in Four Dimensions? - A partial answer is given to this question by observing that in a finite universe the LEFTS is quantum mechanics on moduli space. In the region where more than four spacetime dimensions are large, it is plausible that this is a good approximation, since the correction to the action for homogeneous modes from integrating out other modes, including gravity, is under control in the infrared. Since the volume of moduli space is finite, the theory predicts a probability for large dimensions which falls like a large power of the size of the universe in string units. In four dimensions, this argument may be incorrect. If potentials are generated on moduli space by low energy dynamics, then we may have inflation. Thus, once the universe was a few orders of magnitude above the string scale, inflation would blow it up to very large size. We conclude that, the most probable large universes in string theory have four (or perhaps fewer) dimensions. Some arguments have been given\cite{12} that two dimensions are also improbable so the problem is reduced to understanding why we do not live in three dimensions.

Note that it has not yet been argued that inflation is highly probable. Indeed, the improbability of inflation in string theory might be part of the principle which selects the string vacuum state. That is, the reason that a particular vacuum is chosen may be because
it is the basin of attraction for one of the rare(?) places in moduli space where inflation occurs.

4. **How is the String Vacuum Chosen?** - Here the answer given in [13] is less satisfactory and depends on quantitative phenomenological input rather than the qualitative requirement that the universe be much larger than microscopic size (though note the speculation at the end of the previous answer). Data on the microwave background radiation suggest a potential energy density just a few orders of magnitude below the Planck scale, much larger than that allowed by SUSY breaking. Our most successful explanation of this discrepancy involves a superpotential generated by SUSY preserving strong dynamics at the high scale. This has the added advantage that most of the string moduli (but not the dilaton) will become massive at a high enough scale to avoid problems with reheating. The requirement that the low energy theory have such a SUSY preserving sector becomes a vacuum selection principle because of the observation that (as in 1 above) supersymmetry preserving dynamics can only lead to successful inflation if the superpotential satisfies a nongeneric condition. It must vanish at the minimum of the potential. In string theory this condition is quite restrictive and can only be satisfied at points in moduli space where chiral multiplets charged under the hidden sector gauge group become massless (without enlarging the group). Very few points satisfying this criterion are known (and as yet, none of them have anything resembling satisfactory phenomenology). It is not implausible to suggest that they are isolated points or lie in very low dimension subspaces of moduli space. This means that most of the moduli are massive at these stationary points.

Clearly this is at best the beginning of an answer to the question of vacuum selection. The vacuum selection principle we have invoked so far depends on fitting a quantitative aspect of the known universe, the amplitude of microwave background fluctuations. We would like to be able to show that all vacuum states other than the one which fits the detailed phenomenology of our world lead to clearcut cosmological disasters. Thus one would like to put only the barest phenomenological input into an explanation of vacuum selection. It may be that in the end we will have to accept a more equivocal answer to this question. Quantum string theory may predict a reasonable probability for many different kinds of large universe. In this case fundamental physics would inherit part of the
contingent, historical character of evolutionary biology. Certain features of the world would just be lucky accidents. One might attempt to invoke some sort of anthropic principle to distinguish our world, but I despair of ever being able to understand the physics of one of the “alternate universes” well enough to prove that no sort of intelligent life could ever evolve. In the context of string theory as we understand it today, the best hope for a clearcut vacuum selection principle seems to me to be the problem of the cosmological constant.

5. Why is the Cosmological Constant So Small? - For negative values of the cosmological constant, we have given an answer to this question, but for positive values I believe that the answer goes beyond the effective field theory approximation to string theory. Lowe et al. have shown that in the presence of horizons, the effective low energy theory of string theory is nonlocal, and contains states of long strings stretched along the horizon. This is true even when the spacetime curvature is small. Indeed it is true in flat space in appropriate coordinate systems (there is no contradiction with general covariance here if we believe that string theory is essentially nonlocal). We conclude that if the low energy field theory approximation to string theory predicts a positive cosmological constant, then it is not self-consistent. Truly stringy dynamics will have to be invoked to understand the dynamics of the universe on scales much larger than the DeSitter horizon. Note that this does not contradict the field theoretic treatment of a sufficiently short period of inflation. For a given value of the Hubble parameter, the field theoretic treatment will be valid up to some maximum number of e-foldings. At present I do not understand how to calculate this maximum number. Thus, we may claim that in string theory we are completely ignorant of the fate of vacuum states that are predicted to have positive cosmological constant by low energy field theory approximations.

Ignorance is bliss. The observation of Lowe et al. allows us to imagine that string theory will predict disaster for any vacuum state which, according to low energy field theory, is a DeSitter space. Disaster might mean recollapse, or some more exotic phenomenon which does not even have a local spacetime description. If this is the case, then the fine tuning problem of the cosmological constant could become a virtue. Only vacuum states which had very small cosmological constant would survive, and these would surely be few
and far between. Note the further constraint that the cosmological constant vanish in the weak coupling regime. A strongly coupled vacuum with vanishing cosmological constant would not have any obvious mechanism for producing a hierarchy. In particular, one would expect the scale of primordial $\delta \rho_\rho$ in such a vacuum (assuming it gives rise to some sort of inflation) to be $O(1)$. This leads to black hole formation and cosmological disaster. The idea that the vacuum is selected by requiring vanishing cosmological constant at weak coupling seems sufficiently powerful to pick out a unique state. Indeed, from the point of view of low energy field theory it seems surprising that there is any solution to this constraint.

However, we must remember that the work of Lowe et al. may be taken to support the conjecture[20] that string theory is holographic. If string theory is truly as nonlocal as the holographic principle implies, then our field theoretic understanding of a natural value for the cosmological constant is surely suspect. We presented a wildly speculative calculation of the cosmological constant based on these ideas. With hidden sector SUSY breaking, it gives a value close to the present observational bound.

5. What is The Pattern of Low Energy SUSY Breaking? - In our model for SUSY breaking, almost everything of interest is bound up with the as yet unknown Kahler potential of the low energy fields. The only clear cut prediction is that gaugino masses are universal. Flavor problems will have to be resolved by a horizontal symmetry. We noted that both the dilaton and axion have masses of order 1 TeV, so that the strong CP problem will have to be solved by a vanishing up quark mass. The nonperturbative Kahler potential will also change the predictions for coupling constant unification, in a manner which is at present uncomputable.

More definite phenomenological predictions wait a nonperturbative calculational scheme for string theory. It is possible that something could be done by resumming the perturbation expansion. If we could, by other means, decide on the correct point in moduli space to expand about, it would be worth the effort to try to compute several terms in the expansion of the Kahler potential. Furthermore one can probably also get a more precise determination of the large order behavior of the series than we have at present. Then, some variation of the Borel transform/conformal mapping techniques which have been successful in the theory of critical phenomena might be applied. This program sounds
like an enormous amount of work, but given a likely candidate for the ground state it is quite concrete, and perhaps feasible.

We also noted that low energy SUSY breaking models suffer from an as yet unresolved cosmological reheating problem. This is the consequence of the fact that in such models the dilaton is extremely light. Indeed it is probably light enough to be experimentally discovered in direct terrestrial measurements of the forces between objects. It would be extremely interesting to improve the bounds on the range of coherent forces that are consistent with such experiments.

The dilaton reheating problem is one of two indications that string theory will in the end lead us to a hidden sector mechanism for SUSY breaking. The other is the order of magnitude cancellation of the cosmological constant which is “automatic” in hidden sector theories and requires some tinkering in theories of low energy SUSY breaking.

To summarize, we have proposed cosmological answers to many of the fundamental questions facing superstring theory. The basic cosmological constraint is the existence of a vacuum state of the theory with vanishing cosmological constant, stable dilaton, and broken R-symmetry. SUSY breaking (probably via a hidden sector mechanism) is a consequence of these principles. A further constraint is that the vacuum lie in the basin of attraction of a portion of moduli space where inflation occurs. We have argued that very large numbers of e-foldings are not expected in the LEFTS, and that the low energy theory inevitably loses validity (i.e. is not a good approximation to string theory) when it predicts a large number of e-foldings. We do not yet have a sufficient quantitative understanding of the theory to determine whether the phenomenologically necessary 60 e-foldings is a large or small number from this point of view.

Indeed much of the detailed quantitative information about the theory is obscured in our approach by the necessity of invoking poorly understood nonperturbative phenomena to stabilize the dilaton in the weak coupling regime. This is an inevitable consequence of the Dine-Seiberg argument. Any truly predictive framework for string theory requires us to formulate and solve the theory in a nonperturbative manner.
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