

M Theory As A Matrix Model: A Conjecture

*T. Banks*¹, *W. Fischler*², *S.H. Shenker*¹, *L. Susskind*³

¹Department of Physics and Astronomy

Rutgers University, Piscataway, NJ 08855-0849

`banks, shenker@physics.rutgers.edu`

²Theory Group, Department of Physics

University of Texas, Austin, TX, 78712

`fischler@physics.utexas.edu`

³Department of Physics

Stanford University, Stanford, CA 94305-4060

`susskind@dormouse.stanford.edu`

We suggest and motivate a precise equivalence between uncompactified eleven dimensional M-theory and the $N = \infty$ limit of the supersymmetric matrix quantum mechanics describing D0 branes . The evidence for the conjecture consists of several correspondences between the two theories. As a consequence of supersymmetry the simple matrix model is rich enough to describe the properties of the entire Fock space of massless well separated particles of the supergravity theory. In one particular kinematic situation the leading large distance interaction of these particles is exactly described by supergravity . The model appears to be a nonperturbative realization of the holographic principle. The membrane states required by M-theory are contained as excitations of the matrix model. The membrane world volume is a noncommutative geometry embedded in a noncommutative spacetime.

October 1996

1. Introduction

M theory [1] is the strongly coupled limit of type IIA string theory. In the limit of infinite coupling it becomes an eleven dimensional theory in a background infinite flat space. In this paper M-theory will always refer to this decompactified limit. We know very little about this theory except for the following two facts. At low energy and large distances it is described by eleven dimensional supergravity . It is also known to possess membrane degrees of freedom with membrane tension $\frac{1}{l_p^3}$ where l_p is the eleven dimensional Planck length. It seems extremely unlikely that M-theory is any kind of conventional quantum field theory. The degrees of freedom describing the short distance behavior are simply unknown. The purpose of this paper is to put forward conjecture about these degrees of freedom and about the Hamiltonian governing them.

The conjecture grew out of a number of disparate facts about M-theory , D branes [2], matrix descriptions of their dynamics [3], supermembranes [4,5,6], the holographic principle [7], and short distance phenomena in string theory [8,9]. Simply stated the conjecture is this. M-theory , in the light cone frame, is exactly described by the large N limit of a particular supersymmetric matrix quantum mechanics. The system is the same one that has been used previously used to study the small distance behavior of D0 branes [9]. P.K. Townsend [10] was the first to point out that the supermatrix formulation of membrane theory suggested that membranes could be viewed as composites of D0 branes. Our work is a precise realization of his suggestion.

In what follows we will present our conjecture and some evidence for it. We begin by reviewing the description of string theory in the infinite momentum frame. We then present our conjecture for the full set of degrees of freedom of M-theory and the Hamiltonian which governs them. Our strongest evidence for the conjecture is a demonstration that our model contains the excitations which are widely believed to exist in M-theory , supergravitons and large metastable classical membranes. These are discussed in sections 3 and 5. The way in which these excitations arise is somewhat miraculous, and we consider this to be the core evidence for our conjecture. In section 4 we present a calculation of supergraviton scattering in a very special kinematic region, and argue that our model reproduces the expected result of low energy supergravity. The calculation depends on a supersymmetric nonrenormalization theorem whose validity we will discuss there. In section 6 we argue that our model may satisfy the Holographic Principle. This raises crucial issues about Lorentz invariance which are discussed there.

We emphasize that there are many unanswered questions about our proposed version of M-theory. Nonetheless these ideas seem of sufficient interest to warrant presenting them here. If our conjecture is correct, this would be the first nonperturbative formulation of a quantum theory which includes gravity.

2. The infinite momentum frame and the Holographic Principle

The infinite momentum frame [11] is the old name for the misnamed light cone frame. Thus far this is the only frame in which it has proved possible to formulate string theory in Hamiltonian form. The description of M-theory which we will give in this paper is also in the infinite momentum frame. We will begin by reviewing some of the features of the infinite momentum frame formulation of relativistic quantum mechanics. For a comprehensive review we refer the reader to [11]. We begin by choosing a particular spatial direction x^{11} called the longitudinal direction. The nine dimensional space transverse to x^{11} is labeled x^i or x^\perp . Time will be indicated by t . Now consider a system of particles with momenta (p_\perp^a, p_{11}^a) where a labels the particle. The system is boosted along the x^{11} axis until all longitudinal momenta are much larger than any scale in the problem. Further longitudinal boosting just rescales all longitudinal momenta in fixed proportion. Quantum field theory in such a limiting reference frame has a number of properties which will be relevant to us.

It is convenient to begin by assuming the x^{11} direction is compact with a radius R . The compactification serves as an infrared cutoff. Accordingly, the longitudinal momentum of any system or subsystem of quanta is quantized in units of $1/R$. In the infinite momentum frame all systems are composed of constituent quanta or partons. The partons all carry strictly positive values of longitudinal momentum. It is particularly important to understand what happens to quanta of negative or vanishing p_{11} . The answer is that as the infinite momentum limit is approached the frequency of these quanta, relative to the Lorentz time dilated motion of the boosted system becomes infinite and the zero and negative momentum quanta may be integrated out. The process of integrating out such fast modes may influence or even determine the Hamiltonian of the remaining modes. In fact the situation is slightly more complicated in certain cases for the zero momentum degrees of freedom. In certain situations such as spontaneous symmetry breaking, these longitudinally homogeneous modes define backgrounds whose moduli may appear in the

Hamiltonian of the other modes. In any case the zero and negative momentum modes do not appear as independent dynamical degrees of freedom.

Thus we may assume all systems have longitudinal momentum given by an integer multiple of $1/R$

$$p_{11} = N/R \quad (2.1)$$

with N strictly positive. At the end of a calculation we must let R and N/R tend to infinity to get to the uncompactified infinite momentum limit.

The main reason for the simplifying features of the infinite momentum frame is the existence of a transverse Galilean symmetry which leads to a naive nonrelativistic form for the equations. The role of nonrelativistic mass is played by the longitudinal momentum p_{11} . The Galilean transformations take the form

$$p_i \rightarrow p_i + p_{11} v_i \quad (2.2)$$

As an example of the Galilean structure of the equations, the energy of a free massless particle is

$$E = \frac{p_{\perp}^2}{2p_{11}} \quad (2.3)$$

For the eleven dimensional supersymmetric theory we will consider, the Galilean invariance is extended to the Supergalilean group which includes 32 real supergenerators. The supergenerators divide into two groups of 16, each transforming as spinors under the nine dimensional transverse rotation group. We denote them by Q_{α} and q_A and they obey anticommutation relations

$$\begin{aligned} [Q_{\alpha}, Q_{\beta}]_+ &= \delta_{\alpha\beta} H \\ [q_A, q_B]_+ &= \delta_{AB} P_{11} \\ [Q_{\alpha}, q_A] &= \gamma_{A\alpha}^i P_i \end{aligned} \quad (2.4)$$

The Lorentz generators which do not preserve the infinite momentum frame mix up the two kinds of generators.

Let us now recall some of the features of string theory in the infinite momentum or light cone frame [12]. We will continue to call the longitudinal direction x^{11} even though in this case the theory has only ten space-time directions. The transverse space is of course eight dimensional. To describe a free string of longitudinal momentum p_{11} a periodic parameter σ which runs from 0 to p_{11} is introduced. To regulate the world sheet theory a cutoff $\delta\sigma = \epsilon$ is introduced. This divides the parameter space into $N = p_{11}/\epsilon$ segments,

each carrying longitudinal momentum ϵ . We may think of each segment as a parton but unlike the partons of quantum field theory these objects always carry $p_{11} = \epsilon$. For a multi particle system of total longitudinal momentum $p_{11}(total)$ we introduce a total parameter space of overall length $p_{11}(total)$ which we allow to be divided into separate pieces, each describing a string. The world sheet regulator is implemented by requiring each string to be composed of an integer number of partons of momentum ϵ . Interactions are described by splitting and joining processes in which the number of partons is strictly conserved. The regulated theory is thus seen to be a special case of Galilean quantum mechanics of N partons with interactions which bind them into long chains and allow particular kinds of rearrangements.

The introduction of a minimum unit of momentum ϵ can be given an interpretation as an infrared cutoff. In particular we may assume that the x^{11} coordinate is periodic with length $R = \epsilon^{-1}$. Evidently the physical limit $\epsilon \rightarrow 0, R \rightarrow \infty$ is a limit in which the number of partons N tends to infinity.

It is well known [7] that in this large N limit the partons become infinitely dense in the transverse space and that this leads to extremely strong interactions. This circumstance together with the Bekenstein bound on entropy, has led to the *holographic* speculation that the transverse density of partons is strictly bounded to about one per transverse Planck area. In other words the partons form a kind of incompressible fluid. This leads to the unusual consequence that the transverse area occupied by a system of longitudinal momentum p_{11} can not be smaller than p_{11}/ϵ in Planck units.

The general arguments for the holographic behavior of systems followed from considerations involving the Bekenstein - 't Hooft bound on the entropy of a spatial region [13] and were not specific to string theory. If the arguments are correct they should also apply to 11 dimensional theories which include gravitation. Thus we should expect that in M-theory the radius of a particle such as the graviton will grow with p_{11} according to

$$\rho = \left(\frac{p_{11}}{\epsilon}\right)^{\frac{1}{9}} l_p = (p_{11} R)^{\frac{1}{9}} l_p \quad (2.5)$$

where l_p is the eleven dimensional Planck length. In what follows we will see quantitative evidence for exactly this behavior.

At first sight the holographic growth of particles appears to contradict the boost invariance of particle interactions. Consider the situation of two low energy particles moving past one another with some large transverse separation, let us say of order a

meter. Obviously these particles have negligible interactions. Now boost the system along the longitudinal direction until the size of each particle exceeds their separation. They now overlap as they pass each other. But longitudinal boost invariance requires that the scattering amplitude is still essentially zero. This would seem to require extremely special and unnatural cancellations. We will see below that one key to this behavior is the very special BPS property of the partons describing M-theory. However, we are far from having a complete understanding of the longitudinal boost invariance of our system. Indeed, we view it as the key dynamical puzzle which must be unravelled in understanding the dynamics of M-theory.

3. M-theory and D0 branes

M-theory with a compactified longitudinal coordinate x^{11} is by definition type IIA string theory. The correspondences between the two theories include [1] :

1. The compactification radius R is related to the string coupling constant by

$$R = g^{2/3} l_p = g l_s \quad (3.1)$$

where l_s is the string length scale

$$l_s = g^{-1/3} l_p \quad (3.2)$$

2. The Ramond-Ramond photon of IIA theory is the Kaluza Klein photon which arises upon compactification of eleven dimensional supergravity.
3. No perturbative string states carry RR charge. In other words all perturbative string states carry vanishing momentum along the x^{11} direction. The only objects in the theory which do carry RR photon charge are the D0 branes of Polchinski. D0 branes are point particles which carry a single unit of RR charge. Equivalently they carry longitudinal momentum

$$P_{11}^{D0 \text{ branes}} = 1/R \quad (3.3)$$

The D0 branes carry the quantum numbers of the first massive KK modes of the basic eleven dimensional supergravity multiplet, including 44 gravitons, 84 components of a 3-form, and 128 gravitinos. We will refer to these particles as supergravitons. As 11-dimensional objects these are all massless. As a consequence they are BPS saturated states in the 10-dimensional theory. Their 10-D mass is $1/R$.

4. Supergravitons carrying Kaluza Klein momentum $p_{11} = N/R$ also exist but are not described as elementary D0 branes. As shown in [3] their proper description is as bound composites of N D0 branes.

These properties make the D0 branes candidate partons for an infinite momentum limit description of M-theory. We expect that if, as in quantum field theory, the degrees of freedom with vanishing and negative p_{11} decouple then M-theory in the infinite momentum frame should be a theory whose only degrees of freedom are D0 branes. Anti-D0 branes carry negative Kaluza Klein momenta and strings carry vanishing p_{11} . The decoupling of anti-D0 branes is particularly fortunate because brane anti-brane dynamics is something about which we know very little [14]. The BPS property of zero branes ameliorates the conflict between infinitely growing parton wave functions and low energy locality which we noted at the end of the last section. We will see some partial evidence for this in a nontrivial scattering computation below. We will also discuss below the important point that a model containing only D0 branes actually contains large classical supermembrane excitations. Since the conventional story of the M theoretic origin of strings depicts them as membranes wrapped around the compactified eleventh dimension, we have some reason to believe that strings have not really been left out of the system.

All of these circumstances lead us to propose that M-theory in the infinite momentum frame is a theory in which the only dynamical degrees of freedom or partons are D0 branes. Furthermore it is clear in this case that all systems are built out of composites of partons, each of which carries the minimal p_{11} . We note however that our system does have a set of degrees of freedom which go beyond the parton coordinates. Indeed, as first advocated in [3], the D0 brane coordinates of N partons have to be promoted to matrices. At distance scales larger than the eleven dimensional Planck scale, these degrees of freedom become very massive and largely decouple¹, but their virtual effects are responsible for all parton interactions. These degrees of freedom are BPS states, and are related to the parton coordinates by gauge transformations. Furthermore, when the partons are close together they become low frequency modes. Thus they cannot be omitted in any discussion of the dynamics of D0 branes .

¹ Indeed we will propose that this decoupling is precisely what defines the regime in which the classical notion of distance makes sense.

4. D0 brane Mechanics

If the infinite momentum limit of M-theory is the theory of D0 branes, decoupled from the other string theory degrees of freedom, what is the precise form of the quantum mechanics of the system? Fortunately there is a very good candidate which has been extensively studied in another context in which D0 branes decouple from strings [9].

As emphasized at the end of the last section, open strings which connect D0 branes do not exactly decouple. In fact the very short strings which connect the branes when they are practically on top of each other introduce a new kind of coordinate space in which the nine spatial coordinates of a system of N D0 branes become nine $N \times N$ matrices, $X_{a,b}^i$ [3]. The matrices X are accompanied by 16 fermionic superpartners $\theta_{a,b}$ which transform as spinors under the $SO(9)$ group of transverse rotations. The matrices may be thought of as the spatial components of 10 dimensional Super Yang Mills (SYM) fields after dimensional reduction to zero space directions. These Yang Mills fields describe the open strings which are attached to the D0 branes. The Yang Mills quantum mechanics has $U(N)$ symmetry and is described (in units with $l_s = 1$) by the Lagrangian

$$L = \frac{1}{2g} \left[\text{tr} \dot{X}^i \dot{X}^i + 2\theta^T \dot{\theta} - \frac{1}{2} \text{tr} [X^i, X^j]^2 - 2\theta^T \gamma_i [\theta, X^i] \right] \quad (4.1)$$

Here we have used conventions in which the fermionic variables are 16 component nine dimensional spinors.

In [9] this Lagrangian was used to study the short distance properties of D0 branes in weakly coupled string theory. The 11-D Planck length emerged as a natural dynamical length scale in that work, indicating that the system (4.1) describes some M-theoretic physics. In [9] (4.1) was studied as a low velocity effective theory appropriate to the heavy D0 branes of weakly coupled string theory. Here we propose (4.1) as the most general infinite momentum frame Lagrangian, with at most two derivatives which is invariant under the gauge symmetry and the Supergalilean group [15]². It would be consistent with our assumption that matrix D0 branes are the only degrees of freedom of M-theory, to write a Lagrangian with higher powers of first derivatives. We do not know if any such Lagrangians exist which preserve the full symmetry of the infinite momentum frame. What is at issue

² The gauge invariance is in fact necessary to supertranslation invariance. The supergenerators close on gauge transformations and only satisfy the supertranslation algebra on the gauge invariant subspace.

here is eleven dimensional Lorentz invariance. In typical infinite momentum frame field theories, the naive classical Lagrangian for the positive longitudinal momentum modes is renormalized by the decoupled infinite frequency modes. The criterion which determines the infinite momentum frame Lagrangian is invariance under longitudinal boosts and null plane rotating Lorentz transformations (the infamous angular conditions). Apart from simplicity, our main reason for suggesting the Lagrangian (4.1) is that we have found some partial evidence that the large N limit of the quantum theory it defines is indeed Lorentz invariant. A possible line of argument systematically leading to (4.1) is discussed in Section 9.

Following [9], let us rewrite the action in units in which the 11-D Planck length is 1. Using (3.1) and (3.2) the change of units is easily made and one finds

$$L = \text{tr} \left[\frac{1}{2R} D_t Y^i D_t Y^i - \frac{1}{4} R [Y^i, Y^j]^2 - \theta^T D_t \theta - R \theta^T \gamma_i [\theta, Y^i] \right] \quad (4.2)$$

where $Y = \frac{X}{g^{1/3}}$. We have also changed the units of time to 11-D Planck units. We have restored the gauge field ($\partial_t \rightarrow D_t = \partial_t + iA$) to this expression (previously we were in $A = 0$ gauge) in order to emphasize that the SUSY transformation laws (here ϵ and ϵ' are two independent 16 component anticommuting SUSY parameters),

$$\delta X^i = -2\epsilon^T \gamma^i \theta \quad (4.3)$$

$$\delta \theta = \frac{1}{2} \left[D_t X^i \gamma_i + \gamma_- + \frac{1}{2} [X^i, X^j] \gamma_{ij} \right] \epsilon + \epsilon' \quad (4.4)$$

$$\delta A = -2\epsilon^T \theta \quad (4.5)$$

involve a gauge transformation. As a result, the SUSY algebra closes on the gauge generators, and only takes on the form (2.4) when applied to gauge invariant states.

The Hamiltonian has the form

$$H = R \text{tr} \left\{ \frac{\Pi_i \Pi_i}{2} + \frac{1}{4} [Y^i, Y^j]^2 + \theta^T \gamma_i [\theta, Y^i] \right\} \quad (4.6)$$

where Π is the canonical conjugate to Y . Note that in the limit $R \rightarrow \infty$, all finite energy states of this Hamiltonian have infinite energy. We will be interested only in states whose energy vanishes like $\frac{1}{N}$ in the large N limit, so that this factor becomes the inverse power of longitudinal momentum which we expect for the eigenstates of a longitudinal boost invariant system. Thus, in the correct infinite momentum frame limit, the only relevant

asymptotic states of the Hamiltonian should be those whose energy is of order $\frac{1}{N}$. We will exhibit a class of such states below, the supergraviton scattering states. The difficult thing will be to prove that their S-matrix elements depend only on ratios of longitudinal momenta, so that they are longitudinally boost invariant.

To understand how this system represents ordinary particles we note that when the Y 's become large the commutator term in H becomes very costly in energy. Therefore for large distances the finite energy configurations lie on the flat directions along which the commutators vanish. In this system with 16 supercharges,³ these classical zero energy states are in fact exact supersymmetric states of the system. In contrast to field theory, the continuous parameters which describe these states (the Higgs VEVs in the language of SYM theory) are not vacuum superselection parameters, but rather collective coordinates. We must compute their quantum wave functions rather than freeze them at classical values. They are however the slowest modes in the system, so that we can integrate out the other degrees of freedom to get an effective SUSY quantum mechanics of these modes alone. We will study some aspects of this effective dynamics below.

Along the flat directions the Y^i are simultaneously diagonalizable. The diagonal matrix elements are the coordinates of the D0 branes. When the Y are small the cost in energy for a noncommuting configuration is not large. Thus for small distances there is no interpretation of the configuration space in terms of ordinary positions. Classical geometry and distance are only sensible concepts in this system in regions far out along one of the flat directions. We will refer to this as the long distance regime. In the short distance regime, we have a noncommutative geometry. Nevertheless the full Hamiltonian (4.6) has the usual Galilean symmetry. To see this we define the center of mass of the system by

$$Y(c.m.) = \frac{1}{N} \text{tr} Y \quad (4.7)$$

A transverse translation is defined by adding a multiple of the identity to Y . This has no effect on the commutator term in L because the identity commutes with all Y . Similarly rotational invariance is manifest.

The center of mass momentum is given by

$$P(c.m.) = \text{tr} \Pi = \frac{N}{R} \dot{Y}(c.m.) \quad (4.8)$$

³ The sixteen supercharges which anticommute to the longitudinal momentum act only on the center of mass of the system and play no role in particle interactions.

Using $p_{11} = N/R$ gives the usual connection between transverse velocity and transverse momentum

$$\frac{1}{p_{11}}P(c.m.) = \dot{Y}(c.m.) \quad (4.9)$$

A Galilean boost is defined by adding a multiple of the identity to \dot{Y} . We leave it to the reader to show that this has no effect on the equations of motion. This establishes the Galilean invariance of H . The Supergalilean invariance is also completely unbroken. The alert reader may be somewhat unimpressed by some of these invariances, since they appear to be properties of the center of mass coordinate, which decouples from the rest of the dynamics. Their real significance will appear below when we show that our system possesses multiparticle asymptotic states, on which these generators act in the usual way as a sum of single particle operators.

5. A Conjecture

Our conjecture is thus that M-theory formulated in the infinite momentum frame is exactly equivalent to the $N \rightarrow \infty$ limit of the supersymmetric quantum mechanics described by the Hamiltonian (4.6). The calculation of any physical quantity in M-theory can be reduced to a calculation in matrix quantum mechanics followed by an extrapolation to large N . In what follows we will offer evidence for this surprising conjecture.

Let us begin by examining the single particle spectrum of the theory. For $N = 1$ the states with $p_{\perp} = 0$ are just those of a single D0 brane at rest. The states form a representation of the algebra of the 16 θ 's with 2^8 components. These states have exactly the quantum numbers of the 256 states of the supergraviton. For non-zero p_{\perp} the energy of the object is

$$E = \frac{R}{2}p_{\perp}^2 = \frac{p_{\perp}^2}{2p_{11}} \quad (5.1)$$

For states with $N > 1$ we must study the $U(N)$ invariant Schroedinger equation arising from (4.6). H can easily be separated in terms of center of mass and relative motions. Define

$$Y = \frac{Y(c.m.)}{N}I + Y_{rel} \quad (5.2)$$

where I is the unit matrix and Y_{rel} is a traceless matrix in the adjoint of $SU(N)$ representing relative motion. The Hamiltonian then has the form

$$H = H_{c.m.} + H_{rel} \quad (5.3)$$

with

$$H_{c.m.} = \frac{P(c.m.)^2}{2p_{11}} \quad (5.4)$$

The Hamiltonian for the relative motion is the dimensional reduction of the supersymmetric 10-D Yang Mills Hamiltonian. Although a direct proof based on the Schroedinger equation has not yet been given, duality between IIA strings and M-theory requires the relative Schroedinger equation to have normalizable threshold bound states with zero energy [3]. The bound system must have exactly the quantum numbers of the 256 states of the supergraviton. For these states the complete energy is given by (5.4). Furthermore these states are BPS saturated. No other normalizable bound states can occur. Thus we find that the spectrum of stable single particle excitations of (4.6) is exactly the supergraviton spectrum with the correct dispersion relation to describe massless 11 dimensional particles in the infinite momentum frame.

Next let us turn to the spectrum of widely separated particles. That a simple quantum mechanical Hamiltonian like (4.6) should be able to describe arbitrarily many well separated particles is not at all evident and would certainly be impossible without the special properties implied by supersymmetry. Begin by considering commuting block diagonal matrices of the form

$$\begin{pmatrix} Y_1^i & 0 & 0 & \dots \\ 0 & Y_2^i & 0 & \dots \\ 0 & 0 & Y_3^i & \dots \\ \vdots & & & \end{pmatrix} \quad (5.5)$$

where Y_a^i are $N_a \times N_a$ matrices and $N_1 + N_2 + \dots + N_n = N$. For the moment suppose all other elements of the Y 's are constrained to vanish identically. In this case the Schroedinger equation obviously separates into n uncoupled Schroedinger equations for the individual block degrees of freedom. Each equation is identical to the original Schroedinger equation for the system of N_a D0 branes. Thus the spectrum of this truncated system includes collections of noninteracting supergravitons.

Now let us suppose the supergravitons are very distant from one another. In other words for each pair the relative distance, defined by

$$R_{a,b} = \left| \frac{\text{tr} Y_a}{N_a} - \frac{\text{tr} Y_b}{N_b} \right| \quad (5.6)$$

is asymptotically large. In this case the commutator terms in (4.6) cause the off diagonal blocks in the Y 's to have very large potential energy proportional to $R_{a,b}^2$. This effect

can also be thought of as the Higgs effect giving mass to the broken generators (“W-bosons”) of $U(N)$ when the symmetry is broken to $U(N_1) \times U(N_2) \times U(N_3) \dots$. Thus one might naively expect the off diagonal modes to leave the spectrum of very widely separated supergravitons unmodified. However this is not correct in a generic situation. The off diagonal modes behave like harmonic oscillators with frequency of order $R_{a,b}$ and their zero point energy will generally give rise to a potential energy of similar magnitude. This effect would certainly preclude an interpretation of the matrix model in terms of well separated independent particles.

Supersymmetry is the ingredient which rescues us. In a well known way the fermionic partners of the off diagonal bosonic modes exactly cancel the potential due to the bosons, leaving exactly flat directions. We know this from the nonrenormalization theorems for supersymmetric quantum mechanics with 16 supergenerators [15]. The effective Lagrangian for the collective coordinates along the flat directions must be supersymmetric and the result of [15] guarantees that up to terms involving at most two derivatives, the Lagrangian for these coordinates must be the dimensional reduction of $U(1)^n$ SYM theory, where n is the number of blocks (*i.e.* the number of supergravitons). This is just the Lagrangian for free motion of these particles. Furthermore, since we are doing quantum mechanics, and the analog of the Yang Mills coupling is the dimensional quantity l_p^3 , the coefficient of the quadratic term is uncorrected from its value in the original Lagrangian.

There are residual virtual effects at order p^4 from these heavy states which are the source of parton interactions. Note that the off diagonal modes are manifestly nonlocal. The apparent locality of low energy physics in this model must emerge from a complex interplay between SUSY and the fact that the frequencies of the nonlocal degrees of freedom become large when particles are separated. We have only a limited understanding of this crucial issue, but in the next section we will provide some evidence that local physics is reproduced in the low energy, long distance limit.

The center of mass of a block of size $N^{(a)}$ is defined by equation (5.2). It is easy to see that the Hamiltonian for an asymptotic multiparticle state, when written in terms of center of mass transverse momenta, is just

$$H_{asympt} = \sum_a \frac{R \mathbf{P}^{(a)2}}{N^{(a)}} = \sum_a \frac{\mathbf{P}^{(a)2}}{p_{11}^{(a)}} \quad (5.7)$$

Note that the dispersion relation for the asymptotic particle states has the fully eleven dimensional Lorentz invariant form. This is essentially due to the BPS nature of the

asymptotic states. For large relative separations, the supersymmetric quantum state corresponding to the supersymmetric classical flat direction in which the gauge symmetry is “broken” into n blocks, will be precisely the product of the threshold boundstate wave functions of each block subsystem. Each individual block is a BPS state. Its dispersion relation follows from the SUSY algebra and is relativistically invariant even when (*e.g.* for finite N) the full system is not.

We also note that the statistics of multi supergraviton states comes out correctly because of the residual block permutation gauge symmetry of the matrix model. When some subset of the blocks are in identical states, the original gauge symmetry instructs us to mod out by the permutation group, picking up minus signs depending on whether the states are constructed from an odd or even number of Grassmann variables. The spin statistics connection is the conventional one.

Thus, the large N matrix model contains the Fock space of asymptotic states of eleven dimensional supergravity, and the free propagation of particles is described in a manner consistent with eleven dimensional Lorentz invariance. The field theory Fock space is however embedded in a system which, as we shall see, has no ultraviolet divergences. Particle statistics is embedded in a continuous gauge symmetry. We find the emergence of field theory as an approximation to an elegant finite structure one of the most attractive features of the matrix model approach to M-theory.

6. Long Range Supergraviton Interactions

The first uncanceled interactions in the matrix model occur in the effective action at order \dot{y}^4 where \dot{y} is the velocity of the supergravitons [9]. These interactions are calculated by thinking of the matrix model as SYM theory and computing Feynman diagrams. At one loop one finds an induced quartic term in the velocities which corresponds to an induced $F_{\mu\nu}^4$ term. The precise term for two D0 branes is given by

$$\frac{A[\dot{y}(1) - \dot{y}(2)]^4}{R^3 r^7} \quad (6.1)$$

where r is the distance between the D0 branes and A is a coefficient of order one which can be extracted from the results of [9]. This is the longest range term which governs the interaction between the D0 branes as r tends to infinity. Thus the effective Lagrangian governing the low energy long distance behavior of the pair is

$$L = \frac{\dot{y}(1)^2}{2R} + \frac{\dot{y}(2)^2}{2R} - A \frac{[\dot{y}(1) - \dot{y}(2)]^4}{R^3 r^7} \quad (6.2)$$

The calculation is easily generalized to the case of two well separated groups of N_1 and N_2 D0 branes forming bound states. Keeping only the leading terms for large N (planar graphs) we find

$$L = \frac{N_1 \dot{y}(1)^2}{2R} + \frac{N_2 \dot{y}(2)^2}{2R} - AN_1 N_2 \frac{[\dot{y}(1) - \dot{y}(2)]^4}{R^3 r^7} \quad (6.3)$$

To understand the significance of (6.3) it is first useful to translate it into an effective Hamiltonian. To leading order in inverse powers of r we find

$$H_{eff} = \frac{p_{\perp}(1)^2}{2p_{11}(1)} + \frac{p_{\perp}(2)^2}{2p_{11}(2)} + A \left[\frac{p_{\perp}(1)}{p_{11}(1)} - \frac{p_{\perp}(2)}{p_{11}(2)} \right]^4 \frac{p_{11}(1)p_{11}(2)}{r^7 R} \quad (6.4)$$

From (6.4) we can compute a scattering amplitude in Born approximation. Strictly speaking the scattering amplitude is defined as a 10-D amplitude in the compactified theory. However it contains information about the 11-D amplitude in the special kinematic situation where no longitudinal momentum is exchanged. The relation between the 10-D amplitude and 11-D amplitude at vanishing Δp_{11} is simple. They are essentially the same thing except for a factor of R which is needed to relate the 10 and 11 dimensional phase space volumes. The relation between amplitudes is $A_{11} = RA_{10}$. Thus from (6.4) we find the 11-D amplitude

$$A \left[\frac{p_{\perp}(1)}{p_{11}(1)} - \frac{p_{\perp}(2)}{p_{11}(2)} \right]^4 \frac{p_{11}(1)p_{11}(2)}{r^7} \quad (6.5)$$

The expression in (6.5) is noteworthy for several reasons. First of all the factor r^{-7} is the 11-D green function (in space) after integration over x^{11} . In other words it is the scalar Feynman propagator for vanishing longitudinal momentum transfer. Somehow the simple matrix Hamiltonian knows about massless propagation in 11-D spacetime. The remaining momentum dependent factors are exactly what is needed to make (6.5) identical to the single (super) graviton exchange diagram in 11-D.⁴ Even the coefficient is correct. This is closely related to a result reported in [9] where it was shown that the annulus diagram governing the scattering of two D0 branes has exactly the same form at very

⁴ In [9] the amplitude was computed for D0 branes which have momenta orthogonal to their polarizations (this was not stated explicitly there but was implicit in the choice of boundary state). The spin dependence of the amplitude is determined by the supersymmetric completion of the $\frac{v^4}{r^7}$ amplitude, which we have not computed. In principle, this gives another check of eleven dimensional Lorentz invariance. We suspect that the full answer follows by applying the explicit supersymmetries of the light cone gauge to the amplitude we have computed.

small and very large distances, which can be understood by noting that only BPS states contribute to this process on the annulus. This, plus the usual relations between couplings and scales in Type IIA string theory and M-theory guarantee that we obtain the correct normalization of eleven dimensional graviton scattering in SUGRA. In the weak coupling limit, very long distance behavior is governed by single supergraviton exchange while the ultra short distances are governed by the matrix model. In [9] the exact equivalence between the leading interactions computed in these very different manners was recognized but its meaning was not clear. Now we see that it is an important consistency criterion in order for the matrix model to describe the infinite momentum limit of M-theory.

Let us next consider possible corrections to the effective action coming from higher loops. In particular, higher loops can potentially correct the quartic term in velocities. Since our interest lies in the large N limit we may consider the leading (planar) corrections. Doing ordinary large N counting one finds that the \dot{y}^4 term may be corrected by a factor which is a function of the ratio N/r^3 . Such a renormalization by $f(N/r^3)$ could be dangerous. We can consider several cases which differ in the behavior of f as $\frac{N}{r^3}$ tends to infinity. In the first two the function tends to zero or infinity. The meaning of this would be that the coupling to gravity is driven either to zero or infinity in the infinite momentum limit. Either behavior is intolerable. Another possibility is that the function f tends to a constant not equal to 1. In this case the gravitational coupling constant is renormalized by a constant factor. This is not supposed to occur in M-theory. Indeed supersymmetry is believed to protect the gravitational coupling from any corrections. The only other possibility is that $f \rightarrow 1$. The simplest way in which this can happen is if there are no corrections at all other than the one loop term which we have discussed.

We believe that there is a nonrenormalization theorem for this term which can be proven in the context of SUSY quantum mechanics with 16 generators. The closest thing we have been able to find in the literature is a nonrenormalization theorem for the $F_{\mu\nu}^4$ term in the action of ten dimensional string theory⁵ which has been proven by Tseytlin [16]. In the quantum mechanical context, we believe that it is true and that the scattering of two supergravitons at large transverse distance and zero longitudinal momentum is exactly given in the matrix model by low energy 11-D supergravity perturbation theory. M. Dine[17] has constructed the outlines of an argument which demonstrates the validity of the nonrenormalization theorem.

⁵ We thank C. Bachas for pointing out this theorem to us.

We have considered amplitudes in which vanishing longitudinal momentum is exchanged. Amplitudes with nonvanishing exchange of p_{11} are more complicated. They correspond to processes in 10-D in which RR charge is exchanged. Such collisions involve rearrangements of the D0 branes in which the collision transfers D0 branes from one group to the other. We are studying such processes but we have no definitive results as yet.

We have thus presented some evidence that the dynamics of the matrix model respects eleven dimensional Lorentz invariance. If this is correct then the model reduces exactly to supergravity at low energies. It is clear however that it is much better behaved in the ultraviolet than a field theory. At short distances, as shown extensively in [9] restoration of the full matrix character of the variables cuts off all ultraviolet divergences. The correspondence limit by which M-theory reduces to supergravity indicates that we are on the right track.

7. The Size of a Supergraviton

As we have pointed out in section 2 the holographic principle requires the transverse size of a system to grow with the number of constituent partons. It is therefore of interest to estimate the size of the threshold bound state describing a supergraviton of longitudinal momentum N/R . According to the holographic principle the radius should grow like $N^{1/9}$ in 11-D Planck units. We will use a mean field approximation in which we study the wave function of one parton in the field of N others. We therefore consider the effective Lagrangian (6.3) for the case $N_1 = 1, N_2 = N$. The action simplifies for $N \gg 1$ since in this case the N particle system is much heavier than the single particle. Therefore we may set its velocity to zero. The Lagrangian becomes

$$L_1 = \frac{\dot{y}^2}{2R} - N \frac{\dot{y}^4}{R^3 y^7} \quad (7.1)$$

where y refers to the relative coordinate between the two systems. We can remove all N and R dependence from the action $S = \int L_1 dt$ by scaling

$$\begin{aligned} y &\rightarrow N^{1/9} y \\ t &\rightarrow \frac{N^{2/9}}{R} t . \end{aligned} \quad (7.2)$$

The characteristic length, time and velocity ($v = \dot{y}$) scales are

$$\begin{aligned}
y &\sim N^{1/9} \\
t &\sim \frac{N^{2/9}}{R} \\
v &\sim \frac{y}{t} \sim N^{-1/9} R .
\end{aligned}
\tag{7.3}$$

That the size of the bound state wave function scales like $N^{1/9}$ is an indication of the incompressibility of the system when it achieves a density of order one degree of freedom per Planck area. This is in accord with the holographic principle.

This mean field picture of the bound state, or any other description of it as a simple cluster, makes the problem of longitudinal boost invariance mentioned earlier very concrete. Suppose we consider the scattering of two bound states with N_1 and N_2 constituents respectively, $N_1 \sim N_2 \sim N$. The mean field picture strongly suggests that scattering will show a characteristic feature at an impact parameter corresponding to the bound state size $\sim N^{1/9}$. But this is not consistent with longitudinal boosts which take $N_1 \rightarrow \alpha N_1$, $N_2 \rightarrow \alpha N_2$. Boost invariance requires physics to depend only on the ratio N_1/N_2 , or said another way, only on the ratio of the bound state sizes. This strongly suggests that a kind of scale invariance must be present in the dynamics that is clearly absent in the simple picture discussed above. In the string case the scale invariant world sheet dynamics is crucial for longitudinal boost invariance.

The possibility that partons might form subclusters within the bound state was ignored in mean field discussion. A preliminary discussion of a hierarchical clustering model with many length scales is presented in the Appendix. Note also that wavefunctions of threshold bound states are power law behaved.

Understanding the dynamics of these bound states well enough to check longitudinal boost invariance reliably is an important subject for future research.

8. Membranes

In order to be the strong coupling theory of IIA string theory, M-theory must have membranes in its spectrum. Although in the decompactified limit there are no truly stable finite energy membranes, very long lived large classical membranes must exist. In this section we will show how these membranes are described in the matrix model, a result

first found in [5]⁶. Townsend[10] first pointed out the connection between the matrix description of D0 brane dynamics and the matrix description of membranes, and speculated that a membrane might be regarded as a collective excitation of D0 branes. Our conjecture supplies a precise realization of Townsend's idea.

The formulation we will use to describe this connection is a version of the methods introduced in [5,18,19].

Begin with a pair of unitary operators U, V satisfying the relations

$$\begin{aligned} UV &= e^{\frac{2\pi i}{N}} VU \\ U^N &= 1 \\ V^N &= 1. \end{aligned} \tag{8.1}$$

These operators can be represented on an N dimensional Hilbert space as clock and shift operators. They form a basis for all operators in the space. Any matrix Z can be written in the form

$$Z = \sum_{n,m=1}^N Z_{nm} U^n V^m \tag{8.2}$$

U and V may be thought of as exponentials of canonical variables p and q :

$$\begin{aligned} U &= e^{ip} \\ V &= e^{iq}. \end{aligned} \tag{8.3}$$

where p, q satisfy the commutation relations

$$[q, p] = \frac{2\pi i}{N}. \tag{8.4}$$

From (8.2) we see that only periodic functions of p and q are allowed. Thus the space defined by these variables is a torus. In fact there is an illuminating interpretation of these coordinates in terms of the quantum mechanics of particles on a torus in a strong background magnetic field. The coordinates of the particle are p, q . If the field is strong enough the existence of a large gap makes it useful to truncate the space of states to the finite dimensional subspace of lowest Landau levels. On this subspace the commutation relations (8.4) are satisfied. The lowest Landau wave packets form minimum uncertainty

⁶ We are grateful to M. Green for pointing out this paper to us when a preliminary version of this work was presented at the Santa Barbara Strings 96 conference.

packets which occupy an area $\sim 1/N$ on the torus. These wave packets are analogous to the “Planckian cells” which make up quantum phase space. The p, q space is sometimes called the noncommuting torus, the quantum torus or the fuzzy torus. In fact, for large enough N we can choose other bases of N dimensional Hilbert space which correspond to the lowest Landau levels of a charged particle propagating on an arbitrary Riemann surface wrapped by a constant magnetic field. For example in [5], de Wit *et al.* construct the finite dimensional Hilbert space of lowest Landau levels on a sphere. This connection between finite matrix models and two dimensional surfaces is the basis for the fact that the large N matrix model contains membranes. For finite N , the model consists of maps of quantum Riemann surfaces into a noncommuting transverse superspace, *i.e.* it is a model of a noncommuting membrane embedded in a noncommutative space ⁷.

In the limit of large N the quantum torus behaves more and more like classical phase space. The following correspondences connect the two:

1. The quantum operators Z defined in (8.2) are replaced by their classical counterparts. Eq. (8.2) becomes the classical Fourier decomposition of a function on phase space.
2. The operation of taking the trace of an operator goes over to N times the integral over the torus.

$$\text{tr} Z \rightarrow N \int Z(p, q) dp dq \quad (8.5)$$

3. The operation of commuting two operators is replaced by $1/N$ times the classical Poisson bracket

$$[Z, W] \rightarrow \frac{1}{N} [\partial_q Z \partial_p W - \partial_q W \partial_p Z] \quad (8.6)$$

We may now use the above correspondence to formally rewrite the matrix model Lagrangian. We begin by representing the matrices Y^i and θ as operator functions $Y^i(p, q), \theta(p, q)$. Now apply the correspondences to the two terms in (4.2). This gives

$$L = \frac{p_{11}}{2} \int dp dq Y^i(p, q)^2 - \frac{1}{p_{11}} \int dp dq [\partial_q Y^i \partial_p Y^j - \partial_q Y^j \partial_p Y^i]^2 + \text{fermionic terms} \quad (8.7)$$

⁷ Note that it is clear in this context that membrane topology is not conserved by the dynamics. Indeed for fixed N a given matrix can be thought of as a configuration of many different membranes of different topology. It is only in the large N limit that stable topological structure may emerge in some situations.

and a Hamiltonian

$$H = \frac{1}{2p_{11}} \int dpdq \Pi_i(p, q)^2 + \frac{1}{p_{11}} \int dpdq [\partial_q Y^i \partial_p Y^j - \partial_q Y^j \partial_p Y^i]^2 + \text{fermionic terms} \quad (8.8)$$

Equation (8.8) is exactly the standard Hamiltonian for the 11-D supermembrane in the light cone frame. The construction shows us how to build configurations in the matrix model which represent large classical membranes. To do so we start with a classical embedding of a toroidal membrane described by periodic functions $Y^i(p, q)$. The Fourier expansion of these functions provides us with a set of coefficients Y_{mn}^i . Using (8.2) we then replace the classical Y 's by operator functions of U, V . The resulting matrices represent the large classical membranes.

If the matrix model membranes described above are to correspond to M-theory membranes their tensions must agree. Testing this involves keeping track of the numerical factors of order one in the above discussion. We present this calculation in Appendix B where we show that the matrix model membrane tension exactly agrees with the M-theory membrane tension. This has also been verified by Berkooz and Douglas [20] using a different technique.

We do not expect static finite energy membranes to exist in the uncompactified limit. Nevertheless let us consider the conditions for such a static solution. The matrix model equations of motion for static configurations is

$$[Y^i, [Y^j, Y^i]] = 0. \quad (8.9)$$

It is interesting to consider a particular limiting case of an infinite membrane stretched out in the 8, 9 plane for which a formal solution of (8.9) can be found. We first rescale p and q

$$\begin{aligned} P &= \sqrt{N}p \\ Q &= \sqrt{N}q \\ [Q, P] &= 2\pi i \end{aligned} \quad (8.10)$$

In the $N \rightarrow \infty$ limit the P, Q space becomes an infinite plane. Now consider the configuration

$$\begin{aligned} Y^8 &= R_8 P \\ Y^9 &= R_9 Q \end{aligned} \quad (8.11)$$

with all other $Y^i = 0$. R_i is the length of the corresponding direction, which should of course be taken to infinity. Since $[Y^8, Y^9]$ is a c-number, eq (8.9) is satisfied. Thus we find the necessary macroscopic membranes require by M-theory. This stretched membrane has the requisite “wrapping number” on the infinite plane. On a general manifold one might expect the matrix model version of the wrapping number of a membrane on a two cycle to be

$$W = \frac{1}{N} \text{Tr } \omega_{ij}(X^i(p, q)) [X^i, X^j] \quad (8.12)$$

where ω is the two form associated to the cycle. This expression approaches the classical winding number as we take the limit in which Poisson brackets replace commutators.

Another indication that we have found the right representation of the membrane comes from studying the supersymmetry transformation properties of our configuration⁸. The supermembrane should preserve half of the supersymmetries of the model. The SUSY transformation of the fermionic coordinates is

$$\delta\theta = (P^i \gamma_i + [X^i, X^j] \gamma_{ij}) \epsilon + \epsilon' \quad (8.13)$$

For our static membrane configuration, $P^i = 0$, and the commutator is proportional to the unit matrix, so we can choose ϵ' to make this variation vanish. The unbroken supergenerators are linear combinations of the IMF q_A and Q_α .

It is interesting to contemplate a kind of duality and complementarity between membranes and D0 branes. According to the standard light cone quantization of membranes, the longitudinal momentum p_{11} is uniformly distributed over the area of the p, q parameter space. This is analogous to the uniform distribution of p_{11} along the σ axis in string theory. As we have seen the p, q space is a noncommuting space with a basic indivisible quantum of area. The longitudinal momentum of such a unit cell is $1/N$ of the total. In other words the unit phase space cells that result from the noncommutative structure of p, q space are the D0 branes with which we began. The D0 branes and membranes are dual to one another. Each can be found in the theory of the other.

The two kinds of branes also have a kind of complementarity. As we have seen, the configurations of the matrix model which have classical interpretations in terms of D0 branes are those for which the Y 's commute. On the other hand the configurations of a membrane which have a classical interpretation are the extended membranes of large

⁸ This result was derived in collaboration with N. Seiberg, along with a number of other observations about supersymmetry in the matrix model, which will appear in a future publication.

classical area. The area element is the Poisson bracket which in the matrix model is the commutator. Thus the very classical membranes are highly nonclassical configurations of D0 branes.

In the paper of de Wit, Luscher, and Nicolai [6] a pathology of membrane theory was reported. It was found that the spectrum of the membrane Hamiltonian is continuous. The reason for this is the existence of the unlifted flat directions along which the commutators vanish. Previously it had been hoped that membranes would behave like strings and have discrete level structure and perhaps be the basis for a perturbation theory which would generalize string perturbation theory. In the present context this apparent pathology is exactly what we want. M-theory has no small coupling analogous to the string splitting amplitude. The bifurcation of membranes when the geometry degenerates is expected to be an order one process. The matrix model, if it is to describe all of M-theory must inextricably contain this process. In fact we have seen how important it is that supersymmetry maintains the flat directions. A model of a single noncommutative membrane actually contains an entire Fock space of particles in flat eleven dimensional space time.

Another pathology of conventional membrane theories which we expect to be avoided in M-theory is the nonrenormalizability of the membrane world volume field theory. For finite N , it is clear that ultraviolet divergences on the world volume are absent because the noncommutative nature of the space defines a smallest volume cell, just like a Planck cell in quantum mechanical phase space (but we should emphasize here that this is a classical rather than quantum mechanical effect in the matrix model). The formal continuum limit which gives the membrane Hamiltonian is clearly valid for describing the classical motion of a certain set of metastable semiclassical states of the matrix model. It should not be expected to capture the quantum mechanics of the full large N limit. In particular it is clear that the asymptotic supergraviton states would look extremely singular and have no real meaning in a continuum membrane formalism. We are not claiming here to have a proof that the large N limit of the matrix quantum mechanics exists, but only that the issues involved in the existence of this limit are not connected to the renormalizability of the world volume field theory of the membrane.

There is one last point worth making about membranes. It involves evidence for 11-D Lorentz invariance of the matrix model. We have considered in some detail the Galilean invariance of the infinite momentum frame and found that it is satisfied. But there is more to the Lorentz group. In particular there are generators J^i which in the light cone formalism rotate the light like surface of initial conditions. The conditions for invariance

under these transformations are the notorious angular conditions. We must also impose longitudinal boost invariance. The angular conditions are what makes Lorentz invariance so subtle in light cone string theory. It is clearly important to determine if the matrix model satisfies the angular conditions in the large N limit. In the full quantum theory the answer is not yet clear but at the level of the classical equations of motion the answer is yes. The relevant calculations were done by de Wit, Marquard and Nicolai [21]. The analysis is too complicated to repeat here but we can describe the main points.

The equations for classical membranes can be given in covariant form in terms of a Nambu-Goto type action. In the covariant form the generators of the full Lorentz group are straightforward to write down. In passing to the light cone frame the expressions for the nontrivial generators become more complicated but they are quite definite. In fact they can be expressed in terms of the $Y(p, q)$ and their canonical conjugates $\Pi(p, q)$. Finally, using the correspondence between functions of p, q and matrices we are led to matrix expressions for the generators. The expressions of course have factor ordering ambiguities but these, at least formally, vanish as $N \rightarrow \infty$. In fact according to [21] the violation of the angular conditions goes to zero as $1/N^2$. Needless to say a quantum version of this result would be very strong evidence for our conjecture.

We cannot refrain from pointing out that the quantum version of the arguments of [21] is apt to be highly nontrivial. In particular, the classical argument works for every dimension in which the classical supermembrane exists, while, by analogy with perturbative string theory, we only expect the quantum Lorentz group to be recovered in eleven dimensions. Further, the longitudinal boost operator of [21] is rather trivial and operates only on a set of zero mode coordinates, which we have not included in our matrix model. Instead, we expect the longitudinal boost generator to involve rescaling N in the large N limit, and thus to relate the Hilbert spaces of different SUSY quantum mechanics models. We have already remarked in the previous section that, as anticipated in [7], longitudinal boost invariance is the key problem in our model. We expect it to be related to a generalization of the conformal invariance of perturbative string theory.

9. Towards Derivation and Compactification

In this section we would like to present a line of argument which may lead to a proof of the conjectured equivalence between the matrix model and M-theory. It relies on a stringy

extension of the conjectured nonrenormalization theorem combined with the possibility that all velocities in the large N cluster go to zero as $N \rightarrow \infty$.

Imagine that R stays fixed as $N \rightarrow \infty$. Optimistically one might imagine that finite R errors are as small as in perturbative 11D supergravity, meaning that they are suppressed by powers of $k_{11}R$ or even $\exp(-k_{11}R)$ where k_{11} is the center of mass longitudinal momentum transfer. So for $k_{11} \sim 1/l_p$ we could imagine R fixed at a macroscopic scale and have very tiny errors. The mean field estimates discussed in section 7 give the velocity $v \sim N^{-1/9}R$, which, with R fixed, can be made arbitrarily small at large N . Although it is likely that the structure of the large N cluster is more complicated than the mean field description, it is possible that this general property of vanishing velocities at large N continues to hold. In particular, in Appendix A we present arguments that the velocities of the coordinates along some of the classical flat directions of the potential are small. We suspect that this can be generalized to all of the flat directions. If that is the case then the only high velocities, would be those associated with the “core” wave function of Appendix A. The current argument assumes that the amplitude for the core piece of the wave function vanishes in the large N limit.

Nonabelian field strength is the correct generalization of velocity for membrane type field configurations like those discussed in Section 8. For classical configurations at least these field strengths are order $1/N$ and so are also small.

We have previously conjectured that the v^4 terms in the quantum mechanical effective action are not renormalized beyond one loop. For computing the eleven dimensional supergravity amplitude, we needed this result in the matrix quantum mechanics, but is possible that this result holds in the full string perturbation theory. For example, the excited open string states can be represented as additional non-BPS fields in the quantum mechanics. These do not contribute to the one loop v^4 term because they are not BPS. Perhaps they do not contribute to higher loops for related reasons.

If these two properties hold then the conjecture follows. The scattering of large N clusters of D-0-branes can clearly be computed at small g (small R) using the quantum mechanics. But these processes, by assumption, only involve low velocities independent of g , and so only depend on the v^4 terms in the effective action which, by the stringy extension of the nonrenormalization theorem, would not receive g corrections. So the same quantum mechanical answers would be valid at large g (large R).

This would prove the conjecture.

From this point of view we have identified a subset of string theory processes (large N D-0-brane scattering) which are unchanged by stringy loop corrections and so are computable at strong coupling.

If this line of argument is correct, it gives us an unambiguous prescription for compactification. We take the quantum mechanics which describes 0-brane motion at weak string coupling in the compactified space and then follow it to strong coupling. This approach to compactification requires us to add extra degrees of freedom in the compactified theory. We will discuss an alternative approach in the next section.

For toroidal compactifications, it is clear at weak coupling that one needs to keep the strings which wrap any number of times around each circle. These unexcited wrapped string states are BPS so they do contribute to the v^4 term and hence must be kept. In fact these are the states which, in the annulus diagram, correct the power law in graviton scattering to its lower dimensional value, and are crucial in implementing the various T dualities.

To be specific let us discuss the case of one coordinate X^9 compactified on a circle on radius R_9 . Here we should keep the extra string winding states around X^9 . An efficient way to keep track of them is to T dualize the X^9 circle. This converts D-0-branes to D-1-branes and winding modes to momentum modes. The collection of N D-1-branes is described by a 1+1 dimensional $SU(N)$ Super Yang Mills quantum field theory with coupling $g_{\text{SYM}}^2 = R^2/(R_9 l_p^3)$ on a space of T dual radius $R_{\text{SYM}} = l_p^3/(R R_9)$. The dimensionless effective coupling of the Super Yang Mills theory is then $g_{\text{SYM}}^2 R_{\text{SYM}}^2 = (l_p/R_9)^3$ which is independent of R . For p -dimensional tori we get systems of D- p -branes described by $p+1$ dimensional $SU(N)$ Super Yang Mills theory. Related issues have also recently been discussed in [22].

For more general compactifications the rule would be to keep every BPS state which contributes to the v^4 term at large N . We are currently investigating such compactifications, including ones with less supersymmetry.

The line of argument presented in this section raises a number of questions. Is it permissible to hold R finite, or to let it grow very slowly with N ? Are there nonperturbative corrections to the v^4 term? Large N probably prohibits instanton corrections in the quantum mechanics, but perhaps not in the full string theory. This might be related to the effect of various wrapped branes in compactified theories.

Does the velocity stay low? A key problem here is that in the mean field theory cloud the 0-branes are moving very slowly. If two 0-branes encounter each other, their relative

velocity is much less than the typical velocity in a bound pair ($v \sim R$). It seems that the capture cross section to go into the pair bound state should be very large. Why isn't there clumping into pairs? One factor which might come into play is the following. If the velocity is very low the de Broglie wavelength of the particles might be comparable to the whole cluster (this is true in mean field), so there could be delicate phase correlations across the whole cluster—some kind of macroscopic quantum coherence. Whenever a pair is trying to form, another 0-brane might get between them and disrupt them. This extra coherent complexity might help explain the Lorentz invariance puzzle.

10. Another Approach to Compactification

The conjecture which we have presented refers to an exact formulation of M-theory in uncompactified eleven dimensional spacetime. It is tempting to imagine that we can regain the compactified versions of the theory as particular collections of states in the large N limit of the matrix model. There is ample ground for suspicion that this may not be the case, and that degrees of freedom that we have thrown away in the uncompactified theory may be required for compactification. Indeed, in IMF field theory the only general method for discussing theories with moduli spaces of vacua is implementable only when the vacua are visible in the classical approximation. Then, we can shift the fields and do IMF quantization of the shifted theory. Different vacua correspond to different IMF Hamiltonians for the same degrees of freedom. The proposal of the previous section is somewhat in the spirit of IMF field theory. Different Hamiltonians, and indeed, different sets of degrees of freedom, are required to describe each compactified vacuum.

We have begun a preliminary investigation of the alternative hypothesis, that different compactifications are already present in the model we have defined. This means that there must be collections of states which, in the large N limit, have S matrices which completely decouple from each other. Note that the large N limit is crucial to the possible existence of such superselection sectors. The finite N quantum mechanics cannot possibly have superselection rules. Thus, the only way in which we could describe compactifications for finite N would be to add degrees of freedom or change the Hamiltonian. We caution the reader that the approach we will describe below is very preliminary and highly conjectural.

This approach to compactification is based on the idea that there is a sense in which our system defines a single “noncommuting membrane”. Consider compactification of a membrane on a circle. Then there are membrane configurations in which the embedding

coordinates do not transform as scalars under large diffeomorphisms of the membrane volume but rather are shifted by large diffeomorphisms of the target space. These are winding states. A possible approach to identifying the subset of states appropriate to a particular compactification, is to first find the winding states, and then find all states which have nontrivial scattering from them in the large N limit. In fact, our limited study below seems to indicate that all relevant states, including compactified supergravitons, can be thought of as matrix model analogs of membrane winding states.

Let us consider compactification of the ninth transverse direction on a circle of radius $2\pi R_9$. A winding membrane is a configuration which satisfies

$$X^9(q, p + 2\pi) = X^9(q, p) + 2\pi R_9 \quad (10.1)$$

and the winding sector is defined by a path integral over configurations satisfying this boundary condition. A matrix analog of this is

$$e^{-iNq} X^9 e^{iNq} = X^9 + 2\pi R_9 \quad (10.2)$$

It is easy to see (by taking the trace) that this condition cannot be satisfied for finite N . However, if we take the large N limit in such a way that $q \rightarrow \frac{\sigma}{N} \otimes \mathbf{1}_{M \times M}$, with σ an angle variable, then this equation can be satisfied, with

$$X^9 = \frac{2\pi R_9}{i} \frac{\partial}{\partial \sigma} \otimes \mathbf{1}_{M \times M} + x^9(\sigma) \quad (10.3)$$

where x^9 is an MXM matrix valued function of the angle variable. The other transverse bosonic coordinates, and all of the θ s are MXM matrix valued functions of σ . These equations should be thought of as limits of finite matrices. Thus $2\pi R_9 \mathcal{P} \equiv \frac{2\pi R_9}{i} \frac{\partial}{\partial \sigma}$ can be thought of as the limit of the finite matrices $\text{diag}(-2\pi P R_9 \dots 2\pi P R_9)$, with σ the obvious tridiagonal matrix in this representation. The total longitudinal momentum of such a configuration is $\frac{(2P+1)M}{R}$, and the ratio M/P is an effectively continuous parameter characterizing the states in the large N limit. We are not sure of the meaning of this parameter.

To get a feeling for the physical meaning of this proposal, we examine the extreme limits of large and small R_9 . For large R_9 it is convenient to work in the basis where \mathcal{P} is diagonal. If we take all of the coordinates X^i independent of σ , then our winding membrane approaches a periodic array of $(2P + 1)$ collections of D0 branes, each with longitudinal momentum M/R . We can find a solution of the BPS condition by putting

each collection into the M zero brane threshold bound state wave function. For large R_9 configurations of the X^i which depend on σ have very high frequency, and can be integrated out. Thus, in this limit, the BPS state in this winding sector is approximately a periodic array of supergravitons. We identify this with the compactified supergraviton state. This state will have the right long range gravitational interactions (at scales larger than R_9) in the eight uncompactified dimensions. To obtain the correct decompactified limit it would appear that we must rescale R , the radius of the longitudinal direction by $R \rightarrow (2P + 1)R$ as we take P and M to infinity. With this rescaling, all trace of the parameter M/P seems to disappear in the decompactification limit.

For small R_9 , our analysis is much less complete. However, string duality suggests an approximation to the system in which we keep only configurations with $M = 1$ and $P \rightarrow \infty$. In this case the σ dependence of X^9 is pure gauge, and the X^i all commute with each other. The matrix model Hamiltonian becomes the Hamiltonian of the Green-Schwarz type IIA string:

$$H \rightarrow \int d\sigma (\dot{\mathbf{X}})^2 + \left(\frac{\partial \mathbf{X}}{\partial \sigma}\right)^2 + \theta^T \gamma_9 \frac{\partial \theta}{\partial \sigma} \quad (10.4)$$

As in previous sections, we will construct multi wound membrane states by making large block diagonal matrices, each block of which is the previous single particle construction. Lest such structures appear overly baroque, we remind the reader that we are trying to make explicit constructions of the wave functions of a strongly interacting system with an infinite number of degrees of freedom. For large R_9 it is fairly clear that the correct asymptotic properties of multiparticle states will be guaranteed by the BPS condition (assuming that everything works as conjectured in the uncompactified theory).

If our ansatz is correct for small R_9 , it should be possible to justify the neglect of fluctuations of the matrix variables away from the special forms we have taken into account, as well as to show that the correct string interactions (for multistring configurations defined by the sort of block diagonal construction we have used above) are obtained from the matrix model interactions. In this connection it is useful to note that in taking the limit from finite matrices, there is no meaning to the separation of configurations into winding sectors which we have defined in the formal large N limit. In particular, X^9 should be allowed to fluctuate. But we have seen that shifts of X^9 by functions of σ are pure gauge, so that all fluctuations around the configurations which we have kept, give rise to higher derivative world sheet interactions. Since the $P \rightarrow \infty$ limit is the world sheet continuum limit we should be able to argue that these terms are irrelevant operators in that limit. We have less

understanding about how the sum over world sheet topologies comes out of our formalism, but it is tempting to think that it is in some way connected with the usual topological expansion of large N matrix models. In the appendix we show that in eleven dimensions, dimensional analysis guarantees the dominance of planar graphs in certain calculations. Perhaps, in ten dimensions, the small dimensionless parameter $\frac{R_9}{l_p}$ must be scaled with a power of N in order to obtain the limit of the matrix model which gives IIA string theory.

These ideas can be extended to compactification on multidimensional tori. A wrapping configuration of a toroidal membrane can be characterized by describing the cycles on the target torus on which the a and b cycles of the membrane are mapped. This parametrization is redundant because of the $SL(2, Z)$ modular invariance which exchanges the two membrane cycles. We propose that the analogs of these wrapping states, for a d -torus defined by modding out R^d by the shifts $X^i \rightarrow X^i + 2\pi R_a^i$, is defined by the conditions

$$e^{-i\sigma} X_{(\mathbf{m}, \mathbf{n})}^i e^{i\sigma} = X_{(\mathbf{m}, \mathbf{n})}^i + 2\pi R_a^i n^a \quad (10.5)$$

$$e^{2\pi i \mathcal{P}} X_{(\mathbf{m}, \mathbf{n})}^i e^{-2\pi i \mathcal{P}} = X_{(\mathbf{m}, \mathbf{n})}^i + 2\pi R_a^i m^a \quad (10.6)$$

where \mathbf{n} and \mathbf{m} are d-vectors of integers. The solutions to these conditions are

$$X_{(\mathbf{m}, \mathbf{n})}^i = [2\pi R_a^i n^a \mathcal{P} + R_a^i m^a \sigma] \otimes 1_{MXM} + x_{(\mathbf{m}, \mathbf{n})}^i(\sigma) \quad (10.7)$$

where the x^i are periodic MXM matrix valued functions of σ . The fermionic, and non-compact, coordinates are also matrix valued functions of σ .

In order to discuss more complicated compactifications, we would have to introduce coordinates and find a group of large diffeomorphisms associated with one and two cycles around which membranes can wrap. Then we would search for embeddings of this group into the large N gauge group. Presumably, different coordinate systems would correspond to unitarily equivalent embeddings. We can even begin to get a glimpse of how ordinary Riemannian geometry would emerge from the matrix system. If we take a large manifold which breaks sufficient supersymmetries, the effective action for supergravitons propagating on such a manifold would be obtained, as before, by integrating over the off diagonal matrices. Now however, the nonrenormalization theorem would fail, and the kinetic term for the gravitons would contain a metric. The obvious conjecture is that this is the usual Riemannian metric on the manifold in question. If this is the case, our prescription for compactification in the noncommutative geometry of the matrix model, would reduce to ordinary geometry in the large radius limit.

A question which arises is whether the information about one and two cycles is sufficient to characterize different compactifications. We suspect that the answer to this is no. The moduli of the spaces that arise in string theoretic compactifications are all associated with the homology of the space, but in general higher dimensional cycles (*e.g.* three cycles in Calabi-Yau three folds) are necessary to a complete description of the moduli space. Perhaps in order to capture this information we will have to find the correct descriptions of five branes in the matrix model. If the theory really contains low energy SUGRA, then it will contain solitonic five branes, but it seems to us that the correct prescription is to define five branes as the D-branes of membrane theory. We do not yet understand how to introduce this concept in the matrix model.

Finally, we would like to comment on the relation between the compactification schemes of this and the previous sections. For a single circle, if we take P to infinity, and substitute the formula (10.3) into the matrix model Hamiltonian (as well as the prescription that all other coordinates and supercoordinates are functions of σ), then we find the Hamiltonian of $1 + 1$ dimensionally reduced $10D$ SYM theory in $A_0 = 0$ gauge, with x^9 playing the role of the spatial component of the vector potential. Thus the prescription of the previous section appears to be a particular rule for how the large N limit should be taken in the winding configurations we have studied here. P is taken to infinity first, and then M is taken to infinity. The relation between the two approaches is reminiscent of the Eguchi-Kawai[23] reduction of large N gauge theory. It is clear once again that much of the physics of the matrix model is buried in the subtleties of the large N limit. For multidimensional tori, the relationship between the formalisms of this and the previous section is more obscure.

11. Conclusions

Although the evidence we have given for the conjectured exact equivalence between the large N limit of supersymmetric matrix quantum mechanics and uncompactified 11 dimensional M-theory is not definitive, it is quite substantial. The evidence includes the following :

1. The matrix model has exact invariance under the Supergalilean group of the infinite momentum frame description of 11D Lorentz invariant theories.
2. Assuming the conventional duality between M-theory and IIA string theory, the matrix model has normalizable marginally bound states for any value of N . These states have

exactly the quantum numbers of the 11-D supergraviton multiplet. Thus the spectrum of single particle states is exactly that of M-theory.

3. As a consequence of supersymmetric nonrenormalization theorems, asymptotic states of any number of non-interacting supergravitons exist. These well separated particles propagate in a Lorentz invariant manner in 11 dimensions. They have the statistics properties of the SUGRA Fock space.
4. The matrix model exactly reproduces the correct long range interactions between supergravitons implied by 11-D supergravity, for zero longitudinal momentum exchange. This one loop result could easily be ruined by higher loop effects proportional to four powers of velocity. We believe that a highly nontrivial supersymmetry theorem protects us against all higher loop corrections of this kind.
5. By examining the pieces of the bound state wave function in which two clusters of particles are well separated from each other, a kind of mean field approximation, we find that the longest range part of the wave function grows with N exactly as required by the holographic principle. In particular the transverse density never exceeds one parton per Planck area.
6. The matrix model describes large classical membranes as required by M-theory. The membrane world volume is a noncommutative space with a fundamental unit of area analogous to the Planck area in phase space. These basic quanta of area are the original D0 branes from which the matrix model was derived. The tension of this matrix model membrane is precisely the same as that of the M-theory membrane.
7. At the classical level the matrix model realizes the full 11-D Lorentz invariance in the large N limit.

Of course many unanswered questions remain. Locality is extremely puzzling in this system. Longitudinal boost invariance, as we have stressed earlier, is very mysterious. Resolving this issue, perhaps by understanding the intricate dynamics it seems to require, will be crucial in deciding whether or not this conjecture is correct.

One way of understanding Lorentz invariance would be to search for a covariant version of the matrix model in which the idea of noncommutative geometry is extended to all of the membrane coordinates. An obvious idea is to consider functions of angular momentum operators and try to exploit the connection between spin networks and three dimensional diffeomorphisms. Alternatively one could systematically study quantum corrections to the angular conditions.

It is likely that more tests of the conjecture can be performed. In particular it should be possible to examine the large distance behavior of amplitudes with nonvanishing longitudinal momentum transfer and to compare them with supergravity perturbation theory.

It will be important to try to make precise the line of argument outlined in Section 9 that may lead to a proof of the conjecture. The approaches to compactification discussed in Sections 9 and 10 should be explored further.

If the conjecture is correct, it would provide us with the first well defined nonperturbative formulation of a quantum theory which includes gravitation. In principle, with a sufficiently big and fast computer any scattering amplitude could be computed in the finite N matrix model with arbitrary precision. Numerical extrapolation to infinite N is in principle, if not in practice, possible. The situation is much like that in QCD where the only known definition of the theory is in terms of a conjectured limit of lattice gauge theory. Although the practical utility of the lattice theory may be questioned, it is almost certain that an extrapolation to the continuum limit exists. The existence of the lattice gauge Hamiltonian formulation insures that the theory is unitary and gauge invariant.

One can envision the matrix model formulation of M-theory playing a similar role. It would, among other things, insure that the rules of quantum mechanics are consistent with gravitation. Given that the classical long distance equations of 11-D supergravity have black hole solutions, a Hamiltonian formulation of M-theory would, at last, lay to rest the claim that black holes lead to a violation of quantum coherence.

Appendix A.

In this appendix we will report on a preliminary investigation of the threshold bound state wave function of N zero branes in the large N limit. In general, we may expect a finite probability for the N brane bound state to consist of p clusters of $N_1 \dots N_p$ branes separated by large distances along one of the flat directions of the potential. We will try to take such configurations into account by writing a recursion relation relating the N cluster to a k and $N - k$ cluster. This relation automatically incorporates multiple clusters since the pair into which the original cluster is broken up will themselves contain configurations in which they are split up into further clusters. There may however be multiple cluster configurations which cannot be so easily identified as two such superclusters. We will ignore these for now, in order to get a first handle on the structure of the wave function.

The configuration of a pair of widely separated clusters has a single collective coordinate whose Lagrangian we have already written in our investigation of supergraviton scattering. The Lagrangian is

$$L = \frac{1}{2} \frac{k(N-k)}{N} v^2 + \frac{k(N-k)}{r^7} v^4 \quad (\text{A.1})$$

where r is the distance between the clusters and v is their relative velocity. By scaling, we can write the solution of this quantum mechanics problem as $\phi(\frac{r(k[N-k])^{2/9}}{N^{1/3}})$, where ϕ is the threshold bound state wave function of the Lagrangian

$$\frac{1}{2} v^2 + \frac{v^4}{r^7} . \quad (\text{A.2})$$

This solution is valid when $r \gg l_p$.

We are now motivated to write the recursion relation

$$\Psi_N = \Psi_N^{(c)} + \frac{1}{2} P \sum_{k=1}^{N-1} A_{N,k} \Psi_k \Psi_{N-k} \phi\left(\frac{r(k[N-k])^{2/9}}{N^{1/3}}\right) e^{-r T r W_k^\dagger W_k} \quad (\text{A.3})$$

Here we have chosen a gauge in order to make a block diagonal splitting of our matrices. Ψ_j is the exact normalized threshold bound state wave function for j zero branes. W_k are the off diagonal $k \times N - k$ matrices which generate interactions between the two clusters. P is the gauge invariant projection operator which rotates our gauge choice among all gauge equivalent configurations. The $A_{N,k}$ are normalization factors, which in principle we would attempt to find by solving the Schroedinger equation. $\Psi_N^{(c)}$ is the ‘‘core’’ wave function, which describes configurations in which all of the zero branes are at a distance

less than or equal to l_p from each other. We will describe some of its properties below. In this regime, the entire concept of distance breaks down, since the noncommuting parts of the coordinates are as large as the commuting ones.

The interesting thing which is made clear by this ansatz, is that the threshold bound state contains a host of internal distance scales, which becomes a continuum as $N \rightarrow \infty$. This suggests a mechanism for obtaining scale invariant behavior for large N , as we must if we are to recover longitudinal boost invariance. Note that the typical distance of cluster separation is largest as N goes to infinity when one of the clusters has only a finite number of partons. These are the configurations which give the $N^{1/9}$ behavior discussed in the text, which saturates the Bekenstein bound. By the uncertainty principle, these configurations have internal frequencies of the bound state $\sim N^{-2/9}$. Although these go to zero as N increases they are still infinitely higher than the energies of supergraviton motions and interactions, which are of order $1/N$. As in perturbative string theory, we expect that this association of the large distance part of the wave function with modes of very high frequency will be crucial to a complete understanding of the apparent locality of low energy physics.

As we penetrate further in to the bound state, we encounter clusters of larger and larger numbers of branes. If we look for separated clusters carrying finite fractions of the total longitudinal momentum, the typical separation falls as N increases. Finally, we encounter the core, $\Psi_N^{(c)}$, which we expect to dominate the ultimate short distance and high energy behavior of the theory in noncompact eleven dimensional spacetime.

It is this core configuration to which the conventional methods of large N matrix models, which have so far made no appearance in our discussion, apply. Consider first gauge invariant Green's functions of operators like $Tr X_i^{2k}$, where i is one of the coordinate directions. We can construct a perturbation expansion of these Green's functions by conventional functional integral methods. When the time separations of operators are all short compared to the eleven dimensional Planck time, the terms in this expansion are well behaved. We can try to resum them into a large N series. The perturbative expansion parameter (the analog of g_{YM}^2 if we think of the theory as dimensionally reduced Yang Mills theory) is $\frac{R^3}{l_p^6 E^3}$. Thus, the planar Green's functions are functions of $\frac{R^3 N}{l_p^6 E^3}$.

The perturbative expansion of course diverges term by term as $E \rightarrow 0$. If we imagine that, as suggested by our discussion above, these Green's functions should be thought of as measuring properties of the core wave function of the system, there is no physical origin for such an infrared divergence. If, as in higher dimensions, the infrared cutoff is found

already in the leading order of the $\frac{1}{N}$ expansion, then it must be of order $\omega_c \sim R l_p^{-2} N^{1/3}$. Note that this is much larger than any frequency encountered in our exploration of the parts of the wave function with clusters separated along a flat direction.

Now let us apply this result to the computation of the infrared divergent expectation values of single gauge invariant operators in the core of the bound state wave function. The idea is to evaluate the graphical expansion of such an expression with an infrared cutoff and then insert the above estimate for the cutoff to obtain the the correct large N scaling of the object. The combination of conventional large N scaling and dimensional analysis then implies that planar graphs dominate even though we are not taking the “gauge coupling” $R^3 l_p^{-6}$, to zero as we approach the large N limit. Dimensional analysis controls the otherwise unknown behavior of the higher order corrections in this limit. The results are

$$\langle \frac{1}{N} \text{tr} X^{2k} \rangle \sim N^{2k/3} \quad (\text{A.4})$$

$$\langle \frac{1}{N} \text{tr} [X^i, X^j]^{2k} \rangle \sim N^{4k/3} \quad (\text{A.5})$$

$$\langle \frac{1}{N} \text{tr} (\bar{\theta} [\gamma_i X^i, \theta])^{2k} \rangle \sim N^{4k/3} \quad (\text{A.6})$$

$$\langle \frac{1}{N} \text{tr} \dot{X}^{2k} \rangle \sim N^{4k/3} \quad (\text{A.7})$$

In the first of these expressions, X refers to any component of the transverse coordinates. In the second the commutator refers to any pair of the components. The final expression, whose lowest order perturbative formula has an ultraviolet divergence, is best derived by combining (A.5), (A.6) and the Schroedinger equation which says that the threshold bound state has zero binding energy. Note that these expressions are independent of R , the compactification radius of the eleventh dimension. This follows from a cancellation between the R dependence of the infrared cutoff, that of the effective coupling, and that of the scaling factor which relates the variables X to conventionally normalized Yang Mills fields.

The first of these equations says that the typical eigenvalue of any one coordinate matrix is of order $N^{1/3}$, much larger than the $N^{1/9}$ extension along the flat directions. The second tells us that this spectral weight lies mostly along the non flat directions. In conjunction the two equations can be read as a kind of “uncertainty principle of non-commutative geometry”. The typical size of matrices is controlled by the size of their commutator. The final equation fits nicely with our estimate of the cutoff frequency. The

typical velocity is such that the transit time of a typical distance ⁹is the inverse of the cutoff frequency. It is clear that the high velocities encountered in the core of the wave function could invalidate our attempt to derive the matrix model by extrapolating from weakly coupled string theory. We must hope that the overall amplitude for this part of the wave function vanishes in the large N limit, relative to the parts in which zero branes are separated along flat directions.

It is important to realize that these estimates do not apply along the flat directions, but in the bulk of the N^2 dimensional configuration space. In these directions, it does not make sense to multiply together the “sizes ” along different coordinate directions to make an area since the different coordinates do not commute. Thus, there is no contradiction between the growth of the wave function in nonflat directions and our argument that the size of the bound state in conventional geometric terms, saturates the Bekenstein bound.

Appendix B.

In this appendix we compute the matrix model membrane tension and show that it exactly agrees with the M-theory membrane tension. The useful summary in [24] gives the tension of a D-p-brane T_p in IIA string theory or M-theory as

$$T_p = \frac{(2\pi)^{1/2}}{g_s} (2\pi)^{-p/2} \left(\frac{1}{2\pi\alpha'} \right)^{\frac{1+p}{2}} \quad (\text{B.1})$$

where g_s is the fundamental string coupling and $1/2\pi\alpha'$ is the fundamental string tension.

The membrane tension T_2 is defined so that the mass \mathcal{M} of a stretched membrane of area A is given by $\mathcal{M} = T_2 A$. The mass squared for a light cone membrane with no transverse momentum described by the map $\mathcal{X}^i(\sigma_1, \sigma_2)$, $i = 1 \dots 9$ can be written

$$\mathcal{M}^2 = (2\pi)^4 T_2^2 \int_0^{2\pi} \frac{d\sigma_1}{2\pi} \int_0^{2\pi} \frac{d\sigma_2}{2\pi} \sum_{i < j} \{ \mathcal{X}^i, \mathcal{X}^j \}^2 \quad (\text{B.2})$$

where the Poisson bracket of two functions $A(\sigma_1, \sigma_2), B(\sigma_1, \sigma_2)$ is defined by

$$\{A, B\} \equiv \frac{\partial A}{\partial \sigma_1} \frac{\partial B}{\partial \sigma_2} - \frac{\partial B}{\partial \sigma_1} \frac{\partial A}{\partial \sigma_2}. \quad (\text{B.3})$$

⁹ in the space of eigenvalues, which in this noncommutative region, is not to be confused with the classical geometrical distance between D0 branes .

The coefficients in (B.2) are set by demanding that \mathcal{M}^2 for the map $\mathcal{X}^8 = \frac{\sigma_1}{2\pi}L$, $\mathcal{X}^9 = \frac{\sigma_2}{2\pi}L$ is given by $\mathcal{M}^2 = (T_2 L^2)^2$.

To understand the relation to the matrix model we write, as in section 8, the map \mathcal{X}^i as a fourier series

$$\mathcal{X}^i(\sigma_1, \sigma_2) = \sum_{n_1, n_2} x_{n_1 n_2}^i e^{i(n_1 \sigma_1 + n_2 \sigma_2)}. \quad (\text{B.4})$$

Then the corresponding matrices X^i are given by

$$X^i = \sum_{n_1, n_2} x_{n_1 n_2}^i U^{n_1} V^{n_2} \quad (\text{B.5})$$

where the matrices U, V are elements of $\text{SU}(N)$, have spectrum $\text{spec}(U) = \text{spec}(V) = \{1, \omega, \omega^2 \dots \omega^{N-1}\}$ and obey $UV = \omega VU$ where $\omega = \exp(2\pi i/N)$. In a specific basis, $U = \text{diag}(1, \omega, \omega^2 \dots \omega^{N-1})$ and V is a cyclic forward shift.

The scale of (B.5) is fixed since $\text{spec}(U)$ and $\text{spec}(V)$ go over as $N \rightarrow \infty$ to the unit circle $\exp(i\sigma)$. Note that \mathcal{X}^i is real so X^i is hermitian.

The dynamics of the matrix model is governed by the Lagrangian

$$\frac{T_0}{2} \text{Tr} \left(\sum_i \dot{X}^i \dot{X}^i + C \sum_{i < j} [X^i, X^j]^2 \right). \quad (\text{B.6})$$

The normalizations here are fixed by the requirement that (B.6) describe D-0-brane dynamics. The first term, for diagonal matrices describing D-0-brane motion, is just the nonrelativistic kinetic energy $\frac{m_0}{2} \sum_{a=1}^N v_a^2$ since the 0-brane mass $m_0 = T_0$. The rest energy of the system is just Nm_0 which in the M-theory interpretation is just p_{11} so

$$p_{11} = NT_0. \quad (\text{B.7})$$

The coefficient C is fixed by requiring that the small fluctuations around diagonal matrices describe harmonic oscillators whose frequencies are precisely the masses of the stretched strings connecting the 0-branes. This ensures that (B.6) reproduces long range graviton interactions correctly. Expanding (B.6) to quadratic order we find that $C = (1/2\pi\alpha')^2$.

The energy of a matrix membrane configuration with zero transverse momentum is given by the commutator term in (B.6). We can evaluate this commutator in a semiclassical manner at large N as in section 8 by introducing angular operators q with spectrum the interval $(0, 2\pi)$ and $p = \frac{2\pi}{Ni} \frac{\partial}{\partial q}$ with spectrum the discretized interval $(0, 2\pi)$ so that

$[p, q] = \frac{2\pi i}{N}$. The matrices U, V become $U = e^{ip}, V = e^{iq}$. By Baker-Campbell-Hausdorff we see $UV = \omega VU$. The formal \hbar in this algebra is given by $\hbar = \frac{2\pi}{N}$. Semiclassically we have

$$\begin{aligned} [X, Y] &\rightarrow i\hbar\{\mathcal{X}, \mathcal{Y}\} \\ Tr &\rightarrow \int_0^{2\pi} \int_0^{2\pi} \frac{dpdq}{2\pi\hbar}. \end{aligned} \quad (\text{B.8})$$

So we get

$$Tr[X^i, X^j]^2 \rightarrow -\frac{(2\pi)^2}{N} \int_0^{2\pi} \frac{d\sigma_1}{2\pi} \int_0^{2\pi} \frac{d\sigma_2}{2\pi} \{\mathcal{X}^i, \mathcal{X}^j\}^2. \quad (\text{B.9})$$

This commutator can also be evaluated for a given finite N matrix configuration explicitly with results that agree with (B.9) as $N \rightarrow \infty$.

Now we can perform the check. The value of the matrix model Hamiltonian on a configuration with no transverse momentum is

$$H = \frac{T_0}{2} \left(\frac{1}{2\pi\alpha'}\right)^2 \frac{(2\pi)^2}{N} \int_0^{2\pi} \frac{d\sigma_1}{2\pi} \int_0^{2\pi} \frac{d\sigma_2}{2\pi} \sum_{I < j} \{\mathcal{X}^i, \mathcal{X}^j\}^2. \quad (\text{B.10})$$

The conjecture interprets the matrix model Hamiltonian H as the infinite momentum frame energy $\sqrt{p_{11}^2 + \mathcal{M}^2} - p_{11} \simeq \frac{\mathcal{M}^2}{2p_{11}}$. So the matrix membrane mass squared is $\mathcal{M}_{\text{mat}}^2 = 2p_{11}H$. Using (B.7) we find

$$\mathcal{M}_{\text{mat}}^2 = T_0^2 \left(\frac{1}{2\pi\alpha'}\right)^2 (2\pi)^2 \int_0^{2\pi} \frac{d\sigma_1}{2\pi} \int_0^{2\pi} \frac{d\sigma_2}{2\pi} \sum_{i < j} \{\mathcal{X}^i, \mathcal{X}^j\}^2. \quad (\text{B.11})$$

From (B.2) we can now read off the matrix model membrane tension as

$$(T_2^{\text{mat}})^2 = T_0^2 \left(\frac{1}{2\pi\alpha'}\right)^2 \frac{1}{(2\pi)^2}. \quad (\text{B.12})$$

So we can write

$$\begin{aligned} \frac{T_2^2}{(T_2^{\text{mat}})^2} &= (2\pi)^2 (2\pi\alpha')^2 \left(\frac{T_2}{T_0}\right)^2 \\ &= (2\pi)^2 \left(\frac{1}{2\pi}\right)^2 \\ &= 1. \end{aligned} \quad (\text{B.13})$$

So the M-theory and matrix model membrane tensions exactly agree.

Acknowledgments

This work crystallized at the Strings 96 Conference in Santa Barbara, and a very preliminary version of it was presented there. We would like to thank the organizers and participants of this conference for providing us with a stimulating venue for the production of exciting physics. We would particularly like to thank M. Green for pointing out the relation to previous work on supermembranes, and B. de Wit and I. Bars for discussing some of their earlier work on this topic with us. We would also like to thank C. Bachas and M. Dine for conversations on nonrenormalization theorems, M. Dine for discussing his work with us, C. Thorn for discussions about light cone string theory, N. Seiberg for discussions about supersymmetry, and M. Berkooz, M. Douglas and N. Seiberg for discussions about membrane tension. T.B., W.F., and S.H.S. would like to thank the Stanford Physics Department for its hospitality while some of this work was carried out. The work of T.B. and S.H.S. was supported in part by the Department of Energy under Grant No. DE - FG02 - 96ER40959 and that of W.F. was supported in part by the Robert A. Welch Foundation and by NSF Grant PHY-9511632. L.S. acknowledges the support of the NSF under Grant No. PHY - 9219345.

References

- [1] M.J. Duff, P. Howe, T. Inami, K.S. Stelle, "Superstrings in D=10 from supermembranes in D=11", Phys. Lett. B191 (1987) 70;
M.J. Duff, J. X. Lu, "Duality rotations in membrane theory", Nucl. Phys. B347 (1990) 394 ;
M.J. Duff, R. Minasian, James T. Liu, "Eleven-dimensional origin of string/string duality: a one-loop test", Nucl. Phys. B452 (1995) 261 ;
C. Hull and P. K. Townsend, "Unity of superstring dualities", Nucl. Phys. B438 (1995) 109, hep-th/9410167 ;
P. K. Townsend, "The Eleven-Dimensional Supermembrane Revisited" Phys. Lett. B350 (1995) 184 hep-th/9501068 ;
"String-membrane duality in seven dimensions", Phys. Lett. B354 (1995) 247, hep-th/9504095 ;
C. Hull and P. K. Townsend, "Enhanced gauge symmetries in superstring theories", Nucl. Phys. B451 (1995) 525, hep-th/9505073;
E. Witten "String Theory Dynamics in Various Dimensions", Nucl.Phys. B443 (1995) 85 hep-th/9503124.
- [2] J.Polchinski, "Dirichlet Branes and Ramond-Ramond charges", Phys.Rev.Lett.75 (1995) 4724;
for a review see J.Polchinski, S.Chaudhuri, and C.V. Johnson, "Notes on D-Branes", ITP preprint NSF-ITP-96-003, hep-th/9602052 .
- [3] E.Witten, "Bound States of Strings and p-Branes Nucl.Phys. B 460, 1995, 335 hep-th/9510135.
- [4] E. Bergshoeff, E. Sezgin and P.K. Townsend Phys.Lett. B189 (1987) 75; Ann. Phys. (NY) 185 (1988) 330;
P.K. Townsend in "Superstring '88", proc.of the Trieste Spring School, eds M.B. Green, M.T. Grisaru, R. Iengo and A. Strominger (World Scient., 1989);
M.J. Duff, Class. Quantum Grav. 5 (1988) 189
- [5] B. de Wit, J. Hoppe and H. Nicolai, Nucl.Phys. B 305 [FS 23] (1988) 545
- [6] B. de Wit, M. Luscher and H. Nicolai, Nucl. Phys. B 320 (1989) 135
- [7] G.'t Hooft, "Dimensional Reduction in Quantum Gravity" Utrecht preprint THU-93/26 gr-qc/9310026;
L.Susskind, "The World as a Hologram" J. Math. Phys. 36 (1995) 6377, hep-th/940989.
- [8] S. H. Shenker, "Another Length Scale in String Theory?", hep-th/9509132.
- [9] U.H. Danielsson, G. Ferretti and B. Sundborg, "D-particle Dynamics and Bound States", hep-th/9603081;

- D. Kabat and P. Pouliot, “A Comment on Zero-Brane Quantum Mechanics”, hep-th/9603127;
- M.R. Douglas, D. Kabat, P. Pouliot and S. Shenker, “D-Branes and Short Distance in String Theory”, hep-th/9608024 .
- [10] P.K. Townsend, ”D-branes from M-branes”, Phys. Lett. B373 (1996) 68, hep-th/9512062.
 - [11] S. Weinberg, “Dynamics at Infinite Momentum”, Phys.Rev. 150 (1966) 1313;
J. Kogut and L.Susskind, “The Parton Picture of Elementary Particles”, Physics Reports 8 (1973) 75.
 - [12] C.B. Thorn, Proceedings of Sakharov Conf. on Physics Moscow (1991) 447-454, hep-th/9405069.
 - [13] J. D. Bekenstein, Phys.Rev. D 49 (1994), 6606; G. 't Hooft in ref [7].
 - [14] T. Banks and L.Susskind, “Brane - Antibrane Forces”, hep-th/9511194.
 - [15] M. Baake, P. Reinicke, V. Rittenberg, J. Mathm. Phys., 26, (1985), 1070; R. Flume, Ann. Phys. 164, (1985), 189; M.Claudson, M.B.Halpern, Nucl. Phys. B250, (1985), 689.
 - [16] A.A. Tseytlin, “On $SO(32)$ Heterotic - type I Superstring Duality in Ten Dimensions” Imperial /TP/95-96/6 hep-th/9510173;
“Heterotic - Type I Superstring Duality and Low -Energy Effective Actions,” Imperial /TP/95-96/16 hep-th/9512081.
 - [17] M. Dine, work in progress.
 - [18] D. Fairlie, P. Fletcher, & C. Zachos, J. Math. Phys. **31** (1990) 1088-1094.
 - [19] J. Hoppe, Int. J. Mod. Phys. A4 (1989) 5235.
 - [20] M. Berkooz and M. R. Douglas, “ Five-branes in M(atric) Theory,” hep-th/9610236.
 - [21] B.de Wit, V. Marquard and H. Nicolai “Area Preserving Diffeomorphisms and Supermembrane Lorentz Invariance” Commun. Math. Phys. 128, 39 (1990).
 - [22] W. Taylor, “D-brane field theory on compact spaces,” hep-th/9611042
 - [23] T.Eguchi, H.Kawai, *Phys. Rev. Lett.* **48**, 1063, (1982).
 - [24] S. de Alwis, “A Note on Brane Tension and M Theory,” hep-th/9607011.